Scalar and Pseudoscalar Glueballs Revisited

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Abstract. Using two simple and robust inputs to constrain the mixing matrix of the isosinglet scalar mesons \( f_0(1710) \), \( f_0(1500) \), \( f_0(1370) \), we have shown that in the SU(3) symmetry limit, \( f_0(1500) \) becomes a pure SU(3) octet and is degenerate with \( a_0(1450) \), while \( f_0(1370) \) is mainly an SU(3) singlet with a slight mixing with the scalar glueball which is the primary component of \( f_0(1710) \). These features remain essentially unchanged even when SU(3) breaking is taken into account. We have deduced the mass of the pseudoscalar glueball \( G \) from an \( \eta-\eta'G \) mixing formalism based on the anomalous Ward identity for transition matrix elements. With the inputs from the recent KLOE experiment, we found a solution for the pseudoscalar glueball mass around \( 1.4 \pm 0.1 \) GeV. This affirms that \( \eta(1405) \), having a large production rate in the radiative \( J/\psi \) decay and not seen in \( \gamma\gamma \) reactions, is indeed a leading candidate for the pseudoscalar glueball. It is much lower than the results from quenched lattice QCD (> 2.0 GeV).

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INTRODUCTION

The existence of glueballs is a unique prediction of QCD as a confining theory. The latest lattice calculation of the glueball spectroscopy shows that the lightest glueballs are scalar, tensor and pseudoscalar glueballs with masses of order 1710, 2390 and 2560 MeV, respectively [1]. Since the lattice calculation was done in the quenched approximation, the predicted masses are for pure glueballs in the Yang-Mills gauge theory. The question is what happens to the glueballs in the presence of quark degrees of freedom? First, a glueball will mix with the ordinary meson with the same quantum numbers so that a pure glueball does not exist in nature. Since the glueball is hidden somewhere in the quark sector, this is one of main reasons why it is so elusive. Second, we shall show that the mass of the pseudoscalar glueball could be drastically affected by the dynamic fermion effect.

SCALAR GLUEBALL

It is generally believed that the scalar glueball is hidden somewhere in the isosinglet scalar mesons with masses above 1 GeV. The argument goes as follows. Many scalar mesons with masses lower than 2 GeV have been observed and they can be classified into two nonets: one nonet with mass below or close to 1 GeV, such as \( \sigma, \kappa, f_0(980) \) and \( a_0(980) \) that are generally believed to be composed mainly of four quarks and the other nonet with mass above 1 GeV such as \( K^*_0(1430), a_0(1450) \) and two isosinglet scalar mesons. This means that not all three isosinglet scalars \( f_0(1710), f_0(1500), f_0(1370) \) can be accommodated in the \( q\bar{q} \) nonet picture. One of them should be primarily a scalar glueball [4].

Among the three isoscalar mesons, it has been quite controversial as to which of these is the dominant scalar glueball. It was suggested that \( f_0(1500) \) is primarily a scalar glueball in [4], due partly to the fact that \( f_0(1500) \), discovered in \( p\bar{p} \) annihilation at LEAR, has decays to \( \eta\eta \) and \( \eta\eta' \) which are relatively large compared to that of \( \pi\pi \) and that the earlier quenched lattice calculations predict the scalar glueball mass to be around 1550 MeV [5]. Furthermore, because of the small production of \( \pi\pi \) in \( f_0(1710) \) decay compared to that of \( KK \), it has been thought that \( f_0(1710) \) is primarily \( s\bar{s} \) dominated. In contrast, the smaller production rate of \( KK \) relative to \( \pi\pi \) in \( f_0(1370) \) decay leads to the conjecture

