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A Cicerone for the Physics of Charm

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Summary. — After briefly recapitulating the history of the charm quantum number we sketch the experimental environments and instruments employed to study the behaviour of charm hadrons and then describe the theoretical tools for treating charm dynamics. We discuss a wide range of inclusive production processes before analyzing the spectroscopy of hadrons with hidden and open charm and the weak lifetimes of charm mesons and baryons. Then we address leptonic, exclusive semileptonic and nonleptonic charm decays. Finally we treat $D^0 - \bar{D}^0$ oscillations and CP (and CPT) violation before concluding with some comments on charm and the quark-gluon plasma. We will make the case that future studies of charm dynamics – in particular of CP violation – can reveal the presence of New Physics. The experimental sensitivity has only recently reached a level where this could reasonably happen, yet only as the result of dedicated efforts.
This review is meant to be both a pedagogical introduction for the young scholar and a useful reference for the experienced researcher. We aim for a self-contained description of the fundamental features while providing a guide through the literature for more technical issues.

1. – Preface

"Physicists, colleagues, friends, lend us your ears – we have come to praise charm, not bury it!" We have chosen such a theatrical opening not merely to draw your attention to our review. We feel that charm's reputation – like Caesar's – has suffered more than its fair share from criticisms by people that are certainly honourable. Of course, unlike in Caesar's case the main charge against charm is not that it reaches for the crown; the charge against charm is one of marginality, i.e. that charm can teach us nothing of true consequence any longer: at best it can serve as a tool facilitating access to something of real interest – like beauty; at worst it acts as an annoying background – so goes the saying.

Our contention instead is:

- While charm of course had an illustrious past, which should not be forgotten and from which we can still learn,
- it will continue to teach us important lessons on *Standard Model (SM)* dynamics, some of which will be important for a better understanding of beauty decays, and
- the best might actually still come concerning manifestations of *New Physics*.

The case to be made for continuing dedicated studies of charm dynamics does *not* rest on a single issue or two: there are several motivations, and they concern a better understanding of various aspects of strong and weak dynamics.

In this article we want to describe the present state-of-the-art in experiment and theory for charm studies. We intend it to be a self-contained review in that all relevant concepts and tools are introduced and the salient features of the data given. Our emphasis will be on the essentials rather than technical points. Yet we will provide the truly dedicated reader with a Cicerone through the literature where she can find all the details. We sketch charm's place in the $\boldsymbol{SM}$ – why it was introduced and what its characteristics are – and the history of its discovery. Then we describe the basic features of the experimental as well as theoretical tools most relevant in charm physics. Subsequent chapters are dedicated to specific topics and will be prefaced with more to the point comments on the tools required in that context: production, spectroscopy and weak lifetimes.

We shall then address exclusive leptonic, semileptonic and nonleptonic transitions, before we cover $\boldsymbol{D^0 - \bar{D}^0}$ oscillations, CP violation and the onset of the quark-gluon plasma. This discussion prepares the ground for an evaluation of our present understanding; on that base we will make a case for future studies of charm physics.

2. – A Bit of History

2˙1. *Charm's Place in the Standard Model.* – Unlike for strangeness the existence of hadrons with the quantum number charm had been predicted for several specific reasons and thus with specific properties as well. Nevertheless their discovery came as a surprise to large parts or even most of the community [1].

Strangeness acted actually as a 'midwife' to charm in several respects. Extending an earlier proposal by Gell-Mann and Levy, Cabibbo [2] made the following ansatz in 1963 for the charged current

$$J_\mu^{(+)}[J_\mu^{(-)}] = \cos\theta_C \bar{d}_L\gamma_\mu u_L[\bar{u}_L\gamma_\mu d_L] + \sin\theta_C \bar{s}_L\gamma_\mu u_L[\bar{u}_L\gamma_\mu s_L] \tag{1}$$

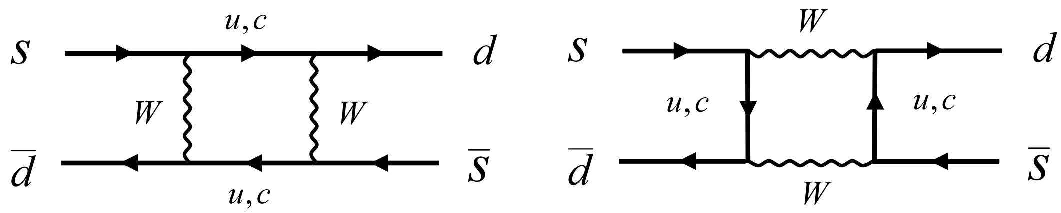

Fig. 1. – The box diagram responsible for $K^0 - \bar{K}^0$ oscillations

(written in today's notation), which successfully describes weak decays of strange and nonstrange hadrons. Yet commuting $J_\mu^{(+)}$ with its conjugate $J_\mu^{(-)}$ yields a *neutral* current that necessarily contains the $\Delta S = \pm 1$ term $\sin\theta_C \cos\theta_C\ (\bar{d}_L \gamma_\mu s_L + \bar{s}_L \gamma_\mu d_L)$. Yet such a strangeness changing neutral current (SChNC) is phenomenologically unacceptable, since it would produce contributions to ΔM_K and $K_L \to \mu^+\mu^-$ that are too large by several orders of magnitude. The match between leptons and quarks with three leptons – electrons, muons and neutrinos – and three quarks – up, down and strange – had been upset already in 1962 by the discovery that there were two distinct neutrinos. Shortly thereafter the existence of charm quarks was postulated to re-establish the match between the two known lepton families (ν_e, e) and (ν_μ, μ) with two quark families (u, d) and (c, s) [3, 4]. Later it was realized [5] that the observed huge suppression of strangeness changing neutral currents can then be achieved by adopting the form

$$\begin{gathered} J_\mu^{(+)} = \bar{d}_{C,L}\gamma_\mu u_L + \bar{d}_{C,L}\gamma_\mu c_L \\ d_C = cos\theta_C\, d + sin\theta_C\, s \quad , \quad s_C = -sin\theta_C\, d + cos\theta_C\, s \end{gathered} \tag{2}$$

for the charged current. The commutator of $J_\mu^{(+)}$ and $J_\mu^{(-)}$ contains neither a $\Delta S \neq 0$ nor a $\Delta C \neq 0$ piece. Even more generally there is no contribution to ΔM_K in the limit $m_c = m_u$; the GIM mechanism yields a suppression $\propto (m_c^2 - m_u^2)/M_W^2$. From the value of ΔM_K one infers $m_c \sim 2$ GeV.

This procedure can be illustrated by the quark box diagram for $K^0 - \bar{K}^0$ oscillations, Fig.(1). It is shown for a two-family scenario, since the top quark contribution is insignificant for Δm_K (though it is essential for ϵ_K).

To arrive at a renormalisable theory of the weak interactions one has to invoke non-abelian gauge theories [6]. In those the gauge fields couple necessarily to the charged currents and their commutators thus making the aforementioned introduction of charm quarks even more compelling. Yet one more hurdle had to be passed. For there is still one danger spot that could vitiate the renormalizability of the Standard Model. The so-called triangle diagram, see Fig.(2), has a fermion loop to which three external spin-one lines are attached – all axial vector or one axial vector and two vector: while by itself finite it creates an anomaly, the Adler-Bell-Jackiw (ABJ) anomaly. It means that the axial vector current even for *massless* fermions ceases to be conserved on the loop, i.e. quantum level (1). The thus induced nonconservation of the axial current even for massless fermions creates infinities in higher orders that cannot be removed in the usual

(1) The term 'anomaly' is generally applied when a *classical* symmetry is broken by quantum corrections.

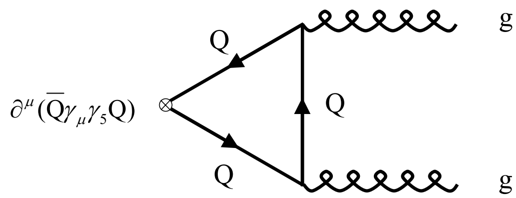

Fig. 2. – An example of a triangle diagram contributing to the ABJ anomaly.

way. The only way out is to have this anomaly, which does not depend on the mass of the internal fermions, cancel among the different fermion loops. Within the SM this requires the electric charges of all fermions – quarks and leptons - to add up to zero. With the existence of electrons, muons, up, down and strange quarks already established and their charges adding up to -2, this meant that a fourth quark with three colours was needed each with charge $+\frac{2}{3}$ – exactly like charm. There is an ironic twist here: as described below, the discovery of open charm hadrons was complicated and therefore delayed, because the charm threshold is very close to the τ lepton threshold; cancellation of the ABJ anomaly then required the existence of a third quark family (which in turn allows for CP violation to be implemented in the SM in charged current couplings).

The fact that charm 'bans' these evils is actually the origin of its name ([2]). It was the first quark flavour *pre*dicted, and even the salient features of charm quarks were specified:

- They possess the same couplings as u quarks,
- yet their mass is much heavier, namely about 2 GeV.
- They form charged and neutral hadrons, of which in the $C = 1$ sector three mesons and four baryons are stable; i.e., decay only weakly with lifetimes of very roughly 10^{-13} sec – an estimate obtained by scaling from the muon lifetime, as explained below.
- Charm decay produces direct leptons and preferentially strange hadrons.
- Charm hadrons are produced in deep inelastic neutrino-nucleon scattering.

Glashow reiterated these properties in a talk at EMS-74, the 1974 Conference on Experimental Meson Spectroscopy and concluded [7]:
"What to expect at EMS-76: There are just three possibilities:

1. Charm is not found, and I eat my hat.
2. Charm is found by hadron spectroscopers, and we celebrate.
3. Charm is found by outlanders, and you eat your hats."

([2]) The name "strangeness" refers to the feature – viewed as odd at the time – that the production rate of these hadrons exceeds their decay rate by many orders of magnitude.

A crucial element in the acceptance of the $SU(2)_L \times U(1)$ theory as the SM for the electroweak forces was the observation of flavour-conserving neutral currents by the Gargamelle collab. at CERN in 1973. Despite this spectacular success in predicting weak neutral currents of normal strength in the flavour-conserving sector together with hugely suppressed ones for $\Delta S \neq 0$ transitions, the charm hypothesis was not readily accepted by the community – far from it. Even after the first sightings of charm hadrons were reported in cosmic ray data [8], a wide spread sentiment could be characterized by the quote: "Nature is smarter than Shelly [Glashow] ... she can do without charm." ([3]) In the preface we have listed three categories of merits that charm physics can claim today. Here we want to expand on them, before they will be described in detail in subsequent sections.

- The production and decays of strange hadrons revealed or at least pointed to many features central to the SM, like parity violation, the existence of families, the suppression of flavour-changing neutral currents and CP violation. Charm physics was likewise essential for the development of the SM: its foremost role has been to confirm and establish most of those features first suggested by strange physics and thus pave the way for the acceptance of the SM. It did so in dramatic fashion in the discovery of charmonium, which together with the observation of Bjorken scaling in deep inelastic electron-nucleon scattering revealed quarks acting as dynamical degrees of freedom rather than mere mathematical entities. The demands of charm physics drove several lines in the development of accelerators and detectors alike. The most notable one is the development of microvertex detectors: they found triumphant application in charm as well as in beauty physics – they represent a conditio sine qua non for the observation of CP violation in $B \to J/\psi K_S$ – and in the discovery of top quarks through b-flavour tagging, to be followed hopefully soon by the discovery of Higgs bosons again through b-flavour tagging. Some might scoff at such historical merits. We, however, see tremendous value in being aware of the past – maybe not surprisingly considering where two of us live and the other two would love to live (we are not referring to South Bend here.).
- The challenge of treating charm physics *quantitatively* has lead to testing and refining our theoretical tools, in particular novel approaches to QCD based on heavy quark ideas. This evolutionary process will continue to go on. The most vibrant examples are lattice QCD and heavy quark expansions described later.
- Charm can still 'come through' as the harbinger or even herald of New Physics. It is actually qualified to do so in a unique way, as explained in the next section.

2·2. *On the Uniqueness of Charm.* – Charm quarks occupy a unique place among up-type quarks. *Top* quarks decay before they can hadronize [9], which, by the way, makes searches for CP violation there even more challenging. On the other end of the mass spectrum there are only two weakly decaying light flavour hadrons, namely the neutron and the pion: in the former the d quark decays and in the latter the quarks of the first family annihilate each other. The charm quark is the only up-type quark whose hadronization and subsequent weak decay can be studied. Furthermore the charm quark mass m_c provides a new handle on treating nonperturbative dynamics through an expansion in powers of $1/m_c$.

([3]) It seems, even Glashow did not out rule this possibility, see item 1 on his list above.

Decays of the down-type quarks s and b are very promising to reveal new physics since their CKM couplings are greatly suppressed, in particular for beauty. This is not the case for up-type quarks. Yet New Physics beyond the SM could quite conceivably induce flavour changing neutral currents that are larger for the up-type than the down-type quarks. In that case charm decays would be best suited to reveal such non-standard dynamics.

2˙3. *The Discovery of Charm.* –

2˙3.1. The heroic period. A candidate event for the decay of a charm hadron was first seen in 1971 in an emulsion exposed to cosmic rays [8]. It showed a transition $X^{\pm} \to h^{\pm}\pi^{0}$ with $h^{\pm}$ denoting a charged hadron that could be a meson or a baryon. It was recognized that as the decaying object $X^{\pm}$ was found in a jet shower, it had to be a hadron; with an estimated lifetime around few$\times 10^{-14}$ sec it had to be a weak decay. Assuming $h^{\pm}$ to be a meson, the mass of $X^{\pm}$ was about 1.8 GeV. The authors of Ref.[10] analyzed various interpretations for this event and inferred selection rules like those for charm. It is curious to note that up to the time of the J/ψ discovery 24 papers published in the Japanese journal *Prog. Theor. Physics* cited the emulsion event versus only 8 in Western journals; a prominent exception was Schwinger in an article on neutral currents [11]. The imbalance was even more lopsided in experimental papers: while about twenty charm candidates had been reported by Japanese groups before 1974, western experimentalists were totally silent [12].

It has been suggested that Kobayashi and Maskawa working at Nagoya University in the early 70's were encouraged in their work – namely to postulate a third family for implementing CP violation – by knowing about Niu's candidate for charm produced by cosmic rays. Afterwards the dams against postulating new quarks broke and a situation arose that can be characterized by adapting a well-known quote that "... Nature repeats itself twice, ... the second time as a farce".

It was pointed out already in 1964 [13] that charm hadrons could be searched in multilepton events in neutrino production. Indeed evidence for their existence was also found by interpreting opposite-sign dimuon events in deep inelastic neutrino nucleon scattering [14] as proceeding through $\nu N \to \mu^- c + ... \to \mu^- D ... \to \mu^- \mu^+ ...$.

2˙3.2. On the eve of a revolution. The October revolution of '74 – like any true one – was preceded by a period where established concepts had to face novel challenges, which created active fermentation of new ideas, some of which lead us forward, while others did not. This period was initiated on the one hand by the realization that spontaneously broken gauge theories are renormalizable, and on the other hand by the SLAC-MIT study of deep inelastic lepton-nucleon scattering. The discovery of approximate Bjorken scaling gave rise to the parton model to be superseded by QCD; the latter's 'asymptotic freedom' – the feature of its coupling $\alpha_S(Q^2)$ going to zero (logarithmically) as $Q^2 \to \infty$ – was just beginning to be appreciated.

Attention was turned to another deep inelastic reaction, namely $e^+e^- \to had$. In some quarters there had been the expectation that this reaction would be driven merely by the tails of the vector mesons ρ, ω and ϕ leading to a cross section falling off with the c.m. energy faster than the $1/E^2_{c.m.}$ dependence of the cross section for the 'point like' or 'scale-free' process $e^+e^- \to \mu^+\mu^-$ does. On the other hand it was already known at that time that within the quark-parton model the transition $e^+e^- \to had$ would show the same scale-free behaviour at sufficiently high energies leading to the ratio $R = \sigma(e^+e^- \to had)/\sigma(e^+e^- \to \mu^+\mu^-)$ being a constant given by the sum of the quark electric charges squared. The three known quarks $\boldsymbol{u}$, $\boldsymbol{d}$ and $\boldsymbol{s}$ yield $\boldsymbol{R = 2/3}$. It was pointed out by theorists that having three colours would raise $\boldsymbol{R}$ to a value $\boldsymbol{2}$. Yet the data seemed to paint a different picture. Data taken at the ADONE storage ring in Frascati yielded $\boldsymbol{R \sim 3 \pm 1}$ at $\boldsymbol{E_{c.m.} = 3}$ GeV. The old Cambridge Electron Accelerator (CEA) in Massachusetts was converted to an $\boldsymbol{e^+e^-}$ machine in 1972. Measurements made there showed no signs of $\boldsymbol{R}$ decreasing: $\boldsymbol{R = 4.9 \pm 1.1}$ and $\boldsymbol{6.2 \pm 1.6}$ at $\boldsymbol{E_{c.m.} = 4}$ and $\boldsymbol{5}$ GeV, respectively. Yet these findings were not widely accepted as facts due to the low acceptance of the detectors. The first measurement of $\boldsymbol{e^+e^-}$ annihilation with a large acceptance detector was performed by the MARK I collaboration at SLAC's SPEAR storage ring for $\boldsymbol{E_{c.m.} \sim 3-5}$ GeV. When their initial results were announced at the end of 1973, they caused quite a stir or even shock. They established that $\boldsymbol{R}$ was indeed in the range of $\boldsymbol{2-4}$ and not falling with energy. The publicly presented data with their sizeable error bars actually seemed to show $\boldsymbol{R}$ rising like $\boldsymbol{E^2_{c.m.}}$ meaning $\boldsymbol{\sigma(e^+e^- \to had)}$ approaching a constant value [15]. This was taken by some, including a very prominent experimentalist, as possible evidence for electrons containing a small hadronic core.

The '74 revolution thus shares more features with other revolutions: In the end it did not produce the effect that had emerged first; furthermore even prominent observers do not own a reliable crystal ball for gazing into the future. Rather than revealing that electrons are hadrons at heart, it showed that quarks are quite similar to leptons at small distances.

The New Physics invoked to induce the rise in $\boldsymbol{R}$ was parameterized through four-fermion operators built from quark and lepton bilinears. Some amusing effects were pointed out [16]: if the new operators involved scalar [pseudoscalar] fermion bilinears, one should see $\boldsymbol{\sigma(e^+e^- \to had)}$ decrease [increase] *with time* from the turn-on of the beams. For in that case the cross section would depend on the transverse polarization of the incoming leptons, and the latter would grow with time due to synchrotron radiation. Later more precise data did away with these speculations. They showed $\boldsymbol{R}$ to change with $\boldsymbol{E_{c.m.}}$ as expected from crossing a production threshold.

Other theoretical developments, however, turned out to be of lasting value. In a seminal 1973 paper [17] M.K. Gaillard and B. Lee explored in detail how charm quarks affect kaon transitions – $\boldsymbol{K^0 - \bar{K}^0}$ oscillations, $\boldsymbol{K_L \to \mu^+\mu^-}$, $\boldsymbol{K_L \to \gamma\gamma}$ etc. – through quantum corrections. Their findings firmed up the bound $\boldsymbol{m_c \leq 2}$ GeV. Together with J. Rosner they extended the analysis in a review, most of which was written in the summer of 1974, yet published in April 1975 [18] with an appendix covering the discoveries of the fall of 1974. At the same time it was suggested [19] that charm and anticharm quarks form unusually narrow vector meson bound states due to gluons carrying colour and coupling with a strength that decreases for increasing mass scales.

The theoretical tools were thus in place to deal with the surprising observations about to be made.

2˙3.3. The October revolution of '74. It is fair to say that the experimental signatures described above did not convince the skeptics – they needed a Damascus experience to turn from 'Saulus' into 'Paulus', from disbelievers into believers. Such an experience was provided by the October revolution of 1974, the discovery of the $\boldsymbol{J/\psi}$ and $\boldsymbol{\psi'}$ viewed as absurdly narrow at the time. It provides plenty of yarn for several intriguing story lines [1]. One is about the complementarity of different experiments, one about the value of persistence and of believing in what one is doing and there are others more. On the conceptual side these events finalized a fundamental change in the whole outlook of the community onto subnuclear physics that had been initiated a few years earlier, as sketched above: it revealed quarks to behave as real dynamical objects rather than to represent merely mathematical entities.

One exotic explanation that the J/ψ represents an $\Omega\bar{\Omega}$ bound state fell by the wayside after the discovery of the ψ'. The two leading explanations for the new threshold were charm production and 'colour thaw'. Since the early days of the quark model there were two types of quarks, namely the Gell-Mann-Zweig quarks with fractional charges and the Han-Nambu [20] quarks with integer charges. Of those there are actually nine grouped into three triplets, of which two contained two neutral and one charged quark and the last one two charged and one neutral quark. The Han-Nambu model was actually introduced to solve the spin-statistics problem of baryons being S-wave configuration of three quarks. The idea of 'colour thaw' is to assume that up to a certain energy each of the three triplets acts *coherently* reproducing results as expected from Gell-Mann-Zweig quarks, i.e. $R = 2$. Above this energy those 'colour' degrees of freedom get liberated to act *incoherently* as nine quarks producing $R = 4$!

Charm gained the upper hand since it could provide a convincing explanation for the whole family of narrow resonances as 'ortho-' and 'para-charmonia' in a dramatic demonstration of QCD's asymptotic freedom. 'Colour thaw' could not match that feat.

Yet the final proof of the charm hypothesis had to be the observation of open charm hadrons. In one of the (fortunately) rare instances of nature being malicious, it had placed the $\tau^+\tau^-$ threshold close to the charm threshold. Typical signatures for charm production – increase production of strange hadrons and higher multiplicities in the final state – were counteracted by $\tau^+\tau^-$ events, the decays of which lead to fewer kaons and lower hadronic multiplicities. It took till 1976 till charm hadrons were observed in fully reconstructed decays.

2˙3.4. The role of colour. The need for the quantum number 'colour' had arisen even before the emergence of QCD as the theory for the strong interactions. On the one hand there was the challenge of reconciling Fermi-Dirac statistics with identifying the Ω^- baryon as an sss system in the symmetric $J = 3/2$ combination: having colour degrees of freedom would allow for the wavefunction being odd under exchange for an S-wave configuration. On the other hand the aforementioned avoidance of the ABJ anomaly implied the existence of three colours for the quarks.

'Colour' is of course central to QCD. Its introduction as part of a non-abelian gauge theory is required by the need for a theory combining asymptotic freedom in the ultraviolet and confinement in the infrared. With three colours qqq combinations can form colour singlets.

It should be noted that studying $e^+e^- \to hadrons$ around the charm threshold revealed several other manifestations of colour:

(i) It had been noted before the discovery of the J/ψ that three colours for quarks are needed to also accommodate the observed value of $R = \frac{\sigma(e^+e^- \to had.)}{\sigma(e^+e^- \to \mu^+\mu^-)}$ within quark dynamics. Yet this argument was not viewed as convincing till data indeed showed that R below and (well) above the charm threshold could be adequately described by two 'plateaus' – i.e. relatively flat functions of the c.m. energy – with their difference in height approximately $N_C \sum_i e_i^2 = 4/3$.

(ii) The amazingly narrow width of the J/ψ resonance can be ascribed naturally to the fact that the decay of this ortho-charmonium state to lowest order already requires the $c\bar{c}$ to annihilate into three gluons making the width proportional to α_S^3. It is amusing to remember that one of the early competitors to the $c\bar{c}$ explanation for the J/ψ was

the speculation that the colour symmetry is actually broken leading to the existence of non-colour singlets in the hadronic spectrum.
(iii) The lifetime of τ leptons is reproduced correctly by scaling it from the muon lifetime $\tau_\tau \simeq \tau_\mu \cdot \left(\frac{m_\mu}{m_\tau}\right)^5 \cdot \frac{1}{2+N_C}$ with $N_C = 3$; $N_C = 2$ or 4 would not do. Likewise for the prediction of the leptonic branching ratio BR$(\tau \to e\nu\bar{\nu}) \simeq \frac{1}{2+N_C} = 0.2$ for $N_C = 3$. This is remarkably close to the experimental number $BR(\tau \to e\nu\bar{\nu}) \simeq 0.1784$ with the difference understood as due to the QCD radiative corrections. **(iv)** Similar estimates were made concerning the lifetime and semileptonic branching ratio for charm. Yet the former is a rather iffy statement in view of $\tau_c \propto m_c^{-5}$ and the complexity of defining a charm *quark* mass. The latter, which argues in favour of $BR(c \to e\nu s) \sim 1/(2+N_C)$ (again modulo QCD radiative corrections) is actually fallacious if taken at face value. These two points will be explained in Sect. **6**˙4.

3. – Experimental Environments and Instruments

The birth of the charm paradigm and its experimental confirmation fostered a time of development in experimental techniques, which has few parallels in the history of high energy. For charm was predicted with a set of properties that facilitate their observation. Its mass was large by the times' standards, but within reach of existing accelerators. It possessed charged current couplings to d and s quarks, and therefore should be visible in neutrino beams available then; e^+e^- colliders had come into operation. Open charm would decay preferentially to final states with strangeness, making them taggable by particle ID detectors able to discriminate kaons from protons and pions. Hidden charm states would have a large decay rate to lepton pairs providing a clean and signature. Charm lifetimes would be small, but within reach experimentally. Charm would decay semileptonically, thus providing chances of observing the relatively easy to detect muon.

In this section we will retrace the historical development, from which we will draw lessons on the production environments - focusing on various colliders versus fixed target set-ups - and then sketch key detector components.

3˙1. *On the history of observing charm.* –

3˙1.1. Hidden charm. The J/ψ was discovered simultaneously 1974 by two experiments, one at the Brookhaven fixed target machine with 30 GeV protons and the other one at SLAC's SPEAR e^+e^- collider, neither of which was actually searching for charm. Ting's experiment studying $pBe \to e^+e^- + X$, after having been rejected at Fermilab and CERN, was approved at BNL to search for the possible existence of a *heavy photon*, i.e., a higher mass recurrence of the ρ, ϕ, and ω mesons. Richter's group at SPEAR on the other hand was interested in the energy dependence of e^+e^- annihilation into hadrons. In 1974 Ting's group observed a sharp enhancement at $M(e^+e^-) = 3.1$ GeV. They did not announce the result waiting some months to confirm it. Finally they went public together with Richter's SLAC-LBL experiment, which observed a sharp resonant peak at the same energy in the interactions $e^+e^- \to \mu^+\mu^-, e^+e^-$. The ADONE e^+e^- collider at Frascati found itself in the unfortunate circumstance of having been designed for a maximum center-of-mass energy of 3.0 GeV. Immediately after the news of the J/ψ observation was received, currents in ADONE magnets were boosted beyond design limits, a scan in the 3.08-3.12 GeV was carried on and the new resonance found and confirmed. Three papers [21],[22], [23], announcing the J/ψ discovery appeared in early

December 1974 in Physical Review Letters ([4]) Within ten days of the announcement of the J/ψ's discovery the SLAC-LBL group at SPEAR found another narrow resonance, the ψ' at 3.7 GeV [25]. Soon thereafter other actors entered the stage, namely DESY's DORIS storage ring, where the DASP collaboration found a resonance just above charm threshold, the ψ'' at 3.77 GeV [26]. Over the years a very rich and gratifying experimental program was pursued at SPEAR and DORIS by a succession of experiments: MARK I - III, Crystal Ball, DASP, PLUTO etc. Their achievements went well beyond mapping out charmonium spectroscopy in a detailed way: a host of new experimental procedures was established – actually a whole style of doing physics at a heavy flavour 'factory' was born that set the standards for the later B factories.

Only charmonium states with $J^{PC} = 1^{--}$ can be produced *directly* in e^+e^- to lowest order in α. A novel technique was developed allowing the formation of other states as well, namely through low energy $\bar{p}p$ annihilation . This was pioneered at CERN by experiment R704 using a $\bar{p}$ beam on a gas jet target. It led to greatly enhanced accuracy in measuring masses and widths of $\chi_{c1,2}$ states [27]. The same technique was later used by Fermilab experiment E760 and its successor E835.

The shutdown of SPEAR and the upgrade of DORIS to study B physics created a long hiatus in this program, before it made a highly welcome comeback with the BES program and now with CLEO-c.

3˙1.2. Open charm. Hadrons with open charm had to be found before charm could be viewed as the established explanation for the J/ψ. Indirect evidence for their existence surfaced in neutrino experiments. An event apparently violating the $\Delta Q = -\Delta S$ rule was detected at Brookhaven [29], and opposite-sign dimuon events were observed as well [14, 30]. At CERN neutrino-induced $\mu^- e^+ V^0$ events were seen [31, 32] indicating that the new resonance was correlated with strangeness in weak reactions as required by the presence of charm.

An intense hunt for finding charm hadrons at accelerators was begun ([5]); the MARK I collaboration found the prey through narrow mass peaks in $K^-\pi^+, K^-\pi^+\pi^+, K^-\pi^+\pi^+\pi^-$ [34, 35] for the iso-doublet D^0 and D^+, i.e. in final states that had been predicted [18]. D mesons were soon thereafter detected also in neutrino- [36], hadron- [37] and photon-induced [38] reactions.

3˙1.3. Measuring charm lifetimes. Not surprisingly, the first experimental evidence for weakly decaying charm hadrons was obtained in an emulsion experiment exposed to cosmic rays [8], Fig. 3. For till after the time of the J/ψ discovery only photographic emulsions could provide the spatial resolution needed to find particles with lifetimes of about 10^{-13} sec. Their resolving power of about 1 micron was a very powerful tool for tracking charm particles; moreover identification of particles and their kinematical properties could be inferred by measuring ionisation and multiple scattering.

Emulsion experiments had become much more sophisticated since their early successes in discovering the pion and the strange particles: in the early 1950's it had been proposed [40] to combine packs of thick metal plates, acting as absorber or target, with thin emulsion layers for tracking. This type of hybrid detector was developed mainly in Japan and successfully used in cosmic ray studies. *"One can say that nuclear emulsion*

([4]) The history of the J/ψ discovery is described in full, including comments of the main actors, in [24].
([5]) The question whether there are more than four quarks was soon raised [33].

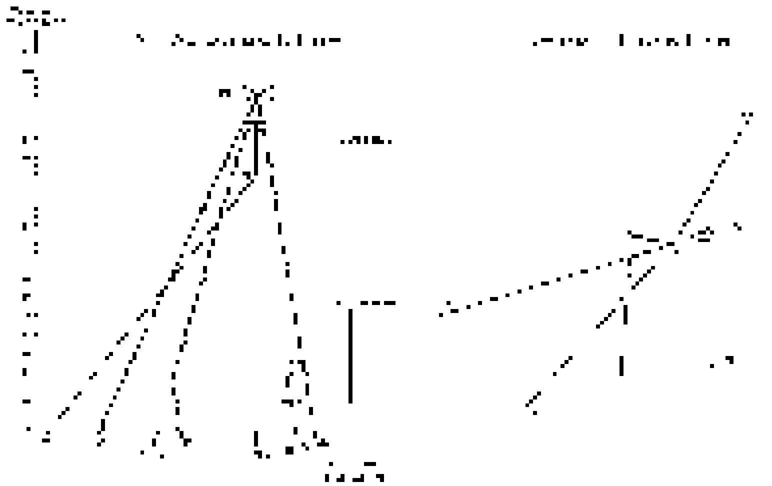

Fig. 3. – First charm candidate event in nuclear emulsions [8]. Figure from Ref. [12].

was the ancestor technique of heavy quark physics" [41]. By 1974 one had already seen lifetime differences between charged and neutral charm hadrons in cosmic ray emulsion data [42], although that was largely ignored outside Japan.

Hybrid detectors, where a forward spectrometer complements emulsions, were then used to study charm at accelerators. Experiments were done at Fermilab from 1975 to 1979 with 205 GeV [43] and 400 GeV [44] proton beams. Those experiments detected the first charm event (and even a charm particle pair) at accelerators. By the end of the seventies, the numbers of charm detected in emulsions at accelerators exceeded the one from cosmic rays. However statistics was still limited to a total of few tens events.

To overcome this limitation, the traditional visual inspection and reconstruction of events in nuclear emulsions was gradually replaced by computer techniques – from semi-automatic scanning machines [45] to fully automatic systems driven by the forward spectrometer tracking information [46] . The new technique saved time in both finding and reconstructing candidate events without introducing a bias in event selection. In 1979,

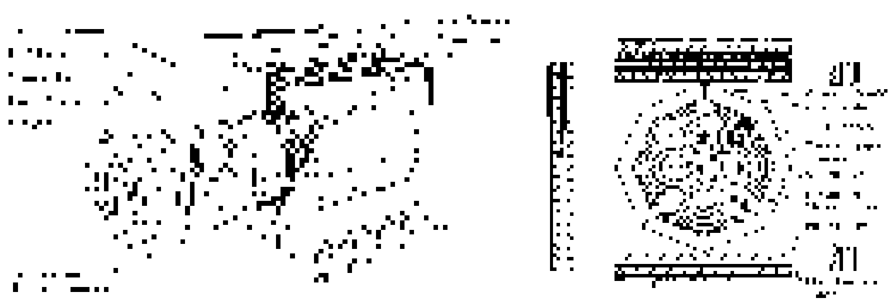

Fig. 4. – MARK2 Detector: exploded and beam view (From [39]).

in a few months five thousand events were analysed in an experiment on negative pion beam of 340 GeV at Cern [47], and the huge (at the time) number of four charm pairs, five charged and three neutral charm particles were detected.

These improved emulsion techniques were applied in full to study charm neutrino production by E531 at Fermilab and by CHORUS at CERN. The E531 collaboration [48] collected more than 120 charm events; among its most notable results was the confirmation of the lifetime differences first seen in cosmic ray data a few years earlier [49].

This new technique contributed also to early beauty searches. WA75 at CERN using a 350 GeV pion beam was the first to detect beauty hadrons[50] in a hadron beam. In a single event both beauty hadrons $\boldsymbol{B}$ and $\bar{\boldsymbol{B}}$ were detected, and their decays into charmed particles observed clearly showing the full sequence of decays from beauty to light quark. WA75 detected about 200 single charm pairs events, among them two peculiar ones with simultaneous production of two pairs of charm.

The CHORUS detector[51] combined a nuclear emulsion target with several electronic devices. By exploiting a fully automated scanning system it localized, reconstructed and analysed several hundred thousand interactions. A sample of about 1000 charm events, a ten-fold increase over E531, was obtained by CHORUS. This big sample should allow the measurement of the, so far never measured, total charmed-particle production inclusive cross-section in antineutrino induced event [52].

The scanning speed achievable with fast parallel processors increases by about one of magnitude every three years. Soon a scanning speed of 20 $\mathrm{cm}^{\mathbf{2}}/\mathrm{s}$ should be possible[53]. These developments assure a continuing presence of emulsion techniques in high energy physics.

Bubble chambers made important contributions as well. Charm decays were seen in the 15 ft bubble chamber at Fermilab [54]. Very rapid cycle bubble chambers coupled with a forward magnetic spectrometer contributed since the early days of charm physics at Fermilab [55] and Cern [56]. LEBC was utilized by NA16 and NA27 searching for charm states at CERN , while SLAC operated the SHF (Slac Hybrid facility). Yet these devices have remained severely limited in the statistics they can generate, due to low repetition rate of 20-40 Hz, the short sensitivity time 200 microseconds, and to the small fiducial volume. Thus they are of mainly historical interest now.

3˙1.4. The silicon revolution. Charm quark physics witnessed in a very distinct fashion the very transition from image to logic[57] which is common to several fields of particle physics. Turning point of such transition was the replacement of emulsions and bubble chambers with electronic imaging devices.

The NA1 experiment at CERN was one of the first experiments that introduced silicon and germanium devices into the field (6). This was soon followed by one of the major breakthroughs in the detector techniques of the last 20 years: the silicon microstrip high-resolution vertex detector .

To measure lifetimes, NA1 used a telescope composed of several silicon detectors (150-300 microns thick) with beryllium sheet targets in between, installed directly in the photon beam Fig.5. The telescope acts as an active target: when an interaction occurs, the silicon device detects the energy released by the recoil system (the nuclei or

(6) The degree to which charm's arrival in the data produced a revolution not merely in our view of fundamental dynamics, but also in detector science can be seen from the fact that experiments converted their objectives *in flight* to new quests. E.g., NA1 at CERN was originally designed to study hadronic fragmentation (as its FRAMM name recalls).

a proton) and by particles emerging from the interaction points. The pattern on the detected energy in the subsequent detectors identifies production and decay locations along the silicon telescope. The recoil fragment or nucleon releases sufficient energy to identify the interaction point even when the emerging particles were neutral. The ideal sequences of energy-deposited steps are shown in Fig.5, for photoproduction of both charged and neutral charmed mesons pairs compared to a typical event configuration. The first determination of the time evolution curve of a charmed particle was obtained by NA1 collaboration with this innovative device proving a lifetime of $\mathbf{9.5^{+3.1}_{-1.9}\,10^{-13}}$ s out of a sample of 98 events [58]. NA1 published data also on the $\boldsymbol{\Lambda_c}$ lifetime and production asymmetry[59].

Finally microstrip vertex detectors were brought to the scene. This new device allowed one to perform tracking of particles trajectories upstream of the forward spectrometer magnetic field, and to reconstruct with precision the primary (production) and the secondary (decay) vertices of short living particles in the events. It moved lifetime determinations to the fully digital state and also opened the field to search and study specific decay channels. Microstrip vertex detectors are composed of several stations, each formed of three microstrip planes typically 200-300 micron thick, with strips running at different orientation. Between the target (passive Cu or Be bulk or active silicon telescope) where the interaction occurs (primary vertex), and the subsequent decay (secondary vertex) there is an empty region were most of the searched decays should happens, whose size must be optimised taking into account expected lifetimes and their relativistic boost. A second series of microstrip detectors is placed at the secondary vertex location and downstream to it. This configuration allows one to reconstruct the sequence of decay vertices, and to link emitted tracks to those reconstructed in the forward spectrometer. The strips typically were 20-100 micron wide, 20-50 micron pitch. Spatial resolutions on the plane perpendicular to the beam of the order of several microns were obtained. The multiple Coulomb scattering limits to 4-5 mm the total thickness allowed. The first examples of this kind of apparatus are ACCMOR[60] on hadron beams and NA14 (Fig.6), E691 (Fig.7) on photon beams.

By the mid-80's fixed target experiments using microstrip vertex detectors had become mature, the technique migrated from CERN to the US, experiments with thousands of channels were built and took data for more than ten years. The two main experiments at fixed target were operated at FNAL: E691 [61] (later running also as E769 and E791) and E687[62], later upgraded to E831-FOCUS. At present the overall largest statistics with more than a million identified charm events has been accumulated by E831-FOCUS, which concluded data taking eight years ago. In the meantime CLEO at Cornell's CESR ring – for a long time the only $\boldsymbol{B}$ factory in the world – passed through several upgrades and developed new methods of analysis. LEP produced a heavy flavour program at the $\boldsymbol{Z^0}$ that had not been foreseen. Finally the second generation $\boldsymbol{B}$ factories at KEK and SLAC arrived on the scene at the end of the millennium. They have obtained charm samples of similar size to FOCUS and will surpass it considerably in the coming years.

The discovery of charm had been largely a US affair, yet CERN experiments made a dramatic entry in the second act with conceptually new detectors and mature measurements.

Semiconductor detector technology migrated from nuclear to high-energy physics experiments where it attained its apogee. It had a truly far reaching impact: (i) The resulting technology that allows tracing lifetimes of about $\mathbf{few} \times \mathbf{10^{-13}}$ s for charm was 'on the shelves' when beauty hadrons were discovered with lifetimes around 1 ps. This was a 'conditio sine qua non' for the success of the $\boldsymbol{B}$ factories. (ii) It is essential for

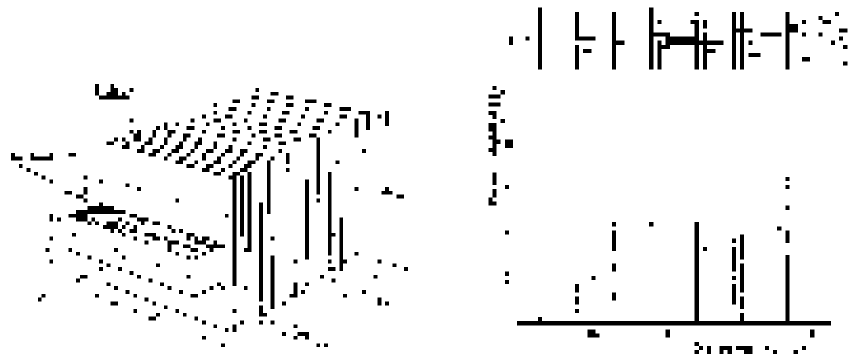

Fig. 5. – Ge-Si active target of CERN NA1 experiment (left). A D^+D^- event and corresponding pulse height pattern in target (right). From Ref. [63].

Fig. 6. – NA14/2 Spectrometer (from Ref. [64]).

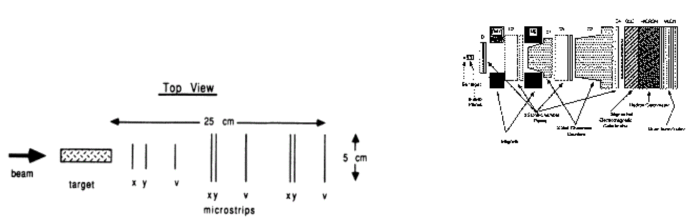

Fig. 7. – E691 Detector: vertex detector region(left); forward spectrometer (right). From ref.[61].

heavy flavour studies at hadronic machines. (iii) The resulting $\boldsymbol{B}$ flavour tagging was essential - and will continue to be so - in finding top production through its decays to beauty hadrons. (iv) It will be an indispensable tool in future Higgs searches.

3˙2. *The past's lessons on the production environment.* – The historical sketch presented above shows that practically the high energy physics' whole pantheon of experimental techniques has contributed to charm physics. We can draw various lessons for the future of heavy flavour physics from the past experiences.

The cleanest environment is provided by $\boldsymbol{e^+e^-}$ annihilation, where threshold machine, $\boldsymbol{B}$ and $\boldsymbol{Z^0}$ factories complement each other. Threshold machines like SPEAR and DORIS in the past, BES in the present and also CLEO-c in the future allow many unique measurements benefiting from low backgrounds and beam-energy constraints. They suffer somewhat from the fact that the two charm hadrons are produced basically at rest thus denying microvertex detectors their power. A $\boldsymbol{Z^0}$ factory, as LEP and SLC have been, on the other hand benefits greatly from the hard charm fragmentation: the high momenta of the charm hadrons and their 'hemispheric' separation allows to harness the full power of microvertex detectors; similar for beauty hadrons. The LEP experiments ALEPH, DELPHI, L3 and OPAL together with SLD have made seminal contributions to our knowledge of heavy flavour physics in the areas of spectroscopy, lifetimes, fragmentation functions, production rates and forward-back asymmetries. The advantage of $\boldsymbol{B}$ factories is their huge statistics with low background level. We are probably not even near the end of the learning curve of how to make best use of it.

Photoproduction experiments have been a crucially important contributor. The charm production rate is about 1/100 of the total rate with a final state that is typically of low multiplicity. A crucial element for their success was the ability to track the finite decay paths of charm hadrons routinely, which has been acquired due to the dedicated R & D efforts described above. Their forte is thus in the areas of time-dependent effects like lifetimes, $\boldsymbol{D^0 - \bar{D}^0}$ oscillations and CP violation there.

The largest cross sections for charm production are of course found in hadroproduction. In high energy *fixed target* experiments one has to deal with a signal to background of a priori about 1/1000 with high multiplicity final states. That this challenge could be overcome again speaks highly of the expertise and dedication of the experiments. At hadron colliders like the TEVATRON the weight of the charm cross section is higher – about 1/500 of the total cross section – yet so is the complexity of the final state. CDF, which previously had surprised the community by its ability to do high-quality beauty physics, is pioneering now the field of charm physics at hadron colliders with its ability to trigger on charm decays. A silicon vertex tracker [65] reconstructs online the track impact parameters, enriching the selected data set of charm events, by triggering on decay vertices. First charm physics results from CDF seem promising[66].

On the novel idea of using cooled antiproton beams impinging on an internal proton jet target was commented in Sect.**3**˙1.1. Such a technique allowed formation studies of charmonium states other than $\mathbf{1^{--}}$, was pioneered at CERN by experiment R704 and further refined by Fermilab experiments E760 and E835 [27, 28, 29].

Charm baryon production at fixed target by means of hyperon beams sees SELEX at Fermilab as probably the last exponent of a technique which is able to provide unique information on production mechanisms (Sect.**5**), as well as on charmed baryon properties.

Studies of charm and beauty production at deep inelastic lepton nucleon scattering as done at HERA primarily act as tools for a better understanding of the nucleon's structure in general and the gluon structure functions in particular.

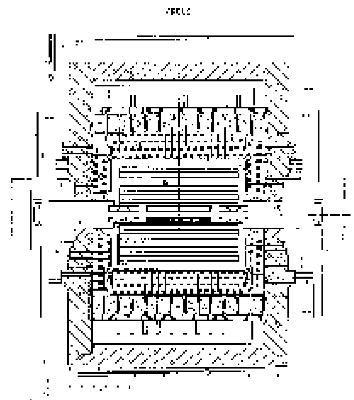

Fig. 8. – ARGUS Detector (elevation view from ref.[68]).

3˙3. *Key detector components*. – The arrival of charm on the market produced a major revolution, not only in physics, but also in detector science. The distinct properties predicted for charm decays (mentioned in Sect.**3**: short but finite lifetime, dominance of kaon decays, relevant branching ratio for semimuonic decays) gave a definite roadmap for the development of new experimental techniques.

Experiments suddenly converted on flight their objectives to the new physics quests, and a big R&D adventure started to conceive new devices able to reach the needed spatial resolutions. This pushed the migration of the semiconductor detector technique from nuclear to high-energy experiments. First ideas relied on silicon active targets, where jumps in silicon pulse height would be a signal for jumps in charged particle multiplicity, i.e., of a charm decay point. Space resolution was limited by the thickness of the silicon targets. Key element for the transition to modern charm lifetime measurements was the silicon microstrip detector. Such a transition could not have been accomplished without the development of DAQ systems able to handle the very large dataflow provided by the huge number of channels in microstrip detectors. The advantage given by the Lorentz boost at fixed-target experiment was immediately realized. It was also realized that, given the statistical essence of the lifetime measurement, a very large data sample was needed to reach high statistical precisions. Porting of silicon microstrips to collider charm experiment is a relatively recent history pioneered by CLEO, and fully embraced by B-factories by the usage of asymmetric beams in order to avail of some Lorentz

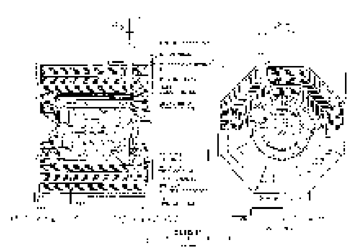

Fig. 9. – CLEO-2 Detector (From ref. [69]).

boost. Silicon pixels were the natural quantum leap from microstrip detectors, providing two-dimensional readout, reduced thickness and therefore less multiple scattering, lower track occupancy and better space resolution, at the cost of a much increased number of readout channels. In the course of R&D for vertex detectors for charm, several good ideas were investigated, such as the use of scintillating fibers as micron-resolution tracking devices[70] which did not last in charm physics but had many applications in HEP, or elsewhere.

Vertex reconstruction for charm decays is intimately linked to the possibility of triggering on it. Charm physics at hadron colliders was born very recently with the success obtained by CDF at Tevatron in exploiting a hadronic trigger based on the online reconstruction of vertex impact parameters. Future experiments such as BTeV at Fermilab (see Sect.**13**˙2.6) plan an aggressive charm physics program based on a first level trigger selecting events with secondary vertices reconstructed in pixel detectors.

Particle identification, namely the rejection of pions and protons against kaons, was immediately recognized as a winner in charm physics. In pioneer e^+e^- experiments this was basically limited to an identification based on $-dE/dx$ measurement with gas tracking devices. Thanks to the favourable geometry, fixed target experiments could make use of threshold Cerenkov counters. Ring-imaging Cerenkov counters only appeared at e^+e^- colliders with CLEO, and have been further improved with B-factories. The unique role of semielectronic and semimuonic decays in understanding the underlying hadron dynamics gave momentum to electron and muon particle identification techniques, with collider experiments traditionally more efficient in identifying the former, and fixed target experiments favoured by the higher muon momentum is deploying muon filters and detectors.

Finally, electromagnetic calorimetry was recognized as a necessity in charm physics by CLEO II, with the operation of a world-class CsI crystal array. Photon and π^0 detection initially provided textboook measurements such as measurements [71] of $BR(D^0 \to \pi^0\pi^0)$ decays to study, when compared to the charged pion modes, isospin amplitudes. Measurements [72] of π^0 decays for D_s^* unveiled isospin-violating processes thus opening the way to exploring the full L=1 excited mesons spectroscopy with neutrals, until the very recent observations by BABAR and CLEO of the enigmatic $D_{sJ}^*(2317)$ states discussed in detail in Sect.**6**. Such a lesson was deeply metabolized by the physics community and translated to B physics and CKM matrix investigations: planned experiments such as BTeV do foresee the use of sophisticated em calorimeters for detection of photons and π^0.

This section cannot be considered complete without mentioning how the invention of WEB in the 1990's by the HEP community soon was devised as a crucial tool for developing online monitoring systems which would actually span borders, oceans and frontiers — the first truly-WEB-based online monitoring system was developed for a charm experiment [73]. As a summary of the last two sections we show in Tab.I and Tab.II features of present and future experiments, reserving a full discussion of future initiatives at the end of this paper.

4. – Theoretical Technologies

The relationship between the world of *hadrons* experiments have to deal with and the world of quarks (and gluons) in which our basic theories are formulated is a highly complex one. Quarks undergo various processes of *hadronization*, namely how they exchange energy with their environment, how they end up asymptotically in hadrons and

TABLE I. – *Charm in today's experiments.* Sample *column shows number of reconstructed events.* $\boldsymbol{\Delta M}$ *is the typical mass resolution,* $\boldsymbol{\Delta t}$ *is the typical proper time resolution.*

	Beam	Sample	$\boldsymbol{\Delta M}$ MeV	$\boldsymbol{\Delta t}$ fs	$\boldsymbol{\sigma_{c\bar{c}}/\sigma_T}$
CLEO	$\boldsymbol{e^+e^-(\Upsilon(4s))}$	$\mathbf{1.5\,10^5}$ D	0.3	200	∼1/5
BABAR	$\boldsymbol{e^+e^-(\Upsilon(4s))}$	$\mathbf{1.5\,10^6}$ D	0.3	200	∼1/5
BELLE	$\boldsymbol{e^+e^-(\Upsilon(4s))}$	$\mathbf{1.5\,10^6}$ D	0.3	200	∼1/5
E791	$\boldsymbol{\pi}$ 500 GeV	$\mathbf{2.5\,10^5}$ D	1	50	1/1000
SELEX	$\boldsymbol{\pi, \Sigma, p}$ 600 GeV	$\mathbf{1.7\,10^3}\ \boldsymbol{\Lambda_c^+}$	1	40	1/1000
FOCUS	γ 200 GeV	$\mathbf{1\,10^6}$ D	1	40	1/100
CHORUS	$\boldsymbol{\nu_\mu}$ 27 GeV	$\mathbf{2\,10^3}$ D			1/20
E835	$\boldsymbol{p\bar{p}}$ <8 GeV	$\mathbf{4\,10^3}\ \boldsymbol{\chi_{c0}}$	2		1/70000
BES	$\boldsymbol{e^+e^-(J/\psi)}$	$\mathbf{6\,10^7}\ \boldsymbol{J/\psi}$	1		∼1
CDF	$\boldsymbol{p\bar{p}}$ 1 TeV	$\mathbf{1.5\,10^6}$ D	2		1/500
HERA Expts.	ep 100 GeV		1		1/100
LEP Expts.	$\boldsymbol{e^+e^-(Z^0)}$	$\mathbf{1\,10^5}$ D	1	100	1/10

specifically in what kinds of hadrons, etc.

Almost all theoretical concepts and tools used in high energy physics are relevant for treating charm physics in particular, albeit often with quite specific features. From the outset it had been realized – or at least conjectured – that hadronization's impact on charm transitions would become more treatable than for ordinary hadrons due to the large charm mass:

- *Producing* charm from a charmless initial state requires an energy concentration that places in into the realm of short distance dynamics, which can be described perturbatively with sufficient accuracy. It is understood here that one considers production well above threshold since complexities associated with the opening of individual channels can invalidate a short-distance treatment, as discussed in Sect.**4**˙11. At such high energies it is expected that (inclusive) hadronic rates can be approximated with rates evaluated at the quark-gluon level, i.e. that *quark-hadron duality* should hold with sufficient accuracy. This topic will be addressed in Sect.**4**˙11.

TABLE II. – *Charm in future experiments.*

	Beam	Lumin. $\boldsymbol{cm^{-2}s^{-1}}$	Cross sect.	$\int \boldsymbol{L}$ in $\mathbf{10^7}$ s	# events $\boldsymbol{c\bar{c}}$ recon'd/y	S/B
BTEV	$\boldsymbol{p\bar{p}}$ 1 TeV	$\mathbf{2\,10^{32}}$	$\mathbf{500}\boldsymbol{\mu}$b $\boldsymbol{c\bar{c}}$	2 fb^{-1}	$\mathbf{10^8}$	fair
LHCB	$\boldsymbol{pp}$ 7 TeV	$\mathbf{2\,10^{32}}$	$\mathbf{1000}\boldsymbol{\mu}$b $\boldsymbol{c\bar{c}}$	2 fb^{-1}	—	—
CLEO-C	$\boldsymbol{\psi(3770)}$	$\mathbf{2\,10^{32}}$	10 nb $\boldsymbol{c\bar{c}}$	2 fb^{-1}	$\mathbf{2\,10^6}$	large
COMPASS	$\boldsymbol{\pi}$Cu FT	$\mathbf{1\,10^{32}}$	10 $\boldsymbol{\mu}$b $\boldsymbol{c\bar{c}}$	1 fb^{-1}	$\mathbf{5\,10^6}$	fair
BABAR	$\boldsymbol{e^+e^-(\Upsilon(4s))}$	$\mathbf{3\,10^{33}}$	1.2 nb $\boldsymbol{b\bar{b}}$	30 fb^{-1}	$\mathbf{4\,10^6}$	large
BELLE	$\boldsymbol{e^+e^-(\Upsilon(4s))}$	$\mathbf{3\,10^{33}}$	1.2 nb $\boldsymbol{b\bar{b}}$	30 fb^{-1}	$\mathbf{4\,10^6}$	large

- To identify charm production experimentally one typically has to deal with charm fragmentation, i.e. the fact that charm quarks give off energy before they hadronize. For asymptotically heavy quarks such fragmentation functions are expected to turn into delta functions [74]. For charm quarks they should already be 'hard' with the final charm hadron retaining a large fraction of the primary charm quark energy. A simple quantum mechanical ansatz yields a single parameter description that describes data quite well [75].

- The lifetimes of weakly decaying charm hadrons were expected to exhibit an approximate 'partonic' pattern with lifetime ratios close to unity, in marked contrast to strange hadrons. We will sketch the reasons for such expectations and explain their shortcomings.

Very significant progress has happened in formalizing these ideas into definite frameworks that allow further refinements.

- Corrections of higher orders in $\boldsymbol{\alpha_S}$ have been computed for cross sections, structure and fragmentation functions.

- Different parameterizations have been explored for the latter.

- Heavy quark expansions have been developed to describe, among other things, weak decays of charm hadrons.

- Considerable efforts have been made to treat charm hadrons on the lattice.

The goal in sketching these tools and some of their intrinsic limitations is to give the reader a better appreciation of the results to be presented later rather than complete descriptions. Those can be found in dedicated reviews we are going to list at the appropriate places.

4·1. *The stalwarts: quark (and bag) models.* – Quark models (actually different classes of them, nonrelativistic as well as relativistic ones) have been developed well before the emergence of charm. They cannot capture all aspects of the quantum world. Their relationship with QCD is actually somewhat tenuous, unlike for the second generation technologies described below. Quark model quantities like quark masses, potential parameters etc. cannot be related reliably to SM quantities defined in a quantum field theory. Varying these model quantities or even comparing predictions from different quark models does not necessarily yield a reliable yardstick for the theoretical uncertainties, and no *systematic* improvement on the error is possible.

Nevertheless considerable mutual benefits arise when quark models are applied to charm physics. Often quark models are the tool of last resort, when tools of choice, like lattice QCD, cannot be applied (yet). They can certainly educate our intuition and help our interpretation of the results obtained from more refined methods. Lastly they can be invoked to at least estimate certain matrix elements arising in heavy quark expansions, QCD sum rules etc.

Quark models on the other hand are trained and improved by the challenges and insights offered by charm physics. Charmonia constitute the most suitable systems for a description based on inter-quark potential. Open charm mesons consisting of a heavy and a light quark represent a more direct analogy to the hydrogen atom than light-flavour hadrons. Charm baryons, in particular those with $\boldsymbol{C = 2}$, offer novel probes for quark dynamics: the two charm quarks move in close orbits around each other like binary stars surrounded by a light quark farther out.

Bag models – in particular their protagonist, the MIT bag model [76] – were very much en vogue in the 1970's and 1980's. The underlying idea actually impresses by its simplicity. One implements the intuitive picture of quarks being free at short distances while permanently confined at long distances in the following way: one describes a hadron at rest as a cavity of fixed shape (typically a spherical one), yet a priori undetermined size; the quark fields are assumed to be free inside the cavity or "bag", while to vanish outside; this is achieved by imposing certain boundary conditions on the quark fields on the interface between the inside and outside of the bag. The resulting wavefunctions are expressed in terms of spherical Bessel functions; they are relativistic and can be used to evaluate matrix elements. Again open charm mesons lend themselves quite readily to a description by a spherical cavity. Bag models have gained a second lease on life in nuclear physics under the names of "cloudy bag models" or "chiral bag models"; clouds of pions and kaons are added to the bag to implement chiral invariance.

4˙1.1. Quarkonium potential. Since QCD dynamics at small distances can be treated perturbatively, one expects the interactions between very heavy flavour quarks to be well approximated by a Coulombic potential. This expectation can actually be proven using NRQCD to be sketched below; the resulting description is an excellent one for top quarks [77] due to their enormous mass and their decay width $\Gamma_t > \Lambda_{QCD}$ [9] providing an infrared cutoff.

The situation is much more involved for charm quarks. Unlike for t quarks, $\bar{c}c$ bound states have to exist. The fact that m_c exceeds ordinary hadronic scales suggests that a potential description might yield a decent approximation for $\bar{c}c$ dynamics as a sequence of resonances with a narrow width, since they possess only Zweig rule (see Sect.**4**˙10.5) violating decays, and with mass splittings small compared to their masses in qualitative analogy with positronium, hence the moniker charmonium. That analogy can be pursued even further: there are s-wave vector and pseudoscalar resonances named ortho- and para-charmonium, respectively, with the former being even narrower than the latter. For while paracharmonia can decay into two gluons, orthocharmonia annihilation has to yield at least three gluons: $\Gamma([\bar{c}c]_{J=1}) \propto \alpha_S^3(m_c)|\psi(0)|^2$ vs. $\Gamma([\bar{c}c]_{J=0}) \propto \alpha_S^2(m_c)|\psi(0)|^2$; $\psi(0)$ denotes the $\bar{c}c$ wavefunction at zero spatial separation, which can be calculated for a given $\bar{c}c$-potential.

For the latter one knows that it is Coulombic at small distances and confining at large ones. The simplest implementation of this scenario is given the ansatz

$$V(r) = \frac{A}{r} + B\,r + V_0. \tag{3}$$

One finds the energy eigenvalues and wavefunctions by solving the resulting Schrödinger equation as a function of the three parameters A, B, V_0, which are then fitted to the data.

4˙2. *Charm Production and fragmentation.* – Producing charm hadrons from a charmless initial state requires an energy deposition of at least $2M_D$ into a small domain (or at least M_D in neutrino production). Such production processes are thus controlled by short distance dynamics – unless one asks for the production of individual species of charm hadrons, considers only a very limited kinematical range or special cases like leading particle effects. It has to be understood also that charm production close to its threshold cannot be described by short distance dynamics since relative momenta between the c and $\bar{c}$ are low and the opening of individual channels can dominate the rate. For a perturbative treatment one has to stay at least a certain amount of energy above threshold, so that the relevant momenta are sufficiently large and a sufficient number of exclusive channels contribute; some averaging or 'smearing' over energy might still be required. This minimal amount of energy above threshold is determined by nonperturbative dynamics. Therefore we refer to it generically as Λ_{NPD}; sometimes we will invoke a more specifically defined energy scale like $\bar{\Lambda}$ denoting the asymptotic mass difference for heavy flavour hadrons and quarks – $\bar{\Lambda} \equiv \lim_{m_Q \to \infty} (M(H_Q) - m_Q)$. On general grounds one guestimates values like $0.5 - 1$ GeV for them (7).

We do not have the theoretical tools to describe reliably charm production close to threshold – a region characterized by resonances and other prominent structures. Yet well above threshold violations of duality will be of no real significance; the practical limitations are then due to uncertainties in the value of m_c and the input structure functions. It is important to keep in mind that m_c has to be defined not merely as a parameter in a quark model, but in a field theoretical sense. Among other things that means that it will be a scale-dependent quantity like the QCD coupling α_S.

Furthermore one cannot automatically use the same value for m_c as extracted from heavy flavour decays, since the impact of nonperturbative dynamics will differ in the two scenarios. The charm quark mass that enters in production and in decay processes is of course related. The tools to identify this relationship are available; however it has not been determined explicitly yet. Similar comments apply to the masses of strange and beauty quarks.

After charm quarks have been produced well above threshold, they move relativistically and as such are the source of gluon radiation: $c \to c + gluons$. Such reactions can be treated perturbatively for which well-defined prescriptions exist based on shower models. This radiation degrades the charm quark energy till its momentum has been lowered to the GeV scale, when nonperturbative dynamics becomes crucial, since the charm quark will hadronize now:

$$Q \to H_Q (= [Q\bar{q}]) + q \tag{4}$$

On very general grounds [74] one expects the fragmentation function for asymptotically heavy quarks to peak at $z \simeq 1$, where $z \equiv p_{H_Q}/p_Q$ denotes the ratio between the momentum of the emerging hadron and of the primary (heavy) quark with a width Λ_{NPD}/m_Q on dimensional grounds. This conjecture has been turned into an explicit ansatz by approximating the amplitude $T(Q \to H_Q + q)$ with the energy denominator $(\Delta E)^{-1}$, where $\Delta E = E_{H_Q} + E_q - E_Q = \sqrt{z^2 P^2 + m_Q^2} + \sqrt{(1-z)^2 P^2 + m_q^2} - \sqrt{P^2 + m_Q^2} \propto 1 - \frac{1}{z} + \frac{\epsilon}{1-z}$ with $\epsilon = m_q^2/m_Q^2$. Hence one arrives at the following expression for the fragmentation function :

$$D(z) \propto \frac{z(1-z)^2}{[(1-z)^2 + \epsilon z]^2} , \tag{5}$$

(7) Strictly speaking they should not be identified with Λ_{QCD} entering in the argument of the running strong coupling – $\alpha_S(Q^2) = 4\pi / \left(\beta_0 \log \frac{Q^2}{\Lambda^2_{QCD}} \right)$ – although they are all related to each other.

which is strongly peaked at $z = 1$.

Naively one expects $z \leq 1$ even though the function in Eq.(5) yields $D(z) \neq 0$ also for $z > 1$. However this might not hold necessarily; i.e., the heavy flavour hadron H_Q might pick up some extra energy from the "environment". In particular in hadronic collisions charm and beauty production is central; the energy of the $Q\bar{Q}$ subsystem is quite small compared to the overall energy of the collision. A very small 'leakage' from the huge amount of energy in the environment into $Q\bar{Q}$ and finally H_Q system can increase the latter's energy – as well as $p_\perp$ – very significantly. Since those primary distributions are steeply falling such energy leakage would 'fake' a larger charm (or beauty) production cross section than is actually the case [78].

4˙3. *Effective field theories (EFTh)*. – Nature exhibits processes evolving at a vast array of different scales. To describe them, we typically need an explicit theory only for the dynamics at 'nearby' scales; this is called the *effective* theory. The impact from a more fundamental underlying theory at *smaller* distance scales is mainly indirect: the fundamental dynamics create and shape certain quantities that appear in the *effective* theory as free input parameters.

This general concept is realized in quantum field theories as well. For illustrative purposes let us consider a theory with two sets of fields Φ_i and ϕ_j with masses M_{Φ_i} and m_{ϕ_j}, respectively, where $\min\{M_{\Phi_i}\} \gg \max\{m_{\phi_j}\}$. Let us also assume that the theory is asymptotically free, i.e. that at ultraviolet scales $\Lambda_{UV} \gg \max\{M_{\Phi_i}\}$ the theory describing the interactions of the 'heavy' and 'light' fields Φ_i and ϕ_j can be treated perturbatively. At lower scales μ s.t. $\min\{M_{\Phi_i}\} > \mu > \max\{m_{\phi_j}\}$ only the light fields ϕ_j remain fully 'dynamical', i.e. can be excited as on-shell fields. The dynamics occurring around such scales μ can then be described by an effective Lagrangian $\mathcal{L}_{eff}$ containing operators built from the light fields only. Yet the heavy fields are not irrelevant: they can contribute to observables as off-shell fields and through quantum corrections. Such effects enter through the c number coefficients, with which the light field operators appear in $\mathcal{L}_{eff}$:

$$(6)\quad \mathcal{L}(\Lambda_{UV} \gg M_\Phi) = \sum_i c_i \mathcal{O}_i(\Phi, \phi) \Rightarrow \mathcal{L}(M_\Phi \gg \mu > m_\phi) \simeq \sum_i \bar{c}_i(\Phi)\overline{\mathcal{O}}_i(\phi)$$

One typically obtains a larger set of operators involving a smaller set of *dynamical* fields.

This factoring is usually referred to as 'integrating out' the heavy fields. We will present two examples explicitly below, namely the effective weak Lagrangian in Sect.**4**˙10.1 and the QCD Lagrangian for static heavy quarks in Sect.**4**˙6.1. The latter example will also illustrate that effective Lagrangian are typically non-renormalizable; this does not pose a problem, though, since they are introduced to tackle low- rather than high-energy dynamics.

4˙4. $1/N_C$ *expansions*. – As described in Sect.**2**˙3.4 there are several reasons why the number of colours N_C has to be three. Yet in the limit of $N_C \to \infty$ QCD's nonperturbative dynamics becomes tractable [79] with the emerging results highly welcome : to leading order in $1/N_C$ only planar diagrams contribute to hadronic scattering amplitudes, and the asymptotic states are mesons and baryons; i.e., confinement can be proven then; also the Zweig rule (also called the OZI rule) holds.

Such expansions are employed as follows to estimate at least the size of nonperturbative contributions: one treats short distance dynamics perturbatively with $N_C = 3$

kept fixed to derive the effective Lagrangian at lower and lower scales, see Sect.**4**˙10.1. Once it has been evolved down to scales, where one wants to evaluate hadronic matrix elements, which are shaped by long distance dynamics, one expands those in powers of $1/N_C$:

$$\langle f|\mathcal{L}_{eff}|i\rangle \propto b_0 + \frac{b_1}{N_C} + \mathcal{O}(1/N_C^2) \tag{7}$$

How this is done, will be exemplified in Sect.**9**˙4. In almost all applications only the leading term b_0 is retained, since the next-to-leading term b_1 is in general beyond theoretical control. In that sense one indeed invokes the $N_C \to \infty$ limit.

While $1/N_C$ expansions offer us novel perspectives onto nonperturbative dynamics, they do *not* enable us to decrease the uncertainties *systematically*, since we have little theoretical control over the nonleading term b_1, let alone even higher order contributions.

4˙5. *Heavy quark symmetry (HQS).* – The nonrelativistic dynamics of a spin 1/2 particle with charge g is described by the Pauli Hamiltonian :

$$\mathcal{H}_{\mathrm{Pauli}} = -gA_0 + \frac{(i\vec{\partial} - g\vec{A})^2}{2m} + \frac{g\vec{\sigma}\cdot\vec{B}}{2m} \tag{8}$$

where A_0 and $\vec{A}$ denote the scalar and vector potential, respectively, and $\vec{B}$ the magnetic field. In the heavy mass limit only the first term survives:

$$\mathcal{H}_{\mathrm{Pauli}} \to -gA_0 \quad \text{as} \quad m \to \infty \; ; \tag{9}$$

i.e., an infinitely heavy ‘electron’ is static: it does not propagate, it interacts only via the Coulomb potential and its spin dynamics have become decoupled. Likewise for an infinitely heavy quark its mass drops out from its dynamics (though not its kinematics of course); it is the source of a static colour Coulomb field independent of the heavy quark spin. This is the statement of heavy quark symmetry of QCD in a nutshell.

There are several immediate consequences for the spectrum of heavy-light systems, namely mesons = $[\bar{Q}q]$ or baryons = $[Qq_1q_2]$:

- In the limit $m_Q \to \infty$ the spin of the heavy quark Q decouples, and the spectroscopy of heavy flavour hadrons can be described in terms of the spin and orbital degrees of freedom of the *light* quarks alone.
- Therefore to leading order one has no hyperfine splitting:

$$M_D \simeq M_{D^*} \; , \; M_B \simeq M_{B^*} \tag{10}$$

- In the baryons $\Lambda_Q = [Qud]$ and $\Xi_Q = [Qsu/d]$ the light diquark system forms a scalar; to leading order in $1/m_Q$ *baryons* accordingly constitute a simpler dynamical system than *mesons*, where the light degrees of freedom carry spin one-half. Among other things this feature reduces the number of *independent* form factors describing semileptonic decays of heavy flavour baryons. We will return to this point in Sect.**8**˙2.

- Some hadronic properties are independent of the mass of the heavy quark flavour. For example, in a transition $Q_1\bar{q} \to Q_2\bar{q} +$ "$W/\gamma/Z^0$" between *two heavy* quarks $Q_{1,2}$ the formfactor, which reflects the response of the cloud of light degrees of freedom, has to be
  - normalized to unity asymptotically for zero-recoil – i.e. when there is no momentum transfer;
  - in general dependent on the velocity $v = p/m_Q$ only.
- There are simple scaling laws about the approach to asymptotia:

$$M_{B^*} - M_B \simeq \frac{m_c}{m_b}\left(M_{D^*} - M_D\right) \tag{11}$$

$$M_B - M_D \simeq m_b - m_c \simeq M_{\Lambda_b} - M_{\Lambda_c} \tag{12}$$

The question how quickly the heavy quark case is approached can be addressed through $1/m_Q$ expansions sketched below. A priori it is not clear to which degree the statements listed above apply to the actual charm hadrons with their marginally heavy mass.

Beyond its intrinsic interest of probing QCD in a novel environment there is also another motivation for studying the spectroscopy of the excitations of charm mesons, namely to enhance our understanding of semileptonic B meson decays and how to extract the CKM parameter $V(cb)$ there. Rigorous sum rules can be derived from QCD that relate basic heavy quark parameters relevant to B decays – like quark masses, hadronic expectation values, the slope of the Isgur-Wise functions – to the observable transition rates for $B \to \ell\nu D^{JLP}$, where the produced charm meson D^{JLP} carries fixed spin J, orbital L and parity P quantum numbers. We will discuss some explicit examples later on.

4˙6. *Heavy quark expansions (HQE).* – With HQS representing an asymptotic scenario, one can ask whether one can evaluate pre-asymptotic effects. The example of the Pauli Hamiltonian already shows that the heavy quark mass constitutes an expansion parameter for describing its dynamics in general and its nonperturbative aspects in particular. There are two variants for the implementation of such expansions, namely for (a) describing the dynamics of heavy quarks purely within QCD and (b) the weak decays of hadrons containing heavy quarks when electroweak forces are included.

4˙6.1. QCD for heavy quarks. One starts by decomposing the QCD Lagrangian at scales larger than m_Q into a part that contains only light degrees of freedom and one that contains the heavy quarks:

$$\mathcal{L}_{QCD}(\mu \gg m_Q) = \mathcal{L}_{light}(\mu \gg m_Q) + \mathcal{L}_{heavy}(\mu \gg m_Q) \tag{13}$$

$$\mathcal{L}_{light}(\mu \gg m_Q) = -\frac{1}{4}G_{\mu\nu}G_{\mu\nu} + \sum_q \bar{q}i(\not{D} - m_q)q \tag{14}$$

$$\mathcal{L}_{heavy}(\mu \gg m_Q) = \sum_Q \bar{Q}i(\not{D} - m_Q)Q \tag{15}$$

with D denoting the covariant derivative. At scales below m_Q – yet still above normal hadronic scales – $\mathcal{L}_{light}$ remains basically the same since its degrees of freedom are still fully dynamical, whereas $\mathcal{L}_{heavy}$ undergoes a fundamental change since the heavy quark cease to be dynamical degrees of freedom:

$$\mathcal{L}^{heavy} = \sum_Q \left[\bar{Q}(i \not{D} - m_Q)Q + \frac{c_G}{2m_Q}\bar{Q}\frac{i}{2}\sigma \cdot GQ + \sum_{q,\Gamma} \frac{d^{(\Gamma)}_{Qq}}{m_Q^2}\bar{Q}\Gamma Q\bar{q}\Gamma q \right] + \mathcal{O}\left(1/m_Q^3\right) \tag{16}$$

where c_G and $d^{(\Gamma)}_{Qq}$ are coefficient functions: the Γ denote the possible Lorentz covariant fermion bilinears and $\sigma \cdot G = \sigma_{\mu\nu} G_{\mu\nu}$ with the gluonic field strength tensor $G_{\mu\nu} = gt^a G^a_{\mu\nu}$. Thus a dimension five operator arises – usually referred to as *chromomagnetic* operator – and various dimension six four-fermion operators. When expressing the heavy quark fields through their static nonrelativistic fields also the so-called kinetic energy operator of dimension five $O_{kin} = \bar{Q}(i\vec{D})^2 Q$ enters. Since it is not Lorentz invariant, it cannot appear in the Lagrangian.

This effective Lagrangian is not renormalizable since it contains operators of dimension higher than four. This is no drawback, though, when treating hadronic spectroscopy.

4·6.2. The Operator Product Expansion (OPE) and weak decays of heavy flavour hadrons. A powerful theoretical tool of wide applicability is provided by the operator product expansion a la Wilson [80]. One can apply it profitably when *inclusive* transitions involving hadrons are driven by short-distance dynamics characterized by a high momentum or energy scale $\sqrt{Q^2}$. 'Classical' examples are provided by deep-inelastic lepton-nucleon scattering with *space-like* Q^2 and by $\sigma(e^+e^- \to had.)$ for *time-like* $Q^2 = s$. Starting in 1984 [81] another application has been developed for decays of a heavy flavour hadron H_Q, where the width for a sufficiently inclusive final state can be expressed as follows [82]:

$$\Gamma(H_Q \to f) = \frac{G_F^2 m_Q^5(\mu)}{192\pi^3}|V_{CKM}|^2 \left[c_3(m_f;\mu)\langle H_Q|\bar{Q}Q|H_Q\rangle|_{(\mu)} + \right.$$

$$\left. c_5(m_f;\mu)\frac{\mu_G^2(H_Q,\mu)}{m_Q^2} + \sum_i c_{6,i}(m_f;\mu)\frac{\langle H_Q|(\bar{Q}\Gamma_i q)(\bar{q}\Gamma_i Q)|H_Q\rangle|_{(\mu)}}{m_Q^3} + \mathcal{O}(1/m_Q^4)\right] \tag{17}$$

with $\mu_G^2(H_Q) \equiv \langle H_Q|\bar{Q}\frac{i}{2}\sigma \cdot GQ|H_Q\rangle$; Γ_i denote the various Lorentz structures for the quark bilinears and V_{CKM} the appropriate combination of CKM parameters. Eq.(17) exhibits the following important features:

- The expansion involves
  - c-number coefficients $c_{...}(m_f;\mu)$ given by short-distance dynamics; they depend on the final state as characterized by quark masses m_f;
  - expectation values of *local* operators controlled by long-distance physics;
  - inverse powers of the heavy quark m_Q scaling with the known dimension of the operator they accompany.
- A central element of Wilson's prescription is to provide a self consistent separation of short-distance and long-distance dynamics implied above. This is achieved by introducing an *auxiliary* energy scale μ demarking the border: short-distance $<$

μ^{-1} < long-distance. The *heavy* degrees of freedom – i.e. with masses exceeding μ – are 'integrated out' meaning that their contributions are lumped into the *coefficients* c_i, which are thus shaped by short-distance dynamics. Only degrees of freedom with masses *below* μ – the 'light' fields – appear as fully dynamical fields in the local *operators*. The one exception from this rule are the heavy quark fields Q; the operators have to be bilinear in $\bar{Q}$ and Q, since the initial state – the decaying hadron H_Q – carries heavy flavour.

