

Absence of gluonic components in axial and tensor mesons

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Abstract. A quarkonium–gluonium mixing scheme previously developed to describe the characteristic of the pseudoscalar mesons is applied to axial and tensor mesons. The parameters of the model are determined by fitting the eigenvalues of a mass matrix. The corresponding eigenvectors give the proportion of light quarks, strange quarks and glueball in each meson. However, the predictions of the model for the branching ratios and electromagnetic decays are incompatible with the experimental results. These results suggest the absence of gluonic components in the states of axial and tensor isosinglet mesons analyzed here.

1 Introduction

The existence of gluon self-coupling in QCD gives rise to the possibility of glueball formation. These states may have the same quantum numbers as those of some quarkonia. A signature of gluonium states is that they have no place in the mesons nonets. If a glueball and one or more quarkonia with the same quantum numbers have nearly the same masses these states may interfere and new states are formed. Thus the physical states formed by gluonia and quarkonia interference needs a mixing scheme to describe them. Several kinds of the mixing schemes have been proposed to give account of the peculiar properties of these mesons.

In some schemes the physical states are written as linear combinations of pure quarkonia and gluonia states. The linear coefficients are generally related to the rotation angles and may be determined by the decay properties of, or into, the physical mesons [1–3].

Another approach, in which the interference is considered at a more fundamental level, consists in writing a mass matrix for the physical states in the basis of the pure quarkonia and gluonia states. The elements of this mass matrix are obtained from a model that describes the process of interference. The mixtures of the basic states are induced by the off-diagonal elements. Thus, these elements must contain the amplitudes for transitions from one to another states of the basis. The eigenvalues of that matrix give the masses of the physical states and the corresponding eigenvectors give the proportion of quarkonia and gluonia in each meson [4–9].

In a previous paper we presented a mixing scheme for the pseudoscalar mesons, based on a mass matrix approach. The flavor-dependent annihilation amplitudes and binding energies are the responsible mechanisms for the quarkonium–gluonium mixing. The properties of the three lowest energy states of the pseudoscalar isosinglet mesons $\eta(547)$, $\eta'(958)$ and $\eta''(1410)$ are well described by a model based on the assumption that these states are mixtures of the light quarks, strange quarks and a glueball [10].

The nonet of axial ($1^{++}, 1^3P_1$) and tensor ($2^{++}, 1^3P_2$) mesons are well established [11]. The axial nonet consists of the isodoublet $K_{1A}(1340)$, the isovector $a_1(1260)$ and the isoscalars $f_1(1285)$ and $f_1(1510)$. The K_{1A} is a mixture of $K_1(1270)$ and $K_1(1470)$ with a close to 45° mixing angle [12]. The tensor nonet is formed by the isodoublet $K_2^*(1430)$, the isovector $a_2(1320)$ and the isoscalars $f_2(1270)$ and $f_2'(1525)$. Nonetheless, there are extra isoscalar states with quantum numbers and masses permitting that they can be interpreted as partners of the nonets of axial and tensor mesons. The axial state $f_1(1420)$, observed in two experiments [13], has been considered by some authors [14] as a possible candidate to be exotic. On the other side, there are two candidates to be exotic tensor states: $f_2(1640)$ [15] and $f_J(1710)$ [16]. There is a controversy about the value of the spin of the $f_J(1710)$: it may be a scalar or a tensor state [17].

In the present paper the candidates to the exotics $f_1(1420)$ and $f_2(1640)$, or $f_2(1710)$, are supposed to be components of quarkonium–gluonium mixing schemes similar to that previously applied to the pseudoscalar mesons [10]. The same mixing scheme is not applied to the scalar states because only the assignment for the scalar isodoublet is well established.

