Identification of a Scalar Glueball

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We perform a coupled channel study of the meson-meson S waves with isospin (I) 0 and 1/2 up to 2 GeV. A new approach is derived that allows one to include the many channels $\pi\pi$, $K\bar{K}$, $\eta\eta$, $\sigma\sigma$, $\eta\eta'$, $\eta'\eta'$, $\rho\rho$, $\omega\omega$, $\omega\phi$, $\phi\phi$, $a_1\pi$, and $\pi^*\pi$ with still few free parameters. It follows that coupled channel dynamics is strong and cannot be neglected in order to study resonance properties in the region 1.4–1.6 GeV. All the resonances with masses below 2 GeV and $I = 0$ and 1/2 are generated. We identify the $f_0(1710)$ and an important contribution to the $f_0(1500)$ as an unmixed glueball. This is based on an accurate agreement of our results with predictions of lattice QCD and the chiral suppression of the coupling of a scalar glueball to $\bar{q}q$. Another pole, mainly corresponding to the $f_0(1370)$, is a pure octet state.

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QCD, the theory of strong interactions, is a non-Abelian Yang-Mills theory so that gluons carry color charge and interact between them. It is generally believed that QCD predicts the existence of mesons without valence quarks, the so-called glueballs. Its confirmation in the spectrum of strong interactions is then at the heart of the theory. In quenched lattice QCD the lightest glueball has the quantum numbers $J^{PC} = 0^{++}$, with a mass of (1.66 ± 0.05) GeV [1]. The closest 0^{++} scalar resonances to this energy range that are listed in the Particle Data Group (PDG) [2] are the $f_0(1500)$ and $f_0(1710)$. Some references favor the former as the lightest scalar glueball [3], while others do so for the latter [4,5].

We analyze the $I = 0$ meson-meson S wave in terms of 13 coupled channels, $\pi\pi(1)$, $K\bar{K}(2)$, $\eta\eta(3)$, $\sigma\sigma(4)$, $\eta\eta'(5)$, $\eta'\eta'(6)$, $\rho\rho(7)$, $\omega\omega(8)$, $K^*\bar{K}^*(9)$, $\omega\phi(10)$ $\phi\phi(11)$, $a_1(1260)\pi(12)$, and $\pi^*(1300)\pi(13)$. The number labeling each state is given between brackets. The multipion states, which play an increasing role for energies above $\sim$1.2 GeV, are mimicked through the $\sigma\sigma$, $\rho\rho$, and $\omega\omega$ channels. We stress that our approach is the first one with such a large number of channels and that a similar scheme could be applied to other controversial meson-meson partial waves. We also study simultaneously the S wave of $K^-\pi^+ \to K^-\pi^+$ (involving $I = 1/2$ and 3/2) with the channels $K\pi$, $K\eta$, and $K\eta'$. As emphasized in Ref. [6], a coupled channel analysis is much more constrained and thus more likely to reveal the number of real resonances required to explain the data.

