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The structure, mixing angle, mass and couplings of the light scalar $f_0(500)$ and $f_0(980)$ mesons

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ABSTRACT

The mixing angle, mass and couplings of the light scalar mesons $f_0(500)$ and $f_0(980)$ are calculated in the framework of QCD two-point sum rule approach by assuming that they are tetraquarks with diquark–antidiquark structures. The mesons are treated as mixtures of the heavy $|H\rangle = ([su][\overline{su}] + [sd][\overline{sd}])/\sqrt{2}$ and light $|L\rangle = [ud][\overline{ud}]$ scalar diquark–antidiquark components. We extract from corresponding sum rules the mixing angles φ_H and φ_L of these states and evaluate the masses and couplings of the particles $f_0(500)$ and $f_0(980)$.

1. Introduction

Light scalar mesons that reside in the region $m < 1$ GeV of the meson spectroscopy are sources of long-standing problems for the conventional quark model. The standard approach when treating mesons as bound states of a quark and an antiquark $q\bar{q}$ meets with evident troubles to include $f_0(500)$ and $f_0(980)$, as well as some other light particles into this scheme: There are discrepancies between predictions of this model for a mass hierarchy of light scalars and measured masses of these particles. Therefore, for instance, the $f_0(980)$ meson was already considered as a four–quark state with $q^2\bar{q}^2$ content [1].

During passed decades physicists made great efforts to understand features of the light scalar mesons: They were treated as meson–meson molecules [2–5], or considered as diquark–antidiquark bound states [6,7]. These models stimulated not only qualitative analysis of the light scalar mesons, but also allowed one to calculate their parameters using different methods. Thus, in Ref. [8] masses of the $f_0(500)$, $f_0(980)$, $a_0(980)$ and $K_0^*(800)$ mesons were evaluated in the context of the relativistic diquark–antidiquark approach and nice agreements with the data were found. There are growing understanding that the mesons from the light scalars' nonet are exotic particles or at least contain substantial multiquark components: lattice simulations and experimental data seem support these suggestions. Further information on relevant theoretical ideas and models, as well as on experimental data can be found in original and review articles [9–13].

Intensive studies of the light scalars as tetraquark states were carried out using QCD sum rules method [14–22]. Essential part of these investigations confirmed assignment of the light scalars as tetraquark states despite the fact that to explain experimental data in some of them authors had to introduce various modifications to a pure diquark–antidiquark picture and to treat the particles as a mixture of diquark–antidiquarks with different flavor structures [18], or as superpositions of diquark–antidiquark and $q\bar{q}$ components [20–22]. There was also the article (see, Ref. [19]), results of which did not support an interpretation of the light scalars as diquark–antidiquark bound states.

As is seen, theoretical analyses performed even within the same method lead to different conclusions about the internal structures of the mesons from the light scalar nonet. One should add to this picture also large errors from which suffer experimental data on the masses and widths of these particles [23] to understand difficulty of existing problems.

2. Mixing schemes

An approach to the nonet of light scalars as mixtures of tetraquarks belonging to different representations of the color group was recently proposed in Ref. [24]. In accordance with this approach the nonet of the light spin-0 mesons can be considered as tetraquarks composed of the color $(\overline{\mathbf{3}}_c)$ and flavor $(\overline{\mathbf{3}}_f)$ antitriplet scalar diquarks. Then, these tetraquarks in the flavor space form a nonet of the scalar particles $\overline{\mathbf{3}}_f \otimes \mathbf{3}_f = \mathbf{8}_f \oplus \mathbf{1}_f$. In order

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to embrace the second nonet of the scalar mesons occupying the region above 1 GeV spin-1 diquarks belonging to the color–flavor representation $(\mathbf{6}_c,\ \overline{\mathbf{3}}_f)$ can be used. The tetraquarks built of the spin-1 diquarks have the same flavor structure as ones constructed from spin-0 diquarks, and therefore can mix with them.

