Light Hadron Spectroscopy

Abstract

A draft for discussion. Contributors:

- Theory
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1 Quark model for hadrons

1.1 General remarks on the quark model

Hadron spectroscopy has been a major platform for probing many dynamic aspects of strong interactions in the non-perturbative regime. It also bridges fundamental approaches such as lattice QCD calculations with phenomenological studies such as quark model, QCD sum rules, etc., from which insights into the non-pQCD phenomena can be gained.

Quarks, as basic building blocks for hadrons, are bound together by the color force generated by gluon exchanges to form color-singlet hadrons, and the underlying dynamics are described by the widely-tested QCD Lagrangian:

$$\begin{aligned} \mathcal{L} &= i\sum_i \bar{q}_i(x)[\partial_\mu - ig_s\sum_a \frac{1}{2}\lambda^a A^a_\mu(x)]\gamma^\mu q_i(x) \\ &\quad -\frac{1}{4}\sum_a F^a_{\mu\nu}(x)F^{\mu\nu a}(x) - \sum_i \bar{q}_i(x)m_i q_i(x) \ , \end{aligned} \tag{1}$$

where g_s is the strong coupling constant, and i is the flavor index for quark $q_i(x)$, which is a Dirac spinor and a three dimensional vector in color space. λ^a denotes the 3×3 Gell-Mann matrices with $a = 1, \ldots, 8$. Thus, a $q\bar{q}$ pair is the minimum number of quark and antiquark to form a color-singlet meson, while qqq is the minimum to form a color-singlet baryon.

Note that QCD quark-gluon interaction conserves flavor, and the interaction strength is flavor-independent. The only dependence on flavor in the QCD Lagrangian is thus through the quark-mass terms.

In the light quark sector, i.e. u, d and s, the mass difference is relatively small, $m_d - m_u \simeq 3$ MeV and $m_s - m_d = 150$ MeV. Therefore, the strong interactions have an approximate global SU(3) flavor symmetry, and quarks (antiquarks) are assigned to representation $\mathbf{3}$ ($\bar{\mathbf{3}}$). Mesons made of $q\bar{q}$ are then irreducible representations given by the following product decomposition:

$$\mathbf{3}\otimes\bar{\mathbf{3}} = \mathbf{1}+\mathbf{8} \ , \tag{2}$$

and baryons are:

$$\mathbf{3} \otimes \mathbf{3} \otimes \mathbf{3} = \mathbf{1}_a + \mathbf{8}_\lambda + \mathbf{8}_\rho + \mathbf{10}_s, \tag{3}$$

where the subscripts a, s, λ and ρ denote antisymmetry, symmetry, and two mixed symmtries for the two-body substates within the three-quark system.

The SU(3) flavor symmetry implies the existence of flavor nonets with the same J^P but difference charges in the meson spectrum. For example, there are 8 pseudoscalars with masses below 500 MeV and one at about 1 GeV, i.e., π^0, $\pi^\pm$, K^0, $\bar{K^0}$, $K^\pm$, η, and η'. These states together are identified as the 0^- flavor nonet in the meson spectrum.

Similarly, for the lowest-mass baryons, one would expect the existence of a flavor singlet, octet and decuplet. Comparing with the baryon spectrum for the low-lying states, one identifies eight baryons for the octet with $J^P = 1/2^+$: $p(uud)$, $n(udd)$, $\Sigma^+(uus)$, $\Sigma^0(uds)$, $\Sigma^-(dds)$, $\Lambda(uds)$, $\Xi^0(uss)$, and $\Xi^-(dss)$, of which the masses are in a range of $0.9 \sim 1.3$ GeV, and 10 states for the decuplet with $J^P = 3/2^+$: $\Delta^{++}(uuu)$, $\Delta^+(uud)$, $\Delta^0(udd)$, $\Delta^-(ddd)$, $\Sigma^+(uus)$, $\Sigma^0(uds)$, $\Sigma^-(dds)$, $\Xi^0(uss)$, $\Xi^-(dss)$, and $\Omega^-(sss)$ in a range of $1.2 \sim 1.7$ GeV. The flavor singlet baryon with $1/2^+$ has a higher mass arising from the spatial and spin degrees of freedom.

By relating the SU(3) flavor symmetry breaking to the mass term in the QCD Lagrangian, a premier evidence for the validity of the quark model solution is that the hadrons within an SU(3) flavor multiplet will differ in mass linearly due to the presence of different number of s ($\bar{s}$) in the hadrons. Namely, one has

$$\begin{aligned}
\Omega^-(1672) - \Xi(1530) &\simeq \Xi(1530) - \Sigma(1385) \\
&\simeq \Sigma(1385) - \Delta(1232) \\
\phi(1020) - K^*(890) &\simeq K^*(890) - \rho(770) \ .
\end{aligned} \tag{4}$$

A compact expression is given by the Gell-Mann-Okubo mass relation:

$$\Sigma + 3\Lambda = 2(N + \Xi) \ . \tag{5}$$

In brief, there are hypotheses made for the naive quark model [2, 3] as required by QCD: i) Chiral symmetry spontaneous breaking leads to the presence of massive constituent quarks within a hadron as an effective degrees of freedom. ii) Hadrons can be viewed as a quark system in which the gluon fields generate effective potentials that depend on the positions and spins of the massive quarks.

These two hypotheses make the non-relativistic treatment an inspiring approach at leading order. Meanwhile, since the quarks are such massive objects compared to the QCD scale, the creation of constituent quark pairs will be suppressed. Thus, the low-lying states would be $q\bar{q}$ for meson and qqq for baryon. By treating the gluon fields as effective potentials the hadron wavefunction will only depend on quark variables.

Based on those simple hypotheses and accommodating the color, flavor, spin, and spatial degrees of freedom in the wavefunction solutions for the quark binding system, quark model provides an efficient and evidently successful classification for a large number of hadrons observed in experiment [21].