Let $T_{i,j}^{(I)}$ be the $i \leftrightarrow j$ S-wave amplitude with isospin I and $i, j = 1, \ldots, n$, with n the number of channels. We use the master formula $T^{(I)} = [I + Ng]^{-1}N$, where N is the symmetric matrix of interaction kernels and g is a diagonal matrix of elements $g_i(s)$. The function $g_i(s)$ is calculated from kinematics in terms of a once subtracted dispersion relation and a subtraction constant a_i [7]. Since SU(3) breaking is milder in the vector sector we take $a_7 = a_8 = a_9 = a_{10} = a_{11}$. The rest of the subtraction constants are fitted to the data. The matrix elements $N_{i,j}$ consist of the sum of two tree level contributions. The first is a contact interaction calculated from the lowest order chiral perturbation theory Lagrangian, $\mathcal{L}_2$. The second is due to the exchange of bare resonances in the s channel with the couplings calculated from the lowest order chiral Lagrangian including an octet and singlet of 0^{++} resonances, $\mathcal{L}_S$ [8]. Explicit expressions of $N_{i,j}$ can be found in Ref. [7] for the simplified case of three channels without including the η_1 field. We extend these Lagrangians from SU(3) to U(3), as the η_1 field is needed to deal with the η and η' mesons, similarly as in Ref. [9]. The matrix $\Phi = \sum_{i=1}^{8} \phi_i\lambda_i/\sqrt{2} + \eta_1/\sqrt{3}$ incorporates in a standard way the nonet of the lightest pseudoscalars. We also employ $U = \exp(i\sqrt{2}\Phi/f)$ and the covariant derivative $D_\mu U = \partial_\mu U - ir_\mu U + iU\ell_\mu$, with $f = f_\pi = 92.4$ MeV. The classical left and right external fields, r_μ and ℓ_μ, respectively, are necessary to gauge the global chiral symmetry to a local one [8]. The field $v_\mu = (r_\mu + \ell_\mu)/2$ plays a special role in our approach since it is identified with λW_μ, where W_μ is the nonet of the lightest 1^{--} vector resonances and $\lambda = 4.3$ is fixed from the width $\rho \to \pi^+\pi^-$. The couplings of the vector-vector states to the pseudoscalar-pseudoscalar and $\sigma\sigma$ ones are then determined by minimal coupling [10]. Our fits require a singlet and two octets of bare resonances. The two octets were already considered in Ref. [9] in the study of $K^-\pi^+ \to K^-\pi^+$. We fix the parameters of the first octet, mass and coupling constants, to those in Ref. [9], $M_8^{(1)} = 1.29$ GeV, $c_d^{(1)} = c_m^{(1)} = 26$ MeV. The bare mass of the second octet is fixed from the same reference, $M_8^{(2)} = 1.90$ GeV. We are then left with three parameters for the singlet, M_1, $\tilde{c}_d^{(1)}$, $\tilde{c}_m^{(1)}$, and two for the second octet, $c_d^{(2)}$ and $c_m^{(2)}$. It results from our fits that $M_1 \lesssim 0.9$ GeV.

Concerning the $\sigma\sigma$ channel we derive a novel method to calculate its transition amplitudes, $N_{i,4}$, without including any new free parameter. This can be done because the σ corresponds to a pole due to the interactions between two pions in the $I = 0$ S wave, $(\pi\pi)_0$ [11]. For the interaction kernel $N_{i,4}$ one starts by calculating from $\mathcal{L}_2$ and $\mathcal{L}_S$ the tree level amplitude $T_{i,4}^{2+S}$ for $i \to (\pi\pi)_0(\pi\pi)_0$. To take into account the pion final state interactions, $T_{i,4}^{2+S}$ is multiplied by the factor $\prod_{k=1}^{m} 1/D(s_k)$, with m the number of σ's in the scattering process (2 or 4) and s_k the total center of mass (c.m.) energy squared of the kth pair. We use here that the rescattering of two $I = 0$ S-wave pions from a production kernel is given by the factor $1/D = 1/(1 + V_2 g_1)$, with $V_2 = (s - m_\pi^2/2)/f^2$ [11]. To isolate $N_{i,4}$ one takes the limit (for definiteness $i \neq 4$)

$$\lim_{s_1,s_2\to s_\sigma} \frac{T_{i,4}^{2+S}}{D_{\mathrm{II}}(s_1)D_{\mathrm{II}}(s_2)} = \frac{N_{i,4} g_{\sigma\pi\pi}^2}{(s_1 - s_\sigma)(s_2 - s_\sigma)}, \qquad (1)$$