In the present Letter we restrict ourselves by considering only the first nonet of the scalar particles. Therefore, in what follows we neglect their possible mixing with tetraquarks composed of the spin-1 diquarks. The flavor singlet and octet components of this nonet have the structures

$$|\mathbf{1}_f\rangle = \frac{1}{\sqrt{3}}\left\{[su][\overline{su}]+[ds][\overline{ds}]+[ud][\overline{ud}]\right\},$$
$$|\mathbf{8}_f\rangle = \frac{1}{\sqrt{6}}\left\{[ds][\overline{ds}]+[su][\overline{su}]-2[ud][\overline{ud}]\right\},$$

that in the exact $SU_f(3)$ symmetry can be directly identified with the physical mesons. But the real scalar particles are mixtures of these states, and in the singlet-octet basis and one-angle mixing scheme have the decomposition

$$\begin{pmatrix} |f\rangle \\ |f'\rangle \end{pmatrix} = U(\theta)\begin{pmatrix} |\mathbf{1}_f\rangle \\ |\mathbf{8}_f\rangle \end{pmatrix},\ \ U(\theta)=\begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix}, \tag{1}$$

where, for the sake of simplicity, we denote $f=f_0(500)$ and $f'=f_0(980)$, and θ is the corresponding mixing angle. Alternatively, one can introduce the heavy-light basis

$$|\mathbf{H}\rangle = \frac{1}{\sqrt{2}}\left\{[su][\overline{su}]+[ds][\overline{ds}]\right\},\ \ |\mathbf{L}\rangle = [ud][\overline{ud}], \tag{2}$$

and for the physical mesons get the expansion

$$\begin{pmatrix} |f\rangle \\ |f'\rangle \end{pmatrix} = U(\varphi)\begin{pmatrix} |\mathbf{H}\rangle \\ |\mathbf{L}\rangle \end{pmatrix}. \tag{3}$$

Here we use φ as the mixing angle in the heavy-light basis. An emerged situation is familiar to one from analysis of the mixing problems in the nonet of the pseudoscalar mesons, namely in the $\eta-\eta'$ system [25–27]. The heavy-light basis in the case under consideration is similar to the quark–flavor basis employed there. The mixing angles in the two basis are connected by the simple relation

$$\tan\theta = \frac{\sqrt{2}\tan\varphi+1}{\tan\varphi-\sqrt{2}}. \tag{4}$$

In general, one may introduce also two-angles mixing scheme if it leads to a better description of the experimental data

$$\begin{pmatrix} |f\rangle \\ |f'\rangle \end{pmatrix} = U(\varphi_H,\varphi_L)\begin{pmatrix} |\mathbf{H}\rangle \\ |\mathbf{L}\rangle \end{pmatrix}, U(\varphi_H,\varphi_L)=\begin{pmatrix} \cos\varphi_H & -\sin\varphi_L \\ \sin\varphi_H & \cos\varphi_L \end{pmatrix}. \tag{5}$$

The couplings in the $f-f'$ system can be defined in the form

$$\langle 0|J^i|f(p)\rangle = F_f^i m_f,\ \ \langle 0|J^i|f'(p)\rangle = F_{f'}^i m_{f'},\ \ i=H,L. \tag{6}$$

We suggest that the couplings follow pattern of state mixing in both one- and two-angles scheme. In the general case of two-angles mixing scheme this implies fulfillment of the equality

$$\begin{pmatrix} F_f^H & F_f^L \\ F_{f'}^H & F_{f'}^L \end{pmatrix} = U(\varphi_H,\varphi_L)\begin{pmatrix} F_H & 0 \\ 0 & F_L \end{pmatrix}, \tag{7}$$

where F_H and F_L may be formally interpreted as couplings of the "particles" $|\mathbf{H}\rangle$ and $|\mathbf{L}\rangle$.

