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IS SCALAR GLUONIUM OBSERVABLE ?

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ABSTRACT

We discuss couplings of scalar gluonium σ on the basis of the low energy theorems of broken chiral symmetry and the anomalous trace of the energy-momentum tensor, implemented using a phenomenological Lagrangian. Taking the ITEP value of the gluon condensate as input, we find $\Gamma(\sigma \rightarrow \pi\pi) \approx 0.6 \text{ GeV} \times (m_\sigma/1 \text{ GeV})^5$ and $\Gamma(\sigma \rightarrow \gamma\gamma) \approx 90 \text{ eV} \times (m_\sigma/1 \text{ GeV})^5$, while m_σ is undetermined. These results suggest that if the scalar gluonium mass is above 1 GeV, it is probably unobservably wide, while production in $\gamma\gamma$ collisions is probably too small to be detectable if $m_\sigma < 1.5 \text{ GeV}$. We comment on the observability of $J/\psi \rightarrow \sigma\gamma$ and on the relevance of our results to other gluonia.

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There is currently great theoretical [1,2] and experimental [3] interest in bound states of gluons, called gluonia or glueballs. Their existence is a clear prediction [1] of QCD, but predictions of their properties are less clear. Lattice QCD calculations [4] may now be approaching definite values for gluonium masses in pure SU(3) gauge theory without quarks, but the effects of including light quarks and the decay properties of glueballs are unknown. Some theorists suggest that gluonia are narrow resonances, while others think they should be wide [2]. There exist several experimental candidates [3] for gluonia in the $J^{PC} = 0^{++}$, 0^{-+} , and 2^{++} channels, some narrow and some wide, but they all have ambiguous interpretations.

In this paper we estimate the couplings of scalar gluonium on the basis of the low energy theorems of broken chiral symmetry and scale invariance. Both symmetries are broken by masses in a well-understood way, while scale invariance is also broken by quantum effects expressed in the anomalous trace of the energy-momentum tensor Θ_{μ}^{μ} [5]. This means that, unlike the pion, scalar gluonium is not expected to have a mass $\ll 1$ GeV, which makes the relevance of low energy theorems for Θ_{μ}^{μ} much less certain. However, we are encouraged by the general success of analogous predictions based on vector meson dominance [6], even though there is no nearby symmetry limit in which $m_{\rho, \omega, \phi} \rightarrow 0$, and the vector mesons have masses $O(1)$ GeV. If $m_{\sigma} = O(1)$ GeV, our results may be as successful. Although m_{σ} is unknown, knowledge [7,8] of the gluon condensate $G_0 \equiv \langle 0 | \alpha_s / \pi F_{\mu\nu} F^{\mu\nu} | 0 \rangle$ enables one to estimate as a function of m_{σ} the scalar analogue f_{σ} of the pion decay constant f_{π} :

$$f_{\sigma} \equiv \langle 0 | \Theta_{\mu}^{\mu} | 0 \rangle / m_{\sigma}^2 = - \frac{3}{m_{\sigma}} \sqrt{\frac{G_0}{2}} \quad (1)$$

As we will show, the $\sigma \rightarrow \pi\pi$, $K\bar{K}$, $\eta\eta$ and $\gamma\gamma$ couplings which are $\propto 1/f_{\sigma}$ can then be estimated using classical low energy theorems [9,10], yielding

$$\Gamma(\sigma \rightarrow \pi\pi) \approx 0.6 \text{ GeV} \times \left(\frac{m_{\sigma}}{1 \text{ GeV}} \right)^5, \quad \Gamma(\sigma \rightarrow \gamma\gamma) \approx 90 \text{ eV} \times \left(\frac{m_{\sigma}}{1 \text{ GeV}} \right)^5 \quad (2)$$

if one takes the ITEP [7] value of $G_0 = 0.012 \text{ GeV}^4$.

