On the $\pi\pi$ contribution to the QCD sum rules for the light tetraquark

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Abstract

We perform a QCD sum rule analysis for the $f_0$ light tetraquark taking into account the contribution arising from the two pion intermediate state. With the interpolating currents of the different chiral combinations of scalar and pseudoscalar diquarks, it is demonstrated that the interpolating current with maximum chirality has a large coupling to the two pion state, but the current with zero chirality interacts only weekly with this state. Taking into account the form factor in the $f_0$–two pion vertex, it is shown that the $f_0$–coupling to the two pion state leads to an increase of the lightest tetraquark mass by a value of about 100 MeV. The analysis of the resulting sum rule shows that the $\sigma(f_0(600))$–meson state might be treated as the four–quark bound state in the instanton field which has a rather strong coupling to the two pion state.

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1 Introduction

One of the direct ways to investigate the properties of the exotic states is the QCD sum rule (SR) approach. Recently, this method was applied to the study of the scalar tetraquark states with different interpolating currents. It was shown that the light $0^{++}$ mesons ($\sigma(f_0(600))$, $\kappa(800)$, $a_0(980)$ and $f_0(980)$) might be interpreted as four–quark exotic states [1, 2, 3, 4, 5]. The investigations give additional support to possible large not $q$–$\bar{q}$ component in these states in the line of their previous study within different constituent quark models (see [6] and references therein) and the approach based on the $1/N_c$ expansion [7]. In most of these SR analysis only the exotic resonance contribution to the phenomenological part of the SR has been considered within the narrow width approximation for the pole. However, it is well known that the tetraquark can couple strongly to two meson colorless states. Such coupling is super–allowed according to the OZI rule and might be responsible for the large observed width of the lightest candidate for tetraquark the $\sigma(f_0(600))$ meson. Therefore one may expect a rather large effect of such coupling on the extracted properties of the tetraquark within both QCD SR and lattice approaches [8]. We should emphasize that the problem of the possible large contribution of intermediate hadronic states to the correlator of the multiquark current is quite general and, for example, has been discussed recently for the pentaquark case in ref. [9].

When the QCD sum rule is applied to the tetraquarks, it is well known that the operators of higher dimensions could yield large contributions in the operator product expansion (OPE) and would spoil convergence in the OPE [10, 11]. In our previous paper [5], it was demonstrated that the interpolating currents with equal weights of scalar and pseudoscalar diquarks yield strong cancelation of the contributions coming from high dimension operators and direct instantons. It was shown that such cancelation is related to the specific chirality structure of the interpolating current. As a result, we can avoid the problem of huge contribution from high dimension operators to the tetraquark SR. In this way, a quite stable SR for the light tetraquark meson with $u\bar{u}d\bar{d}$ quark content has been obtained.

In this here we extend the previous study to include the two pion intermediate contribution to the QCD SR. In section II the contribution to the SR from the two pion state for the interpolating current with arbitrary mixing of scalar and pseudoscalar $ud$–diquark will be obtained by using soft pion theorems. In section III the numerical analysis of the corresponding SR will be done analyzing the effect of the form factor in the tetraquark–two pion vertex, and in section IV we will give the summary of our results.

2 Two pion contribution to QCD sum rule for light tetraquark

The starting point of the QCD sum rule for the scalar meson is the dispersion relation of the correlator:

$$\Pi(q^2) = \frac{1}{\pi} \int_0^\infty ds^2 \frac{\text{Im}\Pi(s^2)}{s^2 - q^2}$$

where the correlator is defined by

$$\Pi(q^2) = i \int d^4x e^{ix} \langle 0 | T J_{f_0}(x) J_{f_0}^\dagger(0) | 0 \rangle$$
where $J_{f_0}(x)$ is the interpolating current for the scalar meson. The imaginary part of the correlator is given by the spectral sum over various intermediate particle states

$$
\frac{1}{\pi} \text{Im}\Pi(q^2) = (2\pi)^3 \sum_n \delta^4(q - P_n) \langle 0 | J_{f_0}(0) | n \rangle \langle n | J_{f_0}^\dagger(0) | 0 \rangle .
$$

(3)

In the single narrow width resonance approximation and with the assumption of hadron–quark duality, the contribution to phenomenological part of the SR coming from the first three diagrams pictured in Fig. 1 is usually considered. In this case the imaginary part is the following,

$$
\frac{1}{\pi} \text{Im}\Pi(s^2) = 2 f_{f_0}^2 m_{f_0}^8 \delta(s^2 - m_{f_0}^2) + \theta(s^2 - s_0^2) \frac{1}{\pi} \text{Im}\Pi^{\text{OPE}}(s^2),
$$

(4)

where $f_{f_0}$ is the residue of the resonance, $m_{f_0}$ is its mass, $s_0$ is the continuum threshold and $\Pi^{\text{OPE}}$ is the correlator within the standard operator expansion (OPE).