Currents $J^H(x)$ and $J^L(x)$ in Eq. (6) that correspond to $|\mathbf{H}\rangle$ and $|\mathbf{L}\rangle$ states are given by the expressions

$$J^H(x) = \frac{\epsilon^{dab}\epsilon^{dce}}{\sqrt{2}}\left\{\left[u_a^T(x)C\gamma_5 s_b(x)\right]\left[\overline{u}_c(x)\gamma_5 C\overline{s}_e^T(x)\right] + \left[d_a^T(x)C\gamma_5 s_b(x)\right]\left[\overline{d}_c(x)\gamma_5 C\overline{s}_e^T(x)\right]\right\}, \tag{8}$$

and

$$J^L(x) = \epsilon^{dab}\epsilon^{dce}\left[u_a^T(x)C\gamma_5 d_b(x)\right]\left[\overline{u}_c(x)\gamma_5 C\overline{d}_e^T(x)\right], \tag{9}$$

where a,b,c,d and e are color indices and C is the charge conjugation operator. Then the interpolating currents for physical states $J^f(x)$ and $J^{f'}(x)$ take the forms

$$\begin{pmatrix} J^f(x) \\ J^{f'}(x) \end{pmatrix} = U(\varphi_H,\varphi_L)\begin{pmatrix} J^H(x) \\ J^L(x) \end{pmatrix}. \tag{10}$$

In the simple case of one-angle mixing scheme Eq. (10) transforms to the familiar superpositions

$$\begin{aligned} J^f(x) &= J^H(x)\cos\varphi - J^L(x)\sin\varphi, \\ J^{f'}(x) &= J^H(x)\sin\varphi + J^L(x)\cos\varphi. \end{aligned} \tag{11}$$

These currents or their more complicated forms in the two-angles mixing scheme may be used in QCD sum rule calculations to evaluate the masses and couplings of the mesons f and f'.

3. Sum rules

At the first stage of our calculations we derive the sum rule for the mixing angle φ of the $f-f'$ system. To this end, we use the heavy-light basis and one-angle mixing scheme and start from the correlation function [28]

$$\Pi(p) = i\int d^4x e^{ip\cdot x}\langle 0|\mathcal{T}\{J^f(x)J^{f'\dagger}(0)\}|0\rangle. \tag{12}$$

The sum rule obtained using $\Pi(p)$ allow us to fix the mixing angle φ. In fact, because the currents $J^f(x)$ and $J^{f'}(x)$ create only $|f\rangle$ and $|f'\rangle$ mesons, respectively, a phenomenological expression for the correlator $\Pi^{\rm Phys}(p)$ equals to zero. Then the second ingredient of the sum rule, namely expression of the correlation function calculated in terms of quark–gluon degrees of freedom $\Pi^{\rm OPE}(p)$ should be equal to zero. Because $\Pi^{\rm OPE}(p)$ depends on the mixing angle φ, it is not difficult to find

$$\tan 2\varphi = \frac{2\Pi_{HL}^{\rm OPE}(p)}{\Pi_{LL}^{\rm OPE}(p)-\Pi_{HH}^{\rm OPE}(p)}, \tag{13}$$

where

$$\Pi_{ij}(p) = i\int d^4x e^{ip\cdot x}\langle 0|\mathcal{T}\{J^i(x)J^{j\dagger}(0)\}|0\rangle. \tag{14}$$

In deriving of Eq. (13) we benefited from the fact that $\Pi_{HL}^{\rm OPE}(p)=\Pi_{LH}^{\rm OPE}(p)$, which can be proved by explicit calculations. After applying the Borel transformation and performing required continuum subtractions one can employ it to evaluate φ.

Having found the mixing angle we proceed and evaluate the spectroscopic parameters of the mesons f and f'. The correlation functions

$$\begin{aligned} \Pi_f(p) &= i\int d^4x e^{ip\cdot x}\langle 0|\mathcal{T}\{J^f(x)J^{f\dagger}(0)\}|0\rangle, \\ \Pi_{f'}(p) &= i\int d^4x e^{ip\cdot x}\langle 0|\mathcal{T}\{J^{f'}(x)J^{f'\dagger}(0)\}|0\rangle, \end{aligned} \tag{15}$$

