Present status of glueball searches

Kamal K. Seth*[a]

[a]Physics Department, Northwestern University,
Evanston, IL 60201, USA

The present status of experimental searches for glueballs is reviewed. A critical examination is made of the criteria used in the identification of glueballs. It is concluded that the mixing of glueball states with normal $q\bar{q}$ states deprives us so far of any 'smoking guns', and at least two states, $f_0(1500)$ and $f_0(1710)$ remain viable contenders for the scalar (0^{++}) glueball. It is also shown that in the proton-antiproton annihilation experiments no evidence is found for the existence of $\xi(2230)$, the much heralded candidate for the tensor glueball.

1. INTRODUCTION

'Gluons', with the property of self-coupling, were proposed long before quarks[1], in order to explain the strong interaction. Fermi and Yang[2] had proposed that the 'elementary' particles were the nucleons and the antinucleons, and that pions were states of $N\bar{N}$ bound by a very strong force; the 'glue'. Now, of course, we accept that the elementary strongly interacting fermions are the colored quarks of six flavours, and the quanta of the strong interaction are the octet of coloured gluons. Already in 1972 Fritzch and Gell-Mann[3] pointed out that "if the quark-gluon field theory indeed yields a correct description of the strong interactions, there must exist glue states in the hadron spectrum", and Fritzch and Minkowski[4] presented a detailed discussion of the phenomenology of the glueball ("glue-states") spectrum.

Glueballs are bosons, with baryon number and isospin equal to zero. They are SU(3) singlets, but only "to the extent to which SU(3) breaking effects can be neglected[4]". As color singlets, two-gluon glueballs have the same quantum numbers as two-photon systems, i.e. $J^{PC} = 2n^{++}, (2n+1)^{++}$, and $2n^{-+}$, with $n = 0, 1, 2....$

For a long time glueballs were a curiosity item. Any non-baryonic enhancement which did not fit the 'standard picture' of a normal quark-antiquark meson was considered as a candidate for a glueball. Early glueball candidates were iota (now called $\eta(1440)$)[5], theta (now called $f_J(1710)$), g(2300) (now called $f_2(2300, 2340)$), and $\xi(2230)$. A number of rather loosely defined criteria were used to advance the candidacy of the observed structures as glueballs, and even as the Higgs.

That era is clearly over. Serious pursuits of glueballs are now on. This makes it necessary now to carefully examine the criteria which have been advanced for the identification

*kseth@nwu.edu

of glueballs. In this review, I intend to do that. I hope that the necessary critique will not offend the advocates of one or the other glueball candidate.

2. THE SPECTRUM OF GLUEBALLS

The expected spectrum of glueballs has been calculated in many QCD based models, and on the lattice. These include: (a) MIT Bag model[6], and its variant, the Constituent Gluon model[7], (b) Flux-Tube Model[8], (c) QCD Sum-Rule calculations[9] , and (d) Lattice calculations[10-18]. Besides giving references to the model calculations, I will not discuss them any further. I will only discuss briefly the current status of lattice calculations of the glueball spectrum. According to the folklore these calculations are the most reliable, and we should look at what their predictions are, and how reliable and stable they are.

As a non-expert I believe that lattice calculations are still in a state of flux. Most calculations are still in the quenched approximation, and substantial uncertainties in obtaining continuum limits, and converting lattice units to physical units remain. Better and improved actions are being continuously proposed, and the predictions of glueball masses continue to oscillate. Whether the latest means the best is by no means certain. However, I quote from some of the latest.

CP-PACS tells us[10] that the quenched hadron mass "spectrum in the continuum limit deviates systematically from experiment on a level of about 10%(5σ)", and that the "study of full QCD is in a less advanced stage". Morningstar and Peardon[11] warn us that special "problems arise with simulations of the lightest scalar glueball state". All in all, I subscribe to the cautious summary of Kenway[12] that "knowing its (quenched QCD's) limitations should give us greater confidence in using it as a phenomenological model for quantities where 10-20% accuracy is useful".

Let us look at the evolution of the predictions from three well-known groups about the mass of the lowest scalar glueball. Using my own name designations:

IBMQCD: (1994) 1740 ± 71 MeV[13] $\rightarrow$ (1999) 1632 ± 49 MeV[14]
UKQCD: (1993) 1568 ± 89 MeV[15] $\rightarrow$ (1998) $1611 \pm 30 \pm 160$ MeV[16]
SDQCD: (1997) 1600 ± 100 MeV[17] $\rightarrow$ (1999) $1730 \pm 50 \pm 80$ MeV[18]

I point out this evolution simply to caution experimentalists not to read too much into agreement or disagreement of their favourite glueball candidate with one or the other lattice prediction. Remember that "10-20 % accuracy" is all we can expect at this point.

To set the stage, however tentative, for our expectations, I reproduce in Fig. 1 the predictions of the masses of the fifteen lowest glueballs as calculated by Morningstar and Peardon[18] in the pure SU(3) gauge theory. Note that the negative charge conjugation glueballs are necessarily three-gluon glueballs. Also, the authors point out that for glueballs with masses greater than ~ 3.65 GeV the J assignment is uncertain, and that the masses in Table 1 "neglect the effects of light quarks and mixing with nearby conventional mesons." This means that even the predicted ratios of masses cannot be used for direct comparisons with experiment, as some people try to do. To summarize: Take lattice with a grain of salt!!

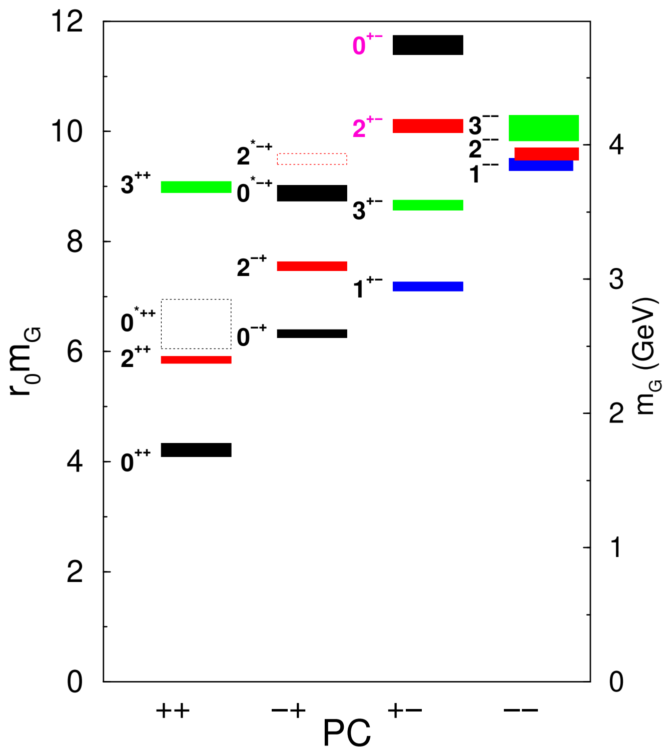

FIG. 1. Lattice predictions of glueball masses.[18]

