Two-photon decay of light scalars: a comparison of tetraquark and quarkonium assignments

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Abstract

Two-photon decays of light scalar mesons are discussed within the quarkonium and tetraquark assignments: in both cases the decay rate of the sigma resonances turns out to be smaller than 1 keV.

1 Introduction

The two-photon decay of light mesons represents an important source of information [1]. In particular, the $\gamma\gamma$ decay of light scalar mesons has been considered as a possible tool to deduce their nature. According to the interpretation of light scalars (quarkonia, tetraquarks or molecules) different $\gamma\gamma$-rates are expected [2, 3]. In this proceeding we first study (Section 2) the quarkonium assignments for the light scalar states by studying $SU(3)$-relations for the two-photon decays. While the quarkonium assignment is disfavored when looking
at mass pattern \[4\], strong decays \[3\ \[5\] and large-\(N_c\) studies \[6\] (see the discussion in the recent proceeding \[7\]), here we adopt a neutral point of view. After a test-study with the well-known pseudoscalar and tensor mesons \[8\] we turn to the \(\gamma\gamma\) decay of light scalars as quarkonia, finding that \(f_0(600)\) has a decay rate well below 1 keV. This is in agreement with the microscopic evaluation in the recent work \[9\], where the two-photon decay of a low-lying quarkonium state with \(\overline{m} = \sqrt{1/2(\overline{u}u + \overline{d}d)}\): contrary to usual results quoted in the literature it is shown that the corresponding two-photon decay rate is well below 1 keV for a mass \(M_{\overline{m}} < 0.8\) GeV. In Section 3 we turn to the two-photon transition within the tetraquark assignment and in Section 4 we briefly summarize our results.

2 Quarkonia into \(\gamma\gamma\)

2.1 Quarkonia into \(\gamma\gamma\): symmetry relations

We first consider the two-photon decay of scalar quarkonia states. However, the formula we introduce are also valid, with simple changes, for the pseudoscalar, tensor and axial vector nonets as we will discuss later on. The charge neutral scalar quarkonia states \(N, S\) and \(a_0\) are introduced as \(N \equiv \overline{m} = \sqrt{1/2(\overline{u}u + \overline{d}d)}\), \(S \equiv \overline{s}s\) and \(a_0 \equiv \sqrt{1/2(\overline{u}u − \overline{d}d)}\). The \(3 \times 3\) diagonal matrix \(S_{[qq]} = diag\{\overline{u}u, \overline{d}d, \overline{s}s\} = diag\{N/\sqrt{2} + a_0/\sqrt{2}, N/\sqrt{2} − a_0/\sqrt{2}, S\}\) plays a central role. In flavor (and large-\(N_c\)) limit the two-photon decay of these states is described by the effective interaction Lagrangian

\[
\mathcal{L}_{\gamma\gamma} = c_{\gamma\gamma} Tr \left[ Q^2 S_{[qq]} \right] F_{\mu\nu}^2
\]

where \(Q = diag\{2/3, −1/3, −1/3\}\) is the quark charge matrix, \(F_{\mu\nu} = \partial_\mu A_\nu − \partial_\nu A_\mu\) the field tensor of the electromagnetic field \(A_\mu\) and \(c_{S,\gamma\gamma}\) a coupling constant. As a result the decay rates of \(N, S,\) and \(a_0\) are given by

\[
\Gamma_{N,\gamma\gamma} = \frac{c_{\gamma\gamma}^2}{4\pi} M_N^3 \left[ \frac{5}{9\sqrt{2}} \right]^2, \quad \Gamma_{S,\gamma\gamma} = \frac{c_{\gamma\gamma}^2}{4\pi} M_S^3 \left[ \frac{1}{9} \right]^2, \quad \Gamma_{a_0,\gamma\gamma} = \frac{c_{\gamma\gamma}^2}{4\pi} M_{a_0}^3 \left[ \frac{3}{9\sqrt{2}} \right]^2.
\]