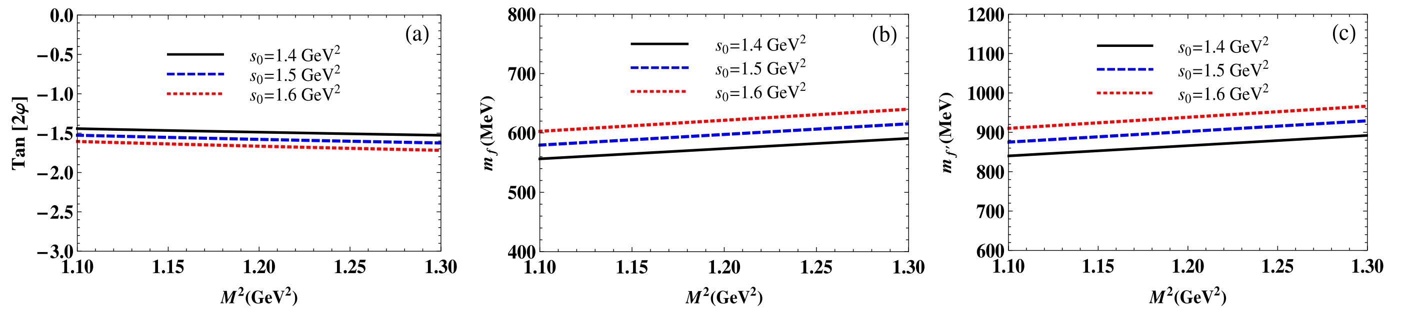

Fig. 1. The $\tan 2\varphi$ (a), and the masses m_f (b) and $m_{f'}$ (c) in the two-angles mixing scheme as functions of the Borel parameter M^2 at fixed s_0.

are appropriate for these purposes and can be utilized to derive the relevant sum rules. The expression of $\Pi_f(p)$ in terms of the physical parameters of the f meson is given by the following simple formula

$$\Pi_f^{\rm Phys}(p) = \frac{\langle 0|J^f|f(p)\rangle\langle f(p)|J^{f\dagger}|0\rangle}{m_f^2 - p^2} + \ldots,$$

where the dots stand for contributions of the higher resonances and continuum states. Using the interpolating current and matrix elements of the f meson from Eqs. (11) and (6) it is a easy task to show that

$$\Pi_f^{\rm Phys}(p) = \frac{m_f^2}{m_f^2 - p^2}\left(F_H \cos^2\varphi + F_L \sin^2\varphi\right)^2 + \ldots.$$

After calculating the correlation function $\Pi_f^{\rm OPE}(p)$ and applying the Borel transformation in conjunction with continuum subtraction one gets the sum rule

$$m_f^2 \left(F_H \cos^2\varphi + F_L \sin^2\varphi\right)^2 e^{-m_f^2/M^2} = \Pi_f(s_0, M^2), \tag{16}$$

where $\Pi_f(s_0, M^2) = \mathcal{B}\ \Pi_f^{\rm OPE}(p)$ is the Borel transformed and subtracted expression of $\Pi_f^{\rm OPE}(p)$ with M^2 and s_0 being the Borel and continuum threshold parameters, respectively. This sum rule and another one obtained from Eq. (16) by means of the standard operation $d/d(-1/M^2)$ can be used to evaluate the mass of the f meson.

The similar analysis for f' yields

$$m_{f'}^2 \left(F_H \sin^2\varphi + F_L \cos^2\varphi\right)^2 e^{-m_{f'}^2/M^2} = \Pi_{f'}(s_0, M^2). \tag{17}$$

From Eqs. (16) and (17) it is also possible to extract $\left(F_H \cos^2\varphi + F_L \sin^2\varphi\right)^2$ and $\left(F_H \sin^2\varphi + F_L \cos^2\varphi\right)^2$ for evaluating of the couplings F_H and F_L, but they may suffer from large uncertainties: We instead evaluate F_H and F_L from sum rules for the scalar "particles" $|\mathbf{H}\rangle$ and $|\mathbf{L}\rangle$, using to this end correlation functions $\Pi_H(p)$ and $\Pi_L(p)$ given by Eq. (15), where $J^f(x)$ and $J^{f'}(x)$ are replaced by $J^H(x)$ and $J^L(x)$, respectively.