This paper is organized as follows: Sect. 2 outlines a brief review of the matrix formalism used formerly for the pseudoscalar mesons. We also fix the notation that will

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be used in the subsequent sections. Section 3 is devoted to the application of the mass matrix formalism for the three lowest energy states of the axial mesons. Afterwards, in Sect. 4, two different mixing configurations for the tensor isosinglet mesons are considered. In both sections the results obtained from the mass matrix formalism are used for calculating some quantities related to branching ratios and decay widths. Finally, in the conclusion in Sect. 5, the results obtained are analyzed and confronted with those ones presented in the literature.

2 The mass matrix formalism

The mass matrix in the basis $|u\bar{u}\rangle$, $|d\bar{d}\rangle$, $|s\bar{s}\rangle$ and $|gg\rangle$, including flavor-dependent binding energies and annihilation amplitudes, has matrix elements given by

$$M_{ij} = (2m_i + E_{ij})\delta_{ij} + A_{ij}, \quad (1)$$

where $i, j = u, d, s, g$. The contribution to the elements of the mass matrix are the rest masses of the quarks and the gluon, the eigenvalues E_{ij} of the Hamiltonian for the stationary bound state (ij) and the amplitudes A_{ij} , that account for the possibility of $q\bar{q} \leftrightarrow gg \leftrightarrow q'\bar{q}'$ and $q\bar{q} \leftrightarrow gg$ transitions. As in the previous paper we assume that E_{ij} and A_{ij} are not SU(3)-invariant quantities. Two other bases will be used in this paper. The first basis consists of the isoscalar singlet and octet of SU(3),

$$|1\rangle = \frac{1}{\sqrt{3}} (|u\bar{u}\rangle + |d\bar{d}\rangle + |s\bar{s}\rangle), \quad (2)$$

$$|8\rangle = \frac{1}{\sqrt{6}} (|u\bar{u}\rangle + |d\bar{d}\rangle - 2|s\bar{s}\rangle). \quad (3)$$

The second basis is chosen assuming a segregation of the strange and the nonstrange quarks,

$$|N\rangle = \frac{1}{\sqrt{2}} (|u\bar{u}\rangle + |d\bar{d}\rangle), \quad (4)$$

$$|S\rangle = |s\bar{s}\rangle. \quad (5)$$

Besides these states we need also the gluonium and the isovector states,

$$|G\rangle = |gg\rangle, \quad (6)$$

$$|\tilde{\pi}^0\rangle = \frac{1}{\sqrt{2}} (|u\bar{u}\rangle - |d\bar{d}\rangle). \quad (7)$$

In these bases the mixing among the isoscalar and isovector states is caused by isospin symmetry breaking terms. Therefore, assuming the exact SU(2)-flavor symmetry, one needs only consider the subspace spanned by the isoscalar states, when the mass matrix reduces to a 3×3 matrix M_0 .

The invariants of the mass matrix under a unitary transformation give the following mass relations for the isoscalar physical states:

$$m_1 + m_2 + m_3 = \text{tr}(M_0), \quad (8)$$

$$m_1 \cdot m_2 \cdot m_3 = \det(M_0), \quad (9)$$

$$m_1 \cdot m_2 + m_1 \cdot m_3 + m_2 \cdot m_3 = \frac{1}{2} \left[(\text{tr}(M_0))^2 - \text{tr}(M_0^2) \right], \quad (10)$$

where M_0 is the 3×3 mass matrix for the isoscalar states, and m_i ($i = 1, 2, 3$) are their eigenvalues.

The eigenvectors of the mass matrix M_0 are the physical states $|m_1\rangle$, $|m_2\rangle$ and $|m_3\rangle$ which are mixtures of $|1\rangle$, $|8\rangle$ and $|G\rangle$:

$$|m_1\rangle = -c_2 s_1 |1\rangle + c_1 c_2 |8\rangle - s_2 |G\rangle, \quad (11)$$

$$|m_2\rangle = (c_1 c_3 + s_1 s_2 s_3) |1\rangle + (c_3 s_1 - c_1 s_2 s_3) |8\rangle - c_2 s_3 |G\rangle, \quad (12)$$

$$|m_3\rangle = (c_1 s_3 - s_1 s_2 c_3) |1\rangle + (c_1 s_2 c_3 + s_1 s_3) |8\rangle + c_2 c_3 |G\rangle. \quad (13)$$

The coefficients of the eigenvectors are written in terms of three Euler angles defining a rotation in a three dimensional space. For brevity, we have defined the notation $c_i \equiv \cos \theta_i$ and $s_i \equiv \sin \theta_i$ ($i = 1, 2, 3$).