These results suggest that it will not be easy to observe scalar gluonium. If $m_\sigma \gtrsim 1 \text{ GeV}$, the σ decay width Γ_σ is probably so large that it is no longer a recognizable particle. Moreover, the $\gamma\gamma$ coupling (2) is so small that the σ is unlikely to be observable in $\gamma\gamma$ collisions. For comparison, $\Gamma(f \rightarrow \gamma\gamma) \approx 3 \text{ KeV}$ while upper limits on the $\gamma\gamma$ decay widths of narrow scalars vary between about 0.5 keV ($m_\sigma \lesssim 1 \text{ GeV}$, $\Gamma_\sigma \lesssim 50 \text{ MeV}$) and 1.4 keV ($m_\sigma \sim 1.3 \text{ to } 1.5 \text{ GeV}$, $\Gamma_\sigma \lesssim 200 \text{ MeV}$) [11].

We obtain these results using the ancient technology of phenomenological effective Lagrangians for broken chiral and scale invariance [12,13]. These were developed more than 10 years ago before the quark model and QCD were generally accepted, and *a fortiori* before the anomalous trace of the energy-momentum tensor in QCD:

$$\Theta^\mu_\mu \supset (\Theta^\mu_\mu)_{an} \equiv \frac{\beta}{2g} F^{(a)}_{\mu\nu} F^{(a)\mu\nu} : a=1,2,\dots,8 \quad (3)$$

was derived [5]. Indeed, the derivation of the result (3) was a sequel to the canonical anomaly calculation of $\sigma \rightarrow \gamma\gamma$ [10,14]. At that time, the existence of a large $\sigma \rightarrow \pi\pi$ coupling was also deduced [9,10], but the σ was identified with the conjectural $\epsilon(700) 0^{++}$ meson. At that time before QCD, *a fortiori* no one had thought about glueballs. Now one can make the more cogent identification of the σ as gluonium, an effective Lagrangian treatment of the trace anomaly (3) is available [15,16], and f_σ can be related (1) to another "known" quantity, the gluon condensate [7,8]. We first review the phenomenological Lagrangian framework incorporating these improvements [17] and chiral symmetry breaking, and then derive and discuss our results (2).

In the first efforts to formulate not only broken chiral but also scale invariance within an effective Lagrangian approach, Salam and Strathdee [12] showed that to arrange the scale properties of effective Lagrangians in a suitable way one has to assume the existence of a dimensionful and chirally invariant scalar

field σ (dimension $d_\sigma = 1$). The field σ can be parametrized as follows

$$\sigma(x) = \sigma_0 \exp\left(\frac{\tilde{\sigma}(x)}{\sigma_0}\right) \quad (4)$$

where $\sigma_0 = \langle 0|\sigma|0\rangle = -f_\sigma$ and $\tilde{\sigma}(x)$ is a physical interpolating field which transforms as $\tilde{\sigma}(x) \rightarrow \tilde{\sigma}(x) - \sigma_0 \ln \rho$ if $x \rightarrow \rho x$. In QCD this σ -particle can be interpreted as scalar gluonium. This is seen most easily in the case of pure gluodynamics if one constructs an effective Lagrangian containing just this field that describes correctly the QCD trace anomaly (3), and satisfies the low-energy theorems of ref. [18]:

$$\left. \begin{aligned} i \int d^4x \langle 0|T(H(x)H(0))|0\rangle &= 4H_0, \\ i^2 \int d^4x \int d^4y \langle 0|T(H(x)H(y)H(0))|0\rangle &= 16H_0, \text{ etc.} \end{aligned} \right\} \quad (5)$$

where the scalar gluonic current $H(x) \equiv -\beta(g)/2g F_{\mu\nu}^{(a)} F^{(a)\mu\nu}$ and $H_0 \equiv \langle 0|H|0\rangle$.