The physical states, denoted as \(f_0(600)\) and \(f_0(980)\) in the low-scalar case, are in general a mixing of \(N\) and \(S\): \(f_0(600) = \cos \varphi S N + \sin \varphi S S\) and orthogonal combination for \(f_0(980)\). The two-photon decay rates of \(f_0(600)\) and \(f_0(980)\)
are given by:

\[
\Gamma_{f_0(600)\gamma\gamma} = \frac{c^2\gamma^2}{4\pi} M_{f_0(600)}^3 \left[ \frac{5}{9\sqrt{2}} \cos \varphi_S + \frac{1}{9} \sin \varphi_S \right]^2,
\]

\[
\Gamma_{f_0(980)\gamma\gamma} = \frac{c^2\gamma^2}{4\pi} M_{f_0(980)}^3 \left[ -\frac{5}{9\sqrt{2}} \sin \varphi_S + \frac{1}{9} \cos \varphi_S \right]^2.
\]

Before studying the scalar case we test these simple expressions on other nonets. In particular, we will be interested to the ratios \( \Gamma_{f_0(600)\gamma\gamma}/\Gamma_{a_0(0)\gamma\gamma} \) and \( \Gamma_{f_0(980)\gamma\gamma}/\Gamma_{a_0(0)\gamma\gamma} \) for which the dependence on the unknown parameter \( c_{\gamma\gamma} \) cancels out under the hypothesis that an eventual momentum dependence is weak, see the following discussions.

**Pseudoscalar nonet** - Here we have \( \pi^0 = \sqrt{1/2}(\bar{u}u - \bar{d}d) \) and the isoscalar states \( \eta \) and \( \eta' \), expressed in terms of bare states as

\[
\begin{pmatrix}
\eta \\
\eta'
\end{pmatrix}
= \begin{pmatrix}
\cos \varphi_P & \sin \varphi_P \\
-\sin \varphi_P & \cos \varphi_P
\end{pmatrix}
\begin{pmatrix}
N \\
S
\end{pmatrix}
= \begin{pmatrix}
\cos \theta_P & -\sin \theta_P \\
\sin \theta_P & \cos \theta_P
\end{pmatrix}
\begin{pmatrix}
P^8 \\
P^0
\end{pmatrix},
\]

where \( P^8 = \sqrt{1/6}(\bar{u}u - \bar{d}d - 2\bar{s}s) \) and \( P^0 = \sqrt{1/3}(\bar{u}u + \bar{d}d + \bar{s}s) \). In the literature the angle \( \theta_P \) is in general discussed. The corresponding Lagrangian is similar to (1): \( L_{P\gamma\gamma} = c_{P\gamma\gamma}Tr \left[ Q^2 F_{\mu\nu} \tilde{F}_{\mu\nu} \right] \), where \( F_{\mu\nu} = \varepsilon_{\mu\nu\rho\sigma} F_{\rho\sigma} \); eq. (2) retain the same form. Out of (3) one finds the relation \( \theta_P = -\arcsin(\sqrt{2/3} + \varphi_P) = -(54.736^\circ + \varphi_P) \). In [10] the values \( \Gamma_{\eta\gamma\gamma} = 7.74 \pm 0.55 \) eV, \( \Gamma_{\eta'\gamma\gamma} = 0.510 \pm 0.026 \) keV and \( \Gamma_{\eta\gamma'\gamma} = 4.30 \pm 0.15 \) keV. The corresponding experimental ratios read \( \frac{\Gamma_{\eta\gamma\gamma}}{\Gamma_{\bar{u}u\gamma\gamma}} = 65.9 \pm 8.1 \), \( \frac{\Gamma_{\eta'\gamma\gamma}}{\Gamma_{\bar{u}u\gamma\gamma}} = 556 \pm 59 \). A fit of \( \varphi_P \) to the latter ratios implies \( \varphi_P = -36.0^\circ \), and thus \( \theta_P = -18.7^0 \). The corresponding ratios evaluated at this mixing angle read \( \frac{\Gamma_{\eta\gamma\gamma}}{\Gamma_{\bar{u}u\gamma\gamma}} = 76.6 \) and \( \frac{\Gamma_{\eta'\gamma\gamma}}{\Gamma_{\bar{u}u\gamma\gamma}} = 661.9 \) with \( \chi^2/2 = 2.48 \).

