September 18, 2006 LBNL-61631

Hunting the Scalar Glueball: Prospects for BES III

Michael S. Chanowitz[1]

Theoretical Physics Group Ernest Orlando Lawrence Berkeley National Laboratory
University of California
Berkeley, California 94720

Presented at CHARM 2006,
International Workshop on Tau-Charm Physics,
June 5 - 7, 2006, Beijing, China
To be published in the proceedings.

Abstract

The search for the ground state scalar glueball G_0 is reviewed. Spin zero glueballs will have unique dynamical properties if the $< G_0|\overline{q}q >$ amplitude is suppressed by chiral symmetry, as it is to all orders in perturbation theory: for instance, mixing of G_0 with $\overline{q}q$ mesons would be suppressed, radiative Ψ decay would be a filter for new physics in the spin zero channel, and the decay $G_0 \to \overline{K}K$ could be enhanced relative to $G_0 \to \pi\pi$. These properties are consistent with the identification of $f_0(1710)$ as the largely unmixed ground state scalar glueball, while recent BES data implies that $f_0(1500)$ does not contain the dominant glueball admixture. Three hypotheses are discussed: that G_0 is 1) predominantly $f_0(1500)$ or 2) predominantly $f_0(1710)$ or 3) is strongly mixed between $f_0(1500)$ and $f_0(1710)$.

[1]Email: chanowitz@lbl.gov

1 Introduction

Glueballs are a dramatic consequence of the local, unbroken, non-Abelian symmetry that is the unique defining property of QCD. Non-Abelian gauge bosons (gluons) carry the non-Abelian charge and therefore interact directly with one another. Because the symmetry is unbroken, charge is confined and singlet combinations of two or more gauge bosons form bound states. In QED the Abelian gauge boson, the photon, has no electric charge, the force does not confine, and there is no "lightball" counterpart of the QCD glueball.

The prediction that glueballs exist is simple and fundamental but has proven difficult to verify. We expect their discovery soon, for two reasons. First, BES III will provide huge J/Ψ data samples — potentially several billion — allowing definitive studies of J/Ψ decay and, especially, partial wave analysis of the glueball-preferred radiative J/Ψ decay channel. Second, in roughly the same time frame, lattice QCD (LQCD) will provide reliable unquenched predictions for the glueball spectrum, mixing, and decays. This powerful combination of theory and experiment should suffice to finally resolve this fundamental and difficult problem.

Quenched LQCD calculations have verified the naive expectation that glueballs exist. The most recent quenched results[1] put the ground state scalar mass at $m_G = 1710 \pm 50 \pm 80$ MeV. Glueballs are hard to identify because they are not easily distinguished from ordinary $\overline{q}q$ mesons with which they can mix, and because dynamical properties, such as decay widths and branching ratios, are not understood. The problem is further complicated by the likely presence of $\overline{q}qg$ hybrids and possibly also $\overline{q}\overline{q}qq$ states, with which they may also be confused and mix.

For now we rely on a few simple ideas:

- Glueballs are extra states, beyond the $\overline{q}q$ spectrum. To exploit this we must understand the "ordinary" $\overline{q}q$ spectrum very well, using data from Ψ, B, and Z decays, and from $\overline{p}p$, πp, $\gamma\gamma$, and γN scattering. It is already clear that there are indeed "extra" $I, J^{PC} = 00^{++}$ states in the mass region where the scalar glueball is expected.

- Glueballs couple strongly to gluons so they are prominent in radiative Ψ decay, which proceeds via $\Psi \to \gamma gg$. They couple weakly to photons so they are not prominent in photon-photon scattering.[2]

- Glueballs are flavor singlets so their decays should be $SU(3)_F$ symmetric. However, this may not be true of spin zero glueball decays because of chiral suppression, as discussed below.

[2] While neither generic nor physically motivated, it is possible to arrange singlet-octet flavor mixing so that a $\overline{q}q$ meson also has a small or vanishing $\gamma\gamma$ coupling.