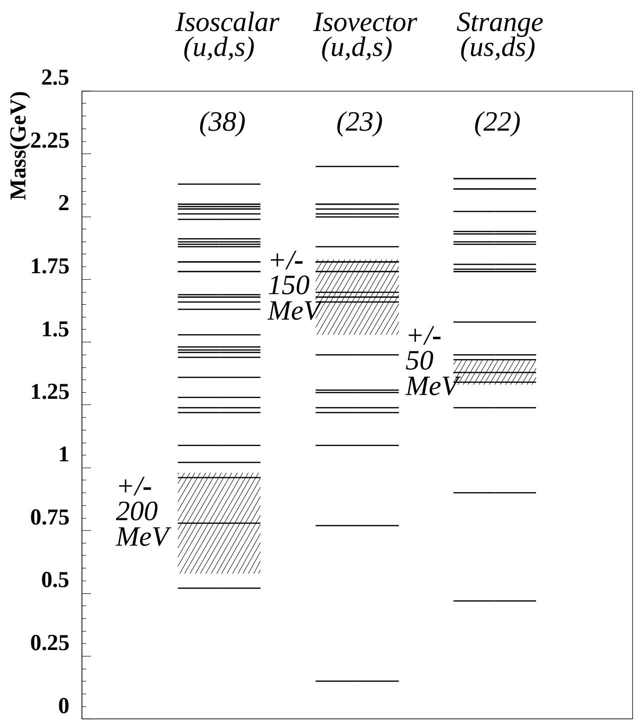

FIG. 2. Spectrum of light $q\bar{q}$ mesons.

Since this review is confined to the discussion of the 0^{++} scalar and the 2^{++} tensor glueballs, I will take to heart my own advice and look for the 0^{++} glueball in the general vicinity of $M(0^{++}) \approx 1600$ MeV and the 2^{++} glueball in the general vicinity of $M(2^{++}) \approx 2300$ MeV. To appreciate what this means in terms of trouble, let us look at Fig. 2, where the spectrum of normal $q\bar{q}$ states is reproduced. Our 0^{++} and 2^{++} glueballs are sitting somewhere in the middle of a forest of normal $q\bar{q}$ states, which themselves overlap because of their large widths (100 - 400 MeV). Even without considering the problem of mixing of the glueballs with $q\bar{q}$ states of same J^{PC}, to which we will turn later, it is obvious that searching for glueballs in the thicket of $q\bar{q}$ states will be a challenging job. Claims and counterclaims will abound, and controversies will rage long and loud. And, so they do!

3. HOW TO PRODUCE GLUEBALLS?

The simple answer to this question is "in glue-rich processes". As illustrated in Fig. 3, two obvious processes are J/ψ radiative decay (Fig (3(a)) which has been exploited at SLAC (Mark II, III), Orsay (DM2), and more recently BEPC (BES), and $p\bar{p}$ annihilation (Figs. 3b, 3c) which has been employed at BNL, CERN, and FNAL. A somewhat less obvious means of producing glueballs is central production reactions at high energies (Fig. 3d) with beams of pions, kaons, and protons. Since the colliding particles maintain their identity in these reactions, it is assumed that the exchanged 'particles' are Pomerons, which are considered to be two-gluon entities in their most naive interpretation. Important contributions in this domain have come from GAMS, WA76, and WA102 at CERN.

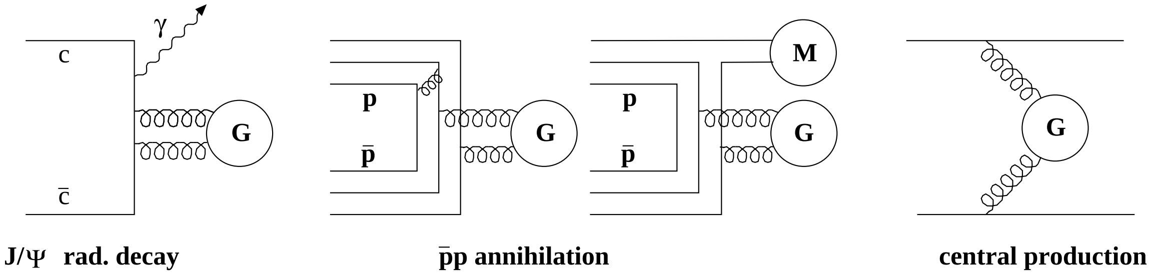

FIG. 3. Glue-rich processes for glueball production.

4. HOW TO SMELL A GLUEBALL?

Obviously, by the characteristics of its decays!

The popular criteria for the production and identification of a glueball can be summarized as:

a) It should be produced in one of the above, favoured reactions.

b) For a given J^{PC}, if it can not be* $q\bar{q}$, it must be exotic, and therefore a candidate glueball. (*cannot be = supernumery, or nonet +)

c) It should be narrower than a typical $q\bar{q}$ meson.

d) It should decay flavour blind, since gluons do not know one flavour from another. Further, a glueball should not like to decay into photons.

Unfortunately, all the above criteria have associated caveats!

5. CAVEATS

Most of the experimental and theoretical investigations to date have been made in relation to the 0^{++} scalar glueball. In the following, I will therefore discuss most of the caveats in relation to the scalar glueball.

5.1. Production reactions

(a) J/ψ Radiative Decay, $J/\psi \to \gamma X$ (Fig 3(a)): Everybody agrees that this is perhaps the cleanest way to produce gg glueballs. We should, however, remember that J/ψ radiative decay produces $q\bar{q}$ mesons far more effectively than any glueballs. This is illustrated by Table I based on Ref.[5].

TABLE I. Radiative decays of J/ψ.[5]

X	$\eta'(958)$	$f_4(2050)$	$f_2(1270)$	$\eta(547)$	$f_1(1285)$	$f_0(1500)$
$\mathrm{B}(J/\psi \to \gamma X) \times 10^3$	4.3(3)	2.7(7)	1.4(1)	0.9(1)	0.7(1)	0.6(1)

The $f_0(1500)$ is, of course, one of the favourite candidates for the scalar glueball. Admitted that the listed branching ratio of $f_0(1500)$ is based only on its decay into 4 pions, and the total branching ratio can be as large as $1.5 \times (0.6(1) \times 10^{-3})$, $f_0(1500)$ production is still more that an order of magnitude smaller than that of the undisputed $q\bar{q}$ mesons, in the table.

(b) $p\bar{p}$ Annihilation (Fig 3(b)): This is also a great way to produce glueballs. Unfortunately, the usually observed final states, consisting of hadrons with u,d quarks can be just as effectively produced by quark rearrangements, without intermediate glueballs. This can enhance (u,d) quark final states (π, η, etc.) compared to the others (K, η', etc.), and upset the expected flavour blindness in glueball decay, unless great care is taken in data analysis.

(c) Central Production and Glueball Filters: Central production reactions may be rich in glue, but they are even richer in quarks! In the original example offered by Close and Kirk[19], it is the $q\bar{q}$ state $f_1(1285)$ which is most strongly produced [Fig. 4.]. Not only that, it is $f_1(1285)$ which has a Δp_T dependence, not $f_0(1500)$. In a recent paper[20], WA102 has provided more details on Δp_T dependence. From their numbers I conclude that there is essentially no difference for any resonance if production with $\Delta p_T > 0.5$ GeV is compared with $\Delta p_t = 0.2$ - 0.5 GeV. Against $\Delta p_T < 0.2$ GeV there is a difference, but it is essentially the same for all well established $J^{PC} = 0^{++}$, I=0 states, f_0(980, 1300, 1500, 1710), and certainly not all are glueballs, or contain equal admixtures of glueballs. So, all in all, I am rather skeptical about the claimed glueball filter - it may turn out to be just an angular momentum filter. We have seen similar effects in our $p\bar{p} \to \pi^0(\pi^0\pi^0)$ reaction at < 6 GeV/c[21].