The eigenvectors (11)–(13) can also be rewritten in the basis $|N\rangle$, $|S\rangle$ and $|G\rangle$:

$$|m_1\rangle = X_1 |N\rangle + Y_1 |S\rangle + Z_1 |G\rangle, \quad (14)$$

$$|m_2\rangle = X_2 |N\rangle + Y_2 |S\rangle + Z_2 |G\rangle, \quad (15)$$

$$|m_3\rangle = X_3 |N\rangle + Y_3 |S\rangle + Z_3 |G\rangle. \quad (16)$$

We adopt an expression for the amplitude of the process $q\bar{q} \leftrightarrow gg \leftrightarrow q'\bar{q}'$ similar to that of Cohen and Lipkin [19] and Isgur [20], where the numerator of the two-gluon annihilation amplitude expression is assumed to be a SU(3)-invariant parameter, which means that we parameterize the annihilation amplitude in the form

$$A_{qq'} = \frac{\Lambda}{m_q m_{q'}}. \quad (17)$$

Analogously the amplitude for the processes $q\bar{q} \leftrightarrow gg$ is parameterized by

$$A_{gg} = \frac{\Lambda_g}{\sqrt{m_q}}. \quad (18)$$

according to the results of Close et al. [3] and Kühn et al. [21]. The phenomenological parameters Λ and Λ_g are to be determined. There is a parameter relating the binding energies which is very convenient in this mass matrix formalism; it is defined by

$$\varepsilon \equiv \frac{1}{2} (E_{uu} + E_{ss}) - E_{us}. \quad (19)$$

This parameter appears in the formalism when one uses the basis $|1\rangle$, $|8\rangle$ and $|G\rangle$ (or the basis $|N\rangle$, $|S\rangle$ and $|G\rangle$) and the mass relation for the non-self-conjugate mesons:

$$m_{I=1/2} = m_u + m_s + E_{us}, \quad (20)$$

$$m_{I=1} = 2m_u + E_{uu}. \quad (21)$$

The mass matrix contains off-diagonal elements involving not only the annihilation amplitudes but also the breaking of the SU(3) symmetry in the binding energies, represented by parameter ε .

The invariants of the mass matrix are functions of m_s/m_u , Λ/m_u^2 , $\Lambda_g/m_u^{1/2}$, ε and m_G . These quantities are not all free. Equations (8)–(10) impose some constraints among them. These equations can be solved for m_s/m_u , Λ/m_u^2 and $\Lambda_g/m_u^{1/2}$, which are functions of ε and m_G . Fixing the values of ε and m_G , the independent parameters of the model, all the remaining quantities become determined.

In the pseudoscalar sector [10] the value of m_G was limited to the interval between the masses of the pseudoscalar mesons η and η' , in order to keep the mass matrix Hermitian, because outside this interval Λ_g becomes a complex number. For a given value of m_G the parameter ε is determined by the minimum of m_s/m_u , consistent with the usual values in the nonrelativistic constituent quark models, which are in the range 1.3–1.8. For the determination of m_G , the remaining free parameter, we searched for the best values for the data from the branching ratios and from electromagnetic decay widths. We found