Since hybrid and $\overline{qq}qq$ states are also "extra" states and some are flavor singlets, the only distinguishing property unique to glueballs is their strong coupling to the color singlet digluon channel. Radiative Ψ decay then plays a very special role because for heavy quarks Q we know reliably from perturbation theory that the leading mechanism is $\Gamma(\Psi(\overline{Q}Q) \to \gamma X) \simeq \Gamma(\Psi(\overline{Q}Q) \to \gamma gg)$, with[2]

$$\frac{\Gamma(\Psi \to \gamma gg)}{\Gamma(\Psi \to ggg)} = \frac{16\alpha}{5\alpha_S} \simeq 0.09. \tag{1}$$

Using $B(\Psi \to ggg) \simeq B(\Psi \to \text{hadrons})_{\text{direct}} = 0.71$, we obtain $B(\Psi \to \gamma X) \simeq 0.06$. This is consistent with the only attempted inclusive measurement,[3] verifying that the perturbative mechanism is at least roughly correct.[3] The leading partial waves of the digluon in perturbation theory are $J^{PC} = 0^{++}, 0^{-+}, 2^{++}$,[4] corresponding precisely to the quantum numbers of the lightest glueballs. Radiative Ψ decay is then a copious source of photon-tagged, color-singlet gluon pairs, perfectly matched to the expected masses and quantum numbers of the lightest glueballs.

Heavier quarkonia cannot compete: e.g., for equal luminosity, the number of events in the glueball mass region for Υ radiative decay is smaller by a factor $\simeq 10^2 \times 4 \times 10 = 4000$, where 10^2 reflects the observed peak cross sections, 4 is from the square of the quark charges, and the final 10 is from the branching ratio into the relevant digluon mass region. Radiative Ψ decay is the ideal glueball hunting ground, for which BEPC II/BES III will be the premier world facility.

Glueballs are sticky because they couple strongly to gluons and weakly to photons. The stickiness of particle X is defined as[5]

$$S_X = \frac{\Gamma(\Psi \to \gamma X)}{\Gamma(X \to \gamma\gamma)} \times \frac{PS(X \to \gamma\gamma)}{PS(\Psi \to \gamma X)}, \tag{2}$$

where PS denotes phase space. We consider stickiness ratios, since glueballs will typically be much stickier than $\overline{q}q$ mesons, $\overline{q}qg$ hybrids, or $\overline{qq}qq$ states. It is worth considering if the high luminosity at BEPC will make it feasible to study $\gamma\gamma$ scattering at BES III despite the low beam energy.

2 Chiral Suppression

If chiral symmetry breaking in glueball decay is dominated by quark masses, then the coupling of a spin zero glueball to light $\overline{q}q$ pairs is chirally suppressed,[6]

$$< G_0|\overline{q}q > \propto m_q/m_G, \tag{3}$$

[3]Although the shape is distorted by resonances, the measured rate for photons with $\geq 60\%$ of the beam energy is consistent with QCD,[3] as expected for "global duality."

like the suppression of $\pi \to e\nu$, though different in detail. This is easily understood: for $m_q = 0$ chiral symmetry requires the quark and antiquark to have equal chirality, hence unequal helicity, implying nonvanishing net angular momentum, so that the amplitude must vanish for $J = 0$ in the chiral limit. Chiral suppression, eq. (3), is valid to all orders in perturbation theory.[6, 7] Explicitly at leading order[6]

$$\mathcal{M}(G_0 \to \overline{q}q) = -f_0\alpha_S\ \frac{16\pi\sqrt{2}}{3}\ \frac{m_q}{\beta}\ \log\frac{1+\beta}{1-\beta}\ \overline{u}_3 v_4 \delta_{ij}. \tag{4}$$

where f_0 is the effective G_0gg coupling and β is the quark velocity in the G_0 cms.

However, there is no limit in which (4) is a reliable estimate of the magnitude. Even for $m_G \to \infty$ the t and u channel quark exchange amplitudes are not under perturbative control. We cannot calculate the magnitude of the amplitude but we know it is suppressed of order m_q/m_G to all orders in perturbation theory.