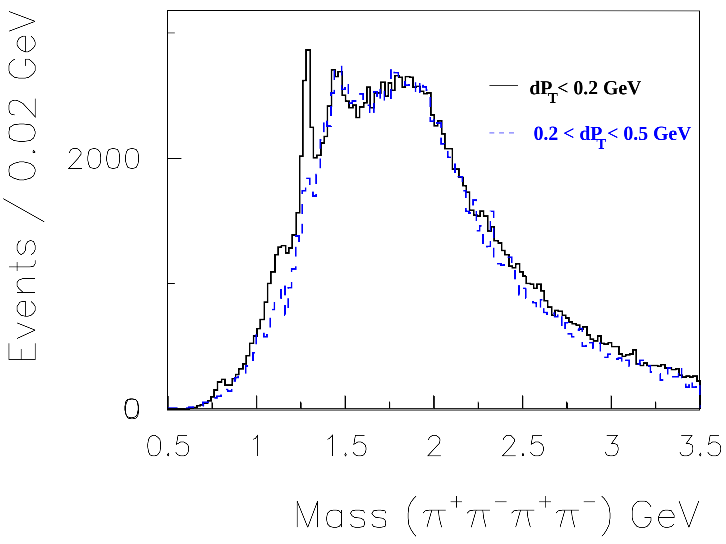

FIG. 4. Central production of $f_1(1285)$ and $f_0(1500)$.[19]

5.2. Supernumery?

In order to judge if we have more states than what is expected for a SU(3) nonet (i.e., we have supernumeries), we have to first identify members of the SU(3) nonet, i.e., one $f_J(n\bar{n})$, one $f_J(s\bar{s})$, one $a_J(n\bar{n})$ - 3 fold, and one $K_J(\bar{n}s, n\bar{s})$ - 4 fold. This is generally not easy, especially for $J^{PC} = 0^{++}$ and 2^{++} states. In Table II, I list the known scalars in the region of interest, and some of their properties.

TABLE II. Parameters of scalar resonances and their decays.[5][21]

Resonance X	$f_0(\sigma)$	$f_0(980)$	$f_0(1370)$	$f_0(1500)$	$f_0(1710)$
$M(X)$ - MeV	800 ± 400	980 ± 10	1350 ± 150	1505 ± 9	1712 ± 5
$\Gamma(X)$ - MeV	800 ± 200	70 ± 30	350 ± 150	111 ± 12	133 ± 14
$\Gamma_{\gamma\gamma}(X)$ - keV	seen (10 ± 6)	0.56 ± 0.11	5.4 ± 2.3	< 3.1	< 11
$B(\pi\pi)(\%)$	dominant	dominant, (75)	26 ± 9	29.0 ± 7.5	15 ± 5
$B(K\bar{K})(\%)$		seen, (25)	35 ± 13	3.5 ± 0.3	60 ± 14
$B(\eta\eta)(\%)$			14 ± 5	4.6 ± 1.3	24 ± 10
$B(\eta\eta')(\%)$				1.2 ± 0.3	
$B(4\pi)(\%)$				61.7 ± 9.6	
$B(J/\psi \to \gamma X) \times 10^3$				0.93 ± 0.13	1.6 ± 0.5

Obviously, there are too many scalars with masses $M \leq 2000$ MeV. We have five f_0's when we needed only two for our ground state nonet. How to decide which two belong, and which are supernumeries. The convenient, but not so convincing, approach taken by some is to throw out $f_0(400-1200, \sigma)$ as an artifact of final state interaction, throw out $f_0(980)$ (along with $a_0(980)$) as a $K\bar{K}$ molecule, and play the game between $f_0(1370)$, $f_0(1500)$, and $f_0(1710)$. (The spin of $f_0(1710)$ has been now convincingly determined to be 0^+[22]). Two of these are considered to be $|n\bar{n}\rangle$ and $|s\bar{s}\rangle$ and the third is the supernumery $|gg\rangle$ glueball.

Having disposed of the 'undesirables', all that remains is to concoct a mixing mechanism between the remaining three to 'fit' the observables, in order to justify ones favourite idea of which one is the dominantly $|gg\rangle$ state. I, for one, do not find this to be an intellectually pleasing situation, irrespective of whether it leads to the Crystal Barrel candidate $f_0(1500)$, or the IBM candidate $f_0(1710)$, as the dominantly scalar glueball state. Indeed, there are very good arguments for $f_0(980)$ to be a member of the g.s. nonet[23][24]. Minkowski and Ochs even argue for $f_0(\sigma)$ to be the glueball[25]. Anisovich and Sarantsev[24] propose a doublet of f_0's (one very wide, $\Gamma \approx 450$ MeV, one very narrow, $\Gamma \approx 65$ MeV) at ≈ 1500 MeV. In other words, the whole premise of putting $f_0(\sigma)$ and $f_0(980)$ aside, and only considering $f_0(1370)$, $f_0(1500)$ and $f_0(1710)$ as one-resonance-too-many (and therefore two nonet members and a glueball) appears to be on rather tenuous grounds.

5.3. Glueballs should be narrower than $q\bar{q}$ resonances.

We do not know how this rumour got started. Perhaps it arose from a very naive interpretation of the Zweig rule, or perhaps from the fact that the χ_{cJ} resonances, which decay hadronically via two gluons (not a bound glueball) have narrow widths $\Gamma(\chi_{c0}, 0^{++}) \approx 15$ MeV, $\Gamma(\chi_{c2}, 2^{++}) \approx 2$ MeV. We do not know, and we should not care.

The fact of the matter is, as already stated, that glueballs reside in a sea of normal $q\bar{q}$ states, and quantum mechanics tells us that if they have the same J^{PC}, they must mix. How strong is the mixing is the central question. Thus, if we have $|n\bar{n}\rangle$, $|s\bar{s}\rangle$, and $|gg\rangle$ progenitor states, the mixed states can have decay properties quite different from those of the progenitor states. It is also worth noting that since the quark-gluon interaction is about twice as strong as the quark-quark interaction, the mixing between glueball states and $q\bar{q}$ states may even be stronger than between two $q\bar{q}$ states. We return to mixing later.

There is a predicted glueball state (see Fig. 1) which can not mix with $q\bar{q}$ states. It is the $J^{PC} = 2^{+-}$ glueball*. Unfortunately, it is predicted to have a mass of ~ 4.1 GeV[18], and it is presently out of reach. [*For a $q\bar{q}$ state $P = (-1)^{L+1}, C = (-1)^{L+S}$.]

5.4. Glueballs must decay flavour blind

In principle, this is the smoking gun for pure glueball identification. Unfortunately, since glueballs are going to be necessarily mixed with $q\bar{q}$ states, the smoke is stolen away from the gun. Further, the flavour blindness expectation is usually evaluated in terms of the SU(3) wave functions for the decay mesons, and SU(3) is a strongly broken symmetry. So, all in all, it is difficult to determine what to expect for the decays of a states containing a fraction (even large) of the glueball state.