i) Pure gluodynamics

Although the effective Lagrangian has already been found [15,16] in this case, in what follows the rederivation of this Lagrangian will be instructive and useful. Let the effective Lagrangian containing just one scalar field $\sigma(x)$ [e.g. (4)] have the general form:

$$\mathcal{L}_\sigma = \frac{1}{2} (\partial_\mu \sigma)^2 - V(\sigma) \quad (6)$$

where $V(\sigma)$ is a potential to be determined. This Lagrangian gives the following trace of the "improved" energy-momentum tensor [19]

$$\left(\Theta_\mu^\mu\right)_{an} \equiv -H = 4V(\sigma) - \sigma \frac{\partial V(\sigma)}{\partial \sigma} \quad (7)$$

Expanding the potential $V(\sigma)$ in the field $\tilde{\sigma}$, i.e.,

$$V(\sigma) = V(\sigma_0) + \frac{1}{2} \langle 0|\frac{\partial^2 V}{\partial \tilde{\sigma}^2}|0\rangle \tilde{\sigma}^2 + \dots \quad (8)$$

where we have already put $\langle 0 | dV/d\tilde{\sigma} | 0 \rangle = 0$ in order to eliminate the term linear in $\tilde{\sigma}$, then eq. (7) reads

$$H(x) = H_0 + H_1 \tilde{\sigma}(x) + H_2 \tilde{\sigma}^2(x) + \dots \quad (9)$$

where

$$\left. \begin{aligned} H_0 &= - \langle 0 | \frac{\beta(g)}{2g} F^2 | 0 \rangle = - 4V(\sigma_0) \\ H_1 &= \sigma_0 \langle 0 | \frac{\partial^2 V}{\partial \tilde{\sigma}^2} | 0 \rangle = \sigma_0 m_\sigma^2, \text{ etc.} \end{aligned} \right\} \quad (10)$$

Combining eq. (9) and the first of eqs. (5), one gets

$$H_1^2 = 4m_\sigma^2 H_0 \quad (11)$$

Analogously, eqs. (5) can be used to calculate all the coefficients H_i ($i = 1, 2, \dots$) in terms of m_σ and H_0 . As a result one finds the following exact relation

$$H(x) = H_0 \exp \left[\frac{2m_\sigma}{\sqrt{H_0}} \tilde{\sigma}(x) \right] \quad (12a)$$

or, using parametrization (4), we have

$$H(x) = \frac{m_\sigma^4}{16 H_0} [\sigma(x)]^4 \quad (12b)$$

We see that the low-energy theorems (5) represent the fact that the dimension of $H(x)$ must be 4 as is needed for the trace anomaly in pure gluodynamics. Since the σ -particle dominates the scalar gluonic current [eqs. (5) and (12)] this particle must be identified as scalar gluonium. Combining (7) and (12b), and solving this differential equation, we get [15]

$$V(\sigma) = \frac{m_\sigma^4}{16 H_0} \sigma^4 \ln \sigma/C \quad (13a)$$

where C is a constant of integration which is fixed by eqs. (10) and (11) as follows: $4 \ln (\sigma_0/C) = -1$. Then eqs. (4) and (10) give $V(\sigma)$ in the useful form [16]

$$V(\sigma) = H_0 \left(-\frac{1}{4} + \frac{\tilde{\sigma}(x)}{\sigma_0} \right) \exp \left(\frac{4\tilde{\sigma}(x)}{\sigma_0} \right) \quad (13b)$$

It is clear that the logarithmic term breaks scale invariance and the potential (13) is necessary and sufficient to reproduce the effects of the trace anomaly of QCD.

