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WIDE SCALAR GLUONIUM AND THE QUARK-ANTIQUARK SCALAR NONET

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1. INTRODUCTION

The existence of bound states of gluons called gluonia or glueballs is a clear prediction [1] of QCD but predictions of their properties are less definite. While lattice QCD calculations of, e.g., the mass m_σ of the scalar gluonium $\sigma \to (gg)$ may now be with some reservation approaching approximately the value $m_\sigma \approx 1.3$ GeV [2] in a reasonable agreement with the independent estimates based on bags [3], QCD sum rules [4a,b] (see, however, [4c] where a possibility of a light and narrow scalar gluonium is advocated) and others [5], the decay properties of σ are more puzzling. In particular, it is not clear whether its decay widths into ordinary hadrons should be small or large [6] compared to its mass m_σ. Within a large N_c colour counting [7] the widths of gluonia are $O(1/N_c^2)$ and thus are expected to be small of an order of ten MeV [6]. However, in a previous paper [8] an exception from the $1/N_c$ -rule has been demonstrated explicitly for the scalar gluonium σ decaying into $\pi\pi$.

In fact, analysing the coupling of σ to $\pi\pi$ on the basis of low-energy theorems [9] of broken chiral symmetry and scale invariance through the anomalous trace of the hadronic energy-momentum tensor [10] implemented by using phenomenological Lagrangians [11], the following partial width has been found [8]

$$\Gamma(\sigma \to \pi\pi) = \frac{3}{16\pi b} \frac{m_\sigma^5}{G_0}, \tag{1}$$

where $G_0 = \langle 0| (\alpha_s/\pi) G^a_{\mu\nu} G_a^{\mu\nu} |0\rangle$ is a familiar gluon condensate [12] of QCD, $G^a_{\mu\nu}$ are the gluonic field strength tensors and $b = (11N_c - 2N_F)/3$ with N_c and N_F being the numbers of colours and flavours, respectively. Although being of $O(1/N_c^2)$ - order, (1) gives a large value of $\Gamma(\sigma \to \pi\pi)$ for $N_c = N_F = 3$ and the ITEP value of $G_0 = 0.012$ GeV^4 [12] if m_σ were around or above 1 GeV. Such a conclusion has also been independently supported recently by Gounaris et al.[13] who have found that a 0^{++} gluonium cannot be narrow if it exists at all. Moreover, this agrees, too, with QCD sum rule analysis in [4b].

The result (1) has been found[8] under the assumption that only a scalar gluonium σ dominates the following low-energy theorems[9]

$$i\int d^4x \langle 0|T(H(x)H(0))|0\rangle = \frac{b}{2} G_0 , \quad (2)$$

etc., where the scalar gluonic current H(x) is given by the anomalous trace of the hadronic energy-momentum tensor of QCD[10]

$$H \equiv -(\theta^{\mu}_{\mu})_{an} = \frac{b}{8} \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^{\mu\nu}_a . \quad (3)$$

Other states like scalar $q\bar{q}$, etc., mesons and eventual mixing of σ with them were ignored and the effective Lagrangian used for coupling of σ to pseudoscalar mesons[8] was constructed so as to satisfy (2) and (3).

Although an inclusion[14] of mixing with the quarkonium mesons alone does not seem to change the above conclusion, it has been noted[15] that if in addition a derivative coupling term is introduced into the usual type of the linear sigma model Lagrangian[11b], then the problem of large widths of the ordinary scalar $q\bar{q}$ mesons as well as a heavy physical state (called here G and resulting after mixing between the pure gluonium σ and a flavour $SU(3)_F$ -singlet (u, d, s) quarkonium $S_0 = (1/3)^{1/2}$ $(u\bar{u} + d\bar{d} + s\bar{s})$) can be solved. Since this looks like a possible contrast with the result (1), we want to investigate these effects in more detail here. Using phenomenological Lagrangians[11] for broken chiral and scale invariance we shall show in section 2 that, in fact, while such effects can suppress the couplings of the heavier state G to the octet of the pseudoscalar Goldstone mesons ϕ_i , the decay of the lighter scalar particle ϵ still remains in agrgement with (1) for a small ratio of the squared masses $M^2_\epsilon / M^2_G << 1$.

On the basis of the large - N_c counting[7] we shall find that just this state ϵ having a large $\epsilon \to \pi\pi$ decay width of the order $O(1/N_c^2)$ in agreement with (1) should play the role of an effective physical gluonium while the narrow state G plays the role of an effective $SU(3)_F$ singlet quarkonium because its coupling $G\phi\phi$ has the same behaviour in $1/N_c$ expansion[7] as the couplings $S_i\phi\phi$ for the nonet of the scalar $q\bar{q}$ mesons $S_i (i = 0, 1, \ldots, 8)$ and so no contrast with (1) appears.

The exact, not dependent on M^2_ϵ/M^2_G expansion results will be presented in section 3. These results, however, are based on some more sophisticated but reasonable assumptions on the bare gluonium σ and quarkonium S_0 masses $M_{\sigma\sigma}$ and M_{oo} , respectively. In particular, we assume and discuss the attractive possibility that these masses are equal to each other, i.e.

$M^2_{\sigma\sigma} = M^2_{oo} = M^2$, with the value of M around 1.3 GeV as may be suggested by the quark model and by the recent QCD lattice estimates[2] taken seriously despite of the existing reservations. Using, moreover, the"standard" values of G_0[12,16] we shall find consistently with the results of the previous section that the heavier meson G has the predicted mass $M_G = (3/2)^{1/2}$ M and suppressed decays into $\pi\pi$ and $K\bar{K}$ while its gluonium companion ϵ is predicted to be the wide state lying probably below 1 GeV with the mass $M_\epsilon = (1/2)^{1/2}$ M. The effective $SU(3)_F$ singlet quarkonium G[17] is in agreement with the GAMS $f_0(1590)$ meson (the old name G(1590)) discovered at the IHEP[18] and the meson ϵ may not be inconsistent[14a] with the broad, not easily observable[19] (see also[20] and references therein) state ϵ(900) below 1 GeV seen probably again recently by analyzing[21] the AFS data[22] obtained at the CERN's ISR.