Nonperturbative chiral symmetry breaking might lift the chiral suppression, as suggested[8] in the context of the liquid instanton model. A reliable, model-independent, nonperturbative method is needed to decide: for now LQCD is the only game in town. Early results are equivocal, as discussed below. The phenomenological proposal that chiral symmetry is restored[9] in the baryon and meson spectra for $\gtrsim$O(2) GeV suggests that nonperturbative chiral symmetry breaking is not large at the glueball mass scale, and in fact motivated the suggestion that a "high lying" scalar glueball would not mix strongly with $\overline{u}u+\overline{d}d$ mesons.[10]

Chiral suppression has important consequences for spin zero glueballs. Mixing with light (u, d, s) mesons is suppressed of order $\mathrm{O(m_q/m_G)}$, so that $J = 0$ glueballs are more likely than $J \neq 0$ to be largely unmixed. (Although mixing amplitudes are suppressed, mixing angles can be large if the quenched glueball and meson states happen to be extremely degenerate.) To the extent $G_0 - M_0(\overline{q}q)$ mixing does occur, it should be dominated by $\overline{s}s$ components. Mixing with hybrids and four-quark states is not suppressed.

A second consequence is that radiative Ψ decay becomes a filter for new physics in the $J = 0$ channel, since at leading order the exclusive amplitude $\Psi \to \gamma X$ is proportional to $< gg|X >$, so that radiative decays to spin zero light quark mesons, $X = M_0(\overline{q}q)$, are suppressed, and, to the extent they do occur, favor $J = 0$ strangeonium, $M_0(\overline{s}s)$, over $M_0(\overline{u}u + \overline{d}d)$. Radiative decays to $J = 0$ glueballs, hybrid, and four quark states are not suppressed.

A third consequence is that $\overline{q}q$ decays of $J = 0$ glueballs favor the heaviest quark q. If the multibody decays have discernible jet structure, decays to two jets will contain leading strange particles if $m_G < 2m_D$ or two charm particles if $2m_D < m_G < 2m_B$, while the leading particles of three jet decays are flavor symmetric. For $m_G < 2m_D$ we could then see an increase in leading strange particles in events with high thrust. For $m_G \simeq 1700$

MeV the partonic preference for $\overline{s}s$ over $\overline{u}u + \overline{d}d$ decays would favor $\overline{K}K$ favored over $\pi\pi$ if hadronization of $G_0 \to \overline{q}q$ is an important short distance mechanism for two meson decays at this mass scale. Another possibility is that the dominant short distance mechanism is $G_0 \to \overline{q}\overline{q}qq$, which would imply that $\overline{K}K$ and $\pi\pi$ are more nearly equal.[7] Or both mechanisms could be important and the ratio could lie between the two predictions.

The existing evidence from LQCD is preliminary and equivocal. A quenched study[11] of scalar glueball decay to two pseudoscalar mesons found that the amplitude decreases with the meson mass m_P at a rate consistent with the m_P^2 dependence expected for chiral suppression (since $m_q \propto m_P^2$). A study of $G_0 - \overline{s}s$ mixing[12] found a small mixing energy, $E_M = 43 \pm 31$ MeV, also as expected, but was not consistent with $E_M \propto m_q$. Given the small $\sim 1\sigma$ "signal," this calculation may have lacked the precision needed to obtain the m_q dependence. Another study found large mixing at the strange quark mass but the lattice granularity was far from the continuum limit.[13] All these studies extrapolated from quark masses near or above the strange quark mass and did not directly probe the region of the up and down quark masses. They should be revisited with today's computing power, to simultaneously explore the chiral and continuum limits. A quenched calculation of mixing would suffice to determine whether chiral suppression occurs or not.

3 Experimental Status of the Scalars

The most recent quenched LQCD calculation obtained $m_G = 1710 \pm 50 \pm 80$ MeV for the scalar glueball mass.[1] Experimentally there are too many isoscalar, scalar mesons between 1.4 and 2 GeV to be explained by the naive quark model alone. I assume $f_0(600)$ and $f_0(980)$ are cryptoexotic $\overline{q}\overline{q}qq$ states.[14] The p-wave $\overline{q}q$ scalar nonet is likely to lie in the region of the other spin-triplet p-wave nonets, with isoscalars roughly between ~ 1250 and ~ 1600 MeV. Between ~ 1400 and 2000 MeV there are five $I, J^{PC} = 0, 0^{++}$ states: $f_0(1370)$, $f_0(1500)$, and $f_0(1710)$ are well known, while $f_0(1790)$ and $f_0(1810)$ were recently discovered by BESII. It seems likely that some of these five states have gluon constituents: we have probably seen the scalar glueball although we cannot yet identify it. I will briefly discuss the possibility that the scalar glueball is predominantly 1) $f_0(1500)$,[15] or 2) $f_0(1710)$[11, 12, 6], or 3) is shared by both in a maximally mixed glueball-strangeonium duo.