where the subscript II indicates that the corresponding function is calculated on the second Riemann sheet, s_σ is the σ pole position, and $g_{\sigma\pi\pi}$ is its coupling to $\pi\pi$. Performing the Laurent expansion around s_σ of $1/D_{\mathrm{II}}(s) = \alpha_0/(s - s_\sigma) + \cdots$ the evaluation of $N_{i,4}$ from Eq. (1) requires the ratio $(\alpha_0/g_{\sigma\pi\pi})^2$. Since $g_{1,\mathrm{II}}(s_\sigma) = -f^2/(s_\sigma - m_\pi^2/2)$ at s_σ, where $1 + V_2 g_{1,\mathrm{II}} = 0$, and taking $T_{\mathrm{II}} \simeq V_2/(1 + V_2 g_{1,\mathrm{II}})$, appropriate for these energies [11], then $(\alpha_0/g_{\sigma\pi\pi})^2 = f^2/(1 - \frac{dg_{1,\mathrm{II}}}{ds}|_{s_\sigma} \frac{(s_\sigma - m_\pi^2/2)^2}{f^2}) \simeq f^2$. In this way, $N_{i,4} = T_{i,4}^{2+S} f^2$, $i \neq 4$, and $N_{4,4} = T_{4,4}^{2+S} f^4$. Using $N_{i,4}$ evaluated with $s_k = s_\sigma$ violates unitarity because s_σ is complex and $N_{i,4}$ must be real. Instead, we interpret the width of the σ resonance as a Lorentzian mass distribution around its nominal mass value $\sim$450 MeV with a width $\sim$500 MeV. In this way the σ masses ($\sqrt{s_k}$) used in the functions $N_{i,4}$ and g_4 are folded with the previous mass distribution. Similarly, for the $\rho\rho$ state, g_7 is also convoluted with a ρ mass distribution.

We fit our 12 free parameters to 370 data points from threshold up to 2 GeV. The data comprise the $I = 0$ S-wave $\pi\pi$ phase shifts δ_0^0, the elasticity $\eta_0^0 = |S_{1,1}|$, the $I = 0$ S-wave $\pi\pi \to K\bar{K}$ phase shifts $\delta_{1,2}$ and modulus $|S_{1,2}|$, the S-wave contribution to the $\pi\pi \to \eta\eta, \eta\eta'$ event distributions, and the phase (ϕ) and modulus (A) of the $K^-\pi^+ \to K^-\pi^+$ amplitude from the LASS data. The S-matrix element $S_{i,j}$ is given by $S_{i,j} = \delta_{ij} + 2i\sqrt{\rho_i} T_{i,j}^{(I)} \sqrt{\rho_j}$, where $\rho_i = q_i/8\pi\sqrt{s}$ and q_i is the c.m. three-momentum for channel i. In order, these data are shown on the first eight panels of Fig. 1 from top to bottom and left to right. For $\sqrt{s} \leq m_K$ in the $\delta_0^0(s)$ panel we have the inset corresponding to the precise data from K_{e4} decays. The reproduction of the data is fair, as shown in the figure. The dashed lines on the first eight panels include the $a_1\pi$ and $\pi^*\pi$ states, while the solid ones do not. The similarity between both curves indicates that these channels give small contributions. The width of the band rep-