- As a matter of principle observables have to be independent of the auxiliary scale μ since nature cannot be sensitive to how we arrange our computational tasks. The μ dependence of the coefficients c_i has therefore to cancel against that of the expectation values due to the operators. In practice, though, the value of μ is not arbitrary, but has to be chosen judiciously for those very tasks: on one hand one would like to choose μ as high as possible to obtain a reliable *perturbative* expression in powers of $\alpha_S(\mu)$; on the other hand one likes to have it as low as possible to evaluate the nonperturbative *expectation values* in powers of μ/m_Q:

$$\Lambda_{QCD} \ll \mu \ll m_Q \tag{18}$$

  For simplicity we will not state the dependence on μ explicitly.

- The expectation values of these local operators are shaped by long-distance dynamics. The nonperturbative effects on the decay width – a *dynamical* quantity – can thus be expressed through expectation values and quark masses. Such *static* quantities are treated more easily; their values can be inferred from symmetry arguments, other observables, QCD sum rules, lattice studies or quark models.

- The same cast of local operators $(\bar{Q}...Q)$ appears whether one deals with nonleptonic, semileptonic or radiative decays of mesons or baryons containing one, two or even three heavy quarks or antiquarks. The weight of the different operators depends on the specifics of the transition though.

- *No* $\mathcal{O}(1/m_Q)$ contribution can arise in the OPE since there is no independent dimension four operator in QCD [8] [83]. A $1/m_Q$ contribution can arise only due to a massive duality violation; this concept will be discussed in Sect.4˙11. Even then it cannot lead to a *systematic* excess or deficit in the predicted rate; for a duality violating contribution has to *oscillate* around the 'true' result as a function of m_Q [84]. Thus one should set a rather high threshold before accepting the need for such a contribution. The absence of such corrections gives rise to the hope that a $1/m_c$ expansion can provide a meaningful description.

- The free quark model or spectator expression emerges asymptotically (for $m_Q \to \infty$) from the $\bar{Q}Q$ operator since $\langle H_Q|\bar{Q}Q|H_Q\rangle = 1 + \mathcal{O}(1/m_Q^2)$, see Eq.(23).

4˙6.3. Heavy Quark Parameters (HQP): Quark masses and expectation values. An internally consistent definition of the heavy quark mass is crucial for $1/m_Q$ expansions

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[8] The operator $\bar{Q}i\not{D}Q$ can be reduced to the leading operator $\bar{Q}Q$ through the equation of motion.

conceptually as well as numerically. While this remark is obvious in hindsight, the theoretical implications were at first not fully appreciated.

In QED one naturally adopts the *pole* mass for the electron, which is defined as the position of the pole in the electron Green function (actually the beginning of the cut, to be more precise). It is gauge invariant, and can be measured, since it represents the mass of an isolated electron. For quarks the situation is qualitatively different due to confinement (except for top quarks since they decay before they can hadronize [9]). Yet computational convenience suggested to use the pole mass for quarks as well: while not measurable per se, it is still gauge invariant and infrared stable *order by order in perturbation theory*. It thus constitutes a useful theoretical construct – as long as one addresses purely perturbative effects. Yet the pole mass is *not* infrared stable in *full* QCD – it exhibits an *irreducible theoretical* uncertainty called a renormalon ambiguity [85] : $\frac{\delta m_Q^{pole}}{m_Q} \sim \mathcal{O}\left(\frac{\bar\Lambda}{m_Q}\right)$. Its origin can be understood intuitively by considering the energy stored in the chromoelectric field $\vec{E}_{Coul}$ in a sphere of radius $R \gg 1/m_Q$ around a static colour source of mass m_Q [82]:

$$\delta\mathcal{E}_{Coul}(R) \propto \int_{1/m_Q \le |\vec{x}| \le R} d^3x \vec{E}^2_{Coul} \propto \text{const.} - \frac{\alpha_S(R)}{\pi}\frac{1}{R} \tag{19}$$

The definition of the pole mass amounts to setting $R \to \infty$; i.e., in evaluating the pole mass one undertakes to integrate the energy density associated with the colour source over *all space* assuming it to have a Coulombic form as inferred from perturbation theory. Yet in the full theory the colour interaction becomes strong at distances approaching $R_0 \sim 1/\Lambda_{QCD}$, and the colour field can no longer be approximated by a $1/R$ field. Thus the long distance or infrared region around and beyond R_0 cannot be included in a meaningful way, and its contribution has to be viewed as an intrinsic uncertainty in the pole mass, which is then estimated as stated above. Using the pole mass in the width $\Gamma \propto m_Q^5$ would generate an uncertainty $\sim\ 5\bar\Lambda/m_Q$ and thus dominate (at least parameterically) the leading nonperturbative contributions of order $1/m_Q^2$ one works so hard to incorporate.

Instead one has to employ a running mass $m_Q(\mu)$ defined at a scale μ that shields it against the strong infrared dynamics. There are two kinds of well defined running masses one can rely on, namely the '$\overline{MS}$' mass $\overline{m}_Q(m_Q)$ ([9]) and the 'kinetic' mass $m_Q^{kin}(\mu)$. The former represents a quantity of computational convenience – in particular when calculating perturbative contributions in dimensional regularization – rather than one that can be measured directly. For $\mu \geq m_Q$ it basically coincides with the running mass in the Lagrangian and is best normalized at $\mu \sim m_Q$. However it diverges logarithmically for $\mu \to 0$. It is quite appropriate for describing heavy flavour *production* like in $Z^0 \to Q\bar{Q}$, but not for treating H_Q *decays*, since there the dynamics are characterized by scales *below* m_Q.

The kinetic mass (so-called since it enters in the kinetic energy of the heavy quark) on the other hand is regular in the infrared regime with [86, 85, 87]

$$\frac{dm_Q^{kin}(\mu)}{d\mu} = -\frac{16}{9}\frac{\alpha_S}{\pi} - \frac{4}{3}\frac{\alpha_S}{\pi}\frac{\mu}{m_Q} + \mathcal{O}(\alpha_S^2) \tag{20}$$

([9]) $\overline{MS}$ stands for 'modified minimal subtraction scheme'.

and is well suited for treating decay processes. It can be shown that for b quarks $\mu \sim 1$ GeV is an appropriate scale for these purposes whereas $\mu \sim m_b$ leads to higher order perturbative corrections that are artificially large [87]. For charm quarks on the other hand this distinction disappears since m_c exceeds the 1 GeV scale by a moderate amount only.

There are four classes of observables from which one can infer the value of m_c, listed in descending order of the achieved theoretical reliability: (i) the spectroscopy of hadrons with hidden or open charm; (ii) the *shape* of spectra in semileptonic B decays driven by $b \to c$; (iii) charm production in deep inelastic lepton nucleon scattering or e^+e^- annihilation; (iv) the weak decays of charm hadrons.

Another approach to the value of m_c is provided by relating the difference $m_b - m_c$ to the spin averaged masses of charm and beauty mesons:

$$m_b - m_c = \langle M_B \rangle - \langle M_D \rangle + \frac{\mu_\pi^2}{2}\left(\frac{1}{m_c} - \frac{1}{m_b}\right) + \mathcal{O}(1/m_{c,b}^2) \tag{21}$$

with $\langle M_{B[D]} \rangle \equiv M_{B[D]}/4 + 3M^*_{B[D]}/4$ and

$$\mu_\pi^2 \equiv \langle H_Q | \bar{Q}(i\vec{D})^2 Q | H_Q \rangle / 2M_{H_Q} \; ; \tag{22}$$

$\vec{D}$ denotes the covariant derivative and $i\vec{D}$ thus the (generalized) momentum; $\mu_\pi^2/2m_Q$ therefore represents the average kinetic energy of the quark Q inside the hadron H_Q. This relation is free of the renormalon ambiguity mentioned above. On the down side it represents an expansion in $1/m_c$, which is of uncertain numerical reliability. Furthermore in order $1/m_{c,b}^3$ *non*local operators appear. Later we will give numerical values for m_c.

The expectation value of the leading operator $\bar{Q}Q$ can be related to the flavour quantum number of the hadron H_Q and operators of dimension five and higher:

$$\langle H_Q | \bar{Q}Q | H_Q \rangle / 2M_{H_Q} = 1 - \frac{1}{2}\frac{\mu_\pi^2}{m_Q^2} + \frac{1}{2}\frac{\mu_G^2}{m_Q^2} + \mathcal{O}(1/m_Q^3) \tag{23}$$

The chromomagnetic moment $\langle H_Q | \bar{Q}\frac{i}{2}\sigma \cdot GQ | H_Q \rangle$ is known with sufficient accuracy for the present purposes from the hyperfine splittings in the masses of vector and pseudoscalar mesons V_Q and P_Q, respectively:

$$\mu_G^2(H_Q, 1 \text{ GeV}) \equiv \frac{\langle H_Q | \bar{Q}\frac{i}{2}\sigma \cdot GQ | H_Q \rangle}{2M_{H_Q}} \simeq \frac{3}{4}(M(V_Q)^2 - M(P_Q)^2) \tag{24}$$

The size of the charm chromomagnetic moment is similar to what is found for beauty hadrons

$$\mu_G^2(D, 1 \text{ GeV}) \simeq 0.41 \text{ (GeV)}^2 \;\; vs. \;\; \mu_G^2(B, 1 \text{ GeV}) \simeq 0.37 \text{ (GeV)}^2 \tag{25}$$

and thus in line what one expects for a heavy quark ([10]).

([10])One should keep in mind though that for reasons we do not understand the hyperfine splittings are quite universal: $M_\rho^2 - M_\pi^2 \sim 0.43$ (GeV)^2, $M_{K^*}^2 - M_K^2 \sim 0.41$ (GeV)^2.

For the $\Lambda_Q \equiv [Qdu]$ and $\Xi_Q \equiv [Qsq]$ baryons we have

$$\mu_G^2(\Lambda_Q, 1 \text{ GeV}) \simeq 0 \simeq \mu_G^2(\Xi_Q, 1 \text{ GeV}) \; , \tag{26}$$

since the light diquark system $q_1 q_2$ in $[Qq_1q_2]$ carries spin 0 in Λ_c and Ξ_c. Yet in the $\Omega_Q \equiv [Qss]$ baryon the ss diquark carries spin one and we have

$$\mu_G^2(\Omega_c, 1 \text{ GeV}) \simeq \frac{2}{3} \left(M^2(\Omega_c^{(3/2)}) - M^2(\Omega_c) \right) . \tag{27}$$

The kinetic energy expectation values are less precisely known beyond the inequality [88, 82]

$$\mu_\pi^2(H_Q) \geq \mu_G^2(H_Q) \tag{28}$$

derived in QCD. To the degree that charm quarks fill the role of heavy quarks one expects very similar values as for B mesons; i.e.

$$\mu_\pi^2(D, 1 \text{ GeV}) \sim 0.45 \pm 0.10 \text{ (GeV)}^2 \; . \tag{29}$$

The largest uncertainties enter in the expectation values for the dimension-six four-fermion operators in order $1/m_Q^3$. In general there are two classes of expectation values, namely for $SU(3)_C$ singlet and octet quark bilinears $\langle H_Q|(\bar{Q}_L\gamma_\mu q_L)(\bar{q}_L\gamma_\mu Q)|H_Q\rangle$ and $\langle H_Q|(\bar{Q}_L\gamma_\mu\lambda_i q_L)(\bar{q}_L\gamma_\mu\lambda_i Q)|H_Q\rangle$, respectively:

$$\langle H_Q|(\bar{Q}_L\gamma_\mu q_L)(\bar{q}_L\gamma_\mu Q)|H_Q\rangle = \frac{1}{4} f_{H_Q}^2 M_{H_Q}^2 B_{H_Q} \tag{30}$$

$$\langle H_Q|(\bar{Q}_L\gamma_\mu\lambda_i q_L)(\bar{q}_L\gamma_\mu\lambda_i Q)|H_Q\rangle = f_{H_Q}^2 M_{H_Q}^2 \epsilon_{H_Q} \tag{31}$$

A natural way to estimate the mesonic expectation values is to assume *factorization* or *vacuum saturation* at a low scale of around 1 GeV; i.e. $B_{H_Q} = 1$ and $\epsilon_{H_Q} = 0$. Such an approximation should be sufficient considering we cannot, as already mentioned, count on more than semi-quantitative predictions about charm decays:

$$\langle P_Q|(\bar{c}_L\gamma_\mu q_L)|0\rangle \cdot \langle 0|(\bar{q}_L\gamma_\mu c)|P_Q\rangle = \frac{1}{4} f_{P_Q}^2 M_{P_Q}^2 \tag{32}$$

$$\langle P_Q|(\bar{Q}_L\gamma_\mu\lambda_i q_L)|0\rangle\langle 0|(\bar{q}_L\gamma_\mu\lambda_i Q)|P_Q\rangle = 0 \tag{33}$$

with the last equation reflecting invariance under colour $SU(3)_C$. One should note that factorizable contributions at a low scale $\sim$ 1 GeV will be partially nonfactorizable at the high scale m_Q. These expectation values are then controlled by the decay constants.

For *baryons* there is no concept of factorization for estimating or at least calibrating the expectation values of four-fermion operators, and we have to rely on quark model results.

Since the moments μ_π^2 and μ_G^2 represent long-distance contributions of order $1/m_Q^2$, one can use their values to estimate the scale characterising nonperturbative dynamics as

$$\Lambda_{NPD} \sim \sqrt{\mu_\pi^2} \sim 700 \text{ MeV} \tag{34}$$

Later we will see that this scale agrees with what one infers from $\bar{\Lambda} \equiv \lim_{m_Q \to \infty} (M(H_Q) - m_Q)$.

These considerations lead to a first resume:

- There is little 'plausible deniability' if a description based on HQE fails for B decays: since m_b is a multiple of Λ_{NPD} given in Eq.(34), Furthermore for the scale μ separating short and long distance dynamics in the OPE one can adopt $\mu \sim 1$ GeV, which satisfies the computational 'Scylla and Charybdis' requirements stated in Eq.(18) for m_b.
- For charm decays, on the other hand, the situation is much more iffy on both counts: with the expansion parameter $\sim \Lambda_{NPD}/m_c$ being only moderately smaller than unity, higher order nonperturbative corrections decrease only slowly, if at all. Furthermore the computational requirement of Eq.(18) is hardly satisfied. The one saving grace might be provided by the absence of an $\mathcal{O}(1/m_c)$ contribution noted above. Finally one expects limitations to quark-hadron duality to be characterized by a factor $e^{-\Lambda_{NPD}/m_c}$ with Λ_{NPD} reflecting the onset of nonperturbative dynamics. It is obviously of essential importance then if this scale is indeed about 700 MeV or 1 GeV (or even higher), which would be bad news.
- For these reasons one cannot count on more than a semi-quantitative description and going beyond $\mathcal{O}(1/m_c^3)$ would then seem pretentious. More generally, it is not clear to which degree charm quarks act dynamically as heavy quarks with respect to QCD. It is unlikely that the answer to the question "Is charm heavy ?" will be a universal 'yes' or universal 'no'. The answer will probably depend on the type of transition one is considering. Yet this uncertainty should not be seen as necessarily evil. For charm transitions allow us to probe the onset of the regime where duality provides a useful concept.

We will adopt as working hypothesis that charm quarks are sufficiently heavy so that HQE can provide a semi quantitative description. We treat it as a learning exercise in the sense that we fully expect the HQE description to fail in some cases. We will apply it to fully inclusive observables like weak lifetimes and integrated semileptonic widths of mesons and baryons. Counting on HQE to describe energy *distributions* in inclusive semileptonic decays is presumably not realistic since the averaging or 'smearing' over the lepton energies etc. required in particular near the end points is such that it amounts to a large fraction of the kinematical range anyway. Furthermore there is no justification for treating strange quarks in the final state of semileptonic decays as heavy.

4·7. *HQET*. – Heavy Quark Effective Theory (HQET) is another implementation of HQS, where one calculates pre-asymptotic contributions as an expansion in $1/m_Q$ [89]. While the core applications of HQET used to be hadronic spectroscopy and the evaluation of form factors for exclusive semileptonic decays of heavy flavour hadrons, the name HQET has been applied to more and more types of observables like lifetimes with varying degrees of justification. Yet we will address here only how HQET deals with spectroscopy and hadronic form factors.

The heavy flavour part of the QCD Lagrangian is expressed with the help of non-relativistic spinor fields $\Phi_Q(x)$ [90]:

$$\mathcal{L}_{HQET} = \sum_Q \left\{ -m_Q \Phi_Q^\dagger \Phi_Q + \Phi_Q^\dagger i D_0 \Phi_Q - \frac{1}{2m_Q} \Phi_Q^\dagger (i\vec{\sigma} \cdot \vec{D})^2 \Phi_Q - \right.$$

$$\frac{1}{8m_Q^2}\Phi_Q^\dagger\left[-\vec{D}\cdot\vec{E}+\vec{\sigma}\cdot(\vec{E}\times\vec{\pi}-\vec{\pi}\cdot\vec{E}\right]\Phi_Q\Bigg\}+\mathcal{O}(1/m_Q^3) \tag{35}$$

where

$$\vec{\pi}\equiv i\vec{D}=\vec{p}-\vec{A}\ ,\ (\vec{\sigma}\cdot\vec{\pi})^2=\vec{\pi}^2+\vec{\sigma}\cdot\vec{B} \tag{36}$$

with $\vec{D}$, $\vec{A}$, $\vec{B}$ and $\vec{E}$ denoting the covariant derivative, the gluon vector potential, the colour magnetic and electric fields, respectively.

On the other hand forces outside QCD – namely the electroweak ones – are given for the full relativistic fields Q. To obtain the relation between the Φ_Q and Q fields one first factors off the time dependence associated with m_Q, which makes up the lion share of Q's energy: $Q(x)=e^{-im_Qt}\hat{Q}(x)$. This can be written covariantly in terms of the four-velocity v_μ:

$$Q(x)=e^{-im_Qx\cdot v}\hat{Q}(x) \tag{37}$$

Yet a consistent separation of the 'large' and 'small' components of the Dirac spinor $\hat{Q}$ cannot be achieved by simply using $h(x)=\frac{1+\gamma_0}{2}\hat{Q}(x)$. A Foldy-Wouthuysen transformation [91] has to be applied yielding [92, 90]:

$$\Phi_Q(x)=\left(1+\frac{\vec{\sigma}\cdot\vec{\pi}}{8m_Q^2}+...\right)\frac{1+\gamma_0}{2}Q(x) \tag{38}$$

There is another complication – both conceptual and technical – in the way HQET is usually defined, namely without introducing an auxiliary scale $\mu>0$ to separate self-consistently heavy and light degrees of freedom as discussed in Sect.4˙6.2. With proper care this problem can be cured, though [90].

HQET has actually become an important tool for inferring lessons on the dynamics of heavy flavour hadrons from lattice QCD results.

4˙7.1. Basics of the spectroscopy. Like for any hadron the mass of a heavy flavour hadron H_Q is given by the properly normalized expectation value of the QCD Hamiltonian; the only difference, which actually amounts to a considerable simplification, is that the latter can be expanded in powers of $1/m_Q$:

$$M_{H_Q}=\langle H_Q|\mathcal{H}|H_Q\rangle$$

$$\mathcal{H}=m_Q+\mathcal{H}_Q+\mathcal{H}_{\rm light}\ ,\ \mathcal{H}_Q=\mathcal{H}_0+\frac{1}{m_Q}\mathcal{H}_1+\frac{1}{m_Q^2}\mathcal{H}_2+... \tag{39}$$

where $\mathcal{H}_{\rm light}$ contains the dynamics for the light degrees of freedom and

$$\mathcal{H}_0=-A_0\ ,\ \mathcal{H}_1=\frac{1}{2}\left(\vec{\pi}^2+\vec{\sigma}\cdot\vec{B}\right)\ ,\ \mathcal{H}_2=\frac{1}{8}\left[-\vec{D}\cdot\vec{E}+\vec{\sigma}\cdot(\vec{E}\times\vec{\pi}-\vec{\pi}\times\vec{E})\right] \tag{40}$$

with the first and second term in $\mathcal{H}_2$ being the Darwin and LS term, respectively, familiar from atomic physics. Hence

$$M_{H_Q} = m_Q + \bar{\Lambda} + \frac{(\mu_\pi^2 - \mu_G^2)_{H_Q}}{2m_Q} + ... \tag{41}$$

with μ_π^2 and μ_G^2 defined in Eqs.(22,24). Eq.(41) has an obvious intuitive interpretation: the mass of H_Q is given by the heavy quark mass m_Q, the 'binding energy' $\bar{\Lambda}$, the average kinetic energy μ_π^2 of the heavy quark Q inside H_Q and its chromomagnetic moment μ_G^2. Since the latter term, which is spin dependent, vanishes in the limit $m_Q \to \infty$, the spin of a heavy flavour hadron can be labeled by the total spin J as well as the spin of the light degrees of freedom j. S-wave pseudoscalar and vector mesons thus form a pair of $j = 1/2$ ground states that in the heavy quark limit are degenerate. Baryons Λ_Q and Ξ_Q carry $j = 0$ and thus represent actually a simpler state than the mesons.

With these expressions one can derive and extend to higher orders the expression for $m_b - m_c$ already stated in Eq.(21):

$$m_b - m_c = \langle M_B \rangle - \langle M_D \rangle + \frac{\mu_\pi^2}{2}\left(\frac{1}{m_c} - \frac{1}{m_b}\right) + \frac{\rho_D^3 - \bar{\rho}^3}{4}\left(\frac{1}{m_c^2} - \frac{1}{m_b^2}\right) + \mathcal{O}(1/m_{c,b}^3) \tag{42}$$

where ρ_D^3 denotes the Darwin term and $\bar{\rho}^3$ the sum of two positive *non*local correlators.

4˙7.2. Semileptonic form factors. Consider a hadron H_Q, where the heavy quark Q is surrounded by – in a terminology coined by Nathan Isgur – the 'brown muck' of the light degrees of freedom in analogy to the situation in an atom. When Q decays weakly into a lighter, yet still heavy quark Q' plus a $l\nu$ pair with invariant mass $\sqrt{q^2}$, the surrounding cloud of light degrees of freedom will not feel this change in its center instantaneously – the hadronization process requires time to adjust. Heavy quark symmetry has two main consequences here, one concerning the normalization of the hadronic formfactors and one their q^2 dependence.

- In the infinite mass limit $m_Q > m_{Q'} \to \infty$ the rate for an exclusive semileptonic transition $H_Q \to H_{Q'} \ell\nu$ at *zero recoil* for the final state hadron $H_{Q'}$ will depend neither on $m_{Q'}$ nor on the heavy quark spin as can be inferred from the Pauli Hamiltonian given in Eq.(8) [81]. I.e., the form factor for $H_Q \to H_{Q'}$ at zero recoil is asymptotically (ignoring also perturbative gluon corrections) unity in 'heavy-to-heavy' transitions for pseudoscalar and vector hadrons $H_{Q^{(\prime)}}$.
- The q^2 of the lepton pair can be expressed through the four-momenta p and p' of H_Q and $H_{Q'}$, respectively, and their four-velocities v and v':

$$q^2 = 2M_{H_Q}^2\left(1 - \frac{p \cdot p'}{M_{H_Q}^2}\right) = 2M_{H_Q}^2(1 - v \cdot v') \tag{43}$$

For both H_Q and $H_{Q'}$ being pseudoscalars one can write

$$\langle H_{Q'}(v')|\bar{Q}'\gamma_\mu Q|H_Q(v)\rangle =$$

$$\left(\frac{M(H_Q)+M(H_{Q'})}{2\sqrt{M(H_Q)M(H_{Q'})}}(p+p')_\mu - \frac{M(H_Q)-M(H_{Q'})}{2\sqrt{M(H_Q)M(H_{Q'})}}(p-p')_\mu\right)\xi(v\cdot v') \ . \tag{44}$$

I.e., there is a single independent form factor $\xi(v \cdot v')$, which is 'universal' in the double sense that it is independent of the heavy quark masses and that it also controls the q^2 dependence, when $H_{Q'}$ is a vector meson; it is usually referred to as the 'Isgur-Wise' function .

At finite values of $m_{Q^{(\prime)}}$ there are corrections to both these features.

Such a scenario is realized with reasonable accuracy for $B \to D^{(*)}\ell\nu$ channels. On the other hand the charm decays $c \to s$ as well as $c \to d$ are of the type '(moderately) heavy to light'. Even then heavy quark symmetry allows to relate the form factors in, say, $D \to \ell\nu\pi$, $D \to \ell\nu\rho$ etc. to those for $B \to \ell\nu\pi$, $B \to \ell\nu\rho$ etc. at the same values of $v \cdot v'$. Yet this relation is not overly useful quantitatively due to the potentially large $1/m_c$ corrections.

4·8. *NRQCD.* – Heavy quark bound states like $[c\bar{c}]$, $B_c = [b\bar{c}]$, $[b\bar{b}]$ etc. are nonrelativistic bound states in the heavy-quark limit. Pre-asymptotic corrections to this limit can conveniently be calculated employing another effective theory, namely nonrelativistic QCD (NRQCD). The local operators that appear in NRQCD are formally very similar as in HQET, Eqs(16. Yet at the same time there is a basic difference in the dynamical stage for $[Q\bar{Q}]$ and $[Q\bar{q}]$ systems: the light antiquark $\bar{q}$ in the latter has to be treated fully relativistically. Formally the same operators can thus appear in different orders in the two schemes. Technically it is easily understood how this comes about: since E_Q/m_Q and p_Q^2/m_Q^2 are of the same order in a nonrelativistic expansion, the primary expansion parameter in NRQCD is the heavy quark velocity $v_Q = p_Q/m_Q$ rather than $1/m_Q$. Among other things this implies that the average heavy quark kinetic energy $\langle H_Q|\bar{Q}|\vec{D}|^2Q|H_Q\rangle/2m_Q$, which enters as a leading order pre-asymptotic *correction* in HQET appears already as part of the asymptotic contribution in NRQCD. Similar to the situation with lattice QCD, see below, there are alternative formulations of NRQCD.

One of the main applications of NRQCD is describing the production of quarkonia in different reactions. The basic picture is the following: one invokes *short*-distance dynamics to produce a $\bar{Q}Q$ pair without restriction on the latter's spin, angular and colour quantum numbers from two initial partons i and j. This $\bar{Q}Q$ pair is then assumed to evolve into the final state quarkonium H – a process involving long-distance dynamics. The analysis is thus based on three main elements [93, 94]:

- One makes a factorization ansatz for the (differential) cross section for producing a quarkonium h from partons i and j:

$$d\sigma = \sum_n d\hat{\sigma}(ij \to \bar{Q}Q(n) + X)\langle O^H(n)\rangle \tag{45}$$

- The quantities $d\hat{\sigma}(ij \to \bar{Q}Q(n) + X)$ are calculated perturbatively and convoluted with parton distribution functions, when necessary.

- The long-distance matrix elements $\langle O^H(n)\rangle$, which encode the hadronization of $\bar{Q}Q(n)$ into H are assumed to be universal, i.e. irrespective of the subprocess $ij \to \bar{Q}Q(n) + X$. On fairly general grounds one can infer how they scale with

> the heavy quark velocity v. One should note that both colour singlet and octet configurations are included. It thus goes well beyond models assuming dominance by colour singlet configurations (or colour octet ones for that matter). One can extract these matrix elements most cleanly from quarkonia production at LEP and apply them to Tevatron or HERA data. While data from LEP have very limited statistics, the predictions for rates at Tevatron and HERA depend sensitively on the parton distribution functions adopted. Some quantitative information on them can also be inferred from quark models and lattice QCD.

The basic philosophy is similar to what was described above for the OPE treatment of charm decays, and the factorization ansatz of Eq.(45) is quite reasonable. However it has not (yet) been proven in a rigorous fashion. One might also be concerned about treating the matrix elements $\langle O^H(n)\rangle$ as universal quantities (^{11}).

Looking beyond these general caveats one expects this formulism to apply to sufficiently heavy quarkonia, like the Υ. Whether it can be applied already for charmonia is another question of course, for which we do not know the answer a priori. As it is with applying HQE to charm decays we should use NRQCD as a tool for *learning* about nonperturbative dynamics and incorporating such lessons rather than ruling out models.

4˙9. *Lattice QCD.* – Monte Carlo simulations of QCD on the lattice or lattice QCD for short provides a very different framework to deal with QCD's complementary features of asymtotic freedom in the ultraviolet and infrared slavery. The four-dimensional space-time continuum is replaced by a discrete lattice with spacing a between lattice sites. This is (usually) viewed not as representing physical reality, but providing the mathematical means to deal with long-distance dynamics through an expansion in the *inverse* coupling. Distances $\sim a$ and smaller obviously cannot be treated in this way. This can be expressed by saying that the finite spacing introduces an ultraviolet cut-off $\sim \pi/a$ for the lattice version of QCD. The short distance dynamics is treated by perturbative QCD and considerable care has to be applied in matching the two theories at a distance scales $\sim a$. One uses the technique of effective field theory sketched in Sect.**4**˙3 to incorporate short-distance dynamics cut off by the finite lattice spacing; the discretization effects are described through an expansion in powers of a:

$$\mathcal{L}_{eff} = \mathcal{L}_{QCD} + a\mathcal{L}_1 + a^2\mathcal{L}_2 + ... \tag{46}$$

With the $\mathcal{L}_i$ containing operators of dimension higher than four, they are nonrenormalizable; this poses no problem since they are constructed to describe long-distance dynamics.

There are actually two measures for the quality of the lattice for our purposes: (i) To get as close as possible to the continuum case one would like to have a as small as possible. (ii) At the same time one wants to have a sufficient number of lattice sites in each dimension to be not overly sensitive to finite size effects. I.e., effectively one has put the particles inside a box to study the response to the forces they experience; yet one does not want having them bounce off the walls of the box too frequently since that is an artifact of the framework.

(^{11})This latter concern could be overcome by including some nonperturbative corrections in the subprocess $ij \to QQ(n) + X$.

Obviously there are practical limitations in the available computing power to achieve these desirable goals. Yet as a condition sine qua non for treating light degrees of freedom one requires am_q, $a\Lambda_{NPD} \ll 1$ for the expansion of Eq.(46) to have practical value.

There are actually a number of different implementations of lattice QCD; they differ mainly in three areas [95]:

- Different expressions for the action defined on the lattice will merge into the same QCD action in the continuum limit. The lattice action can be optimized or 'improved' for example by eliminating $\mathcal{L}_1$; for in that case the continuum case $a \to 0$ is approached like a^2, i.e. much faster.
- Putting fermions on the lattice creates problems between the 'Scylla' of 'fermion doubling' and the 'Charybdis' of vitiating chiral invariance. For very general theorems tell us that in four dimensions chiral invariance is either violated for $a \neq 0$ or maintained at the price of getting too many fermions.
- For heavy quarks one needs actually $am_Q \ll 1$. However with presently available computing power we can achieve merely $am_b \sim 1-2$ and am_c about a third of it. It seems unlikely that in the near future one can achieve $am_b \ll 1$. Several strategies have been suggested to overcome this limitation, namely relying on the static approximation, lattice NRQCD, matching up with HQET and/or extrapolating from m_c up to m_b.

  This is actually another example, where charm hadrons and their decays can provide us with an important bridge on the road towards a deeper understanding of the dynamics of beauty hadrons.
- Including light quarks as fully dynamical degrees of freedom that can be produced and annihilated slows down lattice computations tremendously. This Gordian knot has been treated mostly in the tradition of Alexander the Great, i.e. by 'cutting' or ignoring it. This is called the quenched approximation. The first partially unquenched studies have been presented recently, where two different flavours of light quarks have been fully included in the Monte Carlo simulations.

From the start the primary goal of lattice QCD has been to provide a framework for dealing *quantitatively* with nonperturbative dynamics in all its aspects and in ways that are genuinely based on the first principles of QCD and where the uncertainties can be reduced in a *systematic* way. Indeed no other method has surfaced which can lay claim to a similarly 'universal' validity. On the other hand there are other theoretical technologies that provide a 'first principles' treatment of nonperturbative dynamics, albeit in a more restricted domain; examples are chiral invariance and HQE. Those most definitely benefit from input lattice QCD produces, as described above. Yet lattice QCD benefits also from them, which serve not only as a cross check, but can also provide valuable insights for interpreting findings by lattice QCD.

There are some observables where there is no plausible deniability if lattice QCD failed to reproduce them. Matrix elements involving at most a single hadron each in the initial and final state are in that category. The best developed case history is provided by the decay constants. Studies show an enhancement by 8% in the values for f_{D_s} and f_D when going from quenched to partially unquenched (with $N_f = 2$) while not affecting the ratio f_{D_s}/f_D [96]:

$$f_D(N_f = 0) = 203 \pm 14\,\mathrm{MeV}, \implies f_D(N_f = 2) = 226 \pm 15\,\mathrm{MeV} \tag{47}$$

(48) $f_{D_s}(N_f = 0) = 230 \pm 14\,\mathrm{MeV} \implies f_{D_s}(N_f = 2) = 250 \pm 30\,\mathrm{MeV}$
(49) $(f_{D_s}/f_D)(N_f = 0) = 1.12 \pm 0.02 \implies (f_{D_s}/f_D)(N_f = 2) = 1.12 \pm 0.04$

which can be compared with the experimental findings from $D \to \mu\nu$ as explained later

$$f_D|_{exp} \sim 200 \div 300 \text{ MeV} . \tag{50}$$

Very recently a short paper appeared [97] with the ambitious title "High-Precision Lattice QCD Confronts Experiment" stating that "... realistic simulations are possible now, with all three flavors of light quark" due to a breakthrough in the treatment of light quarks. The authors point out that the treatment in particular of heavy quark physics will benefit greatly.

A note of caution seems appropriate (and it is also sounded by the authors of Ref.[97]. Before a difference in a measured and predicted rate – with the latter based solely on lattice QCD – can be taken as conclusive evidence for the intervention of New Physics, lattice QCD has to be subjected to a whole battery of tests through different types of observables. Charm physics – and this is one of the recurring themes of this review – provides ample opportunity for such a comprehensive program as described later. As an extra bonus, one can, at least in principle, approach the charm scales in both direction, namely from below by using finer and finer lattices and from above by extrapolating from the limit of *static* quarks, for which b quarks provide a good approximation.

4·10. *Special tools*. – In the preceding Subsections we have described theoretical technologies that are most relevant for dealing with heavy flavour hadrons, yet at the same time apply to many other areas as well. Now we sketch some tools with a more limited range of application or more special nature, namely the short distance renormalization of the weak Lagrangian, QCD sum rules, dispersion relations and the concept of final state interactions.

4·10.1. Effective weak Lagrangian. The weak Lagrangian responsible for Cabibbo allowed nonleptonic charm decays is given by a single term at scales just below M_W:

$$\mathcal{L}_W^{\Delta C=1}(\mu < M_W) = (4G_F\sqrt{2})V(cs)V^*(ud)(\bar{s}_L\gamma_\nu c_L)(\bar{u}_L\gamma_\nu d_L) + h.c. \tag{51}$$

Radiative QCD corrections lead to a renormalization at scale m_c, often referred to as *ultraviolet* renormalization , since its scales are larger and thus more ultraviolet than m_c. One-loop contributions generate an operator different from $(\bar{s}_L\gamma_\nu c_L)(\bar{u}_L\gamma_\nu d_L)$, namely $(\bar{s}_L\gamma_\nu \frac{\lambda_i}{2} c_L)(\bar{u}_L\gamma_\nu \frac{\lambda_i}{2} d_L)$ with the λ_i denoting the $SU(3)$ matrices. The renormalization is therefore additive and not multiplicative , i.e. $\mathcal{L}_W^{\Delta C=1}(\mu = m_c) \not\propto \mathcal{L}_W^{\Delta C=1}(\mu < M_W)$. Considering all operators that under QCD renormalization can mix with the original transition operator(s) one can determine which are the multiplicatively renormalized operators and with which coefficients they appear in the effective Lagrangian by diagonalizing the matrix with the one-loop corrections of these operators.

However there is a more direct way to understand why QCD corrections double the number of transition operators and which operators are multiplicatively renormalized. Already on the one-loop level one has two types of couplings in colour space, namely $\mathbf{1}\otimes\mathbf{1}$ and $\lambda_i \otimes \lambda_i$ with λ_i denoting the eight Gell-Mann matrices. Higher loop contributions do not change this pattern since $\lambda_i\lambda_j \otimes \lambda_i\lambda_j$ can again be expressed through a linear combination of $\mathbf{1}\otimes\mathbf{1}$ and $\lambda_i \otimes \lambda_i$. This holds no matter what the Lorentz structure of the coupling is. For couplings involving left-handed quark currents only we have the simplification that the product of two such currents remains a product of two left-handed currents under Fierz transformations [98]. This allows us to write a current product of the form $\lambda_i \otimes \lambda_i$ as linear combination of $1 \otimes 1$ and $[1 \otimes 1]_{Fierz}$, where the second term is the Fierz transformed product. Consider now the interaction described by Eq.(51) $c_L \to s_L \bar{d}_L u_L$. Its operator is purely isovector. Yet under V spin , which groups (u, s) into a doublet with c and d being singlets the final state is a combination of $V = 0$ and $V = 1$. Fermi-Dirac statistics tells us that if u and s are in the antisymmetric $V = 0$ [symmetric $V = 1$] configuration, they have to be in the symmetric [antisymmetric] $SU(3)_C$ 6 $[\bar{3}]$ representation. I.e., the two multiplicatively renormalized operators have to be Fierz even and odd:

$$O_{\pm}^{\Delta C=1} = \frac{1}{2}[(\bar{s}_L \gamma_\nu c_L)(\bar{u}_L \gamma_\nu d_L) \pm (\bar{s}_L \gamma_\nu d_L)(\bar{u}_L \gamma_\nu c_L)] \tag{52}$$

Therefore

$$\frac{\mathcal{L}_W^{\Delta C=1}(\mu = m_c)}{(4G_F\sqrt{2})V(cs)V^*(ud)} = c_+ O_+^{\Delta C=1} + c_- O_-^{\Delta C=1} =$$

$$[c_1(\bar{s}_L \gamma_\nu c_L)(\bar{u}_L \gamma_\nu d_L) + c_2(\bar{u}_L \gamma_\nu c_L)(\bar{s}_L \gamma_\nu d_L)] \tag{53}$$

The coefficients $c_{1,2}$ can be expressed as follows at leading log level:

$$c_1|_{LL} = \frac{1}{2}(c_+ + c_-) \ , \ \ c_2|_{LL} = \frac{1}{2}(c_+ - c_-) \tag{54}$$

$$c_\pm \equiv \left[\frac{\alpha_S(M_W^2)}{\alpha_S(m_c^2)}\right]^{\gamma_\pm} \ , \ \ \gamma_+ = \frac{6}{33 - 2N_f} = -\frac{1}{2}\gamma_- \tag{55}$$

N_f denotes the active flavours. Next-to-leading log corrections are sizeable at the charm scale. They cannot be expressed in a compact analytical way; numerically one finds when including these contributions:

$$c_1|_{LL+NLL} \simeq 1.32 \ , \ \ c_2|_{LL+NLL} \simeq -0.58 \tag{56}$$

Noting that without QCD radiative corrections one has $c_1 = 1$ and $c_2 = 0$, QCD renormalization constitutes a quite sizeable effect. This is not surprising since the leading log result represents an expansion in powers of $\alpha_S \log M_W^2$ rather than just α_S.

With hadronic matrix elements evaluated at ordinary hadronic scales Λ_{NPD} rather than the heavy quark mass, one has to consider also renormalization from m_c down to Λ_{NPD}. This is often called *hybrid* renormalization since its scales are in the infrared relative to m_c and in the ultraviolet relative to Λ_{NPD}. Yet since m_c – unlike m_b – exceeds Λ_{NPD} by a moderate amount only, one does not expect hybrid renormalization to play a major role in most cases for charm. One notable exception is the $D^+ - D^0$ lifetime ratio, which will be discussed later.

There are analogous effects on the Cabibbo once and twice suppressed levels. The analogues of the Fierz even and odd operators of Eq.(52) are multiplicatively renormalized with the coefficients $c_{\pm}$ as in Eq.(55). In addition Penguin operators emerge on the Cabibbo disfavoured level. These renormalization effects with $c_- \simeq 1.9 > c_+ \simeq 0.74$ lead to the enhancement of the $\Delta I = 0$ [$\Delta I = 1/2$] over the $\Delta I = 1$ [$\Delta I = 3/2$] transition operators for once [doubly] Cabibbo suppressed modes. These issues will be addressed further in Sect.**9**.

4˙10.2. Sum Rules. Sum rules are an ubiquitous tool in many branches of physics where sums or integrals over observables – rates, moments of rates etc. – are related to a normalization condition reflecting unitarity etc. or a quantity that can be calculated in the underlying theory. They form an important ingredient in our treatment of deep inelastic lepton-nucleon scattering for example, where moments of structure functions are related to terms in an OPE. Examples are the Adler and the Gross-Llewellyn-Smith sum rules [99, 98].

Another celebrated case are the SVZ QCD sum rules named after Shifman, Vainshtein and Zakharov [100], which allow to express *low energy* hadronic quantities through basic QCD parameters. The starting point is again provided by an OPE in terms of local operators. Nonperturbative dynamics are parameterized through vacuum expectation values – or *condensates* – $\langle 0|\bar{q}q|0\rangle$, $\langle 0|G^2|0\rangle$ etc., since those have to vanish in perturbation theory. Those condensates are treated as free parameters the values of which are fitted from some observables. One typically matches up a quantity calculated on the quark-gluon level through a dispersion relation – see the next Subsection – with an ansatz for the hadronic observables; the stability of the match under variations of input values provides an intrinsic gauge for the theoretical control in this case. Introducing nonperturbative dynamics through condensates represents an approximation of less than universal validity: such an ansatz cannot be counted on to reproduce observables exhibiting rapid variations in, say, energy like narrow resonances and their phase shifts. In such situations one can hope at best for being able to treat 'smeared' hadronic observables, i.e. ones that have been averaged over some energy interval. Manifold experience shows that one has to allow for an irreducible theoretical uncertainty of about 20-30 % due to unknown contributions from higher operators in the OPE, excited states in the dispersion relations and due to the ansatz with condensates. Contrary to the situation with lattice QCD, one cannot hope for a systematic reduction in the theoretical uncertainty.

A very similar approach under the name of "lightcone sum rules" [102] has been developed for describing the formfactors in exclusive semileptonic decays of heavy flavour hadrons to be mentioned later.

There is a second and third class of sum rules that are relevant here, namely the so-called S(mall)V(elocity) [88] and the spin sum rules [103] that have been formulated for semileptonic decays of heavy flavour hadrons. They are based on *systematic* expansions in $1/m_Q$ and – for $b \to c$ transitions – in the velocity of the final state quark. Accordingly they do not exhibit this brickwall of about 20 - 30% in theoretical uncertainty, but can be improved successively.

4˙10.3. Dispersion relations. Dispersion relations are encountered in many branches of physics and in quite different contexts. The common element is that certain fundamental features of general validity can be imposed by requiring that physical quantities have to be analytical functions of their variables, when they are allowed to be complex. One then invokes Cauchy's theorem on path integrals in the complex plane to relate the real and imaginary part of these quantities to each other.

For example in classical electrodynamics causality implies field amplitudes to be analytic. This holds in particular for the dielectric constant ϵ when it is frequency dependent: $\vec{D}(\vec{x},\omega) = \epsilon(\omega)\vec{E}(\vec{x},\omega)$. Causality implies the Kramers-Kronig relation [104]

$$\mathrm{Im}\epsilon(\omega)/\epsilon_0 = -\frac{2\omega}{\pi} P \int_0^\infty d\omega' \frac{\mathrm{Re}\,\epsilon(\omega')/\epsilon_0 - 1}{\omega'^2 - \omega^2} \,, \tag{57}$$

where P denotes the principal part computation of this integral. In S matrix theory one postulates unitarity, Lorentz and crossing symmetry and analyticity. Dispersion relations relate the scattering amplitudes for different reactions through integrals.

Likewise one can relate the values of a two-point function $\Pi(q^2)$ in a quantum field theory at different complex values of q^2 to each other through an integral representation; q denotes a four-momentum. In particular one can evaluate $\Pi(q^2)$ for large *Euclidean* values $-q^2 = q_0^2 + |\vec{q}|^2$ with the help of an OPE and then relate the coefficients I_n^{OPE} of local operators O_n to observables like $\sigma(e^+e^- \to \mathbf{had.})$ and their moments in the physical, i.e. Minkowskian domain through an integral over the discontinuity along the real axis; the integral over the asymptotic arcs vanishes [98]:

$$I_n^{OPE} \simeq \frac{1}{\pi} \int_0^\infty ds \frac{s}{(s+q^2)^{n+1}} \cdot \sigma(s) \tag{58}$$

Such a procedure is based on there being only physical singularities – poles and cuts – on the real axis of q^2: then one can first calculate two-point functions for large Euclidean values of q^2 and secondly one will not pick up extra unphysical contributions from poles etc. This is the basis of the derivation of the celebrated QCD sum rules by the ITEP group [100].

Such dispersion relations are used to calculate transition rates in the HQE and to derive new classes of sum rules [88].

4˙10.4. Final State Interactions (FSI) and Watson's theorem. The mass of charm hadrons places them into an environment populated by many non-charm resonances, hadronic thresholds etc. making FSI quite virulent. This provides for a particularly challenging dynamical environment. Let us consider the decay of a meson. The primary weak force transmogrifies the initially present valence quark and antiquark into two quarks and antiquarks. Yet those will not rearrange themselves immediately into two mesons that emerge as asymptotic states. Typically quarks and antiquarks will be exchanged, they can change their flavour identity thus giving rise to final states that are absent otherwise, and even additional $q\bar{q}$ pairs can be excited. Precisely since the forces driving these processes are strong, those secondary interactions cannot be ignored or treated to first (or any finite) order only. They can induce even spectacular resonance enhancements (or depletions for that matter). This is sometimes described by saying that the initially produced two quark-antiquark clusters can and typically will *rescatter* into different kinds of two-meson or even multi-meson final states.

Fortunately there is a modicum of theoretical guidance for dealing with this quagmire as sketched by the following remarks.

- While FSI can change the nature of the final state dramatically, they mainly rearrange the rate between different channels *without* create overall rate. I.e., typically they do not increase or decrease the total nonleptonic or semileptonic widths.

- However FSI can affect even fully inclusive transitions. As we will discuss later the nearby presence of a hadronic resonance of appropriate quantum numbers can enhance or suppress significantly the width of a charm hadron – an effect that would constitute a violation of quark-hadron duality.
- With the strong interactions conserving isospin and G parity, possible rescatterings are constrained by these quantum numbers.
- The most treatable case after total rates is provided by two-body final states, where we include hadronic resonances in the latter, due to their 'trivial' kinematics. A small number of quark-level diagrams can drive a large number of hadronic transitions. Consider for example $D^0 \to K^-\pi^+$ where two different four-quark operators contribute changing isospin by $1/2$ and $3/2$:

$$T(D^0 \to K^-\pi^+) = e^{i\alpha_{1/2}}T_{1/2} + e^{i\alpha_{3/2}}T_{3/2} \tag{59}$$

  A priori one expects – correctly, as it turns out – that the FSI generate a nontrivial relative phase between the two different isospin amplitudes $T_{1/2,3/2}$ – $\alpha_{1/2} \neq \alpha_{3/2}$ – and affect also their size. As we will discuss later in detail, such relative strong phases are a conditio sine qua non for *direct* CP asymmetries to arise. A well-known theorem is frequently quoted in this context, namely Watson's theorem. Below we will describe it mainly to make explicit the underlying assumptions and corresponding limitations.
- Novel theoretical frameworks have been put forward recently to treat nonleptonic two-body decays of B mesons [105, 106]. However there is no a priori justification for applying such treatments to D decays.
- Three-body final states can be and are subjected to Dalitz plot analyses. Unfortunately theory can provide very little guidance beyond that.

In describing *Watson's theorem* we follow the discussion in Ref.[107]. For reasons that will become clear we consider $K \to n\pi$.

A $\Delta S[C, ...] \neq 0$ process has to be initiated by weak forces which can be treated perturbatively. Yet the final state is shaped largely by strong dynamics mostly beyond the reach of a perturbative description. Nevertheless one can make some reliable theoretical statements based on symmetry considerations – sometimes.

With the strong interactions conserving G-parity a state of an even number of pions cannot evolve *strongly* into a state with an odd number. Therefore

$$K \stackrel{H_{weak}}{\longrightarrow} 2\pi \stackrel{H_{strong}}{\not\longrightarrow} 3\pi \tag{60}$$

On the other hand, the two pions emerging from the weak decay are not asymptotic states yet; due to the strong forces they will undergo rescattering before they lose sight of each other. Deriving the properties of these strong FSI from first principles is beyond our present computational capabilities. However, we can relate some of their properties to other observables.

Let us assume the weak interactions to be invariant under time reversal:

$$TH_WT^{-1} = H_W \tag{61}$$

We will show now that even then the amplitude for $K^0 \to 2\pi$ is complex; the strong FSI generate a phase, which actually coincides with the S wave $\pi\pi$ phase shift δ_I taken at energy M_K [108]. That is, the amplitude is real, except for the fact that the two pions interact before becoming asymptotic states.

At first we allow the phase for the $K^0 \to 2\pi$ amplitude to be arbitrary:

$$\langle (2\pi)_I^{out} | H_W | K^0 \rangle = |A_I| e^{i\phi_I} \tag{62}$$

where the label I denotes the isopin of the 2π state. With T being an antiunitary operator and using $T^\dagger T = \mathfrak{I}^\dagger \mathfrak{I}$ with $\mathfrak{I}$ denoting the complex conjugation operator we have

$$\langle (\pi\pi)_I; out | H_W | K \rangle = \langle (\pi\pi)_I; out | T^\dagger T H_W T^{-1} T | K \rangle^* = \langle (\pi\pi)_I; in | H_W | K \rangle^* \ , \tag{63}$$

since for a single state – the kaon in this case – there is no distinction between an in and out state. After inserting a complete set of out states

$$\langle (\pi\pi)_I; out | H_W | K \rangle = \sum_n \langle (\pi\pi)_I; in | n; out \rangle \langle n; out | H_W | K \rangle^* \ , \tag{64}$$

where the S matrix element $\langle (\pi\pi)_I; in | n; out \rangle$ contains the delta function describing conservation of energy and momentum, we can analyze the possible final states. The only hadronic states allowed kinematically are 2π and 3π combinations. With G parity enforcing

$$\langle (\pi\pi)_I; in | 3\pi; out \rangle = 0 \tag{65}$$

only the 2π out state can contribute in the sum:

$$\langle (\pi\pi)_I; out | H_W | K \rangle = \langle (\pi\pi)_I; \mathrm{in} | (\pi\pi)_I; \mathrm{out} \rangle \langle (\pi\pi)_I; out | H_W | K \rangle^* \tag{66}$$

This is usually referred to as the condition of *elastic* unitarity . With the S matrix for $(\pi\pi)_I \to (\pi\pi)_I$ given by

$$S_{\mathrm{elastic}} = \langle (\pi\pi)_I; \mathrm{out} | (\pi\pi)_I; \mathrm{in} \rangle = e^{-2i\delta_I} \tag{67}$$

we have

$$\langle (\pi\pi)_I; \mathrm{out} | H_W | K^0 \rangle = | \langle (\pi\pi)_I; \mathrm{out} | H_W | K^0 \rangle | e^{i\delta_I} \ ; \tag{68}$$

i.e., as long as H_W conserves T, the decay amplitude remains real after having the strong phase shift factored out. This is Watson's theorem in a nutshell.

FSI also affect the decays of *heavy* flavour hadrons, yet we *cannot* apply Watson's theorem blindly even for T conserving H_W. In particular it would be absurd to assume *elastic* unitarity to apply in two-body or even quasi-two-body beauty decays: strong FSI are bound to generate additional hadrons in the final state. The decays of *charm* hadrons provide a borderline case: while the FSI can change the identity of the emerging particles and can produce additional hadrons, their impact is moderated since the available phase space is less than abundant. This is consistent with the observation that (quasi-)two-body modes constitute the bulk of nonleptonic D decays, although we have not learnt yet how to assign precise numbers to this statement, see Sect.**9**˙4. Introducing the concept of *absorption* – $T_f \to \eta_f T_f$ with $|\eta_f| < 1$ – provides a useful *phenomenological* approximation for parameterising such inelasticities.

4˙10.5. Zweig's rule. The Zweig rule goes back to the earliest days of the quark model [109]. It can be expressed as follows: In scattering or decay processes driven by the *strong* interactions those quark diagrams dominate where all valence quarks and antiquarks from the initial state are still present in the final state; i.e., initially present quarks and antiquarks do not annihilate.

The motivation for this selection rule came from the observation that the ϕ meson interpreted as an $s\bar{s}$ bound state decays mainly into a $K\bar{K}$ pair rather than a kinematically favoured pion pair. The rule was later somewhat extended by stating that all *disconnected* quark diagrams are suppressed.

Obviously such a rule holds only approximately. It was the discovery of the extremely narrow J/ψ resonance that turned the Zweig rule from respectable folklore into a dynamical notion based on colour symmetry and QCD's asymptotic freedom. For it was realized that an ortho[para]-quarkonium state has to annihilation into (at least) three [two] gluons to decay and that their couplings become smaller for increasing quarkonium masses:

$$\Gamma[\bar{Q}Q]_{\mathrm{ortho}} \propto \alpha_S^3(m_Q) < \Gamma[\bar{Q}Q]_{\mathrm{para}} \propto \alpha_S^2(m_Q) \tag{69}$$

Thus one can estimate how much 'Zweig forbidden' transitions are suppressed, and how it depends on the specifics of the decaying state.

4˙11. *On quark-hadron duality.* – Quark-hadron duality – or duality for short – is one of the central concepts in contemporary particle physics. It is invoked to connect quantities evaluated on the quark-gluon level to the (observable) world of hadrons. It is used all the time as it has been since the early days of the quark model and of QCD, more often than not without explicit reference to it. A striking example of the confidence the HEP community has in the asymptotic validity of duality was provided by the discussion of the width $\Gamma(Z^0 \to H_b H_b' X)$. There was about a 2% difference in the predicted and measured decay width, which lead to lively debates on its significance vis-a-vis the *experimental* error. No concern was expressed about the fact that the Z^0 width was calculated on the quark-gluon level, yet measured for hadrons. Likewise the strong coupling $\alpha_S(M_Z)$ is routinely extracted from the perturbatively computed hadronic Z^0 width with a stated theoretical uncertainty of $\pm$ 0.003 which translates into a theoretical error in $\Gamma_{had}(Z^0)$ of about 0.1%.

There are, however, several different versions and implementations of the concept of duality. The problem with invoking duality *implicitly* is that it is very often unclear which version is used. In B physics – in particular when determining $|V(cb)|$ and $|V(ub)|$ – the measurements have become so precise that theory can no longer hide behind experimental errors. To estimate theoretical uncertainties in a meaningful way one has to give clear meaning to the concept of duality; only then can one analyze its limitations. In response to the demands of heavy flavour physics a considerable literature has been created on duality over the last few years, which we want to summarize. We will sketch the underlying principles; technical details can be found in the references we list.

Duality for processes involving time-like momenta was first addressed theoretically in the late '70's in references [110] and [111]. We sketch here the argument of Ref.[110], since it contains several of the relevant elements in a nutshell. The cross section for $e^+e^- \to hadrons$ can be expressed through an operator product expansion (OPE) of two hadronic currents. One might be tempted to think that by invoking QCD's asymptotic freedom one can compute $\sigma(e^+e^- \to \mathbf{hadrons})$ for large c.m. energies $\sqrt{s} \gg \Lambda_{QCD}$ in terms of quarks (and gluons) since it is shaped by short distance dynamics. However production thresholds like for charm induce singularities that vitiate such a straightforward computation. This complication can be handled in the following way: One evaluates the OPE in the (deep) Euclidean region thus avoiding proximity to singularities induced by hadronic thresholds; then one analytically continues it into the Minkowskian domain through a dispersion relation. There is a price to be paid: in general one cannot obtain the cross section as a point-for-point function of s, only averaged – or 'smeared' – over an energy interval, which can be written symbolically as follows:

$$\langle \sigma(e^+e^- \to hadrons) \rangle \simeq \int_{s_0}^{s_0+\Delta s} ds \sigma(e^+e^- \to hadrons) \tag{70}$$

This feature is immediately obvious: for the smooth s dependence that the OPE necessarily yields in Euclidean space has to be compared to the measured cross section $e^+e^- \to$ hadrons as a function of s, which has pronounced structures, in particular close to thresholds for $c\bar{c}$-production.

This simple illustration already points to the salient elements and features of duality and its limitations [112, 84]:

- An OPE description for the observable under study is required in terms of quark *and* gluon degrees of freedom. (12)
- The extrapolation from the Euclidean to the Minkowskian domain implies some loss of information: in general one can calculate only hadronic observables that are averaged over energy.

$$\langle \sigma^{hadronic} \rangle_w \simeq \langle \sigma^{partonic} \rangle_w \tag{71}$$

  where $\langle ... \rangle_w$ denotes the smearing which is an average using a smooth weight function $w(s)$; it generalizes the simplistic use of a fixed energy interval:

$$\langle ... \rangle_w = \int ds \, ... \, w(s) \tag{72}$$

- Some contributions that are quite insignificant in the Euclidean regime and therefore cannot be captured through the OPE can become relevant after the analytical continuation to the Minkowskian domain, as explained below. For that reason we have used the approximate rather than the equality sign in Eq.(71).
- One can make few *universal* statements on the numerical validity of duality. How much and what kind of smearing is required depends on the specifics of the reaction under study.

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(12)The name *parton*-hadron duality is actually more appropriate in the sense that gluon effects have to be included for duality to hold.

The last item needs expanding right away. The degree to which $\langle\sigma^{partonic}\rangle_w$ can be trusted as a theoretical description of the observable $\langle\sigma^{hadronic}\rangle_w$ depends on the weight function w, in particular its width. It can be broad compared to the structures that may appear in the hadronic spectral function, or it could be quite narrow, as an extreme case even $w(s) \sim \delta(s - s_0)$. It has become popular to refer to the first and second scenarios as *global* and *local* duality, respectively. Other authors use different names, and one can argue that this nomenclature is actually misleading.

When one treats distributions rather than fully integrated widths, another complication arises. Consider for example inclusive semileptonic transitions $H_Q \to \ell\nu X_q$. The lepton spectrum is expressed through an expansion in powers of $1/m_Q(1 - x_l)$ rather than $1/m_Q$ where $x_l = 2E_l/m_Q$. It obviously is singular for $x_l \to 1$ and thus breaks down in the endpoint region. One can still make statements on partially integrated spectra; yet for semileptonic *charm* decays the situation becomes somewhat marginal since μ/m_c is not a small number to start with.

A fundamental distinction concerning duality is often drawn between semileptonic and nonleptonic widths. Since the former necessarily involves smearing with a smooth weight function due to the integration over neutrino momenta, it is often argued that predictions for the former are fundamentally more trustworthy than for the latter. However, such a categorical distinction is overstated and artificial. Of much more relevance is the differentiation between distributions and fully integrated rates sketched above.

No real progress beyond the more qualitative arguments of Refs. [110] and [111] occurred for many years. For as long as one has very limited control over nonperturbative effects, there is little meaningful that can be said about duality violations. Yet this has changed for heavy flavour physics with the development of heavy quark expansions.

The possibility of duality violations clearly represents a theoretical uncertainty. However it is not helpful to lump all such uncertainties into a single 'black box'. For proper evaluation and analysis it is useful to distinguish between three sources of theoretical errors:

1. unknown terms of higher order in α_S;
2. unknown terms of higher order in $1/m_Q$;
3. uncertainties in the input parameters α_S, m_Q and the expectation values.

Duality violations constitute uncertainties *over and above* these; i.e. they represent contributions not accounted for due to

- truncating these expansions at finite order and
- limitations in the algorithm employed.

These two effects are not unrelated. The first one means that the OPE in practice is insensitive to contributions of the type $e^{-m_Q/\mu}$ with μ denoting some hadronic scale; the second one reflects the fact that under a analytic continuation the term $e^{-m_Q/\mu}$ turns into an oscillating rather than suppressed term $\sin(m_Q/\mu)$.

Of course we do not have (yet) a full theory for duality and its violations. Yet we know that without an OPE the question of duality is ill-posed. Furthermore in the last few years we have moved beyond the stage, where we could merely point to folklore. This progress has come about because theorists have – driven by the availability of data of higher and higher quality – developed a better understanding of the physical origins of duality violations and of the mathematical portals through which they enter the formalism.

Again charm studies can teach us lessons on duality that are neatly complementary to those from light quark studies on one hand and beauty physics on the other:

- It has been estimated that duality violating contributions to $\Gamma_{SL}(B)$ fall safely below 1/2 % and thus are expected to remain in the 'noise' of other theoretical uncertainties, i.e. to not exceed unknown higher order (in α_S as well as $1/m_Q$) contributions [84].

  The expansion parameter μ/m_c in charm decays on the other hand provides at best only a moderate suppression for higher order contributions, and at the same time limitations to duality will become more relevant and noticeable. This means that while we cannot have confidence in quantitative predictions, we can learn valuable lessons from a careful analysis of the data.

- A duality violating contribution $e^{-m_b/\mu}$ will remain in the theoretical 'noise' level. Yet the charm analogue $e^{-m_c/\mu}$ might become visible, again meaning that a careful study of charm dynamics can teach us lessons on the transition from short- to long-distance dynamics that could not be obtained in beauty decays.

4˙12. *Resume on the theoretical tools.* – The fact that the charm mass exceeds ordinary hadronic scales Λ_{NPD} provides a new expansion parameter – $\Lambda_{NPD}/m_c < 1$ – and thus a very useful handle on treating nonperturbative dynamics. Yet the excess is only moderate. Therefore – unlike the situation for beauty – nonperturbative effects can still be sizeable or even large, and it constitutes a theoretical challenge to bring them under control. However we view the glass as (at least) half full rather than (at most) half empty. Exactly because nonperturbative effects are sizeable, one can learn important lessons on nonperturbative dynamics in a novel, yet still controlled environment by analysing charm interactions in a detailed way.

Encouraging evidence that this is not an idle hope – that we are developing a better understanding of nonperturbative dynamics at the charm scale – is provided by the realization that, as described later in detail, HQS provides an approximate understanding of charm spectroscopy and HQE reproduce correctly – in part even correctly *pre*dicted – the observed pattern of charm lifetimes. NRQCD is yielding complementary new insights, and there is the expectation that lattice QCD will provide us not only with valuable guidance in charm physics, but even with reliable quantitative answers.

4˙13. *On Future Lessons.* – Our intent is not to write a historical review or present a mere status report. We want to emphasize the importance of future charm studies based on a triple motivation:

- As sketched in this section there is a vast array of theoretical technologies that are truly based on QCD, yet require some additional assumptions. They apply to beauty physics with considerably enhanced validity and thus can be *tested* there. Yet we view the fact that nonperturbative effects are larger in charm than in beauty physics as a virtue rather than a vice, at least for the discriminating observer: charm physics constitutes a rich lab to *probe* (rather than test) these methods, to provide new insights into the transition from the nonperturbative to the perturbative domain. We have to be prepared that these methods will occasionally fail; yet we shall be able to obtain valuable lessons even from such failures.

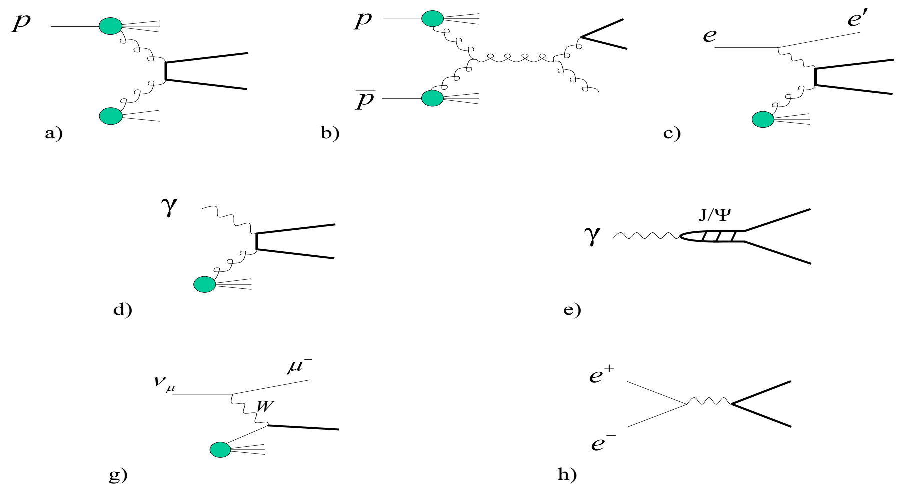

Fig. 10. – Diagrams for charm production: Hadroproduction (a,b); electron-proton production (c); photoproduction point like photon (d), resolved photon (e); neutrino (f); e^+e^- (g).

- A more detailed knowledge and understanding of charm physics than presently available is also essential for a better understanding of beauty physics and for a fuller exploitation of the discovery potential for New Physics there. This starts with the trivial observation that knowing charm branching ratios and decay sequences are important for interpreting beauty decays. Secondly, as indicated above, the theoretical technologies employed in beauty decays can be cross-referenced in charm decays. Lastly, a detailed understanding of charm spectroscopy is important in properly interpreting certain $B \to \ell\nu X_c$ transitions and the information they can yield concerning the underlying QCD treatment. This last more subtle point will be explained later.
- High sensitivity studies of $D^0 - \bar{D}^0$ oscillations, CP violation and rare decays provide a novel window onto conceivable New Physics – actually of *non*-standard extensions of the SM as indicated in the previous subsection.

5. – Production dynamics

Understanding the production processes for hadrons containing charm quarks is of obvious practical importance if one wants to obtain a well-defined sample of those hadrons for studying their decays. Yet new conceptual insights into QCD can be gained as well. In the following we address these two issues for different production reactions. In doing so one has to treat separately the cases of hidden and open charm hadrons, whose scenarios are quite different theoretically as well as experimentally. Rather than give an exhaustive discussion we aim at describing a few telling examples. Recent reviews can be found in [113, 114, 115, 116, 117], and predictions in [118].

As described before in our historical sketch of Sect.**2**, the use of a variety of intense particle beams on a wide variety of nuclear fixed targets dates back to the beginning of the charm adventure, and it constitutes a mature technique for investigating charm production. On the other hand, heavy flavour physics at hadron colliders, after pioneering work at the ISR, has undergone a renaissance at CDF. There are multiple motivations for studying hadroproduction of open and hidden charm states:

- The production of heavy flavour hadrons presents new tests of our quantitative understanding of QCD. Their worth is enhanced by the fact that there are similar ingredients in the theoretical treatment of charm and beauty production.
- It serves as a sensitive and efficient probe for determining gluon distributions inside nucleons.
- Understanding the production mechanisms helps us in fully harnessing the statistical muscle of hadroproduction for studies of weak *decays* of charm hadrons.
- Analyzing charm production inside heavy nuclei provides us with insights into how QCD's dynamics act under exotic or even extreme conditions. Furthermore it can signal the onset of the quark-gluon plasma as discussed later.

We have chosen to organize the vast material in the following way: first we will describe hidden charm production in the different settings, then we will turn to open charm produced through e^+e^- annihilation, at fixed target experiments, hadronic colliders and deep inelastic lepton-nucleon scattering and conclude with charm production inside heavy nuclei.

5˙1. *Charmonium production.* – A priori there are three experimentally distinct scenarios for the production of prompt J/ψ: the *secondary* production via a para-charmonium state χ_c cascading down $e^+e^- \to \chi_c + X \to J/\psi + \gamma + X$ or *primary* production of J/ψ together with the excitation of two charm hadrons – like $e^+e^- \to J/\psi + D\bar{D}' + X$ –, which is a Zweig allowed process, or without such additional charm states, which is not. In 1995 the CDF collaboration [119, 120] discovered that B meson decays are not the major source of J/ψ production in hadronic collisions : many J/ψ are prompt rather than the decay products of an object with a lifetime of around 1 ps. The production of these 'direct' charmonia was found to be enhanced by a factor of about fifty (Fig. 11 with respect to predictions of the theoretical model of that time, the colour-singlet model, Fig.12. In this model it is assumed that charmonium states can get excited only via their $c\bar{c}$ component making the production of para-charmonium – χ_c – to dominate over that for ortho-charmonium – J/ψ. There is no reason beyond simplicity, why the J/ψ cannot be produced via a $c\bar{c}$ octet component. The most radical of such colour octet models is often called the 'colour evaporation model', where the octet sheds its colour with unit probability via soft gluons.

These models can be embedded in NRQCD, see Sect. **4**˙8, which was developed partly in response to the challenge posed by J/ψ production at the TEVATRON. By including charmonium production off colour *octet* $\bar{c}c$ configurations, where colour is shed via soft gluons, NRQCD is able to reproduce these data; the colour octet component Fig.12 thus represents by far the dominant source of prompt charmonia at TEVATRON energies – in clear contrast to the situation at lower energies.

Our understanding can be further tested by measuring the polarization α of the J/ψ and ψ' defined as $(d\Gamma/d\cos\theta) \propto 1 + \alpha\cos^2\theta$ in the angular distribution of decay leptons pairs from charmonium; $\alpha = 1[-1]$ corresponds to pure transverse [longitudinal] polarization. Both charmonia states are predicted [121, 122] to be increasingly transversely polarized with growing $p_\perp$ since one expects the transverse polarization of almost 'on-shell' gluons to be transferred to the $\bar{c}c$ bound state produced from them.

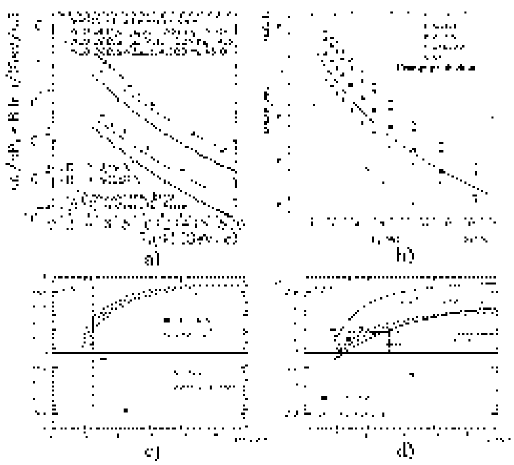

Fig. 11. – CDF results on cross section for J/ψ (a), ψ' (b), and polarization (c,d). Data from [119, 120], theoretical predictions from [121, 122].

However this effect is certainly not apparent in the data, see Fig.11 c), d) [123, 124]. More data has become available from CDF out of RUN II data taking period presently in progress [125]. Again, this might not be a fatal flaw in NRQCD; it might just mean that contributions of higher order in α_S and in v are still sizeable for charmonia and affect polarization more than cross sections, which are more robust against higher order contributions.

Real photoproduction experiments provide also results on hidden charm states, whose diffractive production proceeds via VMD coupling of the beam photon to $J^{PC} = 1^{--}$ mesons such as J/ψ. Such studies are generally limited to dimuon final states, since dielectron decay modes are hindered by the presence of electron-positron pairs copiously produced by Bethe-Heitler mechanisms. Perhaps more importantly, the e^+e^- decay mode is made difficult by the presence of a very long tail in the dielectron invariant mass spectrum, due to bremmsstrahlung, which needs to be corrected for. Very recently, the first observation of $\psi(3770)$ was preliminarily reported by FOCUS [126].

Experiments E760 and E835 at the Fermilab antiproton source have performed precision measurements of all charmonium states, also measuring the χ_{c1} for the first time. Charmonium states are produced by the collisions of antiprotons on a hydrogen jet target, thus providing interaction whose geometry is effectively the one typical of fixed target experiments. Experiment E835 charmonium results are discussed in Sect.**6**.

5˙2. *Charm at LEP (mainly).* – With no hadron being present in the initial state, e^+e^- annihilation represents the simplest scenario. At the same time charm was the first quantum number high enough in mass that one can invoke perturbative QCD to describe $e^+e^- \rightarrow \boldsymbol{hadrons}$ *below* as well as *above* its threshold; i.e., the ratio R had reached a constant value below threshold and a higher one above it.

One of the most intriguing aspects theoretically has been discussed in Sect.**4**˙11,

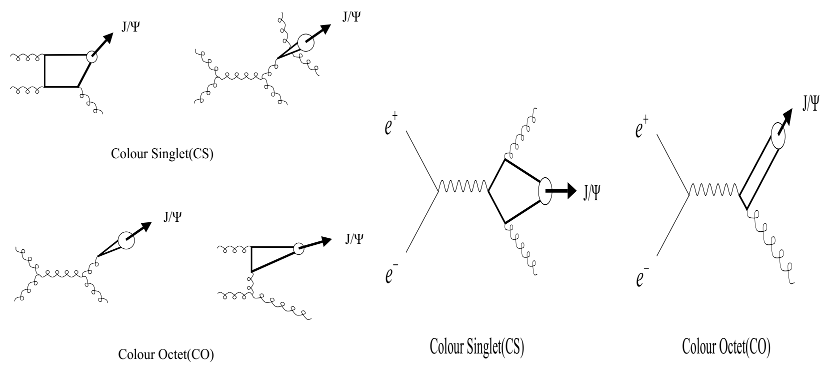

Fig. 12. – Colour-singlet and colour-octet diagrams for hadroproduction (left) and e^+e^- production (right) of charmonium; photons and gluons are denoted by wave and curly lines, respectively.

namely how quark-hadron duality and the approach to it is realized in nature. The conclusion is that starting at about 1 GeV above threshold perturbative QCD can be employed for predicting the total cross section for $e^+e^- \to \mathbf{hadrons}$. The experimental situation just above charm threshold had been somewhat unsettled with some data sets showing an unusually large cross section. Measurements [127] done by the BES collaboration in Beijing have clarified the situation; the value of R does not seem to be excessively large. Future studies from CLEO-c should settle it completely. The transition region around 4 GeV will presumably remain beyond theoretical control, since so many thresholds for exclusive final states open up: $e^+e^- \to D\bar{D}$, $D^*\bar{D} + D\bar{D}^*$, $D^*\bar{D}^*$, $D_s^+D_s^-$, ..., $\Lambda_c\bar{\Lambda}_c$ etc.. One can attempt to describe this highly complex landscape through models involving a coupled-channel approach [128, 129]; yet the predictions based on such models are not reliable, since they are quite unstable under variations of the model parameters. Nevertheless important measurements can be performed there: in particular the *absolute* branching ratios for the different charm hadrons can be measured in a model independent way as explained in Sect.**9**˙3 of this review.

BELLE has shown highly surprising data on double $c\bar{c}$ production: it finds the J/ψ to be accompanied more often than not by an additional $\bar{c}c$ pair [130]:

$$\frac{\sigma(e^+e^- \to J/\psi\bar{c}c)}{\sigma(e^+e^- \to J/\psi X)} = 0.59^{+0.15}_{-0.13} \pm 0.12 \tag{73}$$

There is no good idea from theory (yet) how such a large ratio could be accommodated.

Measurements by HRS, MARK II and TASSO at the PETRA and PEP storage rings at DESY and SLAC, respectively, had provided the first reliable information on the fragmentation functions of charm and beauty quarks described in Sect. **5**˙6. Yet they have been superseded by ARGUS and CLEO measurements with their much higher statistics and by LEP and SLD data, the latter for beauty as well as charm quarks; see [131, 132] for concise reviews. CLEO finds a deviation from a Peterson *et. al.*-type fragmentation function for D_s^+, D_s^{*+} [133]. It is actually more difficult to measure the charm than the beauty fragmentation fragmentation at LEP, since in the case of charm one has to rely on exclusively reconstructed charm hadrons to obtain sufficient purity of the sample, which reduces considerably the statistics [134]. Furthermore an interpretation of the data is less straightforward, among other reasons due to the secondary production of charm via gluon splitting $g \to c\bar{c}$. Nevertheless there exists a strong twofold motivation for determining the charm fragmentation function as accurately as possible: (i) it is an important ingredient in predicting charm production cross sections and distributions to be measured at the TEVATRON; (ii) comparing it to the beauty fragmentation function will shed further light on the nonperturbative dynamics driving it.