For the decay of a pure $|gg\rangle$ glueball, after phase space and form factor effects are removed, one expects for pseudoscalar decays $\pi\pi : K\bar{K} : \eta\eta : \eta\eta' = 1 : 1 : 1 : 0$, per 'charge' state. How realistic is it to expect anything like this for a physical glueball state?

Close has argued that since the decay of the charmonium $^3P_0(\chi_{c0})$ state takes place via two gluons (unbound) the ansatz of flavour-blind decay can be tested in the decays $\chi_{c0} \rightarrow gg \rightarrow (pseudoscalar)^2$. Indeed, with the latest results from BES[26] and Fermilab E835[27] and PDG[5], we find for the decays of χ_{c0} that the branching ratios (units of 10^{-3}) $\pi^+\pi^- : K^+K^- : \eta\eta = 2.3(3) : 2.8(4) : 2.0(10)$. Since phase space corrections are small for χ_{c0} ($m = 3.4$ GeV), this does constitute a verification of the flavour-blindness ansatz for pure gg decays. We note at this point that for $\chi_0 \rightarrow p\bar{p}$ the measured branching ratio is an order of magnitude smaller. The significance of this observation will become clearer in Sec. 7 when we discuss the tensor glueball candidate $\xi(2230)$.

6. THE SCALAR 0^{++} GLUEBALL

There are two very agressively advertized candidates for the ground state scalar glueball. These are the $f_0(1500)$ and the $f_J(1710)$. The $f_0(1500)$ has been extensively studied by the Crystal Barrel collaboration at CERN, and championed by them, and by Amsler and Close[28][29]. The $f_J(1710)$, which is the old $\theta(1700)$, is one of the original glueball candidates, and its case has been mostly advanced by Weingarten *et al.*[13][14]. The candidacy of $f_J(1710)$ was originally not taken seriously because its spin had not been unambiguously determined (0^{++} or 2^{++}). As stated earlier, that problem no longer exists, because WA102 has made the firm determination that its $J^{PC} = 0^{++}$[22]. Thus we now have $f_0(1500)$ and $f_0(1710)$ as the likely candidates for the 0^{++} glueball, with the lattice prediction of a mass $\approx$ 1600 MeV, exactly in between.

With the accepted idea that glueballs will necessarily mix with $q\bar{q}$ states of same J^{PC}, perhaps it is not a terribly meaningful question whether $f_0(1500)$ or $f_0(1710)$ is *the* glueball. The question has come down to which one of the two is *more of the glueball* than the other? To answer this question we have to look into the experimental decay patterns of the two states.

In Table II, I have already presented the totality of the relevant data presently available on the scalar/isoscalar resonances, the f_0's, with mass < 2 GeV. The first thing to notice is that there are large gaps and uncertainties in the data for most resonances. The exception is $f_0(1500)$, for which we should complement the Crystal Barrel. (A word of caution: most of the branching ratios in the table are actually branching fractions, i.e., fractions of the total of all the channels which have been measured so far. Also sometimes there are huge inconsistencies.)

Let us examine the properties of $f_0(1370, 1500, 1710)$, the three states which are generally discussed as $|n\bar{n}\rangle$, $|s\bar{s}\rangle$ and $|gg\rangle$, mixed or unmixed. The $f_0(1500)$ and $f_0(1710)$ are both narrower than $f_0(1370)$ but wider than $f_0(980)$. At first sight the dominant component of $f_0(1500)$ looks like $|n\bar{n}\rangle$, that of $f_0(1710)$ like $|s\bar{s}\rangle$, and $f_0(1370)$ appears to be most ecumenical - ($|gg\rangle$?). One would wish life were this easy, or this trivial! Hunting for glueballs is certainly not that!

At the most naive level we expect that a pure glueball decays flavour-blind into pairs of pseudoscalars with branching ratios in the proportions

$\pi\pi : \eta\eta : \eta\eta' : K\bar{K} = 3 : 1 : 0 : 4$

In the real world, these relative yields have to be modified to take account of phase space and form factors. This corresponds to multiplying the above be $p^{2L+1} exp(-p^2/8\beta^2)$ where p is the cm momentum of the pseudoscalar, L is the relative angular momentum of the pair, and β, the oscillator parameter, is ≈ 0.5[28][29]. The multiplication factor obviously depends on the masses of the decaying state as well as the pseudoscalars. In Table III, I list the final results for the decays normalized to the yield of $\pi\pi$. I also list the results of the experimental measurements. We note that both $f_0(1500)$ and $f_0(1710)$ strongly violate the expectations based on flavour-blind decay in at least one decay; $f_0(1500)$ is deficient by an order of magnitude in $K\bar{K}$, and $f_0(1710)$ is equally deficient in $\pi\pi$. Strangely enough, $f_0(1370)$, which nobody considers as a glueball candidate, comes close to meeting the flavour-blind expectations in all respects!

TABLE III: Decays of scalar resonances into pseudoscalar pairs, and expectations for flavour-blind (FB) decays of pure glueballs. The FB decay $f_0(1500) \to \eta\eta'$ has been omitted because it has a very large phase space factor.

Ratios	$f_0(1370)$ FB decay	$f_0(1370)$ Expt	$f_0(1500)$ FB decay	$f_0(1500)$ Expt	$f_0(1710)$ FB decay	$f_0(1710)$ Expt
$R_1 \equiv \eta\eta/\pi\pi$	0.225	0.54(27)	0.266	0.159(61)	0.297	1.60(85)
$R_2 = \eta\eta'/\pi\pi$			0 (?)	0.041(15)		
$R_3 = K\bar{K}/\pi\pi$	1.056	1.35(85)	1.142	0.121(33)	0.201	4.0(16)

To summarize the above, it is clear that glueballs do not violate quantum mechanics and 'pure' glueballs expectations are not realized. Glueballs mix with $|q\bar{q}\rangle$ states, and whatever else of the same J^{PC} that nature allows.

Now we have the much more difficult, and often subjective, job of what we are going to mix and how. Since most attempts in the literature have been about mixing ground state scalars of $|n\bar{n}\rangle$, $|s\bar{s}\rangle$, and $|gg\rangle$, I will confine myself to these. The radials are, for the present, set aside, although several authors have ventured into that, even more complicated, landscape.

Assuming the SU(3) ideal mixing angle (mixing $|n\bar{n}\rangle$ and $|s\bar{s}\rangle$) of $tan^{-1}(1/\sqrt{2})$, the mixing matrix can be written as

$$\begin{pmatrix} m_g^0 & z & \sqrt{2}z' \\ z & m_s^0 & 0 \\ \sqrt{2}z' & 0 & m_n^0 \end{pmatrix}$$

where m_g^0, m_s^0, and m_n^0 are the progenitor, or 'pure' $|gg\rangle$, $|s\bar{s}\rangle$, and $|n\bar{n}\rangle$ scalar states, respectively, z is the mixing energy $\langle n\bar{n}|V|G_0\rangle$ and z' is the mixing energy $\langle s\bar{s}|V|G_0\rangle$. (often $z = z'$ is assumed.) As Close has pointed out[29] the mixed states can turn out to have their dominant component completely different from their progenitor nature.