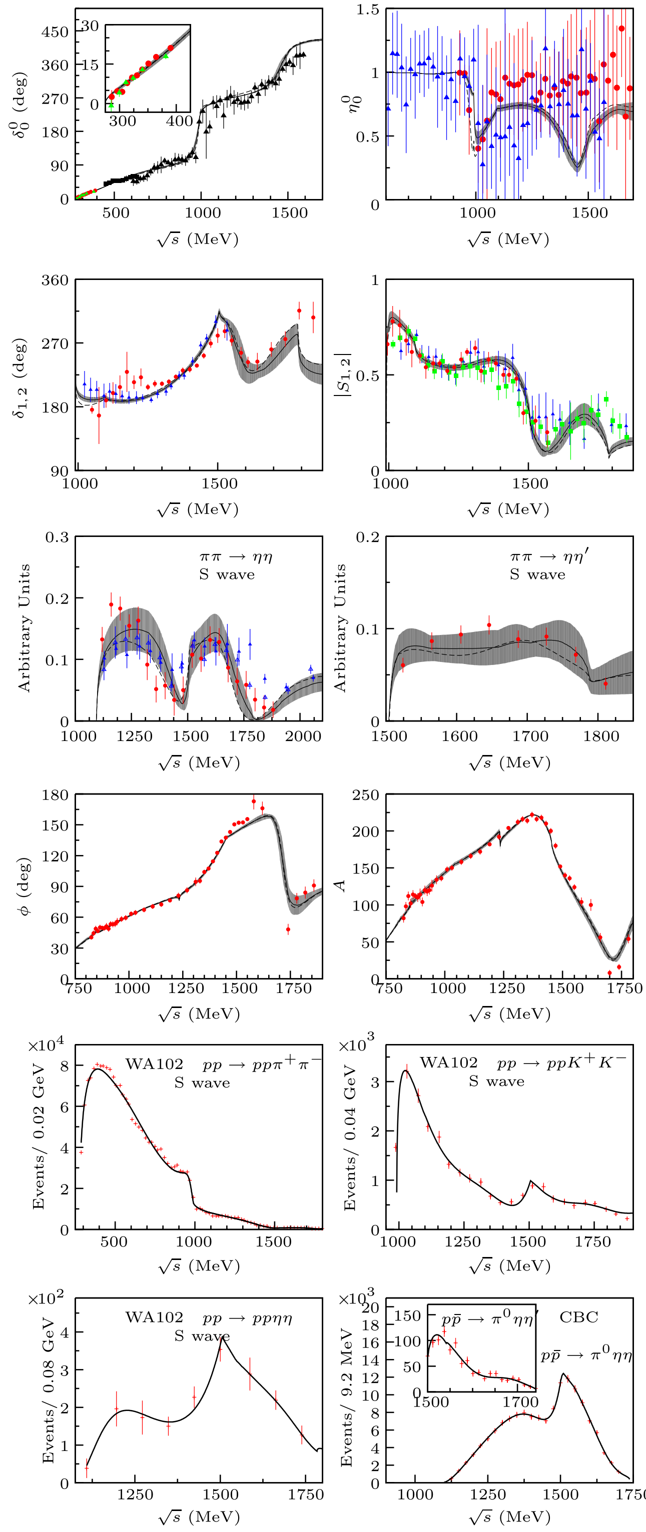

FIG. 1 (color online). Fit to experimental data.

resents our systematic uncertainties at the level of 2 standard deviations, $n_\sigma = \Delta\chi^2/(2\chi^2)^{1/2}$ [12]. Compared with other works we determine the interaction kernels from standard chiral Lagrangians, avoid *ad hoc* parametrizations, include many more channels, and fewer free parameters are used. For $I = 1/2$ the κ pole is located at $(708 \pm 6 - i313 \pm 10)$ MeV, the $K_0^*(1430)$ at $(1435 \pm 6 - i142 \pm 8)$ MeV, and the $K_0^*(1950)$ at $(1750 \pm 20 - i150 \pm 20)$ MeV, similarly to Ref. [9]. For $I = 0$ one has the $f_0(600)$ or σ at $(456 \pm 6 - i241 \pm 7)$ MeV and the $f_0(980)$ at $(983 \pm 4 - i25 \pm 3)$ MeV. There are poles at $(1690 \pm 20 - i110 \pm 20)$ MeV, corresponding to the $f_0(1710)$, and at $(1810 \pm 15 - i190 \pm 20)$ MeV, with mass and width in agreement with those reported for the $f_0(1790)$ by BESII. In the PDG [2] the width for the $f_0(1710)$ is 137 ± 8 MeV, much smaller than 220 ± 40 MeV from our pole position. However, we have checked that on the real axis the value of the width corresponding to the half maximum for the partial waves with prominent $f_0(1710)$ peaks is just 160 MeV [13]. This reduction is due to a Flatté effect because of the opening of new channels along the resonance region [14], and the agreement with the PDG is restored. The other poles at $(1466 \pm 15 - i158 \pm 12)$ MeV and $(1602 \pm 15 - i44 \pm 15)$ MeV are referred in the following as f_0^L and f_0^R, respectively. The f_0^L pole gives rise to the $f_0(1370)$ and its mass is very close to the one preferred by the high statistics study of the Belle Collaboration on $\gamma\gamma \to \pi^0\pi^0$ [15], at 1.47 GeV. On the other hand, the sum of both f_0^R and f_0^L originates the $f_0(1500)$ resonance shape below the $\eta\eta'$ threshold. Note that these two poles are located on a Riemann sheet that connects continuously with the physical axis only below the $\eta\eta'$ threshold. Above it the Riemann sheet that matches continuously with the physical s axis changes. On this sheet the f_0^L and f_0^R poles disappear and the $f_0(1710)$ and $f_0(1790)$ are the poles that matter. As a result, above the $\eta\eta'$ threshold the $f_0(1500)$ comes from the coherent sum of the latter two poles. Similar results on the generation of the $f_0(1500)$ were also obtained in Ref. [6]. This reference performed a coupled channel analysis of considerable data, some of them common to us but many other independent (e.g., it includes J/Ψ decay data) and considered the $\pi\pi$, $K\bar{K}$, $\eta\eta$, and $\eta\eta'$ coupled channels. Note that despite the fact that we have included only three bare resonances in $I = 0$ we have generated six. We can include an arbitrary number of bare resonances, but the unavoidable consequence from these trials is that no good fit results with a single pole corresponding to the $f_0(1500)$ resonance. In any case, our calculation unambiguously shows that a coupled channel analysis should be undertaken in the region 1.4–1.6 GeV, otherwise the results are not realistic. Unfortunately, this is the standard situation for the present analyses and determinations [2] of the $f_0(1500)$ resonance parameters by fitting data.