2 Meson spectrum

The ultimate goal of studying the hadron spectroscopy is to learn the dynamics for the constituent interactions. As we know that in the light hadron sector due to the failure of perturbation expansion for QCD, the course towards such an ultimate goal will rely on either phenomenological approaches or lattice QCD (LQCD) calculations. In the past decade the LQCD has experienced drastic improvements along with the fast development of computing resources. But there still exist a lot of technical difficulties in the simulation of a fully non-perturbative QCD process, e.g. so far, an unquenched calculation is still unavailable. In contrast, the development of phenomenologies has marked the whole course of modern sciences, especially, in the study of hadron spectroscopy. Experimental data will provide necessary constraints on the parameters introduced to the theory.

Focussing on the meson spectroscopy in this subsection, we shall learn how to construct a $q\bar{q}$ meson in the quark model. It is also known that apart from the conventional $q\bar{q}$ mesons, the non-Abelian property of QCD also allows the possible existence of bound states which are made of gluons, i.e. the so-called "glueball", and/or a gluon continuum. Furthermore, it is also possible to form multiquark mesons, such as $qq\bar{q}\bar{q}$, and the so-called "hybrid" which contains both $q\bar{q}$ and gluon (g) as its constituents, $q\bar{q}g$. All these unconventional states, if exist, will enrich greatly the meson spectrum and shed light on the dynamics of strong QCD scenario. Thus, searching for those unconventional mesons in experiment has been a topical subject in modern intermediate high-energy physics.

2.1 Conventional meson spectrum

For conventional $q\bar{q}$ meson, a phenomenological study is, on the one hand, to find an empirically efficient way to describe the meson spectrum, and on the other hand keep some general properties from QCD (see the general remarks on the quark model). In the quark model framework, to study the meson spectrum is to construct the Hamiltonian for a color-singlet $q\bar{q}$ system:

$$H = \int d\mathbf{x} \sum_i q_i^\dagger(\mathbf{x})\beta\left(m_i - \frac{\Delta}{2m_i}\right) q_i(\mathbf{x}) + \frac{1}{2}\int d\mathbf{x}d\mathbf{y} V_0(\mathbf{x}-\mathbf{y}) \sum_{ija} q_i^\dagger(\mathbf{x})\frac{\lambda^a}{2} q_i(\mathbf{x}) q_j^\dagger(\mathbf{y})\frac{\lambda^a}{2} q_j(\mathbf{y}) , \tag{6}$$

where i and j are flavor indices, and λ^a is the Gell-Mann Matrices for the SU(3)-color interactions; $V_0(\mathbf{x}-\mathbf{y})$ is a central potential, i.e. only depends on $|(\mathbf{x}-\mathbf{y})|$. Note that this Hamiltonian is independent of flavor as required by QCD. Also, one can see that a particular assumption for the form of the potential V_0 can reflect the QCD properties such as one-gluon-exchange Coulomb approximation, $V_0(\mathbf{x}) = \alpha_s/|\mathbf{x}|$, and confinement potentials.

The potential part of the Hamiltonian operating on a color-singlet $q\bar{q}$, $|(q\bar{q})_\mathbf{1}\rangle$, gives

$$\langle (q\bar{q})_\mathbf{1}|\hat{V}|(q\bar{q})_\mathbf{1}\rangle = -\frac{4}{3}V_0 , \tag{7}$$

where $\hat{V}$ denotes the second term of Eq. (6). Therefore, one can simplify the Hamiltonian for a color-singlet $q_i\bar{q}_j$ system to be:

$$H = \frac{\mathbf{p}_x^2}{2m_i} + \frac{\mathbf{p}_y^2}{2m_j} - \frac{4}{3}V_0(\mathbf{x}-\mathbf{y}) \ , \tag{8}$$

where $\mathbf{p}_x$ ($\mathbf{p}_y$) and $\mathbf{x}$ ($\mathbf{y}$) are the momentum and position of quark i (j), respectively.

In the center-of-mass frame, the Hamiltonian can be written as

$$H = \frac{\mathbf{P}^2}{2M} + \frac{\mathbf{p}^2}{2m} + V(\mathbf{r}) \ , \tag{9}$$

where $\mathbf{P} = \mathbf{p}_x + \mathbf{p}_y$ is the c.m. momentum of the $q\bar{q}$ system; and $\mathbf{r} = \mathbf{x} - \mathbf{y}$ and $\mathbf{p} = (\mathbf{p}_x - \mathbf{p}_y)/2$ are the relative distance and momentum between these two quarks; $M \equiv m_i + m_j$ and $m \equiv m_i m_j/M$ are the total and reduced masses. As $V(\mathbf{r})$ is assumed to be spin-independent, Eq. (8) is invariant under separate orbital and spin rotations. By defining the radial quantum number N, orbital angular momentum L, and total spin $S = 0, 1$, one can express the eigenstate of Eq. (8) (including the spin wavefunctions) as $N^{2S+1}L_J$ with the total angular momentum $J = L$ for $S = 0$ or $J = |L-1|,\ L,\ L+1$ for $S = 1$. With an explicit form of $V(\mathbf{r})$, one in principle will be able to produce a $q\bar{q}$ spectrum to compare with the experimental data. For instance, the ground state will correspond to $N = 1, L = 0$ with $S = 0$ or 1, i.e. a spin singlet 1^1S_0 and a spin triplet 1^3S_0. In the charmonium spectrum, – a suitable example for spin-independent potential quark model, one can identify $\eta_c(2980)$ and $J/\psi(3097)$.