Figure 1: The contributions to the phenomenological part of the SR for a light tetraquark: (a) resonance, (b) and (c) continuum, and (d) two pion contributions.

In ref. [5] we have shown that, the correlator of the interpolating current consisting of scalar and pseudoscalar diquarks

$$
J_{f_0} = \alpha J_S + \beta J_{PS},
$$

(5)

where

$$
J_S = \epsilon_{abc} \epsilon_{ade} (u_b^T \Gamma_S d_c)(\bar{u}_d \Gamma_S \bar{d}_e^T),
$$

$$
J_{PS} = \epsilon_{abc} \epsilon_{ade} (u_b^T \Gamma_{PS} d_c)(\bar{u}_d \Gamma_{PS} \bar{d}_e^T),
$$

(6)

and $\Gamma_S = C\gamma^5$, $\Gamma_{PS} = C$, $\Gamma_i = \gamma^0 \Gamma^i_\gamma \gamma^0$, has very specific properties for a particular choice of the mixing parameters $\alpha = \pm \beta$. Indeed, it was found that these choice of mixing parameters leads to a cancelation of the contribution from the high dimensions operators in the OPE, as well as that of some dangerous direct instanton contributions.

The contribution from the two pion state to the imaginary part of the correlator (the last diagram in Fig.1) is presented by phase space product of the two pions as

$$
\frac{1}{\pi} \text{Im}\Pi^{2\pi}(q^2) = (2\pi)^3 |\lambda_{2\pi}|^2 \int \frac{d^3p_1}{(2\pi)^3 2E_1} \int \frac{d^3p_2}{(2\pi)^3 2E_2} \delta^4(q - p_1 - p_2)
$$

$$
= \frac{|\lambda_{2\pi}|^2}{16\pi^2} \sqrt{1 - \frac{4m_{\pi}^2}{q^2}} \theta(q^2 - 4m_{\pi}^2),
$$

(7)

\footnote{We do not consider the possible effects of the interaction between two pions for simplicity. This assumption was also used in the paper [12] to estimate the effect in the SR arising from the coupling of two uncorrelated pions to the quark–antiquark $\sigma$–meson state.}
where $\lambda_{2\pi} = \langle 0 | J_{f_0}(0) | \pi\pi \rangle$ is the correspondent residue.

To obtain the value of $\lambda_{2\pi}$, we rewrite the interpolating current in terms of pion fields by using the Fierz transformation with the two flavor quark field $\psi$.

\[
J_{f_0} = \alpha J_S + \beta J_{PS} = \frac{1}{16} \left[ (\alpha - \beta) \left( (\bar{\psi}\gamma_\mu \pi^\mu_\alpha)^2 - (\bar{\psi}\gamma_\mu \pi^\mu_\alpha)^2 \right) + (\alpha + \beta) \left( (\bar{\psi}\gamma^5 \gamma_\mu \pi^\mu_\alpha)^2 - (\bar{\psi}\gamma^5 \gamma_\mu \pi^\mu_\alpha)^2 \right) + \frac{1}{2} (\alpha - \beta) \left( (\bar{\psi} \sigma_{\rho\sigma} \pi^\rho)^2 - (\bar{\psi} \sigma_{\rho\sigma} \pi^\rho)^2 \right) \right].
\] (8)

The relevant part in the Eq. 8 which includes the pionic interpolating currents is

\[
J_{f_0}^\pi = \frac{1}{16} \left[ (\alpha - \beta) \left( (\bar{\psi}\gamma^5 \pi^\mu_\alpha)^2 - (\bar{\psi}\gamma^5 \pi^\mu_\alpha)^2 \right) + (\alpha + \beta) \left( (\bar{\psi} \gamma^5 \gamma_\mu \pi^\mu_\alpha)^2 - (\bar{\psi} \gamma^5 \gamma_\mu \pi^\mu_\alpha)^2 \right) + \frac{1}{2} (\alpha - \beta) \left( (\bar{\psi} \sigma_{\rho\sigma} \pi^\rho)^2 - (\bar{\psi} \sigma_{\rho\sigma} \pi^\rho)^2 \right) \right].
\] (9)