On the basis of these results we suggest in conclusion, section 4, that a mixing pattern for the scalar $q\bar{q}$ nonet is far from the ideal one and could instead be analogous to that for the pseudoscalar mesons π, K, η and η' with probably negligible mixing between the octet and $SU(3)_F$ singlet scalar states. A picture that arises consists of the I = 1, I = 1/2 and I = 0 scalar $q\bar{q}$ octet members corresponding to the experimental candidates[19] a_0(980) and/or a_0(1400)?[23], K^*_0(1350) and f_0(1300) (old names δ(980) and/or δ'(1400)?[23], κ(1350) and ϵ(1300)), respectively, and of other two $SU(3)_F$ singlets G and ϵ representing by themselves lagre, approximately half-and-half mixture of the pure gluonic and quark degrees of freedom and corresponding probably to the experimental states f_0(1590)[18,19] and ϵ(900)[21].

The present scenario of the scalar mesons does not explain the S*(975)-state[19] (or even more states?[21]) and thus we encourage efforts for still other suppositions[24,25,26] made to understand better the effects near the $K\bar{K}$ threshold.

2. THE ANALYSIS OF DERIVATIVE COUPLINGS

Let us recall that a convenient and original way to investigate interactions of the $q\bar{q}$ scalar and pseudoscalar mesons is to use a linear sigma model (LSM)[11b] with the 3x3 field matrix U(x):

$$U(x) = \lambda_j (S_j(x) + i\phi_j(x)), \quad (4)$$

where $S_j(x)$ and $\phi_j(x)$ ($j = 0, 1, \ldots, 8$) are the $q\bar{q}$ nonet scalar and pseudoscalar fields, respectively and λ_j are the Gell-Mann matrices λ normalized to $Tr(\lambda_i \lambda_j) = 2\delta_{ij}$. Neglecting the quark mass term, this model is described by the following Lagrangian

$$\mathcal{L}_{LSM} = \frac{1}{4} Tr(\partial_\mu U \partial^\mu U^+) - V, \qquad (5)$$

where the potential V is an arbitrary chiral invariant function of the fields under consideration. We assume that chiral symmetry is spontaneously broken and reparametrize the fields S_i so as to have $S_i(x) = \langle 0 | S_i | 0 \rangle + \tilde{S}_i(x)$, where $\tilde{S}_i(x)$'s are already correct fields and $\langle 0|S_i|0\rangle = (\sqrt{3}/\sqrt{2}) f_o \delta_{io}$ with $f_o = -f_\pi$, $f_\pi = 93$ MeV being the pion decay constant as one can easily see from the usual definition of f_π through the axial current. At a tree level, Lagrangian (5) gives the couplings between the scalars and pseudoscalar pairs as follows[11b]

$$\mathcal{L}^{nonder}_{S\phi\phi} = -\frac{1}{2} \frac{M^2_{S_k}}{f_0} d_{kij} \tilde{S}_k(x)\phi_i(x)\phi_j(x), \qquad (6)$$

where $d_{kij} = (1/4)Tr(\{\lambda_i, \lambda_j\}\lambda_k)$ are fully symmetrical and $M^2_{S_k}$'s are the squared masses of the decaying scalars, so their widths are much larger than the experimental ones. A possible way out of this discrepancy has been suggested by Gomm et al. in[15] where they have noted that an introduction of chiral invariant derivative terms like, e.g.

$$Tr(\partial_\mu U \partial^\mu U^+ UU^+), \text{ etc.} \qquad (7)$$

into (5) can add a derivative interaction term of the type

$$\frac{A}{f_o} d_{kij} \tilde{S}_k(x)(\partial_\mu \phi_i(x))(\partial^\mu \phi_j(x)), \qquad (8)$$

(A being an arbitrary number) to (6) and this again yields amplitudes proportional to the squared masses of the decaying scalar mesons. They have concluded[15] that the general amplitude for the decay $S \to \phi\phi$ as based on chiral symmetry is proportional to the squared mass of the decaying meson S, or equivalently, the general $S\phi\phi$ coupling is supposed to be of a derivative type[15] :

$$\mathcal{L}^{der}_{S\phi\phi} = \frac{\gamma}{f_o} d_{kij} \tilde{S}_k(x)(\partial_\mu \phi_i(x))(\partial^\mu \phi_j(x)), \qquad (9)$$

where $\gamma = 1 + A$ is the only (numerical) parameter to be specified from the scalar $q\bar{q}$ meson decays.

If $a_o(980)$ is a $q\bar{q}$ state, then the width $\Gamma(a_o(980) \to \eta\pi) = 54$ MeV[19] implies

$$\gamma = \begin{cases} 0.24 \\ 0.27 \\ 0.34 \end{cases} \quad \text{for} \quad \theta_{\eta\eta'} = \begin{cases} -18^\circ \\ -10^\circ \\ 0^\circ \end{cases} \qquad (10a)$$

where $\theta_{\eta\eta'}$ is the $\eta\eta'$ mixing angle. Knowing γ, (9) predicts all other decay widths, e.g., for the $K^*_0(1350) \to K\pi$ decay we obtain

$$\Gamma(K^*_0(1350) \to K\pi) = \begin{cases} 210 \text{ MeV} \\ 260 \text{ MeV} \\ 420 \text{ MeV} \end{cases} \quad \text{for } \gamma = \begin{cases} 0.24 \\ 0.27 \\ 0.34 \end{cases} \qquad (10b)$$

in a good agreement with otherwise not very precise experimental results[19]. In fact, if the partial decay width $\Gamma(K^*_0(1350) \to K\pi)$ is experimentally known more precisely, then (10b) would be more convenient for the determination of γ since (10b) depends neither on $\theta_{\eta\eta'}$, nor on a particular interpretation of $a_0(980)$ as, e.g. $q\bar{q}$[21,27,28], $q^2\bar{q}^2$[25], etc.[26] state. However, the agreement between (10a) and (10b), if not accidental, seems rather to uphold the $q\bar{q}$ assignment of the $a_0(980)$ state.