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1 It has been suggested that \( f_0(600) \) or the \( \sigma \) state is a good candidate for the scalar glueball (see e.g. [3]). Since a pure glueball state (scalar or pseudoscalar) cannot decay into a photon pair or a massless quark pair to the leading order, the identification of \( f_0(600) \) with a scalar glueball seems to be unlikely in view of its broad width, of order \( 600 \sim 1000 \) MeV, and its sizable partial width to \( \gamma\gamma \) of order \( 4\text{keV} \) [6]. In the conventional tetraquark picture of \( f_0(600) \), its broad width can be naturally understood as fall-apart decays.
that \( f_0(1370) \) is governed by the non-strange light quark content. Based on the above observations, a flavor-mixing scheme was proposed \([4]\) to consider the glueball and \( q\bar{q} \) mixing in the neutral scalar mesons \( f_0(1710), f_0(1500) \) and \( f_0(1370) \). Best \( \chi^2 \) fits to the measured scalar meson masses and their branching fractions of strong decays have been performed in several references by Amsler, Close and Kirk \([4]\), Close and Zhao \([6]\), and He et al. \([7]\). A typical mixing matrix in this scheme is \([6]\)

\[
\begin{pmatrix}
(f_0(1370)) \\
(f_0(1500)) \\
(f_0(1710))
\end{pmatrix} =
\begin{pmatrix}
-0.91 & -0.07 & 0.40 \\
-0.41 & 0.35 & -0.84 \\
0.09 & 0.93 & 0.36
\end{pmatrix}
\begin{pmatrix}
|N| \\
|S| \\
|G|
\end{pmatrix}.
\]

A common feature of these analyses is that, before mixing, the \( s\bar{s} \) quarkonium mass \( M_S \) is larger than the glueball mass \( M_G \) which, in turn, is larger than the \( N(\approx (u\bar{u} + d\bar{d})/\sqrt{2}) \) mass \( M_N \), with \( M_G \) close to 1500 MeV and \( M_S - M_N \) of the order of 200 ~ 300 MeV.

However, the above scenario encounters several difficulties: (i) The isovector scalar meson \( a_0(1450) \) is now confirmed to be the \( q\bar{q} \) meson in the lattice calculation \([8]\). As such, the degeneracy of \( a_0(1450) \) and \( K_0^0(1430) \), which has a strange quark, cannot be explained if \( M_S \) is larger than \( M_G \) by \( \sim 250 \) MeV. (ii) The most recent quenched lattice calculation with improved action and lattice spacings extrapolated to the continuum favors a scalar glueball mass close to 1700 MeV \([1] \). (iii) If \( f_0(1710) \) is dominated by the \( s\bar{s} \) content, the decay \( J/\psi \to \phi f_0(1710) \) is expected to have a rate larger than that of \( J/\psi \to \omega f_0(1710) \). Experimentally, it is otherwise. This is true for the \( \Gamma(J/\psi \to \gamma f_0(1500)) \) decay rate to be substantially larger than that of \( \Gamma(J/\psi \to \gamma f_0(1710)) \). Again, experimentally, the opposite is true.

In our work \([9]\), we employed two simple and robust results as the input for the mass matrix which is essentially the starting point for the mixing model between scalar mesons and the glueball. First of all, we know empirically that flavor SU(3) is an approximate symmetry in the scalar meson sector above 1 GeV. The near degeneracy of \( K_0^0(1430), a_0(1470), \) and \( f_0(1500) \) has been observed. In the scalar charmed meson sector, \( D_{s0}^*(2317) \) and \( D_{s0}^*(2308) \) have similar masses even though the former contains a strange quark. It is most likely that the same phenomenon also holds in the scalar bottom meson sector. This unusual behavior is not understood as far as we know and it serves as a challenge to the existing hadronic models. Second, an improved quenched lattice calculation of the glueball spectrum at the infinite volume and continuum limits based on much larger and finer lattices have been carried out \([1] \). The mass of the scalar glueball is calculated to be \( m(0^{++}) = 1710 \pm 50 \pm 80 \) MeV. This suggests that \( M_G \) should be close to 1700 MeV rather than 1500 MeV from the earlier lattice calculations \([5]\).