The important step then is the choice of the energies of the progenitor states and the mixing energies. This is the stage at which subjectiveness and prejudices can enter. In their original paper, Close and Amsler[28] more or less assumed the order of the energies of the states to be $m_n^0 < m_g^0 < m_s^0$ and then adjusted the mixing to fit the existing branching ratio data for $f_0(1500)$. They found a mixing solution in which the $f_0(1500)$ was dominantly $|gg\rangle$, as intended, and the order of magnitude deficiency in its $K\bar{K}$ decay was explained as a consequence of destructive interference between its $|n\bar{n}\rangle$ and $|s\bar{s}\rangle$ contents. This solution prescribed that $f_0(1370)$ is dominantly $|n\bar{n}\rangle$, and that "the mainly $|s\bar{s}\rangle$ state lies in the 1600-1700 MeV region"[28]. This conclusion was at least partially connected to the $f_J(1710)$ spin being uncertain. In fact, the authors inserted a cautionary note, saying "If $f_J(1710)$ is confirmed to have a $J = 0$ component in $K\bar{K}$ but not in $\pi\pi$, this could be a viable candidate for a $G_0 - s\bar{s}$ mixture."[28]

Now that the spin of $f_J(1710)$ has been determined to be 0^+, we must consider $f_0(1710)$ to be a much more serious contender for the scalar glueball than before. Actually, the IBMQCD group of Weingarten has had $f_0(1710)$ as their candidate for glueball ever since their first lattice calculation of 1995[13]. They claim certain improvements over the completely phenomenological model of Amsler and Close. In a highly detailed paper, which is now available[14], Lee and Weingarten show that their lattice calculations lead to two very important results: (1) In the continuum limit, the progenitor $|s\bar{s}\rangle$ lies well below their progenitor glueball state. They obtain its mass as $m_s^0 = 1322(42)$ MeV, (2) The mixing energy turns out to be a strong function of the lattice spacing, and has a value of 43 ± 31 MeV in the continuum limit. With these two important lattice-based conclusions, and the assumption that the progenitor mass of the isoscalar $|n\bar{n}\rangle$, m_n^0, is equal to the experimental mass of the isovector $a_0(1470)$ they adjust the parameters of the matrix to fit the experimental masses $M(f_0(gg)) = 1710$ MeV, $M(f_0(s\bar{s})) = 1505(9)$ MeV, and $M(f_0(n\bar{n})) = 1404(24)$ MeV. They find that the fit leads to the masses of the progenitor states $m_g^0 = 1622(29)$ MeV, $m_s^0 = 1514(11)$ MeV, and mixing energy $z = 64(13)$ MeV. They consider these results to be in excellent agreement with their lattice results $m_g^0 = 1657(47)$ MeV, $z = 43(31)$ MeV, except for $m_s^0 = 1514$ MeV which is about 13% larger. The physical eigenvectors they obtain are

$|f_0(1710)\rangle = 0.86(5)|gg\rangle + 0.30(8)|s\bar{s}\rangle + 0.41(9)|n\bar{n}\rangle$
$|f_0(1500)\rangle = -0.13(5)|gg\rangle + 0.91(4)|s\bar{s}\rangle - 0.40(11)|n\bar{n}\rangle$
$|f_0(1370)\rangle = -0.50(12)|gg\rangle + 0.29(9)|s\bar{s}\rangle - 0.82(9)|n\bar{n}\rangle$

The main difference in this work from the Amsler-Close approach is that these wave functions have not been obtained to fit the experimental data for branching ratios. As such, they can lead to predictions about the branching ratios. Lee and Weingarten state that the opposite signs of the $|s\bar{s}\rangle$ and $|n\bar{n}\rangle$ components in $|f_0(1500)\rangle$ lead to a suppression of $K\bar{K}$ decay by a factor 0.39(16) in comparison to that of an unmixed $|s\bar{s}\rangle$. This corresponds to the flavour blind decay expectation (Table II) becoming 0.44(18) compared to the measured value 0.12(3). It is a pity that not enough experimental data are available to carry the detailed comparison of this type for $f_0(1710)$ and $f_0(1370)$ decays, as well. In the meanwhile, we summarize the above lattice results as indicating that the dominant components are $f_0(1710) = 74\%|gg\rangle$, $f_0(1500) = 83\%|s\bar{s}\rangle$, and $f_0(1370) = 67\%|n\bar{n}\rangle$.

Of course, one can play other games with the mixing matrix. I find, for example, that the results remain essentially unchanged from the above if $z = z'$ is assumed, and that you can distribute glueball strength essentially equally between all three states by making 1622 MeV as the $|s\bar{s}\rangle$ progenitor state. As Toussaint has shown[30], if you are willing to assume arbitrarily large values of the mixing energy you can make $f_0(\sigma)$ or $f_0(980)$ as the glueball also. But we must regard these excercises as essentially frivolous.

I mentioned earlier that the reluctance of glueballs to couple to photons can also be used as a means of identifying glueballs.[29] Unfortunately, as shown in Table II, the presently available information on decays to $\gamma\gamma$ is too incomplete and imprecise to help in this respect.

A fair summary of the present situation is that $f_0(1370)$, $f_0(1500)$, and $f_0(1710)$ all may very well contain substantial fractions of scalar $|gg\rangle$ glueball. Lattice results favour $f_0(1710)$, but the case is likely to remain unproven unless more experimental data on decay branching ratios become available.

7. THE TENSOR 2^{++} GLUEBALL

For the tensor glueball we have mercifully only one agressively advertized candidate, $\xi(2230)$. Its history is almost as old as that of the scalar glueballs, and is no less confusing. In 1986, Mark III at SLAC, investigating radiative decays of $5.8\times10^6 J/\psi$, reported[31] an abnormally narrow enhancement, dubbed ξ, in the K_SK_S channel ($M(\xi) = 2232\pm10$ MeV, $\Gamma(\xi) = 18^{+23}_{-15} \pm 10$ MeV) and K^+K^- channel ($M(\xi) = 2230 \pm 15$ MeV, $\Gamma(\xi) = 26^{+20}_{-16} \pm 17$ MeV). They reported allowed $J^{PC} = (even)^{++}$. Upper limits (90% CL) were also set for decays of ξ into $\pi\pi$ and $p\bar{p}$ at the level $B(J/\psi \to \gamma\xi)B(\xi \to X) < 2\times10^{-5}$. At nearly the same time, Alde *et al.* also reported[32] a relatively broad enhancement ($\Gamma \le 150$ MeV) at $M = 2220$ MeV in the reaction $\pi^- p \to nX, X \to \eta\eta'$. The legend of a 2^{++} glueball was born.

Unfortunately, in 1988 DM2 at ORSAY reported[33] on the radiative decays of a sample of $8.6\times10^6 J/\psi$, and concluded that in "neither channel K^+K^- or K_SK_S is $\xi(2230)$ seen". Their 95% confidence upper limits for both decays were a factor two lower than the branching ratios reported by Mark III. They did, however, observe a broad enhancement ($\Gamma = 201 \pm 51$ MeV) in the region ($M = 2197 \pm 17$ MeV). It appeared that the legend of the narrow $\xi(2230)$ was dead.