The level of production can be computed (with $m_c = 1.5 \pm 0.3$ GeV), and when compared to recent data [132] it is found a couple of standard deviations below.

LEP experiments provide information on charm production through studies of $\gamma\gamma \to c\bar{c}$, where the initial state is realized by initial state radiation off both beams . Recent results are discussed in Ref.[114]. $D^{*\pm}$ production has been measured by LEP experiments at $\sqrt{s} = 183-209$ GeV and found in agreement with NLO QCD predictions. L3 also measured $\sigma(\gamma\gamma \to c\bar{c}X)$ as a function of the $\gamma\gamma$ invariant mass finding reasonable agreement with NLO QCD for $m_c = 1.2$ GeV; for $m_c = 1.5$ GeV one predicts a 50% lower cross section.

Another observable that can be computed in NLO QCD is the aforementioned gluon splitting probability $g_{c\bar{c}}$ of $c\bar{c}$ for $(e^+e^- \to q\bar{q}g, g \to Q\bar{Q})$. An OPAL result yields $g_{c\bar{c}} = (3.20 \pm 0.21 \pm 0.38) \times 10^{-2}$ [135], which is higher than theoretical estimates as well as the L3 measurement [136] $g_{c\bar{c}} = (2.45 \pm 0.29 \pm 0.53) \times 10^{-2}$. Likewise ALEPH and DELPHI find higher than predicted values for the corresponding quantity $g_{b\bar{b}}$.

An area of great importance to the validation of the electroweak sector of the Standard Model is the determination of the forward-backward asymmetries for charm and beauty jets $A_{FB}^{c\bar{c}}, A_{FB}^{b\bar{b}}$, and the ratios of charm and beauty quark partial widths R_b, R_c $R_c = \Gamma(Z^0 \to c\bar{c})/\Gamma(Z^0 \to hadrons)$, see [137, 138] for an extensive review. The charm FB asymmetries had been measured before the LEP era [139], [140], [141], but only the huge data sets gathered at LEP allowed meaningful searches for manifestations of New Physics. The situation has changed considerably over the years, as it can be realized browsing the LEP Electroweak Working group pages [142]. Measurements of R_b, R_c in 1995 [143] differed from SM predictions by $+3.7\sigma, -2.56\sigma$, while they appeared totally consistent in 2002 [144] $+1.01\sigma, -0.15\sigma$. On the other hand, $A_{FB}^{c\bar{c}}, A_{FB}^{b\bar{b}}$ that in 1995 were completely consistent with the SM within their relatively large errors, in 2002 represent a pull of $-0.84\sigma, -2.62\sigma$ respectively, with $A_{FB}^{b\bar{b}}$ being the second largest contribution of pulls in the fit to the SM parameters, after the intriguing NuTeV result on $\sin^2\theta_W$ [145].

5˙3. *Photoproduction.* – The real photon has two components, namely a hadronic one *a la* vector meson dominance , and one coupling directly to quarks via their electric charge. Due to the small weight of the former, photon beams provide a cleaner environment than hadron beams since mostly there is no hadronic jet from beam fragmentation. One still has to contend with a large background of light hadrons. Yet the charm-to-total cross-section ratio of about 1/100 is considerably higher than the 1/1000 for hadroproduction. Also the theoretical treatment of photoproduction is easier than of hadroproduction since only one hadron participates in the collision.

The recent success in gathering sizeable samples of events with both charm and anti-charm hadrons allows novel probes of perturbative QCD and has already lead to new insights on QCD dynamics.

Scattering electrons off protons at small values of Q^2 provides a fluid transitions to *real* photoproduction experiments with $Q^2 = 0$. In real photoproduction high-energy, high-intensity beams impinge on nuclear targets and can produce charm states. The tree-level production mechanism proceeds via the fusion[146] of beam photon off a gluon emitted by the nucleus (Fig. 10).

The charm cross-section is sensitive to the charm quark mass, and it has been thoroughly measured from threshold up to HERA energies, and compared to QCD predictions [118]. A value of 1.5 GeV for the *pole* charm quark mass is favoured by the data with large errors due to the choice of other theory parameters.

Real photoproduction is studied via the very large data samples collected by fixed target experiments [113]. The reconstruction of both D and $\bar{D}$ in the same event allows one to study $D\bar{D}$ correlations that can be predicted in principle by QCD. Important variables are $\Delta\phi$, the angle between the particle and the antiparticle in the plane transverse to the beam, and the transverse P_T^2 momentum squared of the pair. At leading order, with $c\bar{c}$ quarks produced back to back, one expects $\Delta\phi = \pi$ and $P_T = 0$. Recent results from FOCUS [147] show that data disagree with predictions even when taking into account NLO contributions. There is also a small but highly significant excess of data at $\Delta\phi = 0$ suggesting that a small fraction of $D\bar{D}$ pairs are produced collinearly rather than back-to-back. These studies are a valuable tool to tune charm production computational algorithms, such as PYTHIA [148].

Particle-antiparticle asymmetry studies have been carried out in the past by photoproduction experiments NA14/2[149], E691 [150] and E687[151]. E691 and E687 measure a significant asymmetry for D^+, D^0 and D^{*+}, and one compatible with zero within large experimental errors for D_s^+, Λ_c^+. The asymmetry measured is ten times the one predicted by perturbative QCD. Mechanisms based on heavy-quark recombination have been proposed [152]. High statistics results from FOCUS are expected soon.

Finally, FOCUS [153] showed recently a null result on the production of double charm baryons reported by hyperon beam experiment SELEX [154] [155]. The issue is addressed in Sects.**6** and **6**˙3.5.

5˙4. *Fixed target hadroproduction.* – Experiments have been performed with a host of extracted meson as well as baryon beams and also internal beams on gas targets.

The leading particle effect is the most interesting phenomenological feature in charm production studies with extracted beams. This is the enhancement of the production of particles compared to the production of antiparticles, and it is due to the presence of quarks present both in the produced particle, and in the target nucleon or in the beam particle. The enhancement is represented (usually in a differential fashion in x_F and P_T^2) via the asymmetry variable

$$A \equiv \frac{N_{particle} - N_{antiparticle}/R}{N_{particle} + N_{antiparticle}/R} \tag{74}$$

where $R \equiv \bar{\epsilon}/\epsilon$ is the ratio of acceptances for particles and antiparticles.

As a matter of principle perturbative QCD can yield only a very small asymmetry in charm production; in contrast asymmetries as large as 50% have often been reported in the data. In the following we shall limit ourselves to outline the most striking features of charm production at extracted beams, referring the interested reader to recent reviews [156], [157], [113].

Tree-level hadroproduction proceeds via diagrams shown in Fig. 10(a,b). At fixed target, the availability of several kinds of beam allows one to study the leading particle effect. Recent results come from E791 (π^- beam) and SELEX (π^-, proton and Σ^- beams) . Asymmetry data for Λ_c^+ produced by π^- beams in E791 and SELEX do agree, while the asymmetry for proton and Σ^- beams in SELEX is much more pronounced, a clear manifestation of leading particle effect, since baryon beams will produce preferentially Λ_c^+ baryons, not antibaryons.

SELEX recently reported observation of four different $C = 2$ baryons [154] [155]. If confirmed, it would have profound implications for our understanding of charm production. We will discuss this issue in Sect.**6**˙3.5.

5˙5. *Hadroproduction at colliders*. – The study of charm physics at hadronic colliders was pioneered at the CERN ISR (see Sect.**3**˙2). Experiments done there showed evidence for much larger charm cross sections than expected, in particular in the forward region of up to 1.4 mb. It was finally understood that such high values were due to efficiency and acceptance corrections used to get cross sections out of low-acceptance mass-peak observations [158],[159].

The most lasting legacy is maybe the concept of *intrinsic charm* suggested a long tome ago by Brodsky and collaborators [160] to account for the larger charm production in the forward or projectile fragmentation region. It says that protons (and other hadrons) have a $\bar{c}c$ component that unlike in the conventional picture is *not* concentrated in the 'sea' at very small values of fractional momenta x: $|p\rangle \propto |uud\rangle + |uud\bar{c}c\rangle + `sea'$. This has been referred to as a higher Fock state in the proton wave function. There has been and still is an ongoing debate over the validity of this intriguing picture and the danger of double-counting. It seems now that in the framework of the heavy quark expansions one can assign the concept of intrinsic charm an unambiguous meaning [161].

Little new work theoretically as well as experimentally has been done on charm production at hadron colliders until recently. This lull seems to be coming to an end now. Most of our community was quite surprised when in 1991 CDF [162] demonstrated the ability of studying beauty quark physics in a p_T regime totally unsuited for the mainstream W physics the detector was conceived for. CDF's capabilities have been further boosted by the implementation of a detached vertex trigger described in Sect.**3**˙2, which provides online selection of events based on reconstructed decay vertices by the microstrip detector. The vertex trigger might allow reconstruction of charm correlations.

So far we have results on D^* production [66, 114]. In 2000 CDF published the only available measurements of open charm production cross sections in $p\bar{p}$ collisions at $\sqrt{s} = 1.8TeV$ for the process $D^{*+} \to D^0\pi^+ \to (K^-\mu^+X)\pi^+$ for in the rapidity and transverse momentum intervals $|\eta(D^{*+})| < 1.0, p_T(D^{*+} > 10GeV)$. The integrated cross section found $\sigma = 347 \pm 65 \pm 58nb$ exceeds calculations based on both NLO and FONLL. Such a result was confirmed in 2003 by the cross section measurement [163] for exclusive processes $D^0 \to K^-\pi^+$, $D^{*+} \to D^0\pi^+$, $D_s^+ \to \phi\pi^+$, obtained on a dataset selected by the new detached vertex trigger. Cross section found exceeds by 100% central value of theory predictions[164], although theory and experiment are compatible when considering theoretical uncertainties.

5˙6. *Deep inelastic lepton-nucleon scattering*. – As already mentioned in Sect.**2**, the first experimental signal for charm production outside cosmic rays came from dimuon events in deep inelastic neutrino nucleon scattering [14]

$$\nu_\mu N \to \mu^- H_c + X \to \mu^- \mu^+ + \tilde{X} \tag{75}$$

Also some early charm spectroscopy had been done in neutrino induced events. Today's main lessons from charm production by neutrinos are the following:

- It provides important information on $|V(cs)|$ & $|V(cd)|$, as described in Sect.**8**˙3.
- It allows extraction of the structure functions for d and s quarks and antiquarks from open charm and for the gluons from J/ψ production.
- In $\nu X \to l^- \Lambda_c^+ X$ one can measure the form factor of Λ_c baryons in the space-like region [165].
- Since Λ_c is expected to be produced with a high degree of longitudinal polarization, one could search for a T odd correlation $C_{\mathrm{T\ odd}} \equiv \langle \vec{\sigma}_{\Lambda_c} \cdot (\vec{p}_\Lambda \times \vec{p}_l) \rangle$ in semileptonic Λ_c decays $\nu N \to \Lambda_c^+ X \to (l^+ \nu \Lambda)_{\Lambda_c} + X$
- To determine the fundamental electroweak parameters at lower energies in different kinematical domains one has to understand – or at least model reliably – charm production, since it varies with energy and is different in charged vs. neutral current reactions. A very simple ansatz is often used here, namely the 'slow rescaling' model, where one replaces the usual scaling variable x by

$$x \to \xi = x \left(1 + \frac{"m_c"^2}{Q^2}\right)\left(1 - \frac{x^2 M_N}{Q^2}\right) \tag{76}$$

  It should be noted that $"m_c"$ here is merely a quark model parameter. Measuring its value with high accuracy from ν data does *not* mean we know the charm quark mass till we can derive Eq.(76) from QCD.

Charm production occurs in high energy neutrino interactions at the few percent level and to lowest order is described by the diagram in Fig.10 f), with strong dependence to the strange quark sea, since charm production off d quarks is Cabibbo-suppressed. This sensitivity is further enhanced in the case of antineutrino scattering, where only sea $\bar{d}$ and $\bar{s}$ quarks contribute with the latter dominating.

A wealth of results keeps coming from charm neutrino experiments using emulsion and electronic techniques, namely NOMAD and CHORUS at CERN and NuTeV at FNAL [115]:

- Emulsion experiments have been able to measure the inclusive charm production cross-section $\sigma(\nu_\mu N \to c\mu^- X)/\sigma(\nu_\mu N \to \mu^- X)$ which is of order five percent, while electronic experiments measured the inclusive D production rate $\sigma(\nu_\mu N \to D^0 \mu^- X)/\sigma(\nu_\mu N \to \mu^- X)$ about two percent. A recent analysis [166] combines both electronic and emulsion experiments results.
- Insights are gained on the hadronization of charm quark described through fragmentation functions, see Sect.**4**˙2. It should be noted that the fragmentation process is expected to be universal, i.e., independent of the *hard* scattering process under study; i.e., charm quarks emerging from, say, e^+e^- collisions dress into charmed

hadrons in the same fashion as those produced in lepton-nucleon scattering. Cross-section data for neutrino production are parameterized via the usual Peterson form $D(z) \propto [z(1 - z^{-1} - \epsilon_P/(1-z))]^{-1}$, and the customary kinematical variables p_T^2, f_h (the mean multiplicity of charmed hadron h) and z (the fraction of the quark longitudinal momentum carried by the charmed hadron). The fragmentation function $D(z)$ is peaked at $z = 0.8$ which means that the hadronization process is hard, and relatively energetic. Neutrino experiments measure the ϵ_P parameter and compare it to e^+e^- data, generally finding good agreement for D^* production.

Charm production in neutrino physics is thus an alive field, where great interest exists for the huge improvements which are expected at a Neutrino Factory [165], [167].

A new realm of analyzing heavy flavour production – of charm and beauty, open and hidden – has opened up in high energy $\sqrt{s} = 300 - 318\,\mathrm{GeV}$ electron-proton collisions studied at HERA by the H1, Zeus and HERA-C collaborations. Production of charm hadrons and charmonia can occur off gluons, $\bar{c}c$ pairs in the sea at small values of x and off an intrinsic charm component at medium and large values of x [160]. The stage is thus more complex than in e^+e^- annihilation – yet that should be viewed as a virtue, since data allow us access to these parton distribution functions.

Not only the proton *target* adds complexity to the phenomenology, but also the electron *projectile*, which effectively acts either as in deep inelastic lepton-nucleon scattering, (Fig.10 c), or in photoproduction (Fig.10 d), depending on the Q^2 region considered. In the photoproduction regime $(Q^2 \sim 0)$, then, the photon can produce charm via a direct point-like coupling to partons in the proton (Fig.10 d), or it can effectively act as a hadron a la vector meson dominance, Fig.10 e). In the latter case, any intrinsic charm components in the photon or proton particles may give origin to charm excitation processes, such as $cg \to cg$. The variable x_γ^{obs} is normally used to discriminate direct from resolved photon processes.

The experimental panorama is discussed in recent reviews [114, 168, 116]. Photoproduction cross-sections for D^* and D_s generally exceed the next to leading order (NLO) QCD predictions, as well as the fixed order plus next to leading logarithm (FONLL) calculations [169]. The photoproduction cross section is also measured as a function of x_γ^{obs}. This allows to show that a relevant contribution from charm excitation processes needs to be taken into account by theory. In DIS electroproduction regime $(Q^2 > 1\ \mathrm{GeV}^2)$ D^* cross sections are compared to predictions and found in fair agreement, although somehow undershooting data. The NRQCD prediction for J/ψ production yields a rising cross section for $z \equiv E_{J/\psi}/E_\gamma \to 1$, i.e. the kinematic boundary – in conflict with observation.

Another observable predicted by NRQCD predictions is the ratio of diffractive photoproduction rates of J/ψ vs. $\psi(2S)$. New data from H1 are found to be consistent with NRQCD predictions [170].

Resolved photon processes are expected to dominate the low-z inelastic region, while direct photon processes should dominate the region up to about $z \sim 0.9$, with diffractive photoproduction taking over at $z \sim 1$. Recent H1 and ZEUS results are reviewed and compared [116] to colour singlet (CS) and colour singlet + colour octet (CS+CO) predictions. As explained in Sect.4·8, the CO component enters naturally in NRQCD model, and is fitted to the large J/ψ cross section measured by CDF in 1995. HERA data are consistent with CS+CO contributions, although data do not rise as a function of z as rapidly as CS+CO predictions do. On the other hand, in electroproduction regime $Q^2 > 2GeV^2$, the inelastic J/ψ cross section measured by H1 clearly favours CS predictions.

Yet it would be premature to condemn NRQCD for this apparent discrepancy; for in its present level of sophistication it is not applicable in this kinematical domain. Future refinements of NRQCD should enable us to extend its applicability there.

5·7. *Hadroproduction inside heavy nuclei.* – The fabric of QCD is such that it can create an extremely rich dynamical landscape. To explore it fully one has to go beyond observing reactions involving single hadrons. When heavy nuclei collide with hadrons or other heavy nuclei the interactions between individual hadrons take place against the background of nuclear matter; this can lead to highly intriguing phenomena, of which we sketch two examples, namely the *lowering of the D meson mass* and *colour screening induced by the quark-gluon plasma* .

Most of the mass of pions and kaons, which are Goldstone bosons, is due to how approximate chiral symmetry is realized in QCD. Spontaneous chiral symmetry breaking leads to the emergence of non-vanishing quark and gluon condensates. Chiral invariance is partially restored in the medium of nuclear matter. It is expected that the masses pions and kaons exhibit inside nuclei get changed relative to their vacuum values, and that there is even a split between the masses of charge conjugate pairs with the nuclear medium providing an effective CPT breaking; it has been predicted that the masses of π^+ and π^- [K^+ and K^-] get shifted by about 25 [100] MeV. Experimental evidence for such effects has been inferred from the observation of pionic atoms and the study of the onset of K^+ and K^- production in heavy-ion collisions.

The situation is qualitatively different – and richer – in the charm sector since there the mass is due mostly to the c quark mass, and different scales enter the dynamics for the interactions with the nuclear medium. For the J/ψ and η_c only a small mass reduction of around 5 - 10 MeV is predicted, since charmonium masses are affected by mostly gluon condensates. D mesons on the other hand offer the unique opportunity to study the restoration of chiral invariance in a system with a single light valence quark [171]. A lowering of both $D^{\pm}$ masses is predicted with a relative shift of ~ 50 MeV in $M(D^+)$ vs. $M(D^-)$. One of the items in the GSI HESR proposal is to study these effects in detail in $\bar{p}Au$ collisions. Very intriguing effects are expected in the charm threshold region: at normal nuclear density the $D\bar{D}$ thresholds falls below the ψ' resonance; at twice nuclear density this threshold moves below even the χ_{c2}!

The fact that hidden charm states are significantly less extended than open charm states has been invoked as a signature for the quark-gluon plasma, where the correlation between colour sources and sinks is broken up over small distances. If in heavy ion collision a phase transition to the quark-gluon plasma is achieved, one expects a reduction in J/ψ production. The data are intriguing in this respect, yet not conclusive.

Charmonium production is investigated in relativistic heavy ion collisions[172], where the NA50 experiment [173, 174] using 1996 data (158 GeV per nucleon Pb beams on Pb target) provided circumstantial evidence for charmonium suppression, which may be explained by the onset of a quark-gluon plasma regime. They measure J/ψ production relative to Drell-Yan pair production. After accounting for conventional nuclear absorption, their data show evidence for a suddenly lower production, due to the attracting force between the $c\bar{c}$ quarks being screened by gluons, and fewer $c\bar{c}$ pairs hadronizing into J/ψ.

To conclude this section, we discuss a fascinating as much as hypothetical possibility uniquely provided by the study of charm particles in close contact to nuclear media, i.e., the formation of *supernuclei*. In complete analogy to what has been studied in great detail for several decades in Λ-hypernuclei [175], a charm quark produced at rest, or brought to rest, could interact with the nuclear matter, replace a light quark, and form a Λ_C baryon inside the nucleus. The Λ_C would then decay. This is an appealing process because the Λ_C does not need to obey the Pauli exclusion principle, and can occupy nuclear levels forbidden to the nucleons. The lifetime is also expected to differ from that for free Λ_C, and it would be possible to study both mesonic and nonmesonic decays. The only attempts carried out so fare have been in emulsions [176]. Supernuclei studies are foreseen at GSI with the PANDA experiment (see Sect.**13**·2.5).

6. – Spectroscopy and Lifetimes

The minimal information to describe a particle are its mass and its spin ([13]). Under the term 'mass' we can include also the width as the imaginary part of the mass. The width or lifetime of a particle actually characterizes its underlying dynamics in a way that the (real) mass cannot, namely whether they are strong (even if reduced), electromagnetic or weak, and in the latter case whether they are CKM suppressed or not. Beyond these general remarks the situation is different for hidden and open charm hadrons.

Hidden charm states $\bar{c}c$ are characterized by a Compton wave length $\sim 2/m_c \sim 1/3\ fm$, i.e. their extension is somewhat smaller than for light-flavour hadrons. They can decay electromagnetically and even strongly, the latter however with a very reduced width since it is order $\alpha_S^3(m_c)$ (for J/ψ). Powerful algorithms have been and are being developed to obtain very accurate predictions on charmonium spectroscopy from lattice QCD thus turning their experimental study into precision tests of QCD proper .

As already explained in Sect.**4**·5 HQS tells us that for *open* heavy flavour hadrons the two S wave configurations P_Q and V_Q become mass degenerate for $m_Q \to \infty$, while their mass exceeds m_Q by the scale $\sim \Lambda_{NPD}$ ([14]) as do the P wave configuration. I.e.:

$$m_c + \Lambda_{NPD} \sim M_D \simeq M_{D^*}\ ,\ M_{D^{**}} \sim M_D + \Lambda_{NPD} \tag{77}$$

The degree to which the hyperfine splitting $M_{D^*} - M_D$ is small compared to Λ_{NPD} is one measure for whether charm is a heavy flavour. It is, though not by a large factor:

$$M_{D^*} - M_D \sim\ 140\ \text{MeV} < M_{D^{**}} - \langle M_D \rangle \sim\ 480\ \text{MeV} \tag{78}$$

with $\langle M_D \rangle = \frac{1}{4}M_D + \frac{3}{4}M_{D^*}$ denoting the spin averaged meson mass. Also the simple scaling law of Eq.(12) is well satisfied:

$$M_B - M_D \simeq 3.41\ \text{GeV}\ vs.\ M_{\Lambda_b} - M_{\Lambda_c} \simeq 3.34\ \text{GeV} \tag{79}$$

There are further reasons to study the mass spectroscopy of charm resonances:

([13])This can be expressed for the mathematically minded reader by saying that elementary particles are defined by irreducible representations of the Poincare group; for those are labeled by the eigenvalues of two Casimir operators, which happen to be the mass and spin (or helicity for massless particles).
([14])It is usually denoted by $\bar{\Lambda}$.

- For a better understanding of the transition $B \to \ell\nu D^*$ that figures prominently in determinations of $V(cb)$ – and of $B \to \ell\nu X_c$ in general – one needs information on the mass and width of D^{**} and other higher resonances.

- More specifically, the SV sum rules [88] relate the basic HQP to the production of certain charm states in semileptonic B meson decays. E.g. [82]:

$$\begin{aligned} \frac{1}{2} &= -2\sum_n \left|\tau_{1/2}^{(n)}\right|^2 + \sum_m \left|\tau_{3/2}^{(m)}\right|^2 \\ \bar{\Lambda}(\mu) &= 2\left(\sum_n \epsilon_n \left|\tau_{1/2}^{(n)}\right|^2 + 2\sum_m \epsilon_m \left|\tau_{3/2}^{(m)}\right|^2\right) \\ \frac{\mu_\pi^2(\mu)}{3} &= \sum_n \epsilon_n^2 \left|\tau_{1/2}^{(n)}\right|^2 + 2\sum_m \epsilon_m^2 \left|\tau_{3/2}^{(m)}\right|^2 \\ \frac{\mu_G^2(\mu)}{3} &= -2\sum_n \epsilon_n^2 \left|\tau_{1/2}^{(n)}\right|^2 + 2\sum_m \epsilon_m^2 \left|\tau_{3/2}^{(m)}\right|^2 \ ; \end{aligned} \tag{80}$$

  here ϵ_k denotes the excitation energy of the final state D^k beyond the ground states D and D^* ($\epsilon_k = M_{D^k} - M_D$) while $\tau_{1/2}^{(n)}$ and $\tau_{3/2}^{(m)}$ denote the transition amplitudes for producing a state, where the light degrees of freedom carry angular momentum $j_q = 1/2$ or $3/2$, respectively [177]. Obviously, the masses of these charm resonances matter, as does their interpretation in terms of the quantum numbers $1/2$ or $3/2$.

- The mass splittings of baryonic charm resonances provide important cross checks for the evaluation of expectation values of four-quark operators that are highly relevant for predicting charm baryon lifetimes as discussed below.

Beyond classification there are other reasons for measuring total widths as precisely as possible. One needs them as an *engineering* input to translate branching ratios into partial widths. This is needed, for example, to infer the value of CKM parameters from semileptonic decays. On the *phenomenological* level a precise analysis of the D^0 lifetime is a prerequisite for studying $D^0 - \bar{D}^0$ oscillations. Finally on the *theoretical* side the lifetime ratios for the different charm hadrons provide the best, since most inclusive observables to probe hadrodynamics at the charm scale.

From the *raison d'etre for charm quarks*, namely to suppress strangeness changing neutral currents to the observed levels, one infers $m_c \leq 2$ GeV. The lifetime of charm *quarks* can be estimated by relating it to the muon lifetime and the number of colours and lepton flavours $\tau_c \sim \tau_\mu \cdot \left(\frac{m_\mu}{m_c}\right)^5 \cdot \frac{1}{N_C+2} \sim (\text{few}\, 10^{-13} s) \cdot \left(\frac{1.5\ \text{GeV}}{m_c}\right)^5$ with an obviously high sensitivity to the value of m_c.

These very simple estimates have turned out to be remarkably on target. Yet before we describe it, a few comments might be in order on the charm quark mass.

6$\cdot$1. *On the charm quark mass.* – *Within* a given quark *model* a quark mass has a clear meaning as a fixed parameter; however it depends on the specifics of the dynamical treatment adopted there, and therefore differs from model to model. Yet even more importantly the connection between such quark model *parameters* and fundamental quantities appearing in, say, the Lagrangian of the SM is rather tenuous. For example one can model single and double charm production in deep inelastic ν-nucleon scattering by charged and neutral currents with a parton model ansatz, where m_c plays of course a central role. Fitting data can yield a highly constrained value for m_c. Yet such a 'precise' value cannot be taken at face value to describe charm hadroproduction, let alone charmonium physics or charm decays. For that purpose one needs a proper field theoretical definition of the charm quark mass, which takes into account that the dynamical environments for these reactions differ in their perturbative as well as nonperturbative aspects. The resulting quantity has to be a 'running', i.e. scale dependent mass, where one has to specify its normalization scale; these issues have been discussed in Sect.4˙6.3.

The two areas where quark masses have been discussed with considerable care are charmonium spectroscopy and the weak decays of heavy flavour hadrons.

1. The first analysis was based on charmonium sum rules that approximate nonperturbative dynamics through including quark and gluon condensates in the OPE [178]. One finds for the $\overline{MS}$ mass

$$\overline{m}_c(m_c) = 1.25 \pm 0.10 \text{ GeV} \tag{81}$$

More recent analyses find fully consistent values:

$$\overline{m}_c(m_c) = \begin{cases} 1.19 \pm 0.11 \text{ GeV} & \text{Ref.}[179] \\ 1.30 \pm 0.03 \text{ GeV} & \text{Ref.}[180] \end{cases} \tag{82}$$

Lattice studies yield in the quenched approximation [181]

$$\overline{m}_c(m_c) = 1.301 \pm 0.034 \pm 0.13_{quench} \text{ GeV} . \tag{83}$$

2. The expansion for $m_b - m_c$ given in Eq.(21) yields

$$m_b - m_c = 3.50 \text{ GeV} + 40 \text{ MeV} \left(\frac{\mu_\pi^2 - 0.5 \text{ GeV}^2}{0.1 \text{ GeV}^2} \right) \pm 20 \text{ MeV} \tag{84}$$

Using the value for the b quark mass that has been extracted from $e^+e^- \to b\bar{b}$ near threshold by several groups [182]

$$m_b^{kin}(1 \text{ GeV}) = 4.57 \pm 0.08 \text{ GeV} \;\hat{=}\; \overline{m}_b(m_b) = 4.21 \pm 0.08 \text{ GeV} \tag{85}$$

which is in nice agreement with what one infers from a moment analysis of semileptonic B decays [183], and Eq.(84) one arrives at

$$\overline{m}_c(m_c) = 1.13 \pm 0.1 \text{ GeV} . \tag{86}$$

This value is completely consistent with what one obtains directly from the aforementioned moment analysis, namely

$$\overline{m}_c(m_c) = 1.14 \pm 0.1 \text{ GeV} \tag{87}$$

despite the caveats stated in Sect.4˙6.3 about the reliability of this expansion.

3. As will become clear from our discussion below, one cannot infer (yet) a reliable value for m_c from the charm lifetimes.

To summarize: a quite consistent picture has emerged, which supports treating charm as a heavy flavour.

6˙2. *Spectroscopy in the hidden charm sector*. – Charm entered reality in a most dramatic fashion through the discovery of hidden charm mesons and their striking properties, and our knowledge about them increased at breathtaking speed for some time due to very favourable experimental features.

Most of the spectroscopy results have come from e^+e^- storage rings, where $J^{PC} = 1^{--}$ states can be formed directly to lowest order. The three prominent states $J/\psi(3100)$, $\psi'(3700)$ and $\psi''(3770)$ have been well established for a long time as the ${}^3S_1$, 2^3S_1 and ${}^3D_1$ states, respectively, with the last one being broad since above $D\bar{D}$ production threshold. The nonvector states such ${}^3P_J$ (also referred to as χ_{cJ}) and ${}^1S_0$ can be reached by $E1$ and $M1$ transitions from them and thus be observed in two–step processes like $e^+e^- \to \psi' \to (c\bar{c})_{\chi_{cJ}} + \gamma$, see Fig. 13. This area of research pioneered by SPEAR and DORIS has experienced a welcome renaissance due to the operation of the Beijing Spectrometer (BES); in 2002 the BES collaboration has completed a four-month run which yielded 14 million $\psi(2S)$, to be added to the 4 million events previously collected.

A qualitatively new access to charmonium dynamics has been provided by low energy $p\bar{p}$ annihilation, since all J^{PC} quantum numbers then become accessible, in particular also ${}^1P_1$ and ${}^1D_2$ and ${}^3D_2$ states. The idea (pioneered by R704 at the ISR and carried forward by E760, E835 at FNAL) is to study the *formation* of charmonia states in the annihilation of antiprotons on a jet hydrogen target. E835 showed [123, 184] preliminary measurements of masses, widths and branching ratios of the three χ_{cJ} states with an unprecedented level of precision.

Finally a third actor has appeared: the B factories CLEO, BABAR and BELLE have such large statistics that one can study charmonia in $B \to [\bar{c}c]X$. This has been demonstrated quite dramatically by BELLE finding 5 σ and 3.5 σ signals for $\eta_c(1S)$ and $\eta_c(2S)$, respectively [185, 130]. The $\eta_c(2S)$ can boast of quite a saga [186]. Previous simultaneous observations of η_c and $\eta_c(2S)$ date back to conflicting measurements in the 1980's (DASP, Serpukhov, MARK II and Crystal Ball). While the η_c has become well established, the $\eta_c(2S)$ was not confirmed by either DELPHI or E835 in extensive searches ($30\,\mathrm{pb}^{-1}$ in the range 3666 to 3575 MeV). 2000 E835 searched with higher statistics for the $\eta_c(2S)$, with negative results. The $\eta_c(2S)$ was instead spotted in 2002 by BELLE at 3622 ± 12 MeV in B decays, and $3654 \pm 6\ MeV$ in the recoil spectrum of $J/\psi c\bar{c}$ events. Similarly frustrating is the search for the singlet P-state called h_c. Claimed by R704 at the ISR in 1986 and seen by E760 in 1993, the h_c has not, as yet, been confirmed by E760's successor E835 in its 2001 data set.

One expects [187] four charmonium states *below* $D\bar{D}$ threshold (and thus narrow), whose existence has not been established, namely $\eta_c'(2^1S_0)$, $h_c(1^1P_1)$, $\eta_{c2}(1^1D_2)$ and $\psi_2(1^3D_2)$; they can be identified in B decays.

The potential model ansatz pioneered by the Cornell group [188] was successful in describing the charmonium spectroscopy of Fig. 13). The factorization of nonperturbative and perturbative effects into a wave function and α_S corrections, respectively, as mentioned in Sect.**4**˙1.1 can be seen from the theoretical expression for the hadronic width of the J/ψ: $\Gamma(J/\psi \to \mathrm{hadrons}) = \frac{80(\pi^2-9)}{81\pi}\alpha_s(M_{J/\psi})^3(1 + 4.9\frac{\alpha_s}{\pi})|\Psi(0)|^2$.

The ratio of this to the leptonic width can be used to extract a value of α_s: $R_{\mu\mu} = \frac{J/\psi \to \mathrm{hadrons}}{J/\psi \to \mu^+\mu^-} = \frac{5(\pi^2-9)}{81\pi}\frac{\alpha_s(M_{J/\psi})^3}{\alpha_{em}}(1+10.3\frac{\alpha_s}{\pi})$. The experimental value $R_{\mu\mu} \approx 14.9$ leads to a reasonable result: $\alpha_s(M_{J/\psi}) \approx 0.2$. Relativistic effects of order v^2/c^2 can also be included. Yet this method of extracting α_s is not as theoretically sound as others. The first order radiative correction is as large as the lowest order correction calling into question the validity of the perturbative expansion. Furthermore, the expressions for the various widths ultimately depend on the expression that is chosen to describe the quark-antiquark potential, which is based upon phenomenological aspects of QCD rather than rigorously derived from it. By taking ratios, in which the dependence on the wavefunctions vanish, this source of uncertainty can be reduced. It still remains unclear how valid the factorization assumption is for the charmonium system in which m_c is only moderately larger than typical hadronic scales.

Radiative transitions between charmomium states can similarly be described. The radius of the bound state is typically much smaller than the wavelength of the emitted radiation so a multipole expansion is expected to converge quite rapidly. Electric dipole (E1) transitions are responsible for $\Delta S = 0, \Delta L = 1$ processes. The rate for transitions between S- and P- wave states is:

$$\Gamma_\gamma(S \leftrightarrow P) = \frac{4}{9}\left(\frac{2J_f+1}{2J_i+1}\right) Q^2\alpha|E_{if}|^2 E_\gamma^3. \tag{88}$$

Here $J_{f[i]}$ denotes the total angular momentum of the final[initial] state, $Q = 2/3$ is the charge of the charmed quark, E_γ is the photon energy and E_{if} is the matrix element of the transition dipole operator: $E_{if} = \int_0^\infty r^2\Psi_i(r)\, r\, \Psi_f(r)$. Since this matrix element is more sensitive to the exact *shape* of the wavefunction unlike $|\Psi(0)|^2$ that appeared previously, considerable differences emerge among theoretical predictions. Even so, there is reasonable agreement with experiment [129]. Magnetic dipole (M1) transitions are responsible for $\Delta S = 1, \Delta L = 0$ processes and are suppressed by E_γ/m_c with respect to the E1 transitions. The transition rate between spin 0 and 1 S- wave states is given by the following expression:

$$\Gamma_\gamma(^3S_1 \leftrightarrow ^1S_0) = \frac{16}{3}(2J_f+1)\left(\frac{Q^2}{2m_c}\right)\alpha|M_{if}|^2 E_\gamma^3. \tag{89}$$

Here the magnetic dipole moment is the expectation value of the zeroth order spherical Bessel function: $|M_{if}| = \int_0^\infty r^2\Psi_i(r)\, j_0(\frac{1}{2}E_\gamma r)\, \Psi_f(r)$. Since these matrix elements depend quite sensitively on details of the wave functions, it is not surprising that the agreement between theory and experiment for M1 transitions is rather poor.

The lattice community is able now to treat charmonium physics with three flavours of dynamical quarks; from the spin-averaged $1P - 1S$ and $2S - 1S$ splittings one infers for the strong coupling $\alpha_S^{\overline{MS}}(M_Z) = 0.119 \pm 0.004$ [189].

Hadronic transitions like $\psi' \to \psi\pi\pi$ are also treated using a multipole expansion to describe the gluonic radiation. An added complication is the hadronization of the emitted gluonic radiation. By introducing a chiral Lagrangian to describe the effective low energy behaviour of the hadronic state, a semi-quantitative analysis can be carried out for these transitions.

The transition $J\psi \to \gamma X$ driven by $J\psi \to \gamma gg$ provides a gluonic origin for the final state X. Accordingly states with a particular affinity to gluons should figure

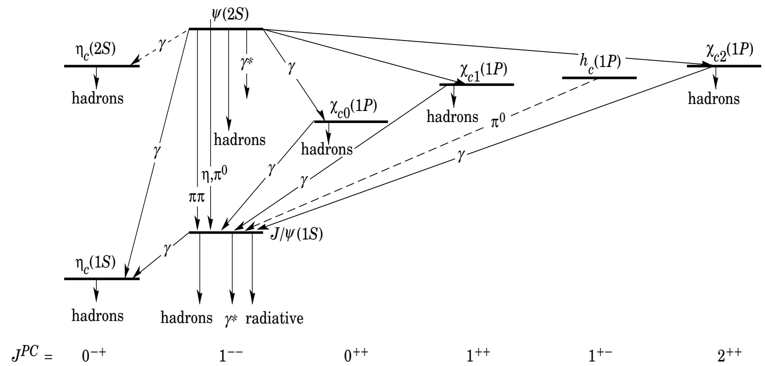

Fig. 13. – *Chart of charmonium states [131].*

prominently in X. Narrow states would show up as mass peaks in the γ recoil spectrum; no prominent signal has been found yet. One can search for them also in exclusive final states, as discussed in Sect.**9**.

An update overview of the experimental panorama can be found in [124]. Breaking news in summer 2003 was the preliminary result by BELLE on the observation of a $J/\psi\pi^+\pi^-$ state in decay $B^+ \to K^+(jp\pi^+\pi^-)$. BELLE finds a clear (8.6σ) signal at 3871.8 ± 0.7 ± 0.4 MeV, which is suggestive of a DD^* molecule.

6˙3. *Spectroscopy in the* $C \neq 0$ *sector*. – Adding charm as the fourth quark leads to a very rich spectroscopy. There are six $C = 1$ pseudoscalar states (plus one $\bar{c}c$ state already discussed) in addition to the familiar $SU(3)$ meson nonet, namely $D^\pm$, $D_s^\pm$ and $D^0/\bar{D}^0$; likewise for the vector mesons with $D^{*\pm}$, $D_s^{*\pm}$ and $D^{*0}/\bar{D}^{*0}$. For baryons even more facets emerge, as described later.

These states can be fitted into $SU(4)$ multiplets. Yet $SU(4)$ breaking driven by $m_c > \Lambda_{NPD} \gg m_s$ is much larger than $SU(3)$ breaking (15). Heavy quark symmetry provides a much more useful classification scheme. As explained in Sect.(**4**˙5) for heavy flavour hadrons H_Q the spin of Q – $\mathbf{S_Q}$ – decouples from the light quark degrees of freedom, and $\mathbf{j_q} \equiv \mathbf{s_q} + \mathbf{L}$ and $\mathbf{S_Q}$ become *separately* conserved quantum numbers. A meson [baryon] can then be characterised by the spin of the light antiquark [diquark] and the orbital angular momentum.

While the masses of the groundstates D/D^* and D_s/D_s^* have been known with 1 MeV precision for a decade now, the experimental information available on other states is still unsatisfactory. In this section we shall discuss open problems and very recent surprises.

6˙3.1. D^* width. Measuring $\Gamma(D^{*+})$ represents an experimental challenge: such widths are predicted in the range of tens or hundreds keV, and must be experimentally deconvoluted to the experimental resolution of detectors. CLEO has presented the first measurement [190, 191] based on the sequence $D^{*+} \to \pi^+ D^0, D^0 \to K^-\pi^+$ from a 9 fb^{-1} sample of e^+e^- data collected with the CLEO II.V detector. $\Gamma(D^{*+})$

(15)Another way to put it is to say that charm mesons – unlike pions and kaons – cannot be viewed as Goldstone bosons.

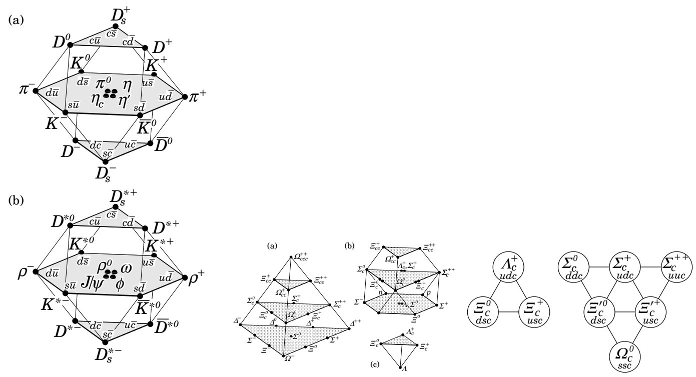

Fig. 14. – *Chart of ground state charmed meson and baryon multiplets [131].*

is controlled by the strong coupling constant, since the electromagnetic contribution $D^{*+} \to \gamma D^+$ produces a very small branching ratio, and can thus be neglected.

Based on an complex analysis of 11,000 reconstructed D^* decays CLEO finds $\Gamma(D^{*+}) = 96 \pm 4 \pm 22\,\mathrm{keV}$. Hence they infer for the $D^*D\pi$ coupling: $g_{D^*D\pi} = 10 \pm 3.5$. Light-cone sum rules have been employed to obtain the prediction [192] $g_{D^*D\pi} = 17.9 \pm 0.3 \pm 1.9$, where the value 13.5 is viewed as a rather firm upper bound [193].

The B factories with their large charm samples should be able to check CLEO's results, once detector simulation will reach the level of accuracy necessary to tame the severe systematic uncertainty.

6˙3.2. Charm mesons - $L = 1$ excited states. For each of the $c\bar{u}$, $c\bar{d}$ and $c\bar{s}$ systems four P-wave and two $n = 2$ radial excitations have been studied. There are four $L = 1$ states, namely two with $j_q = 1/2$ and total spin $J = 0, 1$ and two with $j_q = 3/2$ and $J = 1, 2$. These four states are named respectively D_0^*, $D_1(j_q = 1/2)$, $D_1(j_q = 3/2)$ and D_2^* (Fig.15). Parity and angular momentum conservation force the $(j_q = 1/2)$ states to decay to the ground states via S-wave transitions (broad width), while $(j_q = 3/2)$ states decay via D-wave (narrow width). To be more specific, for the $1/2$ one predicts widths of ~ 100 MeV and for the $3/2$ of about ~ 10 MeV with the exception of the $D_{s1}(j_q = 3/2)(2536)$ which is kinematically forced to a ~ 1 MeV width.

All six $L = 1, j = 3/2$ *narrow* states are well established, with precisions on masses at the 1 MeV level and on widths at the few MeV level. This is due to the fact that excited D states are abundantly produced both at FT experiments, in e^+e^- continuum production, in B decays and at the Z^0 [137]. Common analysis techniques are the selection of a clean sample of D meson candidates typically via a candidate-driven algorithm (see Sect.**6**˙4.1), a cut on the $D\pi$ mass to reject D^* compatible combinations, and the pairing of the D candidate with one (two) soft pion (pions) in the primary vertex to form the D^{**} candidate. A selection on the helicity angle is effective in selecting L=1 states from background, as well as in selecting different states in the same final channel.

At e^+e^- one also invokes the mass constraint of the fixed center of mass energy and kinematical helicity cuts which exploit the constraint of the parent B mass. Dalitz plot and partial wave analyses have also been presented [194].

A review of the data from different experiments can be found in [195]. A common

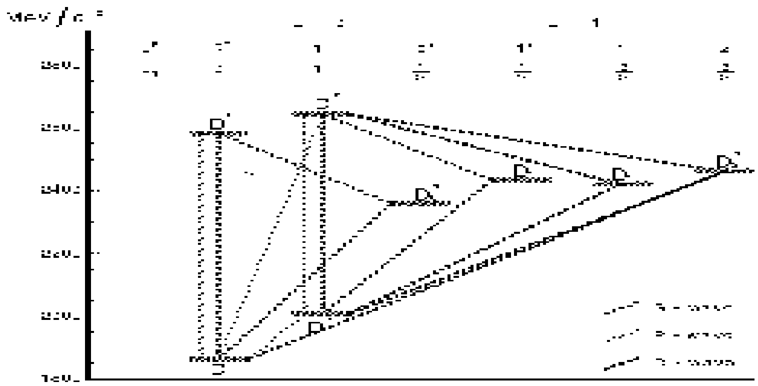

Fig. 15. – *Masses and transitions predicted for the excited D meson states (pre-Spring 2003).*

feature is a prominent peak in the 2460 MeV region, escorted by satellite peaks at one π^0 mass below, due to feeddown from decays where the π^0 escapes detection, such as the decay chain $D_2^{*+} \to D^0\pi^+$ with the D^0 wrongly assigned to $D^{*0} \to D^0\pi^0$ decay.

Generally speaking, narrow L=1 signals at FT have higher statistics but sit on a more prominent background, due to the larger combinatorics coming from larger primary multiplicities. On the other hand photoproduction signals with their lower primary multiplicity have less combinatorics than hadroproduction.

The status of the *broad* L=1 states is much less clear, and the assignments of the quantum numbers are largely based on theory expectations for their masses and widths. In 1998 CLEO [196] showed evidence for the $D_1(j_q = 1/2)$ broad state. The authors of Ref.[197] propose an alternative interpretation of this state as the axial chiral partner. Such an alternative interpretation is actually supported by the SV and spin sum rules describing semileptonic B decays, see Eq.(80): they strongly suggest that the D^{**} states around $2.4 - 2.6\mathrm{GeV}$ have to come mainly from a $3/2$ state – or they do not represent a P wave configuration [198].

Tab.III gives a summary of our knowledge of excited D mesons as it appeared in early 2003. At that time it seemed all one needed was to fill in a few gaps.

A major confirmation for HQS would be the definite observation of the missing $(L = 1, j_q = 1/2, J^P = 0^+)$ D_0^* and $(L = 1, j_q = 1/2, J^P = 1^+)$ D_1 broad states. More recently, FOCUS[201] and BELLE[194] showed new preliminary results and evidence for D_0^*. Errors on both masses and widths are still very large and we expect better measurements from the B-factories due to their larger data sets.

The overall information on L=1 $c\bar{s}$ states is unsatisfactory anyway. The $D_{s1}(j_q = 3/2)$ has been seen in D^*K^0 final state, and not in D^+K^0 or in D^0K^+. The D_{s12} (called $D_{sJ}(2573)$ by PDG) has been seen in D^0K^+ and recently in D^+K^0[202]. Furthermore no candidate for the $D_s(1/2)$ doublet had been seen yet. Since their masses were firmly expected to be about 80 MeV larger than for the corresponding nonstrange states, they would have enough phase space for the decays into $D^{(*)}K$ final states leading to large widths.

An open question remains the first evidence[204] seen by DELPHI of a charm radial excitation $D^{*\prime+}$ in the $D^{*+}\pi^-\pi^+$ final state (called $D*(2640)^\pm$ by PDG); it has not been confirmed by any experiment (OPAL[205], CLEO[206], ZEUS[207]), and questioned by theory predictions [208].

TABLE III. – *Winter '02/'03 status of (L=1, n=1) and (L=0, n=2)* $c\bar{q}$ *and* $c\bar{s}$ *mesons (MeV). Statistical and systematical errors added in quadrature. Experimental results not included in PDG [131] are from BELLE [199], CLEO [196], DELPHI [200], FOCUS [201] [202]. Theory predictions from [203].*

j_q J^P L, n	$1/2$ 0^+ $1, 1$	$1/2$ 1^+ $1, 1$	$3/2$ 1^+ $1, 1$	$3/2$ 2^+ $1, 1$	$1/2$ 0^- $0, 2$	$1/2$ 1^- $0, 2$
Decay Mode	D_0^* $D\pi$	D_1 $D^*\pi$	$D_1(2420)$ $D^*\pi$	$D_2^*(2460)$ $D\pi, D^*\pi$	D'	$D^{*\prime}$ $D^*\pi\pi$
			Mass (MeV)			
PDG **0**			2422 ± 2	2459 ± 2		
PDG ±			2427 ± 5	2459 ± 4		2637 ± 7
FOCUS **0**	∼ 2420			2463 ±**2**		
FOCUS ±	∼ 2420			2468 ±**2**		
BELLE **0**	2308 ± 36	2427 ± 36	2421 ± 2	2461 ± 4		
DELPHI ±		2470 ± 58				
CLEO **0**		2461 ± 51				
Theory	2400	2490	2440	2500	2580	2640
			Width (MeV)			
PDG **0**			19 ± 4	23 ± 5		
PDG ±			28 ± 8	25 ± 7		< 15
FOCUS **0**	∼ 185			30 ± 4		
FOCUS ±	∼ 185			29 ± 4		
BELLE **0**	276 ± 66	384 ± 114	24 ± 5	46 ± 8		
DELPHI ±		160 ± 77				
CLEO **0**		290 ± 100				
Theory	>170	>250	20-40	20-40		40-200
Decay Mode	D_{s0}^*	D_{s1}	$D_{s1}(2536)$ D^*K	$D_{sJ}^*(2573)$ DK	D_s'	$D_s^{*\prime}$
			Mass (MeV)			
PDG ±			2535.3 ± 0.6	2572.4 ± 1.5		
FOCUS ±			2535.1 ± 0.3	2567.3 ± 1.4		
Theory	2480	2570	2530	2590	2670	2730
			Width (MeV)			
PDG ±.			<2.3 90 % cl	15 ± 5		
FOCUS ±			1.6 ± 1.0	28 ± 5		
Theory			**< 1**	**10 − 20**		

6˙3.3. Charm mesons - New $L = 1$ D_s states. Analyses presented by BABAR [209] and CLEO [210] in the spring of 2003 are challenging the whole picture.

1. BABAR reported finding a narrow resonance $D_{sJ}^*(2317)$ with $D_{sJ}^*(2317) \rightarrow D_s^+\pi^0$ in $90 fb^{-1}$ of data. With the observed width consistent with the experimental resolution, the *intrinsic* width has to be below 10 MeV. This discovery has been confirmed by CLEO.

2. CLEO with $13 fb^{-1}$ has observed another similarly narrow state at a mass 2.46 GeV, for which BABAR had found evidence before: $D_{sJ}^*(2463) \rightarrow D_s^{*+}\pi^0$.

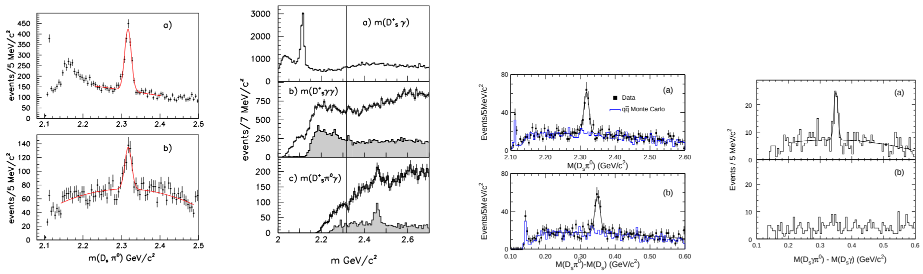

Fig. 16. – *New $D^*_{s0+}(2317)$ and $D^*_{s1+}(2463)$ states observed by BABAR (a,b) [209] and CLEO (c,d) [210].*

3. BELLE[211, 212] out of $86.9 fb^{-1}$ dataset finds evidence for $D^*_{sJ}(2463) \to D^+_s\gamma$ decay, measures relative branching ratio to $D^*_s\pi^0$, and determines $J^P = 1^+$ assignment for $D^*_{sJ}(2463)$.

It seems natural to interpret $D^*_{sJ}(2317)$ and $D^*_{sJ}(2463)$ as 0^+ and 1^+ states, respectively. The decay distributions are consistent with such assignments, yet do not establish them. They together with the mass values would explain the narrow widths: for $D^*_{s1+}(2463) \to DK$ is forbidden by parity, $D^*_{s0+}(2317) \to DK$ and $D^*_{s1+}(2463) \to DK^*$ by kinematics and $D^*_{s0+}(2317) \to D^+_s\pi^0$ and $D^*_{s1+}(2463) \to D^{*+}_s\pi^0$ are isospin violating transition and thus suppressed. Also $D^*_{s0+}(2317) \to D^+_s\gamma$ is forbidden.

There are three puzzling aspects to these states:

- Why have no other decay modes been seen ? In particular CLEO places a low upper bound

$$BR(D^*_{s0+}(2317) \to D^{*+}_s\gamma) < 0.078 \quad 90\%\ C.L. \tag{90}$$

  Why is it not more prominent, when $D^*_{s0+}(2317) \to D_s\pi^0$ is isospin violating ?

- Why are their masses so much below predictions ? One should note that a deficit of ~ 160 and ~ 100 MeV is quite significant on the scale of $M(D^*_{sJ}) - M(D)$. Why is the mass splitting to the previously found narrow states $D_{s1}(2536)$ and $D_{sJ}(2573)$ so much larger than anticipated ?

- A related mystery is the following: where are the corresponding *non*-strange charm resonances ? They should be lighter, not heavier than $D^*_{s0+}(2317)$ and $D^*_{s1+}(2463)$.

Potential model results have been re-analyzed, lattice QCD is being consulted [213] and more exotic scenarios like a DK molecule below threshold [214] and other four-quark interpretations have been put forward. The latter are hard pressed to explain the narrow width in the first place, since the transition would not have to be isospin violating; yet in addition CDF sees no evidence for $D^*_{sJ} \to D_s\pi^\pm$ [215], although the preliminary CDF data of well known L=1 states (such as the D^*_2) reported so far show rather modest signal-to-noise ratios.

Maybe the most intriguing explanation, since it would represent a new paradigm for the implementation of chiral invariance, is the suggestion made in Ref.[216] to combine heavy quark symmetry and chiral invariance in a novel way: the latter is realized through parity doublets pairing (D_s, D_s^*) with $(D_{s0^+}^*, D_{s1^+}^*)$. Chiral dynamics induces a mass splitting ΔM between the heavy quark doublets: invoking a Goldberger-Treiman relation [217] the authors of Ref.[216] estimate

$$\Delta M = M(D_{s0^+}^*) - M(D_s) \simeq M(D_{s1^+}^*) - M(D_s^*) \sim m_N/3 \tag{91}$$

in agreement with the experimental findings. They also find that the radiative width is reduced: $BR(D_{s0^+}^*(2317) \to D_s^{*+}\gamma) \sim 0.08$, i.e. right about CLEO's upper limit. The same chiral symmetry and ΔM apply to B mesons and double charm baryons $[ccq]$.

Last – but certainly not least – the same chiral splitting should arise for nonstrange charm mesons leading to

$$M(D^{\pm}(0^+) \simeq 2217 \text{ MeV} \;,\; M(D^{\pm}(1^+) \simeq 2358 \text{ MeV} \tag{92}$$

$$M(D^0(0^+) \simeq 2212 \text{ MeV} \;,\; M(D^0(1^+) \simeq 2355 \text{ MeV} \tag{93}$$

Of course, these predictions could be modified significantly by order Λ_{NPD}/m_c corrections. These states can decay strongly into $D\pi$ and $D^*\pi$ and thus would be broad. This makes it difficult to distinguish the decays of these resonances from phase space distributions of $D\pi$ and $D^*\pi$. It is very important to search for them; finding them would neccessitate a new interpretation for the previously listed D_0^* and D_1. B factories analysing the final states of B decays have an advantage in such analyses.

These remarks indicate that a fundamental re-evaluation of strange as well as nonstrange charm hadrons might be in store. This would have an important impact not only on charm spectroscopy, but also on the aforementioned sum rules describing semileptonic B decays, Eq.(80), and on the latter's interpretation.

6˙3.4. $C = 1$ baryons. In the framework of SU(4) (Fig.14) we expect nine ground state cqq $J^P = 1/2^+$ baryons (all of them detected after the 1999 2̧ observation of Ξ_c') and six cqq $J^P = 3/2^+$ baryons (with only the Ω_c^{*0} remaining undetected, after the observation of Σ_c^{*+} by 2̧ in 2002 [218]); for a concise and for a more detailed review see [131] and [219], respectively. The Ω_c^{*0} is expected to decay via the experimentally challenging channel $\Omega_c^{*0} \to \gamma\Omega_c^0$) ([16]).

Only a few of the orbitally excited P-wave $(L = 1)$ baryons have been observed. The first doublet $\Lambda_{c1}(2593)$ $(1/2^-)$ and $\Lambda_{c1}^*(2625)$ $(3/2^-)$ was observed several years ago by 2̧, ARGUS, and E687. In 1999 2̧ presented evidence[221] for the charmed-strange baryon analogous to $\Lambda_{c1}^*(2625)$, called $\Xi_{c1}^*(2815)$, in its decay to $\Xi_c\pi\pi$ via an intermediate Ξ_c^* state. The presence of an intermediate Ξ_c^* instead of Ξ_c' supports an $3/2^-$ assignment, while HQS explicitly forbids a direct transition to the $\Xi_c\pi$ ground state due to angular momentum and parity conservation. More recently, CLEO has reported on the observation of both a broad and a narrow state in the $\Lambda_c^+\pi^+\pi^-$ channel, identified as the Σ_{c1} and the Λ_{c0} $1/2^-$ [222].

([16])We adopt for excited baryon states the nomenclature in [220]. Thus, members of 3/2 multiplets are given a $(*)$, the subscript is the orbital light diquark momentum **L**, and $(')$ indicates symmetric quark wavefunctions $c\{q_1q_2\}$ with respect to interchange of light quarks, opposed to antisymmetric wavefunctions $c[q_1q_2]$.

Mass splittings within isospin multiplets are caused by QCD corrections due to $m_d > m_u$ and by electromagnetic effects. Interest in this subject has attracted new interest recently: while for all well-measured isodoublets one increases the baryon mass by replacing a u–quark with a d–quark, the opposite happens in the case of the poorly measured $\Sigma_c^{++}(cuu) - \Sigma_c^0(cdd)$ isosplit [223, 224]. Also, as discussed below, the SELEX candidates for the two isodoublet $C = 2$ baryons Ξ_{cc} show larger than expected mass splitting. FOCUS' new number $M(\Sigma_c^{++}) - M(\Sigma_c^0) = -0.03 \pm 0.28 \pm 0.11$ **MeV** [225] indicates a trend towards a smaller splitting, although still statistically consistent with both E791 one's of $0.38 \pm 0.40 \pm 0.15$ **MeV** as well as PDG02's 0.35 ± 0.18 **MeV**.

There is another use of baryonic spectroscopy in a somewhat unexpected quarter: as explained in Sect. **4**˙6.2 and discussed more specifically below, the weak decay widths of charm baryons can be expressed through the expectation values of local operators. The numerically leading contributions are due to four-quark operators. At present we can compute those only with the help of wave functions obtained in quark models. Yet those wave functions allow us also to calculate baryon mass splittings. One can then relate the needed baryonic expectation values to static observables like $M(\Sigma_c) - M(\Lambda_c)$ or $M(\Sigma_c^*) - M(\Sigma_c)$ [226]. Uraltsev, Phys. Lett. B 376 (1996) 303

Finally, measurements of the Σ_c, Σ_c^{++} natural widths were presented by CLEO[227] and FOCUS[228]. They both agree on a few MeV width, but the level of precision is not enough yet to discriminate among theoretical models.

6˙3.5. $C \geq 2$ baryons. Combining the large charm production rates in hadronic collisions with state-of-the-art microvertex detectors that allow to *trigger* on charm decays opens the window to a more exotic class of hadrons, namely baryons containing two (or even three) charm quarks. There is an $SU(3)$ triplet of such states: $\Xi_{cc}^{++} = [ccu]$, $\Xi_{cc}^{+} = [ccd]$ and $\Omega_{cc}^{+} = [ccs]$ (plus the superheavy $\Omega_{ccc} = [ccc]$).

Like for $C = 1$ baryons, one can employ quark models of various stripes to predict their masses. However there are some qualitative differences: as stated before, in the Λ_c and Ξ_c bound states a light spin-zero diquark surrounds the heavy c quark, whose spin is decoupled to leading order in $1/m_c$ due to QCD's HQS. In Ξ_{cc} and Ω_{cc} on the other hand the light degrees of freedom carry spin 1/2 implying degeneracy among several ground states to leading order in $1/m_c$. It has been suggested to model $C = 2$ baryons as a heavy-light system consisting of a cc-diquark and a light quark. Accordingly there will be two kinds of mass spectra, namely due to excitations of the light quark and of the cc 'core'. Based on such an ansatz the following predictions on the masses were made more than ten years ago [229]:

$$M(\Xi_{cc}) \simeq 3.61 \text{ GeV}, \; M(\Omega_{cc}) \simeq 3.70 \text{ GeV}, \; M(\Omega_{ccc}) \simeq 4.80 \text{ GeV}, \tag{94}$$

These numbers still reflect today's theoretical expectations [230, 231]. Ref.[230] lists various models; their predictions are in the range 3.48 - 3.74 GeV for $M(\Xi_{cc})$ and 3.59 - 3.89 GeV for $M(\Omega_{cc})$.

In 2002, the SELEX Collaboration claimed the observation [154] of the Ξ_{cc}^{+} (ccd) through its decay mode to $\Lambda_c^+ K^- \pi^+$. In this experiment charmed baryons are produced by a 600 GeV charged hyperon beam (Tab.I). Charged tracks are detected and reconstructed by a silicon vertex detector, coupled to a forward magnetic spectrometer. After analyses cuts, a sample of 1630 fully reconstructed $\Lambda_c \to pK^-\pi^+$ events is selected. Double charm baryon candidates are searched for in events which assign a Λ_c to a $K^-\pi^+$ secondary vertex. SELEX found a 15.9 events signal over an expected

TABLE IV. – *Comparison of SELEX and FOCUS results on double-charm baryons (from [235]). Limits are at the 90% cl.*

	$\Xi_{cc}^{+} \to \Lambda_c^+ K^- \pi^+$		$\Xi_{cc}^{++} \to \Lambda_c^+ K^- \pi^+ \pi^+$	
	FOCUS	SELEX	FOCUS	SELEX
Ξ_{cc} Candidate events	<2.21 %	15.8	<2.21 %	8
Reconstructed Λ_c	19 500	1 650	19 500	1 650
Ξ_{cc}/Λ_c Relative Efficiency	5%	10%	13%	5%
Relative Yield Ξ_{cc}/Λ_c	<0.23 %	9.6 %	<0.09%	9.7%
Relative Production Ξ_{cc}/Λ_c	SELEX/FOCUS >42		SELEX/FOCUS >111	

background of $\mathbf{6.1 \pm 0.5}$ events, that they translate to a $\mathbf{6.3\sigma}$ significance, at a mass of $\mathbf{3519 \pm 1}$ MeV, and a width of $\mathbf{3}$ MeV compatible to the experimental resolution. SELEX has also shown [155] preliminary results on an excess of 9 events (over an expected background of 1), at $\mathbf{3460}$ MeV, which they translate to a $\mathbf{7.9\sigma}$ significance in the $\Lambda_c^+ K^- \pi^+ \pi^+$ final state. This was interpreted as evidence for the isodoublet partner, the Ξ_{cc}^{++}. Alternate statistical approaches have been proposed [232] which treat differently the background fluctuations, and provide a signal significance for the SELEX candidates of about 3-4 σ. The very serious trouble with this interpretation is the apparent $\sim \mathbf{60}$ MeV mass splitting between the isospin partners Ξ_{cc}^{+} and Ξ_{cc}^{++}: it causes a major headache for theorists, in particular when one keeps in mind that the proper yardstick for comparing this mass splitting to is *not* the total mass, but the much smaller binding energy of a few hundred MeV. Furthermore as discussed above there is no evidence for such an exotic effect in single charm baryons.

In May 2003 a very intriguing new twist has been added to the story: based on further studies, in particular of the angular distributions of the decay products, SELEX [233] now concludes that there are actually *four* $C = 2$ baryons, namely two $J = 1/2^+$ states with $L = 0$, namely $\Xi_{cc}^{+}(3443)$ and $\Xi_{cc}^{++}(3460)$ decaying isotropically into $\Lambda_c^+ K^- \pi^+$ and $\Lambda_c^+ K^- \pi^+ \pi^+$, respectively, and a heavier pair $\Xi_{cc}^{+}(3520)$ and $\Xi_{cc}^{++}(3541)$ presumably with $J = 1/2^-$ and $L = 1$ decaying nonisotropically into the same final states. The problem with the isospin mass splittings mentioned above has been alleviated now, since the two isodoublets have mass splittings of 17 and 21 MeV, although this is still anomalously large. The heavier doublet could be understood as an excitation of the cc core; preliminary estimates yield for this excitation energy a range of about 70 to 200 MeV and for the isospin mass splitting about 6 MeV [234].

The SELEX evidence has not been confirmed by the photoproduction experiment FOCUS[235]. In FOCUS, following analysis techniques similar to SELEX, a sample of 19,444 $\Lambda_c \to pK^-\pi^+$ events is selected, with neither Ξ_{cc}^{+} nor Ξ_{cc}^{++} candidates. FOCUS concludes (Tab. IV) that this implies a production difference between double charm baryons and Λ_c baryons of $> \mathbf{42}$ for the Ξ_{cc}^{+}, and $> \mathbf{111}$ for the Ξ_{cc}^{++}.

SELEX' interpretation of its data creates very considerable headaches for theory. The observed isospin splitting is larger than expected, yet the main problem concerns the reported lifetimes. None of the four states exhibit a finite lifetime, and the apparent upper bound is about 33 fs. One should note that the upper doublet can decay electromagnetically through $M2$ transitions, for which one can estimate a lifetime of about 1 fs [234]. Therefore if they indeed represent $C = 2$ baryons, their weak lifetimes have to be indeed on the femtosecond level. The poses a very serious challenge to theory, as will be discussed in Sect. **6**.5.

Not unlike BELLE's observation of $J/\psi c\bar{c}$ events [130] SELEX' findings point to unexpectedly large $c\bar{c}c\bar{c}$ production. Yet they present another puzzle as well, namely why two charm quarks each produced in a hard collision presumably incoherently end up in the same $C = 2$ baryon. One would expect that much more often than not they hadronize separately leading to sizeable DD and $\bar{D}\bar{D}$ production.

6.3.6. Production of charm resonances. A very naive estimate for the relative production of different charm states is based on 'spin counting', i.e. assigning equal probability to the production of each spin component of a given charm hadron. Consider the simplest case of vector (V) vs. pseudoscalar (P) production like D^* vs. D. Spin counting suggests a ratio $r = V/(V + P) = 3/4$. There is no justification (beyond its simplicity) for such an ansatz and it fails already for continuum D_s* vs. D_s production in e^+e^- annihilation, where CLEO finds [236] $r = 0.44 \pm 0.04$ which is less than even equal production for D_s^* and D_s. In semileptonic B decays on the other hand one finds $\Gamma(B \to \ell\nu D^*) \sim 2\Gamma(B \to \ell\nu D^)$.

With the increase in data sets, and the simultaneous refinement in the level of sophistication in understanding the sources of systematics, measurements of relative production yields of L=1 states in B decays have become available.

There is a double motivation for understanding charm production in B decays – in particular semileptonic B decays – that goes well beyond testing hadronization models per se.

- The CKM parameter $|V(cb)|$ is best extracted from $B \to \ell\nu X_c$, $B \to \ell\nu D^*$ and $B \to \ell\nu D$. To fully understand detection efficiencies, feed-downs etc. one has to know the quantum numbers of the charm final states produced.

- The sum rules stated in Eq.(80) relate the basic heavy quark parameters to moments of the production rates for $j_q = 1/2$ and $j_q = 3/2$ charm resonances. Those heavy quark parameters form an essential input for the theoretical treatment of semileptonic B decays, beauty lifetimes etc. and also provide a valuable quantitative test ground for lattice QCD.

It has been known for a long time that charm production in B decays is characterized by the dominance of broad over narrow states. Since the former are usually identified with $j_q = 1/2$ and the latter with $j_q = 3/2$ this is again in clear contrast to spin counting.

The BELLE paper [194] discusses the issue. Their measurements *show that narrow resonances compose 33±4% of the* $B \to (D\pi)\pi$ *decays, and 66±7% of* $B \to (D^*\pi)\pi$ *decays.* This trend is consistent with the excess of broad states component in semileptonic B decays $B \to D^{**}\ell\nu$ at LEP [237]

One should keep in mind though that these assignments $1/2$ vs. $3/2$ are typically inferred from theory rather than the data. It was already mentioned that the sum rules of Eq.(80) cast serious doubts on some of these assignments.

Only a few theoretical papers have addressed the issue of relative production rates of L=1 states, from B decays. Neubert[238] predicts the ratio J=2 / J=1(j=3/2) $R = B(B \to D_2^{*0}\pi^-)/B(B \to D_1^0(j = 3/2)\pi^-) \sim 0.35$ while BELLE measures [194] R=0.89±0.14; both numbers actually contradict spin counting.

The Orsay group[239] has developed a model for describing charm production in exclusive semileptonic $\boldsymbol{B}$ decays. The model predicts dominance of narrow states in B semileptonic decays. It can claim reasonable reliability, since it implements the SV and spin sum rules referred to in Eq.(80).

As for L=1 D mesons not produced from B decays, there are no theoretical predictions. The only experimental evidences here for broad states come from FOCUS [201], where the $\boldsymbol{D_0^*}$ is observed relative production yields with respect to the $\boldsymbol{D_2^*}$ of about 3:1.

6˙4. *Weak lifetimes and semileptonic branching ratios of* $\boldsymbol{C=1}$ *hadrons.* – The decay rate of a charm *quark* provides only an order of magnitude estimate for the lifetimes of the weakly decaying charm *hadrons*; their individual lifetimes could differ quite substantially.

To illustrate this point, let us look at strange quarks and hadrons. Very naively one would expect for a strange quark of mass 150 MeV a (Cabibbo suppressed) lifetime of roughly $\boldsymbol{10^{-6}\ s}$. Not surprisingly, such a guestimate is considerably off the mark. Furthermore strange hadrons exhibit huge lifetime differences which are fed by two sources:

$$\frac{\tau(K_L)}{\tau(K_S)} \sim 600\ ,\ \frac{\tau(K^+)}{\tau(K_S)} \sim \mathcal{O}(100) \sim \frac{\tau(K^+)}{\tau(\Lambda)} \tag{95}$$

The first ratio is understood as the combination of (approximate) CP invariance in kaon decays, which forbids $\boldsymbol{K_L}$ to decay into two pions, with the accidental fact that the $\boldsymbol{K_L}$ mass is barely above the three-pion threshold. Such an effect cannot induce a significant lifetime difference among charm hadrons. For the second and third ratio in Eq.(95) a name has been coined – the $\boldsymbol{\Delta I = 1/2}$ rule – yet no conclusive dynamical explanation given. Since all weakly decaying strange baryons benefit from $\boldsymbol{\Delta I = 1/2}$ transitions, no large lifetime differences among them arise. The impact of the $\boldsymbol{\Delta I = 1/2}$ rule is seen directly in the ratio of the two major modes of $\boldsymbol{\Lambda}$ decays: $\boldsymbol{\Gamma(\Lambda \to n\pi^0)/\Gamma(\Lambda \to p\pi^-) \simeq 1/2}$.

Even before charm lifetimes were measured, it had been anticipated that as a ‘first’ for hadronic flavours the lifetime of a few $\boldsymbol{\cdot 10^{-13}}$ sec predicted for $\boldsymbol{c}$ *quarks* provides a meaningful benchmark for the lifetimes of weakly decaying charm *hadrons* and the lifetime ratios for the latter would differ relatively little from unity, certainly much less than for strange hadrons.

The semileptonic branching ratios provide a complementary perspective onto nonperturbative hadrodynamics. The semileptonic widths of charm *mesons* are basically universal meaning that the ratios of their semileptonic branching ratios coincide with the ratios of their lifetimes. Yet the semileptonic widths of charm *baryons* are expected to vary substantially meaning that the ratios of their semileptonic branching ratios yield information over and above what one can learn from the ratios of their lifetimes.

6˙4.1. Brief History, and Current Status of Lifetime Measurements. Unstable particles decay following an exponential law $\boldsymbol{P(t) = \exp(-t/\tau)}$, whose constant slope $\boldsymbol{\tau}$ is defined as the mean decay time in the particle's rest frame, i.e., the lifetime. Charm lifetimes were expected in the range of $\boldsymbol{10^{-12} - 10^{-13}}$. In the lab frame, a particle gets time-dilated by a factor $\boldsymbol{\gamma = E/m}$, and a measurement of the decay length $\boldsymbol{\ell}$ provides determination of the proper decay time $\boldsymbol{\ell = \gamma\beta ct}$ for each decay event. The slope of the exponential distribution of decay lengths $\boldsymbol{\ell}$ is the particle's lifetime. The decay is a probabilistic effect following an exponential distribution, therefore one needs adequate statistics, and precise measurement of decay length and particle momentum.

For an experiment at fixed target, where charm particles are produced with a typical average momentum of $50 - 60\,\mathrm{GeV}$, the expected charm lifetime translates to decay lengths of order one centimeter. At symmetrical colliders, where the charm particle is produced at rest in the lab frame, the decay length is very short, and one cannot determine it by directly observing the separation between primary (production) vertex and secondary (decay) vertex, but needs to use an impact parameter technique. For a detailed review on experimental techniques see [240, 247].

The space resolution at fixed target in the transverse plan is often of a few microns, which translates on a resolution in reconstructing the decay vertex which does depend on the charm particle momentum, and it is typically of order 10 microns in x,y and 300 microns in z. When coupled to a forward magnetic spectrometer with good momentum resolution, the resolution on the decay time is of order 30-50 fs. At collider experiments the decay length is very small, due to the lack of substancial Lorentz boost. High-resolution drift chambers are used for vertex detection, or, recently, microstrip arrays deployed in 4π geometry. The primary vertex is detected as the beam spot in the interaction region. Tracks of secondaries are reconstructed as helices in (ρ, φ) plane. Distribution of proper decay times is not an exponential as in the fixed target case, but a gaussian with an right-hand exponential tail, due to the relatively poorer space resolution. Proper time resolution, thanks to better momentum resolution, of state of the art e^+e^- collider experiments is now comparable to fixed target experiments.

The output signals from the vertex detector are used in track- and vertex-finding algorithms. The principal approach for vertex finding is the candidate-driven algorithm. It consists of determining a set of tracks which have been particle-identified and that are compatible with a charm decay topology. Tracks are requested to be compatible as coming from a common decay vertex. The reconstructed charm particle momentum is then projected to the primary vertex, thus determining the primary-secondary separation ℓ and error σ_ℓ. At e^+e^- (Ref.[241],[242]), candidate tracks are fit to a common vertex, the reconstructed momentum is projected to the interaction region to obtain the decay length.

A selection cut which requires the presence of a parent D^* for D^0 lifetimes (D^*-*tagging*) may be used, albeit at the cost of a reduction in statistics. To cope with the reduced reconstruction efficiency at small decay lengths, one uses the reduced proper time $t' \equiv (\ell - N\sigma_\ell)/\beta\gamma c$ where N is the primary-secondary detachment cut applied. Using t' instead of t corresponds to starting the clock for each event at a fixed detachment significance, and thus the distribution of t' recovers exponential behaviour. Additional bonus is given by the correction function determined by montecarlo simulation which, when espressed in terms of t', provides very small corrections thus reducing the contribution to the systematic error [243].

Main sources of systematics at fixed-target are absorption of both secondary tracks and charm in target, knowledge of D momentum, backgrounds, and montecarlo event sample size. For e^+e^- colliders, the determination of the decay vertex, beam spot, knowledge of D momentum, time-mass correlations, large t outlier events, decay length bias, backgrounds, and montecarlo event sample size.

