Axial Anomaly and Mixing Parameters of Pseudoscalar Mesons

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Abstract. In this work the analysis of mixing parameters of the system involving $\eta, \eta'$ mesons and some third massive state $G$ is carried out. We use the generalized mixing scheme with three angles. The framework of the dispersive approach to Abelian axial anomaly of isoscalar non-singlet current and the analysis of experimental data of charmonium radiative decays ratio allow us to get a number of quite strict constraints for the mixing parameters. The analysis shows that the equal values of axial current coupling constants $f_8$ and $f_0$ are preferable which may be considered as a manifestation of $SU(3)$ and chiral symmetry.

1 Introduction

This work is developing the approach of the papers \cite{1,2} and is devoted to the significant problem of mixing of pseudoscalar mesons. It is especially important with a number of current and planned experiments.

The problem of $\eta-\eta'$ mixing has been studied for many years. The usual approach with one mixing angle dominated for decades, but in the recent years the more elaborated schemes appear to be unavoidable \cite{3-8}. In particular, the theoretical ground of this was based on the recent progress in the ChPT \cite{9-11}. On the other hand, it was shown, that the current experimental data cannot satisfactory describe the whole set of experiments within the one-angle mixing scheme.

The mixing schemes are usually enunciated either in terms of $SU(3)$ or quark basis. In our paper \cite{2} we construct and use the generalization of $SU(3)$ basis similar to the mixing of massive neutrinos. This is because we use the dispersive approach to axial anomaly (\cite{12,13} for a review) to find some model-independent and precise restriction on the mixing parameters.

It was shown that any scheme with more than one angle unavoidably demands an additional admixture of higher mass state. If we restrict ourselves to only one additional state $G$ (denoted as a glueball without really specifying its nature) then the general mixing scheme can be described in terms of 3 angles. In particular cases the number of angles can be reduced to two.

In the paper \cite{2} the analysis of different conventional (and most physically interesting) particular cases was performed (including two-angle mixing schemes) basing on the dispersive representation of axial anomaly from one side and charmonium decays ratio from the other side.

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The main conclusion of the paper \cite{2} is that in all considered cases the only reasonable solutions appear at \( f_8 = f_0 \simeq f_a \). The main aim of this work is to check whether this relation remains valid in the most general case with some specific constraints imposed.

This paper is organized as follows. In the Sec. 2 we introduce our notation and the general approach to the mixing. In Sec. 3 we derive the basic equations relying on the dispersive approach to Abelian axial anomaly of isoscalar non-singlet current \( J_{a\mu}^8 \) and the charmonium radiative decay ratio \( R_{J/\Psi} \), while in Sec. 4 we perform the numerical analysis of these equations. Finally, in Sec. 5 we present the conclusion.

### 2 Mixing scheme

We start with a \((N\)-component) vector of physical pseudoscalar fields consisting of the fields of the lightest pseudoscalar mesons and other fields:

\[
\tilde{\Phi} = \begin{pmatrix}
\pi^0 \\
\eta \\
\eta' \\
G
\end{pmatrix}.
\]

We are not able to specify the physical nature of the other components with higher masses, the lowest of which \( G \) can be either a glueball or some excited state \( ^1 \). Let us also introduce, following \cite{16,17}, a set of \( SU(3) \) fields \( \varphi_3, \varphi_8, \varphi_0 \) (\( \Phi_1, \Phi_2, \Phi_3 \)) and complement them with other (sterile) fields \( g_i \) (\( \Phi_i, i = 4..N \))

\[
\Phi = \begin{pmatrix}
\varphi_3 \\
\varphi_8 \\
\varphi_0 \\
g
\end{pmatrix}.
\]

The three upper fields \( \varphi_3, \varphi_8, \varphi_0 \) are the only ones which define the generalized PCAC relation for axial current \( J_{a\mu}^5 = \frac{q\gamma_\mu\gamma_5}{\sqrt{2}}q \) (no summation over \( a \) contrary to \( j \) and \( k \) is assumed):

\[
\partial_\mu J_{a\mu}^5 = f_a \delta \Delta L = F_{aj} M_{jk} \Phi_k, \quad a = 3, 8, 0, \quad j, k = 1..N,
\]

where \( \Delta L \) is the mass term in the effective Lagrangian with a non-diagonal mass matrix \( M \) (as fields \( \Phi_k \) are not orthogonal to each other):

\[
\Delta L = \frac{1}{2} \Phi^T M \Phi,
\]

and \( F \) is a matrix of decay constants \( ^2 \).