There are also spin-dependent forces between the quarks which result in fine and hyperfine structures in the hadron spectroscopy. In fact, without the spin-dependent forces, the spectrum obtained from the Hamiltonian of Eq. (8) cannot match the pattern observed in experiment. On the other hand, the relativistic effects will break the invariance of Hamiltonian under separate orbital and spin rotations. Therefore, it is natural to introduce spin-dependent forces which are phenomenologically analogous to those in hydrogen atom.

It is attributed to De Rujula, Georgi and Glashow who illustrated that the spin-dependent forces in the quark potential are originated from the short-range QCD one-gluon-exchange (OGE) [1]. In a nonrelativistic expansion, in addition to a colored Coulomb type potential, the OGE leads to the Breit-Fermi interaction:

$$H_{BF} = k\alpha_s \sum_{i<j} S_{ij} \ , \tag{10}$$

where $k = -4/3$ and $-2/3$ are for a $q\bar{q}$ singlet and $q\bar{q}$ in $\bar{\mathbf{3}}$, respectively; and

$$\begin{aligned} S_{ij} \ = \ & \frac{1}{|\mathbf{r}|} - \frac{1}{2m_i m_j}\left(\frac{\mathbf{p}_i\cdot\mathbf{p}_j}{|\mathbf{r}|} + \frac{\mathbf{r}\cdot(\mathbf{r}\cdot\mathbf{p}_i)\mathbf{p}_j}{|\mathbf{r}|^3}\right) - \frac{\pi}{2}\delta(\mathbf{r})\left(\frac{1}{m_i^2} + \frac{1}{m_j^2} + \frac{16\mathbf{S}_i\cdot\mathbf{S}_j}{3m_i m_j}\right) \\ & - \frac{1}{2|\mathbf{r}|^3}\left\{\frac{1}{m_i^2}(\mathbf{r}\times\mathbf{p}_i)\cdot\mathbf{S}_i - \frac{1}{m_j^2}(\mathbf{r}\times\mathbf{p}_j)\cdot\mathbf{S}_j \right. \\ & \left. + \frac{1}{m_i m_j}\left[2(\mathbf{r}\times\mathbf{p}_i)\cdot\mathbf{S}_j - 2(\mathbf{r}\times\mathbf{p}_j)\cdot\mathbf{S}_i - 2\mathbf{S}_i\cdot\mathbf{S}_j + 6\frac{(\mathbf{S}_i\cdot\mathbf{r})(\mathbf{S}_j\cdot\mathbf{r})}{|\mathbf{r}|^2}\right]\right\} + \ldots \end{aligned} \tag{11}$$

Although it is still questionable to apply the OGE picture in the hadron spectrum, their results turn to be in a good agreement with the experimental observations.

We do not review the Goldstone boson exchange model developed in the past decade, but refer readers to recent reviews [4] for details.

2.2 Glueball spectrum

Glueballs are bound states of at least 2 or 3 gluons in a color singlet due to the non-Abelian property of QCD:

$$\begin{aligned} gg: \quad & \mathbf{8} \otimes \mathbf{8} && = \mathbf{1} + \mathbf{8} + \ldots \\ ggg: \quad & \mathbf{8} \otimes \mathbf{8} \otimes \mathbf{8} && = (\mathbf{8} \otimes \mathbf{8})_{\mathbf{8}} \otimes \mathbf{8} = \mathbf{1} + \ldots, \end{aligned} \tag{12}$$

where the charge conjugation is $C = +$ for gg states and $C = -$ for ggg states. Assuming that gluons inside glueballs are massive, for the gg, with the orbital angular momentum $L = 0$, states of 0^{++} and 2^{++} can be formed, and 0^{++} is the groundstate glueball. For ggg, the lowest states are 0^{-+}, 1^{--} and 3^{--}.

Both gg and ggg can form states of which the quantum numbers cannot be produced by the $q\bar{q}$ quark model states. Such states, called "oddballs", will be "smoking gun" in the search for glueball candidates for which, however, experimental evidence is still unavailable. Possible quantum numbers for oddballs are: 0^{--}, 0^{+-}, 1^{-+}, 2^{+-}, 3^{-+}.... However, for gg states, if the gluons inside are massless, $J =$ odd states will be forbidden by Yang's theorem [29] though they may exist in ggg sector.

There is no explicit correlations between the gg and ggg ground state masses though the 0^{++} is expected to be lighter than the 0^{-+}. Theoretical predictions for the glueball masses vary significantly among different approaches.

Early phenomenologies find rather light masses for the scalar glueball, e.g. $M(0^{++}) = 0.65 \sim 1.21$ GeV in the bag model [41, 42, 43], and $M(0^{++}) = 1.15$ GeV in a potential model [44]. Other QCD-based approaches produce larger masses for the scalar such as $M(0^{++}) = 1.52$ GeV in a flux-tube model [33], and $M(0^{++}) \simeq 1.5$ GeV in QCD sum rule calculations [34, 35, 36, 37, 45, 46, 38].

In the last twenty years, extensive numerical studies have been carried out to calculate the glueball spectrum in LQCD. Although earliest LQCD predictions [30, 31] for the glueball masses vary significantly, nowadays, the predictions for several lightest glueballs converge to similar mass region despite of the various approaches being used [32, 7, 8, 39]. The lowest glueball state is the $J^{PC} = 0^{++}$ state (scalar) which has a mass about $1.5 \sim 1.7$ GeV, and the mass ratios of the tensor and psuedoscalar to the scalar is about 1.4 and 1.5, respectively. The newest results [39] of the glueball spectrum from larger and finer lattice are listed in Table 2.2 and Figure 2.2.

The calculations of the glueball spectrum are mostly from quenched lattice QCD. Therefore, it is still an open question, *How large the systematic uncertainty of the quenched approximation would be?* A recent preliminary analysis of the scalar glueball mass based on the MILC dynamical gauge configuration [40] shows that the scalar glueball mass of the dynamic lattice simulation does not change much [40].