When a pure scalar gluonium field $\sigma(x)$ parametrized as[11a,c]

$$\sigma(x) = \sigma_0 \exp\left(\frac{\tilde{\sigma}(x)}{\sigma_0}\right), \qquad (11)$$

with $\sigma_0 = \langle 0 | \sigma | 0 \rangle$ is added to the $q\bar{q}$ scalar $S_i(x)$ and pseudoscalar $\phi_i(x)$ fields, then this system is described by the following Lagrangian

$$\mathcal{L} = \frac{1}{2}(\partial_\mu \sigma)^2 + \frac{1}{4} Tr(\partial_\mu U \partial^\mu U^+) - V + \mathcal{L}', \qquad (12)$$

where the chiral symmetric potential V containing also σ is assumed[8,15] to obey the trace anomaly equation

$$(\theta^{\mu}_{\mu})_{an} = -\frac{b}{8} G_0 \left(\frac{\sigma(x)}{\sigma_0}\right)^4 = 4V - \sigma\frac{\partial V}{\partial \sigma} - S_i\frac{\partial V}{\partial S_i} - \phi_i\frac{\partial V}{\partial \phi_i}, \qquad (13)$$

to guarantee the realization of (2) in the present model. $\mathcal{L}'$ labels derivative terms like (7)/15/ and with a conventional assignment of dimension 1 to U and σ these terms in $\mathcal{L}'$ should be the following, e.g.,

$$K_1 = \frac{3}{2}[\mathrm{Tr}(UU^+)]^{-1}\mathrm{Tr}(\partial_\mu U\partial^\mu U^+ UU^+),$$

$$K_2 = \frac{\sigma_0^2}{2f_0^2\sigma^2}\mathrm{Tr}(\partial_\mu U\partial^\mu U^+ UU^+), \qquad (14)$$

$$K_3 = \frac{\sqrt{3}\,f_0}{2\sigma_0}\frac{\sigma}{[\mathrm{Tr}(UU^+)]^{1/2}}\mathrm{Tr}(\partial_\mu U\partial^\mu U^+), \text{ etc.}$$

in order $\mathcal{L}'$ to be of dimension 4 in agreement with (3) and (13) . In fact, there is an infinite number of possible derivative terms of dimension 4 from which K_1, K_2 and K_3 are the simplest ones (for more discussions, see /15/). The terms (14) are normalized so as to contain a "kinetic" term just in the form $[(\partial_\mu \tilde{S}_i)^2 + (\partial_\mu\phi_i)^2]$.

The guide criterion that we shall use in the construction of $\mathcal{L}'$ from K_i is the requirement that $\mathcal{L}'$ must give, first of all, the term (8) needed to solve the problem of large widths of the nonet of the ordinary scalar $q\bar{q}$ mesons S_i (i = 0, 1,..., 8). Moreover, we demand such an $\mathcal{L}'$ that does not change the correct kinetic term of the fields $\tilde{S}_i$ and ϕ_i in (12) obtained already from $(1/4)\mathrm{Tr}(\partial_\mu U\partial^\mu U^+)$ before adding $\mathcal{L}'$. These requirements together with the demand to use as simplest K_i from (14) as possible lead to the following $\mathcal{L}'$:

$$\mathcal{L}' = \frac{A}{2}\left[\frac{K_1 + K_2}{2} - K_3\right]. \qquad (15)$$

We easily see that, in fact, this simple form of $\mathcal{L}'$ gives the needed suppression term (8) and so (15) is sufficient for our purposes, i.e. to cure the large widths of S_i. Thus, having introduced the suppression effects through $\mathcal{L}'$(15) we deduce their consequences also for the gluonium coupling to the pseudoscalar mesons in what follows. Besides (8) we get:

$$\mathcal{L}^{der}_{\sigma\phi\phi} = -\frac{A}{\sigma_0}\tilde{\sigma}(x)(\partial_\mu\phi_i(x))^2, \qquad (16)$$

from (15) with $A = \gamma - 1$ from (8) and (9).

The potential V satisfying (13) gives the following squared mass sum rules

$$\sigma_0^2 M^2_{\sigma\sigma} - \frac{3}{2} f_0^2 M^2_{oo} = \frac{b}{2} G_0. \qquad (17a)$$

and

$$\sigma_0 M^2_{\sigma o} + \sqrt{\frac{3}{2}}\, f_0 M^2_{oo} = 0, \qquad (17b)$$

where M^2_{ij} are entries in the squared mass matrix for the $\tilde{\sigma}$ and $\tilde{S}_0$ fields. The nonderivative couplings $S\phi\phi$ are given as before in (6) while the nonderivative coupling $\sigma\phi\phi$ is proportional to $M^2_{\sigma o}$. In particular, from (12) and (13) we get:

$$\mathcal{L}^{nonder}_{S_0\phi\phi} = -\frac{1}{2}\sqrt{\frac{2}{3}}\frac{M^2_{oo}}{f_0}\tilde{S}_0(x)\phi_i^2(x) \qquad (18a)$$

in agreement with (6), and

$$\mathcal{L}^{nonder}_{\sigma\phi\phi} = -\frac{1}{2}\sqrt{\frac{2}{3}}\frac{M^2_{\sigma o}}{f_0}\tilde{\sigma}(x)\,\phi_i^2(x). \qquad (18b)$$

In the following we shall concentrate on the couplings (18) and their derivative counterparts (8) and (16). We shall rewrite them in more reliable forms in terms of the $SU(3)_F$ singlet physical mass eigenstates G, and ϵ defined as follows

$$G = \tilde{\sigma}\sin\theta + \tilde{S}_0\cos\theta, \qquad \epsilon = \tilde{\sigma}\cos\theta - \tilde{S}_0\sin\theta, \qquad (19)$$

where the mixing angle θ is given by:

$$\tan 2\theta = -\frac{2M^2_{\sigma o}}{M^2_{\sigma\sigma} - M^2_{oo}}. \qquad (20)$$