ii) Chiral symmetry limit

Here we want to enlarge the Lagrangian (6) so as to incorporate the octet of Goldstone pseudoscalar mesons present in the limit of massless quarks. This inclusion changes the β -function in eq. (3) but the dimension of the scalar gluonic current $H(x)$ remains 4. From the previous exercise of pure gluodynamics we see that it is natural to assume that any effective Lagrangian describing the system of scalar gluonium and the octet of pseudoscalar Goldstone mesons should also lead to a trace anomaly of dimension 4. Moreover, in order to reproduce the important low-energy theorem of refs. [10,20,21], i.e.,

$$\langle \pi^+(p_1), \pi^-(p_2) | \Theta_\mu^\mu | 0 \rangle \Big|_{\text{chiral limit}} = q^2 \quad (14)$$

q^2 being the invariant (mass)² of the $\pi^+\pi^-$ system, it is convenient [9] to start from the conventional effective Lagrangian [13] of the pseudoscalar Goldstone boson system as follows

$$\mathcal{L}_Q = \frac{1}{4} \text{Tr} \left[(\partial_\mu U)(\partial^\mu U^\dagger) \right] \quad (15)$$

where $U(x)$ is parametrized as the unitary matrix:

$$U(x) = \exp \left(i \sum_{j=1}^8 \frac{\lambda_j \phi_j(x)}{f_\pi} \right) \quad (16)$$

Here f_π is the pion decay constant ($f_\pi = 93$ MeV), the ϕ_i 's ($i = 1, \dots, 8$) are fields for the pseudoscalar Goldstone mesons and the Gell-Mann λ matrices are normalized to $\text{Tr}(\lambda_i \lambda_j) = 2 \delta_{ij}$. Because the dimension of the field U is zero [9], the Lagrangian (15) gives a non-zero dimension 2 contribution to the effective trace anomaly [17,22]. In fact, this represents a very strong operator form of

eq. (14), and as such it forbids the coupling of the σ -particle to pseudoscalar Goldstone mesons in the chiral limit. In this case, a $\sigma\pi\pi$ coupling would be possible only through a quark-mass term, and it is not clear whether this unnatural coupling of gluonium to pseudoscalar mesons is realized in the hadronic world [23]. We believe not.

Instead, as we have already mentioned, on the basis of our treatment of pure gluodynamics, it is more natural to assume a Lagrangian that reproduces the trace anomaly of dimension 4. Thus, we shall multiply $\mathcal{L}_Q$ (15) by $[\sigma(x)/\sigma_0]^2$ so that it has dimension 4, and suppose that the effective Lagrangian describing the system of gluonium and pseudoscalar Goldstone mesons is of the form [9]

$$\mathcal{L}_{\text{qq}} = \frac{1}{2} (\partial_\mu \sigma)^2 + \frac{1}{4} \left(\frac{\sigma(x)}{\sigma_0} \right)^2 \text{Tr}[(\partial_\mu U)(\partial^\mu U^\dagger)] - V(\sigma) \quad (17)$$

where σ , $V(\sigma)$ and U are given by eqs. (4), (13), and (16) with the corresponding modification of β -function due to the presence of three massless quarks. This Lagrangian gives

$$\mathcal{L}_{\sigma\pi\pi}(x) = \frac{\tilde{\sigma}(x)}{\sigma_0} \sum_{i=1}^8 (\partial_\mu \phi_i(x))^2 \quad (18)$$

which is the coupling between scalar gluonium σ and a pair of pseudoscalar Goldstone mesons. The interaction term (18) allows us to reproduce the low-energy theorem just in its matrix-element form (14). From eqs. (18), (10), and (11) we obtain, e.g. for the $\sigma \rightarrow \pi\pi$ decay, the following partial width for $m_\pi = 0$:

$$\Gamma(\sigma \rightarrow \pi\pi) = \frac{m_\sigma}{48\pi} \frac{m_\sigma^4}{\langle 0 | \frac{\alpha_s}{\pi} F^2 | 0 \rangle} \quad (19)$$

where we have used [for the $SU(3)_c$ colour group and for three flavours, $N_F = 3$] the approximate relation

$$H_0 = - \langle 0 | \frac{\beta(g)}{2g} F^2 | 0 \rangle = \frac{9}{8} G_0 \quad (20)$$

where $G_0 = \langle 0 | (\alpha_s/\pi) F^2 | 0 \rangle$ is the familiar gluon-condensate term [7] parametrizing non-perturbative effects in QCD. Note that equation (19) is $O(1/N_c^2)$ for large N_c ,

as expected on general grounds.