In order to proceed from initial \( SU(3) \) fields \( \Phi \) to physical mass fields \( \tilde{\Phi} \) the unitary (real, as the CP-violating effects are negligible) matrix \( U \) is introduced

\(^1\) Note, that the mixing with the excited states is usually (e.g. \cite{16,15}) supposed to be suppressed.

\(^2\) Note, that matrix of decay constants \( F \) is non-square expressing the fact that generally the number of \( SU(3) \) currents is less then the number of all possible states involved in mixing. The similar situation takes place (see e.g. \cite{13}) in one of the extensions of the Standard Model – neutrino mixing scenario involving sterile neutrinos.
where \( m_i \) constants are equal, it is reduced to formula (3.40) in [19]. This formula is close to those obtained in [16,17] (in the limit of small mixing). When the decay

respectively.

rents:

can be compared to the standard definition of the "physical" couplings of axial cur-

This expression (recall, that there is no summation over \( a \)) read:

\[
\partial_{\mu} J_{\mu5} = F U^T \tilde{M} \Phi
\]

This formula is close to those obtained in [10-17] (in the limit of small mixing). When the decay constants are equal, it is reduced to formula (3.40) in [19].

The matrix elements of \( \partial_{\mu} J_{\mu5} \) between vacuum state and physical states \(|\tilde{\Phi}_k\rangle\)

\[
\langle 0 | \partial_{\mu} J_{\mu5} | \tilde{\Phi}_k \rangle = F^{a}_i (U^T \tilde{M})_k^a
\]

can be compared to the standard definition of the "physical" coupling constants of axial currents:

\[
\langle 0 | J^a_{\mu5} | \tilde{\Phi}_k \rangle = i f^a_k q_{\mu}.
\]

From (9) and (10) follows the relation

\[
f^a_k = F^{a}_i (U^T)^i_k = f_a (U^T)^i_k.
\]

This expression (recall, that there is no summation over \( a \)) clearly shows that \( f^a_k \) are obtained by multiplication of each line of \( U^T \) by respective coupling \( f_a \) and form a non-diagonal (contrary to \( F \)) matrix.

Taking into account the well-known smallness of \( \pi^0 \) mixing with the \( \eta, \eta' \) sector [16,17,20] and neglecting all higher contributions we restrict our consideration to three physical states \( \eta, \eta', G \) and two currents \( J^8_{\mu5}, J^0_{\mu5} \). Then the divergencies of the axial currents (recall, that \( G \) is a first mass state heavier than \( \eta' \)):

\[
\begin{pmatrix}
\partial_{\mu} J^8_{\mu5} \\
\partial_{\mu} J^0_{\mu5}
\end{pmatrix} = 
\begin{pmatrix}
f_8 & 0 & 0 \\
0 & f_0 & 0
\end{pmatrix}
U^T 
\begin{pmatrix}
m^2_{\pi} & 0 & 0 \\
0 & m^2_{\eta'} & 0 \\
0 & 0 & m^2_G
\end{pmatrix} 
\begin{pmatrix}
\eta \\
\eta' \\
G
\end{pmatrix}.
\]

(12)

Exploring the mentioned similarity of the meson and lepton mixing, we use the Euler parametrization for the mixing matrix \( U \) (we use notation \( c_i \equiv \cos \theta_i, s_i \equiv \sin \theta_i \)):

\[
U = \begin{pmatrix}
c_3 c_8 - c_0 s_3 s_8 - c_3 s_8 - c_8 c_0 s_3 & s_3 s_0 \\
& \\
s_3 c_8 + c_3 c_0 s_8 - s_3 s_8 + c_3 c_0 - c_3 s_8 & c_8 s_0
\end{pmatrix}.
\]

(13)

In the following consideration we will need the divergency of the octet current \( \partial_{\mu} J^8_{\mu5} \), so let us write it out explicitly:

\[
\partial_{\mu} J^8_{\mu5} = f_8 (m^2_{\pi} \eta (c_3 c_0 - c_3 s_0 s_8) + m^2_{\eta'} \eta' (s_3 c_8 + c_3 c_0 s_8) + m^2_G (s_8 s_0)).
\]

(14)

As soon as in the chiral limit \( J^8_{\mu5} \) should be conserved, from Eq. (14) follows that coefficients of the terms \( m^2_{\eta'}, m^2_G \) must decrease at least as \((m_\eta / m_{\eta',G})^2 \). More specifically, we expect the following limits for the terms of Eq. (14):

\[
\frac{|s_8 s_0|}{s_3 c_8 + c_3 c_0 s_8} \lesssim \left( \frac{m_\eta}{m_G} \right)^2.
\]

(15)
3 Abelian axial anomaly and charmonium decays ratio

In our paper the dispersive form of the anomaly sum rule will be extensively used, so we remind briefly the main points of this approach (see e.g. review [13] for details).