New results are shown in Tab.V, with updated world averages with respect to PDG02. Lifetimes ratios significant for comparison to theoretical predictions are listed in Tab.VI. While the accuracy on the lifetimes of long-lived mesons is now at the level of the percent, essentially systematics-dominated, the measurement of very short-lived charm states, such as the Ω_C, still poses relevant challenges. In this case the superior decay times resolution of fixed-target experiments (of order 30 fs), although comparable to the lifetime

TABLE V. – *Summary of world averages from [131], new results, and updated world averages. Statistical and systematical errors are summed in quadrature.*

	Experiment	Lifetime (fs)	Events	Channels	Techn.
D^+	New Average	$\mathbf{1045 \pm 8}$			
	FOCUS [244]	$\mathbf{1039.4 \pm 8}$	110k	$K2\pi$	γN
	BELLE [245]prel	$\mathbf{1037 \pm 13}$	8k	$K2\pi$	e^+e^-
	PDG02	$\mathbf{1051 \pm 13}$			
D^0	New Average	$\mathbf{410.6 \pm 1.3}$			
	FOCUS [244]	$\mathbf{409.6 \pm 1.5}$	210k	$K\pi, K3\pi$	γN
	BELLE[248, 249] prel	$\mathbf{412.6 \pm 1.1}(st)$	448k	$K\pi$	e^+e^-
	PDG02	$\mathbf{411.7 \pm 2.7}$			
D_s^+	New Average	$\mathbf{494 \pm 5}$			
	FOCUS [246] prel	$\mathbf{506 \pm 8}(st)$	6k	$\phi\pi$	γN
	BELLE [245] prel	$\mathbf{485 \pm 9}$	6k	$\phi\pi$	e^+e^-
	PDG02	$\mathbf{490 \pm 9}$			
Λ_c^+	PDG02	$\mathbf{200 \pm 6}$			
Ξ_c^+	PDG02	$\mathbf{442 \pm 26}$			
Ξ_c^0	New Average	$\mathbf{108 \pm 15}$			
	FOCUS [250]	$\mathbf{118 \pm 14}$	$\mathbf{110 \pm 17}$	$\Xi^-\pi^+, \Omega^-K^+$	γN
	PDG02	$\mathbf{98^{+23}_{-15}}$			
Ω_c^0	New Average	$\mathbf{76 \pm 11}$			
	FOCUS [251]	$\mathbf{79 \pm 15}$	64 ± 14	$\Omega^-\pi^+, \Xi^-K^-\ 2\pi^+$	γN
	PDG02	$\mathbf{64 \pm 20}$			

itself, does allow lifetime determination at the level of 15%.

6·4.2. Early phenomenology. All charm hadrons share one contribution, namely the weak decay of the charm quark, which to leading order is not modified by its hadronic environment, see Fig. 17. It is often called the spectator process, since the other partons present in the hadron (antiquarks, quarks and gluons) remain passive bystanders ([17]). This reaction contributes to all charm hadrons equally scaling like

$$\Gamma_{Spect} \propto G_F^2 m_c^5 \tag{96}$$

Originally it had been expected that this term dominates the lifetime already for charm hadrons implying a small difference between $\tau(D^0)$ and $\tau(D^+)$. It caused quite a stir in the community when the lifetime measurements showed the D^+ to be longer lived than the D^0 by a considerable factor [252, 253]. It enhanced the drama that the first data 'overshot the target', i.e. yielded $\tau(D^+)/\tau(D^0) \sim 5$ before 'retreating' to a ratio of

([17])The reader should be warned that some authors use the term 'spectator contribution' for processes where the 'spectators' become active.

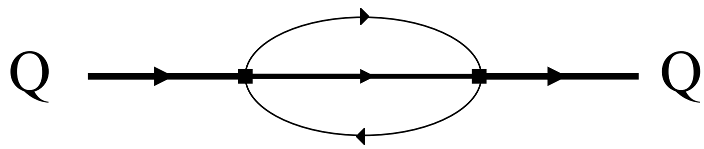

Fig. 17. – The diagram describing the weak decay of the charm quark. This spectator process contributes to the width of all charmed hadrons.

$\sim \mathbf{2.5}$. This surprise caused considerable activities on the theory side to accommodate the data and make new predictions for the other lifetime ratios.

Two mechanisms were quickly put forward as to induce $\tau(D^+)/\tau(D^0) > 1$:

1. On the Cabibbo-favoured level W-exchange can contribute to D^0, as shown in Fig. 18 a), but not to D^+ transitions. The latter are affected only on the Cabibbo suppressed level by W exchange in the s-channel, Fig. 18 b).

2. In the reaction $D^+ = [c\bar{d}] \to s\bar{d}u\bar{d}$ one has two identical quark flavours in the final state, see Fig. 19. Interference effects thus have to be included, which turn out to be destructive for nontrivial reasons. This effect is referred to as Pauli Interference (PI) .

The first mechanism had been known all along, yet its contributions had been discarded for a reason. For it suffers from two suppression factors, namely helicity suppression and wavefunction suppression: (i) With spin-one couplings conserving chirality, a pseudoscalar meson cannot decay into a massless fermion-antifermion pair. Thus the *amplitude* for W-exchange is proportional to the mass of the heaviest quark in the *final* state: $T(D^0 \to s\bar{d}) \propto m_s/M_D$. This is a repetition of the well known tale why $\Gamma(\pi^+ \to \mu^+\nu) \gg \Gamma(\pi^+ \to e^+\nu)$ holds. (ii) Due to the almost zero range of the weak force the c and $\bar{u}$ quark wavefunctions have to overlap to exchange a W boson. The decay constant f_D provides a measure for this overlap: $f_D \simeq \sqrt{12}|\psi_{c\bar{u}}(0)|/\sqrt{M_D}$. Accordingly $T(D^0 \to s\bar{d}) \propto f_D/M_D$.

Putting everything together one obtains $\Gamma_{WX}(D^0) \propto G_F^2|f_D|^2 m_s^2 m_D$ rather than $\Gamma_{Spect} \propto G_F^2 m_c^5$. Since $m_s, f_D \ll m_c$ such a W-exchange contribution is very small. Yet after the data revealed a large lifetime ratio, various mechanisms were suggested that might vitiate helicity suppression and overcome wavefunction suppression. Most of them were quite ad-hoc. One that appeared natural was to invoke gluon emission from the initial light (anti)quark line leading to

$$\Gamma_{WX}(D^0 \to s\bar{d}g) \propto \frac{\alpha_S}{\pi}\left(\frac{f_D}{\langle E_{\bar{u}}\rangle}\right)^2 G_F^2 m_c^5 \tag{97}$$

with $\langle E_{\bar{u}}\rangle$ denoting the average energy of $\bar{u}$ inside the D^0. Despite the double penalty one pays here – it is a term $\sim \mathcal{O}(\alpha_S/\pi)$ controlled by three-body rather than two-body phase space – such a contribution could be sizeable for $\langle E_{\bar{u}}\rangle \sim f_D$.

A distinction had been made between W *exchange* and *weak annihilation* with the weak boson being exchanged in the t and s channels, respectively. Yet since QCD renormalization mixes those two diagrams, it is more meaningful to lump both of them together under the term weak annihilation (WA) as far as *meson* decays are concerned; W

Fig. 18. – a) The Cabibbo-favoured weak exchange contribution to the D^0 width. b) The Cabibbo-suppressed weak annihilation contribution to the D^+ width.

exchange in *baryon* decays, which is not helicity suppressed, is denoted by W scattering (WS) .

While the occurrence of PI in exclusive D^+ decays had been noted, its relevance for inclusive D^+ rates had been discarded for various reasons. Yet in a seminal paper [252] it was put forward as the source of the observed lifetime ratio.

Both mechanisms can raise $\tau(D^+)/\tau(D^0)$ considerably above unity, yet in different ways and with different consequences: (i) While PI reduces $\Gamma(D^+)$ without touching $\Gamma(D^0)$, WA was introduced to enhance $\Gamma(D^0)$ while hardly affecting $\Gamma(D^+)$. As a consequence, since $\Gamma_{SL}(D^+) \simeq \Gamma_{SL}(D^0)$ due to isospin invariance, PI will enhance $\mathrm{BR}_{SL}(D^+)$ while not changing $\mathrm{BR}_{SL}(D^0)$, whereas WA will decrease $\mathrm{BR}_{SL}(D^0)$ while hardly affecting $\mathrm{BR}_{SL}(D^+)$. (ii) For PI one expects $\tau(D^0) \sim \tau(D_s)$, whereas WA should induce a difference in $\tau(D_s)$ vs. $\tau(D^0)$ roughly similar to $\tau(D^+)$ vs. $\tau(D^0)$. (iii) One would expect to see different footprints of what is driving the lifetime differences in the weight of different *exclusive* modes like $D_s \to \pi's$ vs. $D_s \to \phi + \pi's$ vs. $D_s \to K\bar{K} + \pi's$.

Weak baryon decays provide a rich laboratory for these effects: WS is not helicity suppressed in baryon decays, and one can count on it making significant contributions here. Furthermore PI can be constructive as well as destructive, and its weight is more stable under QCD radiative corrections than is the case for meson decays. On fairly general grounds a hierarchy is predicted [254]:

$$\tau(\Omega_c) < \tau(\Xi_c^0) < \tau(\Lambda_c) < \tau(\Xi_c^+) \tag{98}$$

It should be noted that all these analyses invoked – usually implicitly – the assumption that a valence quark description provides a good approximation for computing such transition rates. For if these hadrons contained large 'sea' components, they would all share the same basic reactions, albeit in somewhat different mixtures; this would greatly dilute the weight of processes specific to a given hadron and thus even out lifetime differences.

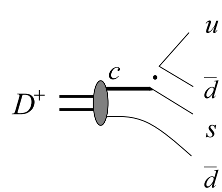

Fig. 19. – The presence of two identical quarks in the final state of the D^+ decay leads to Pauli Interference.

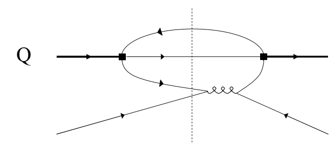

Fig. 20. – The interference of the spectator and WA diagrams that must be included for fully inclusive transitions.

These phenomenological treatments laid important groundwork, in particular in exploring various possibilities. Yet there were some serious shortcomings as well: in particular there was no agreement on the size of WA contributions in D decays, i.e. to which degree the helicity suppression could be overcome; even their scaling with m_c was controversial. One should also be quite surprised – actually mystified – by Eq.(97): for one would expect an inclusive width to be described by short-distance dynamics, which Eq.(97) manifestly is not due to the low scale $\langle E_{\bar{u}}\rangle$ appearing in the denominator. Lastly, as we will explain below, those treatments overlooked one fact of considerable significance concerning semileptonic branching ratios.

6˙4.3. The HQE description. The HQE implemented through the OPE can overcome the shortcomings inherent in the phenomenological models. Yet before it could be fully developed, the just mentioned apparent paradox posed by the expression for $\Gamma_{WX}(D^0 \to s\bar{d}g)$, Eq.(97) had to be resolved. This was achieved in Ref.[255]. For a truly inclusive transition to order α_S one had to include also the interference between the spectator and WA diagrams, where the latter contains an off-shell gluon going into a $q\bar{q}$ pair, see Fig.(20). It was shown that when one sums over all contributions through order α_S, the small energy denominators $1/\langle E_{\bar{u}}\rangle^2$ and $1/\langle E_{\bar{u}}\rangle$ disappear due to cancellations among the different diagrams, as it has to on general grounds. This is just another 'toy model' example that while fully inclusive rates are short-distance dominated, partially inclusive ones – let alone exclusive ones – are not.

The basic method of the HQE has been described in Sect.(**4**˙6); it yields:

$$\Gamma(H_Q \to f) = \frac{G_F^2 m_Q^5}{192\,\pi^3}|KM|^2\Big[A_0 + \frac{A_2}{m_Q^2} + \frac{A_3}{m_Q^3} + \mathcal{O}(1/m_Q^4)\Big] \tag{99}$$

The quantities A_n contain the phase space factors as appropriate for the final state, the QCD radiative corrections and the short-distance coefficients appearing in the OPE and $\langle H_c|\mathcal{O}_n|H_c\rangle$, the hadronic expectation values of local operators $\mathcal{O}_n$ of dimension $n+3$. They scale like μ^n, where μ denotes an ordinary hadronic scale close to and hopefully below 1 GeV. Each term has a transparent physical meaning; let us stress those features that are particularly relevant for lifetimes and semileptonic branching ratios.

- In all cases one has to use the same value of the charm quark mass m_Q properly defined in a field theoretical sense. Furthermore this value is in principle not a free parameter to be fitted to the data on lifetimes, but should be inferred from other observables. Choosing different values for m_Q when describing, say, meson and baryon lifetimes can serve merely as a *temporary crutch* to parameterize an observed difference between mesons and baryons one does not understand at all.

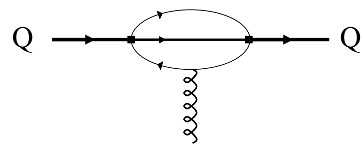

Fig. 21. – The spin interaction of the heavy quark with the light degrees of freedom, a leading nonperturbative effect.

When using quark models to evaluate these expectation values, charm quark masses will also enter through the quark wave functions. Yet those are like 'constituent' masses, i.e. parameters specific to the model used, *not* quantities in the QCD Lagrangian; therefore they can be adjusted for the task at hand.

- The leading term A_0 represents the spectator diagram contribution common to all hadrons, see Fig.(17).
- The *leading* nonperturbative corrections arise at order $1/m_Q^2$. A_2 reflects the motion of the heavy quark inside the hadron and its spin interaction with the light quark degrees of freedom, see Fig.(21). This latter effect had not been anticipated in the phenomenological descriptions of the 1980's.

  These terms in general differentiate between baryons on one side and mesons on the other, yet have practically the same impact on all mesons. However the contribution proportional to μ_G^2 due to the aforementioned spin interaction *unequivocally enhance* the nonleptonic width over the semileptonic one and thus reduce the semileptonic branching ratios of D mesons. This does not happen in Λ_c decays; for $\mu_G^2(\Lambda_c) \simeq 0$, since the light diquark system carries spin zero there.
- Pauli Interference (PI) [252], Weak Annihilation (WA) for mesons and W-scattering (WS) for baryons arise unambiguously and naturally in order $1/m_Q^3$ with WA being helicity suppressed and/or nonfactorisable[255]. They mainly drive the differences in the lifetimes of the various hadrons of a given heavy flavour.
- For the *lifetime ratios* effects of order $1/m_Q^3$ rather than $1/m_Q^2$ are numerically dominant. There are two reasons for that, one straightforward and one more subtle: (i) As illustrated by Figs.(22 a-d) the $\mathcal{O}(1/m_Q^3)$ contributions from four-fermion operators involve integrating over only two partons in the final state rather than the three as for the decay contribution, Fig.(17 a). A_3 is therefore enhanced relative to $A_{0,2}$ by a phase space factor; however the latter amounts effectively to considerably less than the often quoted $16\pi^2$. (ii) As explained in Sect. 4˙6 there are no $\mathcal{O}(1/m_Q)$ contributions. One can actually see that there are various sources for such contributions, but that they cancel exactly between initial and final state radiation. Such cancellations still arise for $\mathcal{O}(1/m_Q^2)$, which are thus reduced relative to their 'natural' scale. In $\mathcal{O}(1/m_Q^3)$ etc., however, they have become ineffective. The latter are thus of normal size, whereas the $1/m_Q^2$ contributions are 'anomalously' small. Accordingly there is no reason to suspect $1/m_Q^4$ terms to be larger than $1/m_Q^3$ ones.
- The HQE is better equipped to predict the *ratios* of lifetimes, rather than lifetimes themselves. For the (leading term in the) full width scales with m_Q^5, whereas

the numerically leading contributions generating lifetime differences are due to dimension-six operators and scale with m_Q^2, i.e. are of order $1/m_Q^3$. Uncertainties in the value of m_Q thus affect lifetime ratios much less.

6˙4.4. Theoretical interpretation of the lifetime ratios. As already stated the three weakly decaying mesons D^+, D_s^+ and D^0 receive identical contributions from the leading term A_0 in Eq.(99) ([18]). This is largely true also for A_2 although less obvious [256]. Yet to order $1/m_Q^3$ their lifetimes get differentiated: on the Cabibbo favoured level PI contributes to D^+ and WA to D^0 and D_s^+ decays. Yet a careful HQE analysis reveals that the WA contributions are helicity suppressed and/or suppressed due to being nonfactorizable etc. [255]. Thus one expects approximate equality between the D^0 and D_s^+ lifetimes: $\Gamma(D^0) \simeq \Gamma_{\text{spect}}(D) \simeq \Gamma(D_s^+)$. In the D^+ width on the other hand PI occurs due to interference between two $\bar{d}$ quark fields, one from the wavefunction, while the second one emerges from the decay. A priori there is no reason for this effect to be small. More specifically one finds

$$\Gamma(D^+) \simeq \Gamma_{\text{spect}}(D) + \Delta\Gamma_{PI}(D^+) \tag{100}$$

$$\Delta\Gamma_{PI}(D^+) \simeq \Gamma_0 \cdot 24\pi^2 \left(f_D^2/m_c^2\right) \kappa^{-4} \cdot$$

$$\cdot \left[(c_+^2 - c_-^2)\kappa^{9/2} + \frac{1}{N_C}(c_+^2 + c_-^2) - \frac{1}{9}(\kappa^{9/2} - 1)(c_+^2 - c_-^2) \right] \tag{101}$$

where

$$\kappa = [\alpha_S(\mu_{had}^2)/\alpha_S(m_c^2)]^{1/b}\, , b = 11 - \frac{2}{3}N_f \tag{102}$$

reflects hybrid renormalization mentioned in Sect.(**4**˙10.1). A few comments are in order: without QCD corrections one has $c_- = c_+ = 1 = \kappa$ and thus $\Delta\Gamma_{PI}(D^+) > 0$, i.e. *constructive* interference meaning $\tau(D^+) < \tau(D^0)$; including UV renormalization flips the *sign* of $\Delta\Gamma_{PI}(D^+)$ and hybrid renormalization makes this effect quite robust.

One arrives then at

$$\frac{\tau(D^+)}{\tau(D^0)} \simeq 1 + (f_D/200 \text{ MeV})^2 \simeq 2.4 \tag{103}$$

for $f_D \simeq 240$ MeV, see Eq.(49); i.e., PI is capable of reproducing the observed lifetime ratio by itself even without WA. Of course this has to be taken as a semi quantitative statement only, since we cannot claim (yet) precise knowledge of f_D, the size of the *non*factorizable contributions in the expectation value for the four-fermion operator or of the $\mathcal{O}(1/m_c^4)$ contributions.

Next one has to compare $\tau(D_s^+)$ and $\tau(D^0)$. The statement underlying $\Gamma(D^0) \simeq \Gamma_{\text{spect}}(D) \simeq \Gamma(D_s^+)$ is actually that $\frac{\tau(D_s^+)}{\tau(D^0)} \ll \frac{\tau(D^+)}{\tau(D^0)} \simeq 2.4$. It is nontrivial since one

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([18])One should note that for inclusive transitions the distinction of *internal* vs. *external* spectator diagrams makes little sense, since in fully inclusive processes – in contrast to exclusive channels – one does not specify which quarks end up in the same hadron.

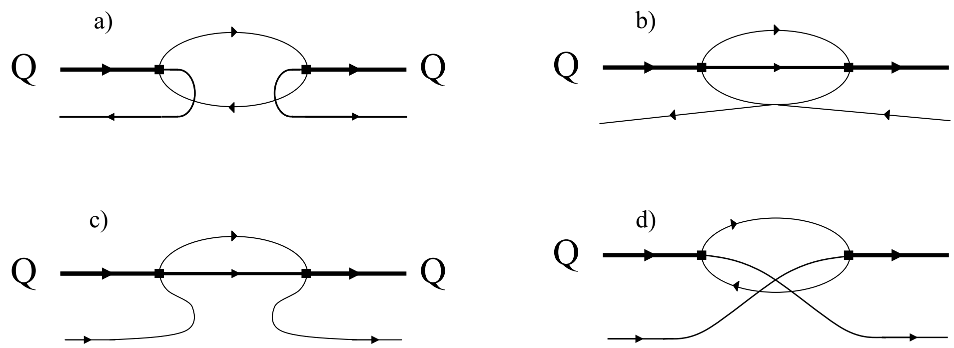

Fig. 22. – Cutting an internal quark line generates order $1/m_Q^3$ four fermion operators. a) Weak annihilation in mesonic decays. d) Pauli Interference in the decays of a meson. c) Weak scattering in baryonic decays. d) Pauli Interference in the decays of a baryon.

would not expect it to hold if WA were the main effect generating the lifetime differences among charm mesons.

A first milestone was reached with the E687 measurement:

$$\frac{\tau(D_s^+)}{\tau(D^0)} = 1.12 \pm 0.04 \; . \tag{104}$$

It provided significant, though not conclusive evidence that $\tau(D_s^+)$ exceeds $\tau(D^0)$. Even more importantly it clearly confirmed the prediction of WA being suppressed.

It has been estimated [256] that even *without* WA D_s can actually be longer lived than D^0, yet by a small amount only due to a combination of various $\mathcal{O}(\ \%\ )$ effects like PI in Cabibbo suppressed transitions:

$$\left.\frac{\tau(D_s^+)}{\tau(D^0)}\right|_{no\ WA} \simeq 1.0 \div 1.07 \tag{105}$$

Furthermore it was stated that WA can modify this ratio by about 20% in *either* direction [256].

Data have conclusively confirmed that $\tau(D_s^+)$ is moderately longer than $\tau(D^0)$.

$$\langle \tau(D_s^+)/\tau(D^0) \rangle = 1.22 \pm 0.02 \tag{106}$$

In summary: the case for PI being the dominant mechanism driving lifetime differences among D mesons rests on the following facts and observations:

1. A careful analysis of the HQE and of the expectation values of four-quark operators shows that WA contributions are either helicity suppressed or non-factorizable and thus suppressed. Accordingly they are too small for being the leading effect. At present this is a purely theoretical argument although it can be checked in the future through measurements of the lepton endpoint spectrum in semileptonic B or D decays [257].

2. The measured $D^+ - D^0$ lifetime ratio can be reproduced.

3. The observation of

$$0.07 \ll |1 - \tau(D_s^+)/\tau(D^0)| \ll |1 - \tau(D^+)/\tau(D^0)| \tag{107}$$

confirms that WA is nonleading, yet still significant. For if WA were the dominant source for the $D^+ - D^0$ lifetime ratio, one would expect $\tau(D_s^+)/\tau(D^0)$ to deviate by a similar amount from unity, which is however clearly not the case.

The observed value for $\tau(D_s)/\tau(D^0)$ is still within range of the general estimate of what can be accommodated with a modest, yet significant contribution from WA [256, 257].

It has been suggested recently that there is no need for invoking PI and WA to reproduce $\tau(D_s)/\tau(D^0)$ [258]. The authors suggest the following simple minded prescription: they pair up all exclusive D^0 channels with their D_s^+ counterparts and compute the strength of the latter by taking the observed strength of the former and applying simple phase space corrections; then they add up all these individual rates and arrive at $\tau(D_s)/\tau(D^0) \simeq 1.17$. This is not an unambiguous prescription, of course, since the appropriate phase space for multibody final states depends on the detailed dynamics of those final states. Yet the main problem with this 'explanation' is that the point is *not* whether one can relate classes of hadronic decay channels to each other approximately with the help of simple prescription. The central theoretical question is whether a quark-based, i.e. short distance treatment genuinely inferred from QCD provides an adequate description of inclusive transitions involving hadrons, which also addresses the issue of quark-hadron duality and its limitations. This question is not even addressed by an ad-hoc ansatz involving individual hadronic modes.

Isospin symmetry tells us that the semileptonic widths of D^+ and D_s^+ mesons have to be equal up to terms $\sim \mathcal{O}(\theta_C^2)$; therefore the ratio of their semileptonic branching ratios has to agree with their lifetime ratio. This is well borne out by the data:

$$\frac{\mathrm{BR}_{SL}(D^+)}{\mathrm{BR}_{SL}(D^0)} = 2.50 \pm 0.27 \; vs. \; \frac{\tau(D^+)}{\tau(D^0)} = 2.54 \pm 0.01 \tag{108}$$

When one considers the absolute values of these branching ratios

$$\mathrm{BR}_{SL}(D^+) = 17.2 \pm 1.9\% \; , \; \mathrm{BR}_{SL}(D^0) = 6.87 \pm 0.28\% \tag{109}$$

there seems to be a fly in the ointment for PI being the main reason for the lifetime difference. As already mentioned by reducing the D^+ nonleptonic width PI *enhances* $\mathrm{BR}_{SL}(D^+)$, while leaving D^0 widths largely unaffected. WA on the other hand is invoked to enhance the D^0 nonleptonic width and thus *reduces* $\mathrm{BR}_{SL}(D^0)$. Then it comes down to the question what one considers to be the 'normal' semileptonic D width, i.e. before hadronization effects differentiating between D^+ and D^0 are included.

In the phenomenological models one infers this value from the decays of quasifree charm quarks:

$$\mathrm{BR}_{SL}(D) \simeq \mathrm{BR}(c \to \ell \nu s) \sim 16\% \tag{110}$$

Comparing this expectation with the data would strongly point to WA as the dominant mechanism for the lifetime difference.

Yet the HQE introduces a new and intriguing twist here: as mentioned before in order $1/m_c^2$ it reduces the semileptonic width common to D^+ and D^0 through the chromomagnetic moment μ_G^2:

$$\mathrm{BR}_{SL}(D^+) \simeq \mathrm{BR}_{SL}(D^0) + \mathcal{O}(1/m_c^3) \sim 8\% \; . \tag{111}$$

The actual numbers have to be taken with quite a grain of salt, of course. Yet this effect – which had been overlooked in the original phenomenological analyses – makes the findings of PI being the main engine consistent with the absolute values measured for the semileptonic branching ratios.

$SU(3)$ symmetry by itself would allow for some still sizeable difference in the semileptonic widths for D_s and D^0 mesons. Yet the HQE yields that $\Gamma_{SL}(D_s)$ and $\Gamma_{SL}(D^0)$ agree to within a few percent. Therefore one predicts

$$\mathrm{BR}_{SL}(D_s) = \mathrm{BR}_{SL}(D^0) \cdot \frac{\tau(D_s)}{\tau(D^0)} \sim 1.2 \cdot \mathrm{BR}_{SL}(D^0) \tag{112}$$

A detailed analysis of charm baryon lifetimes is not a 'deja vu all over again'. While it constitutes a complex laboratory to study hadronization, it will yield novel lessons: (i) There are four weakly decaying baryons: Λ_c, $\Xi_c^{(0,+)}$ and Ω_c. Since we can trust HQE in charm decays on a semi quantitative level only, it makes a considerable difference whether one can reproduce the pattern of seven rather than three hadronic lifetimes. (ii) The theoretical challenge is considerably larger here since there are more effects driving lifetime differences in order $1/m_Q^3$: WS contributions are certainly not helicity suppressed – they are actually enhanced by QCD radiative corrections and by two-body over three-body phase space; also PI can be constructive as well as destructive. Details can be found in Ref.[240]. (iii) While the light diquark system forms a scalar in Λ_c and Ξ_c, it carries spin one in Ω_c. (iv) The semileptonic widths are expected to be strongly modified by constructive PI rather than being universal. (v) There is no unequivocal concept of factorization for the *baryonic* expectation values of four-fermion-operators. Accordingly we have to depend on quark model calculations of those matrix elements more than it is the case for mesons. One can entertain the hope that lattice QCD will provide a reliable handle on those quantities – yet that would require unquenched studies. The one saving grace is that the wave functions employed can be tested by examining whether they can reproduce the observed mass splittings among charm baryons and their resonances.

The Λ_c, Ξ_c and Ω_c nonleptonic and semileptonic widths receive significantly different contributions in order $1/m_c^3$ from WS and constructive as well as destructive PI; for details see the review [240]. The hierarchy stated in Eq.(98) arises very naturally.

When studying *quantitative* predictions we have to be aware of the following complexity: Since the three *non-universal* contributions WS, constructive and destructive PI are of comparable size, the considerable uncertainties in the individual contributions get magnified when cancellations occur between Γ_{WS} & $\Delta\Gamma_{PI,+}$ on one side and $\Delta\Gamma_{PI,-}$ on the other.

Table VI summarizes the predictions and data. From this comparison one can conclude:

- The observed lifetime *hierarchy* emerges correctly and naturally.

	$1/m_c$ expect. [240]	theory comments	data
$\frac{\tau(D^+)}{\tau(D^0)}$	$\sim 1 + \left(\frac{f_D}{200\ \mathrm{MeV}}\right)^2 \sim 2.4$	PI dominant	2.54 ± 0.01
$\frac{\tau(D_s^+)}{\tau(D^0)}$	1.0 - 1.07 0.9 - 1.3	*without* WA [256] *with* WA [256]	 1.22 ± 0.02
$\frac{\tau(\Lambda_c^+)}{\tau(D^0)}$	~ 0.5	quark model matrix elements	0.49 ± 0.01
$\frac{\tau(\Xi_c^+)}{\tau(\Lambda_c^+)}$	$\sim 1.3 \div 1.7$	ditto	2.2 ± 0.1
$\frac{\tau(\Lambda_c^+)}{\tau(\Xi_c^0)}$	$\sim 1.6 \div 2.2$	ditto	2.0 ± 0.4
$\frac{\tau(\Xi_c^+)}{\tau(\Xi_c^0)}$	~ 2.8	ditto	4.5 ± 0.9
$\frac{\tau(\Xi_c^+)}{\tau(\Omega_c)}$	~ 4	ditto	5.8 ± 0.9
$\frac{\tau(\Xi_c^0)}{\tau(\Omega_c)}$	~ 1.4	ditto	1.42 ± 0.14

TABLE VI. – *Lifetime ratios in the charm sector*

- There is *no a priori* justification for a $1/m_Q$ expansion to work already for the moderate charm quark mass, only an *a posteriori* one. For even the quantitative predictions are generally on the mark keeping in mind that one has to allow for at least 30% uncertainties due to contributions of higher order in $1/m_c$. For proper appreciation one should note that the lifetimes span more than an order of magnitude:

(113)
$$\frac{\tau(D^+)}{\tau(\Omega_c)} \sim 20$$

  One should also keep in mind that total widths and width *differences* (among mesons) scale like m_Q^5 and m_Q^2, respectively. The latter is thus more stable under variations in the value of m_Q than the former. It actually has been known for some time that with $m_c \simeq 1.3$ GeV one reproduces merely about two thirds of the observed value of $\Gamma(D \to l\nu X)$ through order $1/m_Q^3$ retaining factorizable contributions only.

- One discrepancy stands out, though: the Ξ_c^+ appears to live considerably longer than predicted, namely by about 50%. If one multiplied $\tau(\Xi_c^+)$ by an *ad-hoc* factor of 1.5, then all the predictions for the baryonic lifetime ratios would be close to the central values of the measurements! One possible explanation is to attribute it to an anomalously large violation of quark-hadron duality induced by the accidental proximity of a baryonic resonance with appropriate quantum numbers near the Ξ_c^+, which interferes destructively with the usual contributions. Since the charm region is populated by many resonances, such 'accidents' are quite likely to happen.

- In computing these ratios one has made an assumption beyond the OPE: one has adopted a valence quark ansatz in evaluating four-quark operators; i.e.,

$$\langle D^+|(\bar c\Gamma u)(\bar u\Gamma c)|D^+\rangle = 0 = \langle D^0|(\bar c\Gamma d)(\bar d\Gamma c)|D^0\rangle \tag{114}$$

$$\langle D^{0,+}|(\bar c\Gamma s)(\bar s\Gamma c)|D^{0,+}\rangle = 0 \tag{115}$$

etc. Such an ansatz cannot be an identity, merely an approximation due to the presence of 'sea' quarks or quark condensates $\langle 0|\bar q q|0\rangle$. The fact that one obtains the correct lifetime patterns shows a posteriori that it is – maybe not surprisingly – a good approximation. Since an expansion in $1/m_c$ is of limited numerical reliability only, there is neither a need nor a value in going beyond this approximation.

Isopin symmetry tells us that $\Gamma(\Xi_c^+ \to \ell\nu X_s) = \Gamma(\Xi_c^0 \to \ell\nu X_s)$ holds. Yet it would be highly misleading to invoke $SU(3)_{Fl}$ to argue for in general universal semileptonic widths of charm baryons. On the contrary one actually expects large differences in the semileptonic widths mainly due to constructive PI in Ξ_c and Ω_c transitions; the lifetime ratios among the baryons will thus not get reflected in their semileptonic branching ratios. One estimates [260, 261]:

$$BR_{SL}(\Xi_c^0) \sim BR_{SL}(\Lambda_c) \leftrightarrow \tau(\Xi_c^0) \sim 0.5 \cdot \tau(\Lambda_c) \tag{116}$$

$$BR_{SL}(\Xi_c^+) \sim 2.5 \cdot BR_{SL}(\Lambda_c) \leftrightarrow \tau(\Xi_c^+) \sim 1.7 \cdot \tau(\Lambda_c) \tag{117}$$

$$BR_{SL}(\Omega_c) < 25\% \tag{118}$$

The conventional way to measure the absolute size of these semileptonic branching ratios is to study

$$e^+e^- \to \Lambda_c\bar\Lambda_c,\ \Xi_c\bar\Xi_c,\ \Omega_c\bar\Omega_c \tag{119}$$

Unfortunately it seems unlikely that the tau-charm factory proposed at Cornell University will reach the Ξ_c threshold. Yet the spectacular success of the B factories BELLE and BABAR has pointed to a novel method for measuring these branching ratios by carefully analyzing $B \to \Lambda_c/\Xi_c + X$ transitions. One would proceed in two steps: One reconstructs one B mesons more or less fully in $\Upsilon(4S) \to B\bar B$ and then exploits various correlations between baryon and lepton numbers and strangeness.
We have pointed out before that the lifetime ratios between charm mesons are much smaller than among kaons. It is curious to note that the lifetime ratios between charm baryons differ somewhat *more* from unity than it is the case for strange baryons:

$$\frac{\tau(\Xi^0)}{\tau(\Sigma^+)} \simeq 3.6 \tag{120}$$

The fact that the lifetimes for strange baryons – unlike those for strange mesons – are comparable is attributed to the fact that all baryons can suffer $\Delta I = 1/2$ transitions. One can add that even the observed hierarchy

$$\tau(\Sigma^+) \simeq \tau(\Omega) < \tau(\Sigma^-) \simeq \tau(\Xi^-) < \tau(\Lambda) < \tau(\Xi^0) \tag{121}$$

runs counter to expectations based on considering PI and WS contributions as the engine behind lifetime differences in the strange sector.

6.4.5. Future prospects. When relying on HQE to describe charm hadron lifetimes we have to allow for several sources of theoretical uncertainties:

1. With the expansion parameter Λ_{NPD}/m_c for nonperturbative dynamics only moderately smaller than unity, unknown higher order terms could be sizeable; likewise for higher order perturbative corrections.

2. The expectation values of four-fermion operators – let alone of higher-dimensional operators – are not well known.

3. Terms of the form $e^{-m_c/\Lambda_{NPD}}$ cannot be captured by the OPE. They might not be insignificant (in contrast to the case with beauty quarks).

4. 'Oscillating' terms $\sim (\Lambda_{NPD}/m_c)^k \cdot \sin(m_c/\Lambda_{NPD})$, $k > 0$, likewise are not under good theoretical control. They reflect phenomena like hadronic resonances, threshold effects etc.

Assuming each of these effects to generate a 10% uncertainty is not conservative, and one cannot count on the overall theoretical uncertainty to fall below 20 %. With the exception of the second item in this list (see its discussion below), the situation is highly unlikely to change decisively anytime soon.

Any of the sources listed above a priori could have produced uncertainties of, say, 30% combining into an overall error of ~ 100%. Our description of charm lifetimes would have clearly failed then – yet *without* causing alarm for the general validity of QCD as the theory of strong interactions. However this failure did not happen – the expected (and partially even predicted) pattern is in at least semi quantitative agreement (except for $\tau(\Xi_c^+)$) with data that are quite mature now. The general lesson is that the transition from nonperturbative to perturbative dynamics in QCD is smoother and more regular than one might have anticipated. This insight is bound to enhance our confidence that we can indeed treat various aspects in the decays of beauty hadrons in a *quantitative* way.

On the other hand the analysis of charm lifetimes is not meant to yield precise quantitative lessons on QCD. With the lifetimes of D^0, D^+ and D_s now known within 1% and of Λ_c within 3% measuring them even more precisely will not help our understanding of charm lifetimes, since that is already limited by theory. One exception to this general statement is the search for $D^0 - \bar{D}^0$ oscillations to be discussed later.

There are two areas where further progress seems feasible:

- There is an explicit calculation of $\tau(D_s^+)/\tau(D^0)$ based on QCD sum rules to estimate the matrix element controlling WA [259]. It finds a ratio that – while larger than unity – appears to fall well below the data:

$$\frac{\tau(D_s)}{\tau(D^0)} \simeq 1.08 \pm 0.04 \tag{122}$$

  It would be premature to view this estimate as conclusive. First we have to deepen our understanding of WA. Its leading impact is due to the expectation value of a four-fermion operators $(\bar{c}_L\gamma_\mu q_L)(\bar{q}_L\gamma_\mu c_L)$ and $(\bar{c}_L\gamma_\mu\lambda_i q_L)(\bar{q}_L\gamma_\mu\lambda_i c_L)$. It would provide an interesting benchmark test to extract their sizes from data and compare them with the predictions from lattice QCD. These quantities will also affect the lepton energy endpoint spectra in inclusive semileptonic D^+ and D_s

decays. Lastly – and maybe most significantly – their b quark counterparts are expected to have a sizeable impact on the lepton energy spectra in B decays with the effects being different for B_d and B^- [257]. This introduces an additional large uncertainty into extracting $V(ub)$ there.

- With WA driving about 20% of all D decays, the further challenge naturally arises whether footprints of WA can be found in *exclusive* modes. I.e., can one show that WA – rather than modifying all nonleptonic decays in a basically uniform way – affects certain channels much more significantly than others. We will return to this issue when discussing exclusive decays.

The data on the lifetimes of the other charm baryons have recently reached a new maturity level. Yet even so, further improvements would be desirable, namely to measure also $\tau(\Xi_c^0)$ and $\tau(\Omega_c)$ to within 10% or even better.

We have identified one glaring problem already, namely that the Ξ_c^+ lifetime exceeds expectation by about 50 %, while the other lifetime ratios are close to expectations. At first sight this might suggest the accidental presence of a baryonic resonance near the Ξ_c^+ mass inducing a destructive interference. Since the charm region is still populated by light-flavour baryon resonances, such an effect might not be that unusual to occur. However Ξ_c^0 (or Λ_c) would be more natural 'victims' for such an 'accident' since there the final state carries the quantum numbers that an $S = -2$ (or $S = -1$) baryon resonance can possess. It would be interesting to see whether future measurements leave the ratios $\tau(\Lambda_c)/\tau(\Xi_c^0)$ and $\tau(\Xi_c^0)/\tau(\Omega_c^0)$ in agreement with expectations.

There is another motivation for analysing charm baryon lifetimes. For several years now there has been a persistent and well publicized problem in the beauty sector, namely that the observed beauty baryon lifetime falls below HQE predictions:

$$\left.\frac{\tau(\Lambda_b)}{\tau(B_d)}\right|_{data} = 0.797 \pm 0.052 \;\; vs. \;\; 0.88 \leq \left.\frac{\tau(\Lambda_b)}{\tau(B_d)}\right|_{HQE} \leq 1 \tag{123}$$

The data are not conclusive yet, and further insights will be gained by measuring $\tau(\Xi_b^0)$ and $\tau(\Xi_b^-)$, which is expected to be done during the present Tevatron run. Theory might still snatch victory of the jaws of defeat.

Yet in any case the lifetimes of charm baryons can act as important diagnostics also for interpreting either success or failures in the HQE predictions for the lifetimes of beauty baryons – as long it makes some sense to treat charm quarks as heavy. The weak link in the HQE analysis of baryon widths are the expectation values of the four-quark operators, since there size is inferred from quark models or QCD sum rules of less than sterling reliability. It has been suggested by Voloshin [260] and Guberina *et al.* [262] to combine HQS, isospin and $SU(3)_{Fl}$ symmetry to relate lifetime differences among charm baryons to those among beauty baryons. From

$$\frac{\Gamma(\Xi_b^-) - \Gamma(\Xi_b^0)}{\Gamma(\Xi_c^+) - \Gamma(\Xi_c^0)} = \frac{m_b^2}{m_c^2}\frac{|V(cb)|^2}{|V(cs)|^2}\left(1 + \mathcal{O}(1/m_c, 1/m_b)\right) \tag{124}$$

and using the observed lifetime difference between Ξ_c^+ and Ξ_c^0 one infers

$$\Gamma(\Xi_b^-) - \Gamma(\Xi_b^0) = (-0.14 \pm 0.06)\; ps^{-1} \;, \tag{125}$$

which is quite consistent with other predictions. Furthermore to the degree one can ignore quark masses in the final state one can approximately equate $\Gamma(\Lambda_b)$ and $\Gamma(\Xi_b^0)$ since both are subject to WS and to destructive PI, Λ_b in the $b \to c\bar{u}d$ and Ξ_b^0 in the $b \to c\bar{c}s$ channels. Using the observed Λ_b lifetime one infers from Eq.(125):

$$\frac{\tau(\Xi_b^-) - \tau(\Lambda_b)}{\tau(\Xi_b^-)} = 0.17 \pm 0.07 \tag{126}$$

6·5. *Masses, weak lifetimes and semileptonic branching ratios of* $C \geq 2$ *baryons*. – The nonleptonic and semileptonic widths of these baryons are even more sensitive probes of the dynamics underlying their structure.

The leading contribution is contained in the quark decay term

$$\begin{aligned} \langle H_{cc}|\bar{c}c|H_{cc}\rangle &= 2 - \frac{1}{2}\frac{\mu_\pi^2(H_{cc})}{m_c^2} + \frac{1}{2}\frac{\mu_G^2(H_{cc})}{m_c^2} + \mathcal{O}(1/m_c^3) \\ \langle \Omega_{ccc}^{++}|\bar{c}c|\Omega_{ccc}^{++}\rangle &= 3 - \frac{1}{2}\frac{\mu_\pi^2(\Omega_{ccc}^{++})}{m_c^2} + \frac{1}{2}\frac{\mu_G^2(\Omega_{ccc}^{++})}{m_c^2} + \mathcal{O}(1/m_c^3)\ , \end{aligned} \tag{127}$$

where the first term of two [three] reflects the fact that there are two [three] valence charm quarks inside H_{cc} [Ω_{ccc}] and the leading nonperturbative corrections are expressed through the kinetic energy moment $\mu_\pi^2(H_{cc})$ and chromomagnetic moment $\mu_G^2(H_{cc})$. The main differences among the widths of the $C = 2$ baryons arise in order $1/m_c^3$ due to WS and destructive as well as constructive PI, similar to the case of $C = 1$ baryons:

$$\begin{aligned} \Gamma_{NL}(\Xi_{cc}^+) &\simeq \left[\Gamma_{decay,NL}(\Xi_{cc}) + \Gamma_{WS}(\Xi_{cc}^+)\right] \\ \Gamma_{NL}(\Xi_{cc}^{++}) &\simeq \left[\Gamma_{decay,NL}(\Xi_{cc}) - \Delta\Gamma_{PI,-}(\Xi_{cc}^{++})\right] \\ \Gamma_{NL}(\Omega_{cc}^+) &\simeq \left[\Gamma_{decay,NL}(\Omega_{cc}) + \Delta\Gamma_{PI,+}(\Omega_{cc})\right] \end{aligned} \tag{128}$$

$$\begin{aligned} \Gamma_{SL}(\Xi_{cc}^+) &\simeq \Gamma_{decay,SL}(\Xi_{cc}),\ \Gamma_{SL}(\Xi_{cc}^{++}) \simeq \Gamma_{decay,SL}(\Xi_{cc}) \\ \Gamma_{SL}(\Omega_{cc}^+) &\simeq \left[\Gamma_{decay,SL}(\Omega_{cc}) + \Delta\Gamma_{PI,+}(\Omega_{cc})\right]\ . \end{aligned} \tag{129}$$

The fact that there are two rather than one charm quark that can decay with or without PI and undergo WS is contained in the size of the expectation values $\langle H_{cc}|\bar{c}c|H_{cc}\rangle$, see Eq.(127), and $\langle H_{cc}|(\bar{c}\Gamma q)(\bar{q}\Gamma c)|H_{cc}\rangle$; i.e., $\Gamma_{decay}(\Xi_{cc}) = 2\Gamma_{decay}(\Xi_c)$ to leading order in $1/m_c$. Based on the expressions in Eq.(128) one expects substantial lifetime differences with

$$\tau(\Xi_{cc}^+),\ \tau(\Omega_{cc}) < \tau(\Xi_{cc}^{++}) \tag{130}$$

One might ask if the PI contribution to the Ξ_{cc}^{++} width could be constructive rather than destructive; in that case the Ξ_{cc}^+ and Ξ_{cc}^{++} lifetimes would be very similar and both short. While the negative sign of PI in the D^+ width is not a trivial matter, since it depends on properly including QCD radiative corrections, it is straightforward (though still not trivial) for baryons. The PI contribution depends on the combination $2c_+(2c_- - c_+)$ of the QCD renormalization coefficients $c_\pm$: since $c_+ < 1 < c_-$ the sign of the effect is stable under radiative QCD corrections.

With the D^0 width given mainly by the decay contribution, it provides an approximate yardstick for the latter's size. To *leading* order in $1/m_c$ one has

$$\frac{1}{2}\Gamma_{decay}(\Xi_{cc}) \simeq \Gamma_{decay}(\Lambda_c) \simeq \Gamma_{decay}(D) \simeq \Gamma(D^0) \tag{131}$$

Since the c quarks move more quickly in a double charm than a single charm hadron, one expects $\Gamma_{decay}(\Xi_{cc})$ to be actually somewhat smaller than $2\Gamma(D^0)$ due to time dilatation that enters in order $1/m_c^2$. *If* the PI term were absent, one would thus have $\tau(\Xi_{cc}^{++}) \geq \frac{1}{2}\tau(D^0) \simeq 2 \cdot 10^{-13} sec$. Including PI, which one confidently predicts to be destructive, one expects Ξ_{cc}^{++} to be considerably longer lived than this lower bound and possibly even longer lived than D^0. The Ξ_{cc}^{+} lifetime on the other hand will be considerably shorter than $\tau(D^0)$. The impact of WS pushing $\tau(\Xi_c^+)$ below $2{\cdot}10^{-13} sec$ is partially offset by the time dilatation effect mentioned above. Nevertheless $\tau(\Xi_c^+) \sim 10^{-13} sec$ would seem to be a reasonable first guess.

For more definite predictions one needs to estimate the relevant Ξ_{cc} expectation values. At present we have to rely on quark models and QCD sum rules. In the future those could be tested and fine tuned through their predictions on the mass splittings of the $C = 2$ baryons and their resonances; estimates based on lattice QCD might become available as well. The authors of Refs.([263],[264]) had the foresight to take on this task before there was any experimental hint for such exotic baryons. The two groups follow a somewhat different philosophy in choosing the range for m_c and m_s.

The authors of Ref.[263] focus on hadrons with two heavy constituents like the meson B_c and the baryons Ξ_{cc}, Ξ_{bc} etc. They have adopted a rather phenomenological attitude in selecting values for m_c (and m_s); from $\tau(B_c)$ they infer $m_c \sim 1.6$ GeV and find

$$\tau(\Xi_{cc}^{++}) \sim 0.46 \pm 0.05\ ps\ ;\ \tau(\Xi_{cc}^{+}) \sim 0.16 \pm 0.05\ ps$$

$$\tau(\Omega_{cc}^{+}) \sim 0.27 \pm 0.06\ ps \tag{132}$$

$$\Rightarrow \frac{\tau(\Xi_{cc}^{++})}{\tau(\Xi_{cc}^{+})} \sim 2.9\ ,\ \frac{\tau(\Xi_{cc}^{++})}{\tau(\Omega_{cc}^{+})} \sim 1.7\ ,\ \frac{\tau(\Omega_{cc}^{+})}{\tau(\Xi_{cc}^{+})} \sim 1.7 \tag{133}$$

The authors of Ref.([264]) on the other hand follow a purer invocation of the OPE and set $m_c = 1.35$ GeV and $m_s = 0.15$ GeV. After some detailed consideration including even Cabibbo suppressed modes they obtain:

$$\tau(\Xi_{cc}^{++}) = 1.05\ ps\ ,\ \tau(\Xi_{cc}^{+}) = 0.20\ ps \tag{134}$$

$$\tau(\Omega_{cc}^{+}) = 0.30\ ps\ ,\ \tau(\Omega_{ccc}^{++}) = 0.43\ ps \tag{135}$$

$$\Rightarrow \frac{\tau(\Xi_{cc}^{++})}{\tau(\Xi_{cc}^{+})} \sim 5.2\ ,\ \frac{\tau(\Xi_{cc}^{++})}{\tau(\Omega_{cc}^{+})} \sim 3.5\ ,\ \frac{\tau(\Omega_{cc}^{+})}{\tau(\Xi_{cc}^{+})} \sim 1.5 \tag{136}$$

and

$$BR_{SL}(\Xi_{cc}^{++}) = 15.8\ \%\ ,\ BR_{SL}(\Xi_{cc}^{+}) = 3.3\ \% \tag{137}$$

$$BR_{SL}(\Omega_{cc}^{+}) = 13.7\ \% \tag{138}$$

Since the semileptonic widths are basically equal for Ξ^{++}_{cc} and Ξ^{+}_{cc}, the ratio of their semileptonic branching ratios reflects the ratio of their lifetimes. On the other hand constructive PI enhances the semileptonic Ω^{+}_{cc} width.

Needless to say, there are substantial differences in these numbers, less so (as expected) for the ratios. For proper evaluation one has to take notice of the following. The authors of Ref.[264] use one value for baryons and a higher one for mesons for phenomenological reasons ([19]). Yet on theoretical grounds this is inadmissible: in the OPE one has to use the same value of m_c for mesons and baryons alike. Using different values can serve only as a *temporary crutch* to parameterize an observed difference between baryons and mesons one does not understand at all!

With neither PI nor WS contributing to Ω^{++}_{ccc} decays, its width is given by the decay of its three charm quarks. Insisting on using the same value for m_c as for D mesons, one would predict roughly the following numbers:

$$\tau(\Xi^{++}_{cc}) = 0.35\ ps\ ,\ \tau(\Xi^{+}_{cc}) = 0.07\ ps \tag{139}$$

$$\tau(\Omega^{+}_{cc}) = 0.10\ ps\ ,\ \tau(\Omega^{++}_{ccc}) = 0.14\ ps \tag{140}$$

Obviously one has to allow for considerable uncertainties in all these predictions due to the unknown higher order $1/m_c$ contributions and our ignorance about the potential controlling the inner dynamics of $C = 2$ baryons.

Yet a certain pattern does emerge, and one can conclude the following:

- It is a very considerable stretch to come up with a Ξ^{+}_{cc} lifetime as short as 0.03 ps.
- The Ξ^{++}_{cc} lifetime is similar to or even larger than that for D^0. There appears no way to push the Ξ^{++}_{cc} lifetime into the "ultrashort" domain ~ 0.1 ps.
- *If* the data forced upon us a scenario with $\tau(\Xi^{+}_{cc})$ well below 0.1 ps and $\tau(\Xi^{++}_{cc})$ near it – let alone below it –, we had to conclude that the successes of the HQE description of $C = 1$ charm hadron lifetimes listed above are quite accidental at least for the $C = 1$ baryons, but probably for the mesons as well. While this is a conceivable outcome, it would come with a hefty price for theory, or at least for some theorists. The HQE offers no argument why $C = 1$ hadrons are treatable, while $C = 2$ are not. Furthermore the double-heavy meson B_c appears to be well described by the HQE [240].

7. – Leptonic and Rare Decays

The simplest final state possible in D decays consists of a lepton or a photon pair, namely

(i) $D^+_q \to \tau^+\nu$, $\mu^+\nu$, $e^+\nu$;
(ii) $D^0 \to \mu^+\mu^-$, e^+e^-, $D^0 \to \gamma\gamma$;
(iii) $D^0 \to e^\pm\mu^\mp$.

While the first two transitions can proceed in the SM, albeit at reduced or even highly suppressed rates, the last one is absolutely forbidden. Transitions (ii) and (iii) thus represent clean, yet quite speculative searches for New Physics. The main motivation for

([19])It is related to the fact mentioned above that with $m_{c,kin}(1\ \text{GeV}) \simeq 1.2 - 1.3$ GeV one fails to reproduce the observed $\Gamma_{SL}(D)$.

measuring accurately transitions (i) on the other hand is to extract the decay constants f_D and f_{D_S} as one test of our theoretical control over nonperturbative QCD.

There is another simple, yet highly exotic final state, namely

(iv) $D^+ \to \pi^+ f^0$, $K^+ f^0$

where f^0 denotes the so-called familon, a neutral scalar that could arise as a Nambu-Goldstone boson in the spontaneous breaking of a continuous global family symmetry. Although it is not a leptonic final state, we will briefly discuss it at the end of this Section. We will discuss also rare threebody modes

(v) $D^+, D_s^+ \to h^+ l^- l^+,\ h^- l^+ l^+,\ h^+ e^\pm \mu^\mp$ (with $h = \pi, K$ and $l = e, \mu$),

which have some favourable experimental signatures; the theoretical interpretation is quite different for the different modes, as explained later.

Diagrams for process (i) to (iv) are shown in Fig. 23.

7˙1. *Expectations on* $D_q^+ \to \ell^+ \nu$. – These transitions are driven by the axial vector component of the hadronic charged current; the transition amplitude reads:

$$\mathcal{T}(D_q^+ \to \ell^+ \nu) = \frac{G_F}{\sqrt{2}} \langle 0|A^\mu|D_q\rangle [\bar{\ell}\gamma_\mu(1-\gamma_5)\nu_\ell],\ q = d, s \tag{141}$$

The hadronic matrix element is conventionally parametrized by the meson decay constant:

$$\langle 0|A^\mu|D_q(p)\rangle = i f_{D_q} p^\mu \tag{142}$$

The total width is then given by

$$\Gamma(D_q^+ \to \ell^+ \nu_l) = \frac{G_F^2}{8\pi} f_{D_q}^2 |V(cq)|^2 m_\ell^2 (1 - \frac{m_\ell^2}{M_{D_q}^2})^2 M_{D_q} \tag{143}$$

These transitions are *helicity suppressed* ; i.e., the amplitude is proportional to m_ℓ, the mass of the lepton ℓ, in complete analogy to $\pi^+ \to \ell^+ \nu$. This property holds in general for any (axial)vector charged current – a point we will return to below. They are also suppressed by $f_{D_q} \ll m_c$. This feature can be understood intuitively: due to the practically zero range of the weak forces the c and $\bar{s}$ or $\bar{d}$ quarks have to come together to annihilate. The amplitude therefore is proportional to $\psi_{c\bar{q}}(0)$, the $c\bar{q}$ wave function at zero separation. This quantity is related to the decay constant as follows: $f_{D_q}^2 = 12|\psi_{c\bar{q}}(0)|^2/M_{D_q}$ [20].

One expects the following branching ratios for the D_s^+ and D^+ modes, which are Cabibbo allowed and forbidden, respectively:

$$\mathrm{BR}(D^+ \to \tau^+ \nu) = 1.0 \cdot 10^{-3} \left(\frac{f_D}{220\ \mathrm{MeV}}\right)^2 \tag{144}$$

$$\mathrm{BR}(D^+ \to \mu^+ \nu) = 4.6 \cdot 10^{-4} \left(\frac{f_D}{220\ \mathrm{MeV}}\right)^2 \tag{145}$$

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[20] One should note that the description through f_{D_q} as defined by Eq.(142) holds irrespective of the existence of a wave function.

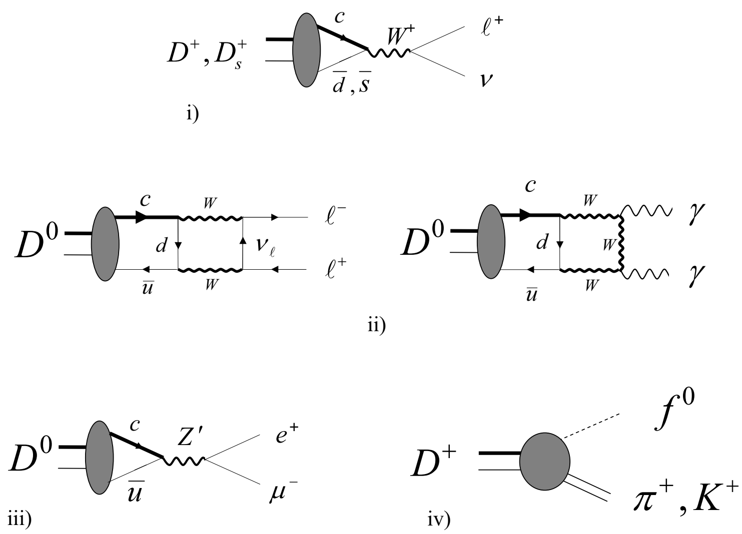

Fig. 23. – Rare and leptonic decay diagrams.

$$\mathbf{BR}(D^+ \to e^+\nu) = 1.07 \cdot 10^{-8} \left(\frac{f_D}{220 \text{ MeV}}\right)^2 \tag{146}$$

$$\mathbf{BR}(D_s^+ \to \tau^+\nu) = 4.5 \cdot 10^{-2} \left(\frac{f_{D_s}}{250 \text{ MeV}}\right)^2 \tag{147}$$

$$\mathbf{BR}(D_s^+ \to \mu^+\nu) = 5.0 \cdot 10^{-3} \left(\frac{f_{D_s}}{250 \text{ MeV}}\right)^2 \tag{148}$$

$$\mathbf{BR}(D_s^+ \to e^+\nu) = 1.2 \cdot 10^{-7} \left(\frac{f_{D_s}}{250 \text{ MeV}}\right)^2 \tag{149}$$

The numbers have been scaled to the lattice estimates of the decay constants.

The fact that there are two different mesons decaying in this way represents an extra bonus. Only lattice QCD could provide us with predictions for the decay constants f_D and f_{D_s} with a theoretical uncertainty not exceeding a few percent. Yet one has to be concerned that maybe not all lattice artefacts can be brought under such excellent control. These systematics should however basically cancel in the ratio f_{D_s}/f_D, which deviates from unity only due to $SU(3)_{Fl}$ breaking.

The motivation for measuring these widths thus lies in extracting the decay constants as a precision measure for the theoretical control lattice QCD can achieve over certain hadronic quantities. Even more importantly lattice QCD successfully passing this test constitutes an important validation for its prediction on $f_{B_{d,s}}$. It has been suggested that even if lattice QCD falls somewhat short of a successful prediction for f_D, one can use the experimentally observed value for f_D together with lattice QCD's prediction for f_B/f_D to arrive at a reliable value for f_B, which obviously is of the greatest benefit for interpreting $B^0 - \bar{B}^0$ oscillations; at the very least success in reproducing f_{D_s}/f_D would tightly constrain f_{B_s}/f_B, which controls $B_s - \bar{B}_s$ vs. $B_d - \bar{B}_d$ oscillations.

These transitions are sometimes advanced as sensitive probes for New Physics in the form of hard pseudoscalar couplings for the charged current, which would modify the ratio $\Gamma(D_q \to \tau\nu)/\Gamma(D_q \to \mu\nu)$ with respect to the helicity factor $(m_\tau/m_\mu)^2$ (plus phase space corrections), see Eq.(143). This is an echo of what happened in the early days of weak interaction studies when the observation of $\Gamma(\pi^+ \to \mu^+\nu) \gg \Gamma(\pi^+ \to e^+\nu)$ pointed to the dominance of spin one over spin zero charged currents.

While such couplings are conceivable on purely phenomenological grounds, one usually views them as somewhat unlikely due to general theoretical arguments. For a spin-zero coupling to the $\ell\nu$ pair independent of the lepton mass m_ℓ, would represent a 'hard' breaking of chiral invariance. It would be surprising if such a feature had not been noted before. There could be Higgs couplings, yet those would again be proportional to m_ℓ (in amplitude) and thus represent an acceptable soft breaking of chiral symmetry. Of course, the dimensionless numerical coefficient in front of such a coupling could differ for the different modes.

7·2. $D_q^+ \to \ell^+\nu$. – Looking at the estimates for the branching ratios one realizes that going after $D_q \to e\nu$ is a hopeless enterprise; yet at the same time observing it would provide spectacular evidence for New Physics. The final state $\tau\nu$ has the largest branching ratios, yet poses the highly nontrivial problem that the τ lepton decays typically into one-prong final states with additional neutrinos. It then makes sense both from an experimental as well as theoretical perspective to make a dedicated effort to measure all four channels $D_q \to (\tau/\mu)\nu$ as precisely as possible.

Small decay rate and difficult event topology make these decays a real challenge to the experimenter (Tab.VII). Good lepton identification is required, as well as acceptable tracking capability, and detector hermeticity to evaluate the neutrino missing energy which is very large in the case of the $D_s^+ \to \tau^+\nu_\tau, \tau^+ \to \mu^+\nu_\mu\bar{\nu}_\tau$ decay chain. At fixed-target experiments, the challenge has been so far overwhelming. Strategies to attempt detecting the $D^+ \to \mu^+\nu_\mu$ decay by tracking the D^+ before its decay have been proposed [265], but the practical realization of the method has encountered a formidable show-stopper in unreducible backgrounds processes, such as $D^+ \to \pi^0\mu^+\nu_\mu$ with the undetected neutral pion.

Model-dependent determinations based on isospin splittings and hadronic decays have also been proposed [266].

For f_{D^+}, a measurement is available from BES, based on one event[267] $f_D^+ = 300 \pm 175$ MeV from a branching ratio value of $BR(D^+ \to \mu^+\nu_\mu) = 0.08^{+0.17}_{-0.05}\%$. The most recent and most precise measurement of $f_{D_s^+} = 285 \pm 19 \pm 40$ MeV is from ALEPH [268], which combines $\mu^+\nu_\mu$ and $\tau^+\nu_\tau$ events using relative production rates taken from lattice computations. For a summary of the experimental scenario see [269, 270, 271], and for comparison with pion and kaon constants see [272].

Great expectations in this field come from CLEO-c and, later, from BES III, where a projected yield of about 1000 reconstructed leptonic decays should allow a 2% error on f_{D,D_s} [271].

7·3. *Adagio, ma non troppo.* – Inclusive or exclusive flavour-changing radiative decays – $D \to \gamma X$ or $D \to \gamma K^*/\rho/\omega/\phi$, respectively – reflect very different dynamics than $B \to \gamma X$ or $B \to \gamma K^*/\rho/\omega/\phi$. While one can *draw* Penguin *diagrams* in both cases, their meaning is quite different. For the leading transition operator in radiative B decays arises from 'integrating out' the top quark in the Penguin loop. On the other hand in the Penguin diagram for $c \to u$ the contribution from an internal b quark is negligible due to the almost decoupling of the third quark family from the first two; the contribution from an internal s quark cannot be treated as a *local* operator, since the s quark is lighter than the external c quark and thus cannot be integrated out. Thus even the inclusive rate $D \to \gamma X$ is not controlled by short-distance dynamics within the SM, let alone exclusive transitions $D \to \gamma V$ with V denoting a vector meson. This can be illustrated through diagrams where weak annihilation is preceded by photon emission from the light (anti)quark line.

The discerning observer can however see a virtue in this complexity. Measuring $D \to \gamma K^*/\rho/\omega/\phi$ can teach us new lessons on QCD's nonperturbative dynamics in general, provide another test challenging lattice QCD's numerical reliability and thus be of considerable value for 'peeling off' layers of long-distance dynamics from the measured rates for $B \to \gamma K^*/\rho/\omega$ when extracting $V(td)/V(ts)$ from the latter. Complementary information can be inferred from $B \to \gamma D^*$, which receives no genuine Penguin contribution either.

Rather detailed analyses exist [273, 274] concerning SM expectations, which – not surprisingly based on what was said above – represent order-of-magnitude predictions only. Typical numbers are

$$\begin{aligned} &\mathrm{BR}(D^0 \to \gamma \bar{K}^{*0}) = (6 \div 36) \cdot 10^{-5} \ , \ \mathrm{BR}(D^0 \to \gamma \rho^0) = (0.1 \div 1) \cdot 10^{-5} \\ &\mathrm{BR}(D^0 \to \gamma \omega) = (0.1 \div 0.9) \cdot 10^{-5} \ , \ \mathrm{BR}(D^0 \to \gamma \phi) = (0.1 \div 3.4) \cdot 10^{-5} \end{aligned} \tag{150}$$

A more speculative motivation to search for these modes is that in some nonminimal SUSY scenarios induce significant contributions due to local Penguin operators for $c \to u$ with SUSY fields appearing in the internal loop [275]. Since they affect neither $D \to \gamma K^*$ nor $D \to \gamma \phi$ those two rates provide a good calibrator for $D \to \gamma \rho$ and $D \to \gamma \omega$, since several systematic uncertainties will largely drop out from the ratios $\Gamma(D^0 \to \gamma \rho/\omega)/\Gamma(D^0 \to \gamma K^*/\phi)$.

BELLE has reported the first observation of any of these modes [276]:

$$\mathrm{BR}(D^0 \to \gamma \phi) = (2.6^{+0.70}_{-0.61}(stat.)^{+0.15}_{-0.17}(syst.)) \cdot 10^{-5} \tag{151}$$

which is consistent with expectations; however the latter cover a particularly wide range.

7·4. *Much rarer still:* $D^0 \to \mu^+\mu^-$ *and* $D^0 \to \gamma\gamma$. – The two effects suppressing $D^+ \to \mu^+\nu$ do likewise in $D^0 \to \mu^+\mu^-$: (i) Any spin-one coupling of the type $[\bar{u}\gamma_\mu\gamma_5 c][\bar{\mu}\gamma_\mu(1-\gamma_5)\mu]$ leads to helicity suppression, i.e. a transition amplitude proportional to m_μ. (ii) Due to the effective zero range also of this coupling the amplitude is suppressed by $\sim f_D/m_c \ll 1$. On top of that one needs a genuine flavour-changing neutral current: (iii) Within the SM such a current does not exist on the tree level – it has to be generated on the one-loop level thus representing a pure quantum effect.

All three effects combine to produce a huge reduction in branching ratio. An extremely crude guestimate invokes BR$(D^+ \to \mu^+\nu)$ to mirror helicity and wavefunction suppression and uses $\frac{\alpha_S}{\pi} \cdot (m_s/M_W)^2$ to reflect the second order GIM effect; i.e.

$$\mathrm{BR}(D^0 \to \mu^+\mu^-) \sim \mathcal{O}\left(\mathrm{BR}(D^+ \to \mu^+\nu) \cdot \frac{\alpha_S}{\pi} \cdot \frac{m_s^2}{M_W^2}\right) \sim \mathcal{O}(10^{-12}) \tag{152}$$

It is straightforward to draw one-loop diagrams describing $\Delta C = 1$ neutral currents and computing them as if they represent a short-distance effect. However there is no a priori reason why $D^0 \to \mu^+\mu^-$ should be controled by simple short-distance dynamics. A detailed treatment yields [277]

$$\mathbf{BR}(D^0 \to \mu^+\mu^-) \simeq \mathbf{3.0 \cdot 10^{-13}} \tag{153}$$

orders of magnitude below the present experimental upper bound:

$$\mathbf{BR}(D^0 \to \mu^+\mu^-)|_{exp} \leq \mathbf{4.1 \cdot 10^{-6}}\ . \tag{154}$$

The authors have employed various prescriptions for estimating possible long distance effects analoguous to what one encounters in $K_L \to \mu^+\mu^-$. Similar to what happens there they actually find that a two-step transition involving long-distance dynamics provides a large or even dominant contribution, namely

$$\mathbf{BR}(D^0 \to \gamma\gamma \to \mu^+\mu^-) \sim \mathbf{2.7 \cdot 10^{-5}} \cdot \mathbf{BR}(D^0 \to \gamma\gamma) \sim (\mathbf{0.3} - \mathbf{1}) \cdot \mathbf{10^{-12}} \tag{155}$$

$$\mathbf{BR}(D^0 \to \gamma\gamma) \sim (\mathbf{1} - \mathbf{3.5}) \cdot \mathbf{10^{-8}} \tag{156}$$

consistent with the estimates of Ref.[278]. The theoretical tools exist for a refined analysis based on the OPE that includes long distance dynamics naturally and self-consistently through quark condensates. Yet that would appear to be a purely academic exercise in view of the immensely tiny branching ratio. For it is unlikely that such a refinement could enhance the branching ratio by three – let alone more – orders of magnitude. Later in discussing $D^0 - \bar{D}^0$ oscillations we will address analoguous issues, where they are much more relevant numerically.

With the SM width so tiny, it could be a promising laboratory to search for manifestations of New Physics. The authors of Ref.([277]) find that New Physics scenarios can produce a wide range in predictions, namely

$$\mathbf{BR}(D^0 \to \mu^+\mu^-)|_{NP} \sim \mathbf{10^{-11}/8 \cdot 10^{-8}/3.5 \cdot 10^{-6}} \tag{157}$$

for models with a superheavy b' quark/ multi-Higgs sector/ SUSY with R parity breaking, respectively.

Since the mode $D^0 \to \mu^+\mu^-$ possesses a clear signature, one can entertain the hope to search for it in hadronic collisions. The best limit in PDG02 actually dates back to 1997 hadroproduction experiment BEATRICE (WA92 at CERN) [279].

CDF has plans to pursue the study of this decay, as has BTeV. CDF showed at Moriond 2003 preliminary results [66, 67] which lower the WA92 best limit by a factor of two (Tab. VII. About the same sensitivity is expected from B-factories at the present level of data collected (about 90 fb^{-1}) [280].

The much more challenging decay $D^0 \to \gamma\gamma$ requires superb em calorimetry and high statistics. CLEO possesses both, and recently presented a first limit [281]. Lower limits are expected from BABAR and BELLE [282].

7.5. *The "forbidden" mode:* $D^0 \to e^{\pm}\mu^{\mp}$. – A transition $D^0 \to e^{\pm}\mu^{\mp}$ violates lepton number conservation ; observing it manifests unequivocally the intervention of New Physics. No signal has been observed so far:

$$\mathrm{BR}(D^0 \to \mu^+ e^-)|_{exp} \leq 8.1 \cdot 10^{-6} \tag{158}$$

Yet as before for $D^0 \to \mu^+\mu^-$ this rate is suppressed by the simultaneous factors $(f_D/m_c)^2 \sim 0.04$ and $(m_\mu/m_c)^2 \sim 0.007$ with the latter arising at least for a spin-1 component in the underlying $c\bar{u}$-coupling.

There are classes of New Physics scenarios that can induce a signal here, namely models with Technicolor, nonminimal Higgs dynamics, heavy neutrinos, horizontal gauge interactions and SUSY with R parity breaking. The last one again provides the most promising – or least discouraging – case allowing for [277]

$$\mathrm{BR}(D^0 \to \mu^+ e^-)|_{SUSY,\not{R}} \leq 1.0 \cdot 10^{-6} \tag{159}$$

One might be inclined to view such searches as bad cases of 'ambulance chasing'. Yet one should keep in mind that the structure and strength of flavour-changing neutral currents could be quite different in the up- and down-type quark sectors. As far as the latter is concerned, very sensitive searches have been and will continue to be performed in K and B decays[283]. Yet the u, c and t quarks provide very different search scenarios: such lepton-number violating flavour-changing neutral currents for u quarks can be probed via $\mu - e$ conversion in deep-inelastic scattering – $\mu N \to eX$ – where the system X might or might not contain a charm hadron; top states on the other hand decay as quarks before they can hadronize [9] and thus present different challenges and promises in such searches. In any case a careful study of D decays is thus complementary to analyses of $\mu - e$ conversion and top decays rather than a repeat effort.

7.6. *Exotic new physics:* $D^+ \to \pi^+/K^+ f^0$. – One of the central mysteries of the SM is the replication of quark-lepton families, which furthermore exhibit a quite peculiar pattern in the masses and CKM parameters. This structure might be connected with the existence of some continuous 'horizontal' or family symmetry that has to be broken. If it is a *global* symmetry broken *spontaneously*, very light neutral Nambu-Goldstone bosons have to exist, the 'familons' [284]. They might actually provide some valuable service by en passant solving the 'strong CP' problem of the SM ([21]).

The effective low-energy interaction of the familon fields f^a with fermion fields ψ_i is given by a non-renormalizable derivative coupling of the form

$$\mathcal{L}_F = \frac{1}{\Lambda_F}\bar{\psi}_i\gamma_\mu(V^a_{ij} + A^a_{ij}\gamma_5)\psi_j\partial_\mu f^a \; ; \tag{160}$$

([21])The 'strong CP problem' refers to the realization that a priori the T-odd operator $G \cdot \tilde{G}$ can appear in the QCD Lagrangian with a coefficient roughly of order unity. Yet bounds on the electric dipole moment of neutrons shows this coefficient to be at most of order 10^{-9}. Peccei-Quinn type symmetries have been invoked to provide a natural solution to this apparent fine-tuning problem. They imply the existence of axions that have turned out to be elusive so far [107]; familons can be their flavour-nondiagonal partners.

the mass scale Λ_F calibrates the strength of these effective forces, while V^a and A^a denote matrices containing the vector and axial vector couplings of the familons to the different fermions.

The analogous decays $K^\pm \to \pi^\pm f^0$ [285], $B^\pm \to \pi^\pm / K^\pm f^0$ and $B_d \to K_S f^0$ [286] have been searched for with no signal found yielding $\Lambda_F^{(2)} \geq 10^{11}$ GeV and $\Lambda_F^{(3)} \geq 10^8$ GeV for the second and third family, respectively.

Surprisingly enough, no bounds have been given so far for familons in D decays. This unsatisfactory situation (familon couplings could a priori be quite different for up- and down-type quarks; searches for $D \to hf$ and $B \to hf$ are thus complementary to each other) should be remedied by BELLE, BABAR and CLEO-c ([22]). In the K sector a new line of attack can be opened by KLOE at DAΦNE ϕ-factory searching for $K_S^0 \to \pi^0 f^0$.

7·7. $D^+, D_s^+ \to h\ell\ell'$ *with* $h = \pi, K$. – Final states with a charged hadron and two charged leptons have in addition to their favourable experimental signatures some phenomenological advantages as well over $D \to l\nu$, ll: their amplitudes suffer from neither helicity nor wavefunction suppression; i.e., the small factors m_l/m_c and f_D/m_c are in general absent. Yet beyond that they represent very different dynamical scenarios.

$D^+, D_s^+ \to h^+\mu^+e^-$ and $D^+, D_s^+ \to h^+e^+\mu^-$ require genuine flavour changing neutral currents and thus would represent unequivocal signals for New Physics; both would actually be natural modes for different variants of 'horizontal' or leptoquark interactions coupling first and second family quarks and leptons to each other.

The modes $D^+, D_s^+ \to h^-l^+l^+,\ h^-e^+\mu^+$ also require the intervention of New Physics; yet from a theoretical rather than purely phenomenological perspective they appear to represent a rather contrived scenario.

Finally the rare transitions $D^+, D_s^+ \to h^-l^+l^-$ could certainly be affected considerably by New Physics; unfortunately they can proceed also within the SM with a strength that cannot be predicted reliably since there they are driven mainly by long distance dynamics [277, 287, 288], which are notoriously difficult to calculate. Recent FOCUS results[289] substancially lower previous limits [290, 291] (Tab.VII) and already allow one to exclude SUSY models with R-Parity breaking [277]. As a technical detail, the FOCUS limit makes use of a new bootstrap-based technique to ascertain the limit[232] confidence level, which has proved to provide a better *coverage* than the standard limit technique of Ref. [292].