iii) The explicit breaking of chiral and scale invariance

Here we shall discuss possible corrections due to the quark mass term in the QCD Lagrangian. For non-zero values of the quark masses the trace of the energy-momentum tensor of QCD contains besides the anomalous contribution (3) also an extra term:

$$(\Theta^\mu_\mu)_{\text{QCD}} = \frac{\beta(g)}{2g} F_{\mu\nu}^{(a)} F^{(a)\mu\nu} - \mathcal{L}_{\text{SB}}^{\text{QCD}} \quad (21)$$

where the quark mass term $\mathcal{L}_{\text{SB}}^{\text{QCD}} = - \sum_i m_{q_i} \bar{q}_i q_i$, i being a flavour index. It has been shown [9,17,22] that consistency with eq. (21) requires the following form of the effective quark mass term

$$\mathcal{L}_{\text{SB}}(\chi) = - [\sigma(\chi)]^3 \text{Tr} [M(U(\chi) + U^\dagger(\chi))] \quad (22)$$

where M is proportional to the (real diagonal) light quark mass matrix and we do not discuss effects due to the anomalous dimension of the $\bar{q}_i q_i$ operator. In this case, to compensate the term linear in $\tilde{\sigma}$ in eq. (22) we must add to eq. (22) a piece of the form

$$\mathcal{L}_{\text{add}}(\chi) = - \frac{3}{8} \sigma_0^4 f_\pi^2 (2m_K^2 + m_\pi^2) [\sigma(\chi)]^4 \quad (23)$$

which is the only possibility consistent with eq. (21). Thus, the complete effective Lagrangian is given as

$$\mathcal{L} = \mathcal{L}_{\text{QCD}} + \mathcal{L}_{\text{SB}} + \mathcal{L}_{\text{add}} \quad (24)$$

where $\mathcal{L}_{\text{QCD}}$, $\mathcal{L}_{\text{SB}}$ and $\mathcal{L}_{\text{add}}$ are given by eqs. (17), (22), and (23).

It is now an easy exercise to evaluate corrections due to non-zero values of the light quark masses. We obtain, for example, that m_σ^2 should be corrected by the factor $[1 + (3/4H_0) f_\pi^2 (m_K^2 + m_\pi^2/2)]$ while the width $\Gamma(\sigma \rightarrow \pi\pi)$ [eq. (19)] has to be corrected by the factor $[1 + m_\pi^2/m_\sigma^2]^2$. These corrections are small and can be neglected along with the phase-space factor $\sqrt{1 - 4m_\pi^2/m_\sigma^2}$, so that eq. (19) is a good approximation to $\Gamma(\sigma \rightarrow \pi\pi)$, also for the physical value of m_π . It is worth noting here that because of small corrections due to the presence of the quark

mass term the present model gives couplings between scalar gluonium and the members of the pseudoscalar meson octet which are almost SU(3) invariant.

We now discuss the phenomenological implications of the scalar gluonium couplings inferred from the effective Lagrangian (24).