$$m_s/m_u = 1.772 \quad (22)$$

and $m_G = 1300$ MeV. With those values for m_s/m_u and m_G we obtained results for the branching ratios and electromagnetic decay widths involving the η , η' and η'' mesons in reasonable agreement with the experimental data. The value for the pseudoscalar glueball mass is to be compared with those predicted by other $\eta - \eta' - \eta''$ mixing schemes: 1369 MeV [22] and 1302 MeV [23]. It must be observed that the mass of the pseudoscalar glueball given by our model, similarly to some other mixing schemes is lower than the mass obtained in the lattice calculations ~ 2300 MeV [24,25]. In fact, there is an incompatibility between these approaches. In contrast to what is obtained in lattice results in the quenched approximation, in the mixing schemes the pseudoscalar glueball is not assumed to be an isolated physical state. The mass of the glueball state is obtained simultaneously with the masses of the $q\bar{q}$ and $s\bar{s}$ pseudoscalar states that are also components of the physical states. This is probably the source of the considerable difference between the masses estimated by these approaches. The ratio m_s/m_u , fixed by the pseudoscalar mesons, will be used as an input in the axial and tensor sectors.

3 Axial mesons

Applying the mixing scheme presented in the previous section to the isoscalar axial mesons, we find, after fitting the eigenvalues to the physical masses, the following eigenvectors:

$$|f_1(1285)\rangle = 0.630|1\rangle + 0.735|8\rangle - 0.250|G\rangle, \quad (23)$$

$$|f_1(1420)\rangle = -0.391|1\rangle - 0.223|8\rangle - 0.920|G\rangle, \quad (24)$$

$$|f_1(1510)\rangle = -0.671|1\rangle + 0.677|8\rangle + 0.302|G\rangle, \quad (25)$$

and

$$|f_1(1285)\rangle = 0.964|N\rangle + 0.090|S\rangle - 0.250|G\rangle, \quad (26)$$

$$|f_1(1420)\rangle = -0.208|N\rangle - 0.332|S\rangle - 0.920|G\rangle, \quad (27)$$

$$|f_1(1510)\rangle = 0.166|N\rangle + 0.939|S\rangle + 0.302|G\rangle. \quad (28)$$

These results suggest that $f_1(1285)$ has 93% of $|N\rangle$, $f_1(1420)$ has 85% of $|G\rangle$ and $f_1(1510)$ has 88% of $|S\rangle$. The independent parameters of the model, corresponding to these eigenvectors, are $\varepsilon = 25$ MeV and $m_G = 1430$ MeV. The remaining parameters are $\Lambda/m_u = 32.4$ MeV and $\Lambda_g/m_u^{1/2} = 0.79$ MeV.

The ratio of J/ψ radiative branching ratios into $f_1(1420)$ and $f_1(1285)$ and the ratio of the two-photon width of $f_1(1420)$ and $f_1(1285)$ are given by [18]

$$\begin{aligned} \frac{B(J/\psi \rightarrow \gamma f_1(1420))}{B(J/\psi \rightarrow \gamma f_1(1285))} &= \left(\frac{\sqrt{2}X_2 + Y_2}{\sqrt{2}X_1 + Y_1} \right)^2 \left(\frac{p_1}{p_2} \right)^2 \\ &= \frac{0.85 \pm 0.25}{B(f_1(1420) \rightarrow \eta\pi\pi)}, \end{aligned} \quad (29)$$

$$\begin{aligned} \frac{\Gamma_{\gamma\gamma}(f_1(1420))}{\Gamma_{\gamma\gamma}(f_1(1285))} &= \left(\frac{X_2 + \frac{\sqrt{2}}{5}Y_2}{X_1 + \frac{\sqrt{2}}{5}Y_1} \right)^2 \left(\frac{M(f_1(1420))}{M(f_1(1285))} \right)^2 \\ &= \frac{0.34 \pm 0.18}{B(f_1(1420) \rightarrow K\bar{K}\pi)}, \end{aligned} \quad (30)$$

where X and Y are the mixing coefficients appearing in (14) and (15) and the labels 1 and 2 stand for $f_1(1285)$ and $f_1(1420)$, respectively. Our results for those ratios are shown in Table 1 and are to be compared with experimental data.