4. Numerical results

In calculations we utilize the light quark propagator

$$S_q^{ab}(x) = i\delta_{ab}\frac{\not x}{2\pi^2 x^4} - \delta_{ab}\frac{m_q}{4\pi^2 x^2} - \delta_{ab}\frac{\langle\overline{q}q\rangle}{12} + i\delta_{ab}\frac{\not x m_q\langle\overline{q}q\rangle}{48} - \delta_{ab}\frac{x^2}{192}\langle\overline{q}g_s\sigma Gq\rangle + i\delta_{ab}\frac{x^2\not x m_q}{1152}\langle\overline{q}g_s\sigma Gq\rangle - i\frac{g_s G_{ab}^{\alpha\beta}}{32\pi^2 x^2}\left[\not x\sigma_{\alpha\beta} + \sigma_{\alpha\beta}\not x\right] - i\delta_{ab}\frac{x^2\not x g_s^2\langle\overline{q}q\rangle^2}{7776} - \delta_{ab}\frac{x^4\langle\overline{q}q\rangle\langle g_s^2 G^2\rangle}{27648} + \ldots, \tag{18}$$

and take into account quark, gluon and mixed operators up to dimension twelve. The vacuum expectations values of the operators used in numerical computations are well known: $\langle\bar{q}q\rangle = -(0.24 \pm 0.01)^3$ GeV^3, $\langle\bar{s}s\rangle = 0.8\ \langle\bar{q}q\rangle$, $\langle\overline{q}g_s\sigma Gq\rangle = m_0^2\langle\overline{q}q\rangle$, $\langle\overline{s}g_s\sigma Gs\rangle = m_0^2\langle\bar{s}s\rangle$, $\langle\alpha_s G^2/\pi\rangle = (0.012 \pm 0.004)$ GeV^4, $\langle g_s^3 G^3\rangle = (0.57 \pm 0.29)$ GeV^6, where $m_0^2 = (0.8 \pm 0.1)$ GeV^2.

The working regions for the Borel and continuum threshold parameters are fixed in the following form

$$M^2 = (1.1 - 1.3)\ \text{GeV}^2,\ \ s_0 = (1.4\text{–}1.6)\ \text{GeV}^2, \tag{19}$$

that satisfy standard requirements of sum rules computations. For example, at the lower limit of the Borel parameter the sum of the dimension-10, 11 and 12 terms in $\Pi_{LL}(s_0, M^2) - \Pi_{HH}(s_0, M^2)$ does not exceed 5% of all contributions. At the upper bound of the working window for M^2 the pole contribution to the same quantity is larger than 0.12 of the whole result, which is typical for multiquark systems. Variation of the auxiliary parameters M^2 and s_0 within the regions (19), as well as uncertainties of the other input parameters generate theoretical errors of sum rules computations. The $\tan 2\varphi$ extracted using Eq. (13), as is seen from Fig. 1 (a), demonstrates a mild dependence on M^2. As a result, it is not difficult to estimate that

$$\varphi = -27^\circ.66 \pm 1^\circ.24. \tag{20}$$

This value of φ in the heavy-light basis is equivalent to $\theta = -33^\circ.00 \pm 1^\circ.17$ in the singlet-octet basis. Using Eq. (20) it is not difficult to evaluate the mesons' masses in the one-angle mixing scheme that read

$$m_f = (597 \pm 81)\ \text{MeV}, \qquad m_{f'} = (902 \pm 125)\ \text{MeV}. \tag{21}$$

As is seen, the one-angle mixing scheme, if take into account the central values from Eq. (21), does not describe correctly the experimental data: it overshoots the mass of the $f_0(500)$ meson and, at the same time, underestimates the mass of the $f_0(980)$ meson. The agreement can be improved by introducing the two mixing angles φ_H and φ_L. It turns out that to achieve a nice agreement with the available experimental data it is enough to vary φ_H and φ_L within the limits (20):

$$\varphi_H = -28^\circ.87 \pm 0^\circ.42,\ \ \varphi_L = -27^\circ.66 \pm 0^\circ.31. \tag{22}$$

For m_f and $m_{f'}$ the sum rules with two mixing angles φ_H and φ_L lead to predictions

$$m_f = (518 \pm 74)\ \text{MeV}, \qquad m_{f'} = (996 \pm 130)\ \text{MeV}, \tag{23}$$

which are compatible with experimental data. The theoretical errors in Eq. (23) accumulate uncertainties connected with M^2 and s_0, as well as arising from other input parameters. The dependence of m_f and $m_{f'}$ on the auxiliary parameters M^2 and s_0 does not exceed limits allowed for such kind of calculations: In Figs. 1(b) and 1(c) we plot m_f and $m_{f'}$ as functions of the Borel parameter to confirm a stability of corresponding sum rules.