Taking into account that we are considering the easiest possible scenario, thus neglecting many possible corrections, these results are very good. In fact, the experimental results range within 3 order of magnitudes: the theoretical dependence on the third power of the mass is well verified. An eventual mass dependence of the effective coupling \( c_{P\gamma\gamma} \) has to be small form \( M_\pi \) up to 1 GeV, a remarkable fact. The mixing angle \( \theta_P = -18.7^0 \) is in the phenomenological range between \(-10^\circ\) and \(-20^\circ\), as found also by more refined studies. This simple exercise shows that SU(3) flavor relations, together with the OZI rule allowing to include the flavor singlet in the game, is well upheld and a good starting point to test the quarkonium assignment in a given nonet.
Tensor nonet- As a further test let’s consider the tensor nonet (see 8 and refs. therein). The resonances under study are the isovector $a_2(1320)$ with $M_{a_2} = 1318.3$ MeV and the isoscalars $f_2(1270)$ and $f_2^*(1525)$ with $M_{f_2} = 1275.1$ and $M_{f_2^*} = 1525$ (we omit tiny errors, see 10). As before, the mixing angle $\varphi_T$ is introduced as $f_2(1270) = \cos \varphi_T N + \cos \varphi_T S$ and orthogonal combination for $f_2^*(1525)$. The interaction Lagrangian is $\mathcal{L}_{T\gamma\gamma} = c_{T\gamma\gamma} \text{Tr} \left[ Q^2 T_{\mu\nu} \right] \Theta^{\mu\nu}$ where $\Theta^{\mu\nu}$ is the energy-momentum tensor of the electromagnetic fields. The experimental values $\Gamma_{a_2\gamma\gamma} = 1.00 \pm 0.06$ keV, $\Gamma_{f_2\gamma\gamma} = 2.60 \pm 0.24$ keV and $\Gamma_{f_2^*\gamma\gamma} = (8.1 \pm 0.9) \times 10^{-2}$ keV lead to the ratios $\frac{\Gamma_{f_2^*\gamma\gamma}}{\Gamma_{a_2\gamma\gamma}} = 2.6 \pm 0.4$ and $\frac{\Gamma_{f_2\gamma\gamma}}{\Gamma_{a_2\gamma\gamma}} = (8.1 \pm 1.4) \times 10^{-2}$. A fit of the angle $\varphi_T$ to these values leads to $\varphi_T = 8.19^\circ$ and a very small $\chi^2/2 = 0.015$, thus the experimental values are reproduced almost exactly. The value of $\varphi_T = 8.19^\circ$ is in good agreement with the study of strong decays 8. Again, the $\gamma\gamma$-ratios can be well described by simple symmetry relations. Indeed, the agreement is even better than in the pseudoscalar case: this is expected because the masses vary in a smaller energy region. In the end, we remind that a quarkonium interpretation works well for vector mesons: here the dominant e.m. transition is into one single photon (i.e. mixing) which is at the basis of the successful phenomenology of the vector meson dominance hypothesis.