3.1 $f_0(1500)$ is the glueball

Since $f_0(1500)$ is produced in "gluon rich" $\overline{p}p$ annihilation and πN central production while $f_0(1710)$ decays prominently to $\overline{K}K$, it was natural to consider the hypothesis that $f_0(1500)$ is the scalar glueball and $f_0(1710)$ is the $\overline{s}s$ scalar nonet partner of $f_0(1370)$.[15] However, the dynamics of $\overline{p}p$ annihilation and πN central production are not as well understood as radiative Ψ decay, which we know from perturbation theory is a copious source of color singlet gluon pairs in the relevant mass region. It is then problematic that $f_0(1500)$ is not strongly produced in radiative Ψ decay, in old Mark III data[16] and more recently in BES II data,[17] while the $f_0(1710)$ is prominent in both data sets.[16, 18] In addition, quenched LQCD calculations find that the glueball is $\simeq$ 200 MeV lighter than scalar strangeonium,[12, 13] while the opposite ordering is required by this hypothesis.[19]

The recently reported BES II partial wave analysis of $\Psi \to \gamma\pi\pi$ is an important result.[17] The rates for $\pi^+\pi^-$ and $\pi^0\pi^0$ agree, a critical check since $\gamma\pi^0\pi^0$ is free of the large $\Psi \to \pi^0\rho^0$ background that afflicts $\gamma\pi^+\pi^-$. This is the best channel to search for $\Psi \to \gamma f_0(1500)$ because of the simplicity of the two pion final state and because $B(f_0(1500) \to \pi\pi) = 0.349 \pm 0.0223$ is large.[20] BES II finds only a small possible signal, $B(\Psi \to \gamma f_0(1500)) \times B(f_0(1500) \to \pi^+\pi^-) = (6.7 \pm 2.8) \cdot 10^{-5}$, implying

$$B(\Psi \to \gamma f_0(1500)) = (2.9 \pm 1.2) \cdot 10^{-4}. \tag{5}$$

This is small, viewed either as a fraction of all radiative decays or compared to the rate for $f_0(1710)$, for which the lower limit is six times larger,

$$B(\Psi \to \gamma f_0(1710)) \geq (16.2 + 3.0 - 2.4) \cdot 10^{-4}, \tag{6}$$

from just the $\overline{K}K$ and $\eta\eta$ decay modes, using $B(\Psi \to \gamma f_0(1710)) \times B(f_0(1710) \to \overline{K}K) = (11.1{+}1.7{-}1.2)\cdot 10^{-4}$ from the BES II bin-by-bin fit[18] and $B(f_0(1710) \to \eta\eta)/B(f_0(1710) \to \overline{K}K) = 0.48 \pm 0.15$.[21, 20] The $f_0(1710)$ probably has other decays, especially multibody modes which are difficult to measure, so the inclusive rate for $\Psi \to \gamma f_0(1710)$ is likely to be appreciably larger.[4] The additional statistical power of BES III may be needed to analyze the multibody modes.