In the two-angle mixing scheme the system of the physical particles $f - f'$ is characterized by four couplings (7). After determining the mixing angles φ_H and φ_L that fix the matrix $U(\varphi_H, \varphi_L)$, quantities which should be found from the relevant sum rules are only the couplings F_H and F_L. As we have mentioned above to this end we consider two additional sum rules by treating basic states $|\mathbf{H}\rangle$ and $|\mathbf{L}\rangle$ as real “particles” and obtain

$$F_H = (1.35 \pm 0.34) \cdot 10^{-3}\ \text{GeV}^4, \quad F_L = (0.68 \pm 0.17) \cdot 10^{-3}\ \text{GeV}^4. \tag{24}$$

The coupling F_H calculated in the present work is comparable with one found in Ref. [16] using the same interpolating current (8) and vacuum condensates up to dimension six and is given by $F_H = (1.51 \pm 0.14) \cdot 10^{-3}\ \text{GeV}^4$.

The mixing angles $(\varphi_H,\ \varphi_L)$, the masses $(m_f,\ m_{f'})$ and the couplings $(F_H,\ F_L)$ complete the set of the spectroscopic parameters of the $f_0(500)$ and $f_0(980)$ mesons.

5. Concluding notes

The investigation performed in the present Letter has allowed us to calculate the mass and couplings of the $f_0(500)$ and $f_0(980)$ mesons by treating them as the mixtures of the diquark–antidiquarks $|\mathbf{H}\rangle$ and $|\mathbf{L}\rangle$. We have demonstrated that by choosing the heavy-light basis and mixing angles $\varphi_H = -28^\circ.87 \pm 0^\circ.42$ and $\varphi_L = -27^\circ.66 \pm 0^\circ.31$ a reasonable agreement with experimental data can be achieved even information on the $f_0(500)$ meson suffers from large uncertainties [23]. The assumption on structures of the light mesons made in the present work determines also their possible decay mechanisms. Indeed, it is known that the dominant decay channels of the $f_0(500)$ and $f_0(980)$ mesons are $f_0(500) \to \pi\pi$ and $f_0(980) \to \pi\pi$ processes. In experiments the decay $f_0(980) \to K\overline{K}$ was seen, as well. The mixing of the $|\mathbf{H}\rangle$ and $|\mathbf{L}\rangle$ diquark–antidiquark states to form the physical mesons implies that all of these decays can run through the superallowed Okubo–Zweig–Iizuka (OZI) mechanism: Without the mixing the decay $f_0(980) \to \pi\pi$ can proceed due to one gluon exchange, whereas $f_0(980) \to K\overline{K}$ is still OZI superallowed process [16]. The another problem that finds its natural explanation within the mixing framework is a large difference between the full width of the mesons $f_0(500)$ and $f_0(980)$, which amount to $\Gamma = 400 - 700$ MeV and $\Gamma = 10 - 100$ MeV [23], respectively. In fact, the strong couplings $g_{f_0(500)\pi\pi}$ and $g_{f_0(980)\pi\pi}$ that determine the width of the dominant partial decays $f_0(500) \to \pi\pi$ and $f_0(980) \to \pi\pi$ depend on the mixing angle φ_L in the form

$$g^2_{f_0(500)\pi\pi} \propto \frac{1}{\sin^2 \varphi_L}, \quad g^2_{f_0(980)\pi\pi} \propto \frac{1}{\cos^2 \varphi_L}. \tag{25}$$

In the model under consideration this dependence is a main source that generates the numerical difference between the partial width of aforementioned processes, and hence between the full width of the mesons $f_0(500)$ and $f_0(980)$.

Analysis of the partial decays of the mesons $f_0(500)$ and $f_0(980)$, as well as calculation of the spectroscopic parameters of other light scalar mesons deserves further detailed investigations results of which will be published elsewhere.

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