Table 1: The final glueball spectrum in physical units. In column 2, the first error is the statistical uncertainty coming from the continuum extrapolation, the second one is the 1% uncertainty resulting from the approximate anisotropy. In column 3, the first error comes from the combined uncertainty of $r_0 M_G$, the second from the uncertainty of $r_0^{-1} = 410(20)\,\mathrm{MeV}$

J^{PC}	$r_0 M_G$	M_G (MeV)
0^{++}	4.16(11)(4)	1710(50)(80)
2^{++}	5.83(5)(6)	2390(30)(120)
0^{-+}	6.25(6)(6)	2560(35)(120)
1^{+-}	7.27(4)(7)	2980(30)(140)
2^{-+}	7.42(7)(7)	3040(40)(150)
3^{+-}	8.79(3)(9)	3600(40)(170)
3^{++}	8.94(6)(9)	3670(50)(180)
1^{--}	9.34(4)(9)	3830(40)(190)
2^{--}	9.77(4)(10)	4010(45)(200)
3^{--}	10.25(4))(10)	4200(45)(200)
2^{+-}	10.32(7)(10)	4230(50)(200)
0^{+-}	11.66(7)(12)	4780(60)(230)

2.3 Glueball signatures

Taking the spectrum of the $q\bar{q}$ nonet of pseudoscalar (0^-) and vector (1^-) as a reference, the scalar nonet should lie in the mass range of 1∼2 GeV [75, 74]. Established states $a_0(1450)$ and $K_0^*(1430)$ in this mass region can be naturally assigned as the $I = 1$ and $I = 1/2$ multiplets, respectively, and thus bring the issue about the $I = 0$ multiplets for the scalar nonet to attention. For $I = 0$, there are more than two states reported in the literature [9, 10, 11, 12, 16], namely, $f_0(1370)$, $f_0(1500)$, $f_0(1710)$, $f_0(1790)$(?), $f_0(1810)$(?). This could be a signal for the existence of scalars beyond conventional quark model classifications such as glueball, hybrid or multiquark states.

There also exists another scalar nonet below 1 GeV, i.e. $f_0(980)$, $a_0(980)$, $\sigma(600)$, and $\kappa(800)$, which are candidates for multiquark or molecule states [5, 6]. Details about this will be discussed in Section XXX.

The up-to-date LQCD result [7, 8, 39] makes the mass region between 1 - 2 GeV extremely interesting for the search for the groundstate scalar glueball ($J^{PC} = 0^{++}$), in particular, due to the observation of more than two f_0 states with similar masses ($f_0(1370)$, $f_0(1500)$ and $f_0(1710)$). However, since the glueball properties are not expected to drastically different from the conventional $q\bar{q}$ mesons, one will be encountered with the difficulty of distinguishing the scalar glueball from conventional $q\bar{q}$ states. Nonetheless, more than two f_0 states with similar masses imply that mixings of the pure gluonic scalar (glueball) with nearby $q\bar{q}$ nonet can occur. This greatly complicates the efforts of identifying the scalar glueball in both experiment and theory. In contrast, signals for oddballs will be decisive evidence for the existence of glueballs. Unfortunately, so far concrete experimental

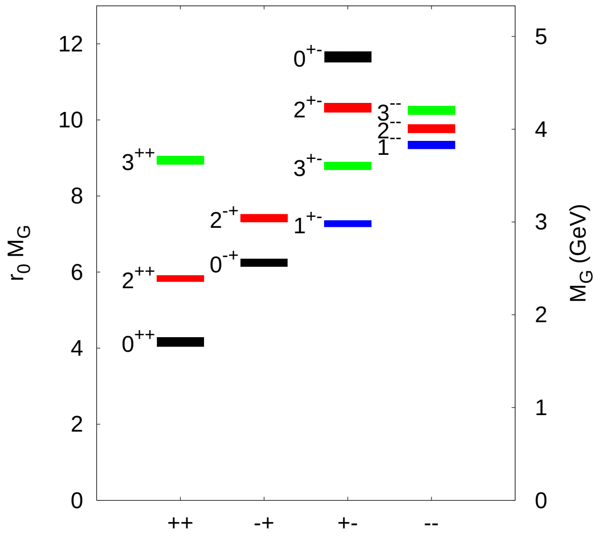

Figure 1: The mass spectrum of glueballs in the pure $SU(3)$ gauge theory. The masses are given both in terms of r_0 ($r_0^{-1} = 410\,\mathrm{MeV}$) and in GeV. The height of each colored box indicates the statistical uncertainty of the mass.

identification of oddballs is still unavailable.

Theoretical expectations for a glueball with conventional quantum numbers have been widely explored in the literature [74]. Although in most cases they are not sufficient for distincting a glueball candidate from conventional $q\bar{q}$ states, they are still useful for providing a guidance for further efforts. We list some of those expectations as a brief review:

i) Flavor-blindness of glueball decays

This character leads to the prediction of the flavor-singlet glueball decay branching ratio fraction:

$$\frac{1}{P.S.}\Gamma(G \to \pi\pi : K\bar{K} : \eta_8\eta_8 : \eta_1\eta_8 : \eta_1\eta_1) = 3 : 4 : 1 : 0 : 1 \ , \tag{13}$$

where P.S. denotes the phase space factors, and η_1 and η_8 are the $I = 0$ flavor singlet and octet of the SU(3) nonet. It can be shown that this relation holds for $\pi\pi : K\bar{K} : \eta\eta : \eta\eta' : \eta'\eta'$ after taking into account the singlet-octet mixing:

$$\begin{aligned} \eta &= \eta_8 \cos\theta - \eta_1 \sin\theta \\ \eta' &= \eta_8 \sin\theta + \eta_1 \cos\theta. \end{aligned} \tag{14}$$