In Fig. 1 we show in panels 9, 10, and 11 data from the WA102 Collaboration on $pp \to pp\pi^+\pi^-$, K^+K^- and $\eta\eta$ at 450 GeV/c. These data are also very convenient theoretically as the two mesons produced interact negligibly with the two protons because the latter are very energetic [16]. To fit the data we employ a coherent sum of Breit-Wigner functions and a nonresonant term, similarly as done by the WA102 Collaboration [17]:

$$
\begin{aligned}
&\text{(i) } \sqrt{s} < m_\eta + m_{\eta'}, \qquad A = \{\sigma, f_0(980), f_0^L, f_0^R\},\\
&\quad A(\sqrt{s})_i = NR(\sqrt{s})_i + \sum_{j\in A} \frac{a_j e^{i\theta_j} g_{j;i}}{M_j^2 - s - iM_j\Gamma_j},\\
&\text{(ii) } \sqrt{s} > m_\eta + m_{\eta'},\\
&\quad B = \{\sigma, f_0(980), f_0(1710), f_0(1790)\},\\
&\quad A(\sqrt{s})_i = NR(\sqrt{s})_i + r_i + \sum_{j\in B} \frac{a_j e^{i\theta_j} g_{j;i}}{M_j^2 - s - iM_j\Gamma_j},\\
&\quad NR(\sqrt{s})_i = \alpha(\sqrt{s} - m_k - m_\ell)^\beta e^{-\gamma\sqrt{s} - \delta s},
\end{aligned} \tag{2}
$$

where a_j and θ_j are the modulus and the phase of the production vertex of the jth resonance, M_j, Γ_j, and $g_{j;i}$ are, respectively, the mass, the width, and the coupling to channel i of the same resonance. In addition, $m_k + m_l$ is the threshold for the channel i, and α, β, γ, δ are real parameters. The form of the nonresonant term is taken from the WA102 Collaboration [17]. The constant r_i is fixed so as the amplitude $A(\sqrt{s})_i$ is continuous at $w_{\eta\eta'} \equiv m_\eta + m'_\eta$. Γ_j in Eq. (2) is the largest between its value from the pole position and the one calculated by summing the partial decay widths $\Gamma_{j;i} = \theta(\sqrt{s} - m_k - m_\ell)\lambda_i |g_{j;i}|^2 q_i/(8\pi M_j^2)$, with $\lambda_i = 1/2$ for identical particles. Equation (2) incorporates important new facts compared to the analyses of the WA102 Collaboration. First, the pole positions, $M_j - i\Gamma_j/2$, and couplings, $g_{j;i}$, for the different resonances are those already determined from our study of the scattering data. Second, the a_j and θ_j parameters are the same for all the reactions. On the other hand, Eq. (2) is a toy model which shows in simple terms how the change of sheet at the $\eta\eta'$ threshold takes place with the corresponding change in the poles involved. Our fitted curves correspond to the solid lines and prominent peaks associated with the $f_0(980)$, $f_0(1710)$, and $f_0(1500)$ are observed. In the last panel we show our good reproduction of the Crystal Barrel data (CBC) for the $\eta\eta$ and $\eta\eta'$ (inset) mass projections from $p\bar{p}$ annihilation into $\pi^0\eta\eta$ and $\pi^0\eta\eta'$, respectively. Equation (2) without $NR(\sqrt{s})$ is used.