Between 1987 and 1993, a number of attempts were made to look for the narrow ξ in $p\bar{p} \to K_SK_S$ and K^+K^- formation experiments[34]. None reported any success.

The legend of $\xi(2230)$ was revived by BES at BEPC in 1996, with a report[35] of the observation of $\xi(2230)$ in the radiative decay of $8.0 \times 10^6 J/\psi$. They claimed observation of $\xi(2230)$ not just in K_SK_S and K^+K^- channels, but also in $\pi^+\pi^-$ and $p\bar{p}$ channels, with masses in all channels between 2230 and 2235 MeV, widths between 15 and 20 MeV, and product branching ratios $B(J/\psi \to \gamma\xi)B(\xi \to X)$ between 1.5×10^{-5} and 5.6×10^{-5}. The statistical level of significance in the different channels was claimed to range from 3.8σ to 4.6σ. In a subsequent publication[36], BES also reported the observation of $\xi(2230) \to \pi^0\pi^0$ at a 3.6σ significance level, with the product branching ratio $(4.5 \pm 2.9) \times 10^{-5}$. BES also reported 90% confidence limits $B(J/\psi \to \gamma\xi)(\xi \to \eta\eta) < 5.6 \times 10^{-5}$ and $B(J/\psi \to \gamma\xi)(\xi \to \eta\eta') < 2.2 \times 10^{-3}$[39]. It appeared that $\xi(2230)$ was very much alive, and perhaps even narrower than 20 MeV, the reported widths in each channel being almost exactly equal to the experimental resolution width.

The most dramatic part of this revival of $\xi(2230)$ was that a definite branching ratio for the $p\bar{p}$ decay was claimed. Thus arose the challenge to find $\xi(2230)$ in $p\bar{p}$ formation experiments in which excellent mass resolution can be achieved. Since then, the JETSET (PS202) experiment at LEAR(CERN) has investigated $p\bar{p} \to \phi\phi$ and $p\bar{p} \to K_SK_S$ reactions[37][38]. They establish 95% confidence upper limits, $B(p\bar{p} \to \xi) \times B(\xi \to K_SK_S) < 7.5 \times 10^{-5}$ for widths > 5 MeV[37], and $B(p\bar{p} \to \xi)(\xi \to \phi\phi) < 6 \times 10^{-5}$ for width $\Gamma = 15$ MeV[38]. At the Crystal Barrel detector a high resolution search for $\xi(2230)$ was made in the $p\bar{p} \to \pi^0\pi^0, \eta\eta$ channels, among other neutral decay channels.

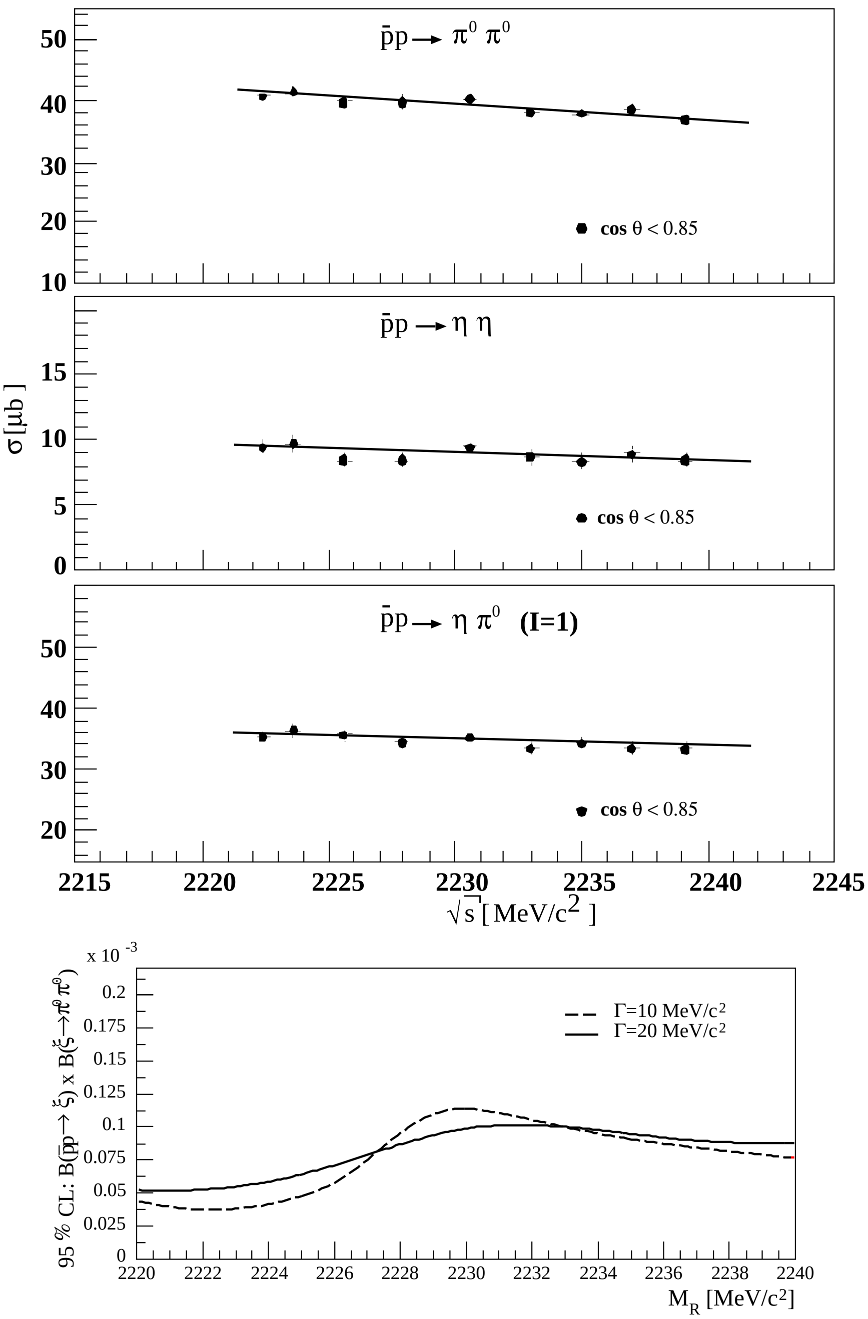

FIG. 5. Top three panels: measured excitation functions for $p\bar{p} \to \pi^0\pi^0$, $\eta\eta$, and $\eta\pi^0$ respectively. Bottom panel: 95% CL upper limits for $B(p\bar{p} \to \xi) \times B(\xi \to \pi^0\pi^0)$ as functions of assumed $M(\xi)$ and $\Gamma(\xi)$.