4 Tensor mesons

The same approach used in the last section is now applied to the tensor mesons. If we consider the candidate to exotic $f_2(1640)$ as the partner of the tensor nonet, the resulting mixtures are

$$|f_2(1270)\rangle = 0.786|1\rangle + 0.480|8\rangle - 0.390|G\rangle, \quad (31)$$

$$|f_2'(1525)\rangle = 0.319|1\rangle + 0.598|8\rangle - 0.801|G\rangle, \quad (32)$$

$$|f_2(1640)\rangle = -0.642|1\rangle + 0.618|8\rangle + 0.454|G\rangle. \quad (33)$$

On the other hand, we can also consider that it is the $f_2(1710)$ that is mixing with the other tensor isosinglets. In this case we obtain

$$|f_2(1270)\rangle = 0.360|1\rangle + 0.634|8\rangle - 0.684|G\rangle, \quad (34)$$

$$|f_2'(1525)\rangle = 0.746|1\rangle + 0.175|8\rangle + 0.580|G\rangle, \quad (35)$$

$$|f_2(1710)\rangle = -0.487|1\rangle + 0.753|8\rangle + 0.442|G\rangle. \quad (36)$$

Changing to the basis $|N\rangle$, $|S\rangle$, $|G\rangle$ we find that

$$|f_2(1270)\rangle = 0.919|N\rangle + 0.062|S\rangle - 0.390|G\rangle, \quad (37)$$

$$|f_2'(1525)\rangle = 0.371|N\rangle - 0.470|S\rangle - 0.801|G\rangle, \quad (38)$$

$$|f_2(1640)\rangle = 0.134|N\rangle + 0.881|S\rangle + 0.454|G\rangle, \quad (39)$$

Table 1. Branching ratios and electromagnetic decay widths involving the $f_1(1285)$, $f_1(1420)$, $f_2(1270)$ and $f_2(1525)$. The results for tensor mesons were obtained in view of $f_2(1640)$ or $f_2(1710)$ as member of the tensor nonet: The results in parentheses refer to $f_2(1710)$. Our results are compared with the experimental data

Observable	Our model	Experiment [11]
$B(J/\psi \rightarrow \gamma f_1(1420))/B(J/\psi \rightarrow \gamma f_1(1285))$	0.16	1.7–8.5
$\Gamma_{\gamma\gamma}(f_1(1420))/\Gamma_{\gamma\gamma}(f_1(1285))$	0.13	0.34–0.68
$B(J/\psi \rightarrow \gamma f'_2)/B(J/\psi \rightarrow \gamma f_2(1270))$	0.0012 (0.06)	0.19
$B(f'_2 \rightarrow \pi\pi)/B(f'_2 \rightarrow K\bar{K})$	17 (127)	0.096

and

$$|f_2(1270)\rangle = 0.726|N\rangle + 0.072|S\rangle - 0.684|G\rangle, \quad (40)$$

$$|f'_2(1525)\rangle = 0.602|N\rangle - 0.549|S\rangle + 0.580|G\rangle, \quad (41)$$

$$|f_2(1710)\rangle = 0.336|N\rangle + 0.833|S\rangle + 0.442|G\rangle. \quad (42)$$

The independent parameters of the model which give the first set of eigenstates, $f_2(1270)$, $f'_2(1525)$ and $f_2(1640)$, are $\varepsilon = 52$ MeV and $m_G = 1510$ MeV. The remaining parameters are $\Lambda/m_u^2 = -1.0$ MeV and $\Lambda_g/m_u^{1/2} = 4.6$ MeV. The main content in these states is 85% of $|N\rangle$ in $f_2(1270)$, 64% of $|G\rangle$ in $f'_2(1525)$ and 78% of $|S\rangle$ in $f_2(1640)$.

The second set of eigenstates correspond to the following parameters, $\varepsilon = -5.0$ MeV, $m_G = 1444$ MeV, $\Lambda/m_u^2 = 4.8$ MeV and $\Lambda_g/m_u^{1/2} = 11$ MeV. The dominant contribution to each state of the second set is 53% of $|N\rangle$ in $f_2(1270)$ and 69% of $|S\rangle$ in $f_2(1710)$. The states $|N\rangle$, $|S\rangle$ and $|G\rangle$ contribute almost with the same proportion to the $f'_2(1525)$.