Using (19) and (20) the couplings (18) can be rewritten in the forms

$$\mathcal{L}^{nonder}_{\epsilon\phi\phi} = \frac{1}{2}\sqrt{\frac{2}{3}}\frac{M_\epsilon^2\sin\theta}{f_0}\epsilon(x)\phi_i^2(x), \qquad (21a)$$

and

$$\mathcal{L}^{\text{nonder}}_{G\phi\phi} = -\frac{1}{2}\sqrt{\frac{2}{3}}\,\frac{M_G^2\cos\theta}{f_o}\,G(x)\phi_i^2(x), \tag{21b}$$

where we have used the physical squared masses M_ϵ^2 and M_G^2 given as follows

$$M_\epsilon^2 = M^2_{\sigma\sigma}\cos^2\theta + M^2_{oo}\sin^2\theta - 2M^2_{\sigma o}\cos\theta\,\sin\theta,$$
$$M_G^2 = M^2_{\sigma\sigma}\sin^2\theta + M^2_{oo}\cos^2\theta + 2M^2_{\sigma o}\cos\theta\,\sin\theta. \tag{22}$$

The amplitudes for the $\epsilon\to\phi\phi$ and $G\to\phi\phi$ decays as obtained from (21) can equivalently be obtained from the following derivative couplings

$$-\sqrt{\frac{2}{3}}\,\frac{\sin\theta}{f_o}\,\epsilon(x)(\partial_\mu\phi_i(x))^2, \quad\text{and}\quad \sqrt{\frac{2}{3}}\,\frac{\cos\theta}{f_o}\,G(x)(\partial_\mu\phi_i(x))^2, \tag{23}$$

or, adding them together and using (19) we get (23) in the following compact form:

$$\sqrt{\frac{2}{3}}\,\frac{1}{f_o}\,\tilde{S}_o(x)(\partial_\mu\phi_i(x))^2. \tag{24}$$

We see that the nonderivative couplings (18) are equivalent to the derivative one (24) in such a sense that both (18) and (24) give the same correct amplitudes for the decays $\epsilon\to\phi\phi$ and $G\to\phi\phi$ of the physical states ϵ and G (19). Summing up (24) and (8) we obtain the complete $S_o\phi\phi$ derivative coupling as follows

$$\mathcal{L}^{\text{der}}_{S_o\phi\phi} = \sqrt{\frac{2}{3}}\,\frac{\gamma}{f_o}\,\tilde{S}_o(x)(\partial_\mu\phi_i(x))^2, \tag{25}$$

which is in agreement with (9) regardless of the presence of σ. Without going into details it is worth to note here that the results (9) and (16) can also be independently obtained from (12)-(15) within the general nonlinear sigma model approach /11a/, when instead of (4) U is parametrized nonlinearly, i.e., $U = \Pi\Sigma\Pi$ where $\Pi = \exp(i\lambda_j\phi_j/2f_\pi)$ and $\Sigma = \lambda_j S_j$ are the field matrices of the pseudoscalar Goldstone mesons ϕ_j and the scalars S_j, respectively.

With the use of (17b) the coupling $\sigma\phi\phi$ (18b) becomes equal to $M^2_{oo}/2\sigma_o$ and so, the corresponding decay $\sigma\to\phi\phi$ amplitude obtained from (18b) is expected to be reduced by the amplitude obtained from (16) if the masses of σ and S_o are comparable and if γ is small, $\gamma\ll 1$, i.e. $A\approx -1$, (10). In contrast to such an expectation, (9) and (16) may suggest that while the $S_i\to\phi\phi$ decays are suppressed when $\gamma\ll 1$ (10), at the same time the pure gluonium $\sigma-\phi\phi$ coupling (16) can be large for a heavy σ. We conclude that because of a probably nonnegligible mixing between $\tilde{\sigma}$ and $\tilde{S}_o$ (20) the amplitudes for the decays $\sigma\to\phi\phi$ and $S_o\to\phi\phi$ are not well defined quantities and any claims concerning them may not be reliable. Instead, we should investigate the correct couplings $\epsilon\phi\phi$ and $G\phi\phi$ of the physical states ϵ and G.

Using (16), (19) and (25) we get the complete $\epsilon\phi\phi$ and $G\phi\phi$ derivative couplings as follows

$$\mathcal{L}_{\epsilon\phi\phi} = g_{\epsilon\phi\phi}\,\epsilon(x)(\partial_\mu\phi_i(x))^2, \tag{26a}$$

and

$$\mathcal{L}_{G\phi\phi} = g_{G\phi\phi}\,G(x)(\partial_\mu\phi_i(x))^2, \tag{26b}$$

where

$$g_{\epsilon\phi\phi} = \frac{1-\gamma}{\sigma_o}\cos\theta - \sqrt{\frac{2}{3}}\,\frac{\gamma}{f_o}\sin\theta, \tag{27a}$$

and

$$g_{G\phi\phi} = \frac{1-\gamma}{\sigma_o}\sin\theta + \sqrt{\frac{2}{3}}\,\frac{\gamma}{f_o}\cos\theta. \tag{27b}$$