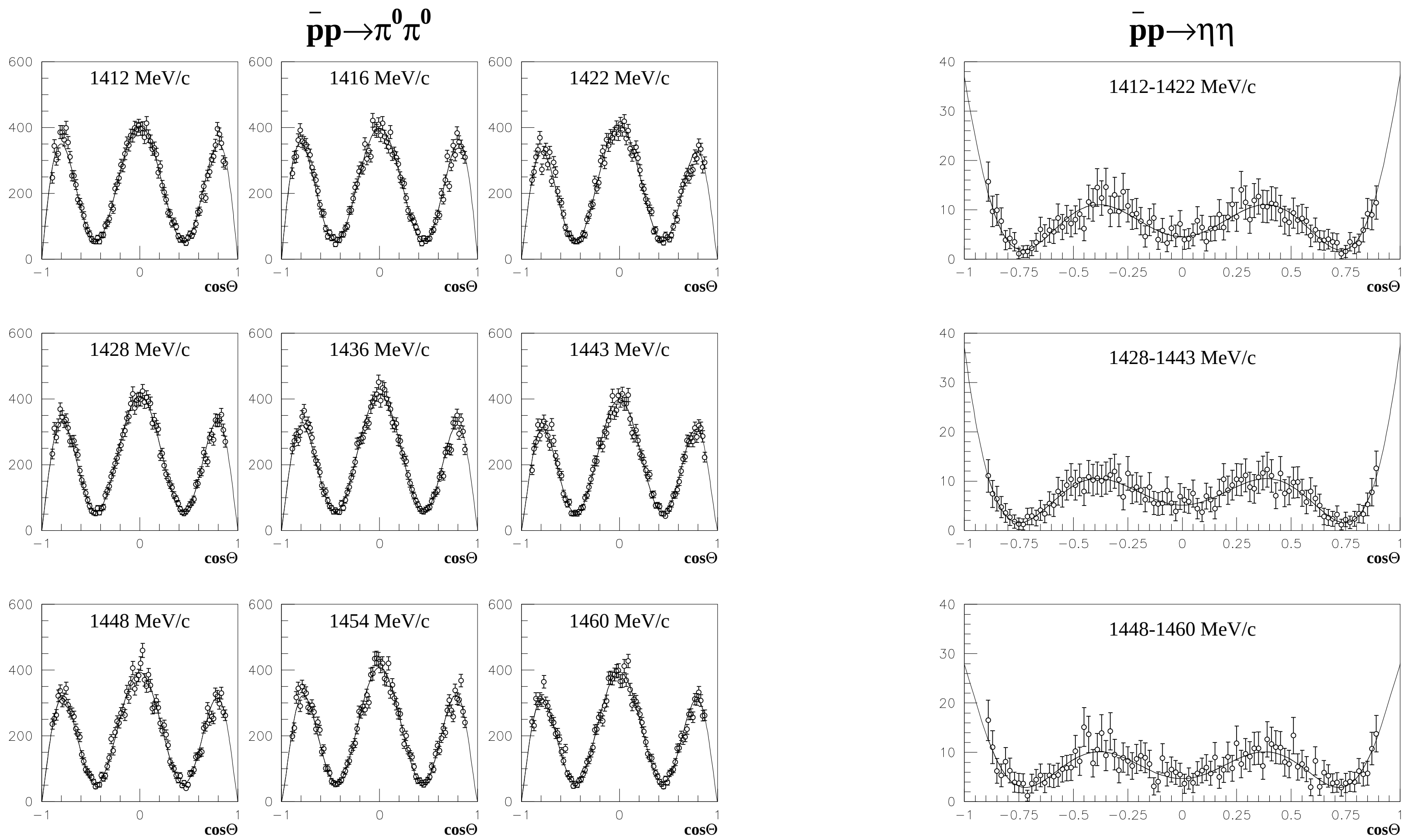

FIG. 6. Measured angular distributions for (left) $p\bar{p} \to \pi^0\pi^0$ and (right) $p\bar{p} \to \eta\eta$. The momenta from 1412 to 1460 MeV/c correspond to energies $\sqrt{s} = 2222$ to 2239 MeV.

Let me describe the results of this channel in some detail, both because I was personally involved in it, and because its rather definitive results have not yet been published.[40] Essentially similar results have been obtained by JETSET.[41]

In the Crystal Barrel experiment the mass region 2222.7 to 2239.7 MeV was scanned in nine nearly equal steps of 2 MeV each. The mass resolution was $\pm$ 0.25 MeV. Between 0.5 and 1.2 million all-neutral triggers were taken at each point, and 6,000-11,000 $\pi^0\pi^0$, and 240-500 $\eta\eta$ at each energy point were identified, with essentially zero background. The results are displayed in Fig. 5, where the yield of the $\eta\pi^0$, which being an isospin 1 channel can not have a glueball, is also shown. No evidence, whatsoever, is seen in any channel for any narrow enhancement in the entire region of our scan. This allows rather strict limits to be placed on the product branching ratios. Our preliminary results for $\Gamma(\xi) \geq 10$ MeV are:

$B(p\bar{p} \rightarrow \xi) \times B(\xi \rightarrow \pi^0\pi^0) \leq 1.2 \times 10^{-4}$,95% CL

$B(p\bar{p} \rightarrow \xi) \times B(\xi \rightarrow \eta\eta) \leq 1.2 \times 10^{-4}$,95% CL

When these results are combined with the BES results for $p\bar{p}$, $\pi^0\pi^0$, and $\eta\eta$ decays, one obtains,

$B(J/\psi \rightarrow \gamma\xi) \geq 2.4 \times 10^{-3}$,95% CL,

which is a very large radiative branch, indeed. One also obtains,

$B(\xi \rightarrow p\bar{p}) \leq 0.63\%$

$B(\xi \rightarrow \pi^0\pi^0) \leq 1.9\%$

and $B(\xi \rightarrow \eta\eta)/B(\xi \rightarrow \pi^0\pi^0) \approx 1$

In addition to the excitation functions, we have also made accurate measurements of the angular distributions for $p\bar{p} \rightarrow \pi^0\pi^0$ and $p\bar{p} \rightarrow \eta\eta$, in the hope that they might contain some tell-tale signs of the presence of $\xi(2230)$. These are displayed in Fig. 6. A Legendre analysis of the angular distributions shows completely smooth evolution of the a_0, a_2, a_4, a_6, a_8 coefficients, confirming that nothing unusual is going on.

In view of these results, I am forced to conclude that none of the two-body decays of $p\bar{p}, (\rightarrow \phi\phi, K_SK_S, \pi^0\pi^0, \eta\eta)$ show any evidence for the existence of the narrow $\xi(2230)$. Either it does not exist, or if it does, its coupling to $p\bar{p}$ is much smaller than suggested by the BES measurements. As a matter of fact, a simple consideration of the Feynman diagrams for the decay of a $|gg\rangle$ state into two pseudoscalar mesons versus its decay into two baryons shows that since the latter must involve the emission of another gluon, decay into two baryons should be an order of magnitude less probable. This is indeed found to be true in $\chi_0 \rightarrow p\bar{p}$ decay.

One more sobering conclusion, which should be drawn from these results for $\xi(2230)$, is that its <u>absence</u> in $\gamma\gamma \rightarrow \pi^+\pi^-$ and K_SK_S[42][43] experiments is no indication of its being a glueball. Most likely, it just does not exist.

The overall summary of my presentation can be stated as following. Serious searches for glueballs are in high gear, but firm conclusions still elude us. The oldest claimant to being the scalar glueball, the θ or $f_0(1710)$, is very much alive, but even the existence of the oldest claimant to being the tensor glueball, the $\xi(2230)$, is in serious question.

I wish to thank David Joffe for the preparation of this manuscript.