8. – Semileptonic Decays

There is no honourable excuse for theory to duck its responsibility for treating semileptonic charm decays. With seven $C = 1$ hadrons decaying weakly – and doing that at the Cabibbo allowed as well as once forbidden level – one is certainly facing a complex challenge here. Yet there exists strong motivation for a dedicated analysis of semileptonic charm decays: while there is little hope for New Physics to surface there, they constitute a clean laboratory for probing our quantitative understanding of QCD; this has both intrinsic merit and serves as preparation for coming to grips with hadronization in the beauty sector. En passant one might learn something novel about light flavour spectroscopy as well through a careful analysis of final states like $D_{(s)}^+ \to (\eta/\eta')\ell\nu$

([22])Top quark decays could produce familons as well.

Decay mode		BR %	Decay mode		BR %
$D^+ \to e^+\nu_\mu$		—	$D_s^+ \to e^+\nu_\mu$		—
$D^+ \to \mu^+\nu_\mu$	PDG02	$0.08^{+0.17}_{-0.05}$	$D_s^+ \to \mu^+\nu_\mu$	PDG02	0.50 ± 0.19
$\Rightarrow f_{D^+} = 177 \div 513$ MeV			$\Rightarrow f_{D_s^+} = 250 \pm 50$ MeV		
$D^+ \to \tau^+\nu_\tau$		—	$D_s^+ \to \tau^+\nu_\tau$	PDG02	6.4 ± 1.5
			$\Rightarrow f_{D_s^+} = 298 \pm 37$ MeV		

Decay mode		BR 10^{-6}	Decay mode		BR 10^{-6}
$D^0 \to \mu^+\mu^-$	PDG02	< 4.1	$D^0 \to \mu^+e^-$	PDG02	< 8.1
	[66, 67]	< 2.4			
$D^0 \to e^+e^-$	PDG02	< 6.2			
$D^0 \to \gamma\gamma$	[281]	< 29			
$D^+ \to K^+\mu^-\mu^+$	PDG02	< 44	$D^+ \to K^-\mu^+\mu^+$	PDG02	< 120
	[289]	< 9.2		[289]	< 13
$D^+ \to \pi^+\mu^-\mu^+$	PDG02	< 15	$D^+ \to \pi^-\mu^+\mu^+$	PDG02	< 17
	[289]	< 8.8		[289]	< 4.8
$D_s^+ \to K^+\mu^-\mu^+$	PDG02	< 140	$D_s^+ \to K^-\mu^+\mu^+$	PDG02	< 180
	[289]	< 36		[289]	< 13
$D_s^+ \to \pi^+\mu^-\mu^+$	PDG02	< 140	$D_s^+ \to \pi^-\mu^+\mu^+$	PDG02	< 82
	[289]	< 26		[289]	< 29

TABLE VII. – *Experimental branching ratios for charm mesons leptonic and rare decays, and pseudoscalar charm meson decay constants. World averages from PDG02 [131]. All limits are 90% cl. Statistical and systematical errors added in quadrature.*

as explained later. There are several features facilitating a theoretical description and making it more interesting at the same time:

- On each Cabibbo level there is only a single quark level transition operator.
- Factorization holds trivially ; i.e., the amplitudes for semileptonic transitions are expressed through the product of a leptonic current and the matrix element of a hadronic current. The latter is given by an expression bilinear in quark fields. While we have not (yet) established full control over quark bilinears, we know a lot about them and in any case considerably more than about the matrix elements of four-quark operators.
- On general, though somewhat handwaving grounds one expects semileptonic charm widths to be dominated by a handful of exclusive modes. Studying how the inclusive width is saturated by exclusive width will provide us with some instructive lessons about how duality can emerge at relatively low scales.
- They allow direct access to $|V(cs)|$ and $|V(cd)|$.
- They constitute a laboratory for studying dynamics similar to those shaping semileptonic beauty decays.
- They provide an important bridge between light flavour and heavy flavour dynamics that can be studied by approaching them from higher mass scales using heavy quark expansions and from lower ones directly through lattice QCD.

8˙1. *Inclusive Transitions.* – In Sect.**6**˙4 we have already discussed the *inclusive* semileptonic branching ratios for the various charm hadrons, since they fit naturally into the treatment of weak lifetimes. In our discussion there we have focussed on *ratios*, since there $|V(cs)|$ as well as the leading term $\propto m_c^5$ drop out.

There is still a fly in the ointment. For one might become emboldened to predict also the absolute size of $\Gamma_{SL}(D)$. The $1/m_Q$ expansion as described in Sect.**6**˙4 yields

$$\Gamma_{\mathrm{SL}}(D \to \ell\nu X) = \frac{G_F^2\, m_c^5}{192\,\pi^3}\, |V_{cq}|^2 \left[z_0(r) \left(1 - \frac{\mu_\pi^2 - \mu_G^2}{2m_c^2} \right) - 2(1-r)^4 \frac{\mu_G^2}{m_c^2} + \mathcal{O}(\alpha_{em}, \alpha_s, \frac{1}{m_c^3}) \right], \tag{161}$$

where $z_0(r)$ is the tree-level phase space factor and $r = m_q^2/m_c^2$:

$$z_0(r) = 1 - 8r + 8r^3 - r^4 - 12r^2 \ln r \; . \tag{162}$$

Comparing the measured width with the theoretical expression using $|V(cs)|$ and $m_c \sim 1.3$ GeV, as inferred from charmonium spectroscopy, one finds that one can reproduce no more than two thirds of the observed $\Gamma_{SL}(D)$ with those input values. It has been conjectured that *non*factorizable contributions might make up the deficit.

Also the lepton energy *spectrum* of inclusive semileptonic charm hadron transitions can teach us valuable lessons. Even among mesons – with practically identical semileptonic widths – one expects significant differences in the spectra of D^0 and D_s^+ [D^+] decays on the $c \to s$ [$c \to d$] level due to WA contributions. While one can count on no more than a semiquantitative description for the integrated widths, even that would be too much to expect for the spectra with their need for 'smearing' [293, 294]. Nevertheless in the spirit of Yogi Berra's dicta ([23]) one can always start an observation of something by looking at it and thus search for evidence of WA in the endpoint spectra; if the observed pattern fits the expectations, then one can try to infer the D meson expectation value of the four-fermion operator; this quantity has both intrinsic and practical interest – the latter since the B meson analogue of this four-fermion operator affects the lepton energy endpoint spectra in B_d and B_u transitions *differently* [257]. OPAL has produced the first measurement [296, 297] of the semileptonic branching ratio averaged over all weakly decaying charm hadrons produced in $Z^0 \to c\bar{c}$ decays. Using the theoretical prediction for $\Gamma(Z^0 \to c\bar{c})$ it finds $B(c \to \ell) = 0.095 \pm 0.006^{+0.007}_{-0.006}$, in good agreement with both expectation based on the dominance of D^0 over D^+ abundance due to D^* production and with ARGUS' data at lower energies.

8˙2. *Exclusive Modes.* – There are three main reasons to study exclusive channels:

1. Exclusive semileptonic decays are easier to measure than inclusive ones, though more difficult to treat theoretically, since they are highly sensitive to long distance dynamics. One can turn this vice into a virtue, though: measuring these transition rates and comparing them with predictions will extend our knowledge and maybe

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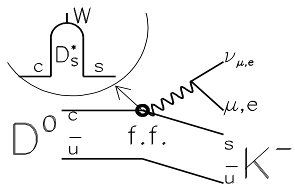
([23])Yogi Berra was an outstanding baseball player, manager and coach, who was voted onto the All-Century team. He is, however, better known as the founder of one of the most popular schools in US philosophy, namely Yogi-ism characterized by immortal quotes like [295]: "You can observe a lot by watching"; "It's deja vu all over again"; "When you come to a fork in the road ... take it"; "If the world were perfect, it wouldn't be"; "The future ain't what it used to be".



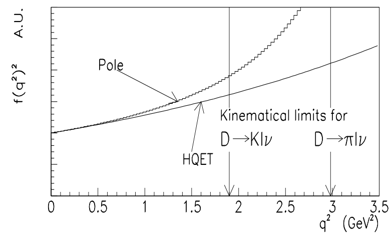


Fig. 24. – Semileptonic D^0 decay due to the coupling of a virtual $D_s^*(c\bar{s})$ vector state in the nearest-pole dominance model (left); $|f_+(q^2)|^2$ as a function of q^2 for pole and HQET parameterizations. The kinematic limits for $K\ell\nu$ and $\pi\ell\nu$ are drawn with vertical lines (adapted from [298]).

even understanding of nonperturbative dynamics. This is a worthy goal in itself – and of obvious benefit for understanding the corresponding decays of beauty hadrons. It is essential in such an endeavour to have a large body of different well measured modes rather than a few ones.

2. Studying how exclusive modes combine for different charm hadrons to saturate the inclusive widths will teach us valuable lessons on how and to which degree quark-hadron duality discussed in Sect.**4**˙11 is implemented in charm decays .

3. The final states of semileptonic charm decays can provide novel insights into light flavour spectroscopy.

8˙2.1. $H_c \to \ell\nu h$. The most tractable case arises when there is a single hadron or resonance h in the final state. Long distance dynamics can then enter only through hadronic form factors, which describe how the charm hadron is transformed into the daughter hadron at the hadron W vertex (Fig.24a) as a function of the momentum transfer. To be more specific: for D decays into a final state with a single pseudoscalar meson P, the matrix element of the hadronic weak current can be expressed in terms of two formfactors

$$\langle P(p')|J_\mu|D(p)\rangle = f_+(q^2)(p+p')_\mu + f_-(q^2)(p-p')_\mu \ , \quad q = p - p' \tag{163}$$

The differential decay rate is given by

$$\frac{d\Gamma(D \to \ell\nu P)}{dq^2} = \frac{G_F^2|V_{cq}|^2 p_P^3}{24\pi^3}\left\{|f_+(q^2)|^2 + |f_-(q^2)|^2\mathcal{O}(m_\ell^2) + ...\right\} \tag{164}$$

$$m_\ell^2 \le q^2 = M_D^2 + m_P^2 - 2M_D E_P \le 2M_D E_\ell + 2m_P^2 E_\ell(2E_\ell - M_D) \tag{165}$$

where p_P and E_P are the momentum and energy, respectively, of P in the rest frame of the charmed meson. The dependance on the second form factor f_- is proportional to the mass squared of the lepton and therefore insignificant in D decays. In decays leading to a vector resonance V – $D \to \ell\nu V$ – there are four formfactors often denoted by V, A_{1-3} and defined by the decomposition of the hadronic weak current:

$$\langle V(p'),\lambda|J_\mu|D(p)\rangle = -i(M_D+M_V)A_1(q^2)\epsilon_\mu^{*(\lambda)} + \frac{iA_2(q^2)}{M_D+M_V}(\epsilon_\mu^{*(\lambda)}\cdot p_D)(p+p')_\mu$$

$$+\frac{iA_3(q^2)}{M_D+M_V}(\epsilon^{*(\lambda)}_\mu \cdot p_D)(p-p')_\mu + \frac{2V(q^2)}{M_D+M_V}\epsilon_{\mu\nu\rho\sigma}\epsilon^{*(\lambda)}_\nu p_\rho p'_\sigma\ ; \tag{166}$$

λ denotes the polarization of V. The A_3 contribution is suppressed by the lepton mass and thus can be ignored here, similar to the f_- term in $D \to \ell\nu P$. The two salient features of the form factors are their q^2 dependence and their normalization at a given value of q^2; a typical, though not unique, choice is $q^2 = 0$. Three different strategies have been suggested for treating these formfactors theoretically:

1. One uses wavefunctions from a specific quark model to calculate the residue $f_\pm(0)$. For the q^2 dependance of $f_\pm(q^2)$ one assumes a parametrization consistent with the quark model used, like nearest pole dominance [299]:

$$f_\pm(q^2) = f_\pm(0)(1 - q^2/M^2_{pole})^{-1} \tag{167}$$

   In a literal application of this ansatz one chooses the value for M_{pole} as the mass of the nearest charm resonance with the same J^P as the hadronic weak current (Fig.24a). This can be an approximation only; one can allow for some flexibility here and fit M_{pole} to the data. A deviation of it from the mass of the nearest resonance indicates contributions from higher states.

   Another model in this class, which is sometimes (although not quite correctly) called the HQET form uses instead [101]

$$f_\pm(q^2) = f_\pm(0)e^{\alpha q^2} \tag{168}$$

   While the two functional dependences on q^2 in Eqs.(167) and (168) look quite different, there is hardly a difference in $D \to K\ell\nu$ due to the small range in q^2 kinematically allowed there (Fig.24b from ref.[298]). On the other hand, for the Cabibbo suppressed channel $D \to \pi\ell\nu$ there is considerable sensitivity to the differences between the two functions. Such approaches represent theoretical engineering. This is not meant as a put-down. For they are certainly useful: they train our intuition, provide tools for estimating detection efficiencies, yield insights into the strong dynamics and allow extrapolations to the corresponding B decays. Yet such lessons cannot be taken too literally on a quantitative level. For one cannot count on quark models to yield truly reliable error estimates or providing quantitative control over the extrapolations to beauty transitions.

2. Right after the emergence of HQET there had been considerable optimism about heavy quark symmetry allowing us to reliably evaluate the formfactors for $B \to \ell\nu\pi/\rho$, which are needed to extract $V(ub)$, by relating them to the ones for $D \to \ell\nu K$, which can be measured (using the known value of $V(cs)$, see our discussion below). Those hopes have largely faded away since unknown corrections of order $1/m_c$ will be quite sizeable. Heavy quark concepts are still useful, yet have to be complemented with other theoretical technologies. One such technology that – unlike quark models – is directly based on QCD and its field theoretical features employs sum rules similar to those introduced by Shifman, Vainshtein and Zakharov and discussed in 4˙10.2. The particular variant used is referred to as light cone sum rules [102]. Ref.[300] finds

$$\mathrm{BR}(D^0 \to e^+\nu\pi^-) = (1.6 \pm 0.34)\cdot 10^{-3} \tag{169}$$

(170) $$\mathbf{BR}(D^0 \to e^+\nu\rho^-) = (4.8 \pm 1.4) \cdot 10^{-3}$$

The prediction of Eq.(169) underestimates the data very significantly:

(171) $$\mathbf{BR}(D^0 \to e^+\nu\pi^-) = (3.6 \pm 0.6) \cdot 10^{-3}$$

The discrepancy appears much larger than the customary 20 - 30 % theoretical uncertainty one has to allow for QCD sum rule as explained in Sect.4˙10.2. A possible explanation for this failure is that the charm mass is too low a scale for the leading terms to provide an accurate description; such reasoning could not be invoked, though, for a similar failure in B decays.

3. Both the shapes and normalizations of the various form factors listed above can be computed based on lattice simulations. Several lattice groups have presented what can be viewed as pilot studies of the differential decay width for processes like $D \to \pi\ell\nu$ in the *quenched* approximation.

   In [301] a lattice version of HQET is implemented to determine the relevent form factors at different pion energies. Many sources of error arise in this analysis leading to statistical and systematic uncertainties of $10 - 20$%. The integrated quantity

(172) $$T_D\left(p_{min}, p_{max}\right) = \int_{p_{min}}^{p_{max}} dp\ p^4 \left|f_+(E)\right|^2/E$$

   can be combined with experimental measurements to obtain V_{cd} using

(173) $$|V_{cd}| = \frac{12\pi^3}{G_F^2 m_D} \frac{1}{T_D\left(p_{min}, p_{max}\right)} \int_{p_{min}}^{p_{max}} dp \frac{d\Gamma_{D\to\pi}}{dp}.$$

   Here p is the magnitude of the pion three-momentum in the D meson rest frame. The lower limit p_{min} is chosen to minimize extrapolation uncertainties in p and the light quark mass. The upper limit is chosen to reduce statistical and discretization uncertainties. With uncertainties of 10-20% in T_B, $|V_{cd}|$ can only be determined at the 15% level using this technique in the quenched approximation.

   Other groups have made similar quenched determinations using other techniques. Reference [302] uses a lattice version of NRQCD, whereas references [303, 304] rely on fermions with light quark normalizations. Some of these results are shown in Fig. 25 demonstrating the agreement between the differing techniques and also displaying the overall size of the uncertainties.

   The high precision group[97] has included semi-leptonic D meson decays in their list of the "golden-plated" modes for which the lattice can provide highly accurate values. The claim is that with *un*quenched lattice simulations the overall uncertainty on the semi-leptonic form factors can be reduced to level of 1-3%, which represents a noble goal. There is the promise that this theoretical accuracy level will be reached on the same time scale as CLEO-c will provide its precise measurements or even before.

All charm experiments have dedicated programs for studying semileptonic decays (for extensive reviews see Refs.[269, 298]). To measure q^2 in exclusive decays one needs

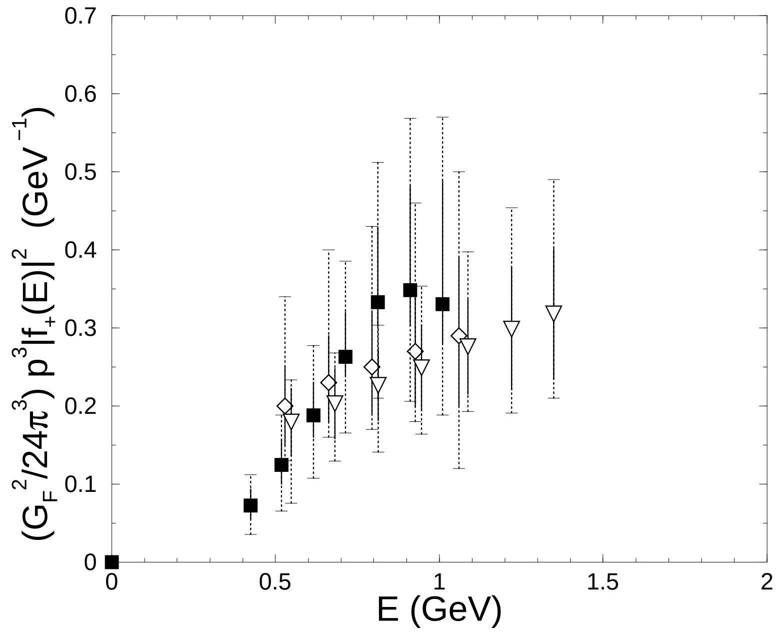


Fig. 25. – A comparison of various lattice predictions for the semi-leptonic differential decay width as a function of the pion energy E: diamonds [303],triangles [304], and squares [301]. This plot is taken from reference [301].

to identify all decay products, determine the H_c rest frame, and compute q^2 via Eq.(165). The determination of this rest frame is not unambiguous due to the neutrino escaping detection. In experiments with good tracking resolution (typically with a fixed-target geometry) the H_c direction can be measured, and the rest frame determined up to a quadratic ambiguity. At colliders with an hermetic detector, one can study q^2 inclusively by estimating the missing neutrino energy and momentum. A must is the capability of selecting semileptonic events by identifying the muon, the electron, or both.

All charm experiments have contributed to the measurement of semileptonic D decays branching ratios, and most of them of formfactors, thanks to generally good tracking, vertexing, and lepton identification capabilities. Results new to PDG02 (Tab.VIII shows 2003 update to [131]) have come from FOCUS [305, 306] and CLEO [307] (see also the recent reviews [270, 308, 282]).

A long-standing issue is in $BR(D^+ \to K^*\ell\nu)$ relative to $K\pi\pi$, where in the electron channel the 2002 CLEO measurement is several σ higher than the 1989 E691 value. The PDG03 average 0.61 ± 0.07 contains a 1.7 error scale factor to reconcile CLEO and E691 incompatible measurements. Quark models predict values close to CLEO's (see [298] for a review). The 2002 FOCUS semimuonic branching ratio is $0.602 \pm .022$ — higher that E691, but about 1σ lower than CLEO.

Most of the efforts have been concentrated on the precise measurement of formfactor ratios (Fig. 26).

The measurements of formfactor ratios have generated in the past considerable controversy, with three measurements (E687, E653, E691) in mild disagreement, albeit with

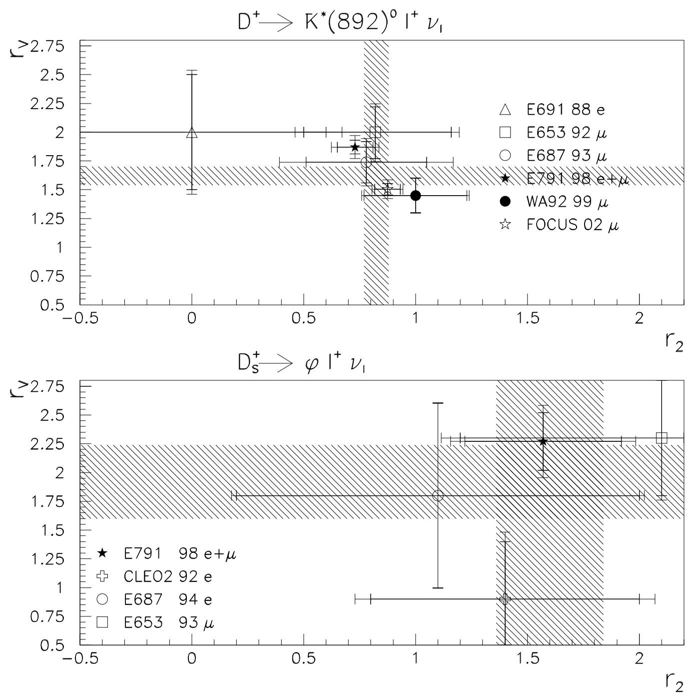


Fig. 26. – Compilation of experimental results on semileptonic vector formfactor ratios. Error bars show the statistical error, and the statistical and systematic added in quadrature. PDG03 updated averages to 2002 [131] are shown in the boxed area. PDG03 applies a 1.5 factor on the error of r_V world average for $D^+ \to \bar{K}^{*0}\ell^+\nu_\ell$.

large statistical errors (Fig.26). A new measurement by FOCUS[305], which collected large samples of clean $D^+ \to \bar{K}^{*0}\mu^+\nu_\mu$ (fifty-thousand events), determines r_V, r_2 with a few-percent error, and confirms the WA92 measurement. SU(3) flavor symmetry between D_s and D^+ semileptonic decays would suggest equal r_V and r_2 for $D^+ \to \bar{K}^{*0}\ell^+\nu_\ell$ and $D_s^+ \to \phi\ell^+\nu_\ell$, where a spectator $\bar{d}$ quark is replaced by a spectator $\bar{s}$ quark. This indeed is the case for r_V, while r_2 appears inconsistent at the 3σ level.

The FOCUS measurement[309] featured the observation of a broad S-wave amplitude component in the $K\pi$ resonant substructure, which evidently accompanies the $K^*(890)$ amplitude in the decay $D^+ \to K^-\pi^+\mu^+\nu_\mu$. This claim is based on the observation of a relevant forward-backward asymmetry in $\cos\theta_V$ (the angle between the pion and the D direction in the $K^-\pi^+$ rest frame of decay), for events with $K^-\pi^+$ mass below the $\bar{K}^{*0}$ pole. FOCUS interprets this as the interference of a broad s-wave amplitude and the $\bar{K}^{*0}$ amplitude. The asymmetry is well represented by a $0.36\exp(i\pi/4)(GeV^{-1})$ s-wave, with a strength of about 7% relative to $\bar{K}^{*0}$.

The S wave component has important consequences in the computation of formfactors, and in the measurement of relative semileptonic branching ratios. The new measurement by FOCUS of branching ratios does include the effect of such a s-wave component and it significantly changes the PDG02 average.

It is desirable to have these findings checked by other experiments. If confirmed, one should study whether this broad S wave structure has some connection to the broad scalar κ state that has emerged recently in a Dalitz plot analysis by E791, see Sect.**9**˙5.2.

There are also consequences for B decays, namely a proper interpretation of the rare $B \to (K\pi)\gamma$, $(K\pi)\ell^+\ell^-$ and $B_s \to (K\pi)\ell\nu$ transitions. FOCUS will soon provide a measurement of formfactor ratios for $D_s^+ \to \phi\ell^+\nu_\ell$ from a ten-thousand events sample. From this measurement it will be interesting to see whether a $\phi - f_0$ interference phenomenon in the $K^-K^+\mu\nu$ decay can be found, analogous to the broad S-wave reported in $K^-\pi^+\mu\nu$. FOCUS have also announced [308] a large sample of clean $D \to \pi\ell\nu$ signals, whose relevance in extracting the q^2 dependance has been discussed previously. BELLE and BABAR have already accumulated much larger statistics than CLEO and shall present new results soon.

Results from B-factories are also highly expected on $D_{(s)}^+ \to (\eta/\eta')\ell\nu$. Only two measurements do exist in PDG02, from CLEO and the neutrino experiment E653. Their rates can shed light on the relative weight of decay mechanisms – spectator vs. WA – and on the strange vs. nonstrange vs. *non-*$q\bar{q}$ components of the η and η' wavefunctions. Very preliminary new results from CLEO have been very recently announced[282]. In Table IX we compare predictions from quark models with the data from a few modes. The predictions show some significant variations, which are – not surprisingly – very sizeable for $D \to \pi$.

We conclude this subsection with the discussion of a few selected topics in charm *baryon* semileptonic decays. If both initial and final state hadrons possess heavy flavour – like in $b \to c\ell\nu$ – then the semileptonic decays of baryons would exhibit some simpler features than those of mesons. Since in Qqq' the light diquark qq' carries spin zero, in the heavy quark limit there is a *single* spin-1/2 groundstate in the baryon sector rather than the spin-0 and -1 meson states. Likewise the three form factors that in general describe $\Lambda_{Q_1} \to \Lambda_{Q_2}\ell\nu$ can be expressed through a single independent Isgur-Wise function. In $\Lambda_c \to \Lambda\ell\nu$ the strange baryon cannot be treated as a heavy flavour object. Nevertheless it would be interesting to study to which degree these heavy quark symmetry expectations hold in Λ_c decays and the total Λ_c semileptonic width is saturated by $\Lambda_c \to \Lambda\ell\nu$.

The experimental scenario for charmed baryons semileptonic decays is at its infancy. Only an handful of branching ratio measurements are reported in PDG02, for decays $\Lambda_c \to \Lambda\ell\nu, \Xi_c \to \Xi\ell\nu$. New results have recently come from CLEO[313] and BELLE[314] on $\Omega_c \to \Omega\ell\nu$ and are reviewed in [282], where new findings from CLEO on Λ_c form factors are also presented. These studies mark just the beginning and we look forward to exciting news.

8˙2.2. Saturating the inclusive width. As explained in Sect.**4**˙11 quark-hadron duality represents a powerful concept at the heart of countless theoretical arguments. With charm being just 'beyond the border' of heavy flavours it can provide us with important insights about the transition to the duality regime. This is true in particular for semileptonic decays. More specifically one can study not only to which degree a quark-gluon based treatment can describe the total semileptonic width of the various charm hadrons, see Sect.**4**˙6, but also how the latter is built up by *exclusive* channels. This is not 'merely' of intrinsic intellectual interest, but might lead to practical lessons helping our understanding of $D^0 - \bar{D}^0$ oscillations and even $B \to \ell\nu X_u$.

Of course, the quantitative features are quite different in $D \to \ell\nu X_{S=-1,0}$ and in $B \to \ell\nu X_u$, let alone $B \to \ell\nu X_c$, and lessons can therefore not be drawn mechanically.

	Norm. Mode	Rel BR	[(1)] Abs. BR %	Scale factor
$D^+ \to$				
$\bar{K}^0 e^+ \nu_e$			6.5 ± 0.9	
	$\bar{K}^0 \pi^+$	2.39 ± 0.31		
	$K^- \pi^+ \pi^+$	0.73 ± 0.09		
$\bar{K}^0 \mu^+ \nu_\mu$			$7.0^{+3.0}_{-2.0}$	
$\bar{K}^- \pi^+ e^+ \nu_e$			$4.4^{+0.9}_{-0.7}$	
$\bar{K}^*(892)^0 e^+ \nu_e$			3.2 ± 0.33	
$K^- \pi^+ \mu^+ \nu_\mu$			3.79 ± 0.33	1.1
$\bar{K}^*(892)^0 \mu^+ \nu_\mu$			3.0 ± 0.4	
$K^- \pi^+ \mu^+ \nu_{\mu\,(NR)}$			0.31 ± 0.12	
	$K^- \pi^+ \mu^+ \nu_\mu$	0.083 ± 0.029		
$\pi^0 \ell^+ \nu_\ell$			0.31 ± 0.15	
	$\bar{K}^0 \ell^+ \nu_\ell$	$0.046 \pm 0.014 \pm 0.017$		
$\bar{K}^*(892)^0 e^+ \nu_e$			5.3 ± 0.7	1.5
	$K^- \pi^+ e^+ \nu_e$	$1.21^{+0.21}_{0.24}$		
	$K^- \pi^+ \pi^+$	0.60 ± 0.07		1.7
$\bar{K}^*(892)^0 \mu^+ \nu_\mu$			$5.2 \pm 0.$	1.1
	$K^- \pi^+ \pi^+$	0.590 ± 0.023		[(2)] 1.1
$\rho^0 e^+ \nu_e$			0.24 ± 0.09	
	$\bar{K}^*(892)^0 e^+ \nu_e$	$0.045 \pm 0.014 \pm 0.009$		
$\rho^0 \mu^+ \nu_\mu$			0.32 ± 0.08	
	$\bar{K}^*(892)^0 \mu^+ \nu_\mu$	0.061 ± 0.014		
$D^0 \to$				
$K^- e^+ \nu_e$			3.58 ± 0.18	1.1
	$K^- \pi^+$	0.94 ± 0.04		
$K^- \mu^+ \nu_\mu$			3.19 ± 0.17	
	$K^- \pi^+$	0.84 ± 0.04		
	$\mu^+ X$	0.49 ± 0.06		
$K^- \pi^0 e^+ \nu_e$			$1.1^{+0.8}_{-0.6}$	1.6
$\bar{K}^0 \pi^- e^+ \nu_e$			1.8 ± 0.8	1.6
$\bar{K}^*(892)^- e^+ \nu_e$			1.43 ± 0.23	
	$K^- e^+ \nu_e$	0.60 ± 0.10		
	$\bar{K}^0 \pi^+ \pi^-$	0.36 ± 0.06		
$\pi^+ e^+ \nu_e$			0.36 ± 0.06	
	$K^- e^+ \nu_e$	0.101 ± 0.017		
$D_s^+ \to$				
$\phi \ell^+ \nu_\ell$			2.0 ± 0.5	
	$\phi \pi^+$	0.55 ± 0.04		
$\eta \ell^+ \nu_\ell + \eta'(958) \ell^+ \nu_\ell$			3.4 ± 1.0	
	$\phi \ell^+ \nu_\ell$	1.72 ± 0.23		
$\eta \ell^+ \nu_\ell$			2.5 ± 0.7	
	$\phi \ell^+ \nu_\ell$	1.27 ± 0.19	0.89 ± 0.33	
$\eta'(958) \ell^+ \nu_\ell$				
	$\phi \ell^+ \nu_\ell$	0.44 ± 0.13		

TABLE VIII. – *Summary of semileptonic branching ratios for D^0, D^+ and D_s^+ [131]. Limits are not shown. Notes: [(1)] Absolute BR are either results of PDG fit, or direct measurements at e^+e^- (see [131] for details). [(2)] This PDG03 new fit includes the FOCUS measurement that finds a small S-wave $K^-\pi^+$ amplitude along with the dominant K^*. Fitting the FOCUS result together with the previous measurements should be taken cum grano salis.*

SL Channel	ISGW2 [101]	Jaus[310]	MS [311]	WWZ [312]	Exp
$\Gamma(D \to K)$	10.0	9.6	9.7	9.6	8.7
$\Gamma(D \to K^*)$	5.4	5.5	6.0	4.8	5.2
$\Gamma(D \to \pi)$	0.47	0.80	0.95	0.73	0.88
$\Gamma(D \to \rho)$	0.24	0.33	0.42	0.34	0.26
$\Gamma(D_s \to K)$	0.43		0.63		
$\Gamma(D_s \to K^*)$	0.21		0.38		

TABLE IX. – *Predictions for various semileptonic decay rates in* $10^{10}s^{-1}$.

Nevertheless it will be instructive to analyze the precise substructure – resonant as well as nonresonant – in final states like $\ell\nu K\pi$, $\pi\pi$, $\eta\pi$, $K\pi\pi$, 3π etc. and the relative weight of semileptonic final states with more than two hadrons. A detailed measurement of $D \to \ell\nu\pi\pi$ provides a lab for studying $\pi - \pi$ interactions in a different kinematical regime than in K_{e4} decays.

8˙2.3. Light flavour spectroscopy in semileptonic decays.. Theoretical predictions for $\Gamma(H_c \to \ell\nu h)$ depend also on the wavefunction of the hadron h in terms of quarks (and even gluons). Considering the spectator diagram one finds that in $D_s^+ \to \ell\nu\eta/\eta'$ the η/η' are excited via their $\bar{s}s$ components, whereas in $D^+ \to \ell\nu\eta/\eta'$ it is done through $\bar{d}d$. Measuring these four widths accurately will provide novel information on the relative weight of their strange and nonstrange $\bar{q}q$ components. Comparing the predictions from two models for $D_s^+ \to \ell\nu\eta/\eta'$ with the data one finds

Yet there is another layer of complexity to it: while WA provides no more than a nonleading contribution to inclusive rates, it could affect exclusive modes very considerably. This can be visualized through the diagram of Fig. 27, where the c and $\bar{s}$ or $\bar{d}$, before they annihilate into a virtual W emit two gluons generating the η or η' through a gg component in the latter's wavefunctions [315]. It could conceivably affect also $D_s^+ \to \ell\nu\phi$.

This mechanism can give rise also to the unusual mode $D_s^+ \to l^+\nu\pi^+\pi^-$. There had even been speculation [315] that a glueball can be produced in such a way, if the latter is sufficiently light.

To conclude this subsection, we remind another example of unexpected surge of connection between charm semileptonics decays, and light flavour hadronic physics (see Sect.8˙2.1). The recent evidence for a broad S-wave resonant component in $D^+ \to (K\pi)\ell\nu$ can possibly be related to the debated observation of the κ resonance in Dalitz

SL Channel	ISGW2 [101]	Jaus[310]	MS [311]	WWZ [312]	Exp
$\Gamma(D_s \to \eta)$	3.5		5.0		5.3
$\Gamma(D_s \to \eta')$	3.0		1.85		1.9
$\Gamma(D_s \to \phi)$	4.6		5.1		4.1

TABLE X. – *Predictions for* D_s *semileptonic decay rates in* $10^{10}s$.

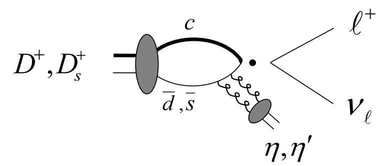


Fig. 27. – Diagrams for $(D^+, D_s^+) \to (\eta, \eta')\ell\nu$

plot analysis of charm decays (Sect.**9**˙5.1).

8˙3. $V(cs)$ *&* $V(cd)$. – The CKM parameters constitute fundamental quantities of the SM. It is believed that they are shaped by dynamics at high scales – a conjecture strengthened by the seemingly non-trivial pattern in their values, a point we will return to in our discussion of CP violation. Therefore one would like to determine them as precisely as possible. At first sight it might seem that of the three CKM parameters involving charm quarks – $V(cd)$, $V(cs)$ and $V(cb)$ – the first two are already known with excellent precision, as stated in PDG '00:

$$|V(cd)| = 0.222 \pm 0.002 \; , \; |V(cs)| = 0.9742 \pm 0.0006$$

$$|V(cb)| = 0.040 \pm 0.002 \tag{174}$$

Yet in arriving at these numbers one has *imposed* three-family unitarity. While there is no evidence for a fourth family and we actually know the Z^0 decays into three neutrino types only, we should not limit ourselves to a three-family scenario. Listing the values obtained from the data *without* three-family unitarity one has according to PDG '98:

$$|V(cd)| = 0.224 \pm 0.016 \; , \; |V(cs)| = 1.04 \pm 0.16$$

$$|V(cb)| = 0.0395 \pm 0.0017 \tag{175}$$

The uncertainties are considerably larger now, in particular for $|V(cs)|$. The latter is obtained primarily from $D \to \ell\nu K$ decays, the description of which suffers from sizeable theoretical uncertainties, as discussed before. On the other hand $|V(cd)|$ is derived mainly from the observed charm production rate in deep inelastic neutrino scattering is much better known – as is $|V(cb)|$ extracted from semileptonic B decays.

With three-family unitarity PDG '02 lists:

$$|V(cd)|_{(3)} = 0.223 \pm 0.002 \; , \; |V(cs)|_{(3)} = 0.9740 \pm 0.0006$$

$$|V(cb)|_{(3)} = 0.041 \pm 0.002 \; ; \tag{176}$$

i.e., no significant change relative to Eq.(174). Allowing for the existence of more families these numbers are relaxed to:

$$|V(cd)|_{(>3)} = 0.218 \pm 0.006 \; , \; |V(cs)|_{(>3)} = 0.971 \pm 0.003$$

$$|V(cb)|_{(>3)} = 0.041 \pm 0.002 \; ; \tag{177}$$

i.e., consistent values, yet again a larger uncertainty in particular for $|V(cs)|$. Relying on direct determinations PDG '02 lists ([24])

$$|V(cd)| = 0.224 \pm 0.016 \; , \; |V(cs)| = 0.996 \pm 0.013$$

$$|V(cb)| = 0.0412 \pm 0.0020 \tag{178}$$

There is a remarkable reduction in the uncertainty for the directly determined value of $|V(cs)|$ relative to the status in 1998. This is due to the usage of the method of extracting it from W leptonic branching fraction $B(W \to \ell\bar{\nu}_\ell)$, a novel method which comes from a combined analysis of about 60 000 WW events collected by four LEP experiments [316]. The leptonic branching fraction of the W is related to the CKM matrix elements without top quark

$$\frac{1}{B(W \to \ell\bar{\nu}_\ell)} = 3\Big[1 + \Big[1 + \frac{\alpha_s(M_W^2)}{\pi}\Big] \sum_{i=u,c;j=d,s,b} |V_{ij}|^2\Big] \tag{179}$$

The above given value of $|V_{cs}|$ is obtained assuming knowledge of $V_{ud,us,ub,cd,cb}$.

It has also been pointed out recently[316] that $|Vcs|$ can be extracted independently of other elements by measuring the $W \to cs$ branching fraction, which requires the reconstruction of $W \to$ ***charm* jet + *strange* jet**. Charm (or beauty) jets can be tagged easily based on peculiar characteristics, such as long lifetimes. Tagging of s jets is performed by tagging high-momentum kaons.The identification of kaons requires relevant particle identification technology, which at LEP was possessed only by DELPHI. They attempted to measure V_{cs} directly [317] using 120 hadronic W decays, and found $|V(cs)| = 0.97 \pm 0.37$.

While the PDG02 value $|V(cs)| = 0.996 \pm 0.013$ seems to be the limit accuracy to-date, it is pointed out[318] how at a Linear Collider with $5 \cdot 10^6$ W decays one could reach a precision of about 0.1%.

Optimism has been expressed that lattice QCD will be able to compute the form factors for $D \to \ell\nu K/\pi$ very considerably in the next very few years. It will be a tall order, though, to reduce them down to even the 5 % level, since it requires a fully unquenched calculation.

9. – Exclusive nonleptonic decays

In Sect.**6** we have already discussed a number of exclusive decays of hidden and open charm hadrons. This was done mainly in the context of spectroscopy and of lifetime measurements. We will treat exclusive modes now in their own right.

([24])Comparing the '98 and '00 errors for $|V(cb)|$ shows that PDG values should not be treated as gospel. Yet the uncertainties stated now are more credible than previously quoted ones.

9˙1. *The $\rho-\pi$ puzzle.* – The annihilation amplitudes for J/ψ and ψ' are proportional to the wavefunction at the origin $\psi(\mathbf{r}=\mathbf{0})$ as discussed in Sect.**6**. The decay widths for these S-wave states to a particular final hadronic state f consisting of light mesons then depends on $|\psi(0)|^2$. The decay width to e^+e^- also depends on this quantity, so one expects the following universal ratio:

$$Q_h = \frac{BR(\psi' \to f)}{BR(J/\psi \to f)} \approx \frac{BR(\psi' \to e^+e^-)}{BR(J/\psi \to e^+e^-)} = (12.3 \pm 0.7)\% \tag{180}$$

This relationship is indeed found to be satisfied for many hadronic final states like $p\bar{p}\pi^0$, $2(\pi^+\pi^-)\pi^0$, and $\pi^+\pi^-\omega$. However, in 1983 the Mark II Collaboration found a startling violation of this '12% rule' in decays to certain vector-pseudoscalar final states [319]. They observed ratios $Q_{\rho\pi} < 0.6\%$ and $Q_{K^*K} < 2\%$. These numbers have recently been confirmed by the BES Collaboration: $Q_{\rho\pi} < 0.23\%$, $Q_{K^{*+}K^-} < 0.64\%$, and $Q_{K^{*0}\bar{K}^0} = (1.7 \pm 0.6)\%$ [320]. The suppression of the ratio Q_h in these states is often referred to as the $\rho - \pi$ puzzle.

Twenty years after the first signal of this puzzle, the issue is still unresolved theoretically. There are several recent discussions of the theoretical situation in the literature [321, 322]. Below we will give a brief description of a few of the proposed solutions and their drawbacks.

One of the earliest proposed solution [323, 324] to the $\rho - \pi$ utilized the vector gluonium state, O, proposed earlier by Freund and Nambu [325]. This vector glueball would decay predominantly to vector pseudoscalar states like $\rho\pi$ and K^*K. If it is sufficiently close in mass to the J/ψ and narrow, it can mix significantly with the J/ψ, but not the ψ' and thus enhance the J/ψ decay rate into these final states relative to the three gluon rate. The current data require $|m_O - m_{J/\psi}| < 80$ MeV and $4\ \mathrm{MeV} < \Gamma_O < 50\ \mathrm{MeV}$[320]. BES has searched for the state and found no evidence[326]. Lattice QCD simulations suggest that these states should be several hundred MeV heavier[327].

It has also been suggested that the ρ meson may contain a large intrinsic charm component [328]. The radial wavefunction of the $c\bar{c}$ component of the ρ is argued to be nodeless and thus to have a much larger overlap with the J/ψ wavefunction than with the wavefunction of the ψ' again enhancing the J/ψ rate. So a sizeable intrinsic charm component in the ρ meson could produce an enhancement in the J/ψ rate to $\rho\pi$ that is not seen in the corresponding ψ' rate. Similar excitations of intrinsic charm in light mesons can be tested using $\eta_c(1S)$ and $\eta_c(2S)$ decays.

Another proposal [329] uses the framework of NRQCD to show that the decay of the J/ψ can proceed predominantly through a higher Fock state in which the $c\bar{c}$ pair is in a color octet which decays via $c\bar{c} \to q\bar{q}$. For ψ' decays this mechanism is suppressed by a dynamical effect arising from the proximity of the ψ' mass to the $D\bar{D}$ threshold. So for ψ' the decay through three gluons is expected to dominate. If the color octet contribution is sizable, we again see a mechanism that could generate an enhancement of the J/ψ rate to states like $\rho\pi$.

The above ideas all implicitly assume that hadron helicity is conserved in these processes. Perturbative QCD requires that decays to states in which the mesons have nonzero helicity are suppressed by at least $1/M_{J/\psi}$ [330]. For the decay to $\rho\pi$, there exists only one possible Lorentz-invariant quantity that can describe the $J/\psi - \rho - \pi$ coupling and the structure of this form factor requires the ρ to have helicity ± 1 violating hadron helicity conservation. For the above analyses, subleading processes are introduced that should otherwise be negligible were it not for the suppression of the leading helicity violating amplitude.

There does exist evidence that the assumption of hadron helicity conservation is questionable for J/ψ and ψ'. The BES Collaboration finds that the rate for $\psi' \to \pi^0\omega$ is not suppressed relative to the J/ψ rate, i.e. $Q_{\pi^0\omega} = (9.3 \pm 5.0)\%$ [320]. This process violates hadron helicity conservation but does not exhibit the behaviour that would be expected if any of the above models were correct. Measuring the decay rates to the hadron helicity conserving final state $\pi^+\pi^-$ should provide insight into the validity of this assumption.

There are many more models that attempt to resolve this intriguing puzzle, many of which do not invoke hadron helicity conservation. Unfortunately, there is no current model which can satisfactorily describe all of the relevant data. Clearly this is a problem that requires more attention both theoretically and experimentally.

9˙2. *Other charmonium decays.* – There are four classes of decays:

1. Electromagnetic decays of higher mass charmonia to lower mass ones – $[\bar{c}c]_N \to [\bar{c}c]_{n<N} + \gamma$ – or of charmonia into on- or off-shell photons: $[\bar{c}c]_N \to \gamma^* \to l^+l^-$, $\gamma\gamma$.

2. Radiative decays into one or several light-flavour hadrons:

$$[\bar{c}c] \to \gamma + \{h_{light}\} \tag{181}$$

3. Hadronic transitions of higher mass to lower mass charmonia:

$$[\bar{c}c]_N \to [\bar{c}c]_{n<N} + \{h_{light}\} \tag{182}$$

4. Hadronic decays into light-flavour hadrons

$$[\bar{c}c] \to \{h_{light}\} \tag{183}$$

A wealth of new data is flowing in thanks to e^+e^- experiments (B-factories BABAR and BELLE, upgraded versions of experiments CLEO and BES), and to final results from established $\bar{p}p$ formation experiment E835 at Fermilab.

An important restructuring of the information contained in the PDG occurred in 2002, and it affected relevantly (up to about 30%) the values of some branching ratios of ψ' and $\chi_{c0,1,2}$. Details are discussed in [331]: with the global fit provided in the 2002 edition, the correlations and vicious circles introduced through the use of self-referencing relative branching ratios should have been corrected.

Great interest is focussed on the $\rho\pi$ puzzle discussed in the previous section. Other recent topical issues are a new determination of $\Gamma(\chi_{c0,2} \to \gamma\gamma)$ by E835, and the complementary measurement of cross-section for charmonium production in two-photon processes $e^+e^- \to e^+e^-\gamma\gamma \to e^+e^-R$ from DELPHI ($R = \eta_c$) and BELLE ($R = \chi_{c0}$) [332]; experiment E835 also presented the first evidence for the decay $\chi_{c0} \to \pi^0\pi^0, \eta\eta$ [333]. First evidence for $\chi_{c0,1,2} \to \Lambda\bar{\Lambda}$ was reported by BES [334].

New results are too numerous to be covered here, the reader is addressed to recent reviews [123, 335, 336] for results on the four classes above, and to the comments already discussed on selected topics interlaced with charmonium spectroscopy in Sect.**6**˙2; here we will add comments on class (2).

9˙2.1. $[\bar{c}c] \rightarrow \gamma + \{h_{light}\}$. The search for glueball candidates is a central motivation in studying these transitions – in particular from the J/ψ –, which to lowest order are driven by $[\bar{c}c] \rightarrow \gamma + gg$: any kinematically accessible glueball or other hadron with a sizeable gluon component in its wave function should figure prominently in the hadronic final state. The most direct way is to measure the *inclusive* hadronic recoil spectrum $d\sigma(J/\psi \rightarrow \gamma X)/dM_X$ and search for peaks there. Alternatively one can analyze *exclusive* hadronic final states for the presence of new states outside the usual $\bar{q}q$ multiplets.

Experiments at e^+e^- colliders accumulating large data samples at the J/ψ have made of glueball search one of the top priorities. No conclusive signal has been found there to-date. The scalar glueball is predicted around 1700 MeV, and the lowest tensor glueball at about 2220 MeV. In 1996 Mark III [337] claimed evidence for the lightest scalar glueball, indicating the $f_0(1700)$ after a Partial Wave Analysis. It was advocated later[338] how this candidacy should be compared to the E791 Dalitz plot studies of $D_s \rightarrow KK\pi$ (Sect.**9**˙5.1). BES claimed candidates for the tensor glueball (the $f_J(2220)$) in 1996, and they are now performing detailed searches of several radiative final states [339] with their total sample of 58 million J/ψ decays. No true glueball candidate is supposed to be observed in $\gamma\gamma$ production, and therefore CLEO and LEP experiments have performed a search, setting limits. Finally, searches for the $f_J(2220)$ have been carried over at CLEO[340] in $\Upsilon \rightarrow \gamma f_J$ decays.

The experimental table is evidently very rich, however dinner is not ready yet. We may very well need to await for the honour guest — CLEO-c.

9˙3. *On absolute charm branching ratios.* – Attempts to measure absolute charm branching ratios date back to the very early stages of the charm adventure at accelerators. In principle one has "merely" to observe all decays of a certain hadron, and then count how often a specific final state appears. In reality things are of course less straightforward. The first attempts [341] relied on estimating the charm cross section at the $\psi(3770)$ to convert the observed signal into D branching fractions. It was assumed that the $D^0\bar{D}^0$ and $D^{\pm}D^{\mp}$ partial width were the same and that they saturate $\Gamma(\psi(3770))$.

Such analyses suffer from the systematic uncertainty on the exact size of the relevant production rate. To overcome this handicap, several approaches have been pursued over the past two decades.

1. The Mark III collaboration at SPEAR first exploited the method of *tagged decays*. Since the $\psi(3770)$ is just barely above charm threshold, its decays into charm cannot produce more than $D^0\bar{D}^0$ and D^+D^- pairs. Absolute branching ratios for D mesons can then be obtained [342] by comparing the number of single and double tags, i.e., fully reconstructed $D\bar{D}$ events . Single tag events were more abundant, but double tags benefitted from additional kinematic constraints, such as beam energy and momentum conservation, etc. This method does *not* have to assume that the total $\psi(3770)$ width is saturated by $\Gamma(\psi(3770)) \rightarrow D\bar{D})$. Typical errors were of order 10%.

2. The method pionereed by ALEPH [343] and HRS [344] exploits the fact that the pion from the strong decay $D^{*+} \rightarrow D^0\pi^+$ is very soft due to the small Q-value. In particular, the D^0 is not reconstructed, but its direction is well approximated by the thrust axis of the event, and the soft pion characteristic kinematical signature tags the event.

3. ACCMOR[345] and LEBC-EHS[346] measured the ratio of $D^0 \to K^-\pi^+$ relative to the total number of even-prong decays. The absolute branching ratio is then calculated using topological branching fractions to correct for missing zero-prong decays.

4. ARGUS[347] and then CLEO[348] determined the inclusive number of D^0's by partial reconstruction of $\bar{B}^0 \to D^{*+}\ell^-\bar{\nu}$ with $D^{*+} \to D^0\pi^+$ where only the soft pion and the lepton are detected. The D^* direction is, as always, approximated by the direction of the soft pion. The missing mass squared $(E_B - E_\ell - E_{D^*})^2 - |\vec{P}_B - \vec{P}_\ell - \vec{P}_{D^*}|^2$ shows a prominent peak at zero: the number of events in the peak provides the normalizing factor of the number of D^* decays.

5. For the absolute branching ratios of the D_s^+ meson, E691[349] and later E687[350] determined it by measuring $\Gamma(D_s^+ \to \phi\mu^+\nu)/\Gamma(D_s^+ \to \phi\pi^+)$, by using the D_s^+ lifetime, and assuming the theoretical expectation that $\Gamma(D_s^+ \to \phi\mu^+\nu) \sim \Gamma(D^+ \to \bar{K}^{*0}\mu^+\nu)$.

6. Also for the D_s^+ meson, low-statistics samples were collected by BES[351] by using double-tag $D_s^+D_s^-$ pairs exclusively produced in e^+e^- at energy just below the D_s^* production threshold.

7. In the charmed baryon sector, the only measurements available for absolute branching ratios refer to the Λ_c^+. A concise but complete review is found in [352]. No model-independent measurements exist. ARGUS[353] and CLEO[354] measure $B(\bar{B} \to \Lambda_c^+X) \cdot B(\Lambda_c^+ \to pK^-\pi^+)$ and, assuming that $\Lambda_c X$ channel saturate the B meson decays into baryons and that $\Lambda_c^+ X$ final states other than $\Lambda_c^+\bar{N}X$ can be neglected, they also measure $B(\bar{B} \to \Lambda_c^+X)$. Hence, $B(\Lambda_c^+ \to pK^-\pi^+)$ is extracted. ARGUS[355] and CLEO[356] also measure $\sigma(e^+e^- \to \Lambda_c^+X) \cdot B(\Lambda_c^+ \to \Lambda\ell^+\nu_\ell)$. The PDG group combines these measurements with $\sigma(e^+e^- \to \Lambda_c^+X) \cdot B(\Lambda_c^+ \to pK^-\pi^+)$, also estimating $B(\Lambda_c^+ \to pK^-\pi^+)$. The model-dependent systematic error estimated is of order 30%. A different approach is attempted by ? [357], that tags charm events with the semielectronic decay of a D^*-tagged $\bar{D}$, and the Λ_c^+ production with a $\bar{p}$. The assumption here is that having a $\bar{D}$ meson in one hemisphere, and a $\bar{p}$ in the opposite hemisphere, is a tagging for a Λ_c^+ in the hemisphere of the antiproton. Their final value is $B(\Lambda_c^+ \to pK^-\pi^+) = (5.0 \pm 0.5 \pm 1.2)$ %, which coincides with the PDG02 average from the two older measurements described above.

8. A novel method analyzes B meson decays into charm hadrons and then extracts the latter's branching ratio by relying on various correlations in the overall B decay. This method has been pioneered by CLEO. The measurement[358] is performed using partially reconstructed decays $\bar{B}^0 \to D^{*+}D_s^{*-}$. This decay is peculiar since it contains a soft pion from D^* decay, and a soft photon from the D_s^* decay. Another example has been discussed in Sect.9˙3: utilizing the huge data samples to be accumulated by the B factories one reconstructs one B meson in $\Upsilon(4S) \to B\bar{B}$ and then exploits correlations between baryon and lepton numbers and strangeness among the decay products of the other B meson to infer the absolute semileptonic branching ratios of charm baryons [359].

Tab.XI shows the present world averages for the measured absolute branching ratios [131] along with the average χ^2 as computed by the PDG group. The 20% error on the Λ_c is

Abs. BR	PDG02	Error scale factor
$D^+ \to$		
$\bar{K}^0\pi^+$	0.0277 ± 0.0018	—
$K^-\pi^+\pi^+$	0.091 ± 0.006	—
$\bar{K}^0\pi^+\pi^0$	0.097 ± 0.030	1.1
$K^-\pi^+\pi^+\pi^0$	0.064 ± 0.011	—
$\bar{K}^0\pi^+\pi^+\pi^-$	0.070 ± 0.009	—
$K^*(892)^-\pi^+\pi^+$	0.021 ± 0.009	—
$\phi\pi^+\pi^0$	0.023 ± 0.010	—
$K^+K^-\pi^+\pi^0$	0.015 ± 0.007	—
$K^*(892)^+\bar{K}^*(892)^0$	0.026 ± 0.011	—
$D^0 \to$		
$K^-\pi^+$	0.0380 ± 0.0009	—
$\bar{K}^0\pi^+\pi^-$	0.0592 ± 0.0035	1.1
$K^-\pi^+\pi^0$	0.131 ± 0.009	1.3
$\bar{K}^0\pi^+\pi^-\pi^0$	0.108 ± 0.013	—
$D_s^+ \to \phi\pi^+$	0.036 ± 0.009	—
$\Lambda_c^+ \to pK^-\pi^+$	0.050 ± 0.013	—

TABLE XI. – *World averages for charm mesons and baryons absolute branching ratios from [131].*

clearly unsatisfactory, and it should be noted that there are no absolute branching ratios for the other charmed baryons.

It is legitimate to ask why one wants to measure charm branching ratios so accurately, why, say, a 10% accuracy is not satisfactory, when one has to allow for about 30% uncertainties in theoretical predictions. There are several motivations of somewhat different nature.

- The absolute values of the $D^{0,+}$, D_s, Λ_c etc. branching ratios are an important 'engineering' input for a sizeable number of B decay analyses like the following:
  - When extracting the CKM parameter $V(cb)$ from the semileptonic B width and from the formfactor at zero recoil for the exclusive transitions $B \to \ell\nu D^*$ and $B \to \ell\nu D$ they are needed to translate the observed rate for, say, $B \to \ell\nu(K\pi)_D\pi$ into a width for $B \to \ell\nu D^*$ etc.
  - They are also of very direct importance for evaluating the charm content in nonleptonic B decays in particular and to compare the $b \to c\bar{u}d$ and $b \to c\bar{c}s$ widths of B mesons. It had been pointed out first almost ten years ago that the observed semileptonic B branching ratio is somewhat on the low side of theoretical expectation [360]. This could be understood if the $b \to c\bar{c}s$ width were larger than expected. The charm content of B decays, i.e. the average number of charm hadrons in the final states of B mesons is thus an observable complementary to BR$(B \to \ell\nu X_c)$. Its determination obviously depends on the absolute size of D^0, D^+, D_s and Λ_c branching ratios.

It turns out that the accuracy, with which absolute charm branching ratios are known, is about to become – or will be in the foreseeable future – a major bottle neck in the analysis of beauty decays.

- As discussed in Sect.**4**˙7 the HQE predicts quite spectacular variations in the semileptonic widths and branching ratios of charm baryons.

- The one-prong decays $D^+,\ D_s^+ \to \mu^+\nu$, $\tau^+\nu$ are measured to extract the decay constants f_D and f_{D^*} from the *widths* and compare the results with predictions in particular of lattice QCD, see Sect.**7**.

- Analogous statements apply to exclusive semileptonic charm decays. It is an often repeated promise of lattice QCD that it will calculate the semileptonic form factors accurately both in their normalization and q^2 dependence.

Fortunately help is in sight: (i) The new tau-charm factory CLEO-c will measure the absolute scale of D and D_s absolute branching ratios with a 1-2% error, and possibly the Λ_c absolute branching ratio significantly more accurately as well. (ii) Novel methods have been proposed to get at the branching ratios for charm baryons [359].

9˙4. *Two-body modes in weak nonleptonic decays*. – Nonleptonic *weak* decays pose a much stiffer challenge to a theoretical description than semileptonic ones: there are more colour sources and sinks in the form of quarks and antiquarks, more different combinations for colour flux tubes to form, and the energy reservoir is not depleted by an escaping lepton pair. There are only two classes of nonleptonic decays where one can harbour reasonable hope of some success at least, namely fully inclusive transitions like lifetimes etc. already discussed and channels with a two-body final state. Once one goes beyond two hadrons in the final state, the degrees of freedom and the complexities of phase space increase in a way that pushes them beyond our theoretical control.

Therefore we will discuss only two-body channels in detail including those with resonances.

For two-body final states the phase space is trivial and the number of formfactors quite limited. Yet even so such transitions present a formidable theoretical challenge, since they depend on long-distance dynamics in an essential way. As an optimist one might point to some mitigating factors:

- Two-body final states allow for sizeable momentum transfers thus hopefully reducing the predominance of long-distance dynamics.

- It is not utopian to expect lattice QCD to treat these transitions some day in full generality. Such results will however be reliable only, if obtained with incorporating fully dynamical fermions – i.e. without "quenching" – and without relying on a $1/m_c$ expansion.

- Watson's theorem can still provide some useful guidance.

In addition there are motivations for taking up this challenge:

- Carefully analysing branching ratios and Dalitz plots can teach us novel lessons on light-flavour hadron spectroscopy, like on characteristics of some resonances like the σ or on the η and η' wavefunctions and a possible non-$\bar{q}q$ component in them.

- Analogous B decay modes are being studied also as a mean to extract the complex phase of $V(ub)$. One could hope that D decays might serve as a validation analysis. For honesty's sake one has to add that in the two frameworks presently available for treating such B decays – usually refered to as "QCD factorization" and "PQCD" – such a connection cannot be exploited largely due to technical reasons. For contrary to the HQE treatment of *inclusive* rates, where the leading nonperturbative corrections arise only in order $1/m_Q^2$, these frameworks for *exclusive* widths allow for corrections already $\sim \mathcal{O}(1/m_Q)$. Those, which for D decays potentially are very large, cannot be controled.

- As discussed later a most sensitive probe for New Physics is provided by searches for CP violation in the decays of charm hadrons. Two-body nonleptonic modes provide good opportunities for such searches. Signals of *direct* CP violation typically require the intervention of FSI in the form of phaseshifts. For designing search strategies and for properly interpreting a signal (or the lack of it) one needs independent information on such FSI phases. A comprehensive analysis of two-body charm decays can provide such information.

- It is not unreasonable to ask whether New Physics – rather than novel features of SM dynamics – could manifest itself also by enhancing or suppressing some exclusive widths, as discussed recently [361]. However one is facing a 'Scylla and Charybdis' dilemma here: due to our limited theoretical control only very sizeable deviations from SM expectations can be viewed as significant. But then one has to wonder why such a discrepancy has not been noticed before in other transitions. Doubly Cabibbo suppressed modes presumably offer the best 'signal-to-noise' ratio.

We can easily see that FSI produce generally large phase shifts by using isospin decompositions of decay amplitudes. Consider the modes $D \to K\pi$: there are three channels – $D^0 \to K^-\pi^+$, $D^0 \to \bar{K}^0\pi^0$ and $D^+ \to \bar{K}^0\pi^+$ – yet only two independent amplitudes, namely with isospin 1/2 and 3/2 in the final state, $T_{\frac{1}{2}}$ and $T_{\frac{3}{2}}$, respectively. Thus there has to be a relation between the three (complex) transition amplitudes; i.e., they have to form a triangle relation, which is easily worked out:

$$
\begin{aligned}
T(D^0 \to K^-\pi^+) &\equiv T_{-+} = \frac{1}{\sqrt{3}}\left(\sqrt{2}T_{\frac{1}{2}} + T_{\frac{3}{2}}\right) \\
T(D^0 \to \bar{K}^0\pi^0) &\equiv T_{00} = \frac{1}{\sqrt{3}}\left(-T_{\frac{1}{2}} + \sqrt{2}T_{\frac{3}{2}}\right) \\
T(D^+ \to \bar{K}^0\pi^+) &\equiv T_{0+} = \sqrt{3}T_{\frac{3}{2}}
\end{aligned}
\tag{184}
$$

and thus

$$
T_{-+} + \sqrt{2}T_{00} - \frac{2}{3}T_{0+} = 0\ . \tag{185}
$$

Using the most recent PDG values for the branching ratios one infers

$$
\begin{aligned}
|T_{\frac{1}{2}}| &= (3.05 \pm 0.06) \times 10^{-3} MeV \\
|T_{\frac{3}{2}}| &= (7.67 \pm 0.25) \times 10^{-4} MeV \\
\delta_{\frac{3}{2}} - \delta_{\frac{1}{2}} &= (95.6 \pm 6.3)^0 .
\end{aligned}
\tag{186}
$$

where $T_{\frac{j}{2}} \equiv |T_{\frac{j}{2}}| e^{i\delta_{\frac{j}{2}}}$, $j = 1, 3$. I.e., the phase shift $|\delta_{\frac{3}{2}} - \delta_{\frac{1}{2}}|$ is indeed very large in this case.

While this turns out as expected and can encourage us to search for direct CP asymmetries in D decays, it also implies that describing nonleptonic two-body modes will be challenging, to put it euphemistically.

9˙4.1. Early phenomenology. The first two pieces of the "charm puzzle" , i.e. of evidence that charm decays do not proceed quite to the *original* expectations actually emerged in D^0 two-body modes ([25]):

$$\left.\frac{\Gamma(D^0 \to K^+K^-)}{\Gamma(D^0 \to \pi^+\pi^-)}\right|_{\mathrm{data}} \sim 3 \quad vs. \quad \sim 1.4 \quad \text{"originally expected"} \tag{187}$$

$$\left.\frac{\Gamma(D^0 \to \bar{K}^0\pi^0)}{\Gamma(D^0 \to K^-\pi^+)}\right|_{\mathrm{data}} \sim \tfrac{1}{2} \quad vs. \quad \ll 1 \quad \text{"originally expected"} \tag{188}$$

Subsequently it was suggested [362] that another channel would be very telling:

$$\mathrm{BR}(D^0 \to \bar{K}^0\phi) \simeq 0 \qquad \text{"naively expected without WA"} \tag{189}$$

The success or failure of a theoretical description will depend on its ability to reproduce or even predict the *whole pattern* of decay modes. Nevertheless the three observable ratios of Eqs.(187,188,189) provide a first orientation for evaluating how well a certain model does in describing the data.

To understand what underlies the "expectations" and what one learns from these discrepancies we have to describe how one arrives at such predictions. In doing so we will *not* follow the historical sequence.

The starting point is always provided by the effective weak $\Delta C = 1$ Lagrangian, as discussed in Sect.**4**˙10.1. For our subsequent discussion it is instructive to write it in terms of the multiplicatively renormalized operators. For Cabibbo allowed modes we have:

$$\mathcal{L}_{eff}^{\Delta C=1}(\mu = m_c) = (4G_F\sqrt{2})V(cs)V^*(ud) \cdot [c_-O_- + c_+O_+] \tag{190}$$

$$O_\pm = \frac{1}{2}\left[(\bar{s}_L\gamma_\nu c_L)(\bar{u}_L\gamma_\nu d_L)\right] \pm (\bar{u}_L\gamma_\nu c_L)(\bar{s}_L\gamma_\nu d_L)] \tag{191}$$

$$c_- \simeq 1.90\ ,\ c_+ = 0.74 \tag{192}$$

Both operators $O_\pm$ carry isospin $(I, I_3) = (1, 1)$, yet are distinguished by their V spin quantum numbers, as described in Sect.**4**˙10.1: the $\Delta V = 0$ O_- is enhanced in $\mathcal{L}_{eff}^{\Delta C=1}$, the $\Delta V = 1$ O_+ reduced. This explains why the *width* for $D^+ \to \bar{K}^0\pi^+$ is reduced relative to $D^0 \to K^-\pi^+$. For D^+ carrying $V = 0$ can decay via O_- only into $V = 0$ final state. Bose statistics, however, tells us that the two V spinors $\bar{K}^0$ and π^+ have to

([25])The third piece was the observation that $\tau(D^+)$ exceeds $\tau(D^0)$ considerably.

be symmetric under exchange and thus form a $V = 1$ configuration. Yet the latter can be reached only through O_+. $D^0 \to K^-\pi^+$ on the other hand can be driven by O_-.

The quark level diagrams for nonleptonic and semileptonic decays do not look so different. Yet with having two quark and two antiquark fields in the final state rather than one each there are more different routings for the colour flux tubes between the fields and ways in which those arrange themselves into hadrons. To illustrate this point consider again $D^0 \to K^-\pi^+$: Fig.28 a [28 b] shows the diagram for this transition driven by the operator O_1 [O_2] through a topology usually called "external [internal] W emission". The mode $D^0 \to \bar{K}^0\pi^0$ is driven by diagrams with the same topologies, yet with O_1 and O_2 having switched places. For inclusive decays this distinction between "internal" and "external" W emission does not exist, yet for exclusive ones it does.

Drawing quark diagrams is one thing, attaching numerical values to them quite another. For the different fields interact strongly with each other. On the one hand this is a blessing since these strong interactions generate a large number of possible hadronic final states from the limited set of quark level configurations. On the other hand it is quite difficult to make rigorous predictions regarding the evolution from quarks and gluons to asymptotic hadronic states. The problem lurks in calculating hadronic matrix elements, since those are controled by long-distance dynamics. Lattice QCD is widely expected to yield reliable answers – someday; $1/N_C$ expansions provide a compact classification scheme, but not quantitative answers as explained below; $1/m_Q$ cannot make reliable statements on exclusive charm decays. Thus we have a situation, where quark models are called to the front as tools of last resort.

Stech and coworkers tackled the problem not unlike Alexander the Great did the Gordian knot. They relied on an unabashedly phenomenological approach where they invoked simplifying assumptions as much as practically possible while allowing for complexities only when unavoidable. Their approach turned out to be quite successful. In describing it we will point out its differences to earlier attempts.

Their prescription involves the following rules:

- One employs the effective weak Lagrangian of Eq.(190), yet leaves the coefficients $c_{1,2}$ as free parameters at first.
- One ignores WA diagrams completely.
- One allows for contributions both from external and internal W emission. In the former, Fig.(28 a), the quark-antiquark pair already forms a colour singlet, in the latter, Fig.(28 b), it does not. Accordingly one assigns a colour factor ξ to the matrix element of the latter relative to the former. Naively, i.e. by just counting the number of the different colour combinations, one would guestimate $\xi \simeq 1/N_C = 1/3$. Here instead one leaves the numerical value of ξ a priori completely free, yet maintains it to possess a *universal* value for all channels, whatever it is. This is a critical assumption, to which we will return.

The first consequence of these three rules is that all decay amplitudes can be expressed as linear combinations of two terms:

$$T(D \to f) \propto a_1 \langle f | J_\mu^{(ch)} J_\mu^{(ch)\prime} | D \rangle + a_2 \langle f | J_\mu^{(neut)} J_\mu^{(neut)\prime} | D \rangle \tag{193}$$

$$a_1 = c_1 + \xi c_2 \ , \ a_2 = c_2 + \xi c_1 \tag{194}$$

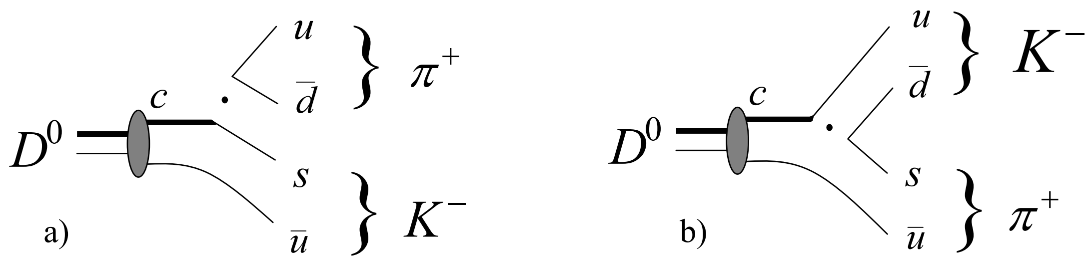


Fig. 28. – a) The external emission of the quark-antiquark pair forms a colour singlet. b) Internal emission of the quark-antiquark pair does not result in a colour singlet.

where $J_\mu^{(ch)}$ and $J_\mu^{(ch)\prime}$ [$J_\mu^{(neut)}$ and $J_\mu^{(ch)\prime}$] denote the charged [neutral] currents appearing in Fig.28 a [28 b]. It should be kept in mind that the quantities c_1 and c_2 on one hand and ξ on the other are of *completely different* origin despite their common appearance in a_1 and a_2: while $c_{1,2}$ are determined by short-distance dynamics, ξ parametrizes the impact of long distance dynamics on the size of matrix elements.

- One adopts a *factorization* ansatz for evaluating the matrix elements, i.e. one approximates the matrix element for $D \to M_1 M_2$ by the products of two simpler matrix elements:

$$\langle M_1 M_2 | J_\mu J_\mu | D\rangle \simeq \langle M_1 | J_\mu | 0\rangle \langle M_2 | J_\mu | D\rangle \ , \ \langle M_2 | J_\mu | 0\rangle \langle M_1 | J_\mu | D\rangle \tag{195}$$

- For $M_{1,2}$ being pseudoscalar mesons P one has

$$\langle P(p) | A_\mu | 0\rangle = -i f_P p^\mu \tag{196}$$

$$\langle P(p_P) | V^\mu | D(p_D)\rangle = (p_P^\mu + p_D^\mu - \frac{M_D^2 - M_P^2}{q^2} q^\mu) f_+(q^2)$$

$$+ \frac{M_D^2 - M_P^2}{q^2} q^\mu) f_0(q^2) \ ; \ q = p_D - p_P \ , \ f_+(0) = f_0(0) \tag{197}$$

The decay constants are known for the low mass mesons. For the formfactors $f_{0,+}$ BSW assumed dominance by the nearest t-channel pole:

$$f_{+,0} = \frac{h_{+,0}}{1 - \frac{q^2}{M^*}}. \tag{198}$$

For the values of the residues $h_i's$ one can rely on quark model calculations or take them from the available data on semi-leptonic decays.

Similar, though lengthier, expressions apply, when vectormesons are involved.

- FSI in the form of rescattering, channel mixing and conceivably even resonance enhancements are bound to affect charm decays. An optimist can entertain the hope that a large part or even most of such effects can be lumped into the quantity ξ, which reflects long-distance dynamics; however there is no good reason why such a *global* treatment of FSI should work even approximately. Stech et al. included residual FSI effects on a case-by-case basis; in the spirit of Ockham's razor FSI effects (in the form of phase shifts and absorption) were included only as much as really needed.
- It has to be kept in mind that certain final states cannot be produced by such factorizable contributions. One example is $D^0 \to \bar{K}^0 \phi$.

One can distinguish three classes of two-body modes [299]:

- class I: $D^0 \to M_1^+ M_2^-$
- class II: $D^0 \to M_1^0 M_2^0$
- class III: $D^+ \to M_1^+ M_2^0$

Class I[II] transitions receive contributions from a_1 [a_2] amplitudes only, whereas class III modes involve the interference between a_1 and a_2 amplitudes. It is therefore the latter that allow us to determine the relative sign between a_1 and a_2.

In summary: the BSW model put forward almost twenty years ago contains two free parameters – a_1 and a_2 – plus some considerable degree of poetic license in the size assumed for the residue factors h_i and the amount of explicit FSI that has been included.

With such limited freedom BSW were able to obtain a decent fit to twenty-odd two-body modes of D^0, D^+ and D_s^+ mesons. This is quite remarkable even keeping in mind that their description was helped by the sizeable errors in most measurements then. Their fit to the data yielded

$$a_1(exp) \simeq 1.2 \pm 0.1 \ , \ a_2(exp) \simeq -0.5 \pm 0.1 \tag{199}$$

Obviously one wants to compare this with what one might infer theoretically from Eq.(194):

$$a_1(QCD) \simeq 1.32 - 0.58\xi \ , \ a_2(QCD) \simeq -0.58 + 1.3\xi \tag{200}$$

It is again remarkable that these numbers are in the right 'ballpark' – and even more so that $\xi \simeq 0$ brings the numbers in Eq.(200) in full agreement with those in Eq.(199).

As already stated, *naively* one guestimates $\xi \simeq 1/N_C = 1/3$ leading to $a_1 \simeq 1.1$ and $a_2 \simeq -0.15$, which would greatly reduce expectations for class II and III transition rates. This and other related issues can be illustrated by considering $\Gamma(D^0 \to \bar{K}^0 \pi^0)$ vs. $\Gamma(D^0 \to K^- \pi^+)$. Ignoring QCD effects – in particular setting $c_1 = 1$, $c_2 = 0$ – one expects $\Gamma(D^0 \to \bar{K}^0 \pi^0)/\Gamma(D^0 \to K^- \pi^+) \simeq \left(\frac{1}{\sqrt{2}}\frac{1}{3}\right)^2 = \frac{1}{18}$. Using $c_1 \simeq 1.3$ and $c_2 \simeq -0.6$ reduces this ratio even further. However one should not have trusted this result in the first place! For the suppression of $T(D^0 \to \bar{K}^0 \pi^0)$ is due to an accidental cancellation of the $T_{I_f=1/2}$ and $T_{I_f=3/2}$ amplitudes in the simple prescription used, see Eq(184). Any FSI is quite likely to vitiate this cancellation. To say it differently: the mode $D^0 \to \bar{K}^0 \pi^0$ is particularly sensitive to the intervention of FSI. The BSW fit can reproduce the observed ratio only by invoking sizeable FSI as inferred from the observed $K\pi$ phase shifts. It is also amusing to note that this ratio was at first seen as clear evidence for WA playing a dominant role in all D decays. For WA produces a pure $I = 1/2$ final state where $\frac{\Gamma(D^0 \to \bar{K}^0\pi^0)}{\Gamma(D^0 \to K^-\pi^+)} = \frac{1}{2}$ has to hold.

The mode $D_s^+ \to \pi^+\pi^0$, which might be seen as a signal for WA, is actually forbidden by isospin invariance: due to Bose statistics the pion pair with charge 1 has to carry $I = 2$, which cannot be reached from the isoscalar D_s^+ via a $\Delta I = 1$ Lagrangian. The analogous Cabibbo forbidden $D^+ \to \pi^+\pi^0$ can be reached via the $\Delta I = 3/2$ Lagrangian, yet *not* by WA.