An early attempt to study the four pion channel, $\Psi \to \gamma + 4\pi$, was made with the 8M event BES I data sample.[22] Given the complexity of the analysis, including assignment of the four pions to two-isobar intermediates, this would be challenging even with the 58M BES II data set. The BES I results from $f_0(1500) \to 4\pi$ are not consistent with the BES II $f_0(1500) \to 2\pi$ results. From BES I, $B(\Psi \to \gamma + f_0(1500)) \times B(f_0(1500) \to \pi^+\pi^-\pi^+\pi^-) = (3.1 \pm 0.2 \pm 1.1) \cdot 10^{-4}$, which implies $B(\Psi \to \gamma + f_0(1500)) \times B(f_0(1500) \to 4\pi) = (7.0 \pm 2.5) \cdot 10^{-4}$ using isospin (with the decay chain $f_0 \to \sigma\sigma \to 4\pi$ used in the BES I analysis) to include neutral pion modes. Using $B(f_0(1500) \to 4\pi) = 0.495 \pm 0.033$,[20] the BES I measurement implies the inclusive rate $B(\Psi \to \gamma + f_0(1500)) = (14 \pm 5.1) \cdot 10^{-4}$, a factor

[4]The lower limit (6) would increase to $20.2 \cdot 10^{-4}$ if the BES II result for $\Psi \to \gamma f_0(1710) \to \gamma\pi\pi$ were included. I have not included it here because a smaller value for the $\pi\pi$ mode is implied by the BES II 95% upper limit from $\Psi \to \omega\pi\pi$ as reviewed in Section 3.2.

5 larger (and 2σ higher) than eq. (5) from the BES II $f_0(1500) \to 2\pi$ measurement. The $\pi\pi$ measurement must be given greater weight, since it considers a much simpler final state, uses the upgraded BES II detector, and is based on seven times more statistics.

Another problem is posed by hadronic Ψ decay data. The $f_0(1710)$ is produced prominently in $\Psi \to \omega f_0(1710) \to \omega \overline{K}K$[23] but not in $\Psi \to \phi f_0(1710) \to \phi \overline{K}K$,[24] contrary to the OZI rule if $f_0(1710)$ is an $\overline{s}s$ state. But the OZI rule does correctly describe the pattern of the four decays to the ideally mixed tensor mesons, $\Psi \to \omega/\phi + f_2(1270)/f_2(1525)$.[20] If $f_0(1710) = \overline{s}s$, it is necessary to assume that dynamics in the $J = 0$ channel somehow makes the doubly OZI suppressed rate not just comparable to the singly suppressed one but ~ 5 times larger (see Close and Zhao[19]).

3.2 $f_0(1710)$ is the glueball

Another possibility is that $f_0(1710)$ is the scalar glueball.[11, 12, 6] This is consistent with its prominence in radiative Ψ decay, eq. (6), and its mass, in the middle of the range of the most recent quenched LQCD prediction.[1] Chiral suppression could then explain the absence of strong glueball-meson mixing. It is also clearly seen in $\Psi \to \omega f_0(1710) \to \omega \overline{K}K$ with virtually the same mass and width as in $\Psi \to \gamma f_0(1710) \to \gamma \overline{K}K$ but it is not seen in $\Psi \to \omega f_0(1710) \to \omega\pi\pi$ despite the much greater statistics of the $\omega\pi\pi$ channel, yielding a robust 95% CL upper limit, $B(f_0(1710) \to \pi\pi)/B(f_0(1710) \to \overline{K}K) < 0.11$.[23] However, a possible indication of $f_0(1710) \to \pi\pi$ appears in $\Psi \to \gamma\pi\pi$ where BES II finds a scalar at $1765^{+4}_{-3} \pm 12$ MeV with $\Gamma = 145 \pm 8 \pm 69$ MeV. If attributed to $f_0(1710)$ it implies $B(f_0(1710) \to \pi\pi)/B(f_0(1710) \to \overline{K}K) = 0.41^{+0.11}_{-0.17}$, which is 1.8σ above the 95% upper limit from $\Psi \to \omega + \pi\pi/\overline{K}K$. The signal at 1765 could also be due to the $f_0(1790)$ seen in $\Psi \to \phi\pi\pi$[24] or it could be the result of interference between $f_0(1710)$ and $f_0(1790)$. If the stronger upper limit from $\Psi \to \omega + \pi\pi/\overline{K}K$ prevails, chiral suppression could explain the suppression of the $\pi\pi$ mode,[6] which is also consistent with a quenched LQCD calculation.[11] The problem then would be to find the strangeonium component of the scalar $\overline{q}q$ nonet, since neither $f_0(1370)$ nor $f_0(1500)$ seem to have much $\overline{s}s$ content.