The most significant feature for a pure glueball decays to PP is the vanishing branching ratio for $G \to \eta\eta'$. However, an observed vanishing ratio for $X \to \eta\eta'$ is not necessarily leads to the conclusion that $X = G$ since interferences between different components in a $q\bar{q}$ state can also lead to vanishing $\eta\eta'$ branch ratio [18, 19, 27].

ii) Glueball couplings to $\gamma\gamma$

Since glue is charge neutral, glueballs production in $\gamma\gamma$ collision are suppressed. Similarly, their decays into $\gamma\gamma$ should also be suppressed. Therefore, one would expect that a glueball candidate should have small decay branching ratio to $\gamma\gamma$. However, again, a small b.r. for $X \to \gamma\gamma$ does not necessarily lead to the conclusion that X is a glueball due to interferences from possible mixed components. For instance, in Ref. [15], it is shown that the branching ratio fractions for $f_0^i \to \gamma\gamma$, where $i = 1, 2, 3$ denotes $f_0(1710)$, $f_0(1500)$, and $f_0(1370)$, respectively, are found to be $f_0^1 : f_0^2 : f_0^3 \simeq 1 : 2 : 12$ with larger $\gamma\gamma$ couplings for $f_0(1500)$ than $f_0(1370)$ of which the $q\bar{q}$ component is strongly favored.

Alternatively, the scalar's $q\bar{q}$ component can be probed in e^+e^- annihilation via two virtual photon intermediate states. VEPP and DAΦNE have access to the production of scalars in 1$\sim$ 2 GeV, while the BEPC will be accessible to heavier scalar search at about 3 GeV [76].

iii) Glueball production in heavy quarkonium radiative decays

The J/ψ radiative decay is a gluon rich process and ideal for the search for glueballs as the intermediate resonance in $J/\psi \to \gamma G \to \gamma + \text{all}$. It gives access to all isoscalar states with charge conjugation $C = +$ and forbids all $C = -$ states. Those allowed quantum numbers include the low-lying glueballs and hybrids for which the production phase space are generally large enough. Thus, a systematic study of J/ψ radiative decays with high statistics at BESIII will be extremely important for clarifying some long-standing puzzles.

Information about the intermediate resonances is obtained by measuring their hadronic or radiative decays. Two measures were proposed in the literature to quantify the gluonic contents of the resonances. "Stickiness",

$$S = \frac{\Gamma(J/\psi \to \gamma X) \times P.S.(X \to \gamma\gamma)}{\Gamma(X \to \gamma\gamma) \times P.S.(J/\psi \to \gamma X)} \tag{15}$$

defined by Chanowitz [47], is designed to measure the color-to-electric-charge ratio with the phase space factored out, and maximize the effects from the glue dominance inside a glueball. If X is a glueball, one would expect that its production is favored in J/ψ radiative decays, while its decays into $\gamma\gamma$ are strongly suppressed. Therefore, a glueball should have large stickiness in comparison with a $q\bar{q}$ state.

Cakir and Farrar [48] propose another quantitative measure of the gluonic content of a resonance by calculating its branching ratio to gluons, i.e. $b_{R\to gg}$. This quantity can be related to the production b.r. of resonance R in heavy quarkonium radiative decays. Its value is in a range of $b_{R\to gg} = O(\alpha_s^2) \simeq 0.1 - 0.2$ for a $q\bar{q}$ state, and ~ 1 for a glueball. Interestingly, for most of the well-established $q\bar{q}$ states $b_{R\to gg}$ is found rather small, while for the scalar glueball candidates $f_0(1500)$ and $f_0(1710)$ the value is rather large, in particular for $f_0(1710)$ [15].

These two measures, however, still cannot provide unambiguous evidence for glueballs. The "stickiness" works in a world that glueball and $q\bar{q}$ must not have mixings, i.e. for pure glueballs. In case that glueball-$q\bar{q}$ mixing occurs, the complicated interferences from glueball and $q\bar{q}$ components can violate the expectation of large stickiness corresponding to glueball states. The scheme of Ref. [48, 15] seems to be more sensible because it measures the coupling of a resonance to total gluons.

iv) Chiral suppression

Another criteria pointed out by Carlson [49] and recently developed by Chanowitz [50] is the chiral suppression mechanism for $J = 0$ glueballs. Due to the fact that in pQCD the amplitude is proportional to the current quark mass in the final states, the $J = 0$ glueballs will have larger couplings into e.g. $K\bar{K}$ rather than $\pi\pi$. For $J \neq 0$, the decay amplitude is flavor symmetric. However, due to the complexity of the non-pQCD phenomena and unclear $G - q\bar{q}$ mixings, the observation of relatively larger b.r. to $K\bar{K}$ for a candidate does not necessarily lead to its being a glueball [51, 52].

v) Charmonium hadronic decays

The Charmonium hadronic decays has great advantages of performing a systematic analysis of light hadrons, i.e. both meson and baryon. For instance, the SU(3) flavor symmetry breakings can be investigated by decays of $J/\psi \to VP, VS, VT$; $\chi_{cJ} \to VV$, PP, SS; $\eta_c \to VV$, PP, etc. Nonetheless, the decays of $J/\psi \to Vf_1$ also give access to the axial vectors, $f_1(1285)$ and $f_1(1420)$. An important issue here is to study the properties of the final state mesons and look for evidence for exotics, such as $f_0(980)$ and $a_0(980)$ as four quark states [5] or $K\bar{K}$ molecules; scalar glueball in $f_0(1370)$, $f_0(1500)$, and $f_0(1710)$; and hybrids. Specifically, by recoiling ω and ϕ in $J/\psi \to \omega X$ and ϕX, one can gain information about the flavor components of resonance X. Since ω and ϕ are ideally mixed, i.e. $|\omega\rangle = |u\bar{u} + d\bar{d}\rangle/\sqrt{2}$ and $|\phi\rangle = |s\bar{s}\rangle$, the flavor contents of resonance X can be probed based on the OZI rule [54]. However, it should cautioned that so far the role played by the empirical OZI rule has not been fully understood at the charmonium energy region. Dynamical studies of the possible OZI rule violations should be carried out, which may be essentially important for understanding the nature of many final state mesons [27, 53, 55, 56].

vi) Other glueball-favored processes

The glueball signals are also searched in other reactions, such as $p\bar{p}$ annihilation, and central collisions of $pp \to ppG$. In $p\bar{p}$ reaction, due to the competition of the quark and antiquark rearrangement to form the $q\bar{q}$ meson, identification of glueball signals will be contaminated. Similarly, competitions from the $q\bar{q}$ production in pp central collisions will mix with any possible signals of the glueballs.