We begin by considering exact SU(3) symmetry as a first approximation for the mass matrix. In this case, two of the mass eigenstates are identified with \( a_0(1450) \) and \( f_0(1500) \) which are degenerate with the mass \( M \). Taking \( M \) to be the experimental mass of 1474 \( \pm \) 19 MeV \([3]\), it is a good approximation for the mass of \( f_0(1500) \) at 1507 \( \pm \) 5 MeV \([3]\). Thus, in the limit of exact SU(3) symmetry, \( f_0(1500) \) is an SU(3) isosinglet octet state and is degenerate with \( a_0(1450) \). In the absence of glueball-quarkonium mixing, \( f_0(1370) \) becomes a pure SU(3) singlet and \( f_0(1710) \) the pure glueball. When the glueball-quarkonium mixing is turned on, there will be some mixing between the glueball and the SU(3)-singlet \( q\bar{q} \). The mass shift of \( f_0(1370) \) and \( f_0(1710) \) due to mixing is of order 10 MeV. Since the SU(3) breaking effect is expected to be weak, it can be treated perturbatively.

**Chiral suppression**

If \( f_0(1710) \) is primarily a pseudoscalar glueball, it is naively expected that \( \Gamma(G \to \pi\pi)/\Gamma(G \to K\bar{K}) \approx 0.9 \) after phase space correction due to the flavor independent coupling of \( G \) to \( PP \). However, experimentally there is a relatively large suppression of \( \pi\pi \) production relative to \( K\bar{K} \) in \( f_0(1710) \) decay, \( R(f_0(1710)) \equiv \Gamma(f_0(1710) \to \pi\pi)/\Gamma(f_0(1710) \to K\bar{K}) = 0.41 \pm 0.11 \) \([10]\) or even smaller. To explain the large disparity between \( \pi\pi \) and \( K\bar{K} \) production in scalar glueball decays, it was noticed long time ago by Carlson et al. \([11]\), by Cornwall and Soni \([12]\) and revitalized recently by Chanowitz \([13]\) that a pure scalar or pseudoscalar glueball cannot decay into a quark-antiquark pair in the chiral limit, i.e., \( A(G \to q\bar{q}) \approx m_q \). Since the current strange quark mass is an order of magnitude larger than \( m_u \) and \( m_d \), decay to \( K\bar{K} \) is largely favored over \( \pi\pi \). However, chiral suppression for the ratio \( \Gamma(G \to \pi\pi)/\Gamma(G \to K\bar{K}) \) at the hadron level should not be so strong as the current quark mass ratio \( m_u/m_s \). It has been suggested \([14]\) that \( m_q \) should be interpreted as the scale of chiral symmetry breaking.

Whether or not \( G \to \pi\pi \) is subject to chiral suppression is a controversial issue because of the hadronization process from \( G \to q\bar{q} \) to \( G \to \pi\pi \) and the possible competing \( G \to q\bar{q}q\bar{q} \) mechanism \([11, 14, 15, 16]\). The only
where $f$ basis have been proposed. For the latter, the include the pseudoscalar glueball $R$ measured masses and branching fractions. The mixing matrix obtained in our model has the form:

$$ f \begin{pmatrix} f_0(1370) \\ f_0(1500) \\ f_0(1710) \end{pmatrix} = \begin{pmatrix} 0.78 & 0.51 & -0.36 \\ -0.54 & 0.84 & 0.03 \\ 0.32 & 0.18 & 0.93 \end{pmatrix} \begin{pmatrix} |\eta\rangle \\ |\eta'\rangle \\ |G\rangle \end{pmatrix}. \tag{2} $$