A possible solution is suggested by data from charmless B meson decays, $B \to K\overline{K}K$ and $B \to K\pi\pi$.[5] Belle[25] and BABAR[26] both see a strong signal for a scalar meson near 1500 MeV which decays to $\overline{K}K$ but not to $\pi\pi$, although the amplitude analysis of the $\overline{K}K$ channel has significant model dependent ambiguities requiring further study. This object cannot be the previously observed $f_0(1500)$, for which $B(\overline{K}K)/B(\pi\pi) = 0.241 \pm 0.028$,[20] but it could be the missing $\overline{s}s$ scalar. If it is the $\overline{s}s$ scalar and $f_0(1370)$ is its $\overline{u}u + \overline{d}d$ nonet partner, then an explanation is needed for the previously observed $f_0(1500)$.

[5] I thank Alex Bondar for telling me of these results.

3.3 $f_0(1500)$ and $f_0(1710)$ share the glue

Since the amplitude $\mathcal{M}(gg \to \overline{q}q)_{J=0} \propto m_q$ is dominated by forward and backward scattering, the effective running mass m_q must be evaluated at a low energy scale of order $\Lambda_{\rm QCD}$. The effective masses are larger than their "current quark" values but the hierarchy $m_u, m_d \ll m_s$ is maintained. The effective value of m_s might then be large enough that $\mathcal{M}(gg \to \overline{s}s)_{J=0}$ is not chirally suppressed while $\mathcal{M}(gg \to \overline{u}u + \overline{d}d)_{J=0}$ is. The scalar glueball could then mix strongly with strangeonium but not with $\overline{u}u + \overline{d}d$ mesons. The quenched scalar glueball, $G_0(gg)$, and the strangeonium meson, $f_0(\overline{s}s)$, which are expected to have masses near one another, could then mix maximally, yielding the eigenstates

$$f_0^{\pm} \simeq \frac{1}{\sqrt{2}}[G_0(gg) \pm f_0(\overline{s}s)]. \tag{7}$$

Now consider the amplitudes $< gg|f_0^{\pm} >$ and $< \overline{s}s|f_0^{\pm} >$. The first determines the rate for $\Psi \to \gamma f_0^{\pm}$, while I will assume the latter is the dominant partonic mechanism for $f_0^{\pm} \to \overline{K}K$. We choose the phases of the wave functions so that the "elastic" amplitudes, $< gg|G_0(gg) >$ and $< \overline{s}s|f_0(\overline{s}s) >$, are real and positive. If the "inelastic" amplitudes, $< \overline{s}s|G_0(gg) >$ and $< gg|f_0(\overline{s}s) >$, have equal phase and that phase is real relative to the "elastic" amplitudes, the rates for $\Psi \to \gamma f_0^{\pm}$ and $f_0^{\pm} \to \overline{K}K$ would replicate the experimentally observed pattern. One state, say $f_0^{+} \simeq f_0(1710)$, would be produced prominently in radiative Ψ decay and would decay prominently to $\overline{K}K$, because of constructive interference of the gg and $\overline{s}s$ components, while the corresponding $f_0^{-} \simeq f_0(1500)$ amplitudes would be suppressed by destructive interference. This modified chiral suppression scenario could be tested in quenched LQCD studies of mixing between the quenched scalar glueball and the $f_0(\overline{u}u + \overline{d}d)$ and $f_0(\overline{s}s)$ scalar mesons.[6] In this connection it is amusing that chiral symmetry restoration is seen clearly in the spectrum of u, d-quark baryons but not in the strange baryon spectrum.[27]

4 Discussion

BES III at BEPC II will begin operation in 2007. At design luminosity it will accumulate several billion Ψ decays in a single year, enabling definitive partial wave analysis of the decay products. We can look forward to better understanding of the scalar glueball candidates, including their multibody decays. During the BEPC II lifetime LQCD should begin to contribute reliable unquenched calculations of the spectrum, mixing and decays. In particular, LQCD can determine if chiral suppression survives nonperturbative effects and, if so, how it

[6] In this scenario, radiative Ψ decay filters out $J = 0$ $\overline{u}u + \overline{d}d$ mesons but not $\overline{s}s$.

effects mixing and decays. The combination of BES III and LQCD should allow us to finally identify and study the scalar glueball, as well as glueballs and hybrids of other quantum numbers.