Consider a matrix element of a transition of the axial current to two photons with momenta \( p \) and \( p' \)

\[
T_{\mu\alpha\beta}(p,p') = \langle p',p|J_{\mu5}|0 \rangle.
\]  

The general form of \( T_{\mu\alpha\beta} \) for a case \( p^2 = p'^2 \) can be represented in terms of structure functions (form factors):

\[
T_{\mu\alpha\beta}(p,p') = F_1(q^2)q_\mu\epsilon_{\alpha\beta\rho\sigma}p_\rho p'_\sigma + \frac{1}{2} F_2(q^2) \left[ \frac{p_\alpha}{p'^2} \epsilon_{\mu\beta\rho\sigma}p_\rho p'_\sigma - \frac{p'_\alpha}{p'^2} \epsilon_{\mu\alpha\rho\sigma}p_\rho p'_\sigma - \epsilon_{\mu\alpha\beta\rho}(p - p')_\rho \right],
\]  

where \( q = p + p' \). The functions \( F_1(q^2) \), \( F_2(q^2) \) can be described by dispersion relations with no subtractions and anomaly condition in QCD results in the sum rule:

\[
\int_0^\infty \text{Im} F_1(q^2) dq^2 = 2\alpha N_c \sum c_q^2,
\]  

where \( c_q \) are quark electric charges and \( N_c \) is the number of colors. This sum rule [21] was developed by Jiří Horejší [22] and later generalized [23]. Notice that in QCD this equation does not have any perturbative corrections [24], and it is expected that it does not have any non-perturbative corrections as well due to ’t Hooft’s consistency principle [25]. It will be important for us that as \( q^2 \to \infty \) the function \( \text{Im} F_1(q^2) \) decreases as \( 1/q^4 \) (see discussion in Ref. [2]). Note also that the relation (18) contains only mass-independent terms, which is especially important for the 8th component of the axial current \( J_{\mu5}^8 \) containing strange quarks:

\[
J_{\mu5}^8 = \frac{1}{\sqrt{6}} (\bar{u} \gamma_\mu \gamma_5 u + \bar{d} \gamma_\mu \gamma_5 d - 2\bar{s} \gamma_\mu \gamma_5 s).
\]  

The general sum rule (18) takes the form:

\[
\int_0^\infty \text{Im} F_1(q^2) dq^2 = \frac{2\sqrt{6}}{3} \alpha (e_u^2 + e_d^2 - 2e_s^2) N_c = \frac{2\sqrt{3}}{3} \alpha,
\]  

where \( e_u = 2/3 \), \( e_d = e_s = -1/3 \), \( N_c = 3 \).

In order to separate the form factor \( F_1(q^2) \), multiply \( T_{\mu\alpha\beta}(p,p') \) by \( q_\mu/q^2 \). Then, taking the imaginary part of \( F_1(q^2) \), using the expression for \( \partial_\mu J_{\mu5}^8 \) from Eq. (12) and unitarity we get:

\[
\text{Im} F_1(q^2) = \text{Im} \left[ q_\mu \frac{1}{q^2} (2\gamma | J_{\mu5}^8 | 0) = -\frac{f_8}{q^2} (2\gamma | m_\eta^2 \eta(s_8c_3 - c_0s_3s_8) + m_\eta^2 \eta'(s_3c_8 + c_3c_0s_8) + m_\eta^2 G s_8s_0 | 0) = \pi f_8 | A_\eta \delta(q^2 - m_\eta^2)(c_8c_3 - c_0s_3s_8) + A_\eta \delta(q^2 - m_\eta^2)(s_3c_8 + c_3c_0s_8) + A_\eta \delta(q^2 - m_\eta^2)(s_8s_0) \rangle. \]
\]  

If we employ the sum rule (20), we obtain a simple equation:

\[
(c_8c_3 - c_0s_3s_8) + \beta(s_3c_8 + c_3c_0s_8) + \gamma(s_8s_0) = \xi,
\]  

where
\[ \beta \equiv \frac{A_G}{A_\eta} = \sqrt{\frac{\Gamma_\eta \to \gamma}{16\pi I_0}} \frac{m_0^3}{m_{\eta}^3}, \quad \gamma \equiv \frac{A_G}{A_\eta} = \sqrt{\frac{\Gamma_\eta \to \gamma}{16\pi I_0}} \frac{m_0^3}{m_{\eta}^3}. \]

\[ \xi \equiv \sqrt{\frac{a^2 m_0^4}{96\pi^2 I_0}} \frac{1}{f_0}, \quad \Gamma_\eta \to \gamma = \frac{m_0^3}{64\pi} A_\eta^2. \]