It is evident that $f_0(1710)$ is composed primarily of the scalar glueball, $f_0(1500)$ is close to an SU(3) octet, and $f_0(1370)$ consists of an approximate SU(3) singlet with some glueball component ($\sim 10\%$). Unlike $f_0(1370)$, the glueball content of $f_0(1500)$ is very tiny because an SU(3) octet does not mix with the scalar glueball. Because the $\bar{s}s$ content is more copious in $f_0(1710)$, it is natural that $J/\psi \rightarrow \omega f_0(1710)$ has a rate larger than $J/\psi \rightarrow \phi f_0(1710)$. Our prediction of $\Gamma(J/\psi \rightarrow \omega f_0(1710))/\Gamma(J/\psi \rightarrow \phi f_0(1710)) = 4.1$ is consistent with the observed value of $3.3 \pm 1.3 [3]$. Moreover, if $f_0(1710)$ is composed mainly of the scalar glueball, it will be expected that $\Gamma(J/\psi \rightarrow \eta f_0(1710))/\Gamma(J/\psi \rightarrow \eta f_0(1500))$, a relation borne out by experiment. It is interesting to compare the mixing matrices [1] and [2]. In the model of Close and Zhao [6], $f_0(1710)$ is dominated by the $\bar{s}s$ quarkonium in order to explain the suppression of $R(f_0(1710))$. In our scheme, $f_0(1710)$ has the smallest content of $\bar{s}s$ and the smallestness of $R(f_0(1710))$ arises from the chiral suppression of scalar glueball decay. Although $f_0(1500)$ in our model has the largest content of $\bar{s}s$, the $K\bar{K}$ production is largely suppressed relative to $\pi\pi$. $R(f_0(1500)) = 3[(\alpha/\beta)^2(p_\pi/p_K)$, where $f_0(1500) = \alpha(|u\bar{u}| + |d\bar{d}|) + \beta|s\bar{s}|$ and $p_h$ is the c.m. momentum of the hadron $h$. In SU(3) limit, $\beta = -2\alpha$ and this leads to $R(f_0(1500)) = 3.9$, in agreement with experiment.

**PSEUDOSCALAR GLUEBALL**

In 1980, Mark II observed a resonance with a mass around 1440 MeV in the radiative $J/\psi$ decay [18] which was subsequently named $\Upsilon(1440)$ by Mark II and Crystal Ball Collaborations [19]. Shortly after the Mark II experiment, $\Upsilon(1440)$ now known as $\eta(1405)$ was a leading candidate for the pseudoscalar glueball. (For an excellent review of the $\Upsilon(1420)$ and $\Upsilon(1440)$ mesons, see [20].) Indeed $\eta(1405)$ behaves like a glueball in its productions and decays because it has a large production rate in the radiative $J/\psi$ decay and is not seen in $\gamma\gamma$ reactions. Besides $\eta(1405)$, other states with masses below 2 GeV have also been proposed as the candidates, such as $\eta(1760)$ and $\Upsilon(1835)$.

However, the pseudoscalar glueball interpretation for $\eta(1405)$ is not favored by most of the theoretical calculations. For example, quenched lattice gauge calculations predict the mass of the $0^{-}$ state to be above 2 GeV in [21] and around 2.6 GeV in [22, 1]. It is not favored by the sum-rule analysis with predictions higher than 1.8 GeV [23, 24] either (for a review of the glueball spectroscopy in various approaches, see [25]). Hence, we are encountering an embarrassing situation that although experimentally $\eta(1405)$ is a favored candidate for the pseudoscalar glueball, theorists seem to prefer to have a $0^{-}$ state heavier than the scalar glueball. The motivation of our recent work [26] is to see if we can learn something about the glueball mass by studying the $\eta - \eta'$ mixing.

The $\eta - \eta'$ mixing has been well studied by Feldmann, Kroll and Stech [27]. We extend the FKS formalism to include the pseudoscalar glueball $G$. In the FKS scheme, the conventional singlet-octet basis and the quark-flavor basis have been proposed. For the latter, the $eG \equiv (u\bar{u} + d\bar{d})/\sqrt{2}$ and $s\bar{s}$ flavor states, labeled by the $\eta_g$ and $\eta_s$ mesons, respectively, are defined. The physical states $\eta$, $\eta'$ and $G$ and their decay constants are related to that of the octet, singlet, and unmixed glueball states $\eta_h$, $\eta_1$ and $g$, respectively, through the rotation

$$ U(\phi_G, \phi_G) \begin{pmatrix} \eta_h \\ \eta_1 \\ g \end{pmatrix} = \begin{pmatrix} f_q \\ f_q \\ f_q \end{pmatrix} G \tag{3}, $$