Scalar nonet below 1 GeV- Let us now turn back to the scalar states below 1 GeV within a quarkonium assignment. We identify the neutral isovector $a_0(980)$ with $a_0^0(980)$ and, as described above, the isoscalars with $f_0(600)$ and $f_0(980)$. The experimental results for the decay width of $a_0$ and $f_0(980)$ are given by 10. $\Gamma_{f_0(980)\gamma\gamma} = 0.39^{+0.10}_{-0.13}$ KeV, $\Gamma_{a_0^0\gamma\gamma} = 0.30 \pm 0.10$ KeV. Thus, the experimental ratio reads then $\frac{\Gamma_{f_0(980)\gamma\gamma}}{\Gamma_{a_0^0\gamma\gamma}} = 1.30 \pm 0.8$ where an average error of 0.115 KeV for $\Gamma_{f_0(980)\gamma\gamma}$ has been used. As noticeable, the error for this ratio is large. Unfortunately, the experimental situation concerning $f_0(600) \rightarrow 2\gamma$ is even worse: no average or fit is presented in 10, however two experiments listed in 10 find large $\gamma\gamma$ decay widths: 3.8 $\pm 1.5$ keV and 5.4 $\pm 2.3$ keV, respectively. In a footnote it is then state that these values could be assigned to $f_0(1370)$ (actually, in a older version of PDG 11 these values were assigned to the resonance $f_0(1370)$). It is not clear if the $\gamma\gamma$ signal comes from the high mass tail of the broad $\sigma$ state or from $f_0(1370)$ (or even from both). We determine the mixing angle $\varphi_S$ by using the experimental result $\frac{\Gamma_{f_0(980)\gamma\gamma}}{\Gamma_{a_0^0\gamma\gamma}} = 1.30 \pm 0.8$. 

Due to the large error we report the possible ranges for $\varphi_S$ compatible with it: $-105.9^\circ \leq \varphi_S \leq -47.4^\circ$ (central value $\varphi_S = -56.9^\circ$) and $15.9^\circ \leq \varphi_S \leq 74.2^\circ$ (central value $\varphi_S = 25.3^\circ$). Indeed, as we saw previously the angle $\varphi_P$ in the pseudoscalar sector is negative: the components of $N$ and $S$ are in phase for $\eta'$ and out of phase for $\eta$ ($\varphi_P = -36.0^\circ$). Within the NJL model the mixing strength is generated by the 't Hooft term, and turns out to have opposite sign with respect to the pseudoscalar sector: this would favour a positive value of $\varphi_S$. Furthermore, studies without the strange meson can reproduce the $f_0(600)$ resonance: this favors small mixing angles. Then, a value $\varphi_S \sim 25^\circ$ is favoured. The corresponding two-photon decay rate is $\Gamma_{f_0(980)\gamma\gamma} \sim 0.4$ (which in turn means $\Gamma_{f_0(980)\gamma\gamma} \sim 0.12$ KeV, considerably smaller than the above mentioned (but not established) experimental result). Surely, when including finite width effects the decay rate $\Gamma_{f_0(600)\gamma\gamma}$ increases: in fact, the kinematical factor $M^3_{f_0(600)}$ makes the right-tail of the broad distribution of $f_0(600)$ important. However, the rate of increase is not dramatic: it can at most double the above quoted results but not reach values of about 2-5 keV. This result is contrary to the usual belief that a quarkonium decay rate should be well above 1 keV: indeed, as we discuss in the next subsection a careful microscopic calculation of the two-photon decay rate shows that results below 1 keV are expected.

For the discussion of scalar states above 1 GeV we refer to 14, where the inclusion of a glueball state mixing with quarkonia also influences the $\gamma\gamma$ decay rate: in fact, a glueball is expected to have a small $\gamma\gamma$-transition amplitude, thus if a resonance will have a consistent glueball component the $\gamma\gamma$ decay rate is small. No $\gamma\gamma$-signal is found for $f_0(1500)$ pointing to a large gluonic amount in its wave function.