Equation (19) takes the following numerical values for popular values of $\langle 0 | (\alpha_s / \pi) F^2 | 0 \rangle$

$$\Gamma(\sigma \rightarrow \pi\pi) = \left(\frac{m_\sigma}{1 \text{ GeV}} \right)^5 \times \begin{cases} 550 \text{ MeV} \\ 220 \text{ MeV} \end{cases} \text{ for } \langle 0 | \frac{\alpha_s}{\pi} F^2 | 0 \rangle = \begin{cases} 0.012 \text{ GeV}^4 \text{ (case A)} \\ 0.03 \text{ GeV}^4 \text{ (case B)} \end{cases} \quad (25)$$

Equations (19) and (25) give a σ width increasing rapidly with increasing m_σ . If it is light enough, scalar gluonium might have escaped detection [24] by being very narrow:

$$\Gamma_\sigma \approx \Gamma_{\sigma \rightarrow \pi\pi} < 2 \text{ MeV for } m_\sigma < \begin{cases} 330 \text{ MeV (case A)} \\ 390 \text{ MeV (case B)} \end{cases} \quad (26)$$

On the other hand, if it is heavy enough, scalar gluonium could escape detection by being very wide:

$$\Gamma(\sigma \rightarrow \pi\pi) > \frac{1}{2} m_\sigma \text{ for } m_\sigma > \begin{cases} 980 \text{ MeV (case A)} \\ 1230 \text{ MeV (case B)} \end{cases} \quad (27)$$

In fact, a scalar gluonium with $m_\sigma > 1 \text{ GeV}$ would have additional decay modes into $K\bar{K}$ and $\eta\eta$ available, with couplings which are approximately SU(3) symmetric.

Hence Γ_σ would be even larger than (25). On the basis of equations (25) to (27) we believe that present data [24] already rule out the following range of scalar gluonium masses:

$$m_\sigma = 400 \text{ to } 1000 \text{ MeV excluded.} \quad (28)$$

We doubt that m_σ could be smaller than the range (28), and therefore believe that scalar gluonium must have a mass larger than 1 GeV, and hence is probably too wide to be recognized as a well-defined resonance. In this case our results could only be qualitatively correct, at best. It is important to bear in mind that our

results are based on single particle dominance of the anomalous trace of the energy-momentum tensor, which can be queried for a wide resonance. Moreover, mixing with scalar $\bar{q}q$ or $\bar{q}q\bar{q}q$ mesons [24] could be important, and higher mass excited gluonium states may also not be negligible. Their inclusion might tend to decrease our estimate (1) of f_σ and hence increase the glueball decay width. We note in passing that a QCD sum rule estimate [25] of f_σ is smaller than we estimate on the basis of (1), and hence would also correspond to a larger value of $\Gamma(\sigma \rightarrow \pi\pi)$. Nevertheless, it is interesting to note that our estimate (25) is qualitatively consistent with the estimated width of the $\epsilon(1300)$ meson [24].

We can also discuss the $\sigma \rightarrow \gamma\gamma$ coupling using the canonical trace anomaly [14], which relates it to $R \equiv \sigma(e^+e^- \rightarrow \gamma^* \rightarrow q\bar{q})/\sigma(e^+e^- \rightarrow \gamma^* \rightarrow \mu^+\mu^-)$ and has been proven [5] true in QCD:

$$\Gamma(\sigma \rightarrow \gamma\gamma) = \left(\frac{\alpha R}{8\pi^2 f_\sigma} \right)^2 \frac{m_\sigma^3}{16\pi} \quad (29)$$

Using (1) this becomes

$$\Gamma(\sigma \rightarrow \gamma\gamma) = \left(\frac{R\alpha}{8\pi^2} \right)^2 \frac{m_\sigma^5}{64\pi \langle 0 | \frac{\alpha_s}{\pi} F^2 | 0 \rangle} \quad (30)$$

Taking $R = 5$ for $q = u, d, s, c, b, t$ we find

$$\Gamma(\sigma \rightarrow \gamma\gamma) = \left(\frac{m_\sigma}{1 \text{ GeV}} \right)^5 \times \begin{cases} 90 \text{ eV} & (\text{case A}) \\ 35 \text{ eV} & (\text{case B}) \end{cases} \quad (31)$$