By transforming well-known PCAC relations at the limit $m_u = m_d = m_q$,

\[
\partial^\mu A^1_\mu = i m_q (\bar{\psi}\gamma^5 \pi^\mu_\alpha), \quad \partial^\mu A^2_\mu = m_q (\bar{\psi}\gamma^5 \pi^\mu_\alpha - \bar{\psi}\gamma^5 d\gamma^5 u), \quad \partial^\mu A^3_\mu = i m_q (\bar{\psi}\gamma^5 d\gamma^5 u)
\]
into operator forms, \( \partial^\mu A^a_\mu = f_\pi m_q \phi^a \) and \( A^a_\mu = i p^a_\mu f_\pi \phi^a \) with the pion field of isospin index \( a, \phi^a \), we can rewrite the above interpolating current in terms of pion fields as

\[
J_{f_0}^\pi = \frac{1}{16} \left[ \frac{f_\pi^2 m_q^4}{m_q^2} (\alpha - \beta) \left( 2 \phi^+ \phi^- \phi^+ \phi^- + \phi^0 \phi^0 \right) + 4(\alpha + \beta) f_\pi^2 \left( 2 p_{\pi^+} \cdot p_{\pi^-} \phi^+ \phi^- + p_{\pi^0} \cdot p_{\pi^0} \phi^0 \phi^0 \right) \right].
\]

Therefore, we have

\[
|\lambda_{2\pi}|^2 = 6 \left[ (\alpha - \beta)^2 \left( \frac{\langle q \bar{q} \rangle^2}{4 f_\pi^2} \right)^2 + (\alpha + \beta)^2 \left( \frac{f_\pi^2}{4} \right)^2 \left( q^2 - 2m_{\pi}^2 \right)^2 \right],
\] (10)

where the Gell–Mann–Oakes–Renner relation

\[
f_\pi^2 m_q^2 = -2m_q \langle \bar{q}q \rangle
\] (11)

has been used. The imaginary part of the correlator coming from the two intermediate pion state becomes

\[
\frac{1}{\pi} \text{Im} \Pi^{2\pi}(q^2) = \frac{6}{16^2 \pi^2} \left[ \frac{\langle q \bar{q} \rangle^4}{f_\pi^4} (\alpha - \beta)^2 + \frac{f_\pi^4}{4} (q^2 - 2m_{\pi}^2)^2 (\alpha + \beta)^2 \right] \times \sqrt{1 - \frac{4m_{\pi}^2}{q^2}} \theta(q^2 - 4m_{\pi}^2).
\] (12)

\(^2\) A similar approach was used in ref. \[13\] to obtain the instanton contribution to the weak decay $K \to \pi\pi$. 

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\[ \]
Incorporating this contribution from the two pion state, the QCD sum rule for the light tetraquark becomes

$$\frac{1}{\pi} \int_0^{s_0} ds^2 \, e^{-s^2/M^2} \text{Im} \Pi^{\text{OPE}}(s^2) + \hat{B}[\Pi^{I+I}(q)] - \frac{1}{\pi} \int_{4m_q^2}^{s_0} ds^2 \, e^{-s^2/M^2} \text{Im} \Pi^{2\pi}(s^2) = 2f_0^8 m_0^8 e^{-m_0^2/M^2},$$

where $\hat{B}[\Pi^{I+I}(q^2)]$ means the Borel transformed instanton effect. Here the contributions from the two pion state and the continuum are transferred to the OPE side in the sum rule.

In the numerical analysis below we will use the results for the standard OPE and direct instanton contributions to the correlator obtained in our previous paper [5]. For massless $u$– and $d$–quarks the relevant parts of the correlator are given by

$$\frac{1}{\pi} \text{Im} \Pi^{\text{OPE}}(q^2) = (\alpha^2 + \beta^2) \left[ \frac{(q^2)^4}{2^{12} \cdot 5 \cdot 3 \pi^6} + \frac{\langle g^2 G^2 \rangle}{2^{11} \cdot 3 \pi^6} (q^2)^2 \right]$$