In Table I we give the couplings of the f_0^L (identified as the $f_0(1370)$), f_0^R, and $f_0(1710)$ poles to the two pseudoscalar channels. We observe that the couplings of the f_0^R and $f_0(1710)$ are quite similar. This is so because the two poles coalesce in the same one when moving continuously from the sheet of one of them to the sheet of the other. They correspond to the same underlying resonance, but split in two due to the interaction in coupled channels. From

TABLE I. Couplings of the $f_0(1370)$, f_0^R and $f_0(1710)$.

GeV	$f_0(1370)$	f_0^R	$f_0(1710)$
$\lvert g_{\pi^+\pi^-}\rvert$	3.59 ± 0.16	1.30 ± 0.22	1.21 ± 0.16
$\lvert g_{K^0\bar{K}^0}\rvert$	2.23 ± 0.18	2.06 ± 0.17	2.0 ± 0.3
$\lvert g_{\eta\eta}\rvert$	1.7 ± 0.3	3.78 ± 0.26	3.3 ± 0.8
$\lvert g_{\eta\eta'}\rvert$	4.0 ± 0.3	4.99 ± 0.24	5.1 ± 0.8
$\lvert g_{\eta'\eta'}\rvert$	3.7 ± 0.4	8.3 ± 0.6	11.7 ± 1.6

the couplings of the $f_0(1710)$ one can calculate the branching ratios $\Gamma(K\bar{K})/\Gamma_{\rm total} = 0.36 \pm 0.12(0.38^{+0.09}_{-0.19})$, $\Gamma(\eta\eta)/\Gamma_{\rm total} = 0.22 \pm 0.12(0.18^{+0.03}_{-0.13})$, and $\Gamma(\pi\pi)/\Gamma(K\bar{K}) = 0.32 \pm 0.14(0.41^{+0.11}_{-0.17})$, where the values of the PDG are given between brackets. The agreement is excellent. We also obtain that the $f_0(1790)$ has a small $K\bar{K}$ coupling, and this is a major difference with respect to the $f_0(1710)$ as stressed by BESII. The couplings of the $f_0^L - f_0(1370)$ in Table I correspond to the pure $I = 0$ octet member $(\bar{u}u + \bar{d}d - 2\bar{s}s)/\sqrt{6}$ because they are very close to the tree level ones $|g_{\pi^+\pi^-}| = 3.9$, $|g_{K^0\bar{K}^0}| = 2.3$, $|g_{\eta\eta}| = 1.4$, $|g_{\eta\eta'}| = 3.7$, $|g_{\eta'\eta'}| = 3.8$ GeV, calculated from $\mathcal{L}_S$ [8], with $c_d^{(1)}$, $c_m^{(1)}$, and $M_8^{(1)}$ given above. We have also checked that this is the case for the $K_0^*(1430)$ resonance which is the $I = 1/2$ member of the same octet. It follows then that the first octet is a pure one, not mixed with the nearby f_0^R and $f_0(1710)$. The $f_0^L - f_0(1370)$ couplings imply a large width to $\pi\pi$ with $\Gamma(f_0(1370) \to 4\pi)/\Gamma(f_0(1370) \to \pi\pi) = 0.30 \pm 0.12$, in agreement with the interval 0.10–0.25 of Ref. [18].