where $\phi = \theta + 54.7^\circ$ with $\theta$ being the $\eta - \eta'$ mixing angle in the octet-singlet basis, and $\phi_G$ is the mixing angle between $G$ and $\eta_1$; that is, we have assumed that $\eta_8$ does not mix with the glueball. The decay constants for the
physical and flavor states are defined by 
\[
\langle 0 | i \gamma^\mu \gamma_5 q | \eta_q (P) \rangle = - \frac{i}{\sqrt{2}} f_q P^\mu , \quad \langle 0 | \bar{q} \gamma^\mu \gamma_5 s | \eta_s (P), g (P) \rangle = - \frac{i}{\sqrt{2}} f_{q,s}^0 P^\mu ,
\]
\[
\langle 0 | \bar{s} \gamma^\mu \gamma_5 s | \eta_s (P) \rangle = - i f_s P^\mu , \quad \langle 0 | \bar{s} \gamma^\mu \gamma_5 s | \eta_s (P), g (P) \rangle = - i f_{s,s}^0 P^\mu .
\]

Applying the equations of motion for the anomalous Ward identity 
\[
\partial_\nu (\bar{q} \gamma^\mu \gamma_5 q) = 2 i m_q \bar{q} \gamma_5 q + \frac{\alpha_s}{4 \pi} G_{\mu \nu} \tilde{G}_{\mu \nu} ,
\]

between vacuum and |\eta\rangle, |\eta'\rangle and |G\rangle, we derive six equations for many unknowns. Hence we have to reply on the large $N_c$ counting rules to solve the equations step by step. To the leading order of $1/N_c$ expansion, we found that the ratio of two of the equations leads to 
\[
c \theta (s \phi - \sqrt{2/3} c \theta \Delta_G)m_{\eta'}^2 - s \theta (c \phi + \sqrt{2/3} s \theta \Delta_G)^2 m_\eta^2 - \sqrt{2/3} c \theta \phi m_\eta^2 - \sqrt{2/3} s \theta \phi m_\eta^2 = \frac{\sqrt{2} f_s}{f_\eta} ,
\]

where $\Delta_G = 1 - \cos \phi_G$ and $s \phi$ ($c \phi$) is the shorthand notation for $\cos \phi$ ($\sin \phi$) and similarly for others. This simple equation tells us that the pseudoscalar glueball mass $m_G$ can be determined provided that the mixing angle $\phi_G$ and the ratio $f_s/f_\eta$ are known. Note that the $\phi_G$ dependence appears at order of $\Delta_G \approx \phi_G^2$ for small $\phi_G$. So the solution for $m_G$ is stable against the most uncertain input $\phi_G$. The mixing angles $\phi$ and $\phi_G$ have been measured by KLOE $^{28}$ from the $\phi \to \gamma \eta, \gamma \eta'$ decays. Using the updated results $\phi = (40.4 \pm 0.6)^\circ$ and $\phi_G = (20 \pm 3)^\circ$ obtained by KLOE $^{29}$ in conjunction with $f_s/f_\eta = 1.352 \pm 0.007$, we derive the pseudoscalar glueball mass from Eq. (6) to be 
\[
m_G = (1.4 \pm 0.1) \text{ GeV}.
\]

The proximity of the predicted $m_G$ to the mass of $\eta(1405)$ and other properties of $\eta(1405)$ make it a very strong candidate for the pseudoscalar glueball. Using the above-mentioned values for $\phi$ and $\phi_G$, we obtain the $\eta - \eta' - G$ mixing matrix 
\[
\begin{pmatrix}
|\eta\rangle \\
|\eta'\rangle \\
|G\rangle
\end{pmatrix}
= 
\begin{pmatrix}
0.749 & -0.657 & 0.085 \\
0.600 & 0.728 & 0.331 \\
-0.279 & -0.197 & 0.940
\end{pmatrix}
\begin{pmatrix}
|\eta_8\rangle \\
|\eta_1\rangle \\
|g\rangle
\end{pmatrix} .
\]

Our next task is to check the stability and robustness of our prediction when higher order effects in $1/N_c$ are included. It turns out that the above simple formula Eq. (6) still holds even after keeping the OZI-correcting decay constants, as long as they obey the relations $f_{q,s}^0 = \sqrt{2} f_s$ and $f_{q,s}^0 = f_\eta$. Therefore, if excluding the solutions with large and negative $m_{G_{\eta}}^2$, the range $(1.4 \pm 0.1)$ GeV of the pseudoscalar glueball mass obtained in Eq. (7) will be more or less respected.