Acknowledgments

I wish to thank Alex Bondar, Robert Cahn, and Zoltan Ligeti for discussions.

This work was supported in part by the Director, Office of Science, Office of High Energy and Nuclear Physics, Division of High Energy Physics, of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

References

[1] Y. Chen *et al.*, *Phys. Rev.* **D73**, 014516 (2006).

[2] M. Chanowitz, *Phys. Rev.* **D12**, 918 (1975); L. Okun and M. Voloshin, ITEP 95, 1976 (unpublished).

[3] Mark II Collab. (G. S. Abrams *et al.*), *Phys. Rev. Lett.* **44**, 114 (1980).

[4] A. Billoire *et al.*, *Phys.Lett.* **B80** 381 (1979).

[5] M. Chanowitz, Resonances in Photon Photon Scattering, *Proc. VI Int. Workshop on Photon-Photon Collisions*, Lake Tahoe, CA, Sep 10-13, 1984, ed. R. L. Lander (World Scientific, Singapore, 1985), (QCD161:I405:1984).

[6] M. Chanowitz, *Phys. Rev. Lett.* **95**, 172001 (2005).

[7] K.T. Chao, Xiao-Gang He, J.P. Ma, Dec 2005, e-Print Archive: hep-ph/0512327

[8] Z.F. Zhang, H.Y. Jin, Nov 2005, e-Print Archive: hep-ph/0511252

[9] L. Glozman, *Phys. Lett.***B475**, 329 (2000), e-Print Archive: hep-ph/0301012.

[10] L. Glozman, *Eur.Phys.J.***A19**, 153 (2004), e-Print Archive: hep-ph/0301012.

[11] J. Sexton, A. Vaccarino, and D. Weingarten, *Phys. Rev. Lett.***75**, 4563 (1995).

[12] W. Lee and D. Weingarten, *Phys. Rev.***D61**, 014015 (1999).

[13] C. McNeile and C. Michael, *Phys. Rev.***D63**, 114503 (2000).

[14] R. L. Jaffe, *Phys. Rev.***D15**, 267 (1977).

[15] C. Amsler and F. Close, *Phys. Lett.* **B342**, 433 (1995).

[16] MARK-III Collaboration (W. Dunwoodie for the collaboration), in *Proc. 7th International Conference on Hadron Spectroscopy*, Upton, NY, 25-30 Aug 1997. Published in *AIP Conf.Proc.* **432**, 753 (1998).

[17] BES Collab. (M. Ablikim *et al.*), e-Print Archive: hep-ex/0603048 (2006).

[18] BES Collab. (J. Z. Bai *et al.*), *Phys. Rev.* **D68**, 052003 (2003).

[19] F. Close and Q. Zhao, *Phys. Rev.* **D71**, 094022 (2005); X-G. He *et al.*, *Phys. Rev.* **D73**, 051502 (2006).

[20] Particle Data Group (W.-M. Yao *et al.*), *Journal of Physics* **G 33**, 1 (2006).

[21] WA102 Collab. (D. Barberis *et al.*), *Phys. Lett.* **B479**, 59 (2000).

[22] BES Collab. (J. Z. Bai *et al.*), *Phys. Lett.* **B472**, 207 (2000).

[23] BES Collab. (M. Ablikim *et al.*), *Phys. Lett.* **603**, 138 (2004).

[24] BES Collab. (M. Ablikim *et al.*), *Phys. Lett.* **607**, 243 (2005).

[25] Belle Collab. (A. Garmash *et al.*), *Phys. Rev.* **D71**, 092003 (2005).

[26] BABAR Collab. (B. Aubert *et al.*), *Phys. Rev.* **D74**, 032003 (2006); *Phys. Rev.* **D72**, 072003 (2005).

[27] R. L. Jaffe, D. Pirjol, and A. Scardicchio, e-Print Archive: hep-ph/0602010 (2006).