To make clear the connection of the present picture with our previous results (eq. (1)) on the gluonium decay, we shall analyze (27) in the limit of large squared mass M^2_{oo}. Labelling

$$t = \sqrt{\frac{3}{2}}\,\frac{f_o}{\sigma_o} \tag{28}$$

for given values of f_o and σ_o we get

$$\tan 2\theta = \frac{2t}{t^2 - 1 + \frac{b}{2}\,\frac{G_o}{\sigma_o^2}\,\frac{1}{M^2_{oo}}} \tag{29}$$

from (17) and (20). Solving (29) for $\tan\theta$ we obtain two solutions written as expansions in $G_o/M_{oo}^2\sigma_o^2$:

$$\tan\theta_1 = -t[1 + \frac{1}{t^2+1}\,\frac{b}{2}\,\frac{G_o}{\sigma_o^2}\,\frac{1}{M_{oo}^2} + O(\frac{G_o^2}{\sigma_o^4 M_{oo}^4})], \qquad (30)$$

and $\tan\theta_2 = -(\tan\theta_1)^{-1}$, i.e., $\sin\theta_2 = \pm\cos\theta_1$ and $\cos\theta_2 = \mp\sin\theta_1$. We see that use of θ_2 instead of θ_1 in (22) and (27) does not change the squared masses and decay widths of ϵ and G if they are interchanged simultaneously, i.e. $\epsilon \leftrightarrow G$ and thus we can put the angle $\theta = \theta_1$ (30) without loss of generality. With the value of θ (30) the squared masses M_ϵ^2 and M_G^2 (22) as well as the couplings $\epsilon\phi\phi$ and $G\phi\phi$ (27) can easily be evaluated within the $G_o/M_{oo}^2\sigma_o^2$ expansion. We get

$$M_\epsilon^2 = \frac{1}{t^2+1}\,\frac{b}{2}\,\frac{G_o}{\sigma_o^2}[1 + O(\frac{G_o}{\sigma_o^2 M_{oo}^2})], \qquad (31a)$$

and

$$M_G^2 = (t^2+1)M_{oo}^2 + \frac{t^2}{t^2+1}\,\frac{b}{2}\,\frac{G_o}{\sigma_o^2}[1 + O(\frac{G_o}{\sigma_o^2 M_{oo}^2})]. \qquad (31b)$$

The couplings (27) then become

$$g_{\epsilon\phi\phi} = \frac{1}{\sigma_o\sqrt{t^2+1}}\left\{1 + [\gamma(t^2+1) - t^2]\frac{M_\epsilon^2}{M_G^2} + O(\frac{M_\epsilon^4}{M_G^4})\right\}, \qquad (32a)$$

$$g_{G\phi\phi} = \frac{1}{\sigma_o t\sqrt{t^2+1}}\left\{[\gamma(t^2+1) - t^2] - t^2\frac{M_\epsilon^2}{M_G^2} + O(\frac{M_\epsilon^4}{M_G^4})\right\}. \qquad (32b)$$

We see that the coupling $\epsilon\phi\phi$ (32a) of the lighter particle ϵ does not depend on γ in the leading O(1) order of the expansion in $G_o/\sigma_o^2 M_{oo}^2$ (or, due to (31), in M_ϵ^2/M_G^2). The width $\Gamma(\epsilon\to\pi\pi)$ calculated in this approximation (i.e., neglecting $O(M_\epsilon^2/M_G^2)$ corrections) on the basis of (26a), (31) and (32a) is given just by (1) with m_σ replaced by M_ϵ. The difference between (1) and $\Gamma(\epsilon\to\pi\pi)$ occurs at the $O(M_\epsilon^2/M_G^2)$ level where also a dependence on γ appears. Thus, our previous result (eq. (1)) is a good approximation for the decay width $\Gamma(\epsilon\to\pi\pi)$ of the relatively light particle ϵ (when compared to its heavier companion G), i.e., if $M_\epsilon^2 \ll M_G^2$, and as expected on general grounds [11a] $\Gamma(\epsilon\to\pi\pi)$ becomes exactly (1) in the limit $M_G^2 \to \infty$ (or $M_{oo}^2 \to \infty$).

For given values of σ_o and f_o (and t, see (28)) the coupling $g_{\epsilon\phi\phi}$ (32a) depends on γ beginning from the next - to - leading order $O(M_\epsilon^2/M_G^2)$ while the coupling $g_{G\phi\phi}$ (32b) is γ-dependent already in the leading order O(1) of the M_ϵ^2/M_G^2 expansion. Remarkably, the coefficients of these dependences are equal to each other, i.e. to the factor $[\gamma(1+t^2)-t^2]$. Thus, the suppression of the $G\phi\phi$ coupling by requiring γ to obey the following equation

$$\gamma(t^2+1) - t^2 = 0, \qquad (33)$$

implies independence of $g_{\epsilon\phi\phi}$ of γ up to the order $O(M_\epsilon^2/M_G^2)$ and vice versa.

Since in the large N_c counting [7] $f_o \sim O(N_c^{1/2})$ and $\sigma_o \sim O(N_c)$, i.e. $t \sim O(1/N_c^{1/2})$ (see (28)), we have $\gamma \sim O(1/N_c)$ from (33) and thus we expect small γ, which is, in fact, in agreement with the fit (10). Then the $S_i\to\phi\phi$ and $G\to\phi\phi$ decay amplitudes calculated from (9), (26b) and (32b) have anomalous behaviour $O(1/N_c\sqrt{N_c})$ instead of $O(1/\sqrt{N_c})$ [6] as expected on general grounds [7] for the amplitude of the OZI allowed decay of a quarkonium meson into two $q\bar{q}$ mesons. The conventional $O(1/\sqrt{N_c})$ behaviour of (9) and (32b) would be realized if $\gamma = 1$ (or A = = 0), i.e., no $\mathcal{L}'$ term (15) is present in (12). However, the presence of $\mathcal{L}'$ in (12) (i.e. $A \neq 0$ and $\gamma \neq 1$ as given by (33)) provides the cancellation of the terms of the conventional order $O(1/\sqrt{N_c})$ in (9) and (32b), and results in the unusual, anomalous $O(1/N_c\sqrt{N_c})$ behaviour of the amplitudes for decays $S_i\to\phi\phi$ and $G\to\phi\phi$, while the coupling $\epsilon\phi\phi$ (32a) still remains of the order $O(1/N_c)$ as it should be for a gluonic state [6,7]. We mention here that a gluonium interpretation of the lighter state ϵ is also supported by the fact that just this particle almost dominates (2), and the neglected contribution again is $O(M_\epsilon^2/M_G^2)$. Thus, this suggests that the roles of the effective, physical $SU(3)_F$ singlet gluonium and quarkonium are played by the lighter ϵ and heavier G states, respectively.