Note that if we include higher resonances in this equation, they will be suppressed as \(1/m_{\text{res}}^2\) by virtue of the mentioned above asymptotic behavior of \(F_1(q^2) \propto 1/q^4\). For the last two terms in (22) we can specify this constraint as follows:

\[ \frac{|s_s s_0|}{|s_3 c_8 + c_3 s_0 s_8|} \lesssim \frac{\beta}{\gamma} \left( \frac{m_{\eta}}{m_\eta} \right)^2. \]

As an additional experimental constraint we use, following [26],[27], the data of the decay ratio \(R_{J/\Psi} = (\Gamma(J/\Psi) \rightarrow \eta'\gamma)/(\Gamma(J/\Psi) \rightarrow \eta\gamma)\).

As it was pointed out in [28], the radiative decays \(J/\Psi \rightarrow \eta(\eta')\gamma\) are dominated by non-perturbative gluonic matrix elements, and the ratio of the decay rates \(R_{J/\Psi} = (\Gamma(J/\Psi) \rightarrow \eta'\gamma)/(\Gamma(J/\Psi) \rightarrow \eta\gamma)\) can be expressed as follows:

\[ R_{J/\Psi} = \left| \frac{\langle 0 | G \bar{G} | \eta' \rangle}{\langle 0 | G \bar{G} | \eta \rangle} \right|^2 \left( \frac{p_{\eta'}}{p_\eta} \right)^3, \]

where \(p_{\eta(\eta')} = M_{J/\Psi}(1 - m_{\eta(\eta')}^2/M_{J/\Psi}^2)/2\). The advantage of this ratio is expected smallness of perturbative and non-perturbative corrections.

The divergencies of singlet and octet components of the axial current in terms of quark fields can be written as:

\[ \partial_{\mu} J_{\mu 5}^s = \frac{1}{\sqrt{6}} (m_u \bar{u} \gamma_5 u + m_d \bar{d} \gamma_5 d - 2m_s \bar{s} \gamma_5 s), \]

\[ \partial_{\mu} J_{\mu 5}^o = \frac{1}{\sqrt{3}} (m_u \bar{u} \gamma_5 u + m_d \bar{d} \gamma_5 d + m_s \bar{s} \gamma_5 s) + \frac{1}{3} \frac{3\alpha_s}{2\sqrt{3} 4\pi} G \bar{G}. \]

Following [26], neglect the contribution of u- and d- quark masses, then the matrix elements of the anomaly term between the vacuum and \(\eta, \eta'\) states are:

\[ \frac{\sqrt{3} \alpha_s}{8\pi} \langle 0 | G \bar{G} \mid \eta \rangle = \langle 0 | \partial_{\mu} J_{\mu 5}^{(0)} \mid \eta \rangle + \frac{1}{\sqrt{2}} \langle 0 | \partial_{\mu} J_{\mu 5}^{(s)} \mid \eta \rangle. \]

\[ \frac{\sqrt{3} \alpha_s}{8\pi} \langle 0 | G \bar{G} \mid \eta' \rangle = \langle 0 | \partial_{\mu} J_{\mu 5}^{(0)} \mid \eta' \rangle + \frac{1}{\sqrt{2}} \langle 0 | \partial_{\mu} J_{\mu 5}^{(s)} \mid \eta' \rangle. \]

Using Eq. (12), (26), (29), (30) we deduce:

\[ R_{J/\Psi} = \left[ \frac{f_0(-s_3 s_8 + c_3 c_8 c_0) + \frac{1}{\sqrt{2}} f_8(s_3 c_8 + c_3 c_0 s_8)}{f_0(-c_3 s_8 - s_3 c_8 c_0) + \frac{1}{\sqrt{2}} f_8(c_3 c_8 - c_3 c_0 s_8)} \right]^2 \times \left( \frac{m_{\eta'}}{m_\eta} \right)^4 \left( \frac{p_{\eta'}}{p_\eta} \right)^3. \]