Adding up the widths for all these (quasi-)two-body channels results in a number close to the total nonleptonic widths for the charmed mesons. I.e., D decays largely proceed through the hadronization of two quark-antiquark pairs in the final state.This means also that the destructive PI mechanism prolonging the D^+ lifetime has to emerge in the two-body D^+ modes as well. It should be noted that this correlation does *not* exist for B decays, where two-body channels represent merely a small fraction of all B decays. It is then quite conceivable – and it has indeed been observed – that two-body B^- modes exhibit *constructive* PI *contrary* to its full width!

9˙4.2. Cabibbo forbidden channels. Several further complications and complexities arise on the once Cabibbo forbidden level. First, there are two transition operators before QCD corrections are included, namely for $c \to s\bar{s}u$ and $c \to d\bar{d}u$. The multiplicatively renormalized operators are

$$O_{\pm}^{(ss)} = \frac{1}{2}\left[(\bar{s}_L\gamma_\nu c_L)(\bar{u}_L\gamma_\nu s_L) \pm (\bar{s}_L\gamma_\nu s_L)(\bar{u}_L\gamma_\nu c_L)\right] \tag{201}$$

$$O_{\pm}^{(dd)} = \frac{1}{2}\left[(\bar{d}_L\gamma_\nu c_L)(\bar{u}_L\gamma_\nu d_L) \pm (\bar{u}_L\gamma_\nu c_L)(\bar{d}_L\gamma_\nu d_L)\right] \tag{202}$$

The operator $O_-^{(ss)}$ carries $\Delta V = 1/2$ only, $O_+^{(ss)}$ also $\Delta V = 3/2$; both are isospinors. The quantum numbers of $O_\pm^{(dd)}$ are quite analogous to those driving strange decays: $O_-^{(dd)}$ is purely $\Delta I = 1/2$, whereas $O_+^{(dd)}$ contains also $\Delta I = 3/2$. Their renormalization coefficients $c_\pm$ are as described in Sect.**4**˙10.1. One should note that to a very good approximation one has $V(cd) \simeq -V(us)$. In addition one has the Penguin-like operator

$$O_P = (\bar{u}_L\gamma_\nu c_L)\left[\sum_{q=u,d,s}(\bar{q}_L\gamma_\nu q_L) + \sum_{q=u,d,s}(\bar{q}_R\gamma_\nu q_R)\right] \tag{203}$$

One should note that the usual Penguin diagram does not yield a local operator since the internal quarks d and s are lighter than the external c quark and not even a short-distance operator since $m_s < \Lambda_{NPD}$. This is in marked contrast to the situation with Penguin diagrams for s and b decays.

In the limit of $SU(3)_{Fl}$ symmetry the rates for the channels $D^0 \to K^+K^-, \pi^+\pi^-$ have to coincide. Pure phase space favours $D^0 \to \pi^+\pi^-$ over $D^0 \to K^+K^-$ somewhat. Yet this is expected to be more than compensated by the larger form factors and decay constants for the $K\bar{K}$ final state. Within the factorization ansatz one finds

$$\frac{\Gamma(D^0 \to K^+K^-)}{\Gamma(D^0 \to \pi^+\pi^-)} \simeq \left|\frac{f_K}{f_\pi}\right|^2 \simeq 1.4 \,, \tag{204}$$

which goes in the right direction, yet not nearly far enough. It has been suggested that Penguin operators can close the gap since they contribute constructively to $D^0 \to K^+K^-$, yet destructively to $D^0 \to \pi^+\pi^-$ because of $V(us) \simeq -V(cd)$. Yet, as indicated above, the Penguin operator is shaped essentially by long distance dynamics, this suggestion has remained a conjecture.

The stand-by culprit for discrepancies between predictions and data are FSI. The preponderance of $D^0 \to K^+K^-$ over $D^0 \to \pi^+\pi^-$ could a priori be due to a large fraction of $\pi^+\pi^-$ being rescattered into $\pi^0\pi^0$ final states. Yet this is clearly not the case, see Table XVII. There is a simple intuitive argument why $\Gamma(D^0 \to K^+K^-)$ should exceed $\Gamma(D^0 \to \pi^+\pi^-)$: Two kaons eat up already a large fraction of the available phase space – unlike two pions; the probability for another quark-antiquark pair to be created should be lower in the former than in the latter case. To say it differently: $D^0 \to K\bar{K}$ represents a larger fraction of $D^0 \to K\bar{K} + n\ \pi's$ than $D^0 \to \pi\pi$ of $D^0 \to n\ \pi's$. This feature has indeed emerged in the data:

$$BR(D^0 \to K^+K^-\pi^+\pi^-) = (2.52 \pm 0.24) \cdot 10^{-3} \tag{205}$$
$$BR(D^0 \to \pi^+\pi^-\pi^+\pi^-) = (7.4 \pm 0.6) \cdot 10^{-3} \tag{206}$$

again exhibit very large $SU(3)_{Fl}$ symmetry breaking

$$\frac{BR(D^0 \to K^+K^-\pi^+\pi^-)}{BR(D^0 \to \pi^+\pi^-\pi^+\pi^-)} \simeq 0.34 \pm 0.04 \tag{207}$$

yet in a direction opposite to what happens in the two-body modes. Adding these two- and four-body rates one obtains

$$\frac{BR(D^0 \to K^+K^-, K^+K^-\pi^+\pi^-)}{BR(D^0 \to \pi^+\pi^-, \pi^+\pi^-\pi^+\pi^-)} \simeq 0.8 \pm 0.1 \tag{208}$$

These numbers illustrate the general feature that individual exclusive modes can exhibit very large symmetry violations that average out when summing over exclusive rates. The ratio in Eq.(208) is actually fully consistent with approximate $SU(3)_{Fl}$ invariance, as expected for the inclusive widths $\Gamma(D_q \to s\bar{s}u\bar{q})$ vs. $\Gamma(D_q \to d\bar{d}u\bar{q})$. These issues will be addressed again in our discussion of $D^0 - \bar{D}^0$ oscillations below.

There is a subtlety in relating the observable rate for $D^0 \to K_SK_S$ to that for $D^0 \to K^0\bar{K}^0$. Since $|K^0\bar{K}^0\rangle$ is equivalent to $|K_SK_S\rangle - |K_LK_L\rangle$ we have

$$\Gamma(D^0 \to K^0\bar{K}^0) = 2\Gamma(D^0 \to K_SK_S) \tag{209}$$

rather than $\Gamma(D^0 \to K^0\bar{K}^0) = 4\Gamma(D^0 \to K_SK_S)$.

The dynamics of *doubly Cabibbo suppressed decays* simplifies again since there is only one operator before QCD corrections are included: $c_L \to d_L\bar{s}_Lu_L$. The multiplicatively renormalized operators yet again are

$$O_{\pm}^{\Delta C=\Delta S=-1} = \frac{1}{2}\left[(\bar{d}_L\gamma_\nu c_L)(\bar{u}_L\gamma_\nu s_L) \pm (\bar{u}_L\gamma_\nu c_L)(\bar{d}_L\gamma_\nu s_L)\right] \tag{210}$$

with $O_{-}^{\Delta C=\Delta S=-1}$ being purely $\Delta I = 0$ and $O_{+}^{\Delta C=\Delta S=-1}$ $\Delta I = 1$.

The relative weight to the corresponding Cabibbo allowed modes is controlled by $(tg\theta_C)^4 \simeq 2.3\cdot 10^{-3}$ *on average*. Channel-by-channel there can be considerable deviations from this number due to differences in formfactors and FSI. For Cabibbo allowed modes are purely $\Delta I = 1$, while doubly Cabibbo suppressed channels are mainly $\Delta I = 0$. Furthermore doubly Cabibbo suppressed D^+ channels are *not* suppressed by PI. Therefore they are *enhanced on average* by the lifetime ratio $\tau(D^+)/\tau(D^0)$. Doubly Cabibbo suppressed D_s^+ rates on the other hand are reduced by PI; therefore they are *further decreased* by about the same factor on average.

The first evidence of the DCS decay $D^+ \to K^+K^-K^+$ had been reported by FOCUS [363], which measures

$$\Gamma(D^+ \to K^+K^-K^+)/\Gamma(D^+ \to K^-\pi^+\pi^+) = (9.5 \pm 2.2)\times 10^{-4} \tag{211}$$

Such a decay can proceed via the quark decay reactions only if coupled with a FSI like the rescattering $\bar{K}^0K^0 \Rightarrow K^+K^-$, which is quite conceivable. Comparing the branching ratio in Eq.(211) with the naive guestimate of $\Gamma_{DCSD}/\Gamma_{CF} \propto \tan^4\theta_C \simeq 2\times 10^{-3}$ shows there is ample room for such a rescattering to take place and still reproduce the observed width.

9·4.3. The $1/N_C$ ansatz. The fit result $\xi \simeq 0$ lead to the intriguing speculation that these weak two-body decays can be described more rigorously through $1/N_C$ expansions sketched in Sect.**4**·4 [364]. They are invoked to calculate hadronic matrix elements. The procedure is the following: One employs the effective weak transition operator $\mathcal{L}_{eff}(\Delta C = 1)$ given explicitly in Eq.(190); since it describes short distance dynamics, one has kept $N_C = 3$ there. Then one expands the matrix element for a certain transition driven by this operators in $1/N_C$:

$$T(D\to f) = \langle f|\mathcal{L}_{eff}(\Delta C=1)|D\rangle = \sqrt{N_C}\left(b_0 + \frac{b_1}{N_C} + \mathcal{O}(1/N_C^2)\right) \tag{212}$$

There are straightforward rules for determining the colour weight of the various quark-level diagrams: (i) assign a colour weight N_C to each *closed* quark loop and $1/\sqrt{N_C}$ to each quark-gluon coupling; (ii) treat every gluon line as a quark and antiquark line in colour space; (iii) normalize the meson wave functions in colour space, which amounts to another factor $1/\sqrt{N_C}$ per meson in the final or initial state.

Quark-gluon dynamics is treated here in a nonperturbative way, as can be seen by considering a quark loop: any "ladder" diagram – i.e. any planar diagram where gluon lines form the rungs between the quark and antiquark lines – has the same colour factor since every such gluon line creates a new loop in colour space, rule (ii), and thus a factor N_C, rule (i), which is compensated by the two quark-gluon couplings, rule (i).

Using these rules it is easy to show that the following simplifying properties hold for the contributions leading in $1/N_C$:

- one has to consider *valence* quark wave functions only;

- *factorization* holds;
- *WA* has to be ignored as have *FSI*.

To leading order in $1/N_C$ only the term b_0 is retained; then one has effectively $\xi = 0$ since $\xi \simeq 1/N_C$ represents a higher order contribution. However the next-to-leading term b_1 is in general beyond theoretical control. $1/N_C$ expansions therefore do not enable us to decrease the uncertainties *systematically*.

The $N_C \to \infty$ prescription is certainly a very compact one with transparent rules, and it provides not a bad first approximation – but not more. One can ignore neither FSI nor WA completely.

9˙4.4. Treatment with QCD sum rules. In a series of papers [365] Blok and Shifman developed a treatment of $D_q \to PP$, PV decays based on a judicious application of QCD sum rules. They analyzed four-point correlation functions between the weak Lagrangian $\mathcal{L}_{weak}(\Delta C = 1)$ and three currents – one a pseudoscalar one generating D_q mesons and two axialvector or vector currents for the mesons in the final state. As usual an OPE is applied to the correlation function in the Euclidean region; nonperturbative dynamics is incorporated through condensates $\langle 0|m\bar{q}q|0\rangle$, $\langle 0|G \cdot G|0\rangle$ etc., the numerical values of which are extracted from other *light*-quark systems. Blok and Shifman extrapolate their results to the Minkowskian domain through a (double) dispersion relation . They succeed in finding a stability range for matching it with phenomenological hadronic expressions; hence they extract the decay amplitudes. For technical reasons their analysis does not extend to $D_q \to VV$ or axialvector resonances.

Their analysis has some nice features:
⊕ It has a clear basis in QCD, and includes, in principle at least, nonperturbative dynamics in a well-defined way.
⊕ It incorporates different quark-level processes – external and internal W emission, WA and PI – in a natural manner.
⊕ It allows to include nonfactorizable contributions systematically.

In practice, however, it suffers from some shortcommings:
⊖ The charm scale is not sufficiently high that one could have full confidence in the various extrapolations undertaken.
⊖ To make these very lengthy calculations at all managable, some simplifying assumptions had to be made, like $m_u = m_d = m_s = 0$ and $SU(3)_{Fl}$ breaking beyond $M_K > m_\pi$ had to be ignored; in particular $\langle 0|\bar{s}s|0\rangle = \langle 0|\bar{d}d|0\rangle = \langle 0|\bar{u}u|0\rangle$ was used. Thus one cannot expect $SU(3)_{Fl}$ breaking to be reproduced correctly.
⊖ *Prominent* FSI that vary rapidly with the energy scale – like effects due to narrow resonances - cannot be described in this treatment; for an extrapolation from the Euclidean to the Minkowskian domain amounts to some averaging or 'smearing' over energies.

A statement that the predictions do not provide an excellent fit to the data on about twenty-odd D^0, D^+ and D_s^+ modes – while correct on the surface, especially when $SU(3)_{Fl}$ breaking is involved – misses the main point:

- No a priori model assumption like factorization had to be made.
- The theoretical description does not contain any free parameters in principle, though in practice there is leeway in the size of some decay constants.

In summary: Overall a decent phenomenological description was achieved, yet realistically this framework provides few openings for *systematic* improvements.

Exp.	Decay Mode f	Norm. Mode f_{norm}	$\frac{BR(f)}{BR(f_{norm})}$ exp %	$\frac{BR(f)}{BR(f_{norm})}$ PDG02 %
		2 BODY		
	$D^0 \to$	$D^0 \to$		
FOCUS [366]	K^-K^+	$K^-\pi^+$	9.93 ± 0.20	10.83 ± 0.27
FOCUS [366]	$\pi^-\pi^+$	$K^-\pi^+$	3.53 ± 0.13	3.76 ± 0.17
CDF [367]	K^-K^+	$K^-\pi^+$	9.38 ± 0.20	
CDF prel.[367]	$\pi^-\pi^+$	$K^-\pi^+$	3.686 ± 0.084	
BELLE [368]	$K^+\pi^-$	$K^-\pi^+$	0.372 ± 0.027	0.39 ± 0.06
BELLE [369]	$K_L^0\pi^0$	$K_S^0\pi^0$	0.88 ± 0.13	–
	$\Lambda_c^+ \to$	$\Lambda_c^+ \to$		
FOCUS [370]	$\Sigma^+K^{*0}(892)$	$\Sigma^+\pi^+\pi^-$	7.8 ± 2.2	–
FOCUS [370]	$\Sigma^+\phi$	$\Sigma^+\pi^+\pi^-$	8.7 ± 1.7	8.7 ± 1.6
FOCUS [370]	$\Xi(1690)^0K^+$	$\Sigma^+\pi^+\pi^-$	2.2 ± 0.8	2.3 ± 0.7
		> 2 BODY		
	$D^+ \to$	$D^+ \to$		
FOCUS [371]	$K^-3\pi^+\pi^-$	$K^-2\pi^+$	5.8 ± 0.6	8.0 ± 0.9
FOCUS [371]	$3\pi^+2\pi^-$	$K^-3\pi^+\pi^-$	29.0 ± 2.0	–
FOCUS [371]	$K^+K^-2\pi^+\pi^-$	$K^-3\pi^+\pi^-$	4.0 ± 2.1	–
FOCUS [363]	$K^-K^+K^+$	$K^-\pi^+\pi^+$	$(9.49 \pm 2.18)\, 10^{-2}$	–
	$D_s^+ \to$	$D_s^+ \to$		
FOCUS [371]	$3\pi^+2\pi^-$	$K^+K^-\pi^+$	14.5 ± 1.4	15.8 ± 5.2
FOCUS [371]	$K^+K^-2\pi^+\pi^-$	$K^+K^-\pi^+$	15.0 ± 3.1	18.8 ± 5.4
FOCUS [371]	$\phi 2\pi^+\pi^-$	$\phi\pi^+$	24.9 ± 3.2	33 ± 6
FOCUS [363]	$K^+K^-K^+$	$K^+K^-\pi^+$	0.895 ± 0.310	< 1.6
	$\Lambda_c^+ \to$	$\Lambda_c^+ \to$		
FOCUS [370]	$\Sigma^+K^+K^-$	$\Sigma^+\pi^+\pi^-$	7.1 ± 1.5	8.1 ± 1.0
FOCUS [370]	$\Sigma^-K^+\pi^+$	$\Sigma^+K^{*0}(892)$	< 35(90% cl)	–
FOCUS [370]	$(\Sigma^+K^+K^-)_{nr}$	$\Sigma^+\pi^+\pi^-$	< 2.8 (90% cl)	< 1.8
CLEO [372]	$\Lambda\pi^+\omega$	$pK^-\pi^+$	24 ± 8	–
CLEO [372]	$\Lambda\pi^+\eta$	$pK^-\pi^+$	41 ± 20	35 ± 8
CLEO [372]	$\Lambda\pi^+(\pi^+\pi^-\pi^0)_{nr}$	$pK^-\pi^+$	< 13 (90% cl)	–
CLEO [372]	$\Lambda\pi^+(\pi^+\pi^-\pi^0)_{tot}$	$pK^-\pi^+$	36 ± 13	–

TABLE XII. – *Recent results not included in PDG02 on 2-body and >2-body charm mesons and baryons branching ratios. Errors are added in quadrature.*

9˙4.5. Status of the data. All charm experiments have important analyses projects ongoing, which address both twobody and multibody nonleptonic decays.

An often overlooked, yet relevant source of systematic errors arises from the — at times poor — knowledge of *absolute* branching ratios of the *normalizing* modes. The status of measurements of *absolute* branching ratios was discussed in Sect.**9**˙3. In Table XII we list some twobody and multibody modes together with their normalizing mode, recent data on the *relative* rates and the PDG02 values; in Table XIII we give values of their *absolute* branching ratios including our estimate of the uncertainty arising from the error in the branching ratio for the normalizing mode.

Exp.	Decay Mode	BR_{exp} %	$\frac{\sigma(f_{norm})}{\sigma(f)+\sigma(f_{norm})}$
	2 BODY		
	$D^0 \to$		
FOCUS [366]	K^-K^+	0.394 ± 0.013	1.2
FOCUS [366]	$\pi^-\pi^+$	0.138 ± 0.005	0.7
BELLE [368]	$K^+\pi^-$	0.014 ± 0.001	0.3
BELLE [276]	$\phi\pi^0$	0.0801 ± 0.0052	0.3
BELLE [276]	$\phi\eta$	0.0148 ± 0.0048	0.1
	$\Lambda_c^+ \to$		
FOCUS [370]	$\Sigma^+K^{*0}(892)$	0.28 ± 0.10	1.0
FOCUS [370]	$\Sigma^+\phi$	0.31 ± 0.08	1.5
FOCUS [370]	$\Xi(1690)^0K^+$	0.08 ± 0.03	0.8
	> 2 BODY		
	$D^+ \to$		
FOCUS [371]	$K^-3\pi^+\pi^-$	0.53 ± 0.06	0.7
FOCUS [371]	$3\pi^+2\pi^-$	0.212 ± 0.036	1.4
FOCUS [371]	$K^+K^-2\pi^+\pi^-$	0.029 ± 0.017	0.2
FOCUS [363]	$K^-K^+K^+$	$(0.86 \pm 0.19)\ 10^{-2}$	0.3
	$D_s^+ \to$		
FOCUS [371]	$3\pi^+2\pi^-$	0.64 ± 0.18	2.7
FOCUS [371]	$K^+K^-2\pi^+\pi^-$	0.66 ± 0.22	1.3
FOCUS [371]	$\phi 2\pi^+\pi^-$	0.90 ± 0.26	2.0
FOCUS [363]	$K^-K^+K^+$	0.039 ± 0.016	0.8
	$\Lambda_c^+ \to$		
FOCUS [370]	$\Sigma^+K^+K^-$	0.281 ± 0.069	1.3
FOCUS [370]	$(\Sigma^+K^+K^-)_{nr}$	< 0.06 (90% cl)	
CLEO [372]	$\Lambda\pi^+(\pi^+\pi^-\pi^0)_{tot}$	1.80 ± 0.64	0.7
CLEO [372]	$\Lambda\omega\pi^+$	1.20 ± 0.43	0.8
CLEO [372]	$\Lambda\eta\pi^+$	1.75 ± 0.49	0.5
CLEO [372]	$\Lambda\pi^+(\pi^+\pi^-\pi^0)_{nr}$	< 0.65 (90% cl)	

TABLE XIII. – *Recent results not included in PDG02 on 2-body and >2-body charm mesons and baryons absolute branching ratios. Last column shows ratio of relative errors on branching fractions of normalizing mode and decay mode.*

9˙4.6. Modern models. As the data improved, the BSW prescription became inadequate, however almost every subsequent attempt to describe two-body nonleptonic decays in the D system uses the assumption of naive factorization as a starting point.

At this point the following question arises: Why not just wait for lattice QCD (or some other calculational breakthrough) to gain us theoretical control over exclusive decays? For refining quark model predictions could be viewed somewhat unkindly like adding epicycles to a Ptolemaic system: while producing more accurate numbers, it would not deepen our understanding.

There are several reasons for not waiting idly:

- The wait might be quite long considering a quenched (or even partially unquenched) approximation is inadequate to include FSI.

- Even a merely phenomenological description of exclusive modes will be of great help for estimating the strength of $D^0 - \bar{D}^0$ oscillations in the SM, as explained later.
- As already stated, information on the strong phases due to FSI is instrumental for understanding *direct* CP violation.

Improvements (hopefully) and generalizations of the BSW description are made in three areas:

1. Different parametrizations for the q^2 dependence of the form factors are used and different evaluations of their normalization. This is similar to what was addressed in our discussion of exclusive semileptonic decays. One appealing suggestion has been to use only those expressions for form factors that asymptotically – i.e. for m_c, $m_s \to \infty$ – exhibit heavy quark symmetry.
2. Contributions due to WA and Penguin operators have been included.
3. Attempts have been made to incorporate FSI more reliably.

One can easily see that the form factors of BSW do not agree with QCD's HQS in the heavy quark limit. In this limit, the various form factors can be expressed in terms of one universal form factor. For example, Eq. (197) then reads as follows

$$\langle P(p_P)|V^\mu|D(p_D)\rangle = \xi(v \cdot v')(v + v')^\mu$$

with $v[v']$ refering to the velocity of the $D[P]$ meson. This expression allows us to relate the form factors F_1 and F_0:

$$\xi(v \cdot v') = \frac{2\sqrt{M_D M_P}}{M_D + M_P} F_1(q^2) = \frac{2\sqrt{M_D M_P}}{M_D + M_p} \frac{F_0(q^2)}{1 - q^2/(M_D + M_P)^2} \tag{213}$$

Assuming a simple pole dependence for both F_1 and F_0 is clearly inconsistent with the preceding expression. Yet it turns out that the q^2 dependence of the form factors – whether it is described by a pole, double pole or exponential – has a rather limited impact on the results – not surprisingly, since in D (unlike B) decays the q^2 range is quite limited. This is displayed in Table XIV.

	$F_0^{DK}(m_\pi^2)$	$F_0^{D\pi}(m_K^2)$	$F_1^{DK}(m_\rho^2)$	$F_1^{D\pi}(m_{K^*}^2)$	$A_0^{DK^*}(m_\pi^2)$	$A_0^{D\rho}(m_K^2)$
pole	0.76	0.72	0.84	0.79	0.74	0.69
multipole	-	-	0.91	0.92	0.74	0.72
exponential	0.75	0.78	0.85	0.94	0.83	0.96

TABLE XIV. – *Numerical values of pole, double pole and exponential form factors*

The much greater challenge is provided by items 2 and 3. As discussed in Sect.**6**˙4 WA is not the dominant engine driving the $D^0 - D^+$ lifetime difference. It had been predicted [256] to contribute about 10 - 20 % of the overall D^0 and D_s^+ widths, and their observed lifetime ratio shows that it indeed does. However it does not generate a $\sim$ 10 - 20 % contribution *uniformly* to all channels. Actually it could substantially enhance or suppress the widths of individual channels or even dominate them [255]. It would be interesting to see whether such 'exclusive footprints' of WA could be identified in the data.

There is a further complication beyond WA's strength varying greatly from channel to channel. WA always produces non-exotic final states, i.e. those with quantum numbers possible for a $\bar{q}_1 q_2$ $[q_1 q_2 q_3]$ combination in D $[\Lambda_c]$ decays. Exactly those channels are sensitive to the most prominent FSI, namely resonance effects. As far as *fully inclusive* widths, which are not subject to FSI, are concerned, WA has a well-defined meaning as an independent process (and enters as an independent $1/m_Q^3$ term in the HQE). Yet for exclusive nonleptonic modes the distinction between WA and FSI becomes blurred. Till one has established full theoretical control over FSI, its contributions and those of WA are indistinguishable in an *individual* channel. Only by carefully analyzing the whole pattern can one hope to arrive at some meaningful conclusions.

There exists a wide variety of ansätze in the literature differing in their parametrizations, treatment of non-factorizable contributions etc. The first detailed model of the post-BSW generation was presented in Ref.[373]. It is based on a modified (??) factorization prescription, fits the colour suppression factor ξ to the data and allows for prominent WA terms. FSI are implemented by allowing for rescattering among final states belonging to the same $SU(3)$ representation.

Using these assumption the authors fitted the then available data. The results are shown in Tables (XV-XVIII). We give there the data as listed in PDG96 (relevant for the time the analysis of Ref.[373] was performed) and PDG02. On the experimental side it is intriguing to see that the quoted errors do not always decrease in time and that changes in the central values by two sigma do occur.

On the theoretical side we would like to note that the overall agreement is quite good. However a comparison of the second and third columns shows that the inclusion of well chosen FSI is essential for this success; for proper perspective one should keep in mind that including FSI at present involves a considerable amount of 'poetic license'. There is a good side to the prominence of FSI as well: they are a conditio sine qua non for direct CP violation revealing itself in partial width asymmetries. Anticipating the later discussion of CP violation we have listed also the asymmetries predicted by Ref.[373] for Cabibbo forbidden modes in Tables (XVI,XVII). However there are some glaring discrepencies. In particular, the predicted branching ratio for $D_s^+ \to \rho^+ \eta$ is well above the *presently* measured value, when FSI are included. The predicted value for BR$(D_s^+ \to \rho^+ \eta')$ on the other hand appears well below the data.

A less ambitious approach has also been undertaken [374, 375]. One can decompose the various amplitudes for a given class of decays into four *topological* amplitudes. These diagrams account for final state interactions and are not actual Feynman graphs. The four diagrams are: (1) a colour-favoured external W-emission tree diagram T, (2) a colour-suppressed internal W-emission tree diagram C,(3) an exchange amplitude E, and (4) an annihilation amplitude A.

Using the available data, these amplitudes can be fitted for the Cabibbo allowed modes. At first glance, this may appear to be an empty analysis. However, by making various phenomenological assumptions one can begin to discuss nonfactorizable effects and final state interactions. For example, in Ref. ([375]) the various topological amplitudes are defined in a similar vein to the BSW approach, e.g.

$$T = \frac{G_F}{\sqrt{2}} V_{ud} V_{cs}^* a_1 f_\pi \left(M_D^2 - M_K^2\right)\big) F_0^{DK}\left(M_\pi^2\right) \tag{214}$$

Decay Channel		BR_{th}(no FSI) $\times 10^2$	BR_{th}(FSI) $\times 10^2$	BR_{exp} $\times 10^2$	Exp Ref
$D^0 \to K^-\pi^+$		5.35	3.85	3.83 ± 0.12	PDG96
				3.80 ± 0.09	PDG03
$D^0 \to K_s\pi^0$		0.44	0.76	1.05 ± 0.10	PDG96
				1.14 ± 0.11	PDG03
$D^0 \to K_s\eta$		0.13	0.45	0.35 ± 0.05	PDG96
				0.38 ± 0.05	PDG03
$D^0 \to K_s\eta'$		0.54	0.80	0.85 ± 0.13	PDG96
				0.94 ± 0.14	PDG03
$D^0 \to \bar{K}^{*0}\pi^0$		1.66	3.21	3.10 ± 0.40	PDG96
				2.8 ± 0.4	PDG03
$D^0 \to \rho^0 K_s$		1.01	0.45	0.60 ± 0.085	PDG96
				0.74 ± 0.15	PDG03
$D^0 \to K^{*-}\pi^+$		1.12	4.66	5.0 ± 0.4	PDG96
				6.0 ± 0.5	PDG03
$D^0 \to \rho^+ K^-$		9.62	11.2	10.8 ± 1.0	PDG96
				10.2 ± 0.8	PDG03
$D^0 \to \bar{K}^{*0}\eta$		1.21	0.47	1.90 ± 0.50	PDG96
				1.8 ± 0.9	PDG03
$D^0 \to \bar{K}^{*0}\eta'$		0.001	0.004	< 0.11	PDG96
				< 0.11	PDG03
$D^0 \to \omega K_s$		0.26	0.97	1.05 ± 0.20	PDG96
				1.1 ± 0.2	PDG03
$D^0 \to \phi K_s$		0.059	0.414	0.43 ± 0.05	PDG96
				0.47 ± 0.06	PDG03
$D^+ \to K_s\pi^+$		1.14	1.35	1.37 ± 0.15	PDG96
				1.39 ± 0.09	PDG03
$D^+ \to \bar{K}^{*0}\pi^+$		1.46	2.00	1.92 ± 0.19	PDG96
$D^+ \to \rho^+ K_s$		1.71	5.82	3.30 ± 1.25	PDG96
$D_s^+ \to K^+ K_s$		1.35	2.47	1.80 ± 0.55	PDG96
$D_s^+ \to \pi^+\eta$		4.53	1.13	2.0 ± 0.6	PDG96
				1.7 ± 0.5	PDG03
$D_s^+ \to \pi^+\eta'$		2.60	5.44	4.9 ± 1.8	PDG96
				3.9 ± 1.0	PDG03
$D_s^+ \to \rho^+\eta$		4.42	8.12	10.3 ± 3.2	PDG96
				3.80 ± 0.09	PDG03
$D_s^+ \to \rho^+\eta'$		1.08	2.46	12.0 ± 4.0	PDG96
				10.8 ± 3.1	PDG03
$D_s^+ \to \phi\pi^+$		2.51	4.55	3.6 ± 0.9	PDG96
				3.80 ± 0.09	PDG03
$D_s^+ \to \bar{K}^{*0} K^+$		5.27	4.81	3.4 ± 0.9	PDG96
				3.6 ± 0.9	PDG03
$D_s^+ \to K^{*+} K_s$		0.87	1.10	2.15 ± 0.7	PDG96
$D_s^+ \to \rho^0\pi^+$		0.012	0.01	< 0.29	PDG96
				< 0.07	PDG03
$D_s^+ \to \rho^+\pi^0$		0.023	0.01	—	
$D_s^+ \to \omega\pi^+$		0.023	0.20	0.27 ± 0.12	PDG96
				0.28 ± 0.11	PDG03

TABLE XV. – *Theoretical predictions and experimental values for the branching ratios of various Cabibbo-allowed decays.*

$D^+ \to$		BR_{th}(no FSI) $\times 10^2$	BR_{th}(FSI) $\times 10^2$	BR_{exp} $\times 10^2$	Exp Ref	a_{CP_3} (10^{-3})
$\pi^+\pi^0$		**0.186**	**0.185**	0.25 ± 0.07	PDG96	—
$\pi^+\eta$		**0.38**	**0.38**	0.75 ± 0.25	PDG96	-0.77
				0.3 ± 0.06	PDG03	
$\pi^+\eta'$		**0.058**	**0.768**	< 0.9	PDG96	+0.90
				0.5 ± 0.1	PDG03	
$K^+\bar{K}^0$		**1.49**	**0.763**	0.72 ± 0.12	PDG96	-0.52
				0.58 ± 0.06	PDG03	
$\pi^+\rho^0$		**0.012**	**0.104**	< 0.14	PDG96	-1.96
				0.104 ± 0.018	PDG03	
$\rho^+\pi^0$		**0.208**	**0.451**	—		+0.89
$\rho^+\eta$		**0.695**	**0.064**	< 1.2	PDG96	-1.60
				< 0.7	PDG03	
$\rho^+\eta'$		**0.004**	**0.122**	< 1.5	PDG96	~ 0
				< 0.5	PDG03	
$\pi^+\omega$		**0.124**	**0.038**	< 0.7	PDG96	-0.60
$\pi^+\phi$		**0.146**	**0.619**	0.61 ± 0.06	PDG96	-0.09
$K^+\bar{K}^{*0}$		**1.99**	**0.436**	0.42 ± 0.05	PDG96	+0.68
$K^{*+}\bar{K}^0$		**1.52**	**0.86**	3.0 ± 1.4	PDG96	-0.19
				3.1 ± 1.4	PDG03	
$D_s^+ \to$		BR_{th}(no FSI) $\times 10^2$	BR_{th}(FSI) $\times 10^2$	BR_{exp} $\times 10^2$	Exp Ref	a_{CP_3} (10^{-3})
$K^+\pi^0$		**0.222**	**0.146**	—		+1.07
$K^+\eta$		**0.046**	**0.299**	—		-0.05
$K^+\eta'$		**0.318**	**0.495**	—		-0.64
π^+K^0		**0.586**	**0.373**	< 0.8	PDG96	+0.48
$K^+\rho^0$		**0.952**	**0.198**	< 0.29	PDG96	+0.25
$K^0\rho^+$		**0.384**	**1.29**	—		+0.36
$K^{*+}\pi^0$		**0.0004**	**0.076**	—		-0.92
$K^{*0}\pi^+$		**0.191**	**0.444**	0.65 ± 0.28	PDG96	-0.75
$K^{*+}\eta$		**0.200**	**0.146**	—		-0.41
$K^{*+}\eta'$		**0.044**	**0.029**	—		-0.09
$K^+\omega$		**0.252**	**0.178**	—		-0.34
$K^+\phi$		**0.103**	**0.008**	< 0.05	PDG96	+1.79

TABLE XVI. – *Branching ratios for Cabibbo-forbidden decays of* D^+ *and* D_s^+ *mesons along with predictions for CP asymmetries*

Here a_1 is related to the coefficient of the same name discussed above. It consists of the naive factorization piece along with a piece describing nonfactorizable contributions.

$$a_1 = c_1 + c_2 \left(\frac{1}{N_c} + \chi_1 \right) \tag{215}$$

This analysis is again consistent with $\chi_1 = -1/N_C$ or $\xi \approx 0$. Final state rescattering

$D^0 \rightarrow$		BR_{th}(no FSI) $\times 10^2$	BR_{th}(FSI) $\times 10^2$	BR_{exp} $\times 10^2$	Exp Ref	a_{CP} (10^{-3})
$\pi^+\pi^-$		**0.505**	**0.152**	$\mathbf{0.152 \pm 0.011}$	PDG96	$\mathbf{-0.10}$
				$\mathbf{0.143 \pm 0.007}$	PDG03	
$\pi^0\pi^0$		**0.106**	**0.115**	$\mathbf{0.084 \pm 0.022}$	PDG96	+0.51
K^+K^-		**0.589**	**0.427**	$\mathbf{0.433 \pm 0.027}$	PDG96	-0.10
				$\mathbf{0.412 \pm 0.014}$	PDG03	
$K^0\bar{K}^0$		**0**	**0.108**	$\mathbf{0.13 \pm 0.04}$	PDG96	+0.26
				$\mathbf{0.071 \pm 0.019}$	PDG03	
$\pi^0\omega$		**0.013**	**0.003**	—		-0.01
$\pi^0\phi$		**0.127**	**0.105**	$\mathbf{< 0.14}$	PDG96	-0.04
$\eta\phi$		**0.080**	**0.080**	$\mathbf{< 0.28}$	PDG96	-0.15
$\pi^0\phi$				$\mathbf{0.0801 \pm 0.0052}$	[276]	
$\eta\phi$				$\mathbf{0.0148 \pm 0.0048}$	[276]	
$K^0\bar{K}^{*0}$		**0.031**	**0.052**	$\mathbf{< 0.16}$	PDG96	-0.56
				$\mathbf{< 0.17}$	PDG03	
$\bar{K}^0K^{*0}$		**0.031**	**0.062**	$\mathbf{< 0.08}$	PDG96	-0.65
				$\mathbf{< 0.09}$	PDG03	
K^-K^{*+}		**0.542**	**0.431**	$\mathbf{0.35 \pm 0.08}$	PDG96	-0.04
				$\mathbf{0.38 \pm 0.08}$	PDG03	
K^+K^{*-}		**0.178**	**0.290**	$\mathbf{0.18 \pm 0.08}$	PDG96	0.27
				$\mathbf{0.20 \pm 0.11}$	PDG03	
$\pi^0\eta$		**0.055**	**0.054**	—		-1.44
$\pi^0\eta'$		**0.174**	**0.175**	—		+0.89
$\eta\eta$		**0.171**	**0.093**	—		-0.51
$\eta\eta'$		**0.011**	**0.186**	—		-0.31
$\eta\rho^0$		**0.010**	**0.020**	—		-0.53
$\eta'\rho^0$		**0.007**	**0.008**	—		+0.01
$\eta\omega$		**0.002**	**0.209**	—		-0.02
$\eta'\omega$		**0.017**	**0.0002**	—		-3.66
$\pi^0\rho^0$		**0.199**	**0.216**	—		-0.01
$\pi^+\rho^-$		**0.442**	**0.485**	—		-0.43
$\pi^-\rho^+$		**1.45**	**0.706**	—		+0.34

TABLE XVII. – *Branching ratios for Cabibbo suppressed decays of D^0 mesons along with predictions for CP asymmetries*

is also discussed in this approach. Using the same resonance based rescattering model of Ref.[373] ([26]), it is found that the fit value for the exchange topology E is consistent with a vanishing quark level exchange contribution as expected from helicity suppression arguments.

This analysis can be extended to include singly and doubly Cabibbo-suppressed modes. In Ref. [376] it was shown that by simply rescaling the Cabibbo allowed topological amplitudes by the relevent CKM factor one finds relatively good agreement with the current experimental data. A notable exception is the mode $D^+ \rightarrow \bar{K}^0 K^{*+}$ discussed

([26])There is a disputed factor of 2 between these analyses in the definition of the rescattering phase.

Decay Channel	BR_{th}(no FSI)$\times 10^2$	BR_{th}(FSI) $\times 10^2$	$BR_{exp} \times 10^2$	Exp Ref
$D^0 \to K^+\pi^-$	0.028	0.033	0.029 ± 0.014	PDG96
			0.0148 ± 0.0021	PDG03
$D^0 \to K^{*0}\pi^0$	0.004	0.004	—	
$D^0 \to K^{*+}\pi^-$	0.033	0.038	—	
$D^+ \to K^+\pi^0$	0.044	0.055	—	
$D^+ \to K^{*+}\pi^0$	0.054	0.057	—	
$D^+ \to K^{*0}\pi^+$	0.040	0.027	< 0.019	PDG96
			0.036 ± 0.016	PDG03
$D^+ \to \phi K^+$	0.003	0.003	< 0.013	PDG96
$D^+ \to \rho^0 K^+$	0.027	0.029	< 0.06	PDG96
			0.025 ± 0.012	PDG03

TABLE XVIII. – *Theoretical predictions and experimental values for the branching ratios of various Double-Cabibbo forbidden decays.*

below. A large annihilation contribution is required to reach the current experiment value. Such a sizable contribution appears to be ruled out by bounds from other processes.

9˙4.7. On manifestations of New Physics. Up to now we have focussed on two motivations for detailed studies of charm decays:

- They help prepare us to search for New Physics in $D^0 - \bar{D}^0$ oscillations, the CP phenomenology in charm decays and in B decays.
- They can shed a novel light on the formation of low-mass hadronic states, as described in Sect.**9**˙5.1.

However one can raise the question whether exclusive D decays can by themselves reveal the intervention of New Physics. The rarer the mode, the better the a priori chance for such an effect. Doubly Cabibbo suppressed channels would offer the best chance, once Cabibbo suppressed ones the next best one. Of course we have to be cognizant of our limited theoretical control over hadronization and not jump to conclusion.

This debate has recently been joined by two groups [361, 377]. Close and Lipkin start by pointing at the "anomalously high branching ratios" for two Cabibbo suppressed transitions

$$\begin{aligned} \mathrm{BR}(D^+ \to K^*(892)^+\bar{K}^0) &= 3.2 \pm 1.5\% \\ \mathrm{BR}(D^+ \to K^*(892)^+\bar{K}^*(892)^0) &= 2.6 \pm 1.1\%\ , \end{aligned} \tag{216}$$

which are amazingly similar to the corresponding Cabibbo allowed branching ratios:

$$\begin{aligned} \mathrm{BR}(D^+ \to \rho^+\bar{K}^0) &= 6.6 \pm 2.5\% \\ \mathrm{BR}(D^+ \to \rho^+\bar{K}^*(892)^0) &= 2.1 \pm 1.3\%\ . \end{aligned} \tag{217}$$

In D^0 decays on the other hand the expected Cabibbo hierarchy is apparent; e.g.,

$$\mathrm{BR}(D^0 \to K^*(892)^+K^-) = 0.35 \pm 0.08\%$$

(218) $$\mathbf{BR}(D^0 \to \rho^+ K^-) = \mathbf{10.8 \pm 0.9\%}\ .$$

A note of skepticism is appropriate here: while the central values in Eq.(216) are certainly large, the size of the error bars does not allow a firm conclusion. But it is intriguing to follow the authors of Ref.[361] and speculate 'what if' the error bars were to shrink substantially, yet the central value to stay basically the same. It can serve as a case study for how such arguments go.

Branching ratios also are not the best yardstick here. For it was pointed out more than twenty years ago [378] that D^+ modes generated by $c \to s\bar{s}u$ – like $D^+ \to \bar{K}^0 K^+$ – are *less* than Cabibbo suppressed relative to the corresponding $c \to s\bar{d}u$ modes – like $D^+ \to \bar{K}^0 \pi^+$ – since the former in contrast to the latter (and to $c \to d\bar{d}u$ modes like $D^+ \to \bar{\pi}^0 \pi^+$) do *not* experience destructive PI. With such PI the dominant mechanism driving $\tau(D^+)/\tau(D^0) \simeq \mathbf{2.5}$ the branching branching *ratios* for the channels in Eq.(216) are enhanced by a factor $\sim \mathbf{2.5}$ due to 'known' physics. To have a fairer comparison of D^+ and D^0 branching ratios, one should 'recalibrate' $D^+ \to \bar{K}^{(*)} K^{(*)}$ branching ratios *downward* by a factor 2.5, which puts it around the 1% level for the two modes in Eq.(216). Even then they are still larger than the branching ratios for the corresponding D^0 modes and also for some similar D^+ modes, namely

(219) $$\mathbf{BR}(D^+ \to \bar{K}^0 K^+) = \mathbf{0.58 \pm 0.06\%}$$

The latter is indeed considerably larger than BR$(D^+ \to \bar{\pi}^0 \pi^+) = \mathbf{0.25 \pm 0.07}$, yet not by as much as one would expect based on comparing quark-level diagrams. The fact that some *individual* branching ratios fall outside the general pattern should be seen as prima facie evidence for the intervention of FSI in the form of resonance effects. It is hard to see how truly New Physics could effect D^0 and D^+ decays very differently, let alone individual D^+ modes on the same Cabibbo level.

If the central values of the branching ratios in Eq.(216) were substantially confirmed, the by far most likely explanation would be to 'blame' it on somewhat accidentally strong FSI. The authors of Ref.[377] arrive at the same conclusion based on their detailed and quite successful model for nonleptonic D decays. The reader can be forgiven if she feels that theorists tend to call on FSI as a very convenient 'deus ex machina' to bail them out of trouble. Yet the burden of the proof has to rest on the shoulders of those arguing in favour of New Physics. It is a worthwhile effort, though, to probe for violations of G parity in D decays as suggested in Ref.[361], more to further our education on QCD's dynamics than to establish New Physics.

9˙5. *Light-flavour spectroscopy from charm hadronic decays*. – Charm decays can be utilized also as a novel probe of *light*-flavour spectroscopy. We have already touched upon this point in our discussion of semileptonic decays. The situation is both richer and more complex in nonleptonic decays requiring special tools.

9˙5.1. Dalitz plot techniques. Dalitz plots were invented in 1953 [379] for treating the full complexity of decays into three body final states, specifically $K^+ \to \pi^+\pi^-\pi^+$. The main goal originally was to infer the spin and parity of the decaying particle from its decay products. Dalitz' work constituted an essential step in resolving the '$\tau - \theta$' puzzle, i.e. that $\tau^+ \to \pi^+\pi^-\pi^+$ and $\theta^+ \to \pi^+\pi^0$ represented decays of the same particle, namely the K^+ – at the 'price' that parity was violated in weak decays [380]. Dalitz showed that three body final states could conveniently be described through a two-dimensional scatter plot

$$d\Gamma(P \to d_1 d_2 d_3) \propto \frac{1}{M_P^3} |\mathcal{M}|^2 ds_{12} ds_{23} \; , \; s_{ij} \equiv (p_{d_i} + p_{d_j})^2 \; . \tag{220}$$

The density of the Dalitz plot will be constant if and only if the matrix element $\mathcal{M}$ is constant; any variation in $\mathcal{M}$ reveals a *dynamical* rather than kinematical effect.

Charm decays proceed in an environment shaped by many resonances and virulent FSI. Our previous discussion has already benefitted handsomely from Dalitz plot studies when listing branching ratios for $D \to PV$ channels like $D \to \bar{K}^* \pi$, $\bar{K} \rho$ etc. and comparing them with theoretical predictions. Yet the benefits – real and potential – extend further still. For better appreciation we sketch the relevant analysis tools. Two main approaches are used in charm physics to formalize the Dalitz plot fit function parametrization, namely the *isobar model* and the *K-matrix model*. Both have been strongly criticized by theorists, as we shall discuss later.

In the *isobar model* the overall amplitude is parametrized via a coherent sum of Breit-Wigner amplitudes for modelling the interfering resonances, each with amplitude $\mathcal{A}^{BW}$, and of a constant amplitude $\mathcal{A}_{\mathcal{NR}}$ for the nonresonant amplitude

$$\mathcal{A} = a_{NR} e^{i\delta_{NR}} + \sum_{j=1}^{n} a_j e^{i\delta_j} \mathcal{A}_j^{BW} \tag{221}$$

where the a parameters give the various relative contributions, and the phases δ account for FSI. Measurements of the phase shifts between different resonant components allow one to gauge the role of FSI and thus to shed some light onto the underlying weak decay dynamics. At tree level, the weak amplitudes are treated as real; phase shifts in the decay are due to FSI. In this picture, the decay to quasi-two body final states is viewed as an s-channel process (Fig. 30), whose propagator is represented by the complex amplitude Eq. (221). The explicit appearance of resonant amplitudes, as well as details of fit strategies, vary with the experimental group. The E791 collaboration (see, as instance, [381]) uses a combination of Breit-Wigner functions, complemented by momentum-dependent form factors which reflect the non-pointlike nature of the D meson and of the resonance, and an angular momentum term. An additional complication arises when the j-th resonance is kinematically allowed to decay to different channels. This is the typical case of f_0 which can decay to both $\pi\pi$ and KK. In this case one uses a coupled-channel Breit-Wigner (Flatte' formula) [384].

The isobar model is the tool normally used for extracting information from the Daliz plots, and a wealth of updated reviews can be found in the literature [298, 387, 388].

The use of isobar models for Dalitz plots in charm decays has, however, been subjected to multiple criticisms: doubts have been cast whether overlapping broad resonances can be represented correctly (which is closely connected to the issue of formulating the unitarity constraint in three and four-body final states), on how to relate the observations from charm to those from scattering, to formulate a coupled-channel treatments, etc.

Recently the use of *K-matrix*-inspired approaches has been advocated, in close connection to the experimental puzzle of the observation of the scalar σ resonance in $D_s \to \pi\pi\pi$ (for an experimentalist's point of view see [387, 388, 389]).

The K-matrix is a representation of the scattering matrix S, where the resonances are defined as poles of S. The K-matrix, originally developed in the context of scattering problems can be extended to cover the case of more complex resonance formation through the P-vector approach [385]. The propagator of Fig.30 is written $(I - iK \cdot \rho)^{-1}$ where K is the matrix for scattering of particles 1 and 2, I the identity matrix, and ρ the phase space matrix. Consequently, the amplitude (Eq. 221) is written as a coherent sum of a nonresonant term, Breit-Wigner terms for narrow, well isolated resonances, and K-matrix terms for broad overlapping resonances

$$\mathcal{A} = a_{NR} e^{i\delta_{NR}} + \sum_{j=1}^{n} a_j e^{i\delta_j} \mathcal{A}_j^{BW} + \sum_{k=n+1}^{m} a_k e^{i\delta_k} \mathcal{A}_k^{K} \tag{222}$$

Experimental information on poles is taken from $\pi\pi$ scattering data, and this allows one to treat coupled-channel decays such as $f_0 \to \pi\pi, KK$ straightforwardly.

9˙5.2. Results from Dalitz analyses. *The recent experimental studies of charm decays have opened up a new experimental window for understanding light meson spectroscopy and especially the controversial scalar mesons, which are copiously produced in these decays* Ref. [386]. This statement may be regarded as overly emphatic and excessively optimistic, but the plain truth is that the study of charm mesons hadronic decays via Dalitz plot formalisms is a field hosting a rich 'ecosystem' of diverse concepts involving heavy and light quarks, gluonia, hybrids or hermaphrodites etc. enlivened by seemingly endless debates on the proper formalism to treat the data.

For the light quark *aficionado*, charm decays have unique features that in principle make them ideally suited for light quark spectroscopy, i.e., a J=0, well defined D meson initial state, small nonresonant component, small background, large coupling to scalars, and independence from isospin and parity conservation [387, 388]. For the charm quark zealot, the study of nonleptonic quasi-two body decays via Dalitz plots is an essential tool to study the extent of final-state interaction effects. This is accomplished by means of the study of the interference among amplitudes which describe conflicting resonant channels. Furthermore finding a state in vastly different environments like charm decays and low energy hadronic collisions or photoproduction with consistent values for its mass and width strengthens considerably its claim for being a genuine resonance rather than a 'mere' threshold enhancement.

The data scenario on Dalitz plots analyses is quite rich, with contributions from all high-resolution, high-statistics charm experiment both at fixed target (E687, E791, FOCUS) and at low-energy e^+e^- colliders (CLEO, BABAR, BELLE). We shall outline the main features of the experimental scenario for $K\pi\pi$, $KK\pi$ and $\pi\pi\pi$, reporting a summary of measurements in Tables XIX,XX,XXI.

(i) Among the major channels, the mode $D^+ \to K^-\pi^+\pi^+$ is the only with an important nonresonant component. The fit fraction was also shown to largely exceed 100%, thus signalling very large constructive interference. A recent result from E791 claims that very much of the nonresonant fraction can be better represented by a broad scalar resonance, which they interpret as the κ. A possible link can be thought with the recent observation by FOCUS [309] of an anomalous interference effect in D^+ semileptonic decay which could be interpreted with the interference of a broad scalar (the κ ?) with the $K^*(890)$. Recent measurement by CLEO [397], however, fails in finding support for the κ hypothesis in Dalitz plot.

(ii) The mode $D^0 \to K_S\phi$ has played an interesting role in our efforts to identify the mechanisms driving D decays. Its existence was predicted with a branching ratio of $\sim 0.1 - 0.5$ % based on the assumption that WA creates most of the excess in the D^0

Decay Mode	Fit Fraction %	Amplitude	Phase ^{o}
	$D^+ \to K^-\pi^+\pi^+$ E791 2002 [381]		
$\bar{K}^*(892)^0\pi^+$	12.3 ± 1.3	1.0 fixed	0 fixed
NR	13.0 ± 7.3	1.03 ± 0.34	-11 ± 16
$\kappa\pi^+$	47.8 ± 13.2	1.97 ± 0.37	187 ± 20
$\bar{K}_0^*(1430)^0\pi^+$	12.5 ± 1.5	1.01 ± 0.13	48 ± 12
$\bar{K}_2^*(1430)^0\pi^+$	0.5 ± 0.2	0.20 ± 0.06	-54 ± 11
$\bar{K}^*(1680)^0\pi^+$	2.5 ± 0.8	0.45 ± 0.16	28 ± 0.20
	$D^0 \to K_s^0\pi^+\pi^-$ E687 1992 [382]		
$K^{*-}\pi^+$	64 ± 9	1.0 fixed	0 fixed
NR	26 ± 9	0.41 ± 0.09	-252 ± 34
$\bar{K}^0\rho^0$	20 ± 7	0.39 ± 0.07	-275 ± 57
	$D^0 \to K_s^0\pi^+\pi^-$ E687 1994 [394]		
$K^{*-}\pi^+$	62.5 ± 4.1		0 fixed
$K_0^*(1430)^-\pi^+$	10.9 ± 3.9		-166 ± 11
$\bar{K}^0\rho^0$	35.0 ± 7.3		-136 ± 6
$\bar{K}^0 f_0(975)$	6.8 ± 2.3		38 ± 11
$\bar{K}^0 f_2(1270)$	3.7 ± 2.2		-174 ± 23
$\bar{K}^0 f_0(1400)$	7.7 ± 3.6		-45 ± 24
	$D^+ \to K^-\pi^+\pi^+$ E687 1994 [394]		
$\bar{K}^{*0}\pi^+$	13.7 ± 1.0		48 ± 2
$\bar{K}_0^*(1430)^0\pi^+$	28.4 ± 3.9		63 ± 4
$\bar{K}^*(1680)^0\pi^+$	4.7 ± 0.6		73 ± 17
NR	99.8 ± 5.9		0 fixed
	$D^0 \to K^-\pi^+\pi^0$ E687 1994 [394]		
$K^-\rho^+$	76.5 ± 4.6		0 fixed
$K^{*-}\pi^+$	14.8 ± 5.6		162 ± 12
$\bar{K}^{*0}\pi^0$	16.5 ± 3.3		-2 ± 26
NR	10.1 ± 4.4		-122 ± 23
	$D^+ \to K^-\pi^+\pi^+$ E691 1993 [395]		
NR	83.8	1.0 fixed	0 fixed
$\bar{K}^*(892)^0\pi^+$	17.0 ± 0.9	0.78 ± 0.02	-60 ± 3
$\bar{K}_0^*(1430)^0\pi^+$	24.8 ± 1.9	0.53 ± 0.2	132 ± 2
$\bar{K}^*(1680)^0\pi^+$	3.0 ± 0.4	0.47 ± 0.03	-51 ± 4
	$D^0 \to K^-\pi^+\pi^0$ E691 1993 [395]		
NR	3.6	1.0 fixed	0 fixed
$\bar{K}^*(892)^0\pi^0$	14.2 ± 1.8	3.19 ± 0.20	167 ± 9
$\bar{K}^*(892)^-\pi^+$	8.4 ± 1.1	2.96 ± 0.19	-112 ± 9
$K^-\rho^+$	64.7 ± 3.9	8.56 ± 0.26	40 ± 7
	$D^0 \to K_s^0\pi^+\pi^-$ E691 1993 [395]		
NR	26.3	1.0 fixed	0 fixed
$K^*(892)^-\pi^+$	48.0 ± 9.7	2.3 ± 0.23	109 ± 9
$\bar{K}^0\rho^0$	21.5 ± 5.1	1.59 ± 0.19	-123 ± 12
	$D^+ \to \pi^-\pi^+\pi^+$ E791 2001 [390]		
$\rho^0(770)\pi^+$	33.6 ± 3.9	1.0 fixed	0 fixed
$\sigma\pi^+$	46.3 ± 9.2	1.17 ± 0.14	205.7 ± 9.0
NR	7.8 ± 6.6	0.48 ± 0.20	57.3 ± 20.3
$f_0(980)\pi^+$	6.2 ± 1.4	0.43 ± 0.05	165 ± 11
$f_2(1270)\pi^+$	19.4 ± 2.5	0.76 ± 0.07	57.3 ± 8.0
$f_0(1370)\pi^+$	2.3 ± 1.7	0.26 ± 0.09	105.4 ± 17.8
$\rho^0(1450)\pi^+$	0.7 ± 0.8	0.14 ± 0.07	319.1 ± 40.5
	$D^+ \to K^-K^+\pi^+$ E687 1995 [383]		
$\bar{K}^*(892)^0K^+$	30.1 ± 3.2		0 fixed
$\phi\pi^+$	29.2 ± 4.3		-159 ± 14
$\bar{K}_0^*(1430)^0K^+$	37.0 ± 3.9		70 ± 8

TABLE XIX. – *Dalitz plot coherent analyses results on* $K\pi\pi$ *and* $\pi\pi\pi$ *decays of* D. *Errors are summed in quadrature.*

Decay Mode	Fit Fraction %	Amplitude	Phase ^{o}
	$D^+ \to \pi^-\pi^+\pi^+$ E687 1997 [409]		
NR	58.9 ± 13.3	1 fixed	0 fixed
$\rho^0(770)\pi^+$	28.9 ± 8.0	0.70 ± 0.11	27 ± 18
$f_0(980)\pi^+$	2.7 ± 4.9	0.22 ± 0.13	197 ± 37
$f_2(1270)\pi^+$	5.2 ± 4.9	0.30 ± 0.11	207 ± 17
	$D^0 \to \bar{K}^0\pi^+\pi^-$ CLEO 2002 [397]		
$K^*(892)^+\pi^-$	0.34 ± 0.22	$(11 \pm 4)\,10^{-2}$	321 ± 14
$\bar{K}^0\rho^0$	26.4 ± 1.3	1.0 fixed	0 fixed
$\bar{K}^0\omega$	0.72 ± 0.20	$(37 \pm 7)\,10^{-3}$	114 ± 9
$K^*(892)^-\pi^+$	65.7 ± 3.1	1.56 ± 0.11	150 ± 4
$\bar{K}^0 f_0(980)$	4.3 ± 1.0	0.34 ± 0.05	188 ± 9
$\bar{K}^0 f_2(1270)$	0.27 ± 0.29	0.7 ± 0.5	308 ± 42
$\bar{K}^0 f_0(1370)$	9.9 ± 3.1	1.8 ± 0.2	85 ± 20
$K_0^*(1430)^-\pi^+$	7.3 ± 2.2	2.0 ± 0.3	3 ± 11
$K_2^*(1430)^-\pi^+$	1.1 ± 0.5	1.0 ± 0.24	155 ± 17
$K^*(1680)^-\pi^+$	2.2 ± 1.6	5.6 ± 4	174 ± 17
NR	0.9 ± 1.2	1.1 ± 0.9	160 ± 57
	$D^0 \to K^-\pi^+\pi^0$ MARK III 1987 [391]		
$K^-\rho^+$	81 ± 7		0 fixed
$K^{*-}\pi^+$	12 ± 4		154 ± 11
$\bar{K}^{*0}pi^0$	13 ± 4		7 ± 7
NR	9 ± 4		52 ± 9
	$D^0 \to \bar{K}^0\pi^+\pi^-$ MARK III 1987 [391]		
$\bar{K}^0\rho^0$	12 ± 7		93 ± 30
$K^{*-}\pi^+$	56 ± 6		0 fixed
NR	33 ± 11		$--$
	$D^+ \to \bar{K}^0\pi^+\pi^0$ MARK III 1987 [391]		
$\bar{K}^0\rho^+$	68 ± 14		0 fixed
$\bar{K}^{*0}\pi^+$	19 ± 8		43 ± 23
NR	13 ± 11		250 ± 19
	$D^+ \to \bar{K}^-\pi^+\pi^+$ MARK III 1987 [391]		
$\bar{K}^{*0}\pi^+$	13 ± 7		105 ± 8
NR	79 ± 16		0.0
	$D^0 \to K^0K^-\pi^+$ BABAR 2002 [426]		
$\bar{K}_0^*(1430)^0K^0$	4.8 ± 2.1		52 ± 27
$\bar{K}_1^*(892)^0K^0$	0.8 ± 0.5		175 ± 22
$\bar{K}_1^*(1680)^0K^0$	6.9 ± 1.6		-169 ± 16
$\bar{K}_2^*(1430)^0K^0$	2.0 ± 0.6		51 ± 18
$K_0^*(1430)^+K^-$	13.3 ± 5.2		-41 ± 25
$K_1^*(892)^+K^-$	63.6 ± 5.7		0 fixed
$K_1^*(1680)^+K^-$	15.6 ± 3.3		-178 ± 10
$K_2^*(1430)^+K^-$	13.8 ± 8.3		-52 ± 7
$a_0(980)^-\pi^+$	2.9 ± 2.4		-100 ± 13
$a_0(1450)^-\pi^+$	3.1 ± 2.1		31 ± 16
$a_2(1310)^-\pi^+$	0.7 ± 0.4		-149 ± 27
NR	2.3 ± 5.6		-136 ± 23
	$D^0 \to \pi^+\pi^-\pi^0$ CLEO 2003 [396]		
$\rho^+\pi^-$	76.5 ± 5.1	1 fixed	0 fixed
$\rho^0\pi^0$	23.9 ± 5.0	0.56 ± 0.07	10 ± 4
$\rho^-\pi^+$	32.3 ± 3.0	0.65 ± 0.05	-4 ± 5
NR	2.7 ± 1.9	1.03 ± 0.35	77 ± 14

TABLE XX. – *Dalitz plot coherent analyses results on $K\pi\pi$, $KK\pi$ and $\pi\pi\pi$ decays of D. Errors are summed in quadrature.*

Decay Mode	Fit Fraction %	Amplitude	Phase ^{o}
	$D^0 \to \bar{K}^0 K^+ \pi^-$ BABAR 2002 [426]		
$K_0^*(1430)^0 \bar{K}^0$	26 ± 16.3		-38 ± 22
$K_1^*(892)^0 \bar{K}^0$	2.8 ± 1.5		-126 ± 19
$K_1^*(1680)^0 \bar{K}^0$	15.2 ± 12.1		161 ± 9
$K_2^*(1430)^0 \bar{K}^0$	1.7 ± 2.5		53 ± 38
$K_0^*(1430)^- K^+$	2.4 ± 8		-142 ± 115
$K_1^*(892)^- K^+$	35.6 ± 8		0 fixed
$K_1^*(1680)^- K^+$	5.1 ± 6		124 ± 27
$K_2^*(1430)^- K^+$	1 ± 1		-26 ± 38
$a_0^+(980)\pi^-$	15.1 ± 12.1		-160 ± 42
$a_0^+(1450)\pi^-$	2.2 ± 2.9		148 ± 25
NR	37 ± 26		-172 ± 13
	$D^0 \to \bar{K}^0 K^+ K^-$ BABAR 2002 [426]		
$\bar{K}^0 \phi$	45.4 ± 1.9		0 fixed
$\bar{K}^0 a_0(980)^0$	61 ± 15		109 ± 15
$\bar{K}^0 f_0(980)$	12.2 ± 9.1		-161 ± 14
$a_0(980)^+ K^-$	34.3 ± 7.5		-53 ± 4
$a_0(980)^- K^+$	3.2 ± 3.9		-13 ± 15
NR	0.4 ± 0.8		40 ± 44
	$D_s^+ \to \pi^-\pi^+\pi^+$ E791 2000 [393]		
$f_0(980)\pi^+$	54.1 ± 4.0	1.0 fixed	0 fixed
$\rho^0(770)\pi^+$	11.1 ± 2.5	0.45 ± 0.06	81 ± 15
NR	5.0 ± 3.8	0.30 ± 0.12	149 ± 25
$f_2(1270)\pi^+$	20.8 ± 3.0	0.62 ± 0.05	124 ± 11
$f_0(1370)\pi^+$	34.7 ± 7.2	0.80 ± 0.11	159 ± 14
$\rho^0(1450)\pi^+$	0	0	0
	$D_s^+ \to \pi^-\pi^+\pi^+$ E687 1997 [409]		
$f_0(980)\pi^+$	107.4 ± 14.6	1 fixed	0 fixed
NR	12.1 ± 12.3	0.34 ± 0.14	235 ± 22
$\rho^0(770)\pi^+$	2.3 ± 2.9	0.15 ± 0.09	53 ± 45
$f_2(1270)\pi^+$	12.3 ± 5.9	0.34 ± 0.09	100 ± 19
$S(1475)\pi^+$	27.4 ± 11.5	0.50 ± 0.13	234 ± 15
	$D_s^+ \to K^- K^+ \pi^+$ E687 1995 [383]		
$\bar{K}^*(892)^0 K^+$	47.8 ± 6.1		0 fixed
$\phi\pi^+$	39.6 ± 5.7		178 ± 31
$f_0(980)\pi^+$	11.0 ± 4.4		159 ± 27
$f_J(1710)\pi^+$	3.4 ± 4.2		110 ± 26
$\bar{K}_0^*(1430)^0 K^+$	9.3 ± 4.5		152 ± 56

TABLE XXI. – *Dalitz plot coherent analyses results on $KK\pi$ and $\pi\pi\pi$ decays of D_s. Errors are summed in quadrature.*

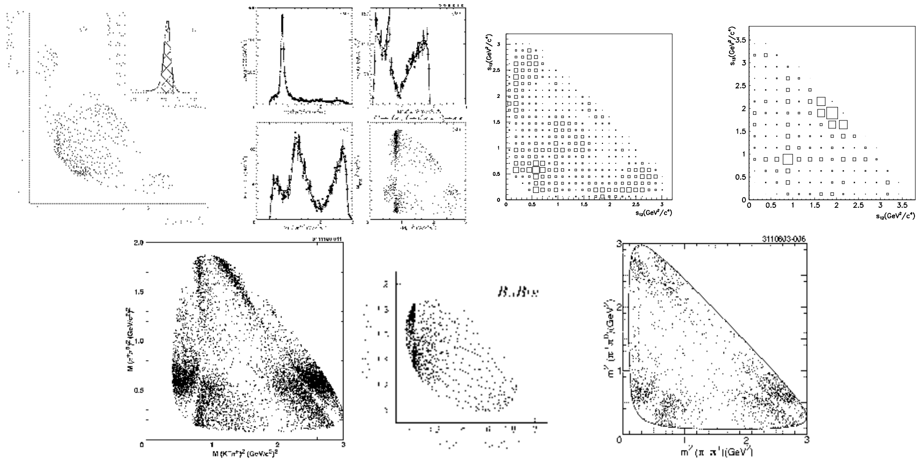


Fig. 29. – *A selection of Dalitz plots of charm meson decays: (a) E791 $D^+ \to K^-\pi^+\pi^+$ [381] ; (b) CLEO $D^0 \to K^0_s\pi^+\pi^-$ [397] ; (c) E791 $D^+ \to \pi^-\pi^+\pi^+$ [390] ; (d) E791 $D_s^+ \to \pi^-\pi^+\pi^+$ [393] ; (e) CLEO $D^0 \to K^-\pi^+\pi^0$ [398]; (f) BABAR $D^0 \to K^0_s K^+K^-$ [426]; (g) CLEO $D^0 \to \pi^-\pi^+\pi^0$ [396].*

over the D^+ width [362]. It was claimed that finding it would amount to a 'smoking gun' evidence supporting this assumption. Afterwards it was indeed found – with close to the predicted branching ratio: BR$(D^0 \to K_S\phi) = (0.47 \pm 0.06)$ %. In retrospect, however, it should be viewed as evidence for the impact FSI can have.

There is, though, some lesson of future relevance we can learn. The state $K_S\phi$ is an odd CP eigenstate; as such it provides intriguing ways to search for CP violation in $D^0 \to K_S\phi$ as explained later. The corresponding beauty mode $B_d \to K_S\phi$ is under active study for similar reasons at the B factories. Yet when extracting $K_S\phi$ from $K_SK^+K^-$ final states one has to contend with the scalar $\bar{K}K$ resonance $f_0(980)$ close to the ϕ mass. This problem is aggravated by the fact that $K_Sf_0(980)$ has the *opposite* CP parity of $K_S\phi$. A CP asymmetry in $D^0 \to K_SK^+K^-$ coming from $D^0 \to K_S\phi$ would then be (partially) cancelled by one due to $D^0 \to K_Sf_0(980)$. Likewise for $B_d \to K_SK^+K^-$.

(iii) The $D^+ \to KK\pi$ Dalitz plot shows the presence of various K^* states, with the asymmetry between $\bar{K}^*$ 'lobes' being interpreted [298] as an interference of a broad S=0 resonance with the $\bar{K}^{*0}K^+$. The $D_s^+ \to KK\pi$ Dalitz plot is strongly dominated by the $\phi(1020)$. However we know that $D_s^+ \to \pi\pi\pi$ is dominated by $f_0(980)$ (see below), which also decays to KK. It is clear that the ϕ contribution overlaps with the $f_0(980)$ contributions, and the two should be disentangled. This has far reaching consequences and implications, also in B physics, as we shall discuss later on in this section.

(iv) The issues involved here – resonance vs. threshold enhancement, Breit-Wigner vs. *non*-Breit-Wigner form of the excitation curve, isobar vs. K matrix model – are being debated with particular passion in connection with the role of the σ resonance in $D^+, D_s^+ \to 3\pi$ modes.

The σ has a checkered past: after being introduced as an s-wave scalar resonance in $\pi\pi$ scattering its mass and width – even its name – have changed over time. It has been never *conclusively* observed in $\pi\pi$ scattering, and its relevance is connected to the role

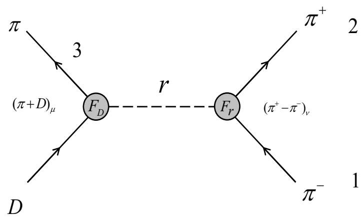


Fig. 30. – *Charm meson decay represented via an s-channel process via resonance propagator* $\boldsymbol{r}$.

of Higgs-like particle in Nambu–Jona-Lasinio models. Evidence for the $\boldsymbol{\sigma}$'s existence was recently obtained by BES in $\boldsymbol{J/\psi \to \omega\pi}$ data. Finding it in $\boldsymbol{D_s \to 3\pi}$ would provide a nontrivial boost to its resonance status.

One has to keep in mind that the dynamical stage for $\boldsymbol{D_s^+ \to 3\pi}$ and $\boldsymbol{D^+ \to 3\pi}$ is actually quite different even beyond the fact that the former is Cabibbo allowed and the latter Cabibbo suppressed:

- $\boldsymbol{D_s^+ \to 3\pi}$ can be generated by WA as well by the leading decay process coupled with $\boldsymbol{K\bar{K}\pi \Rightarrow 3\pi}$ rescattering. Its final state carries $\boldsymbol{I = 1}$.

- $\boldsymbol{D^+ \to 3\pi}$ on the other hand can proceed via (i) WA, (ii) the $\boldsymbol{\Delta I = 1/2}$ and (iii) $\boldsymbol{\Delta I = 3/2}$ components of the quark decay process. In the first two cases the final state isospin is $\boldsymbol{1}$, in the last one $\boldsymbol{I = 2}$.

These general facts can be translated into more channel-specific statements. Given the relevant strange quark content of $\boldsymbol{D_s}$, resonances involved in its decays should couple to both $\boldsymbol{KK}$ and $\boldsymbol{\pi\pi}$, and obvious candidates are the $\boldsymbol{f_0}$'s. Figure 31 shows the diagrams such a decay can proceed through. The spectator-like diagram involves the resonance able to couple simultaneously to a couple of strange quarks, and to a couple of pions. The other two processes are WA-like; while those are reduced in inclusive rates, they can be quite strong in some exclusive modes. Evidence of such a WA component has been recently put forward by E791 which measures a relative fit fraction of about 10% for $\boldsymbol{\rho\pi}$. Interference between the WA processes has also been discussed. In general, the $\boldsymbol{D_s^+ \to \pi\pi\pi}$ is dominated by the $\boldsymbol{f_0(980)}$ narrow resonances, with very little nonresonant component left over. On the other hand, the $\boldsymbol{D^+ \to \pi\pi\pi}$ is characterized by a very large either nonresonant, or broad resonant component. Thus one expects to find a quite different Dalitz plot population in $\boldsymbol{D_s^+ \to 3\pi}$ and $\boldsymbol{D^+ \to 3\pi}$ – and this is indeed borne out by the data. They reveal for D and $\boldsymbol{D_s}$ charm mesons a rich wealth of resonance substructures (Fig. 29). The evidence can be visualised by the cartoon of a quasi-two body charm meson decay proceeding effectively via a $\boldsymbol{D \to \pi^+}$ current (Fig.30), interacting to a $\boldsymbol{P_1P_2}$ current (such as $\boldsymbol{K^+K^-}$, $\boldsymbol{\pi^+\pi^-}$, $\boldsymbol{K^-\pi^+}$) through a strongly-decaying resonance propagator. The interpretation is, however, still in dispute. E687 was able to satisfactorily fit the Dalitz plot area with a 80% fit fraction of nonresonant. E791 claims the observation of a large $\boldsymbol{\sigma\pi^+}$ component in $\boldsymbol{D^+ \to \pi\pi\pi}$. Analyses based on the K-matrix ansatz do recognize the broad, low-mass excess, but claim to be able to explain it purely via $\boldsymbol{\pi\pi}$ scattering amplitudes [399].

There are several possible reasons for this discrepancy:

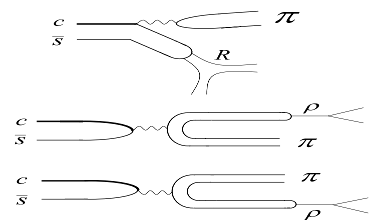


Fig. 31. – Diagrams for $D_s^+ \to \pi\pi\pi$

- It might be due to the excess in the E791 and BES data simply reflecting an attractive enhancement (rather than a full fledged resonance) in $\pi\pi$ scattering.
- The K-matrix approach might be of limited validity at small invariant energies.
- General caveats have been put forward towards extracting resonance parameters from charm decay data when referring to scalars. Only the low energy tails of the resonance phase shifts are visible in charm data and noone has ever observed a complete σ nor κ Breit-Wigner phase motion through 180^o. Novel approaches for testing the phase shift over the Dalitz plot area have been suggested[400], but so far they are not applicable due to severe statistics limitations of datasets. We also remind the reader that Breit-Wigner parameters and pole positions can fairly easily shift by hundred of MeVs, in particular for broad resonance candidates such as σ or κ.
- A Breit-Wigner resonance form is an approximation of varying accuracy; in particular for scalar di-meson resonances a lot of information exists from data that a Breit-Wigner ansatz provides in general a poor description [401]. More work is needed here, and theorists are asked to spend some quality time on this problem.

The last item has an important consequence also for B decays [401]: Whatever the structure of σ is, either a genuine resonance or a strong attractive S wave enhancement at low di-pion masses, it will affect the important mode $B \to \rho\pi$ to be used for extracting the angle ϕ_2/α in the unitarity triangle for the scalar form factor that one should adopt to represent $\pi\pi$ resonances such as the σ is very different from the normally used Breit-Wigner parametrization with a running width.

Finally, a far-reaching effect of the σ puzzle is under the eyes of low-energy hadronic community, and it regards the recent measurement of radiative decay $\Gamma(\phi \to \gamma f_0)$ at the DAΦNE ϕ-factory [402] in the $\pi^0\pi^0\gamma$ final state. In this analysis the presence of a σ is essential ingredient to the fit of data. A recent analysis [403] refits the KLOE data by using the K-matrix formalism and thus avoiding to employ explicity the σ, finding an order of magnitude smaller branching ratio.

There are also ulterior reasons for a detailed understanding of Dalitz plots: as discussed in Sect. **11** larger CP asymmetries might surface in them than in fully integrated widths.

9·6. *Baryon decays*. – All the issues addressed above for meson decays have their analogue in the decays of charm baryons.

- One wants to know absolute branching ratios for Λ_c and Ξ_c for the purpose of *charm counting* in B decays.

- Knowing the relative importance of $\Lambda_c \to \Lambda + X$ vs. $\Lambda_c \to N\bar{K} + X$ and $\Xi_c \to \Xi + X$ vs. $\Xi_c \to \Lambda\bar{K} + X$ is of great help in identifying Λ_b and Ξ_b decays.
- The final states in Λ_c and Ξ_c decays can shed new light on the spectroscopy of light flavour baryons.

Quasi-two-body modes of Λ_c and Ξ_c can be described with the same tools as D decays, namely quark models, QCD sum rules and lattice QCD. Yet the baryon decays pose even more formidable theoretical challenges. There are more types of contributions all on the same level, namely destructive as well as constructive PI and even WS with the latter *not* suffering from helicity suppression. Evaluating matrix elements of the relevant operators for baryons is even harder than for mesons.