REFERENCES

1. J. J. Sakurai, Ann. Phys. (N.Y.) **11** (1960); M. Gell-Mann, Phys. Rev. **125** (1962) 1067.
2. E. Fermi and C. N. Yang, Phys. Rev. **76** (1949) 1739.
3. H. Fritzsch and M. Gell-Mann, Proc. XVI Int. Conf. on High Energy Physics, Fermilab, 1972, vol. 2, p. 135.
4. H. Fritzsch and P. Minkowski, Nuovo Cimento, **30A** (1975) 393.
5. Particle Data Group, 'Review of Particle Properties', [PDG98], Euro. Phys. J. **3** (1998) 1.
6. Barnes, T., Ph. D. dissertation, Caltech (1977), *unpublished*; Barnes, T., F. E. Close and F. de Viron, Nucl. Phys. **B224** (1983) 241; Chanowitz, M. and S. Sharpe, Phys. Lett. **B132** (1983) 413; DeTar, C. and J. Donoghue, Ann. Rev. Nucl. Part. Sci. **33** (1983) 325; Hasenfratz, P., R. R. Horgan, J. Kuti, and J. M. Richard, Phys. Lett. **94B** (1980) 401; C. E. Carlson, T. H. Hansson, and C. Peterson, Phys. Rev. **D27** (1983) 1556.
7. Horn, D. and J. Mandula, Phys. Rev. **D17** (1978) 898; T. Barnes, Z. Phys. **C10** (1981) 275; A. Szczepaniak *et al.*, Phys. Rev. Lett. **76** (1996) 2011.
8. Isgur, N. and J. Paton, Phys. Lett. **B124** (1983) 247; Isgur, N. and J. Paton, Phys. Rev. **D31** (1985) 2910; Isgur, N., R. Kokoski, and J. Paton, Phys. Rev. Lett. **54** (1985) 869.
9. Latorre, J. L., *et al.*, Phys. Lett. **B147** (1984) 169; M. Shifman, Z. Phys. **C9** (1981) 347; P. Pascual and R. Tarrach, Phys. Lett. **113B** (1982) 495; S. Narison, Z. Phys. **C26** (1984) 209; C. A. Dominguez and N. Power, Z. Phys. **C31** (1988) 591; J. Bordes, V. Gimenez, and J. A. Penarrocha, Phys. Lett. **B223** (1989) 251; S. Narison, Nucl. Phys. **B509** (1998) 312.
10. R. Burkhalter, Nucl. Phys. B (Proc. Suppl.) **73** (1999) 3.
11. C. Morningstar and Peardon, Nucl. Phys. B. (Proc. Suppl.) **73** (1999) 927.
12. R. D. Kenway, Nucl. Phys. B. (Proc. Suppl.) **73** (1999) 16.
13. IBMQCD, H. Chen *et al.* Nucl. Phys. B. (Proc. Suppl.) **34** (1994) 357.
14. IBMQCD, W. Lee and D. Weingarten, Nucl. Phys. B (Proc. Suppl.) **73** (1999) 249; also to be published in Phys. Rev. D.
15. UKQCD, G. S. Bali, *et al.* Phys. Lett. **B389** (1993) 378.
16. UKQCD, C. Michael, Proc. Hadron 97, AIP Conference Proceedings **432** (1998) 657.
17. SDQCD, C. Morningstar and M. Peardon, Nucl. Phys. B. (Proc. Suppl.) **53** (1997) 917.
18. C. Morningstar and M. Peardon, hep-lat/9901004 - v2 (1999).
19. F. E. Close and A. Kirk, Phys. Lett. **B397** (1997) 333.
20. A. V. Singovsky (WA102 Collaboration), Proc. ICHEP 97, Vancouver, p. 1263.
21. T. Armstrong *et al.* (E760 Collaboration), Phys. Lett. **B307** (1993) 399; C. Amsler, Revs. Mod. Phys. **70** (1998) 1293.
22. D. Barberis *et al.* (WA102 Collaboration), Phys. Lett. **B453** (1999) 305.
23. E. Klempt *et al.*, Phys. Lett. **B361** (1995) 160; N. A. Tornqvist and M. Roos, Phys. Rev. Lett. **76** (1996) 1575, **77** (1996) 2333.
24. A. Anisovich, V. V. Anisovich, and A. Sarantsev, Phys. Lett. **B95** (1997) 123.

25. P. Minkowski and W. Ochs, Proc. Workshop on Hadron Spectroscopy, Prascati, 1999, *to be published*, hep-ph/9905250.
26. BES Collaboration: J. Z. Bai *et al.*, Phys. Rev. **D60** (1999) 072001; Y. F. Gu and S. F. Tuan, elsewhere in these proceedings.
27. E835 Collaboration, M. Ambrogiani *et al.*, Phys. Rev. Lett. **83** (1999) 2902.
28. C. Amsler and F. Close, Phys. Rev. **D53** (1996) 295.
29. F. Close, G.R. Farrar, and Z. Li, Phys. Rev. **D55** (1997) 5749; F. Close, Nucl. Phys. B (Proc. Suppl.) **63 A - C** (1998) 28.
30. D. Toussaint, Lattice 99, to be published in Nucl. Phys. B (Proc. Suppl.).
31. R. M. Baltrusaitis *et al.*, Mark III Collaboration, Phys. Rev. Lett. **56** (1986) 107.
32. D. Alde *et al.*, IHEP-IISN-LANL-LAPP Collaboration, Phys. Lett **177B** (1986) 120.
33. J. E. Augustin *et al.*, DM2 Collaboration, Phys. Rev. Lett. **60** (1988) 2238.
34. J. Sculli *et al.*, Phys. Rev. Lett. **58** (1987) 1715; G. Bardin *et al.*, Phys. Lett. **B195** (1987) 292; P. D. Barnes *et al.*, Phys. Lett. **B309** (1993) 469.
35. J. Z. Bai *et al.*, BES Collaboration, Phys. Rev. Lett. **76** (1996) 3502.
36. J. Z. Bai *et al.*, BES Collaboration, Phys. Rev. Lett. **81** (1998) 1179.
37. C. Evangelista *et al.*, Phys. Rev. **D56** (1997) 3803.
38. C. Evangelista *et al.*, Phys. Rev. **D57** (1998) 5370.
39. X. Shen, Proc. Hadron 97, edited by S. U. Chung and H. Willutzki, AIP Conf. Report **432** (1998) 47.
40. For preliminary reports of the Crystal Barrel results see: Kamal K. Seth, Proc. 32nd Rencontres de Moriond ('97 QCD), edited by J. Tran Thanh Van, Editions Frontieres (Paris) 1997, pp. 599-606; W. Roethel, Proc. LEAP '98, edited by C. Cicalo *et al.*, Nucl. Phys. **A 655** (1999) 137c.
41. J. B. Ritter, Ph. D. dissertation, Univ. of Illinois, Urbana, 1998, unpublished.
42. R. Godang *et al.*, CLEO Collaboration, Phys. Rev. Lett. **79** (1997) 3829; M. S. Alam *et al.*, CLEO Collaboration, Phys. Rev. Lett. **81** (1998) 3328.
43. A. Favara, Proc. ICHEP '98, edited by A. Astbury *et al.*, World Scientific (1999) 1282.