4 Analysis

For further analysis it is convenient to rewrite the equations (22), (31) in terms of angles \(\theta_1 \equiv \theta_8 + \theta_3, \theta_8\) and \(\theta_0\):

\[ \frac{1}{2}(c_1 + c_2 - c_0(c_2 - c_1)) + \frac{\beta}{2}(s_1 - s_2 + c_0(s_1 + s_2)) + \gamma(s_8 s_0) = \xi. \]
\[
R_{J/\Psi} = \left[ \frac{f_0(c_1 - c_2 + c_0(c_1 + c_2)) + \frac{1}{\sqrt{2}} f_8(s_1 - s_2 + c_0(s_1 + s_2))}{f_0(-s_1 - s_2 - c_0(s_1 - s_2)) + \frac{1}{\sqrt{2}} f_8(c_1 + c_2 - c_0(c_1 - c_2))} \right]^2 \left( \frac{m_{\eta}}{m_{G}} \right)^4 \left( \frac{p_{\eta}}{p_G} \right)^3 ,
\]
where \( \theta_2 \equiv 2\theta_8 - \theta_1 \).

The angles \( \theta_1, \theta_8, \theta_0 \) have the explicit physical meaning. From the definition of the mixing matrix \( U \) one can see that the angle \( \theta_1 \) describes the overlap in the \( \eta - \eta' \) system with an accuracy \( \sim \theta_0^2/2 \) and coincides with their mixing angle as \( \theta_0 \to 0 \). At the same time \( \theta_0 \) is responsible for the glueball admixture to \( \eta - \eta' \) system, and \( s_8s_0 \) describes the contribution of the glueball state \( G \) to the octet component of axial current \( \partial J_{\mu 5}^A \) only.

In the further analysis we will use the following assumptions:

I) As we discussed in Sec. 2, the last term in (14) should be suppressed as \( (m_{\eta}/m_G)^2 \). So we impose the following constraint:

\[
\frac{|s_8s_0|}{|s_3c_8 + c_3c_0s_8|} \lesssim \left( \frac{m_{\eta}}{m_G} \right)^2 .
\]

II) In see 3 we found another constraint, which follows from the asymptotic behavior of \( ImF_1 \) (see [25]):

\[
\frac{|s_8s_0|}{|s_3c_8 + c_3c_0s_8|} \lesssim \beta \left( \frac{m_{\eta}}{m_G} \right)^2 .
\]

III) In our numerical analysis we suppose that \( \gamma \) cannot exceed 1 (i.e., \( \Gamma_{G\to 2\gamma}/m_G^3 \lesssim \Gamma_{\eta\to 2\gamma}/m_\eta^3 \)). This restriction corresponds to the assumption that 2-photon decay widths of pseudoscalar mesons grow like the third power of their masses, or in other words, the glueball coupling to quarks is of the same order as for the meson octet states.

IV) We accept that the decay constants obey the relation \( f_8 \gtrsim f_0 \gtrsim f_\pi \) (for various kinds of justification see, e.g., [3, 9]).

For the purposes of numerical analysis, the values of \( R_{J/\Psi} \) \( (R_{J/\Psi} = 4.8 \pm 0.6) \), masses and two-photon decay widths of \( \eta, \eta' \) mesons are taken from PDG [29]. Using the values \( m_{\eta}, m_{\eta'}, \Gamma_{\eta\to 2\gamma}, \Gamma_{\eta'\to 2\gamma} \), we see that the relation for the constraint (44) is more strict than the constraint (35). Supposing the minimal mass of the glueball to be of order \( m_G \simeq 3m_\eta \simeq 1.5 \) GeV, we get the estimation:

\[
|s_8s_0|/|s_3c_8 + c_3c_0s_8| \lesssim 0.1 .
\]
This work was supported in part by RFBR (Grants 09-02-00732, 09-02-01149), by the Y. K. and O. T. gratefully acknowledge the organizers of the workshop for hospitality and support. The analysis demands $f_\pi < 1.2 f_\pi$ which deviates at 10% level from the results of calculations within the chiral perturbation theory ($f_\pi = 1.34 f_\pi$) \[10\].

The value of the mixing angle $\theta_0$, which is responsible for the glueball admixture to the $\eta - \eta'$, is limited to $\theta_0 < 25^\circ$ for $(f_8, f_0) = (1.0, 1.0) f_\pi$ and to $\theta_0 < 20^\circ$ for $(f_8, f_0) = (1.0, 1.0) f_\pi$.

The improvement of the experimental data of $R_{J/\Psi}$ can significantly limit the constraints for the parameters $\theta_0$, $\theta_8$ and $f_8$, $f_0$.

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