$$+ (\alpha^2 - \beta^2) \left[ \frac{\langle qq \rangle^2}{12 \pi^2} q^2 - \frac{\langle qq \rangle \langle igq \sigma \cdot Gq \rangle}{12 \pi^2} + \frac{59 (\langle igq \sigma \cdot G_q \rangle)^2}{2^9 \cdot 3^2 \pi^2} \delta(q^2) + \frac{7 \langle g^2 G^2 \rangle \langle \bar{q}q \rangle^2}{2^{5} \cdot 3^3 \pi^2} \delta(q^2) \right],$$

and

$$\Pi^{I+I}(q) = (\alpha^2 - \beta^2) \frac{32n_{\text{eff}} \rho_c^4}{\pi^8 m_q^2} f_0(q)$$

$$+ [19(\alpha^2 + \beta^2) - 6\alpha \beta] \frac{n_{\text{eff}} \rho_c^4 \langle \bar{q}q \rangle^2}{18 \pi^4 m_q^{2}} f_0(q),$$

where $n_{\text{eff}}$ is the effective instanton density, $m_q^s$ is the mass parameter in the quark zero mode Green’s function in the instanton field, and $\rho_c$ is average instanton size. The functions $f_0(q)$, $f_0(q)$ are defined by

$$f_0(q) = \int d^4z_0 \int d^4x \frac{e^{iq \cdot x}}{x^6 [z_0^2 + \rho_c^2]^3 [(x - z_0)^2 + \rho_c^2]^3},$$

$$f_0(q) = \int d^4z_0 \int d^4x \frac{e^{iq \cdot x}}{[z_0^2 + \rho_c^2]^3 [(x - z_0)^2 + \rho_c^2]^3}.$$

There are two types of singularities in these functions. One arises from the origin and another from a finite distance from the origin. Note that we subtract the contribution from the pole at the origin in order to avoid double counting with the contributions from the condensates in the standard OPE [14, 15].

It is evident from Eq (14) and Eq (15) that if the relation $\alpha^2 = \beta^2$ holds true, then most part of the high dimension operators in OPE and part of direct instanton contribution disappear from the SR. In this case only the perturbative and gluon condensate chirality conserving contributions remain in the OPE due to the specific chirality structure of the tetraquark interpolating current Eq. [5]. Indeed, this current can be decomposed into two parts with different chirality structures [5]

$$J_{f_0} \sim -(\alpha - \beta)(u_R^TCd_R^\dagger \bar{u}_L C d_L^\dagger + u_R^TCd_R^\dagger \bar{u}_R C d_R^\dagger)$$

$$+ (\alpha + \beta)(u_R^TCd_R^\dagger \bar{u}_L C d_L^\dagger + u_L^TCd_L^\dagger \bar{u}_R C d_R^\dagger).$$
For the values $\alpha = \beta$ or $\alpha = -\beta$ the current carries zero value of chirality or four units of the chirality, respectively. We call the first(second) case as the minimum(maximum) chirality current. It is easy to verify that for these particular cases, most of the high dimensional condensates and direct instanton contributions are forbidden. As result, one has a good convergence of the OPE and the possible stability of the SR (see discussion in [5]). Therefore, below we consider the two pion contribution to the SR only for these values of the mixing parameters.

3 Numerical analysis of the sum rule with two pion contribution

For the numerical analysis of the SR for $\alpha^2 = \beta^2$ we use the following value for the gluon condensate and the average size of the instanton

$$\langle g^2 G^2 \rangle = 0.5 \text{ GeV}^4, \quad \rho_c = 1.6 \text{ GeV}^{-1},$$

and the relation between instanton parameters given by the simplest version of Shuryak’s instanton liquid model [16,17]

$$\frac{2 n_{eff}}{m_q} = \frac{3}{2\pi^2 \rho_c^2}.$$  \hspace{1cm} (19)

We fix the value of the threshold by $s_0 = 1 \text{ GeV}$ since our results below show only weak dependence on this parameter within the interval $s_0 = 1 \div 1.5 \text{ GeV}$. In Figs. 2 and 3, the contribution from the OPE together with direct instantons (the first two terms on the left hand side in Eq.(13)) and the contribution from two pion state (dashed line) are shown as functions of the Borel mass $M$ for the maximum chirality case of $\alpha = -\beta = 1$. As result of the huge contribution coming from the two pion state, the left hand side (LHS) of the SR Eq.(13) becomes negative and therefore it is impossible to obtain information

![Figure 2: The OPE and direct instanton contribution to the SR with $\alpha = -\beta = 1$.](image)

![Figure 3: The two pion contribution to the phenomenological side of the SR with $\alpha = -\beta = 1$. The solid (dashed) line contains (does not contain) the effect of the form factor.](image)
about the resonance state with such SR. The opposite situation we observe for the zero chirality case, $\alpha = \beta = 1$. In Figs. 4 and 5, the contributions from the OPE with direct instantons and the two pions are shown for this case. Now the contribution from the two pion state is very small and the correspondent SR shows the signal for a possible bound state, as shown in Fig. 6 by the dashed line.