3. HALF- AND- HALF MIXTURE OF THE PURE GLUONIUM AND QUARKONIUM?

One may question the applicability of the M_ϵ^2/M_G^2 expansion especially if difference between M_{oo}^2 and $M_{\sigma\sigma}^2$ is negligible, and so, we shall here analyze such a case separately, without using this expansion. In particular, we shall discuss the case based on the following assumption

$$M^2_{\sigma\sigma} = M^2_{oo} = M^2 \tag{34}$$

with M lying around 1.3 GeV in the interval 1.2 - 1.4 GeV. This assumption is interesting not only theoretically to verify the results of the previous section when mixing is large ($\theta = \pm 45^o$) but it might turn out to be approximately realized in the hadronic world.

On the one hand, (34) may be suggested by the quark model providing $M^2_{oo} = (2M^2(I=1/2) + M^2(I=1))/3$ [or, in a linear version, $M_{oo} = (2M(I=1/2) + M(I=1))/3$] with $M(I=1/2)$ and $M(I=1)$ being the masses of the I=1/2 and I=1 $q\bar{q}$ scalars, respectively (see also [28]). When these states are $K^*_0(1350)$ and $a_0(980)$ the quadratic formula predicts M_{oo}= 1.24 GeV while the linear one gives M_{oo}= 1.23 GeV, both in agreement with (34). If $a_0(980)$ is not a $q\bar{q}$ state and the I=1 $q\bar{q}$ scalar state is e.g. $a_0(1400)$[23] , then both the formulae suggest for M_{oo} a higher value, namely, M_{oo}= 1.37 GeV still, however, in coincidence with (34). Thus, regardless of the interpretation of $a_0(980)$ the value of M_{oo} probably lies in the interval 1.2 - 1.4 GeV (34). On the other hand, recent lattice QCD calculations [2] when taken (despite the existing reservations) seriously predict $M_{\sigma\sigma}$ to lie also around 1.3 GeV in the interval 1.2 - 1.4 GeV. The other independent estimates of the gluonium σ mass as based on the bags[3] , QCD sum rules [4a, 4b] , and others[5] are reasonably consistent with (34), too. On this basis we believe that the assumption (34) is not only plausible but may also be successful phenomenologically. Combining (17a), (28) and (34) we get

$$\frac{1}{t^2} = \frac{2}{3}\frac{\sigma_o^2}{f_o^2} = 1 + 3\frac{G_o}{M^2 f_o^2}, \tag{35}$$

where we have put b = 9 for $N_c = N_F = 3$. Having taken M = 1.2 - 1.4 GeV (34) we obtain $M^2 f_o^2 = 0.012-0.017\ GeV^4$ which coincides with the interval of the "known standard" values [12,16] of G_o, i.e., we have approximately $G_o = M^2 f_o^2$. Then (35) gives

$$t = \sqrt{\frac{3}{2}}\frac{f_o}{\sigma_o} = \frac{1}{2}, \tag{36}$$

where we have chosen the positive sign for t and, correspondingly, we choose $\theta = -45^o$ in order to label again the lighter state as ϵ (compare with (30)).

It is here worth mentioning a possible reliability of the present treatment. First of all, we have based our considerations on a particle dominance of, e.g., H in (2) and (3), but the assumption (34) that masses of scalars are O(1) GeV can make the relevance of the low-energy theorems like (2) for H much less certain. However, we are encouraged by the general success of analogous predictions based on vector meson dominance[29] even though vector mesons have also masses O(1) GeV. Thus, our results may be successful as well since another approximation that we have used, namely, the neglect of pseudoscalar meson masses is generally well controlled, too. Moreover, the reliability of our investigations depends on the validity of (34) as well as on the knowledge of the value of the gluon condensate G_o. We have argued above that the present "knowledge" of $M_{\sigma\sigma}$, M_{oo} and G_o also suggests our results may be successful.

From (17b), (22), (34) and (36) we get

$$M_\epsilon = \frac{1}{\sqrt{2}} M \quad \text{and} \quad M_G = \sqrt{\frac{3}{2}} M \tag{37}$$

in a good agreement with (31). Using in (37) M from the interval (1.2-1.4) GeV (34) one obtains the values of M_ϵ and M_G lying in the intervals (850-990) MeV and (1470-1710) MeV, respectively, the average values being M_ϵ = 920 MeV and M_G = = 1590 MeV for the average M = 1300 MeV (34). With (36) and $\theta = -45^o$ the couplings (27) become

$$g_{\epsilon\phi\phi} = \frac{1}{\sqrt{2}}\frac{1+\gamma}{\sigma_o} \tag{38a}$$

and

$$g_{G\phi\phi} = \frac{1}{\sqrt{2}}\frac{3\gamma - 1}{\sigma_o}. \tag{38b}$$

For t = 1/2 (36) and γ = 0.2-0.4 (10) the difference between the exact result (38) and the approximate one (32) is small and this again testifies to the applicability of the M^2_ϵ/M^2_G expansion with the reasonable value of expansion parameter $M^2_\epsilon/M^2_G = 1/3$ (37). Moreover, (32b) suggests that for $\gamma > 0.2$ and t = 1/2 (36) the coupling $g_{G\phi\phi}$ can even be more suppressed since small (due to the approximate validity of (33) in this case) leading order contribution $[\gamma(t^2+1) - t^2]$ is expected to cancel the next-to-leading order in M^2_ϵ/M^2_G contributions (see (32b)). This also agrees with (38b) showing the complete suppression of the $G\phi\phi$ coupling, i.e. $g_{G\phi\phi}= 0$, if $\gamma = 1/3$ while $g_{\epsilon\phi\phi}$ (38a) is still large and unsuppressed giving the width $\Gamma(\epsilon(920) \to \pi\pi) = 380$ MeV for the lighter state ϵ with the mass around 920 MeV.