Let us see that the pattern of the couplings of the $f_0(1710)$ (and hence also for the f_0^R) corresponds to the chiral suppression of the coupling of a scalar glueball, G_0, to $\bar{q}q$ [5]. According to Ref. [5] this coupling is proportional to the quark mass, which then implies a strong suppression in the production of $\bar{u}u$ and $\bar{d}d$ relative to $\bar{s}s$ from G_0. With a pseudoscalar mixing angle $\sin\beta = -1/3$ one has that $\eta = -\eta_s/\sqrt{3} + \eta_u\sqrt{2/3}$ and $\eta' = \eta_s\sqrt{2/3} + \eta_u/\sqrt{3}$ with $\eta_s = \bar{s}s$ and $\eta_u = (\bar{u}u + \bar{d}d)/\sqrt{2}$. Denoting by g_{ss} the production of $\eta_s\eta_s$, g_{sn} that of $\eta_s\eta_u$, and g_{nn} for $\eta_u\eta_u$,

$$\begin{aligned} g_{\eta'\eta'} &= 2g_{ss}/3 + g_{nn}/3 + 2\sqrt{2}g_{ns}/3, \\ g_{\eta\eta'} &= -\sqrt{2}g_{ss}/3 + \sqrt{2}g_{nn}/3 + g_{ns}/3, \\ g_{\eta\eta} &= g_{ss}/3 + 2g_{nn}/3 - 2\sqrt{2}g_{ns}/3. \end{aligned} \tag{3}$$

If the chiral suppression of Ref. [5] operates then $|g_{ss}| \gg |g_{nn}|$. This together with the Okubo-Zweig-Iizuka rule suppresses the coupling g_{ns}. Taking, e.g., the couplings of f_0^R one obtains $g_{ss} = 11.5 \pm 0.5$, $g_{ns} = -0.2$, and $g_{nn} = -1.4$ GeV, and the strong suppression is clear. We now consider the $K\bar{K}$ coupling. A K^0 in terms of valence quarks corresponds to $\sum_{i=1}^{3} \bar{s}_i u^i/\sqrt{3}$, summing over the color indices, and analogously for the $\bar{K}^0$. The production of a color singlet $\bar{s}s$ from the $K^0\bar{K}^0$ requires then the combination $\bar{s}_i s^j = \delta_i^j \bar{s}s/3 + (\bar{s}_i s^j - \delta_i^j \bar{s}s/3)$, and similarly for $\bar{u}_j u^i$. As the production occurs from the color singlet $\bar{s}s$ source, only the configuration $\bar{s}s\bar{u}u$ contributes, picking up a suppression factor of 1/3. In addition, the coupling g_{ss} has an extra factor 2 compared to that of a $\bar{s}s\bar{u}u$, because the former contains two $\bar{s}s$. One then expects for the coupling to $K^0\bar{K}^0$ an absolute value $g_{ss}/6$, $|g_{K^0\bar{K}^0}| \simeq 2$ GeV, as in Table I. Another resonance with a known enhanced coupling to $\bar{s}s$ is the $f_0(980)$. However, the sizes of its couplings to $\eta\eta$, $\eta\eta'$, and $\eta'\eta'$ follow the opposite order to the $f_0(1710)$ and f_0^R cases, and all of them are much smaller than the coupling to $K\bar{K}$. Note that quenched lattice QCD [4] establishes that the couplings of the lightest scalar glueball to pseudoscalar pairs in the SU(3) limit scale as the quark mass, in support of the chiral suppression mechanism of Ref. [5]. This mechanism also implies that the glueball should remain unmixed. This accurately fits with our previous result that both the f_0^R and $f_0(1710)$ do not mix with the nearby f_0^L. In addition, the masses of the f_0^R and $f_0(1710)$ poles are in excellent agreement with the quenched latticed QCD result [1].

In summary, we have presented a coupled channel study of the $I = 0, 1/2$ meson-meson S waves up to 2 GeV with 13 coupled channels. Coupled channel dynamics is particularly relevant in the region between 1.4–1.6 GeV and should be taken into account in analyses of data. The $f_0(1710)$ is identified as an unmixed scalar glueball by the study of their couplings. This is also the case for the f_0^R pole but slightly shifted because of the coupled channel dynamics. The $f_0(1500)$ is made up of the sum of the poles f_0^L and f_0^R below the $\eta\eta'$ threshold and of the $f_0(1710)$ above it. This naturally explains why both the $f_0(1710)$ and the $f_0(1500)$ are copiously produced in gluon rich processes. The pole f_0^L, from which the $f_0(1370)$ mainly originates, is shown to be a pure octet member.

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