However, we should point out that in the above calculation the point–like local tetraquark–two pions vertex has been considered. The simple way to include the effect of nonlocality into our consideration is to introduce a form factor in the tetraquark–two pions vertex. Such form factor for the scalar tetraquark $\sigma(f_0(600))$ meson can be chosen of the monopole form

$$F_{f_0\pi\pi}(s^2) = \frac{1}{1 + \langle r_\pi^2 \rangle s^2 / 6}$$

with the slope given by scalar radius of $\pi$ meson, $\langle r_\pi^2 \rangle \approx 0.75 \text{ fm}^2$ [18]. Before fitting the mass, it is necessary to fix the Borel window with the pole contribution dominance [11]. These windows for the $\alpha = \beta$ and for the $\alpha = -\beta$ cases lie in $2m_\pi < M < 0.93 \text{ GeV}$ and $0.5 \text{ GeV} < M < 0.85 \text{ GeV}$, respectively. Our final result for mass of the tetraquark extracted from the SR with the two pion contribution and the form factor effect is presented in Figs. 6 and 7. It follows from the previous figures that the effect from the form factor in the two pion state is rather small for the interpolating current with $\alpha = \beta = 1$. Unfortunately, there is no good plateau of stability for the mass in Fig.6. Therefore, it is rather difficult to extract the information about $f_0(600)$ for this current. On the other hand, for the interpolating current of $\alpha = -\beta = 1$, the effect from the form factor suppresses the two pion contribution and makes the SR have physical meaning with a very good stability plateau presented in Fig.7. The fitted mass of resonance for the $\alpha = -\beta$ case, $m_{f_0} \approx 800 \text{ MeV}$, is a little bit larger than the value obtained recently from analysis of the Roy equations [19], but still lies within the interval of mass for the $\sigma(f_0(600))$–meson, $m_{f_0}(600) = 400 \div 1200 \text{ MeV}$ given by PDG [20]. A more exact
The mass obtained from the SR for $\alpha = \beta = 1$ including (solid line) and not including (dashed line) the effect of the form factor as a function of the Borel parameter.

Figure 6: The mass obtained from the SR for $\alpha = \beta = 1$ including (solid line) and not including (dashed line) the effect of the form factor as a function of the Borel parameter. The dashed line corresponds to the mass obtained from the SR not including the two pion contribution.

To clarify the origin of $f_0(600)$ state, in Fig. 8 we present separately the contributions coming from standard OPE, direct instantons and two pion state for $\alpha = -\beta$ case. It is evident that direct instantons give the dominant contribution to the SR. In spite of the fact that the two pion state contributes smaller than direct instantons, its contribution is bigger than the standard OPE within practically the full Borel window. It turns out that we may treat in this case the $f_0(600)$ state as the tetraquark bound state in the instanton field with a rather strong coupling to the two pion state. The average size of instanton in the QCD vacuum is rather small $\approx 0.3$ fm \cite{16}, therefore, the size of the lightest tetraquark should be also very small. It would be interesting to find a possible experimental signature for such small size for the $f_0(600)$. Finally, we should emphasize that the crucial role of the direct instantons in the tetraquark structure, as shown above, does not allow us to agree with the results of many tetraquark studies, where the direct instanton contribution was not included.
4 Conclusion

In summary, the estimate of the two pion contribution to the QCD sum rule for a light tetraquark by using the soft pion PCAC relations has been obtained for the different tetraquark interpolating currents. It has been demonstrated that such contribution depends crucially on the chirality structure of the interpolating current. We show that the SR for the interpolating current with maximum chirality provides a very good stable plateau for the mass of the tetraquark when the corresponding form factor in the tetraquark–two pion vertex is introduced. The important role of the direct instantons in light tetraquark dynamics has been demonstrated.

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