Thus, consistently with (10), for $\gamma \leq 1/3$ the decays $G \to \pi\pi$ and $G \to K\bar{K}$ are strongly suppressed and this property of G together with the predicted mass $M_G \approx 1.59$ GeV forces us to identify G with the GAMS $f_0(1590)$ meson, recently discovered at the IHEP [18]. Although the state G is an approximate half-and-half mixture of the pure quarkonium S_0 and gluonium σ states, nevertheless, G plays the role of the effective physical quarkonium $SU(3)_F$ singlet state, as we have shown on the basis of the large N_c counting in the previous section. In this sense the scalar G is analogous to the approximately $SU(3)_F$ singlet pseudoscalar quarkonium η'. It is amusing to note that in both the cases the difference between the actual physical squared masses and the squared masses of the corresponding pure $q\bar{q}$ $SU(3)_F$ singlets as predicted by the quark model are large and approximately equal, i.e. $M_G^2 - M_{00}^2 \approx 0.85$ GeV^2 and $m_{\eta'}^2 - (2m_K^2 + m_\pi^2)/3 \approx 0.75$ GeV^2. The dominance of the $G \to \eta\eta$ and $G \to \eta'\eta$ decays [18] can then naturally be explained as an enhancement [17,30] of couplings between the scalar and pseudoscalar $SU(3)_F$ singlets due to the unsuppressed transitions between the quark and gluon degrees of freedom in 0^+ and 0^- channels [9]. We also note here that the present interpretation of $f_0(1590) \approx G$ may offer us a possibility to have the decay $J/\Psi \to \gamma f_0(1590)$ suppressed several times in comparison with the case [31] when $f_0(1590)$ is either a pure gluonium [30b] or a pure quarkonium [17]. Thus, we might get $BR(J/\Psi \to \gamma G) = O(10^{-4})$ as the experimental results [32] about the $J/\Psi \to \gamma\eta\eta$ and $J/\Psi \to \gamma\eta'\eta$ decays seem to indicate.

The existence of the state G implies, however, the existence of the lighter and very wide state ϵ with the mass M_ϵ below 1 GeV playing the role of the effective gluonium. The wide meson ϵ is not probably so easy to observe and, in fact, there are no such scalar mesons listed in recent issues of the particle data listings [19]. However, very recently Au et al. [21] have analyzed the AFS data [22] obtained at the CERN ISR "gluonium - search experiment" and they have claimed [21] to see a state $\epsilon(900)$ (besides other three states $S_1(991)$, $S_2(988)$, and $\epsilon(1.43)$) with the mass and width around 910 MeV and 350 MeV, respectively, in a good agreement with the predicted gluonium ϵ. If the $\epsilon(900)\pi\pi$ coupling were definitely by a factor of 2 larger than the $\epsilon(900)K\bar{K}$ one [21], then the gluonium interpretation of $\epsilon(900)$ may be in troubles, but as is seen from Table VI of ref. [21] such a result is inconclusive. For example, the solution K_1' in [21] suggests that these couplings are approximately equal to each other thus upholding a gluonium assignment for $\epsilon(900)$. So, the wide $\epsilon(900)$ meson [21] may not be inconsistent with the predicted gluonium ϵ, but any definite claims need a more precise experimental determination of its parameters.

We would also like to mention that there were already suggestions to interpret the old broad $\pi\pi$ s-wave state below 1 GeV as a gluonium. On the basis of the analysis in [20], Mennessier et al. [14a] have concluded, too, that such a state couples almost universally to $\pi\pi$ and $K\bar{K}$ as required for the gluonium. Novikov et al. [31] have suggested to search for a gluonium lying below 1 GeV in the radiative decay $J/\Psi \to \gamma\pi\pi$, but we are not very optimistic about the possibility of observing $J/\Psi \to \gamma\epsilon$ for a wide ϵ even if $BR(J/\Psi \to \gamma\epsilon) = O(10^{-3})$ [31].

On the other hand, Au et al. [21] have announced the narrow state $S_1(991)$ with the width of 21 MeV to be a gluonium candidate while other two states $\epsilon(900)$ and $S_2(988)$ are interpreted by these authors as the I=0 scalar $(1/2)^{1/2}(u\bar{u} + d\bar{d})$ and $s\bar{s}$ quarkonia, respectively, the interpretation being not trouble - free in the quark model. Moreover, unlike the case of a very wide scalar gluonium, the decay $J/\Psi \to \gamma S_1(991)$ could be rather restrictive for a narrow gluonium candidate $S_1(991)$ and the lack of such a decay may represent a serious problem for this state.

4. CONCLUSION

In the previous sections a picture of two $SU(3)_F$ singlet scalar states G and ϵ has been presented and compared with experiment. These states are half-and-half mixtures of the pure gluonium σ and pure quarkonium S_0. Also, we have shown that the heavier state G playing the role of the effective physical quarkonium is in good agreement with the GAMS $f_0(1590)$ meson [18] while the lighter ϵ being an effective gluonium is not probably inconsistent with otherwise inconclusive data on a broad $\pi\pi$ s-wave state $\epsilon(900)$ below 1 GeV [14a,21].