This decay rate is much smaller than those characteristic of conventional $q\bar{q}$ resonances [11] such as $\Gamma(f \rightarrow \gamma\gamma) \approx 3 \text{ keV}$. It is also considerably smaller than the experimental upper limits [1] on $\gamma\gamma$ couplings of scalar resonances: $\Gamma(\sigma \rightarrow \gamma\gamma) < 0(0.5) \text{ keV}$ for light σ with $\Gamma_\sigma < 50 \text{ MeV}$ [relevant in our case to $m_\sigma < 620 \text{ MeV}$ (case A) or 740 MeV (case B)], while the limit $\Gamma(\sigma \rightarrow \gamma\gamma) < 1.4 \text{ keV}$ for a resonance between 1.3 and 1.5 GeV is presumably irrelevant if scalar gluonium is as wide (25) as suggested here. We therefore believe scalar gluonium is essentially unobservable in this channel.

Finally, we make some comments about the $J/\psi \rightarrow \sigma\gamma$ coupling. We assume that scalar gluonium couples to c quarks through the anomalous part $(\Theta_{\mu}^{\mu})_{\text{an}}$ of the trace (3) of the energy-momentum tensor. The $(\Theta_{\mu}^{\mu})_{\text{an}} - c\bar{c}$ vertex can be calculated via a gluon loop diagram. Using the coupling (1) of σ to $(\Theta_{\mu}^{\mu})_{\text{an}}$ one estimates

$$g_{\sigma c\bar{c}} = O\left(\frac{\alpha_s^2}{2\pi^2}\right) \frac{m_{\sigma}}{f_{\sigma}} \quad (32)$$

with the numerical value of the coefficient depending on the ratio $m_{\sigma}/m_c = O(1)$. Since the typical momenta in the gluon loop are $O(1)$ GeV, we are in a non-perturbative régime, so it is not clear the calculation makes any sense, and even if it did we would not know what value of α_s to use in equation (32). If we ignore all these problems and take $\alpha_s = O(0.5)$ for purposes of orientation, we find the branching ratio

$$B(J/\psi \rightarrow \sigma + \gamma) \equiv \frac{\Gamma(J/\psi \rightarrow \sigma + \gamma)}{\Gamma(J/\psi \rightarrow \text{all})} = O(1) \quad (33)$$

for $m_{\sigma} = 1.2$ GeV. This is to be compared with the experimental branching ratio $B(J/\psi \rightarrow f + \gamma) \approx 1.5 \times 10^{-3}$. Despite the fact that our estimate (33) is above this value, we are not optimistic about the possibility of observing $J/\psi \rightarrow \sigma + \gamma$, because a σ in this mass range is expected to be very wide. Moreover, this calculation involves more (and more unreliable?) assumptions than our previous results. For example, varying α_s from 0.2 to 1 varies $B(J/\psi \rightarrow \sigma + \gamma)$ in a range $O(600)$ around the estimate (33). Nevertheless, the expected large decay width (25) poses a severe problem even if our value (33) of $B(J/\psi \rightarrow \sigma + \gamma)$ is an underestimate.

Is there any reason to think that our pessimistic conclusions about scalar gluonium could also apply to gluonia with other spin-parities? Our basic theoretical assumption of scalar gluonium dominance of the anomalous trace of the energy-momentum tensor is clearly questionable, and highly specific to the 0^{++} gluonium channel. However, it is not possible to exclude the possibility that other gluonia might also be very wide, have couplings which are almost $SU(3)$ invariant, and have small $\gamma\gamma$ couplings. We have listed these properties in the

sequence of increasing confidence. We expect that gluonia would all have small $\gamma\gamma$ couplings because they do not contain charged constituents. There are no completely compelling arguments [2] for suppressing or enhancing gluonium couplings to strange particles relative to non-strange particles. However, our estimates of Γ_σ rely entirely on our specific theoretical assumptions, and have no good reason to apply to gluonia with different spin-parities. Therefore, even if we discourage hunters for scalar gluonium, we have no intention of discouraging searches for other gluonia.

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