Let us merely comment here on attempts to search for footprints of WS in certain exclusive channels. An early example is provided by $\Lambda_c^+ \to \Delta^{++}K^-$. Having an only slightly reduced branching ratio, it might be cited [404] as specific evidence for WS since that mechanism produces it naturally, whereas quark decay does not lead to it *directly*. Yet one encounters the usual problem with interpreting a transition as due to WS: it can be produced by the quark decay reaction followed by FSI. Until we gain full control over the nonperturbative dynamics in exclusive transitions we cannot distinguish the two explanations on a case-by-case is.

A similar comment applies to the Cabibbo suppressed (CS) mode $\Xi_c^+ \to pK^-\pi^+$ first seen by SELEX [405] (confirmed shortly thereafter [406]):

$$B(\Xi_c^+ \to pK^-\pi^+)/B(\Xi_c^+ \to \Sigma^+(pn)\,K^-\pi^+) = 0.22 \pm 0.06 \pm 0.03\ ; \tag{223}$$

once corrected for phase space this number is compatible with the branching ratio for the only other CS decay well measured, $\Lambda_c^+ \to pK^-K^+$, relative to three– body CF decay $\Lambda_c^+ \to pK^-\pi^+$.

9˙7. *Resume*. – A decent theoretical description of nonleptonic (quasi-)two-body modes in D decays has emerged. Yet we owe this success not completely – maybe not even mainly – to our ingenuity: nature was kind enough to present us with a relatively smooth dynamical environment considering how disruptive FSI could have been. On the other hand we seem to have reached a point of quickly diminishing returns. Treatments based on the quark model and on QCD sum rules probably have been pushed as far as they go without any qualitatively new theoretical insights on the limits of factorization, on how to go beyond it and on FSI. One is entitled to appeal to lattice QCD for more definitive answers, yet it might be quite a while before lattice QCD can deal with exclusive nonleptonic decays involving pions and kaons *without* the quenched approximation.

There are two main motivations beyond professional pride why we would like to do better in our understanding of nonleptonic charm decays: (i) It is desirable to obtain a reliable estimate for the strength of $D^0 - \bar{D}^0$ oscillation generated by SM dynamics to interprete experimental findings; (ii) for similar reasons we want to obtain reliable predictions on CP asymmetries in charm decays due to SM and New Physics dynamics. These issues will be discussed in more detail later on.

A less profound, yet still interesting question to study is which fraction of the weak decays of a given hadron are of the two-body or quasi-two-body type. This is however much easier said than done due to interference effects between different resonant final states. If one merely adds up the quoted widths for the different contributions to a final state, one typically either over- or under-saturates the overall widths, because the

TABLE XXII. – *Fraction (%) of partial widths for charmed particles identified in exclusive modes, and fraction of exclusive modes identified as two or quasitwo body (hadronic) and three or quasithree (semileptonic). Data from PDG2002. No absolute branching fractions have been measured for Ξ_C^+. No branching fractions have been measured for Ξ_C^0 nor Ω_C^0.*

	inclusive $(e+\mu+h)X$	excl. fract. (semi)leptonic	excl. fract. 3B or quasi-3B	excl. fract. nonleptonic nonleptonic	excl. fract. 2B or quasi-2B
D^+	106 ± 21	21± 3	24± 3	42± 3	40± 4
D^0	113 ± 23	10 ± 1	9.3± 0.4	54 ± 2	57 ± 3
D_s^+	160 ± 70	12± 2	5± 1	64 ± 8	57 ± 6
Λ_C^+	150 ± 50	4.2± 0.8	4.2± 0.8	44 ± 9	9±2

different components in general induce destructive or constructive interference between them. Examples can be found in Tables XIX - XXI. There is some discussion on this point in [407, 408].

It would make sense for PDG to list not merely the branching ratios, but also the amplitudes including their phases for Dalitz plot analyses. A way of addressing the issue, if not of solving it, is to analyze for each final state and resonant structure the magnitude of interference.

10. – $D^0 - \bar{D}^0$ Oscillations

Very often the two notions 'oscillations' and 'mixing' are employed in an interchangable way. Both are concepts deeply embedded in quantum mechanics, yet we want to distinguish them in our review. *Mixing* means that classically distinct states are not necessarily so in quantum mechanics and therefore can interfere. For example in atomic physics the weak neutral currents induce a 'wrong' parity component into a wavefunction, which in turn generate parity odd observables. Or mass eigenstates of quarks (or leptons) contain components of different flavours giving rise to a non-diagonal CKM (or PMNS) matrix. Such mixing creates a plethora of observable effects. The most intriguing ones arise when the violation of a certain flavour quantum number leads to the *stationary* or *mass* eigenstates not being *flavour* eigenstates. This induces *oscillations*, i.e. particular transitions, of which those between neutral mesons and anti-mesons of a given flavour (as discussed below), neutrinos of different flavours (including neutrino-antineutrino transitions due to Majorana terms) and neutron and antineutrons are the most discussed examples: a beam initially containing only one flavour 'regenerates' other flavours at later times through oscillatory functions of time, namely sine and cosine terms. One furthermore distinguishes between *spontaneous* regeneration or oscillations in vacuum and regeneration when the oscillating state has to transverse matter.

Mixing is thus a necessary, though not sufficient condition for oscillations with their peculiar time dependence to occur.

Flavour-changing $\Delta S, \Delta C, \Delta B \neq 0$ weak interactions have a two-fold impact on neutral heavy mesons. In addition to driving their decays, they induce $K^0 - \bar{K}^0$, $D^0 - \bar{D}^0$, $B^0 - \bar{B}^0$ oscillations. I.e., the two mass eigenstates of neutral heavy flavour mesons are linear combinations of the flavour eigenstates; while no longer carrying a definite flavour quantum number, they possess *non*degenerate masses and lifetimes. Its most striking signature is that an initially pure beam of, say, B^0 mesons will not only lose intensity due to B^0 decaying away, but also change its flavour identity over to $\bar{B}^0$, then go back to B^0 and so on. CP violation in the underlying dynamics can manifest itself in a variety of ways as described later. These generic statements apply to $K^0 - \bar{K}^0$, $B^0 - \bar{B}^0$ as well as $D^0 - \bar{D}^0$.

On the practical there are large differences, though. Unlike for strange and B mesons, $D^0 - \bar{D}^0$ oscillations are predicted to proceed quite slowly within the Standard Model. Searching for them has been recognized for a long time as a promising indirect probe for New Physics. Yet with the experimental sensitivity having reached the few percent level, we have to carefully evaluate the reliability of our theoretical treatment of charm transitions. For on one hand charm hadrons – in contrast to kaons – are too heavy to have only a few decay channels; on the other hand – unlike B mesons – we cannot be confident of the applicability of the Heavy Quark Expansion (HQE) to charm decays, although it has been proposed. This problem is particularly serious for $D^0 - \bar{D}^0$ oscillations; over the last twenty years vastly differing predictions have appeared in the literature [410]. A model-independent treatment becomes highly desirable even if it provides us with a mainly qualitative understanding rather than precise numbers.

10˙1. *Notation.* – The time evolution of neutral D mesons is obtained from solving the (free) Schrödinger equation

$$i\frac{d}{dt}\begin{pmatrix} D^0 \\ \bar{D}^0 \end{pmatrix} = \begin{pmatrix} M_{11} - \frac{i}{2}\Gamma_{11} & M_{12} - \frac{i}{2}\Gamma_{12} \\ M_{12}^* - \frac{i}{2}\Gamma_{12}^* & M_{22} - \frac{i}{2}\Gamma_{22} \end{pmatrix}\begin{pmatrix} D^0 \\ \bar{D}^0 \end{pmatrix} \tag{224}$$

CPT invariance imposes

$$M_{11} = M_{22} \quad , \quad \Gamma_{11} = \Gamma_{22} \; . \tag{225}$$

The mass eigenstates obtained through diagonalising this matrix are given by (for details see [411, 107])

$$|D_L\rangle = \frac{1}{\sqrt{|p|^2 + |q|^2}}\left(p|D^0\rangle + q|\bar{D}^0\rangle\right)$$

$$|D_H\rangle = \frac{1}{\sqrt{|p|^2 + |q|^2}}\left(p|D^0\rangle - q|\bar{D}^0\rangle\right) \tag{226}$$

$$\frac{q}{p} = \sqrt{\frac{M_{12}^* - \frac{i}{2}\Gamma_{12}^*}{M_{12} - \frac{i}{2}\Gamma_{12}}} \tag{227}$$

with differences in mass and width:

$$\Delta M_D \equiv M_H - M_L = -2\mathrm{Re}\left[\frac{q}{p}(M_{12} - \frac{i}{2}\Gamma_{12})\right] \tag{228}$$

$$\Delta\Gamma_D \equiv \Gamma_L - \Gamma_H = -2\mathrm{Im}\left[\frac{q}{p}(M_{12} - \frac{i}{2}\Gamma_{12})\right] \tag{229}$$

The labels $\boldsymbol{H}$ and $\boldsymbol{L}$ are chosen such that $\Delta M_D > 0$. Once this *convention* has been adopted, it becomes a sensible question whether

(230) $$\Gamma_H > \Gamma_L \quad \text{or} \quad \Gamma_H < \Gamma_L$$

holds, i.e. whether the heavier state is shorter or longer lived. Note that the subscripts $\boldsymbol{H}$ and $\boldsymbol{L}$ have been swapped in going from ΔM_D to $\Delta\Gamma_D$; this is done to have the analogous quantities positive for kaons.

The time evolution of the 'stationary' states $|D_{H,L}\rangle$ is then as follows

(231) $$|D_{H,L}(t)\rangle = e^{-\frac{1}{2}\Gamma_{H,L}t}e^{-iM_{H,L}t}|D_{H,L}\rangle$$

The probability to find a $\bar{D}^0$ at time t in an initially pure D^0 beam is given by

(232) $$|\langle\bar{D}^0|D^0(t)\rangle|^2 = \frac{1}{4}\left|\frac{q}{p}\right|^2 e^{-\Gamma_H t}\left(1 + e^{-\Delta\Gamma_D t} - 2e^{-\frac{1}{2}\Delta\Gamma_D t}\cos\Delta M_D t\right)$$

While the flavour of the initial meson is tagged by its production, the flavour of the final meson is inferred from its decay.

There are two dimensionless ratios describing the interplay between oscillations and decays:

(233) $$x_D = \frac{\Delta M_D}{\bar{\Gamma}_D}\ ,\quad y_D = \frac{\Delta\Gamma_D}{2\bar{\Gamma}_D}\ ,\quad \text{with}\bar{\Gamma}_D \equiv \frac{1}{2}(\Gamma_1 + \Gamma_2)$$

In the limit of CP invariance $\frac{q}{p} = 1$ and the mass eigenstates are CP eigenstates as well; we can then ask whether the heavier state is CP odd (as for kaons) or even. With the *definitions* $CP|D^0\rangle = |\bar{D}^0\rangle$ and $CP|D_\pm\rangle = \pm|D_\pm\rangle$ we have

(234) $$|D_\pm\rangle = \frac{1}{\sqrt{2}}\left(|D^0\rangle \pm |\bar{D}^0\rangle\right)$$

(235) $$M_{odd} - M_{even} = M_H - M_L = -2\mathrm{Re}\left[\frac{q}{p}\left(M_{12} - \frac{i}{2}\Gamma_{12}\right)\right] = -2M_{12}$$

It is instructive to follow how expressions change, yet observables remain the same, when different conventions for q/p and the CP operator are adopted [107].

10·2. *Phenomenology.* – $D^0 - \bar{D}^0$ oscillations can reveal themselves through four classes of observables:

1. 'wrong-sign' decays, i.e. a 'global' or 'time integrated' violation of a selection rule in the final state of D decays;
2. D^0 and $\bar{D}^0$ after being produced exclusively in $e^+e^- \to D^0\bar{D}^0$ decaying into two seemingly identical final states;
3. different lifetimes in different channels;
4. non-exponential decay rate evolutions.

We now discuss the four cases.

1. The final states of D decays are characterized by certain selection rules like $\Delta C = \Delta S = -\Delta Q_l$ for semileptonic and $\Delta C = \Delta S$ for Cabibbo allowed nonleptonic transitions. Having established 'correct sign' modes ([27]), which satisfy these selection rules, one can then search for corresponding 'wrong sign' channels violating these rules and compare the two rates

$$r^D_{WS}(f) = \frac{\Gamma(D^0 \to f_{WS})}{\Gamma(D^0 \to f_{CS})} \tag{236}$$

A related quantity of phenomenological convenience is

$$\chi^D_{WS}(f) \equiv \frac{\Gamma(D^0 \to f_{WS})}{\Gamma(D^0 \to f_{CS}) + \Gamma(D^0 \to f_{WS})} \tag{237}$$

Within the SM a clean signal for $D^0 - \bar{D}^0$ oscillations is the emergence of 'wrong-sign' leptons:

$$r_D \equiv \frac{\Gamma(D^0 \to l^- X)}{\Gamma(D^0 \to l^+ X)} = \frac{|q/p|^2 (x_D^2 + y_D^2)}{2 + x_D^2 - y_D^2} \simeq \frac{1}{2}(x_D^2 + y_D^2) \; , \tag{238}$$

where we have used the usual short-hand notation $r_D = r^D_{WS}(l^\mp)$; we have also anticipated that x_D^2, $y_D^2 \ll 1$ and $|q|^2 \simeq |p|^2$ hold. A signal is produced by either $x_D \neq 0$ or $y_D \neq 0$ or both. One should note that while $r_D \neq 0$ is an unambiguous sign of New Physics within the SM, it does not require very specific *time* features of $D^0 - \bar{D}^0$ oscillations *per se*, as illustrated by the next example with 'wrong-sign' kaons. For there is a SM background – doubly Cabibbo suppressed decays (DCSD). Thus there are three classes of contributions, namely purely from DCSD, from $D^0 \bar{D}^0$ oscillations followed by a Cabibbo allowed transition and from the interference between the two:

$$\tilde{r}^D_{WS}(K^\pm \pi^\mp) \equiv \frac{\Gamma(D^0 \to K^+ \pi^-)}{\Gamma(D^0 \to K^- \pi^+)} = \frac{|T(D^0 \to K^+\pi^-)|^2}{|T(D^0 \to K^-\pi^+)|^2} \cdot$$

$$\left[1 + \frac{1}{2}\frac{x_D^2 + y_D^2}{\mathrm{tg}^4\theta_C} |\hat{\rho}(K\pi)|^2 + \frac{y_D}{\mathrm{tg}^2\theta_C} \mathrm{Re}\left(\frac{q}{p}\hat{\rho}_{K\pi}\right) + \frac{x_D}{\mathrm{tg}\theta_C^2} \mathrm{Im}\left(\frac{q}{p}\hat{\rho}_{K\pi}\right) \right] \; , \tag{239}$$

where

$$\frac{1}{-\mathrm{tg}^2\theta_C}\hat{\rho}_{K\pi} \equiv \frac{T(\bar{D}^0 \to K^+\pi^-)}{T(D^0 \to K^+\pi^-)} \tag{240}$$

denotes the ratio of *instantaneous* transition amplitudes with $\hat{\rho}_{K\pi} \sim \mathcal{O}(1)$. We have retained terms of first and second order in the small parameters x_D, y_D that are enhanced by $1/\mathrm{tg}^2\theta_C$ and $1/\mathrm{tg}^4\theta_C$, respectively. $|T(D^0 \to K^+\pi^-)|^2 / |T(D^0 \to K^-\pi^+)|^2$ is

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([27]) We prefer the term 'correct sign' over 'right sign', since we do not share the often unreflected preference of the right over the left.

controled by $\mathrm{tg}^4\theta_C \sim 3 \cdot 10^{-3}$. The quark decay process actually enhances the DCS mode by $(f_K/f_\pi)^2 \sim 1.5$, and thus one expects $|T(D^0 \to K^+\pi^-)|^2/|T(D^0 \to K^-\pi^+)|^2$ to be in the range of $\mathbf{few} \cdot 10^{-3}$; yet one has to allow for some significant uncertainty in this prediction. Therefore one can infer $D^0 - \bar{D}^0$ oscillations from the time integrated ratio $\tilde{r}^D_{WS}(K^\pm\pi^\mp)$ only if it is relatively sizable.

2. Producing charm in e^+e^- annihilation on the vector meson resonance $\psi''(3770)$ leads to an exclusive particle-antiparticle pair. Consider the case where both charm mesons decay into seemingly identical final states:

$$e^+e^- \to \psi''(3770) \to D^0\bar{D}^0 \to f_D f_D \tag{241}$$

It might appear that even without $D^0 - \bar{D}^0$ oscillations such final states are possible. E.g., $f_D = K^-\pi^+$ could arise due to the Cabibbo allowed [doubly Cabibbo suppressed] mode $D^0[\bar{D}^0] \to K^-\pi^+$; likewise for the CP conjugate channel $f_D = K^+\pi^-$. Or $f_D = K^+K^-$ would be driven by the Cabibbo suppressed channels $D^0[\bar{D}^0] \to K^+K^-$. However the two charm mesons have to form a P-wave and C odd configuration; Bose-Einstein statistics then does not allow them to decay into identical final states. This is one example of exploiting EPR correlations named after Einstein, Podolsky and Rosen, who first pointed them out as a consequence of quantum mechanics' intrinsically *non*-local features [412].

The situation becomes more complex in the presence of $D^0 - \bar{D}^0$ oscillations. Bose-Einstein statistics still forbid the original $D^0\bar{D}^0$ pair to evolve into a D^0D^0 or $\bar{D}^0\bar{D}^0$ configuration at one time t. The EPR correlation tells us that if one neutral D meson reveals itself as a, say, D^0, the other one has to be a $\bar{D}^0$ at *that* time. Yet at *later* times it can evolve into a D^0, since the coherence between the original D^0 and $\bar{D}^0$ has been lost. Let us consider specifically $f_D = K^-\pi^+$. We then find

$$\frac{\sigma(e^+e^- \to D^0\bar{D}^0 \to (K^\mp\pi^\pm)_D(K^\mp\pi^\pm)_D)}{\sigma(e^+e^- \to D^0\bar{D}^0 \to (K^\mp\pi^\pm)_D(K^\pm\pi^\mp)_D)} = r_D \; ; \tag{242}$$

i.e., this process can occur only through oscillation – $r_D \neq 0$ – and the EPR correlation forces the pair of charm mesons to act like a single meson.

The requirements of Bose-Einstein statistics are satisfied in a subtle way, as best seen when describing the reaction in terms of the *mass* eigenstates D_H and D_L:

$$e^+e^- \to \psi''(3770) \to D_H D_L \to (K^-\pi^+)_{D_H}(K^-\pi^+)_{D_L} \; ; \tag{243}$$

i.e., the two decay final states $K^-\pi^+$ are *not* truly identical – their energies differ by ΔM_D. This difference is of course much too tiny to be measurable directly.

In principle the same kind of argument can be applied to more complex final states like $f_D = K^\mp\rho^\pm$, however it is less conclusive there, in particular when the ρ is identified merely by the dipion mass.

The final state $f_D = K^+K^-$ can occur in the presence of $D^0 - \bar{D}^0$ oscillations, but only if CP invriance is violated. For the CP parity of the initial $J^{PC} = 1^{--}$ state is even, yet odd for the final state $[(K^+K^-)_D(K^+K^-)_D]_{l=1}$. We will return to this point later on.

The fact that the $D^0\bar{D}^0$ pair has to form a coherent $C = -1$ quantum state in $e^+e^- \to \psi''(3770) \to D^0\bar{D}^0$ has other subtle consequences as well. Since the EPR effect (anti)correlates the time evolutions of the two neutral D mesons as sketched above, they act like a single charm meson as far as *like*-sign dileptons are concerned, i.e.,

$$\frac{\sigma(e^+e^- \to D^0\bar{D}^0 \to l^\pm l^\pm X)}{\sigma(e^+e^- \to D^0\bar{D}^0 \to l\,l\,X)} = \chi^D_{WS}(ll) \tag{244}$$

rather than the expression found for an *in*coherent pair of $D^0\bar{D}^0$

$$\frac{\sigma(D^0\bar{D}^0|_{incoh} \to l^\pm l^\pm X)}{\sigma(D^0\bar{D}^0|_{incoh} \to l\,l\,X)} = 2\chi^D_{WS}(ll)[1-\chi^D_{WS}(ll)] \tag{245}$$

3. Lifetime measurements in different D^0 channels represent a direct probe for $y_D \neq 0$. In the limit of CP invariance – which holds at least approximately in D decays, see below – mass eigenstates are CP eigenstates as well. There are three classes of final states: CP even and odd states – $D^0 \to \pi\pi,\ K\bar{K}....$ and $D^0 \to K_S\eta,\ K_S\phi$(28)$,...$, respectively – and mixed ones – $D^0 \to l^+\nu K^-,\ K^-\pi^+,...$ with $\tau_{D^0\to l^+X} \simeq \left(\tau_{D_+} + \tau_{D_-}\right)/2$.

4. The most direct signal for oscillations is the observation that decay rate evolutions in time do not follow a strictly exponential law. For semileptonic transitions we have

$$\mathrm{rate}(D^0(t) \to l^\pm \nu X) \propto e^{-\Gamma_H t} + e^{-\Gamma_L t} \pm 2e^{\bar{\Gamma}t}\cos\Delta M_D t \tag{246}$$

The most dramatic demonstration is provided by $D^0 \to K^+\pi^-$ which without oscillations is doubly Cabibbo suppressed:

$$\frac{\mathrm{rate}(D^0(t) \to K^+\pi^-)}{\mathrm{rate}(D^0(t) \to K^-\pi^+)} = \frac{|T(D^0 \to K^+\pi^-)|^2}{|T(D^0 \to K^-\pi^+)|^2}\cdot$$

$$\cdot\left[1 + \frac{(\Delta m_D t)^2 + \frac{1}{4}(\Delta\Gamma_D t)^2}{4\mathrm{tg}^4\theta_C}\left|\frac{q}{p}\right|^2|\hat{\rho}_{K\pi}|^2 + \frac{\Delta\Gamma_D t}{2\mathrm{tg}^2\theta_C}\mathrm{Re}\left(\frac{q}{p}\hat{\rho}_{K\pi}\right) + \right.$$

$$\left. + \frac{\Delta m_D t}{\mathrm{tg}^2\theta_C}\mathrm{Im}\left(\frac{q}{p}\hat{\rho}_{K\pi}\right)\right] \tag{247}$$

where we have used the notation of Eq.(239). We see that the dependence on the (proper) time of decay t differentiates between DCSD, $D^0\bar{D}^0$ oscillations and their interference.

10·3. *Theory expectations.* – Within the SM two structural reasons combine to make x_D and y_D small in contrast to the situation for $B^0 - \bar{B}^0$ and $K^0 - \bar{K}^0$ oscillations:

- The amplitude for $D^0 \leftrightarrow \bar{D}^0$ transitions is twice Cabibbo suppressed and therefore x_D, $y_D \propto \sin^2\theta_C$. The amplitudes for $K^0 \leftrightarrow \bar{K}^0$ and $B^0 \leftrightarrow \bar{B}^0$ are also twice Cabibbo and KM suppressed – yet so are their decay widths.

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(28)One has to keep in mind that $D^0 \to K_S[K^+K^-]_\phi$ has to be distinguished against $D^0 \to K_S[K^+K^-]_{f_0}$ since the latter in contrast to the former is CP even.

- Due to the GIM mechanism one has $\Delta M = 0 = \Delta\Gamma$ in the limit of flavour symmetry. Yet $K^0 \leftrightarrow \bar{K}^0$ is driven by $SU(4)_{Fl}$ breaking characterised by $m_c^2 \neq m_u^2$, which represents no suppression on the usual hadronic scales. In contrast $D^0 \leftrightarrow \bar{D}^0$ is controled by $SU(3)_{fl}$ breaking typified by $m_s^2 \neq m_d^2$ (or in terms of hadrons $M_K^2 \neq M_\pi^2$), which on the scale M_D^2 provides a very effective suppression.

These general considerations can be illustrated by considering transitions to two pseudoscalar mesons, which are common to D^0 and $\bar{D}^0$ decays and can thus communicate between them:

$$D^0 \stackrel{CS}{\Rightarrow} K^+K^-,\ \pi^+\pi^- \stackrel{CS}{\Rightarrow} \bar{D}^0\ , \tag{248}$$

$$D^0 \stackrel{CA}{\Rightarrow} K^-\pi^+ \stackrel{DCS}{\Rightarrow} \bar{D}^0 \ \text{or}\ D^0 \stackrel{DCS}{\Rightarrow} K^+\pi^- \stackrel{CA}{\Rightarrow} \bar{D}^0\ . \tag{249}$$

where CA, CS and DCS denotes the channel as Cabibbo allowed, Cabibbo suppressed and doubly Cabibbo suppressed, respectively. We have used the symbol "$\Rightarrow$" to indicate that these transitions can be real on-shell ones – for $\Delta\Gamma_D$ – as well as virtual off-shell ones – for Δm_D. Since

$$\begin{aligned} &T(D^0 \Rightarrow K^-\pi^+/K^+\pi^- \Rightarrow \bar{D}^0) \propto -\sin^2\theta_C\cos^2\theta_C \\ &T(D^0 \Rightarrow K^-K^+/\pi^+\pi^- \Rightarrow \bar{D}^0) \propto \sin^2\theta_C\cos^2\theta_C \end{aligned} \tag{250}$$

one obviously has in the $SU(3)$ limit $\Delta\Gamma(D^0 \to K\bar{K}, \pi\pi, K\pi, \pi\bar{K}) = 0$; for the amplitudes for Eqs. (249) would then be equal in size and opposite in sign to those of Eq. (248). Yet the measured branching ratios [131]

$$\begin{aligned} &\mathrm{BR}(D^0 \to K^+K^-) = (4.12 \pm 0.14)\cdot 10^{-3},\ \mathrm{BR}(D^0 \to \pi^+\pi^-) = (1.43 \pm 0.07)\cdot 10^{-3} \\ &\mathrm{BR}(D^0 \to K^-\pi^+) = (3.80 \pm 0.09)\cdot 10^{-2},\ \mathrm{BR}(D^0 \to K^+\pi^-) = (1.48 \pm 0.21)\cdot 10^{-4} \end{aligned}$$

show very considerable $SU(3)$ breakings:

$$\frac{\mathrm{BR}(D^0 \to K^+K^-)}{\mathrm{BR}(D^0 \to \pi^+\pi^-)} \simeq 2.88 \pm 0.18 \tag{251}$$

$$\frac{\mathrm{BR}(D^0 \to K^+\pi^-)}{\mathrm{BR}(D^0 \to K^-\pi^+)} \simeq (1.5 \pm 0.2)\cdot \tan^4\theta_C \tag{252}$$

compared to ratios of unity and $\tan^4\theta_C$, respectively, in the symmetry limit.

One would then conclude that the $K\bar{K}, \pi\pi, K\pi, \pi\bar{K}$ contributions to $\Delta\Gamma$ should be merely Cabibbo suppressed with flavor $SU(3)$ providing only moderate further reduction – similar to the general expectation of Eq. (257):

$$\left.\frac{\Delta\Gamma}{\Gamma}\right|_{D\to K\bar{K},\pi\pi,K\pi,\pi\bar{K}} \sim \mathcal{O}(0.01)\ . \tag{253}$$

Yet despite these large $SU(3)$ breakings an almost complete cancellation takes place between their contributions to $D^0-\bar{D}^0$ oscillations ([29]):

$$\mathrm{BR}(D^0 \to K^+K^-)+\mathrm{BR}(D^0 \to \pi^+\pi^-)-2\sqrt{\mathrm{BR}(D^0 \to K^-\pi^+)\mathrm{BR}(D^0 \to K^+\pi^-)} \simeq$$

$$\left(-8^{+12}_{-10}\right) \cdot 10^{-4} \tag{254}$$

to be compared to

$$\mathrm{BR}(D^0 \to K^-\pi^+)+\mathrm{BR}(D^0 \to K^+K^-)+\mathrm{BR}(D^0 \to \pi^+\pi^-)+\mathrm{BR}(D^0 \to K^+\pi^-) \simeq$$

$$(4.46 \pm 0.01) \cdot 10^{-2} \tag{255}$$

Having two Cabibbo suppressed classes of decays one concludes for the overall oscillation strength:

$$\frac{\Delta M_D}{\bar{\Gamma}_D},\ \Delta\Gamma_D \sim\ SU(3)\ \mathrm{breaking} \times 2\sin^2\theta_C \tag{256}$$

The proper description of $SU(3)$ *breaking* thus becomes the central issue – see our discussion in Sect.**9**˙4. The lesson to be drawn from the example given above is that one *cannot count* on the GIM mechanism to reduce the $D^0 \Rightarrow \bar{D}^0$ transition by more than a factor of three (in particular for $\Delta\Gamma_D$), yet cannot rule it out either. Thus

$$\frac{\Delta M_D}{\Gamma_D} \lesssim \frac{\Delta\Gamma_D}{\Gamma_D} \lesssim \frac{1}{3} \times 2\sin^2\theta_C \ \sim\ \mathrm{few} \times 0.01 \tag{257}$$

represents a conservative bound – for ΔM_D maybe overly conservative – based on general features of the SM.

The vastness of the literature on $D^0 - \bar{D}^0$ oscillations makes it difficult to track where a certain idea originated. There can be little doubt that many people knew about the statement that the $D^0 \leftrightarrow \bar{D}^0$ amplitude is at least $\mathcal{O}(m_s^2)$, i.e. of second order in $SU(3)$ breaking, and even mentioned it in their papers. Without any claim to originality it is given in a few lines following Fig. 2 in Ref. [413] together with a concise proof. We want to repeat this elementary reasoning since it will elucidate subsequent points.

The m_s dependence of the $D^0 \to \bar{D}^0$ transition amplitude can be inferred by considering U spin, which relates s and d quarks. For its analysis it is advantageous to work in the basis of *interaction* eigenstates

$$s' = s\cos\theta_c - d\sin\theta_c, \qquad d' = d\cos\theta_c + s\sin\theta_c\ , \tag{258}$$

([29])Since Eq. (254) is meant only as a qualitative illustration of our general argument we have ignored that $SU(3)$ *breaking* final state interactions can generate a strong phase shift $\delta_{K\pi}$ between $D^0 \to K^-\pi^+$ and $D^0 \to K^+\pi^-$, which would induce a factor $\cos\delta_{K\pi}$ in the last interference term.

Their corresponding quantum numbers 'current strangeness' S' and 'current downness" D' are conserved by the strong and electromagnetic forces. D^0 and $\bar{D}^0$ carry $C=+1$ and -1, respectively, and both have $S'=D'=0$. Since the relevant transition operator $(\bar{u}c)(\bar{s}'d')$ has the *exact* selection rules $\Delta C = -\Delta S' = \Delta D' = 1$, any amplitude $D^0 \to \bar{D}^0$ has $\Delta S' = 2$. For $m_s = m_d$ QCD dynamics strictly conserves S'. The only term violating the conservation of S' or D' is the mass term $\delta\mathcal{H} = \sin\theta_c cos\theta_c (m_s - m_d)\bar{d}'s' + \mathrm{h.c.}$ and *any* $D^0 \to \bar{D}^0$ amplitude involves a second iteration of $\delta\mathcal{H}$; q.e.d. One should note that this reasoning is based on U spin considerations alone rather than full flavour $SU(3)$: it holds irrespective of the mass of the u quark, whether it is light or heavy, degenerate or not with d or s.

Nothwithstanding this observation, there can be contributions to $T(D^0 \to \bar{D}^0)$ that are of first order in m_s [413] – a point claimed by the authors of Ref. [414] to be wrong. The main point to note is that conventional perturbation theory breaks down when transitions can occur between *degenerate* states; for the energy denominators become singular then. Such degeneracies arise here due to the presence of the pseudogoldstone bosons (PGB) K^0, π^0 and η. A contribution $\sim \mathcal{O}(m_s)$ emerges due to an IR singularity in the PGB loop. In the SM with purely left-handed charged currents such effects cancel out, yet they are present in a more general theory. ([30]) Numerical estimates have usually been obtained as follows: (**i**) Quark-level contributions are estimated by the usual quark box diagrams; they yield only insignificant contributions to ΔM_D and $\Delta\Gamma_D$ (see below). (**ii**) Various schemes employing contributions of selected hadronic states are invoked to estimate the impact of long distance dynamics; the numbers typically resulting are $x_D\,,\ y_D \sim 10^{-4} - 10^{-3}$ [277, 373]. (**iii**) These findings lead to the following widely embraced conclusions: An observation of $x_D > 10^{-3}$ would reveal the intervention of New Physics beyond the SM, while $y_D \simeq y_D|_{SM} \leq 10^{-3}$ has to hold since New Physics has hardly a chance to enhance it.

Both ΔM_D and $\Delta\Gamma_D$ have to vanish in the $SU(3)$ limit, yet the dynamics underlying them have different features: ΔM_D receives contributions from *virtual* intermediate states whereas $\Delta\Gamma$ is generated by *on*-shell transitions. Therefore the former represents a more robust quantity than the latter; actually it has often been argued that quark diagrams *cannot* be relied upon to even estimate $\Delta\Gamma$. Yet despite these differences there is no *fundamental* distinction in the theoretical treatment of ΔM_D and $\Delta\Gamma_D$: *both can be described through an OPE in terms of the expectation values of local operators and condensates incorporating short distance as well as long distance dynamics.* Only the numerical aspects differ between ΔM_D and $\Delta\Gamma_D$, as does their sensitivity to New Physics. Finally the evaluation relies on local quark-hadron duality for both ΔM_D and $\Delta\Gamma_D$; the latter is, however, more vulnerable to limitations to duality since it involves less averaging [413].

The formally leading term in the OPE for $\Delta C = -2$ transitions comes from dimension-6 four-fermion operators of the generic form $(\bar{u}c)(\bar{u}c)$ with the corresponding Wilson coefficient receiving contributions from different sources; it coresponds to the quark box diagram.

([30])In a footnote in Ref. [414] it is claimed that such an effect cannot arise "because the π, K, and η are coupled derivatively".

(**a**) Effects due to intermediate b quarks are evaluated in a straightforward way since they are far off-shell:

$$\Delta M_D^{(b\bar{b})} \simeq -\frac{G_F^2 m_b^2}{8\pi^2} \left|V_{cb}^* V_{ub}\right|^2 \frac{\langle D^0|(\bar{u}\gamma_\mu(1-\gamma_5)c)(\bar{u}\gamma_\mu(1-\gamma_5)c)|\bar{D}^0\rangle}{2M_D} ; \tag{259}$$

however they are highly suppressed by the tiny CKM parameters. Using factorization to estimate the matrix element one finds:

$$x_D^{(b\bar{b})} \sim \text{ few} \times 10^{-7} . \tag{260}$$

Loops with one b and one light quark likewise are suppressed.

(**b**) For the light intermediate quarks – d, s – the momentum scale is set by the *external* mass m_c. However, it is highly GIM suppressed [31]

$$\Delta M_D^{(s,d)} \simeq -\frac{G_F^2 m_c^2}{8\pi^2} \left|V_{cs}^* V_{us}\right|^2 \frac{\left(m_s^2 - m_d^2\right)^2}{m_c^4} \times$$

$$\frac{\langle D^0|(\bar{u}\gamma_\mu(1-\gamma_5)c)(\bar{u}\gamma_\mu(1-\gamma_5)c) + (\bar{u}(1+\gamma_5)c)(\bar{u}(1+\gamma_5)c)|\bar{D}^0\rangle}{2M_D} . \tag{261}$$

The contribution to $\Delta\Gamma_D$ from the bare quark box is greatly suppressed by a factor m_s^6. The GIM mass insertions yield a factor m_s^4. Contrary to the claim in Ref. [414] the additional factor of m_s^2 is *not* due to helicity suppression – the GIM factors already take care of that effect; it is of an accidental nature: it arises because the weak currents are purely $V - A$ and only in four dimensions. Including radiative QCD corrections to the box diagram yields contributions $\propto m_s^4 \alpha_S/\pi$. Numerically one finds:

$$\Delta\Gamma_D^{\text{box}} < \Delta M_D^{\text{box}} \sim \text{few } \times 10^{-17} \text{ GeV } \hat{=} \; x_D^{\text{box}} \sim \text{few } \times 10^{-5} \tag{262}$$

With the leading Wilson coefficient so highly suppressed, one has to consider also formally non-leading contributions from higher dimensional operators. It turns out that the $SU(3)$ GIM suppression is in general not as severe as $(m_s^2 - m_d^2)/m_c^2$ per fermion line: it can be merely m_s/μ_{hadr} if the fermion line is soft [413]. In the so-called practical version of the OPE this is described by condensates contributing to higher orders in $1/m_c$. To be more specific: having a condensate induces a suppression factor $\sim \mu_{\text{hadr}}^3/m_c^3$; yet the GIM suppression now becomes only m_s/μ_{hadr} yielding altogether a factor of order $\mu_{\text{hadr}}^2/(m_s m_c)$ which can actually result in an enhancement, since μ_{hadr}/m_c is not much smaller than unity.

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[31] This contribution is obviously saturated at the momentum scale $\sim m_c$, and thus refers to the Wilson coefficient of the $D = 6$ operator. Therefore they are not long-distance contributions despite being proportional to m_s^2.

Analyzing the contributions coming from higher-dimensional operators with the help of condensates one estimates [413] ([32])

$$x_D\ ,\ y_D\ \sim \mathcal{O}(10^{-3}) \tag{263}$$

with the realization that these estimates involve high powers of the ratio of comparable scales implying considerable numerical uncertainties that are very hard to overcome almost as a matter of principle.

Yet despite the similarities in numbers for x_D and y_D the dynamics driving these two $\Delta C = 2$ observables are quite different:

- Δm_D being generated by contributions from virtual states is sensitive to New Physics which could raise it to the percent level. At the same time it necessarily involves an integral over energies thus making it rather robust against violations of local duality.

- $\Delta\Gamma_D$ being driven by on-shell transitions can hardly be sensitive to New Physics. At the same time, however, it is very vulnerable to violations of local duality: a nearby narrow resonance could easily wreck any GIM cancellation and raise the value of $\Delta\Gamma_D$ by an order of magnitude!

The authors of Ref. [414] claim to have shown in a model-independent way that the SM indeed generates x_D, $y_D \sim 1\%$. Yet even a simple minded model – they use basically a phase space ansatz – is still a model. Their analysis can be viewed as illustrating that x_D and y_D indeed could reach the 1% level – but certainly no proof.

If data revealed $y_D \ll 1\% \leq x_D$, we would have an intriguing case for the presence of New Physics. Yet considering the theoretical uncertainties basing the case for New Physics solely on the observation of $D^0 - \bar{D}^0$ oscillations cannot be viewed as conservative.

If data revealed $y_D \ll x_D \sim 1\%$ we would have a strong case to infer the intervention of New Physics. If on the other hand $y_D \sim 1\%$ – as hinted at by the FOCUS data – then two scenarios could arise: if $x_D \leq \mathbf{few} \times 10^{-3}$ were found, one would infer that the $1/m_c$ expansion within the SM yields a correct semiquantitative result while blaming the "large" value for y_D on a sizeable and not totally surprising violation of duality. If, however, $x_D \sim 0.01$ would emerge, we would face a theoretical conundrum: an interpretation ascribing this to New Physics would hardly be convincing since $x_D \sim y_D$.

10·4. *Experiments and data.* – Traditionally, $D^0\bar{D}^0$ oscillations have been searched for by means of event-counting techniques, i.e., measurements of the wrong vs. right sign events for a given final state in D decays, which are then related to the parameter $r^D_{WS}(f)$ (or $\chi^D_{WS}(f)$) defined in Eq.(238). The first search was undertaken in deep inelastic neutrino nucleon scattering. *Opposite*-sign dimuon events reflect charm production: $\nu N \to \mu^- D^0 X \to \mu^-\mu^+ X'$. *Like*-sign dimuon events then are a signature of charm production accompanied by $D^0\bar{D}^0$ oscillations: $\nu N \to \mu^- D^0 X \Rightarrow \mu^- \bar{D}^0 X \to \mu^-\mu^- X'$. Background due to associated charm production is suppressed and furthermore leads to a very different spectrum for the secondary muon: $\nu N \to \mu^- D\bar{D} X \to$

([32])The huge lifetime ratio $\tau(K_L)/\tau(K_S) \sim 600$ is due to the accidental fact that the kaon mass is barely above the three pion threshold.

$\mu^-\mu^- D X'$. Tantalizing evidence for $D^0\bar{D}^0$ oscillations was seen by MARK III [415], where one event consistent with

$$e^+e^- \to \psi''(3770) \to (K^+\pi^-)_D(K^+\pi^-\pi^0)_D \tag{264}$$

and two events consistent with

$$e^+e^- \to \psi''(3770) \to (K^+\pi^-\pi^0)_D(K^+\pi^-\pi^0)_D \tag{265}$$

emerged from a 162 events sample. Background due to *doubly*-misidentified decays was estimated in 0.4 ± 0.2 events. The event in Eq.(264) could be attributed to DCS decay $D^0 \to K^+\pi^-$. Due to the fact that an analysis of $K(n\pi)$ resonant substructure showed that the nonresonant component was very small, i.e., all threebody D decays effectively lead to quasi-twobody pseudoscalar-vector final states, the $K^+\pi^-\pi^0$ combinations could reasonably be assumed as coming from $K^{*0}\pi^0$ and, as such, could not be attributed to DCSD. Instead, they were considered as candidates for $D^0\bar{D}^0$ mixing, although with a vanishing statistical significance: taking into account the estimated background, removing the DCS-compatible decay and accounting for fluctuations of the nonzero background in the invariant mass spectrum, one was left with a little more than a single signal event, corresponding to an alleged mixing rate of $1.6\,\%$. Alas – oscillations could not be established with one event.

Immediately after the MARK III claim, on the other side of the ocean, ARGUS at DESY [416] was able to set a $< 1.4\%$(90% cl) limit on charm mixing or DCS decay rates, based on a sample of 162 correct-sign events, and zero wrong-sign events.

In these searches the important ingredient missing was the ability to vertex D^0 decays and to determine the time evolution. Without such capability, one can distinguish between DCS decays and genuine oscillations only by employing quantum correlations as described above, which requires more statistics.

Recent advances in event statistics as well as vertexing D^0 decays have allowed to look for the specific time signature of oscillations and to probe x_D and y_D separately. First measurements exploiting lifetime information came from fixed-target experiment E691 [417] and from ALEPH at LEP [418]. With further improvements in detector performances, and most importantly with the geometrical increase in reconstructed charm sample size, attempts were performed in measuring mixing via the y parameters, i.e., by measuring lifetime asymmetries of CP-conjugate eigenstates. See Tab.XXIII for a synopsis of mixing results, and [419, 420] for excellent recent reviews. Present results from B-factories have pushed the limits on mixing parameter r down to the 0.1% limit, and the sensitivity on y beyond the 1% level.

We now discuss briefly the experimental techniques for unveiling oscillation parameters and conclude with a discussion of the data presently available, and an outlook of the future.

10˙4.1. Wrong sign vs right sign counting . Searches for $r^D_{WS}(f)$, Eq.(236), benefit greatly from the 'D^* tag trick' , which allows to tag the flavour of the neutral D meson originating from a D^* decay by the accompanying 'slow' pion; the flavour at decay is tagged by a charged kaon or lepton:

$$\begin{aligned} D^{*+} \to\ & D^0\pi^+ \\ & D^0 \Rightarrow \bar{D}^0 \end{aligned}$$

$$\bar{D}^0 \to K^+\pi^-, K^+\pi^-\pi^+\pi^-, K^+\ell^-\bar{\nu}_\ell \tag{266}$$

In the case of nonleptonic decays, the interpretation of the data is complicated by the fact that the selection rule $\Delta C = \Delta S$ can be violated also by doubly Cabibbo suppressed transitions, whose relative branching ratio is given by $\tan^4\theta_C \sim 3\cdot 10^{-3}$:

$$\begin{aligned} &D^{*+} \to D^0\pi^+ \\ &\qquad D^0 \to K^+\pi^-, K^+\pi^-\pi^+\pi^- \end{aligned} \tag{267}$$

Fig.32 shows a pictorial quark level illustration of the interplay between oscillations and DCSD processes. The wrong-sign events receive contributions from DCSD, $D^0\bar{D}^0$ oscillations (followed by a Cabibbo allowed decays) and the interference between the two [421, 422] and that those can be distinguished and thus measured separately by their dependence on the (proper) time of decay, as stated above in Eq.(247):

$$\frac{\mathrm{rate}(D^0(t)\to K^+\pi^-)}{\mathrm{rate}(D^0(t)\to K^-\pi^+)} = \frac{|T(D^0\to K^+\pi^-)|^2}{|T(D^0\to K^-\pi^+)|^2}\cdot \tag{268}$$

$$[X_{K\pi} + Y_{K\pi}(t\Gamma_D) + Z_{K\pi}(t\Gamma_D)^2)] \tag{269}$$

$$\begin{aligned} X_{K\pi} &\equiv 1 \\ Y_{K\pi} &\equiv \frac{y_D}{\mathrm{tg}^2\theta_C}\mathrm{Re}\left(\frac{q}{p}\hat{\rho}_{K\pi}\right) + \frac{x_D}{\mathrm{tg}^2\theta_C}\mathrm{Im}\left(\frac{q}{p}\hat{\rho}_{K\pi}\right) \\ Z_{K\pi} &\equiv \left|\frac{q}{p}\right|^2 \frac{x_D^2+y_D^2}{4\mathrm{tg}^4\theta_C}|\hat{\rho}_{K\pi}|^2 \end{aligned} \tag{270}$$

X and Z represent the DCSD and $D^0\bar{D}^0$ terms, respectively, and Y their interference. The latter receives a nonzero contribution from $\mathrm{Im}\left(\frac{p}{q}\frac{\hat{\rho}_{K\pi}}{|\hat{\rho}_{K\pi}|}\right)$, if there is a *weak* phase, which leads to CP violation as discussed later, and/or if a *strong* phase is present due to different FSI in $D^0 \to K^+\pi^-$ and $\bar{D}^0 \to K^+\pi^-$. One has to allow for such a difference since the latter is a pure $\Delta I = 1$ transition, while the former is given by a combination of an enhanced $\Delta I = 0$ and a suppressed $\Delta I = 1$ amplitude.

Assuming CP conservation, i.e. the absence of weak phases, implies $T(D^0 \to K^+\pi^-) = T(\bar{D}^0 \to K^-\pi^+)$ and $|q/p| = 1$. As explained later the phase of q/p is actually unphysical and can be absorbed into the definition of $\bar{D}^0$ in $\frac{q}{p}\frac{T(\bar{D}^0\to K^+\pi^-)}{T(D^0\to K^+\pi^-)}$; the latter quantity can then be written as

$$\frac{q}{p}\frac{T(\bar{D}^0\to K^+\pi^-)}{T(D^0\to K^+\pi^-)} = e^{i\delta}\frac{1}{\mathrm{tg}^2\theta_C}\hat{\rho}(K\pi) \tag{271}$$

leading to the simplified expression

$$r_{WS}(t) = \frac{|T(D^0\to K^+\pi^-)|^2}{|T(D^0\to K^-\pi^+)|^2}\left[1 + \frac{(x_D^2+y_D^2)(\Gamma_D t)^2}{\frac{|T(D^0\to K^+\pi^-)|^2}{|T(D^0\to K^-\pi^+)|^2}} - \frac{y'_D(\Gamma_D t)}{\frac{|T(D^0\to K^+\pi^-)|}{|T(D^0\to K^-\pi^+)|}}\right] \tag{272}$$

where

(273) $$y'_D \equiv y_D \cos\delta - x_D \sin\delta\ ,\ x'_D \equiv x_D \cos\delta + y_D \sin\delta$$

with $x_D^2 + y_D^2 \equiv (x'_D)^2 + (y'_D)^2$. The observable ratio of wrong- to correct-sign events as a function of the (proper) time of decay t is thus expressed in terms of the branching ratio for the doubly Cabibbo suppressed mode, the oscillation parameters x_D, y_D and the strong phase δ.

Finding $Y_{K\pi}$ and/or $Z_{K\pi}$ to differ from zero unequivocally establishes the presence of oscillations. It would also allow us to extract $x_D^2 + y_D^2$ and y'_D – yet not x_D and y_D separately due to our ignorance concerning the strong phase δ. One can extract δ in a clean way, namely by exploiting the *coherence* of a $D^0\bar{D}^0$ pair produced in e^+e^- annihilation close to threshold. We already stated that

(274) $$\frac{\sigma(e^+e^- \to D^0\bar{D}^0 \to (K^\pm\pi^\mp)_D(K^\pm\pi^\mp)_D}{\sigma(e^+e^- \to D^0\bar{D}^0 \to (K^\pm\pi^\mp)_D(K^\mp\pi^\pm)_D} \simeq \frac{x_D^2 + y_D^2}{2}\ ;$$

i.e., this transition can occur only due to $D^0\bar{D}^0$ oscillations. EPR correlations produce a quite different ratio between these final states when the underlying reaction is $e^+e^- \to D^0\bar{D}^0\gamma$ (due to $D^{0*}\bar{D}^0$, $\bar{D}^{0*}D^0 \to D^0\bar{D}^0\gamma$), since now the $D^0\bar{D}^0$ pair forms a *even* configuration:

$$\frac{\sigma(e^+e^- \to D^0\bar{D}^0\gamma \to (K^\pm\pi^\mp)_D(K^\pm\pi^\mp)_D\gamma}{\sigma(e^+e^- \to D^0\bar{D}^0\gamma \to (K^\pm\pi^\mp)_D(K^\mp\pi^\pm)_D\gamma} \simeq$$

(275) $$\frac{3}{2}\left(x_D^2 + y_D^2\right) + 4\left|\frac{T(D^0 \to K^+\pi^-)}{T(D^0 \to K^-\pi^+)}\right|^2 + 8y_D\cos\delta\left|\frac{T(D^0 \to K^+\pi^-)}{T(D^0 \to K^-\pi^+)}\right|$$

Alternatively one can measure

$$\sigma(e^+e^- \to D^0\bar{D}^0 \to (K^\pm\pi^\mp)_D(K^\mp\pi^\pm)_D) \propto$$

(276) $$\left(1 - \frac{x_D^2 - y_D^2}{2}\right)\left(1 - 2\cos\delta\left|\frac{T(D^0 \to K^+\pi^-)}{T(D^0 \to K^-\pi^+)}\right|^2\right)$$

Eq.(275) or Eq.(276) coupled with the previously described analyses of single D decays allow to extract x_D, y_D and δ.

In the search for oscillationns from wrong-sign events, the event selection procedure at fixed-target experiments requires a good candidate secondary vertex consistent with the D^0 mass, a suitable primary vertex consisting of a minimum number of other tracks, and well isolated from the secondary vertex. The main background to the WS signal is due to doubly misidentified $K^+\pi^-$ pairs from D^0 decays which form a broad peak directly under the D^0 signal in $K^+\pi^-$ and a narrow peak in the $D^* - D$ mass difference signal region. The mass difference background is indistinguishable from the real WS tagged signal. To eliminate this background, the $K\pi$ invariant mass is computed with the kaon and pion particle hypotheses swapped. Any candidate whose swapped mass is within

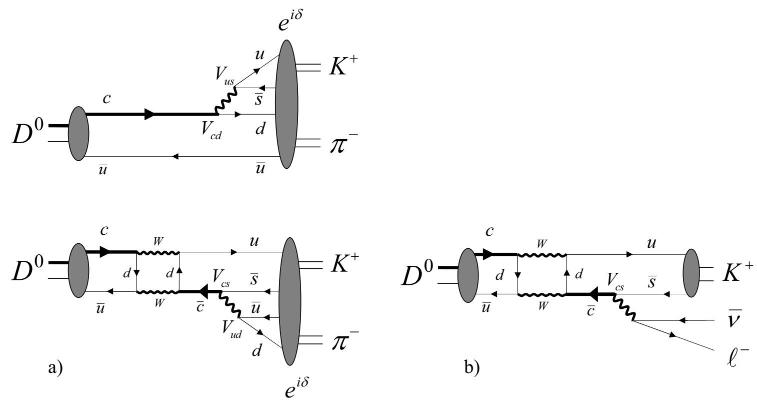

Fig. 32. – *Cartoon which explains at a quark-level (a) the two routes (DCS decays and mixing) to get to* $K^+\pi^-$ *starting from a* D^0*; (b) the one route (mixing) which is possible for semileptonic final state.*

some distance (a few σ's) of the D^0 mass is subjected to a cut on the sum of the $K\pi$ separations for both tracks. Finally, all tracks in the production vertex are tested as potential π candidates, and are accepted if within a narrow (typically, $\pm 50\,\mathrm{MeV}/c^2$) window of the nominal $D^* - D^0$ mass difference, and if they satisfy a loose particle identification cut.

Experiments at e^+e^- colliders follow similar reconstruction strategies, with the important difference being the superior resolution on the $D^* - D^0$ mass difference (typically 200 keV). A cut on the angle of the pion candidate in the D^0 rest frame with respect to the D^0 boost rejects asymmetric D^0 decays, where the pion candidate has low momentum. Another relevant difference with fixed-target experiments is the fact that the primary vertex is reconstructed over blocks of data as the centroid of the luminous e^+e^- interaction region. The vertical extent of such region is only about $10\mu m$, and this permits to reconstruct the proper time t using only the vertical component of the flight distance of the D^0 candidate.

For both FT and e^+e^- experiments, the WS mass peak is fitted in the Q-M plane, i.e., the scatter plot $M(D^* - D^0) \quad vs \quad M(D^0)$. Lifetime information is extracted on the proper time distribution of WS candidates within a few sigmas from the CFD signal value. Due to the superior lifetime resoultion, the distribution of proper times is an exponential at FT, and an exponential convoluted with a gaussian resolution function for e^+e^- experiment. The dominant contribution to the systematic error comes from uncertainties in the shapes and acceptances of backgrounds.

The alternative option in counting techniques is to use semileptonic, D^*-tagged final states $K\ell\nu$; they do not suffer from DCSD pollution, but are harder experimentally. For semileptonic final states, Eq.(268) reduces to

$$r^D_{WS}(lX) \propto \frac{r_D}{2} t^2 e^{-t} \tag{277}$$

Event selection and vertex reconstruction techniques follow guidelines common to all semileptonic studies, where the presence of an undetected neutrino prevents one from reconstructing the D^0 momentum directly. By exploiting the information on the primary and secondary vertices, the K and lepton momenta, and assuming a D^0 parent mass, one can reconstruct the neutrino momentum modulo a two-fold ambiguity. Finally the invariant mass of D^* and proper decay time of D^0 are computed, and events are selected in the Q-M plane in analogy to the hadronic case. Feedthrough of hadronic modes is mainly by $K^-\pi^+\pi^0$ with an undetected π^0 faking the neutrino, and a pion is misidentified as a lepton.

10˙4.2. Lifetime difference measurements. In the presence of $D^0 - \bar{D}^0$ oscillations a single lifetime does not suffice to describe all transitions. The situation can most concisely be discussed when CP invariance is assumed – $[H_{\Delta C \neq 0}, CP] = 0$ – since the mass eigenstates are then CP eigenstates as well (CP violation will be addressed in the next Section). Transitions $D \to K^+K^-$, $\pi^+\pi^-$ are controlled by the lifetime of the CP even state [423] and $D \to K_S\pi^0$, $K_S\eta$ by that of the CP odd state.

The channel $D^0 \to K_S\phi$ has a CP odd final state; yet in $D^0 \to K_SK^+K^-$ one has to disentangle the ϕ and f^0 contributions in K^+K^-, since the latter leads to a CP even state. This complication and its implications for $B \to K_S\phi$ will be addressed again in our discussion of CP violation. As an example of the problem, it has been shown [424], out of photoproduction data, the fraction of f_0 component in $D^0 \to K_SK^+K^-$ being as large as 37.8 ± 3.0% vith a relatively loose mass cut $M(KK)^2 < 1.1 GeV^2$, which is reduced to 8% with a narrower mass cut $(1.034 < M^2 < 1.042)\, GeV^2$, with the obvious penalty of a very large reduction in statistics. BABAR has shown [426] preliminary results where the f_0/ϕ ratio is about 25% integrated over the entire range of M^2.

The final state $D \to K^-\pi^+$ on the other hand has no definite CP parity. Up to terms of order $(\Delta\Gamma/\bar{\Gamma})^2$ its time dependance is controlled by the average lifetime $\bar{\tau} = \bar{\Gamma}^{-1}$[423] and thus

$$y_D \simeq y^D_{CP} = \frac{\Gamma_{K^+K^-} - \Gamma_{K^\mp\pi^\pm}}{\Gamma_{K^\mp\pi^\pm}} = \frac{\tau(D \to K\pi)}{\tau(D \to KK)} - 1 \tag{278}$$

Several sources of systematic errors common to KK and $K\pi$ final states cancel in the ratio of their lifetimes and thus in y^D_{CP}. The strategy common to fixed target and e^+e^- experiments is to select high-statistics, clean modes, for which $D \to K^+K^-$ and $K^-\pi^+$ are prime examples. For the latter requirement, a D^*-tag is of paramount importance. Many of the analysis cuts and the fitting strategies have close analogies with those employed in lifetime measurements (see Sect. **6**). In all measurements of y^D_{CP}, the systematic error only refers to the lifetimes asymmetry and not to the absolute lifetime of the CP eigenstates. As such, lifetimes of, i.e., $\tau(D \to K\pi)$ have a statistical error only.

Measuring y at fixed target experiments: The cuts used to obtain a clean signal are designed to produce a nearly flat efficiency in reduced proper time $t' \equiv (\ell - N\sigma)/(\gamma\beta c)$. The charm secondary is selected by means of a candidate driven algorithm, with stringent requests on particle identification, as well as requiring a minimum σ_ℓ detachment between primary and secondary vertex. To select a clean sample, either a D^* tag is required, or a set of more stringent cuts, such as more stringent Cerenkov requirements on kaons and pions, momenta of decay particles balancing each other, primary vertex inside the target material, and resolution of proper time less than 60 fs. The D^* tagged sample has a

TABLE XXIII. – *Synopsis of recent $D^0\bar{D}^0$ mixing results. CPV phase is φ, strong phase is δ, interference angle is $\phi = \arg(ix + y) - \varphi - \delta$. As instance, $\varphi \neq 0$ stands for no CP conservation is assumed.*

	Assumptions	Mode	N_{RS}	Result (%)
E691 88 [427]	$\varphi = 0, \cos\phi = 0$	$K\pi$	1.5k	$r < 0.5$
(95% CL)		$K\pi\pi\pi$		$r < 0.5$
				combined: $r < 0.37$
ALEPH[418]	No mix	$K\pi$	1.0k	$r_{DCS} = 1.84 \pm .59 \pm .34$
(95% CL)	$\varphi = 0, \cos\phi = 0$			$r < 0.92$
	$\varphi = 0, \cos\phi = +1$			$r < 0.96$
	$\varphi = 0, \cos\phi = -1$			$r < 3.6$
E791[428]	$\varphi = 0$	$K\ell\nu$	2.5k	$r = 0.11^{+0.30}_{-0.27}$
(90% CL)				$r < 0.50$
E791[429]	No mix	$K\pi$	5.6k	$r_{DCS} = 0.68^{+0.34}_{-0.33} \pm 0.07$
(90% CL)	No mix	$K3\pi$	3.5k	$r_{DCS} = 0.25^{+0.36}_{-0.34} \pm 0.03$
	$\varphi \neq 0$ in y			$r = 0.39^{+0.36}_{-0.32} \pm 0.16$
				$r < 0.85$
E791[430]	$\varphi = 0, \delta = 0$	KK	6.7k	$\Delta\Gamma = 0.04 \pm 0.14 \pm 0.05\,\mathrm{ps}^{-1}$
(90% CL)		$K\pi$	60k	$(-0.20 < \Delta\Gamma < 0.28)\,\mathrm{ps}^{-1}$
				$y = 0.8 \pm 2.9 \pm 1.0 \quad (-4 < y < 6)$
CLEO 00[431]	No mix	$K\pi$	14k	$r_{DCS} = 0.332^{+0.063}_{-0.065} \pm 0.040$
(95% CL)	$\varphi \neq 0, \delta \neq 0$			$r_{DCS} = 0.47^{+0.11}_{-0.12} \pm 0.01$
$9\,\mathrm{fb}^{-1}$				$y' = -2.3^{+1.3}_{-1.4} \pm 0.3\,(-5.2 < y' < 0.2)$
				$x' = 0 \pm 1.5 \pm 0.2\,(-2.8 < x' < 2.8)$
				$(x')^2/2 < 0.038$
FOCUS 00[432]	$\varphi = 0, \delta = 0$	$K\pi$	120k	
		KK	10k	$y = 3.42 \pm 1.39 \pm 0.74$
FOCUS 01[433]	$\varphi \neq 0$	$K\pi$	36.7k	$-12.4 < y' < -0.6$
prelim.				$\lvert x'\rvert < 3.9$
FOCUS 02[434]	$\varphi \neq 0$	$K\ell\nu$	60k	stat.err. only $r < 0.12$
prelim.				
CLEO 02[435]	$\varphi = 0, \delta = 0$	$\pi\pi$	710	$A_{CP}(KK) = 0.0 \pm 2.2 \pm 0.8$
		KK	1.9k	$A_{CP}(\pi\pi) = 1.9 \pm 3.2 \pm 0.8$
$9\,\mathrm{fb}^{-1}$		$K\pi$	20k	$y = -1.2 \pm 2.5 \pm 1.4$
BELLE 02[241]	$\varphi = 0, \delta = 0$	$K\pi$	214k	
$23.4\,\mathrm{fb}^{-1}$		KK	18k	$y = -0.5 \pm 1.0^{+0.7}_{-0.8}$
BABAR 03[436]	$\varphi = 0, \delta = 0$	$\pi\pi$	13k	$y = 0.8 \pm 0.4^{+0.5}_{-0.4}$
		KK	26k	$\Delta y = -0.8 \pm 0.6 \pm 0.2$
$91\,\mathrm{fb}^{-1}$		$K\pi$	265k	
BABAR 03[437]	No mix	$K\pi$	120k	$r_{DCS} = 0.357 \pm 0.022 \pm 0.027$
(95% CL)				$A_{CP}(K^+\pi^-) = 9.5 \pm 6.1 \pm 8.3$
$57.1\,\mathrm{fb}^{-1}$	Nomix, $\varphi \neq 0, \delta \neq 0$			$r_{DCS} = 0.359 \pm 0.020 \pm 0.027$
	$\varphi \neq 0, \delta \neq 0$			$-5.6 < y' < 3.9$
				$x'^2 < 0.22, \quad r < 0.16$
	$\varphi = 0, \delta \neq 0$			$-2.7 < y' < 2.2$
				$x'^2 < 0.2, \quad r < 0.13$
BELLE 03[248]	$\varphi = 0, \delta = 0$	KK	36.5k	$y = 1.15 \pm 0.69 \pm 0.38$
$158\,\mathrm{fb}^{-1}$		$K\pi$	448k k	

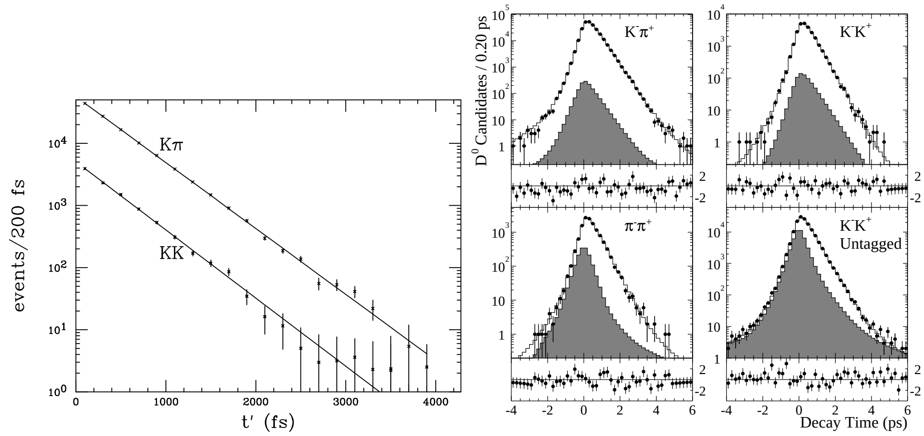

Fig. 33. – Lifetime distributions used to extract lifetime asymmetry y_{CP} (left) Fixed-target experiment FOCUS [432] (right) e^+e^- experiment BABAR [436].

better signal-to-noise ratio, while the inclusive sample accommodates larger sample size. Both samples are treated as systematics checks.

The $D^0 \to K^-K^+$ sample is characterized by a prominent reflection background coming from misidentified $D^0 \to K^-\pi^+$ decays (Fig. 33). The amount of $D^0 \to K^-\pi^+$ reflection is obtained by a mass fit to the K^-K^+ sample and the shape of the reflection is deduced from a high-statistics montecarlo sample. It is assumed that the time evolution of the reflection is described by the lifetime of $D^0 \to K^-\pi^+$ and a fit is performed of the reduced proper time distributions of the $D^0 \to K^-\pi^+$ and $D^0 \to K^-K^+$ samples at the same time. The fit parameters are the $D \to K\pi$ lifetime, the lifetime asymmetry y_{CP}, and the number of background events under each $D^0 \to K^-\pi^+$ and $D^0 \to K^-K^+$ signal region. The signal contributions for the $D^0 \to K^-\pi^+$, $D^0 \to K^-K^+$ and the reflection from the misidentified $D^0 \to K^-\pi^+$ in the reduced proper time histograms are described by a term $f(t') \exp(-t'/\tau)$ in the fit likelihood function. The function $f(t')$, determined from montecarlo, covers any deviation of the reduced proper time distribution from a pure exponential due to acceptance. The background yield parameters are either left floating, or fixed to the number of events in mass sidebands using a Poisson penalty term in the fit likelihood function.

The systematic error on the lifetime asymmetry is determined by calculating the shifts in y_{CP} for a set of detachment cuts, kaon identification cuts, background normalizations, and lifetime fit ranges.

Measuring y_{CP} at e^+e^- B factories: Event selection and fitting are carried out by CLEO, BABAR and BELLE in a very similar fashion to each other and also to fixed-target experiments. A D^*-tag is normally used to select a clean sample of D^0 decays, with the noteworthy exception of BELLE which manages to get a sample of equivalent purity thanks to claimed superior particle identification.

Lifetime fits are performed via unbinned maximum likelihood method, where the fit function contains proper time signal and background terms. An additional penalty term in the likelihood function proportional to the difference to the nominal D^0 invariant

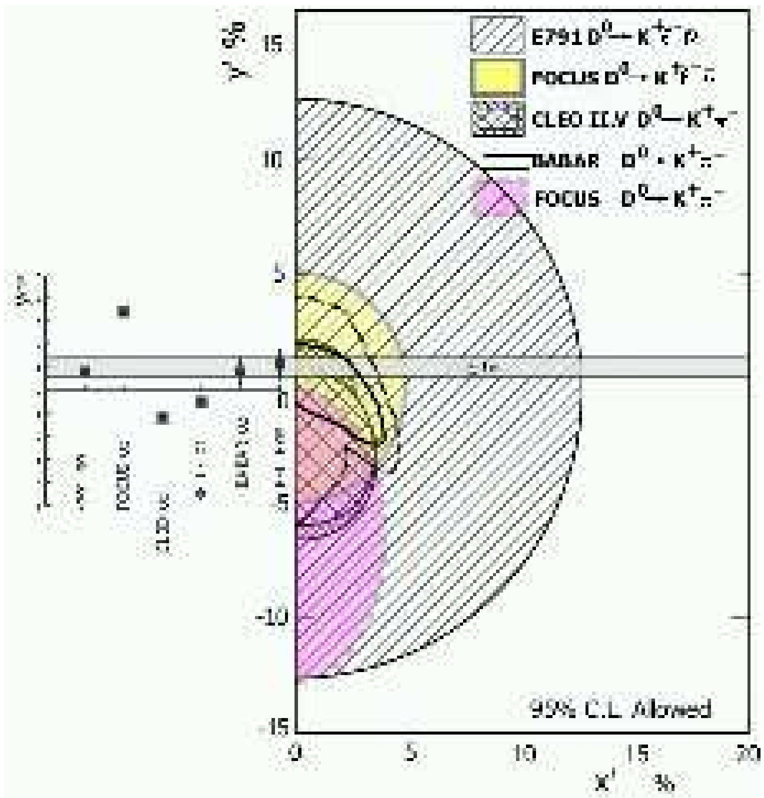

Fig. 34. – Summary of x', y', y_{CP} measurements. The BABAR limits (solid contour: CP conservation assumed; dashed contour: CP conservation not assumed) were scaled to x' from Ref.[437], and are superimposed for qualitative comparison.

mass is also used.

Main sources of systematic error on the lifetime asymmetry are the uncertainty in the MC correction functions $f(t')$, and the background contribution to the signals.

10˙5. *Where do we stand today, and what next?*. – The oscillation industry is being revamped by a steady flux of results (for recent reviews, see [419, 420]). All results so far are consistent with no oscillations. FOCUS had stirred up excitement with a $\sim 2\sigma$ evidence for positive y_{CP}. After several ups and downs (Tab.XXIII), y_{CP} is dominated by the 2003 preliminary measurements [436, 248]

$$y_{CP}^{BABAR03} = 0.8 \pm 0.6 \pm 0.2\% \qquad y_{CP}^{BELLE03} = 1.15 \pm 0.69 \pm 0.38\% \tag{279}$$

We average the results in Tab.XXIII and we get the world average

$$< y_{CP} >^{2003} = (1.0 \pm 0.5)\% \tag{280}$$

In year 2000 CLEO published a limit on WS counting rates of $r < 0.041\%$ based on $9\,fb^{-1}$ data set. BABAR's new limit is about three times higher, although it is based on a tenfold dataset (Fig.34). Besides, BABAR contour limits have same size as CLEO contour limits. Possible reasons for such inconsistent results have been discussed [420]; they are likely due to differences in fit and limit computation techniques. Particularly intriguing is the effect of positive y' fit results providing larger limits than those provided by negative y' fit results.

Although everything is consistent with zero, y_{CP} seems to have a tendency to prefer positive values, while hadronic mixing measurements sort of favour laying in the $y' < 0$ semiplane. Very important measurements will be performed in counting techniques using semileptonic channels, where a preliminary FOCUS result shows a $r < 0.12\%$ limit and results from B-factories should be coming up. A critical issue is how we can determine experimentally the strong phase shift δ. We have already described how it can be extracted from $e^+e^- \to \psi(3770) \to D^0\bar{D}^0$ as will be studied by CLEO-c. Alternatively one can infer it from $D \to K_L\pi$ decays [438].

The future will bring a rich harvest of new data. B factories plan to multiply by five their dataset by 2006. The statistical error on y_{CP} will then be negligible with respect to the systematic error, and a 0.2% precision will be reached. Similarly for hadronic mixing which should get to below the 0.5% level. At CLEO-c oscillation will be investigated under novel coherence conditions, where δ can be measured independently [439, 440].

10˙6. *Resume*. – Table XXIV summarizes our present experimental knowledge on meson-antimeson oscillations. In the second column we have indicated the relative weight *within the SM* between contributions to $\Delta M(P^0)$ from short distance (SD) and long distance (LD) dynamics. We have also translated x into $\Delta M(P^0)$ for an absolute yardstick. $K^0 - \bar{K}^0$ and $B_d - \bar{B}_d$ oscillations have been established and well measured. The observed values for $\Delta M(K)$ and $\Delta M(B_d)$ can be reproduced by SM dynamics within reasonable theoretical uncertainties.

Experimental evidences in nonleptonic D^0 decays can bear on the discussion of the experimental determination of strong phase δ in hadronic $D^0\bar{D}^0$ mixing. Isospin decomposition of $D^0 \to KK, \pi\pi$ decays is used to determine the phase shift between isospin amplitudes[366]. Phase shift is large, this being a signal for large FSI. A further attempt is made by authors of [366] to estimate the relative importance of elastic and inelastic FSI. They find that the elastic FSI cannot account for all the discrepancy found between experiment and theory predicted value of the branching ratio, and they argue that the explanation may be the presence of an inelastic FSI component acting in the transition $KK \to \pi\pi$. The inelastic FSI component is a smoking gun for a sizable strong angle δ which affects also the $K\pi$ final states

The experimental sensitivity for $D^0 - \bar{D}^0$ oscillations has been increased considerably in the last few years with no signal established. Yet for proper perspective one should note that D^0 is the only meson with up-type quarks in this list and that furthermore the upper bound on ΔM_D is not much smaller than ΔM_K and $\Delta M(B_d)$. The best theoretical estimate yields y_D, $x_D \sim \mathcal{O}(10^{-3})$; yet a value ~ 0.01 cannot be ruled out nor is there a clear prospect for reducing the theoretical uncertainties significantly. Yet rather then resigning ourselves to 'experimental nihilism' we advocate exercising sound judgment. It is important to measure x_D and y_D separately as accurately as possible. Finding, say, $y_D \sim 10^{-3}$ together with $x_D \sim 0.01$ would represent a tantalizing 'prima facie' case for the presence of New Physics; on the other hand $x_D \sim y_D \sim 0.01$ would suggest that the OPE treatment involving a $1/m_Q$ expansion can provide little quantitative guidance here; $x_D \sim 10^{-3}$ together with $y_D \sim 0.01$ would point towards a violation of quark-hadron duality in y_D.

TABLE XXIV. – *Compilation of oscillation parameters (95% cl) [131].*

P^0	SD/LD	x	$\Delta M(P^0)$	y
$K^0\,(d\bar{s})$	$SD \sim LD$	0.946 ± 0.003	$(3.5 \pm 0.003) \cdot 10^{-6}$ eV	0.9963 ± 0.0036
$D^0\,(c\bar{u})$	$SD \ll LD$	< 0.03	$\leq 4.6 \cdot 10^{-5}$ eV	-0.002 ± 0.011
$B^0_d\,(d\bar{b})$	$SD \gg LD$	0.771 ± 0.012	$(3.30 \pm 0.05) \cdot 10^{-4}$ eV	?
$B^0_s\,(s\bar{b})$	$SD \gg LD$	> 20.0	$\geq 9.3 \cdot 10^{-3}$ eV	< 0.15

The observation of $D^0 - \bar{D}^0$ oscillations thus might or might not signal the intervention of New Physics. However as explained in the next Section they can create a new stage for CP violation in D^0 decays that would unequivocally reveal New Physics. Furthermore they can have an important or even crucial impact on the CP phenomenology in B decays. One of the most promising ways to extract the angle $\phi_3 = \gamma$ of the CKM unitarity triangle is to compare $B^\pm \to D^{neut} K^\pm$ transition rates, where D^{neut} denotes a neutral D meson that either reveals its flavour through its decay or it does not. It has been pointed out [441] that even moderate values $x_D,\ y_D \sim 10^{-2}$ – which are not far from the border that SM dynamics can reach – could have a sizable impact on the value of $\phi_3 = \gamma$ thus extracted. Therefore $D^0 - \bar{D}^0$ oscillations could either hide the impact New Physics had on $B^\pm \to D^{neut} K^\pm$ – or fake such an impact!

11. – CP violation

The SM predicts quite small CP asymmetries in charm transitions. Therefore searches for CP violation there are mainly motivated as probes for New Physics, as it is with $D^0 - \bar{D}^0$ oscillations. Yet again the experimental sensitivity has reached the very few percent level, and one has to ask how small is small – 10^{-2}, 10^{-3}, 10^{-4}?

Due to CPT invariance (later we add a few comments on CPT violation) CP violation can be implemented only through a complex phase in some effective couplings. For it to become observable two different, yet coherent amplitudes have to contribute to an observable. There are two types of scenarios for implementing this requirement:

1. Two different $\Delta C = 1$ amplitudes of fixed ratio – distinguished by, say, their isospin content – exist leading *coherently* to the same final state.
2. $D^0 - \bar{D}^0$ oscillations driven by $\Delta C = 2$ dynamics provide the second amplitude, the weight of which varies with time.

11·1. *Direct CP violation.* – CP violation appearing in $\Delta C = 1$ amplitudes is called *direct* CP violation. It can occur in the decays of charged and neutral charm meson and baryons.