This suggests that the singlet-octet mixing for the $q\bar{q}$ scalar mesons is probably negligible, too, and the unmixed $q\bar{q}$ scalar octet members S_i $(i = 1,\dots,8)$ are approximately realized in the real world. Choosing the I=1/2 and I=1 members of this octet as $K_0^*(1350)$ and $a_0(980)$ (or $a_0(1400)$ [23] if $a_0(980)$ is not a $q\bar{q}$ state), and using the Gell-Mann-Okubo mass formula we easily find the mass M_8 of the state $S_8 \sim (1/6)^{1/2}$ $(u\bar{u}$ +

$+ d\bar{d} - 2 s\bar{s})$. The GMO mass formula used in both versions, quadratic and linear, predicts the masses $(M_8)_{quad}$ and $(M_8)_{lin}$, respectively, as follows

$$(M_8)_{quad} = 1.45 \text{ GeV}, \qquad (M_8)_{lin} = 1.47 \text{ GeV}, \tag{39}$$

when $a_0(980)$ is a $q\bar{q}$ state, or we have

$$(M_8)_{quad} \approx (M_8)_{lin} = 1.33 \text{ GeV}, \tag{40}$$

if $a_0(980)$ is not a $q\bar{q}$ state, but instead, a $q\bar{q}$ state is, e.g. $a_0(1400)$ [23]. The predictions (39) (or (40)) can be compared to the mass of the meson $f_0(1300)$ [19] which, unfortunately, is not known very precisely. The mass and width of $f_0(1300)$ vary from experiment to experiment and lie [19] in intervals (1.25-1.45) GeV and (150-400) MeV, respectively, and so we see a reasonable agreement with both (39) and (40), but we cannot conclude reliably which of the predictions (39) and (40) is satisfied better. It is amusing to note that the results of Au et al. [21] being in a better agreement with (39) may mildly uphold a $q\bar{q}$ assignment of $a_0(980)$. The effective coupling (9) gives the dominant decay rate of $S_8 \equiv f_0(1300)$ just into $\pi\pi$ and for γ from (10) the estimated values of the decay widths are also in a reasonable agreement with otherwise inconclusive experimental data [19] on the decays of $f_0(1300)$. For example, using $\gamma = 1/3$ as in the previous section and $M_8 = 1.3$ GeV, we estimate $\Gamma(S_8(1.3) \to \pi\pi) = 275$ MeV and $\Gamma(S_8(1.3) \to K\bar{K}) \doteq 60$ MeV from (9). The decay $S_8 \to \eta\eta$ is even more suppressed than the decay $S_8 \to K\bar{K}$ if the $\eta'\eta$ mixing is taken into account.

We see that the scenario for the $q\bar{q}$ scalar mesons presented here is reasonably consistent with experiment and these states are analogous to the pseudoscalar mesons with negligible singlet-octet mixings. The scalar $q\bar{q}$ octet mesons are $a_0(980)$ and/or $a_0(1400)$, $K_0^*(1350)$ and $f_0(1300)$, and are analogous to the pseudoscalars π, K and η while an analogue of η' is the scalar meson $f_0(1590)$. Although on the basis of the large N_c counting (section 2) the state $G \equiv f_0(1590)$ [18] is interpreted as an effective $SU(3)_F$ singlet quarkonium [17] nevertheless, being approximately a half-and-half mixture of pure gluonium and quarkonium degrees of freedom its actual nature is rather exotic, providing deviation from the quark-model prediction for G. In particular, there is a considerable difference between M_G^2 and M_{00}^2 like between m_η^2, and $(2m_K^2 + m_\pi^2)/3$. Another exotic state of our scenario is the effective scalar gluonium ϵ identified here with a rather hardly observable wide state $\epsilon(900)$ below 1 GeV.

Although reasonable both theoretically and experimentally, the present picture does not explain the $S^*(975)$ state (or more states? [21]) near the $K\bar{K}$ threshold as being made of $q\bar{q}$ or two gluons [14a,15,21,27,33], and thus we here support still other exotic, e.g. $q\bar{q}q\bar{q}$ [25], $K\bar{K}$-molecule [26], etc., explanations of the $K\bar{K}$-threshold effect. So, to have more complete description of the scalar mesons the introduction of such exotic states seems to be necessary. However, quite reasonable and rather successful description of gluonium and $q\bar{q}$ states presented here suggests that mixing with scalar $q^2\bar{q}^2$, etc., states is probably negligible.

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Ланик Й. E2-87-483
Широкий скалярный глюоний и кварк-антикварковый скалярный нонет

Обсуждаются связи скалярного глюония и кварк-антикварковых скалярных состояний нонета с псевдоскалярными мезонами на основе низкоэнергетических теорем нарушенной киральной симметрии и масштабной инвариантности с использованием феноменологических лагранжианов. Рассматривается смешивание между чистым глюонием σ и $SU(3)_F$-синглетным кварконием S_o. Беря для масс S_o и σ значения, получающиеся из оценок соответственно кварковой модели и недавних КХД расчетов на решетке и используя "стандартные" значения глюонного конденсата, мы предсказываем, что это смешивание в физических состояниях ε и G приблизительно составляет половина на половину. Предсказывается, что G своими свойствами соответствует ГАМС $f_o(1590)$ - мезону, а эффективный глюоний является широким состоянием лежащим ниже 1 ГэВ. Предлагается возможность последовательного описания всего скалярного $q\bar{q}$-нонета. Мезон $S^*(975)$ не входит в предложенную картину ни как $q\bar{q}$-, ни как gg-состояние, что дает возможность для его еще более экзотической интерпретации.

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Wide Scalar Gluonium and the Quark-Antiquark Scalar Nonet

The couplings of scalar gluonium as well as quark-antiquark scalar nonet states to the pseudoscalar mesons are discussed on the basis of the low-energy theorems of broken chiral symmetry and scale invariance implemented using phenomenological Lagrangians. Mixing between pure gluonium σ and the $SU(3)_F$ singlet quarkonium S_o is considered. Taking for the masses of S_o and σ the values based on estimates of the quark model and of recent QCD lattice calculations, respectively, and using the "standard" values of the gluon condensate, the mixture of S_o and σ in the physical states G and ε is predicted to be approximately half-and -half. We predict G to have properties consistent with the GAMS $f_o(1590)$ meson, while ε is predicted to be the wide effective gluonium state below 1 GeV. On this basis we suggest a possible consistent description of the whole scalar $q\bar{q}$ nonet. The picture contains $S^*(975)$ neither as $q\bar{q}$ nor as gg state and we support thus still more exotic interpretation of the $S^*(975)$ effect.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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