2.2 Quarkonium into $\gamma\gamma$: a microscopic evaluation

In this subsection we refer to 9, where the $\gamma\gamma$-decay rate has been carried out within a local and nonlocal microscopic model. Here we discuss only the latter. The relevant nonlocal interaction Lagrangian, involving the mesonic quarkonium field $\sigma(x)$ and the quark fields $q^i = (u, d)$, reads 9,15

$$L_{\text{int}}(x) = \frac{g_\sigma}{\sqrt{2}} \sigma(x) \int d^4 y \bar{\Phi}(y^2) \bar{q}(x + y/2)q(x - y/2),$$

(4)
where the delocalization takes account of the extended nature of the quarkonium state by the covariant vertex function \( \Phi(y^2) \). The (Euclidean) Fourier transform of this vertex function is taken as \( \tilde{\Phi}(k^2_{E}) = \exp(-k^2_{E}/\Lambda^2) \), also assuring UV-convergence of the model. The cutoff parameter \( \Lambda \) is varied between 1 and 2 GeV, corresponding to an extension of the \( \sigma \) of about \( l \sim 1/\Lambda \sim 0.5 \) fm. The coupling \( g_\sigma \) is determined by the so-called compositeness condition [9, 15]. In the calculation the quark mass varies between 0.3 and 0.45 GeV.

The two-photon decay occurs via a triangle-diagram of quarks. Notice that due to the presence of the vertex function \( \Phi(y^2) \) inclusion of the electromagnetic interaction is achieved by gauging the nonlocal interaction Lagrangian (4): in addition to the usual photon-quark coupling obtained by minimal substitution new vertices arise, where the photon couples directly to the \( \sigma\gamma\gamma \) interaction vertex, see [15] for details. Their contribution, important on a conceptual level to assure gauge invariance, is numerically suppressed. In Table 1 we summarize our results for \( M_\sigma = 0.6 \) GeV varying \( m_q \) both for \( \Lambda = 1 \) GeV and, in parenthesis, for \( \Lambda = 2 \) GeV.

### Table 1: \( \Gamma_{\sigma\gamma\gamma} \) for \( m_q = 0.31 - 0.45 \) GeV, \( \Lambda = 1(2) \) GeV at \( M_\sigma = 0.6 \) GeV.

| \( m_q \) (GeV) | 0.31 | 0.35 | 0.40 | 0.45 |
|----------------|------|------|------|------|
| \( \Gamma_{\sigma\gamma\gamma} \) (keV) at \( M_\sigma = 0.6 \) GeV | 0.529 | 0.458 | 0.361 | 0.294 |
| & (0.512) | (0.415) | (0.327) | (0.267) |

The decay widths decrease slowly for increasing quark mass while the dependence on the cutoff \( \Lambda \) is very weak. The numerical analysis shows that \( \Gamma_{\sigma\gamma\gamma} < 1 \) keV for \( M_\sigma < 0.7-0.8 \) GeV. This result is indeed in agreement with that of the previous subsection: a light quarkonium state has a \( \gamma\gamma \) decay rate smaller than 1 keV. However, this doesn’t prove that the resonance \( f_0(600) \) is a quark-antiquark state. It rather tells us that, being the \( \gamma\gamma \) decay width of a quarkonium smaller than what usually believed, care is needed when using \( \gamma\gamma \)-rates to discuss the nature of light scalars. We also refer to [16] for the evaluation of these diagrams for quarkonia states above 1 GeV.

### 3 Tetraquarks into \( \gamma\gamma \)

We consider now the \( \gamma\gamma \)-transition of tetraquark states [3], whose effective Lagrangian reads

\[
\mathcal{L}_{em} = c_1^\gamma S_{ij}^{[4q]} \langle A^i QA^j Q \rangle F^2_{\mu\nu} - c_2^\gamma S_{ij}^{[4q]} \langle A^i A^j Q^2 \rangle F^2_{\mu\nu},
\] (5)
Within the tetraquark assignment the isoscalars are a term, proportional to the next-to-leading order correction. As discussed in detail in 3, of the parameters form the large-$N_c$ expansion (switch of a quark with an antiquark), while the second term, proportional to $c_2^{\gamma\gamma}$ (annihilation of a quark-antiquark pair), represents the next-to-leading order correction. As discussed in detail in [3][17], the latter mechanism can be relevant because it occurs with only one gluon as intermediate state. The decay width into two photons reads $\Gamma_{\gamma\gamma} = \frac{M^2_{\gamma\gamma}}{4\pi} g_{\gamma\gamma}^2$, where $i = a_0^{[4q]}, \sigma_B^{[4q]}, f_B^{[4q]}$. The coupling constants for $a_0^{[4q]}$ and for the bare states $\sigma_B$ and $f_B$ are deduced from (5) and read:

$$g_{a_0^{[4q]},\gamma\gamma} = \frac{2c_1^{\gamma\gamma} + c_2^{\gamma\gamma}}{3\sqrt{2}}, \quad g_{\sigma_B^{[4q]},\gamma\gamma} = \frac{4c_1^{\gamma\gamma} + 5c_2^{\gamma\gamma}}{9}, \quad g_{f_B^{[4q]},\gamma\gamma} = \frac{2c_1^{\gamma\gamma} + 7c_2^{\gamma\gamma}}{9\sqrt{2}} \quad (6)$$

The mixed physical states $f_0(600)$ and $f_0(980)$ are expressed in the tetraquark framework as $f_0(600) = \cos \varphi_S \sigma_B^{[4q]} + \sin \varphi_S f_B^{[4q]}$ and orthogonal combination for $f_0(980)$. Let us first consider $c_2^{\gamma\gamma} = 0$. When determining the mixing angle $\varphi_S$ by using the experimental ratio \( \frac{\Gamma_{f_0(980)\gamma\gamma}}{\Gamma_{a_0^{[4q]}\gamma\gamma}} = 1.30 \pm 0.8 \) one obtains very large values: $|\varphi_S| \gtrsim 70^\circ$ (indeed $\Gamma_{f_0(980)\gamma\gamma}/\Gamma_{a_0^{[4q]}\gamma\gamma} \leq 1$ for each $\varphi_S$). One of the main advantages of the tetraquark assignment is the explanation of the mass degeneracy of $a_0(980)$ and $f_0(980)$ in the limit $\varphi_S = 0$. However, a large mixing angle would completely spoil the mass degeneracy. We thus consider this possibility disfavored, see discussion in [3]. When $c_2^{\gamma\gamma} \neq 0$ a determination of the parameters form the $\gamma\gamma$ data is no longer possible. However, the mixing angle $\varphi_{SS}$ can be fixed from strong decays [3]: $\varphi_S = -12.8^\circ$. Then we find $0.15 \leq c_2^{\gamma\gamma}/c_1^{\gamma\gamma} \leq 1.39$. Notice that even a small but non vanishing $c_2^{\gamma\gamma}$ can improve a lot the phenomenology: in fact, $c_2^{\gamma\gamma}$ strongly enhances the amplitude $g_{f_B^{[4q]},\gamma\gamma}$, see eq. (8). For $0.15 \leq c_2^{\gamma\gamma}/c_1^{\gamma\gamma} \leq 1.39$ one has $\frac{\Gamma_{f_0(600)\gamma\gamma}}{\Gamma_{a_0^{[4q]}\gamma\gamma}} \leq 0.35$, again pointing to a small $\gamma\gamma$-rate of $f_0(600)$ as in the quarkonium case. More work is needed but one result is stable: the $\Gamma_{f_0(600)\gamma\gamma}$ is well below 1 keV also in the tetraquark assignment and is indeed of the same order of magnitude of the quarkonium interpretation. One could indeed go further by including mixing of tetraquark below 1 GeV and quarkonia above 1 GeV: however, as found in [17] the latter turns out to be small, thus not changing much the present results.
4 Conclusions

In this work we discussed the two-photon transition of light scalar mesons within the quarkonium and the tetraquark assignments. In both cases the decay rate $\Gamma_{f_0(600)\gamma\gamma}$ is smaller than 1 keV, as confirmed by a microscopic calculation in the quarkonium assignment. These results render a possible identification of the nature of scalar states using two-photon decay widths more difficult.

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