It is proposed by Close and Kirk [77] that in the pp central collisions, the production of the S-wave resonances (e.g. 0^{++} and 2^{++} glueball made of vector gluons, or S-wave tetraquark states or $K\bar{K}$ molecules) will be favored the small recoiled transverse momentum difference (dP_T) of the final state protons, while the $q\bar{q}$ production will favor the larger dP_T region (d$P_T \geq O(\Lambda_{QCD})$). Therefore, different kinematic regions serve as a production filter for the S-wave resonances. Following this, experimental analysis at WA102 [78] indeed reveals a clear azimuthal dependence as a function of J^{PC} and P_T at the proton vertices, and scalars appear to divide into two classes: $f_0(980)$, $f_0(1500)$ and $f_0(1710)$ strongly peak at small ϕ angle (corresponding to small dP_T) and $f_0(1370)$ at large ϕ.

In brief, although one could be frustrated by lacking indisputable proof for a glueball candidate with conventional quantum number, putting all together the above expectations and criteria for a glueball candidate one might still be able to identify glueball states from the conventional $q\bar{q}$ background. In the following section, we will analyze some of the controversial states and look for evidence for them to-be, or-not-to-be exotics.

2.4 Glueball candidates

2.4.1 Scalar glueball candidates

In the mass region of $1 \sim 2$ GeV, the abundance of isoscalar scalars, i.e. $f_0(1370)$, $f_0(1500)$, and $f_0(1710)$ ($f_0(1790)$ and $f_0(1810)$ need to be confirmed in further experiment), makes them natural scalar glueball candidates. As follows, we shall brief the experimental information for these states and examine the theoretical expectations of their natures. Controversies will be addressed.

i) $f_0(1370)$

The $f_0(1370)$ turns to be broader than $f_0(1500)$ and $f_0(1710)$, and strongly coupled to $\pi\pi$. Its decays into $K\bar{K}$ were also observed by Crystal Barrel in $p\bar{p}$ annihilation [9], and then confirmed by WA102 in pp scattering [10] in the production channels for $\pi\pi$, $K\bar{K}$ and $\eta\eta$. Meanwhile, its absence in $\eta\eta'$ implies that its configuration is mainly $q\bar{q}$. Taking into account the large ratio of $BR_{f_0(1370)\to\pi\pi}/BR_{f_0(1370)\to K\bar{K}} = 2.17 \pm 0.9$ [10], $f_0(1370)$ can be assigned as an $n\bar{n}$ candidate of $q\bar{q}$ scalar.

The new data from BES [11, 12] also reported observation of $f_0(1370)$ in J/ψ hadronic decays. It appeared in both $J/\psi \to \omega\pi\pi$ and $\phi\pi\pi$. Interestingly, the data show that the branch ratio for $J/\psi \to \phi f_0(1370)$ is larger than that for $J/\psi \to \omega f_0(1370)$. This certainly raised concerns about the $f_0(1370)$ configurations since one would expect that $f_0(1370)$ of $n\bar{n}$ is more likely to recoil against ω instead of ϕ due to the application of the OZI rule.

ii) $f_0(1500)$

The $f_0(1500)$ was observed in many experiments, such as pion induced reaction $\pi^- p$, $p\bar{p}$ annihilation [?, ?], pp central collisions [?, ?] and J/ψ radiative decays [?, ?]. Most of the data on $f_0(1500)$ was from Crystal Barrel collaboration, who resolved two scalar states in this mass region, and determined its decay branching ratios to a number of final states, including $\pi^0\pi^0$, $\eta\eta$, $\eta\eta'$, $K_L K_L$ and $4\pi^0$, using $p\bar{p}$ annihilation at rest. It is also observed that in glueball suppressed processes of $\gamma\gamma$ collision to $K_s K_s$ and $\pi^+\pi^-$, $f_0(1500)$ is absent. All of these favor $f_0(1500)$ being a non-$q\bar{q}$ state.

J/ψ radiative decays have been suggested as promising modes for glueball searches. If $f_0(1500)$ is a scalar glueball, it should be copiously produced in J/ψ radiative decays. The $J/\psi \to \gamma\pi^+\pi^-$ process was analyzed previously in the Mark III [?], DM2 [?] and BES I [?] experiments, in which there was evidence for $f_2(1270)$ and an additional $f_2(1720)$. However, the high mass shoulder of the $f_2(1270)$, at about 1.45 GeV/c^2, was unsettled. A revised amplitude analysis of Mark III data assigned the shoulder to be a scalar at $\sim$ 1.43 GeV/c^2, and, in addition, found the peak at $\sim$ 1.7 GeV/c^2 to be scalar rather than tensor [?]. The $J/\psi \to \gamma\pi^0\pi^0$ process was also studied by the Crystal Ball [?] and BES I experiments [?], but no partial wave analysis has yet been performed on this channel.