The ratio of J/ψ radiative branching ratios into $f_2(1270)$ and $f_2(1525)$ and the ratio of $f(1525)$ branching ratios into $\pi\pi$ and $K\bar{K}$ are given by [18]

$$\frac{B(J/\psi \rightarrow \gamma f'_2)}{B(J/\psi \rightarrow \gamma f_2(1270))} = \left(\frac{\sqrt{2}X_2 + Y_2}{\sqrt{2}X_1 + Y_1} \right)^2 \left(\frac{p_2}{p_1} \right)^3, \quad (43)$$

$$\frac{B(f'_2 \rightarrow \pi\pi)}{B(f'_2 \rightarrow K\bar{K})} = \frac{3X_2^2}{2 \left(\frac{X_2}{\sqrt{2}} + Y_2 \right)^2} \left(\frac{p_\pi}{p_K} \right)^5. \quad (44)$$

Here X and Y are the mixing coefficients appearing in (14) and (15) and the labels 1 and 2 refer to $f_2(1270)$ and $f'_2(1525)$, respectively. Our results for those ratios are shown in Table 1 and are to be compared with experimental data.

5 Conclusion

Birkel and Fritzsche [9] have used SU(3)-invariant annihilation amplitudes in a quadratic mass matrix formalism for describing the mixing in the axial sector. They found that the candidate to exotic $f_1(1420)$ has a gluonic content of 58% and the gluonic component with a mass of 1432 MeV.

These results are to be compared to the ones found in our mixing schemes. We found a gluonic content of 85% in the $f_1(1420)$ and a gluonic component at 1430 MeV.

In the tensor sector we found that the $f_2(1270)$ could be mainly an $|N\rangle$ state. Nevertheless, we found that the candidate to exotic $f_2(1640)$ is predominantly a $|S\rangle$ state, whereas the $f'_2(1525)$ is mainly a $|G\rangle$ state. These results are in contrast with those found in the literature that indicate the $f'_2(1525)$ and $f_2(1640)$ are mainly $|S\rangle$ and $|G\rangle$ states, respectively [2, 6–8]. Our results were obtained for a gluonic component at 1510 MeV. On the other side, if the physical state is the $f_2(1710)$ we found that it is mainly a $|S\rangle$ state and $f'_2(1525)$ is nearly an equiprobable distribution among $|N\rangle$, $|S\rangle$ and $|G\rangle$. For this set of eigenstates a gluonic component at 1444 MeV was obtained. In the first set of eigenstates, the mass of the gluonic component is comparable to the mass found in the range 1536–1590 MeV obtained by other mixing schemes [5].

Here, as in the case of the pseudoscalar sector, we have obtained masses for a glueball state lower than the one obtained in the quenched lattice (~ 2000 – 2300 MeV) [25, 26]. The source of the substantial difference among the masses is probably the same as that in the case of the pseudoscalar mesons: The mass of the glueball states in the mass matrix formalism is obtained regarding the glueball as being a component of a physical state, whereas in the lattice calculations the glueball is a physical state itself. Nevertheless, the results given by the present quarkonium–gluonium mixing scheme for branching ratios and electromagnetic decay widths involving the axial and tensor mesons $f_1(1285)$, $f_1(1420)$, $f_2(1270)$ and $f_2(1525)$ are in clear contradiction with the experimental ones. The theoretical and experimental results are compared in Table 1. These results show that the quarkonium–gluonium mixing model, which works well for scalar isosinglet mesons [10], is not compatible with the constraints coming from decays concerning the axial and tensor isosinglet mesons considered in this work.

The incompatibility above mentioned is a clue that the presence of gluonic components in the axial and tensor isosinglet meson states considered here may be a wrong assumption. On the other hand the interpretations of the states $f_1(1420)$, $f_1(1510)$, $f_2(1640)$ and $f_2(1710)$ are controversial and, moreover, some of them need confirmation [11].

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