One may feel very uncomfortable with our solution for $m_G$ as both lattice QCD and QCD sum rules indicate a pseudoscalar glueball heavier than the scalar one. The point is that lattice calculations so far were performed under the quenched approximation without the fermion determinants. It is believed that dynamical fermions will have a significant effect in the pseudoscalar channel, because they raise the singlet would-be-Goldstone boson mass from that of the pion to $\eta$ and $\eta'$. Indeed, it has been argued that the pseudoscalar glueball mass in full QCD is substantially lower than that in the quenched approximation $^{24}$. In view of the fact that the topological susceptibility is large (of order $10^{-3}$ GeV$^4$) in the quenched approximation, and yet is of order $10^{-3}$ GeV$^4$ in full QCD and zero in the chiral limit, it is conceivable that full QCD has a large effect on the pseudoscalar glueball as it does on $\eta$ and $\eta'$.

Our conclusion of a lighter pseudoscalar glueball is also supported by a recent analysis based on the chiral Lagrangian with instanton effects $^{30}$. Two scenarios for the scalar and pseudoscalar glueball mass difference $\Delta m_G = m(0^{++}) - m(0^{-+})$ with the fixed $m(0^{-+})$ were considered in $^{30}$. For $0^{++} = f_0(600)$ and $f_0(1710)$, it is found that $\Delta m_G = -(0.1 \sim 0.3)$ GeV and $0.2 \sim 1.0$ GeV, respectively. The first scenario with $0^{++} = f_0(600)$ cannot be realized since there is no any $0^{-+}$ glueball candidate at energies between $0.7 \sim 1.0$ GeV. The second scenario indicates the possible candidates of the $0^{-+}$ glueball are $\eta(1295)$, $\eta(1405)$ and $\eta(1475)$. The fact that $\eta(1405)$ has a large production rate while $\eta(1295)$ has not been seen in the radiative $J/\psi$ decays and that $\eta(1405)$ has not been observed in $\gamma \gamma$ reactions while $\eta(1475)$ has supports the proposal that $\eta(1405)$ is indeed a good pseudoscalar glueball candidate. The decay properties of $\eta(1405)$ has been recently studied in $^{31,32}$. 

\[\text{\ldots}\]
CONCLUSIONS

We employed two simple and robust inputs to constrain the mixing matrix of the isoscalar mesons $f_0(1710)$, $f_0(1500)$, $f_0(1370)$. In the SU(3) symmetry limit, $f_0(1500)$ becomes a pure SU(3) octet and is degenerate with $a_0(1450)$, while $f_0(1370)$ is mainly an SU(3) singlet with a slight mixing with the scalar glueball which is the primary component of $f_0(1710)$. These features remain essentially unchanged even when SU(3) breaking is taken into account. From the analysis of the $\eta - \eta'$ - $G$ mixing together with the inputs from the recent KLOE experiment, we found a solution for the pseudoscalar glueball mass around $(1.4 \pm 0.1)$ GeV, suggesting that $\eta(1405)$ is indeed a leading candidate for the pseudoscalar glueball. Contrary to the mainstream, we thus conjecture that the low-lying pseudoscalar glueball is lighter than the scalar one owing to the dynamic fermion or QCD anomaly effect.

For the lattice community, it will be extremely important to revisit and check the chiral suppression effect in scalar glueball decays. The previous lattice calculation in this respect was done almost 15 years ago [17]. If the feature of chiral suppression is confirmed by lattice QCD, it will rule out the possibility of $f_0(1500)$ and $f_0(600)$ as scalar glueball candidates. Although it is a difficult and time-consuming task, a full lattice QCD calculation of the $0^{-+}$ glueball is desperately needed in order to see if the dynamic fermions will affect the pseudoscalar glueball spectroscopy dramatically as they do on $\eta$ and $\eta'$.

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