Recently, BES reported the results on J/ψ radiative decays to $\pi^+\pi^-$ and $\pi^0\pi^0$ based on a sample of 58M J/ψ events taken with the BES II detector. Partial wave analyses (PWA) are carried out using the relativistic covariant tensor amplitude method in the 1.0 to 2.3 GeV/c^2 $\pi\pi$ mass range. There are conspicuous peaks due to the $f_2(1270)$ and two 0^{++} states in the 1.45 and 1.75 GeV/c^2 mass regions. The first 0^{++} state has a mass of $1466 \pm 6 \pm 20$ MeV/c^2, a width of $108^{+14}_{-11} \pm 25$ MeV/c^2, and a branching fraction $\mathcal{B}(J/\psi \to \gamma f_0(1500) \to \gamma\pi^+\pi^-) = (0.67 \pm 0.02 \pm 0.30) \times 10^{-4}$, which is considered as

$f_0(1500)$. Spin 0 is strongly preferred over spin 2. Fig. 1 shows the $\pi^+\pi^-$ and $\pi^0\pi^0$ invariant mass distributions from $J/\psi \to \gamma\pi^+\pi^-$ and $\gamma\pi^0\pi^0$, where, the corsses are data and the histogram show the PWA fit projection.

Figure 2: The $\pi^+\pi^-$ invariant mass distribution from $J/\psi \to \gamma\pi^+\pi^-$. The crosses are data, the full histogram shows the PWA fit projection, and the shaded histogram corresponds to the background.

Searching for $f_0(1500)$ from more J/ψ decays, such as $J/\psi \to \gamma\eta\eta$, $\gamma\eta\eta'$ etc., and studying its spin-parity are crucial in clarifying the nature of $f_0(1500)$.

iii) $f_0(1710)$

The $f_0(1710)$ is a main competitor of $f_0(1500)$ for being the lightest 0^{++} glueball candidate due to its large production rate in gluon rich processes, such as J/ψ radiative decays, *pp* central production etc., and the predictions of lattice QCD. Table 1 lists the results of $f_0(1710)$ from different experiments before BES2 era. Apparently, different experiments gave different masses, widths and spin-parities. The spin-parity of $f_0(1710)$ in the observed processes is crucial in clarifying whether $f_0(1710)$ is a $q\bar{q}$ or non-$q\bar{q}$ state. If $J = 0$, then the $f_0(1710)$ and $f_0(1500)$ might well represent the glueball and the $q\bar{q}$ state, or more likely each is a mixture of both. However, if $J = 2$, it will be difficult to assign a glueball status to $f_2(1710)$, since that would be at odds with all current lattice gauge calculations.

To study the structures in 1.7 GeV/c^2 mass region, the partial wave analysis is carried out to $J/\psi \to \gamma K\bar{K}$, $\gamma\pi\pi$, $\omega\pi^+\pi^-$, ωK^+K^-, $\phi\pi^+\pi^-$ and ϕK^+K^-, based on $5.8 \times 10^7 J/\psi$ events collected at BES2.

In the $K\bar{K}$ invariant mass spectra from $J/\psi \to \gamma K\bar{K}$, as shown in Fig. 2, the resonant structure in the 1.7 GeV/c^2 mass region is very clearly visible in both decay modes. The partial wave analysis shows that spin 0 is strong preferred over spin 2. The $f_0(1710)$ peaks at a mass of $1740 \pm 4^{+10}_{-25}$ MeV/c^2 with a width of 166^{+5+15}_{-8-10} MeV/c^2.

For $J/\psi \to \omega K^+K^-$, the K^+K^- invariant mass shows a conspicuous signal for $f_0(1710)$, as is shown in Fig. 4. The fittd mass and width are: $M = 1738 \pm 30$ MeV/c^2, $\Gamma = 125 \pm 20$ MeV/c^2.

The ϕ signal and the $\pi^+\pi^-$, K^+K^- invariant mass spectra which recoil against ϕ are shown in Figs. 5 and 6. A peak at around 1.79GeV/c^2 is evident in the $\pi^+\pi^-$ mass spectrum. A simultaneous PWA fit to $J/\psi \to \phi K^+K^-$ and $\phi\pi^+\pi^-$ shows that the peak at around 1.79GeV/c^2 has a mass and width of with $M = 1790^{+40}_{-30}$ MeV/c^2 and $\Gamma = 270^{+60}_{-30}$

Figure 3: Invariant mass spectra of a) K^+K^-, b) $K^0_SK^0_S$ for $J/\psi \to \gamma K\bar{K}$ events, where the shaded histogra correspond to the estimated background contributions.

Figure 4: K^+K^- invariant mass spectrum in $J/\psi \to \omega K^+K^-$. The crosses are data and the histogram is PWA fit projection.

MeV/c^2; spin 0 is preferred over spin 2. This state, $f_0(1790)$, is distinct from $f_0(1710)$, seen in $J/\psi \to \gamma K^+K^-$ and ωK^+K^- channels.

Figure 5: The $\pi^+\pi^-$ invariant mass spectrum in $J/\psi \to \phi\pi^+\pi^-$. The crosses are data and the histogram is PWA fit projection.

For $J/\psi \to \gamma\pi\pi$, the production of a 0^{++} is also observed in 1.7 GeV/c^2 mass region. Its mass and width are $1765^{+4}_{-3} \pm 13$ MeV/c^2 and $145 \pm 8 \pm 69$ MeV/c^2. This 0^{++} state may be $f_0(1710)$ which is observed in $J/\psi \to \gamma K\bar{K}$, or may be $f_0(1790)$ which is shown in $J/\psi \to \phi\pi^+\pi^-$, or may also be a superposition of $f_0(1710)$ and $f_0(1790)$.

The light-meson spectroscopy of scalar states in the mass range of 1-2 GeV/c^2, which has long been a source of controversy, is still very complicated. Overlapping states interfere with each other differently in different production and decay channels. Therefore more

Figure 6: K^+K^- invariant mass spectrum in $J/\psi \to \phi K^+K^-$. The crosses are data and the histogram is PWA fit projection.