11·1.1. Partial widths. Consider a final state f that can be reached coherently via two different quark level transition amplitudes $\mathcal{M}_1$ and $\mathcal{M}_2$:

$$T(D \to f) = \lambda_1 \mathcal{M}_1 + \lambda_2 \mathcal{M}_2 \tag{281}$$

We have factored out the weak couplings $\lambda_{1,2}$ while allowing the amplitudes $\mathcal{T}_{1,2}$ to be still complex due to strong or electromagnetic FSI. For the CP conjugate reaction one has

$$T(\bar{D} \to \bar{f}) = \lambda_1^* \mathcal{M}_1 + \lambda_2^* \mathcal{M}_2 \tag{282}$$

It is important to note that the reduced amplitudes $\mathcal{M}_{1,2}$ remain unchanged, since strong and electromagnetic forces conserve CP. Therefore we find

$$\Gamma(\bar{D} \to \bar{f}) - \Gamma(D \to f) = \frac{2\mathrm{Im}\lambda_1\lambda_2^* \cdot \mathrm{Im}\mathcal{M}_1\mathcal{M}_2^*}{|\lambda_1|^2|\mathcal{M}_1|^2 + |\lambda_2|^2|\mathcal{M}_2|^2 + 2\mathrm{Re}\lambda_1\lambda_2^* \cdot \mathrm{Re}\mathcal{M}_1\mathcal{M}_2^*} \tag{283}$$

i.e. for a CP asymmetry to become observable, two conditions have to satisfied simultaneously irrespective of the underlying dynamics:

- Im $\lambda_1\lambda_2^* \neq 0$, i.e. there has to be a relative phase between the weak coulings $\lambda_{1,2}$.
- Im$\mathcal{M}_1\mathcal{M}_2^* \neq 0$, i.e. FSI have to induce a phase shift between $\mathcal{M}_{1,2}$.

For a larger asymmetry one would like to have also $|\lambda_1\mathcal{M}_1| \simeq |\lambda_2\mathcal{M}_2|$.

The first condition requires *within the SM* that such an effect can occur in singly Cabibbo suppressed, yet neither Cabibbo allowed nor doubly suppressed channels. There is a subtle exception to this general rule in $D^+ \to K_S\pi^+$, as explained later.

The second condition implies there have to be nontrivial FSI , which can happen in particular when the two transition amplitudes differ in their isospin content. We know, as discussed in Sect.**9**˙4 that FSI are virulent in the charm region leading to in general sizeable phase shifts. While we cannot calculate them and therefore cannot predict the size of direct CP asymmetries, even when the weak phases are know, CPT symmetry implies some relations between CP asymmetries in different channels. For CPT invariance enforces considerably more than the equality of lifetimes for particles and antiparticles; it tells us that the widths for *sub*classes of transitions for particles and antiparticles have to coincide already, either identically or at least practically. Just writing down strong phases in an equation like Eq.(281) does *not automatically* satisfy CPT constraints .

We will illustrate this feature first with two simple examples and then express it in more general terms.

- CPT invariance already implies $\Gamma(K^- \to \pi^-\pi^0) = \Gamma(K^+ \to \pi^+\pi^0)$ up to small electromagnetic corrections, since in that case there are no other channels it can rescatter with.
- While $\Gamma(K^0 \to \pi^+\pi^-) \neq \Gamma(\bar{K}^0 \to \pi^+\pi^-)$ and $\Gamma(K^0 \to \pi^0\pi^0) \neq \Gamma(\bar{K}^0 \to \pi^0\pi^0)$ one has $\Gamma(K^0 \to \pi^+\pi^- + \pi^0\pi^0) = \Gamma(\bar{K}^0 \to \pi^+\pi^- + \pi^0\pi^0)$.
- Let us now consider a scenario where a particle P and its antiparticle $\bar{P}$ can each decay into two final states only, namely a, b and $\bar{a}, \bar{b}$, respectively [442, 443]. Let us further assume that strong (and electromagnetic) forces drive transitions among a and b – and likewise for $\bar{a}$ and $\bar{b}$ – as described by an S matrix $\mathcal{S}$. The latter can then be decomposed into two parts

$$\mathcal{S} = \mathcal{S}^{diag} + \mathcal{S}^{off-diag} \,, \tag{284}$$

where $\mathcal{S}^{diag}$ contains the diagonal transitions $a \Rightarrow a$, $b \Rightarrow b$

$$\mathcal{S}^{diag}_{ss} = e^{2i\delta_s} \,, s = a, b \tag{285}$$

and $\mathcal{S}^{off-diag}$ the off-diagonal ones $a \Rightarrow b$, $b \Rightarrow a$:

$$\mathcal{S}_{ab}^{off-diag} = 2i\mathcal{T}_{ab}^{resc} e^{i(\delta_a + \delta_b)} \tag{286}$$

with

$$\mathcal{T}_{ab}^{resc} = \mathcal{T}_{ba}^{resc} = (\mathcal{T}_{ab}^{resc})^* \ , \tag{287}$$

since the strong and electromagnetic forces driving the rescattering conserve CP and T. The resulting S matrix is unitary to first order in $\mathcal{T}_{ab}^{resc}$. CPT invariance implies the following relation between the weak decay amplitude of $\bar{P}$ and P:

$$T(P \to a) = e^{i\delta_a} \left[T_a + T_b i \mathcal{T}_{ab}^{resc} \right] \tag{288}$$

$$T(\bar{P} \to \bar{a}) = e^{i\delta_a} \left[T_a^* + T_b^* i \mathcal{T}_{ab}^{resc} \right] \tag{289}$$

and thus

$$\Delta\gamma(a) \equiv |T(\bar{P} \to \bar{a})|^2 - |T(P \to a)|^2 = 4\mathcal{T}_{ab}^{resc} \mathrm{Im} T_a^* T_b \ ; \tag{290}$$

likewise

$$\Delta\gamma(b) \equiv |T(\bar{P} \to \bar{b})|^2 - |T(P \to b)|^2 = 4\mathcal{T}_{ab}^{resc} \mathrm{Im} T_b^* T_a \tag{291}$$

and therefore as expected

$$\Delta\gamma(b) = -\Delta\gamma(b) \tag{292}$$

Some further features can be read off from Eq.(290):

1. If the two channels that rescatter have comparable widths – $\Gamma(P \to a) \sim \Gamma(P \to b)$ – one would like the rescattering $b \leftrightarrow a$ to proceed via the usual strong forces; for otherwise the asymmetry $\Delta\Gamma$ is suppressed relative to these widths by the electromagnetic coupling.
2. If on the other hand the channels command very different widths – say $\Gamma(P \to a) \gg \Gamma(P \to b)$ – then a large *relative* asymmetry in $P \to b$ is accompagnied by a tiny one in $P \to a$.

This simple scenario can easily be extended to two sets A and B of final states s.t. for all states a in set A the transition amplitudes have the same weak coupling and likewise for states b in set B. One then finds

$$\Delta\gamma(a) = 4 \sum_{b \in B} \mathcal{T}_{ab}^{resc} \mathrm{Im} T_a^* T_b \tag{293}$$

The sum over all CP asymmetries for states $a \in A$ cancels the correponding sum over $b \in B$:

$$\sum_{a \in A} \Delta\gamma(a) = 4 \sum_{b \in B} \mathcal{T}_{ab}^{resc} \mathrm{Im} T_a^* T_b = - \sum_{b \in B} \Delta\gamma(b) \tag{294}$$

These considerations tell us that the CP asymmetry averaged over certain classes of channels defined by their quantum numbers has to vanish. Yet these channels can still be very heterogenous, namely consisting of two- and quasi-two-body modes, three-body channels and other multi-body decays. Hence we can conclude:

- If one finds a direct CP asymmetry in one channel, one can infer – based on rather general grounds – which other channels have to exhibit the compensating asymmetry as required by CPT invariance. Observing them would enhance the significance of the measurements very considerably.
- Typically there can be several classes of rescattering channels. The SM weak dynamics select a subsect of those where the compensating asymmetries have to emerge. QCD frameworks like generalized factorization can be invoked to estimate the relative weight of the asymmetries in the different classes. Analyzing them can teach us important lessons about the inner workings of QCD.
- If New Physics generates the required weak phases (or at least contributes significantly to them), it can induce rescattering with novel classes of channels. The pattern in the compensating asymmetries then can tell us something about the features of the New Physics involved.

Direct CP violation can affect transitions involving $\boldsymbol{D^0 - \bar{D}^0}$ oscillations in two different ways, as described later.

11˙1.2. Asymmetries in final state distributions. For channels with two pseudoscalar mesons or a pseudoscalar and a vector meson a CP asymmetry can manifest itself only in a difference between conjugate partial widths. If, however, the final state is more complex – being made up by three pseudoscalar or two vector mesons etc. – then it contains more dynamical information than expressed by its partial width, and CP violation can emerge also through asymmetries in final state distributions. One general comment still applies: since also such CP asymmetries require the interference of two *different* weak amplitudes, within the SM they can occur in Cabibbo suppressed modes only.

In the simplest such scenario one compares CP conjugate *Dalitz plots.* It is quite possible that different regions of a Dalitz plot exhibit CP asymmetries of varying signs that largely cancel each other when one integrates over the whole phase space. I.e., subdomains of the Dalitz plot could contain considerably larger CP asymmetries than the integrated partial width. Once a Dalitz plot is fully understood with all its contributions, one has a powerful new probe. This is not an easy goal to achieve, though, in particular when looking for effects that presumably are not large. It might be more promising as a practical matter to start out with a more euristic approach. I.e., one can start a search for CP asymmetries by just looking at conjugate Dalitz plots. One simple strategy would be to focus on an area with a resonance band and analyze the density in stripes *across* the resonance as to whether there is a difference in CP conjugate plots.

For more complex final states containing four pseudoscalar mesons etc. other probes have to be employed. Consider, e.g., $\boldsymbol{D^0 \to K^+K^-\pi^+\pi^-}$, where one can form a T-odd correlation with the momenta: $\boldsymbol{C_T \equiv \langle \vec{p}_{K^+} \cdot (\vec{p}_{\pi^+} \times \vec{p}_{\pi^-}) \rangle}$. Under time reversal T one has $\boldsymbol{C_T \to -C_T}$ hence the name 'T-odd'. Yet $\boldsymbol{C_T \neq 0}$ does not necessarily establish T violation. Since time reversal is implemented by an *anti*unitary operator, $\boldsymbol{C_T \neq 0}$ can be induced by FSI [107]. While in contrast to the situation with partial width differences FSI are not required to produce an effect, they can act as an 'imposter' here, i.e. induce a T-odd correlation with T-invariant dynamics. This ambiguity can unequivoally be resolved by measuring $\bar{C}_T \equiv \langle \vec{p}_{K^-} \cdot (\vec{p}_{\pi^-} \times \vec{p}_{\pi^+}) \rangle$ in $\bar{D}^0 \to K^+K^-\pi^+\pi^-$; finding $C_T \neq -\bar{C}_T$ establishes CP violation without further ado. Outline of a search carried out at fixed target experiment FOCUS [424] was presented recently in decay $D^0 \to K^-K^+\pi^-\pi^+$, with a preliminary asymmetry $A_T = 0.075 \pm 0.064$ out of a sample of 400 decays.

Decays of *polarized* charm baryons provide us with a similar class of observables; e.g., in $\Lambda_c \Uparrow \to p\pi^+\pi^-$, one can analyse the T-odd correlation $\langle \vec{\sigma}_{\Lambda_c} \cdot (\vec{p}_{\pi^+} \times \vec{p}_{\pi^-}) \rangle$ [444]. Probing $\Lambda_c^+ \to \Lambda l^+ \nu$ for $\langle \vec{\sigma}_{\Lambda_c} \cdot (\vec{p}_\Lambda \times \vec{p}_l) \rangle$ or $\langle \vec{\sigma}_\Lambda \cdot (\vec{p}_\Lambda \times \vec{p}_l) \rangle$ is a particularly intriguing case; for in this reaction there are not even electromagnetic FSI that could fake T violation. The presence of a net polarization transverse to the decay plane depends on the weak phases and Lorentz structures of the contributing transition operators. Like in the well-known case of the muon transverse polarization in $K^+ \to \mu^+\pi^0\nu$ decays the T-odd correlation is controlled by $\mathrm{Im}(T_-/T_+)$, where T_- and T_+ denote the helicity violating and conserving amplitudes, respectively. Since the former are basically absent in the SM, a transverse polarization requires the intervention of New Physics to provide the required helicity violating amplitude. This can happen in models with multiple Higgs fields, which can interfere with amplitudes due to W exchange. For $b \to u\ell\nu$ the two relevant amplitudes are:

$$\mathcal{T}_{W-X} = \frac{G_F}{\sqrt{2}} V_{ub} (\bar{u}\gamma^\alpha(1-\gamma_5)b)(\bar{\ell}\gamma_\alpha(1-\gamma_5)\nu_\ell) \tag{295}$$

$$\mathcal{T}_{H-X} = \frac{G_F}{\sqrt{2}} V_{ub} \sum_i \frac{C_i m_b m_\ell}{\langle v \rangle^2} (\bar{u}(1-\gamma_5)b)(\bar{\ell}(1-\gamma_5)\nu_\ell). \tag{296}$$

with $\langle v \rangle$ denoting the average of the vacuum expectation values for the (neutral) Higgs fields. For an order of magnitude estimate one can calculate the transverse polarization on the quark level. Unfortunately one finds tiny effects: $\mathcal{O}(10^{-5})$ for $b \to u\tau\nu$ even taking $\mathrm{Im}(C_i) \approx 1$. For charm such effects are even further suppressed since decays into final states with τ leptons are not allowed kinematically.

11˙2. *CP asymmetries involving oscillations.* – In processes involving oscillations there are actually two gateways through which CP violation can enter. In the notation introduced in Sect.**10** one can have

$$\left| \frac{q}{p} \right| \neq 1 \,, \tag{297}$$

which unequivocally constitutes CP violation in $\Delta C = 2$ dynamics, and

$$\mathrm{Im} \frac{q}{p} \bar{\rho}_f \neq 0 \,, \tag{298}$$

which reflects the interplay of CP violation in $\Delta C = 1\&2$ dynamics.

The first effect can be cleanly searched for in semileptonic transitions:

$$\mathrm{rate}(D^0(t) \to l^+\nu X) \propto e^{-\Gamma_H t} + e^{-\Gamma_L t} + 2e^{\bar{\Gamma} t} \cos \Delta M_D t \tag{299}$$

$$\mathrm{rate}(D^0(t) \to l^-\nu X) \propto \left| \frac{q}{p} \right|^2 \left(e^{-\Gamma_H t} + e^{-\Gamma_L t} - 2e^{\bar{\Gamma} t} \cos \Delta M_D t \right) \tag{300}$$

and therefore

$$a_{SL}(D^0) \equiv \frac{\Gamma(D^0(t) \to l^- \nu X) - \Gamma(\bar{D}^0(t) \to l^+ \nu X)}{\Gamma(D^0(t) \to l^- \nu X) + \Gamma(\bar{D}^0(t) \to l^+ \nu X)} = \frac{|q|^4 - |p|^4}{|q|^4 + |p|^4} \,. \tag{301}$$

While the rates depend on the time of decay in a non-trivial manner, the asymmetry does not.

The second effect can be looked for in decays to final states that are CP eigenstates, like $D \to K^+K^-$, $\pi^+\pi^-$, in close qualitative – though not quantitative – analogy to $B_d \to \psi K_S$ or $\pi^+\pi^-$ ([33]). The general expressions for the decay rate as a function of (proper) time are lengthy:

$$\begin{aligned}\mathrm{rate}(D^0(t) \to K^+K^-) &\propto e^{-\Gamma_1 t}|T(D^0 \to K^+K^-)|^2 \times \\ &\Big[1 + e^{\Delta\Gamma t} + 2e^{\frac{1}{2}\Delta\Gamma t}\cos\Delta m_D t + \\ &\left(1 + e^{\Delta\Gamma t} - 2e^{\frac{1}{2}\Delta\Gamma t}\cos\Delta m_D t\right)\left|\frac{q}{p}\right|^2 |\bar{\rho}(K^+K^-)|^2 + \\ &2\left(1 - e^{\Delta\Gamma t}\right)\mathrm{Re}\frac{q}{p}\bar{\rho}(K^+K^-) - 4e^{\frac{1}{2}\Delta\Gamma t}\sin\Delta m_D t \,\mathrm{Im}\frac{q}{p}\bar{\rho}(K^+K^-)\Big]\end{aligned} \tag{302}$$

$$\begin{aligned}\mathrm{rate}(\bar{D}^0(t) \to K^+K^-) &\propto e^{-\Gamma_1 t}|T(\bar{D}^0 \to K^+K^-)|^2 \times \\ &\Big[1 + e^{\Delta\Gamma t} + 2e^{\frac{1}{2}\Delta\Gamma t}\cos\Delta m_D t + \\ &\left(1 + e^{\Delta\Gamma t} - 2e^{\frac{1}{2}\Delta\Gamma t}\cos\Delta m_D t\right)\left|\frac{p}{q}\right|^2 |\rho(K^+K^-)|^2 + \\ &2\left(1 - e^{\Delta\Gamma t}\right)\mathrm{Re}\frac{p}{q}\rho(K^+K^-) - 4e^{\frac{1}{2}\Delta\Gamma t}\sin\Delta m_D t \,\mathrm{Im}\frac{p}{q}\rho(K^+K^-)\Big]\end{aligned} \tag{303}$$

To enhance the transparency of these expressions we simplify them by assuming there is no direct CP violation – $|\bar{\rho}(K^+K^-)| = 1$ – and no purely superweak CP violation – $|q| = |p|$:

$$\mathrm{rate}(D^0(t) \to K^+K^-) \propto 2e^{-\Gamma_1 t}|T(D^0 \to K^+K^-)|^2 \times$$

([33])In principle one can also use $D \to K_S\phi$; however one has to make sure that the observed final state $K_S K^+K^-$ is really due to $K_S\phi$ rather than $K_S f_0(980)$; for in the latter case it would constitute a CP even rather than a CP odd state like the former, and therefore the signal would be washed out.

$$\left[1 + e^{\Delta\Gamma t} + \left(1 - e^{\Delta\Gamma t}\right)\mathrm{Re}\frac{q}{p}\bar{\rho}(K^+K^-) - 2e^{\frac{1}{2}\Delta\Gamma t}\sin\Delta m_D t\,\mathrm{Im}\frac{q}{p}\bar{\rho}(K^+K^-)\right] \tag{304}$$

$$\mathrm{rate}(\bar{D}^0(t) \to K^+K^-) \propto 2e^{-\Gamma_1 t}|T(D^0 \to K^+K^-)|^2 \times$$

$$\left[1 + e^{\Delta\Gamma t} + \left(1 - e^{\Delta\Gamma t}\right)\mathrm{Re}\frac{q}{p}\bar{\rho}(K^+K^-) + 2e^{\frac{1}{2}\Delta\Gamma t}\sin\Delta m_D t\,\mathrm{Im}\frac{q}{p}\bar{\rho}(K^+K^-)\right] \tag{305}$$

With those simplifications one has for the relative asymmetry as a function of t:

$$\frac{\mathrm{rate}(D^0(t) \to K^+K^-) - \mathrm{rate}(\bar{D}^0(t) \to K^+K^-)}{\mathrm{rate}(D^0(t) \to K^+K^-) + \mathrm{rate}(\bar{D}^0(t) \to K^+K^-)} =$$

$$-\frac{2e^{\frac{1}{2}\Delta\Gamma t}\sin\Delta m_D t\,\mathrm{Im}\frac{q}{p}\bar{\rho}(K^+K^-)}{1 + e^{\Delta\Gamma t} + \left(1 - e^{\Delta\Gamma t}\right)\mathrm{Re}\frac{q}{p}\bar{\rho}(K^+K^-)} \tag{306}$$

As mentioned before if CP is conserved, then the mass eigenstates have to be CP eigenstates as well. It is interesting to note how this comes about in these expressions: CP invariance implies $|q| = |p|$, $|\bar{\rho}(K^+K^-)| = 1$ and $\frac{q}{p}\bar{\rho}(K^+K^-) = \pm 1$ and thus $\mathrm{rate}(D^0(t) \to K^+K^-) \propto e^{-\Gamma t}|T(D^0 \to K^+K^-)|^2$.

Comparing doubly Cabibbo suppressed modes $D^0(t) \to K^+\pi^-$, Eq.(306) with $\bar{D}^0(t) \to K^-\pi^+$ allows particularly intriguing CP tests, since SM effects there are highly suppressed. Assuming again $|q/p|^2 = 1$ and also $|T(D^0 \to K^+\pi^-)|^2 = |T(\bar{D}^0 \to K^-\pi^+)|^2$, $|T(D^0 \to K^-\pi^+)|^2 = |T(\bar{D}^0 \to K^+\pi^-)|^2$ to enhance the transparency of the expressions we arrive at:

$$\frac{\mathrm{rate}(D^0(t) \to K^+\pi^-) - \mathrm{rate}(\bar{D}^0(t) \to K^-\pi^+)}{\mathrm{rate}(D^0(t) \to K^+\pi^-) + \mathrm{rate}(\bar{D}^0(t) \to K^-\pi^+)} =$$

$$\frac{-x'_D\frac{t}{\tau_D}\sin\phi_{K\pi}|\hat{\rho}_{K\pi}|}{\mathrm{tg}^2\theta_C + \frac{(x_D^2+y_D^2)\frac{t^2}{\tau_D^2}|\hat{\rho}_{K\pi}|^2}{4\mathrm{tg}^2\theta_C} + y'_D\frac{t}{\tau_D}\cos\phi_{K\pi}|\hat{\rho}_{K\pi}|} \tag{307}$$

This provides such a promising lab since the SM contribution is highly suppressed by $\mathrm{tg}^2\theta_C$ in *amplitude*.

As mentioned above, *direct* CP violation can manifest itself in two different ways, when $D^0\bar{D}^0$ oscillations are involved:

- there can be a difference in the absolute size of the CP conjugate amplitudes

$$|T(D^0 \to f)|^2 \neq |T(\bar{D}^0 \to \bar{f})|^2 \tag{308}$$

This produces a $\cos\Delta M_D t$ term.

- It can induce a difference in the quantity $\mathrm{Im}\frac{q}{p}\bar{\rho}_f$ for two different CP eigenstates $f_{1,2}$

$$\mathrm{Im}\frac{q}{p}\bar{\rho}_{f_1} \neq \mathrm{Im}\frac{q}{p}\bar{\rho}_{f_2} \tag{309}$$

  leading to different coefficients for the $\sin\Delta M_D t$ term.

These CP asymmetries involving $D^0 - \bar{D}^0$ oscillations depend on the time of decay in an essential manner. Producing neutral D mesons in a symmetric e^+e^- machine just above threshold precludes measuring such asymmetries. For the time evolution of the difference in, say,

$$\begin{aligned} &e^+e^- \to D^0\bar{D}^0 \to (l^+X)_{t_1}(K^+K^-)_{t_2} \\ vs.\ \ &e^+e^- \to D^0\bar{D}^0 \to (l^-X)_{t_1}(K^+K^-)_{t_2} \end{aligned} \tag{310}$$

is proportional to $\sin\Delta m_D(t_1 - t_2)$ since the $D^0 - \bar{D}^0$ pair is produced as a C odd state; it vanishes upon integration over the times of decay t_1 and t_2. This quantum mechanical effect is of course the very reason why the e^+e^- B factories are asymmetric.

Nevertheless there are two ways to search for CP violation involving $D^0 - \bar{D}^0$ oscillations at such a machine, at least in principle:

- In $e^+e^- \to D^0\bar{D}^0\gamma$, the $D\bar{D}$ pair is produced in a C even state, and the dependance on the times of decay $t_{1,2}$ is $\sin\Delta m_D(t_1 + t_2)$; one finds for the asymmetry in $(l^+X)_D(K^+K^-)_D$ vs. $(l^-X)_D(K^+K^-)_D$ integrated over all times of decay $2x_D\mathrm{Im}\frac{q}{p}\bar{\rho}(K^+K^-)$. This result is actually easily understood: averaging the asymmetry in a coherent C odd and even $D^0\bar{D}^0$ pair yields $\frac{1}{2}\left[0 + 2x_D\mathrm{Im}\frac{q}{p}\bar{\rho}(K^+K^-)\right] = x_D\mathrm{Im}\frac{q}{p}\bar{\rho}(K^+K^-)$, which coincides with what one finds for a incoherently produced neutral D meson.
- The reaction

$$e^+e^- \to \psi'' \to D^0\bar{D}^0 \to f_a f_b\ , \tag{311}$$

  where f_a and f_b are CP eigenstates that are either both even or both odd, can occur *only if* CP is violated. For the initial state is CP even, the combined final state CP odd since f_a and f_b have to form a P wave:

$$CP[\psi''] = +1 \neq CP[f_a f_b] = (-1)^{l=1}\eta_a\eta_b = -1 \tag{312}$$

  It is thus the mere existence of a reaction that establishes CP violation. It is *not* neccessary for the two states $f_{a,b}$ to be the same.

  By explicit calculation one obtains

$$BR(D^0\bar{D}^0|_{C=-} \to f_a f_b) \simeq BR(D \to f_a)BR(D \to f_b)\cdot$$

$$\left[2\left|\bar{\rho}(f_a) - \bar{\rho}(f_b)\right|^2 + x_D^2\left|1 - \frac{q}{p}\bar{\rho}(f_a)\frac{q}{p}\bar{\rho}(f_b)\right|^2\right] \tag{313}$$

  There are several intriguing subscenarios in this reaction.

– In the absence of CP violation – $\frac{q}{p}\bar{\rho}(f_a) = \pm 1 = \frac{q}{p}\bar{\rho}(f_b)$ – the reaction cannot proceed, as already stated.

– In the absence of $D^0 - \bar{D}^0$ oscillations – $x_D = 0$ – it can proceed only if $\bar{\rho}(f_a) \neq \bar{\rho}(f_b)$, i.e. if $D \to f_a$ and $D \to f_b$ show a different amount of direct CP violation.

– Without such direct CP violation the transition requires $x_D \neq 0$. In contrast to the situation with $D^0(t) \to K^+K^-$ it is actually quadratic in the small quantity x_D.

– Eq.313 can actually be applied also when f_a and f_b are not CP eigenstates, yet still modes common to D^0 and $\bar{D}^0$. Consider for example $f_a = K^+K^-$ and $f_b = K^\pm\pi^\mp$ or $f_a = f_b = K^+\pi^-$. Measuring the rate will then yield information also on the strong phase shifts.

11˙3. *Theory expectations and predictions* . – As outlined above, there are three classes of quantities representing CP violation: (i) $|T(D^0 \to f)| \neq |T(\bar{D}^0 \to \bar{f})|$; (ii) $|q| \neq |p|$; (iii) $\mathrm{Im}\frac{q}{p}\bar{\rho}_f \neq 0$.

The Wolfenstein representation of the CKM matrix reveals one of the best pieces of evidence that the SM is incomplete.

$$V_{\mathrm{CKM}} = \begin{pmatrix} 1 - \frac{1}{2}\lambda^2 & \lambda & A\lambda^3\left(\rho - i\eta + \frac{i}{2}\eta\lambda^2\right) \\ -\lambda & 1 - \frac{1}{2}\lambda^2 - i\eta A^2\lambda^4 & A\lambda^2\left(1 + i\eta\lambda^2\right) \\ A\lambda^3\left(1 - \rho - i\eta\right) & -A\lambda^2 & 1 \end{pmatrix} \tag{314}$$

up to higher orders in $\lambda = \sin^2\theta_C$. On general grounds V_{CKM} has to be unitary; yet Eq.(314) exhibits a peculiar pattern: V_{CKM} is 'almost' diagonal – even close to the identity matrix –, almost symmetric with elements that get smaller further away from the diagonal; this can hardly be accidental.

On a practical level it shows that up to higher order in λ mixing with the third family induces only an imaginary part for the charged current couplings of charm with light quarks. The most relevant source for CP violation is the phase in $V(cs)$, which is $\simeq \eta A^2\lambda^4 \simeq \eta|V(cb)|^2 \sim 10^{-3}$ which provides a very rough benchmark number. One can easily draw quark box and penguin diagrams where this phase enters. That does not mean, though, that one knows how to calculate contributions from these diagrams. For since the internal quarks – the strange quarks – are lighter than charm, these diagrams do *not* represent local or even short distance contributions. We have no theoretical description reliably based on QCD to calculate such quantities. The usual panacea – lattice QCD – has first to mature to a unquenched state before it has a chance to yield reliable results. This is meant as a call for a healthy dose of skepticism rather than negativism.

There is one contribution to q/p which can be calculated reliably, namely the one from the quark box diagram with $b\bar{b}$ as internal quarks, since it reflects a local operator. Yet as already mentioned it makes a tiny contribution to $D^0 - \bar{D}^0$ oscillations.

Discuss CP violation in the condensate contribution to D^0 oscillations ...

Within the SM no direct CP violation can emerge in Cabibbo favoured and doubly suppressed modes, since in both cases there is but a single weak amplitude. Observing direct CP violation there would establish the intervention of New Physics. There is one exception to this general statement [445]: the transition $D^\pm \to K_S\pi^\pm$ reflects the interference between $D^+ \to \bar{K}^0\pi^+$ and $D^+ \to K^0\pi^+$ which are Cabibbo favoured

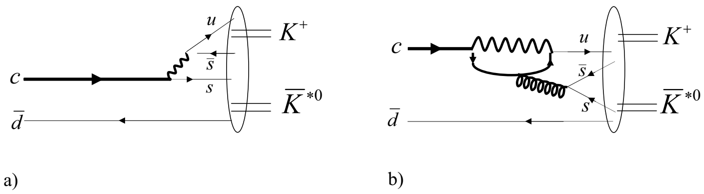

Fig. 35. – Tree and penguin graphs for $D^+ \to \bar{K}^{*0} K^+$. As explained before the leading term in the penguin diagram here represents a nonlocal operator, since the internal s quark in the loop is lighter than the external charm quark.

and doubly Cabibbo suppressed, respectively. Furthermore in all likelihood those two amplitudes will exhibit different phase shifts since they differ in their isospin content. The known CP impurity in the K_S state induces a difference without any theory uncertainty:

$$\frac{\Gamma(D^+ \to K_S \pi^+) - \Gamma(D^- \to K_S \pi^-)}{\Gamma(D^+ \to K_S \pi^+) + \Gamma(D^- \to K_S \pi^-)} = -2\mathrm{Re}\epsilon_K$$

$$\simeq -3.3 \cdot 10^{-3} \tag{315}$$

In that case the same asymmetry both in magnitude as well as sign arises for the experimentally much more challenging final state with a K_L.

If on the other hand New Physics is present in $\Delta C = 1$ dynamics, most likely in the doubly Cabibbo transition, then both the sign and the size of an asymmetry can be different from the number in Eq.(315), and by itself it would make a contribution of the opposite sign to the asymmetry in $D^+ \to K_L \pi^+$ vs. $D^- \to K_L \pi^-$.

In singly-Cabibbo suppressed modes on the other hand there are two amplitudes driven by $c \to d\bar{d}u$ and $c \to s\bar{s}u$, respectively. How they can interfere and generate CP asymmetries can be illustrated by the two diagrams in (Fig.35: the one in a) is a tree diagram, the one in b) of the Penguin type. The latter not only has a different weak phase than the former, but – being a one-loop diagram with internal quarks s, d being lighter than the c quark – induces also a strong phase. One cannot rely, though, on this diagram to obtain a quantitative prediction. This penguin diagram with $m_{s,d} < m_c$ does not represent a local operator and cannot reliably be treated by short-distance dynamics.

Searching for *direct* CP violation in Cabibbo suppressed D decays as a sign for New Physics would however represent a very complex challenge: the CKM benchmark mentioned above points to asymmetries of order 0.1 %. One can perform an analysis like in Ref. [373] based on theoretical engineering of hadronic matrix elements and their phase shifts as described above. There one typically finds asymmetries $\sim \mathcal{O}(10^{-4})$, i.e. somewhat smaller than the rough benchmark stated above. Yet 10^{-3} effects are conceivable, and even 1% effects cannot be ruled out completely. Observing a CP asymmetry in charm decays would certainly be a first rate discovery even irrespective of its theoretical interpretation. Yet to make a case that a signal in a singly Cabibbo suppressed mode reveals New Physics is quite iffy. In all likelihood one has to analyze at least several channels with comparable sensitivity to acquire a measure of confidence in one's interpretation.

The interpretation is much clearer for a CP asymmetry involving oscillations, where one compares the time evolution of transitions like $D^0(t) \to K_S\phi$, K^+K^-, $\pi^+\pi^-$ [446] and/or $D^0(t) \to K^+\pi^-$ [422, 421] with their CP conjugate channels. A difference for a final state f would depend on the product

$$\sin(\Delta m_D t) \cdot \mathrm{Im}\frac{q}{p}[T(\bar{D} \to f)/T(D \to \bar{f})] \; . \tag{316}$$

With both factors being $\sim \mathcal{O}(10^{-3})$ in the SM one predicts a practically zero asymmetry $\leq 10^{-5}$. Yet New Physics could generate considerably larger values, namely $x_D \sim \mathcal{O}(0.01)$, $\mathrm{Im}\frac{q}{p}[T(\bar{D} \to f)/T(D \to \bar{f})] \sim \mathcal{O}(0.1)$ leading to an asymmetry of $\mathcal{O}(10^{-3})$. One should note that the oscillation dependant term is linear in the small quantity x_D (and in t) – $\sin\Delta m_D t \simeq x_D t/\tau_D$ – in contrast to r_D which is quadratic: $r_D \equiv \frac{D^0 \to l^- X}{D^0 \to l^+ X} \simeq \frac{x_D^2 + y_D^2}{2}$. It would be very hard to see $r_D = 10^{-4}$ in CP insensitive rates. It could well happen that $D^0 - \bar{D}^0$ oscillations are first discovered in such CP asymmetries!

Predictions about CP asymmetries in Dalitz plot and other final state distributions are at present even less reliable beyond the general statement that within the CKM description they can arise only in Cabibbo suppressed modes. However it seems conceivable that progress could be made there by a more careful evaluation of the available tools rather than having to hope for a theoretical breakthrough.

11˙4. *Data.* – Experiments so far have basically searched for direct CP asymmetries in partial widths and quite recently in Dalitz plot distributions. Yet experiments have accumulated data sets sufficiently large for meaningful searches for CP asymmetries that can occur only in the presence of $D^0\bar{D}^0$ oscillations. As just argued the second class of CP asymmetries, if observed, would quite unequivocally establish the intervention of New Physics. Recent reviews are in [447].

11˙4.1. Direct CP asymmetries in partial widths. The data analysis strategy is quite straightforward. The asymmetry between CP conjugate partial widths stated in Eq.(**11**˙1) has the advantage that most systematic uncertainties cancel out in the ratio. The particle/antiparticle nature is determined by D^*-tagging for D^0 and by the charge of the final state for D^+, D_s^+ and Λ_c^+. .

As mentioned above, probing for direct CP asymmetries in *Cabibbo favoured* modes represents a search for New Physics, which is required for providing the second weak phase. It has, though, some elements of "searching for lost keys in the night under the lamp posts": one has not necessarily lost the keys there, but it is the first place one can look for them. DCSD channels with their highly suppressed SM contributions are much more promising in that respect.

In fixed target experiments (E791, FOCUS) [220, 157, 113] one has to correct for the notorious particle/antiparticle *production* asymmetries which occur in hadronization due to associated production, leading particle effects etc. This is done by measuring the asymmetry normalized to a copious mode that is not likely to present a CP asymmetry, such as $K^-\pi^+$ or $K^-\pi^+\pi^+$. The main sources of systematic errors are then the tiny differences in reconstruction efficiency, particle identification cuts, and absorption of secondaries in the target and spectrometer.

At e^+e^- colliders no production asymmetry needs to be accounted for. Standard vertexing techniques are used with D^*-tagging for D^0, and refit of soft pion to the D^0

production point. *Ad-hoc* analysis is employed for neutral modes such as $\pi^0\pi^0$ or $K_s\pi^0$ where D^0 daughter neutral tracks do not have sufficient precision to reconstruct the direction of D^0.

The D^0 and D^+ sections of [131] list more than twenty measurements of A_{CP} (Tab.XXV), all consistent with zero, the best ones at the percent level. SM direct CP violation is searched for by experiments by investigating accessable twobody SCS decays channels such as $K^+K^-, \pi^+\pi^-, \pi^0\pi^0$. Fixed target experiments and e^+e^- experiment CLEO have about the same sensitivity for charged channels (several thousand events samples), while CLEO has an obvious advantage in channels with neutral pions. Limits on A_{CP} are at the 1% level, with the important recent limit of the mode $K^0_S\pi^+$ discussed in Sect.**11**˙3. These searches for charm decay rate asymmetries can profitably be done at B factories. The measurements are relatively simple, since most of systematics cancels in the ratio, and do not require any lifetime information. We should expect a factor of two-three improvement in A_{CP} limits, once BABAR and BELLE come onto this scene having gathered a few 10^6 charm event sample, which will bring the present limits below the 1 % level. A similar level of sensitivity can be expected from CDF.

A further improvement below the 0.1% level is expected from BTeV with 10^7 charm event sample. The only other player in this game could be hadroproduction experiment COMPASS[448] at CERN in its eventual phase II, envisioning several 10^6 reconstructed charm meson decays.

11˙4.2. Dalitz plot distributions. The analysis of three-body charged final states is complicated by the possibility of intermediate resonant states such as $K^{*0}K^+$ and $\phi\pi^+$, and thus requires a Dalitz plot analysis. As a byproduct of their Dalitz plot resonant analysis, CLEO has published measurements of A_{CP} *integrated* over the entire Dalitz plot for several channels[397, 396], which do not add much additional information to the decay rate asymmetry measurement. CLEO has also looked for a New Physics CP-violating phase in CF mode $K^-\pi^+\pi^0$, as well as for CP violation in DCS decays under the assumption that the second phase necessary for CP violating effects to materialize be provided by $D^0\bar{D}^0$ mixing.

Dalitz plots are promising fields to search for direct CP asymmetries, as already mentioned. Since FSI are required it is actually quite likely that integrating over the whole Dalitz plot will tend to wash out CP asymmetries. Differences in amplitudes and phases of resonant structures for three-body particle and antiparticle final states are regarded as a sensitive portal for accessing CP violation. No experiment so far has dared publishing results. FOCUS has only shown [434] preliminary results on measured amplitudes and phases $\theta \equiv arg(A_i) + \delta_i$ separately for particle and antiparticle, with statistical errors only, finding zero asymmetry. CLEO [396] only mention that no difference is observed in amplitudes and phases without quantitative statements. The plethora of arguments on the interpretation of Dalitz plots in charm decays discussed in Sect.**9**˙5.1 makes it a complex challenge to apply the Dalitz plot formalism to CP studies. Yet we would like to re-iterate that it could bear precious fruit in charm decays while preparing us for similar studies in beauty decays like $B \to 3\pi \to \rho\pi,\ \sigma\pi$.

11˙4.3. Indirect CP asymmetries. No measurements on time-dependent asymmetries have been published so far, although data sets of hefty size exist now for some of the interesting channels.

We will limit ourselves to briefly discussing the experimental reach in three case studies of indirect CP asymmetries. They all require as essential ingredients superb vertexing for precise lifetime resolution, and large statistics; particle identification and lepton tag capabilities are also a must. In principle such studies can be performed in hadronic collisions, photoproduction and continuum production at e^+e^- B factories. At τ-charm factories like BES and CLEO-c one does not have access to the lifetime information; yet employing quantum correlations one can infer the same information from comparing $e^+e^- \to \psi(3770) \to D^0\bar{D}^0$ and $e^+e^- \to D^{*0}\bar{D}^0 + h.c. \to D^0\bar{D}^0 + \gamma/\pi^0$ as explained in Sect.**11**˙2.

1. $D^0(t) \to l^-X$ vs. $\bar{D}^0(t) \to l^+X$ — Although present experiments such as FOCUS have a 20,000 event sample of semileptonic decays in hand, it is surprising that no progress has occurred in semileptonic measurements concerning oscillations or CP violation beyond the ancient measurement by E791 [428]. The only new finding is FOCUS' tantalizing (yet unpublished) promise of a 0.2% ?check sensitivity for oscillations in semileptonic D^0 decays. When determining $A_{CP}(t)$ in a few lifetime bins, one should expect a 50-10% statistical error, for each of the FOCUS, BABAR and BELLE data sets.

2. $D^0(t) \to K^+K^-$ or $\pi^+\pi^-$ — In this case, the observation of *any* deviation of the proper time distribution from pure single exponential establishes CP violation. The distributions from FOCUS and BABAR are shown indeed in Fig.33. The proper time distribution from an e^+e^- collider has to be deconvoluted from the resolution function, thus seemingly indicating that fixed target have a distinct advantage in this case. However since the CP asymmetry involves oscillations, it needs 'time' to build up. The maximal effect occurs around $t/\tau_D \sim x_D^{-1}\pi/2 > 10$; therefore one wants to go to as large lifetimes as practically possible. The large-lifetime region is critical because a) low-statistics and b) plagued by outliers events which constitute a source of systematics for the very measurement of lifetimes.

3. $D^0(t) \to K^+\pi^-$ vs. $\bar{D}^0(t) \to K^-\pi^+$ — Only a few hundred events have been gathered so far of this DCS decay, therefore this case looks unfeasible right now below the level of a 50% limit or so.

When looking at experimental possibilities in the medium term the same comments presented in the previous section do apply. FOCUS and CLEO are able to provide first limits right away. BABAR and BELLE can eventually improve by a factor of two-three. CDF could enter the game once they have shown their lifetime resolution, and COMPASS when entering their charm phase. After that, a quantum leap is expected from BTeV with 10^8 reconstructed charm decays, $5\,10^6$ reconstructed semileptonic decays and great lifetime measurement capabilities.

11˙5. *Searching for CPT violation in charm transitions.* – CPT symmetry is a very general property of quantum field theory derived by invoking little more than locality and Lorentz invariance. Despite this impeccable pedigree, it makes sense to ask whether limitations exist. Precisely because the CPT theorem rests on such essential pillars of our present paradigm, we have to make *reasonable* efforts to probe its universal validity. While only rather contrived models of CPT violation have been presented, we should keep in mind that super-string theories – suggested as more fundamental than quantum field theory – are intrinsically *non*-local and thus do not satisfy one of the basic axioms of the CPT theorem; they might thus allow for it, though not demand it. In any case, a veritable industry has sprung up [449].

TABLE XXV. – *CP-violating asymmetries. Data from PDG03 [131] unless noted. Statistical and systematic errors are summed in quadrature.*

$D^0 \to$		Events	A_{CP}
K^+K^-	PDG03 Avg	7k	$+0.005 \pm 0.016$
K^+K^-	BELLE03	36.5k	-0.002 ± 0.007
$K^0_S K^0_S$	CLEO	65	-0.23 ± 0.19
$\pi^+\pi^-$	PDG03 Avg	2.5k	$+0.021 \pm 0.026$
$\pi^0\pi^0$	CLEO	810	$+0.001 \pm 0048$
$K^0_S\phi$	CLEO		-0.028 ± 0.094
$K^0_S\pi^0$	CLEO	9.1k	$+0.001 \pm 0.013$
$K^\pm\pi^\mp$	CLEO - assumes no $D^0\bar{D}^0$ mixing	45	$+0.02 \pm 0.19$
$K^\mp\pi^\pm\pi^0$	CLEO - integr. Dalitz pl.	7k	-0.031 ± 0.086
$K^\pm\pi^\mp\pi^0$	CLEO - assumes no $D^0\bar{D}^0$ mixing	38	$+0.09^{+0.25}_{-0.22}$
$\pi^\mp\pi^\pm\pi^0$	CLEO [396], integr. Dalitz pl.		$+0.01 \pm 0.12$
$D^+ \to$			
$K^0_S\pi^+$	FOCUS	10.6k	-0.016 ± 0.017
$K^0_S K^\pm$	FOCUS	949	$+0.071 \pm 0.062$
$K^+K^-\pi^\pm$	PDG03 Avg	14k	$+0.002 \pm 0.011$
$K^\pm K^{*0}$	PDG03 Avg		-0.02 ± 0.05
$\phi\pi^\pm$	PDG03 Avg		-0.014 ± 0.033
$\pi^+\pi^-\pi^\pm$	E791		-0.017 ± 0.042

11˙5.1. Experimental limits. While no evidence for CPT breaking has been found, one should note that the purely empirical underpinning of CPT invariance – in contrast to its theoretical one – is *not* overly impressive [107]. The 'canonical' tests concerning the equality in the masses and lifetimes of particles and antiparticles yield bounds of typically around 10^{-4} for light flavour states. Some very fine work has been done on CPT tests in K_L decays. However even the often quoted bound $|M_{\bar{K}^0} - M_{K^0}|/M_K < 9 \cdot 10^{-19}$ looks much more impressive than it actually is. For there is little justification of using M_K as yardstick (unless one wants to blame CPT violation on gravity); a better – yet still not well motivated – calibrator is the K_S width leading to $|M_{\bar{K}^0} - M_{K^0}|/\Gamma_{K_S} < 7 \cdot 10^{-5}$. Intriguing future tests have been suggested for neutral K and B meson transitions [450]. Their sensitivity is enhanced by involving $K^0 - \bar{K}^0$ and $B^0 - \bar{B}^0$ oscillations and making use of EPR correlations in $e^+e^- \to \phi(1020) \to K^0\bar{K}^0$ and $e^+e^- \to \Upsilon(4S) \to B_d\bar{B}_d$, respectively. Their analyses will be performed at DAΦNE and by BELLE and BABAR.

Masses and lifetimes of D mesons agree to about the 10^{-3} level, i.e. an order of magnitude worse than for light flavour hadrons. One can also search for CP asymmetries that are also CPT asymmetries like $D^+ \to l^+\nu K_S$ vs. $D^- \to l^-\bar{\nu}K_S$ or $D^0 \to l^+\nu K^-$ vs. $\bar{D}^0 \to l^-\bar{\nu}K^+$. Oscillation phenomena again might be the least unlikely place for CPT violation to surface. They present here a much less favourable stage, though, than for K and B mesons, since $D^0 - \bar{D}^0$ oscillations have not been observed. The *phenomenology* is a straightforward generalization of the formalism presented in Eq.(224) of Sect. **10**˙1. Without CPT invariance we can have $\Lambda_{11} \equiv M_{11} - \frac{i}{2}\Gamma_{11} \neq \Lambda_{22} \equiv M_{22} - \frac{i}{2}\Gamma_{22}$. Then one has a complex parameter for CPT violation: $\xi \equiv$

$\Lambda_{11} - \Lambda_{22}/\Delta\lambda$ with $\Delta\lambda$ being the difference in the eigenvalues of the matrix $M - \frac{i}{2}\Gamma$. It can be probed by comparing the rates for the Cabibbo favoured modes $D \to K\pi$: $A_{CPT}(t') \equiv [\mathrm{rate}(\bar{D}^0(t') \to K^+\pi^-) - \mathrm{rate}(D^0(t') \to K^-\pi^+)]/[\mathrm{rate}(\bar{D}^0(t') \to K^+\pi^-) + \mathrm{rate}(D^0(t') \to K^-\pi^+)]$; t' is the reduced proper time. For slow oscillations – $x_D, y_D \ll 1$ – one derives

$$A_{CPT}(t) = [y_D \mathrm{Re}\xi - x_S \mathrm{Im}\xi] \cdot \frac{t}{\tau_{D^0}} \tag{317}$$

Measurement of the slope for $A_{CPT}(t')$ thus returns the quantity $y_D\mathrm{Re}\xi - x_D\mathrm{Im}\xi$. As a first step a loose limit has been published recently [451] based on a sample of 17,000 $D^0 \to K^-\pi^+$ events. One finds

$$y_D\mathrm{Re}\xi - x_D\mathrm{Im}\xi = 0.0083 \pm 0.0065 \pm 0.0041 \tag{318}$$

For $x = 0, y = 1\%$ this translates into $\mathrm{Re}\xi = 0.83 \pm 0.65 \pm 0.41$ corresponding to a 50% limit. One order of magnitude increase in statistics expected from B-factory will allow to improve the limit to the 10-20% level.

11·6. *Resume.* – *Within SM dynamics* the breaking of CP invariance manifests itself only in small ways:

- The main stage is in *singly Cabibbo suppressed* modes. Direct CP asymmetries in partial widths could be 'as large as' 10^{-3}. There is no theorem, though, ruling out SM effects of 1%.
- In *Cabibbo allowed and doubly forbidden* channels *no direct* CP violation can occur with the exception of modes like $D^\pm \to K_S\pi^\pm$, where interference between a Cabibbo allowed and a doubly forbidden amplitude takes place (and where CP violation in the $K^0 - \bar{K}^0$ complex induces an asymmetry of $2\mathrm{Re}\epsilon_K \simeq 3.3 \cdot 10^{-3}$.
- CP violation involving $D^0 - \bar{D}^0$ oscillations is practically absent.

These expectations should be viewed as carrying an optimistic message: the CP phenomenology is a sensitive probe for New Physics. This is further strengthened by the availability of the D^* trick to flavour tag the decaying charm meson and by the strength of FSI in the charm region, which are a pre-requisite for many asymmetries.

While so far no CP asymmetry has emerged in charm decays, 'the game has only now begun', when one has acquired a sensitivity level of down to very few percent. CP asymmetries of $\sim 10\%$ or even more – in particular for Cabibbo favoured channels – would have been quite a stretch even for New Physics models.

Analyzing Dalitz plots and other final state distributions is a powerful method to probe for CP violation. It is also a challenging one, though, with potential pitfalls, and it is still in its infancy. We hope it will be pushed to maturity.

Ongoing experiments at B factories and the planned program at BTeV will provide data samples that should reveal several effects as small as predicted with the SM.

12. – Charm and the Quark-Gluon Plasma

QCD has been introduced to describe the structure of hadrons and their interactions; as such it has to confine quarks and gluons. Yet that is only one of QCD's 'faces': based on general considerations as well as lattice QCD studies [452] one confidently predicts it to exhibit also a *non*-confining phase, where quarks and gluons are not arranged into individual hadrons, but roam around freely . This scenario is referred to as quark-gluon plasma. Its simple – or simplified – operational characterization is to say the *non*perturbative dynamics between quarks and gluons can be ignored; among other things it leads to the disappearance of the condensate terms for quark and gluon fields mentioned in Sect.**4**˙10.2.

Great experimental efforts are being made to verify its existence. There is much more involved than a 'merely academic' interest in fulling understanding all aspects of QCD: studying the onset of the quark-gluon plasma and its properties can teach us essential lessons on the formation of neutron stars and other exotic celestial objects and even on the early universe as a whole.

Following the analogy with QED one undertakes to create a high energy density of roughly $\mathbf{1\ GeV}\boldsymbol{fm^{-3}}$ (compared to $\mathbf{0.15\ GeV}\boldsymbol{fm^{-3}}$ for ordinary cold nuclear matter) by colliding heavy ions against each other. If this energy density is maintained for a sufficiently long time, a phase transition to the de-confined quark-gluon plasma is expected to occur with ensuing thermalization.

The crucial question is which experimental signatures unequivocally establish this occurrence. One of the early suggestions [453] has been that a *suppression* in the production of heavy quarkonia like J/ψ, ψ', χ_c etc. can signal the transition to the quark-gluon plasma. This is based on a very intuitive picture: the high gluon density prevalent in the plasma 'Debye'-screens the colour interactions between the initially produced c and $\bar{c}$ quarks and other quarks as well. After the matter has cooled down re-establishing the confining phase, hadronization can take place again. Yet the c and $\bar{c}$ quarks have 'lost sight' of each other thus enhancing the production of charm hadrons at the expense of charmonia even more than under normal conditions, i.e. when hadronization is initiated *without* delay. More specific predictions can be made [452]: the more loosely bound ψ' is even more suppressed than J/ψ, one expects certain distributions in $E_\perp$ and Feynman x_F [454] etc.

Gross features of these predictions have indeed been observed in heavy ion collisions at the CERN SPS and at RHIC by the NA38 [455], NA50 [456] and PHENIX [457] collaborations, respectively. These findings were indeed interpreted as signaling the transition to the quark-gluon plasma. However this conclusion has been challenged by authors [458], who offered alternative explanations for the observed suppression that does not invoke the onset of the quark-gluon plasma. A better understanding of charmonia production in ordinary proton-nucleus would certainly help to clarify this important issue. Alternatively one can probe more detailed features of charmonium production, like the polarization of J/ψ and ψ' [459]. As described in Sect.**5**˙1 the application of NRQCD to this problem has so far been met with less than clear success, which could mean that *non*perturbative dynamics cannot be treated by NRQCD, at least not for charmonia. The quark-gluon plasma might offers a scenario where NRQCD's *perturbative* features can be tested and vindicated. This might not only provide a clear signature for the quark-gluon plasma itself, but could also serve as a valuable diagnostics of NRQCD and its limitations.

13. – Summary and Outlook

13˙1. *On Charm's future entries into High Energy Physics' Hall of Fame aka. Pantheon In "Old Europe's" Romanic parlance aka. Valhalla In "Old Europe's" Germanic* *parlance.* – It is a highly popular game in many areas to guess which persons – athletes, artists, writers, actors, politicians – will be judged by posterity as having acquired such relevance due to outstanding achievements that they deserve a special place in history. Such an exercise is often applied also to events, even merely conceivable future ones. We succumb to this playful urge here, since it can act as a concise summary for how future discoveries in charm physics could have a profound impact on our knowledge of fundamental dynamics. Such a list is bound to be subjective, of course, yet that is the charm - - pun intended – of this exercise. There is *no* implication that items left out from the list below are unimportant – merely that we sense no Pantheon potential in them.

We divide candidate discoveries into three categories, namely those that will certainly make it onto the 'Valhalla' list, those that are likely to do so and those with no more than an outside chance; a necessary condition for all three categories of discoveries is of course that they happen.

13˙1.1. Sure bets. Any measurement that unequivocally reveals the intervention of New Physics falls into this category. Finding any mode with lepton number violation like $D \to he^{\pm}\mu^{\mp}$ or $D^0 \to e^{\pm}\mu^{\mp}$, Sect.**7**˙5, or with a familon $D^+ \to h^+ f^0$, Sect.**7**˙6, qualifies– as would $D \to \mu^+\mu^-, e^+e^-$ in clear excess of the SM prediction, Sect.**7**˙4. The last case appears the least unlikely one.

Maybe the best chance to find a clear manifestation of New Physics is to observe an CP asymmetry in Cabibbo allowed or doubly suppressed decays or one that involves $D^0 - \bar{D}^0$ oscillations, Sect.**11**˙3.

13˙1.2. Likely candidates. If lattice QCD's predictions on charmonium spectroscopy, on the decay constants f_D, f_{D_s} and on the form factors in exclusive semileptonic charm decays are fully confirmed by future data with uncertainties not exceeding the very few percent level and if its simulations can be shown to have their systematics under control to this level – both nontrivial 'ifs' – then these results would deserve elite status, since it would demonstrate *quantitative* control over strong dynamics involving both heavy and light flavours.

Finding direct CP violation in once Cabibbo suppressed charm decays would be a first-rate discovery. To decide whether it requires New Physics or could still be accommodated within the SM is, however, a quantitative issue and thus much more iffy. It would be overshadowed, though, if CP violation were also found in Cabibbo allowed and doubly suppressed modes, as listed above.

13˙1.3. On the bubble. Being "on the bubble" means that one hopes to attain a high goal while being conscious that one might loose one's footing very quickly. In our context here it means that discoveries in this category can hope for elevation to Valhalla only if they receive an *unequivocal theoretical* interpretation and/or provide important input to other measurements of Pantheon standing. This is best illustrated by the first example.

In Sect.**10** we have expressed skepticism about our ability to base a conclusive case for New Physics on the observation of $D^0 - \bar{D}^0$ oscillations. This might change, though. Yet in addition such a discovery – whether through $x_D \neq 0$ or $y_D \neq 0$ or both – would certainly be an important one complementing that of $K^0 - \bar{K}^0$ and $B^0 - \bar{B}^0$ oscillations, completing the chapter on meson oscillations and as such represent a text book measurement. Lastly, as explained in Sect.**10**˙6, it could have a significant impact on *interpreting* the transitions $B \to D^{neut}K$ with their anticipated CP asymmetries and how the CKM angle ϕ_3/γ is extracted; ignoring $D^0 - \bar{D}^0$ oscillations could fake a signal for New Physics or alternatively hide it.

For other candidates it is even more essential that unambiguous measurements are matched with a clear-cut theoretical interpretation.

Novel insights into light-flavour hadronic spectroscopy inferred from the final states in semileptonic and nonleptonic charm decays, Sects.**8**˙2.3,**9**˙5 and **9**˙6 could qualify for elite status, if a *comprehensive* picture with good theoretical control emerges for a body of data rather than individual states.

Relying on hidden versus open charm production as signal for the quark-gluon plasma, Sect.**12**, might belong here and maybe the weak lifetimes of $C = 1$ and $C = 2$ hadrons, Sects.**6**˙4, **6**˙5. We are not suggesting that all items listed in this third category carry the same intellectual weight. One should keep in mind that a 'Pantheon' is meant to include all gods and goddesses, not only those of Olympic standing.

13˙2. *On the Future of Charm Physics*. – The discussion in Sects. **4 - 11** has hopefully convinced the reader that relevant new insights into SM dynamics can be gained in charm physics and quite possibly even physics beyond the SM be observed. The pursuit of those goals requires even larger samples of high quality data. Fortunately existing machines have an ongoing program in that direction augmented by novel ideas on triggering and analysing; even new initiatives have been suggested as sketched below. A synopsis of existing and planned initiatives was shown in Sect. **3**˙3, Tables I and II.

13˙2.1. Photoproduction. With the end of fixed target programs at the Fermilab Tevatron and the CERN SPS, no photoproduction experiments with real photons are planned. FOCUS still has significant new analyses in progress and under active consideration, and will provide results in channels such as semileptonics, all-charged hadronic decays, and spectroscopy. The two experiments H1 and ZEUS at HERA will refine and extend their analysis of charm (and even beauty) production and fragmentation in the next few years, until their shutdown planned for 2007. Collisions of on- and off-shell photons (and off-shell weak bosons) with protons provide a dynamical environment more involved than e^+e^- annihilation, yet less complex than hadronic collisions. As an extra bonus one can even study the transition to hadroproduction by going from the deep inelastic regime in momentum transfers – $|q^2| \gg 1$ GeV^2 – to $|q^2| \simeq 0$.

13˙2.2. Hadronic collisions. After having proved to the scientific community the feasibility of doing first-class B physics at a hadron collider, CDF is now opening a new chapter in the study of charm hadrons at the TEVATRON by triggering on them. This new ability will give us access to huge statistics. We can expect – or at least hope – that, even after severe cuts, the available data samples will be of unprecedented size, thus allowing us to study single and multiple charm production including their correlations, and the excitations and decays of single and double charm baryons and their decays. It might even extend considerably our sensitivity for $D^0 - \bar{D}^0$ oscillations, as well as CP violation accompanying them and in Dalitz plot distributions, where the normalization is not essential. The ultimate surprise from CDF would be the ability to decently reconstruct in their electromagnetic calorimeters photons and neutral pions from charm decays.

The COMPASS experiment [448] at CERN was proposed in 1996 with a "phase II" devoted to charm hadroproduction and a full spectrum of intended measurements, from lifetimes to spectroscopy, measurement of f_D and searches for $D^0 - \bar{D}^0$ oscillations and CP violation. The experiment is currently taking data in its muonproduction mode for structure functions measurements. The hadroproduction "phase II" foreseen for 2006 will be devoted to production of exotics and glueballs, and to some charm physics items

[460]. They plan to collect $\mathbf{5\,10^6}$ fully recostructed charm events in a one-year running of $\mathbf{10^7}\boldsymbol{s}$.

The most emphasized item is $\boldsymbol{C=2}$ baryon spectroscopy. Assuming $\boldsymbol{C=2}$ baryons are produced at the SELEX rate with respect to $\boldsymbol{\Lambda_c}$, Sect.**6**˙3.5, scaling rates up and considering efficiencies and acceptances via simulation, they end up with an estimate of $\mathbf{10^4}$ reconstructed $\boldsymbol{C=2}$ baryons. A more pessimistic estimate based on measured hadroproduction cross section for $\boldsymbol{C=1}$ baryons scaled to $\boldsymbol{C=2}$ production yields about 100 reconstructed $\boldsymbol{C=2}$ baryons.

13˙2.3. Beauty Factories. The $\boldsymbol{B}$ factories at KEK and SLAC have had a spectacularly successful start-up with respect to both their technical performance (34) and the impact of their measurements. For those have already promoted the CKM description of heavy flavour dynamics from an ansatz to a tested theory. They are also highly efficient factories of charm hadrons. There are three sources for their production: (i) Continuum production; (ii) production of single charm hadrons in $\boldsymbol{B}$ decays driven by $\boldsymbol{b\to c\bar{u}d,\ c\ell\nu}$ and (iii) production of a pair of charm hadrons or of charmonium in $\boldsymbol{B}$ decays due to $\boldsymbol{b\to c\bar{c}s}$.

There are further advantages beyond the sample size: the multiplicities in the final states are relatively low (though not as low as at a tau-charm factory); good vertexing is available and the two $\boldsymbol{B}$ decays can in general be separated on an event by event basis; in particular for charm emerging from $\boldsymbol{B}$ decays several correlations with quantum numbers like beauty, strangeness, lepton number etc. can be exploited.

Such methods will become especially powerful once one has accumulated sample sizes of about 500 $\boldsymbol{fb^{-1}}$ for $\boldsymbol{e^+e^-\to\Upsilon(4S)\to B\bar{B}}$, which might be the case by 2005, and has succeeded in fully reconstructing one of the beauty mesons in about $\mathbf{10^6}$ events [461]. One can then study the decays and decay chains of the other beauty mesons with exemplary cleanliness. This should allow us to study $\boldsymbol{D^0-\bar{D}^0}$ oscillations with excellent control over systematics and enable us to measure certain quantities for the first time like the absolute values for charm *baryon* branching ratios [359].

13˙2.4. Tau-Charm Factories. Studying charm decays at threshold in $\boldsymbol{e^+e^-}$ annihilation offers many unique advantages:

- Threshold production of charm hadrons leads to very clean low multiplicity final states with very low backgrounds.
- One can employ tagged events to obtain the *absolute* values of charm hadron branching ratios in a model independent way.
- Likewise one can measure the widths for $\boldsymbol{D^+\to\mu^+\nu}$ and $\boldsymbol{D_s^+\to\mu^+\nu}$ with unrivaled control over systematics.
- With the charm hadrons being produced basically at rest the time evolution of $\boldsymbol{D^0}$ decays cannot be measured directly. Yet by comparing EPR correlations in $\boldsymbol{D}$ decays produced in $\boldsymbol{e^+e^-\to D^0\bar{D}0}$, $\boldsymbol{e^+e^-\to D^0\bar{D}^0\gamma}$ and $\boldsymbol{e^+e^-\to D^0\bar{D}^0\pi^0}$ one can deduce whether oscillations are taking place or not, as explained in Sect.**10**.

(34)By the summer of 2002 PEP-II had achieved a luminosity of $\mathbf{4.6\times10^{33}}\,\boldsymbol{cm^{-2}s^{-1}}$ thus exceeding its design goal of $\mathbf{3\times10^{33}}\,\boldsymbol{cm^{-2}s^{-1}}$, while KEK-B had established a new world record of $\mathbf{7.4\times10^{33}}\,\boldsymbol{cm^{-2}s^{-1}}$.

One such machine has been operating since 1990, namely the BEPC collider with the BES detector in Beijing. Presently they are running with a luminosity of $5 \times 10^{30}\, cm^{-2}s^{-1}$ at the J/ψ. There are plans for upgrades leading to considerably increased luminosities in the near future.

To fully exploit the advantages listed above one cannot have enough luminosity. Several proposals have been discussed over the last 15 years for tau-charm factories with the ambitious goal of achieving luminosities of up to the $10^{33} - 10^{34}\, cm^{-2}s^{-1}$ range for the c.m. energy interval of 3 - 5 GeV.

One such project is realized at Cornell University. They plan to operate CESR-c with luminosities $(1.5 - 4.4) \times 10^{32}\, cm^{-2}s^{-1}$ in the range $\sqrt{s} = 3- \leq 5$ GeV for three years starting in 2003 using a modified CLEO-III detector. The goal is to accumulate $\sim 1.3 \times 10^9$ J/ψ, 1×10^9 ψ', 3×10^7 $D\bar{D}$, 1.5×10^6 $D_s^+ D_s^-$ and 4×10^5 $\Lambda_c \bar{\Lambda}_c$ events [462],[463].

Data samples of that size and cleanliness would provide ample material for many important studies of the SM:

- They will presumably complete the charmonium spectroscopy, provide authoritative answers concerning charmonium hybrids and clarify the situation with respect to candidates for glueballs and hybrids in charmonium decays like $J/\psi \to \gamma X$.

- CLEO-c's measurements of the widths for $D^+ \to \mu^+\nu$, $D_s^+ \to \mu^+\nu$ and $D_s^+ \to \tau^+\nu$ will be in a class by themselves: with an integrated luminosity of 3 fb^{-1} the uncertainties are expected to be in the 3 - 5 % regime, about an order of magnitude better than what is achievable at the B factories with an integrated luminosity of 400 fb^{-1}!

- The absolute branching ratios for nonleptonic decays like $D^0 \to K\pi$ and $D^+ \to K\pi\pi$ [$D_s^+ \to \phi\pi$] will be measured with uncertainties not exceeding the 1% [2%] level. The improvement over the present situation would be even greater for charm baryon decays.

- Significant improvements in the direct determination of $|V(cd)|$ and $|V(cs)|$ from $D \to \ell\nu\pi$ and $D \to \ell\nu K$ modes could be obtained.

- One could measure the lepton spectra in inclusive semileptonic decays separately for D^0, D^0 and D_s^+ mesons.

Absolute charm branching ratios and decay sequences represent important engineering inputs for B decays, and the present uncertainties in them are becoming a bottleneck in the analysis of beauty decays.

Extracting precise numbers for the decay constants f_D and f_{D_s} provides important tests for our theoretical control over QCD as it is achievable through lattice QCD as described before. The same motivation applies to *exclusive* semileptonic D decays: to which degree one succeeds in extracting the values of $|V(cs)|$ and $|V(cd)|$ – assumed to be known by imposing three-family unitarity on the CKM matrix – provides a sensitive test for the degree to which lattice QCD can provide a quantitative description of non-perturbative dynamics and can be trusted when applied to *exclusive* semileptonic decays of beauty mesons.

Comparing the measured lepton spectra in *inclusive* semileptonic decays separately for D^0, D^+ and D_s^+ mesons with expectations based on the $1/m_c$ expansion can provide us with novel lessons on the on-set of quark-hadron duality. All of this can be summarized by stating that a tau-charm factory provides numerous and excellent opportunities to probe or even test our mastery over QCD.

It is not clear, however, whether searches for $D^0 - \bar{D}^0$ oscillations and CP violation can be extended in a more than marginal way, as long as the luminosity stays below $10^{33}\, cm^{-2}s^{-1}$.

13˙2.5. Gluon-charm Factory at GSI. The GSI laboratory in Germany has submitted a proposal for building a new complex of accelerators on its site. One of its elements is HESR, a storage ring for antiprotons with energies up to 15 GeV. It might allow an experimental program contributing to three areas of charm physics:

- In $p\bar{p}$ annihilation – implemented with internal targets – one can exploit the superior energy calibration achievable there – $\Delta E \sim$ 100 KeV vs. $\Delta E \sim$ 10 MeV in e^+e^- annihilation – to add to our knowledge on charmonium spectroscopy and to search for charm(onium) hybrids through formation as well as production processes. The methodology employed would be an extension of what was pioneered by the ISR experiment R 704 and the FNAL experiments E 760/E 835.

- One can study open charm production in $\bar{p}$ collisions with heavy nuclei like gold. The goal is to analyze how basic properties of charm quarks like their mass are affected by the nuclear medium, i.e. whether it really lowers their masses, as discussed in Sect.**5**˙7.

- Since $p\bar{p}$ annihilation is driven by the strong forces, it leads to huge data samples of charm hadrons, which could be employed to look for novel phenomena like CP violation in the charm meson sector. Of course, one needs very clean signatures since there is a huge background of non-charm events. For the same reason this is not an environment for precision measurements.

13˙2.6. BTeV. The BTeV experiment [464] at the Fermilab Tevatron was proposed to measure oscillations, CP violation and rare decays in both charm and beauty decays. The goal is *to perform an exhaustive set of measurements of beauty and charm particle properties in order to overdetermine the parameters of the CKM matrix, and to either precisely determine SM parameters or to discover inconsistencies revealing new physics.*

Production of beauty and charm hadrons in $\bar{p}p$ interactions at 2 TeV center of mass energy is peaked in forward-backward regions within a few hundreds milliradians. The experiment thus has effectively a fixed-target geometry. Initially conceived as a two-arm spectrometer, BTeV, due to funding constraints, is now proposed and awaiting approval with a one-arm detector.

While the BTeV detector is made of the standard set of subdetectors for charm and beauty physics (including state-of-the-art em calorimetry), the feature that should allow it to be a protagonist in charm physics is the first level trigger on detached vertices, which makes BTeV efficient for fully hadronic final states. In one year of data taking $(10^7 \mathrm{s})$ at a luminosity of $2 \cdot 10^{32}\ \mathrm{cm}^{-2}\mathrm{s}^{-1}$ $2 \cdot 10^{12}$ $c\bar{c}$ pairs will be produced. With a 1% trigger efficiency and 10% reconstruction efficiency, BTeV expects 10^9 *reconstructed* charm decays in one year (in addition to $4 \cdot 10^4$ reconstructed $B^0 \to J/\psi K^0_s$ and 1500 $B^0 \to \pi^+\pi^-$).

BTeV will be able to perform all studies done so far by fixed-target experiments with the higher hadronic background supposedly tamed by superior detector solutions both in tracking (pixel detector), and in neutral reconstruction (crystal em calorimeter). As far as charm is concerned major emphasis is placed on probing high-impact physics topics such as $D^0\bar{D}^0$ oscillations, CP violation and doublecharm spectroscopy.

With 10^9 reconstructed charm decays, BTeV expects to reach a sensitivity of $(1-2)\cdot 10^{-5}$ on the oscillation parameter r_D in both semileptonic and hadronic final states. With 10^6 reconstructed, background-free SCS decays, the sensitivity to CP violating asymmetries will be of order 10^{-3}; BTeV will thus enter the range of SM predictions for direct CP violation there. Particular care will be taken to keep the systematics under sufficient control. The definition of the primary vertex and secondary interactions in the target, which were the main sources of systematic errors at fixed-target experiments, will not be an issue here.

BTeV is anticipated to measure semileptonic formfactors with a percent level precision – similar to what is expected from CLEO-c, yet with very different systematics. Finally BTeV's photon reconstruction capability will allow high sensitivity probes of $D \to \gamma X$. Within the SM they are – unlike $B \to \gamma X$ – dominated by long-distance dynamics with predictions for their branching ratios ranging from 10^{-6} up to 10^{-4}. These rates can be significantly enhanced in SUSY scenarios with the comparison of $D \to \gamma\rho/\omega$ versus $D \to \gamma K^*$ providing an important diagnostic.

13·2.7. Lattice QCD. The tau-charm factory with the precise measurements it will allow has been called a "QCD machine". Such a label implies that a quantitatively precise theoretical treatment of charm(onium) transitions can be given. Lattice QCD represents our best bet in this respect. This theoretical technology has made considerable progress through the creation of more efficient algorithms together with the availability of ever increasing computing power. We have certainly reason to expect such progress to continue. One can also benefit from two complementary approaches to treating charm physics on the lattice: one can approach the charm scale from below largely by 'brute force', i.e. by using finer and larger lattices; or one can approach it from above by scaling down from the heavy quark limit and from b quarks. In either case it will be essential to perform fully unquenched lattice simulations of QCD, which seems to be within reach. In particular the recent 'manifesto' of Ref. [97] expresses considerable optimism that a real breakthrough in this direction is happening right now. In any case there is good hope that the anticipated progress in our experimental knowledge of charm dynamics will be matched by progress in our theoretical understanding with *defendable* theoretical uncertainties in the decay constants and semileptonic formfactors on the percent level.

14. – 'Fabula docet'

The discovery of charm states achieved much more than 'merely' establish the existence of a second complete quark family – it marked a true paradigm shift in how the community viewed quarks: before the observation of Bjorken scaling in deep inelastic lepton-nucleon scattering quarks were regarded by many as objects of mathematical convenience; certainly after the discoveries of J/ψ and ψ' they were seen as real physical objects albeit confined ones. Many important lessons of philosophical, historical and sociological relevance on progress in general and in the sciences in particular can be drawn from this transition

It also gave nature an opportunity to show its kindness, which is much more than its customary lack of malice: the discovery of charm mesons with lifetimes $\sim 4\cdot 10^{-13}$ sec provided the impetus for developing a new *electronic* technology to resolve such lifetimes, namely silicon microvertex detectors . Those were ready 'just in time' to take on another challenge, namely measuring lifetimes of beauty hadrons. Resolving track lengths of a particle with a lifetime $\sim 4 \cdot 10^{-13}$ sec is to first order equivalent to that of a particle with three times the lifetime and mass. The silicon technology has been and is still experiencing spectacular success in tracking $B^0 - \bar{B}^0$ oscillations, CP asymmetries in $B_d \rightarrow J/\psi K_S$ and tagging top quark decays through beauty hadrons in their final state. B tagging is also an essential tool in searching for Higgs bosons. More generally the search for and analysis of charm hadrons gave rise to new detector components, trigger devices and experimental setups and strategies that are now firmly established in the mainstream of HEP. They have met with spectacular successes in beauty physics.

We have also described how new combinations of previously existing theoretical technologies as well as novel ones contributed to our progress in understanding charm and subsequently beauty dynamics.

Every truly good story actually points to the future as well and in more than one way, and this is certainly the case here as well. The thesis of this review is that a strong case can be made for continuing dedicated studies of charm dynamics to be pursued at existing facilities like CLEO-c, the B factories at KEK and SLAC and FNAL's TEVATRON collider. This case rests on three pillars, one experimental and two more of a theoretical nature:

- More precise data on charm spectroscopy and decays are needed as *engineering input* for refining our analysis of beauty decays. To cite but one non-trivial example concerning the determination of $V(cb)$: to extract $\Gamma(B \rightarrow \ell\nu X_c)$ including its uncertainty accurately from real data, one has to know the masses, quantum numbers and decays of the various charm resonances and combinations occurring there. This applies also when measuring $B \rightarrow \ell\nu D^*$, $l\nu D$ at zero recoil.

- Likewise the theoretical tools for treating beauty decays rely on input from the charm sector. Consider the just mentioned example: sum rules derived from QCD proper relate the contributions of different charm resonances to semileptonic B decays to the basic heavy quark parameters that in turn control the theoretical expressions for B decay rates.

  Furthermore the tools for describing beauty decays have not only been developed first for charm decays, but can still be calibrated and refined with more precise and comprehensive charm data. This will enhance considerably their credibility in treating beauty decays. This is true for considerations based on quark models, light cone sum rules, HQE and in particular for lattice QCD as described in the previous section. In particular for the validation of the latter charm can act as an important bridge, since lattice QCD should be able to approach charm dynamics from higher as well as lower scales.

- The existence of charm hadrons and their basic properties has provided essential confirmation for the SM. Yet at the same time it offers a unique angle to searches for physics beyond the SM, which is 'orthogonal' to other approaches. In contrast to s and b quarks charm is an up-type quark; unlike top it hadronizes and can thus exhibit coherent phenomena that enhance CP asymmetries ([35]) This will however

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([35])Hadronization is typically decried as an evil feature curtailing our ability to treat CP violation in strange decays quantitatively. This is however a short-sighted view: indeed it makes it harder to extract the microscopic quantities describing CP violation; yet without it there not happen 'automatically' – dedicated efforts will be required. The fact that no New Physics has shown up yet in charm transitions should not at all deter us from continuing our searches there. On the contrary it has been only very recently that we have entered the experimental sensitivity level in $D^0 - \bar{D}^0$ oscillations and CP violation, where a signal for New Physics would be believable. Finally, experience has shown time and again that when one gives HEP groups data and time, they will find novel ways to formulate questions to nature and understand its replies.

An artist once declared that true art is due to 10% inspiration and 90% perspiration, i.e. committed long term efforts. This aphorism certainly applies to progress in fundamental physics as demonstrated by charm's tale.

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