Table 2: Table 1

Process	Collaboration	M(MeV)	Γ(MeV)	J^{PC}
$J/\psi \to \gamma\eta\eta$	C. B.(82)	1640 ± 50	200^{+100}_{-70}	2^{++}
$\pi^- p \to K^0_S K^0_S n$	BNL(82)	1771^{+77}_{-53}	200^{+156}_{-9}	0^{++}
$\pi^- N \to K^0_S K^0_S n$	FNAL(84)	1742 ± 15	57 ± 38	——
$\pi^- p \to \eta\eta N$	GAMS(86)	1755 ± 8	< 50	0^{++}
$J/\psi \to \gamma K^+ K^-$	MARK3(87)	1720 ± 14	130 ± 20	2^{++}
$J/\psi \to \gamma K^+ K^-$	DM2(88)	1707 ± 10	166 ± 33	——
$pp \to p(K^+ K^-)p$ $\to p(K^0_S K^0_S)p$	WA76(89)	1713 ± 10 1706 ± 10	181 ± 30 104 ± 30	2^{++}
$J/\psi \to \gamma K\bar{K}$	MARK3(91)	1710 ± 20	186 ± 30	0^{++}
$p\bar{p} \to \pi^0\eta\eta$	E760(93)	1748 ± 10	264 ± 25	$even^{++}$
$J/\psi \to \gamma 4\pi$	MARK3 data D. Bugg(95)	1750 ± 15	160 ± 40	0^{++}
$J/\psi \to \gamma K\bar{K}$	MARK3 data Dunwoodie	1704^{+16}_{-23}	124^{+52}_{-44}	0^{++}
$pp \to p(K\bar{K})p$	WA102(99)	1730 ± 15	100 ± 25	0^{++}
$J/\psi \to \gamma 4\pi$	BES(2000)	1740^{+20}_{-25}	135^{+40}_{-25}	0^{++}

experimental data are needed to clarify the properties of these scalar states.

Scalar $f_0(1790)$ was also reported by BES which turned to be distinct from $f_0(1710)$ [17]. More recently, BES reported a subthreshold enhancement in $J/\psi \to \gamma\omega\phi$ at 1.81 GeV which seemed to favor a 0^{++} assignment for its J^{PC}, i.e. $f_0(1810)$. These states, if confirmed, may raise tremendous interests in the study of their natures [79, 80, 81, 82, 83, 84]. It will also raise essential questions about their production mechanism in the J/ψ hadronic and radiative decays [84].

In additional to the above, recent analyses suggest the existence of a broad scalar $f_0(1200-1600)$ with a half width 500-900 GeV [13, 14]. Nonetheless, the sign of such a broad state leads to a rearrangement of the scalar nonets, where $\sigma(600)$ and $\kappa(800)$ are no longer physical states. In contrast to those two nonets, above/below 1 GeV, the new nonets are arranged to be: i) $f_0(980)$, $f_0(1300)$, $a_0(980)$, $K_0(1415)$ for the $q\bar{q}$ with radial quantum number $n=1$; and ii) $f_0(1500)$, $f_0(1750)$, $a_0(1520)$, $K_0(1820)$ for $n=2$. The broad $f_0(1200-1600)$ is then regarded as the descendant of the scalar glueball.

In brief, at the energies of 1$\sim$2 GeV, at least three isoscalar scalars, $f_0(1370)$, $f_0(1500)$ and $f_0(1710)$, are established in experiment. However, there appear quite unexpected behaviors of these states in different processes, which raises a lot of questions about their nature. Those questions include: What are the constituent structures of these scalars? Is any one of these scalars a glueball state? Is the glueball a pure state? What does the present experimental information tell us about the scalar production and decay mechanisms? ...

2.4.2 Pseudoscalar glueball candidates

In PDG [21], $\eta(1475)$ and $\eta(1295)$ are assigned to the radially excited 0^{-+} of 2^1S_0 corresponding to η and η', while the third state $\eta(1405)$ is identified to be a pseudoscalar glueball candidate based on its strong coupling to $K\bar{K}\pi$ and $\eta\pi\pi$.

PDG2006 gives a very detailed review of $\eta(1405)$ and $\eta(1475)$.

- $\eta(1295)$
- $\eta(1405)$
- $\eta(1475)$

2.4.3 Tensor glueball candidates

Lattice QCD predicted 2^{++} tensor glueball as the second lowest state with a mass around 2.3 GeV, which makes it interesting to search for this state at the mass region of around 2.3 GeV in experiment.

Mark III first presented signals for a narrow state ($\Gamma \sim 20$ MeV), so-called $\xi(2230)$, at 2.2 GeV in $J/\psi \to \gamma K^+K^-$ [65], and later in $\gamma K_s^0 K_s^0$ [66]. However, there was no clear signal seen at DM2 in the same channels [67]. In hadron scattering experiments, GAMS Collaboration found a structure at $m = 2220 \pm 10$ MeV with a width of $\Gamma \sim 80$ MeV in $\pi^- p \to \eta\eta' n$ in the $\eta\eta'$ invariant mass spectrum [68], while the LASS group reported a structure at 2.2 GeV in the invariant mass of $K_s^0 K_s^0$ in $K^- p \to K_s^0 K_s^0 \Lambda$ [69].

Further evidence for this state was from BES-I with $\sim 8^6$ J/ψ events. A structure was reported in $\pi^+\pi^-$, K^+K^-, $K_s^0K_s^0$, $p\bar{p}$, and $\pi^0\pi^0$ [70, 71], which however showed controversies to the data from $p\bar{p}$ annihilations to $K_s^0K_s^0$ [72], $\eta\eta$ and $\pi^0\pi^0$ [73].

2.5 Accessible processes probing the glueball contents at BESIII

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