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Analysis the $f_0(980)$ and $a_0(980)$ mesons as four-quark states with the QCD sum rules

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Abstract. In this article, we take the point of view that the scalar mesons $f_0(980)$ and $a_0(980)$ are the diquark–antidiquark states $(qq)_{\bar{3}}(\bar{q}\bar{q})_3$, and we devote our attention to the determination of their masses in the framework of the QCD sum rule approach with the interpolating currents constructed from scalar–scalar type and pseudoscalar–pseudoscalar type diquark pairs respectively. The numerical results indicate that the scalar mesons $f_0(980)$ and $a_0(980)$ may have two possible diquark–antidiquark substructures.

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1 Introduction

The light flavor scalar mesons present a remarkable exception for the constituent quark model and the structures of those mesons have not been unambiguously determined yet [1]. Experimentally, the strong overlaps with each other and the broad widths (for the $f_0(980)$, $a_0(980)$ et cetera, the widths are comparatively narrow) make that their spectra cannot be approximated by the Breit–Wigner formula. The numerous candidates with the same quantum numbers $J^{PC} = 0^{++}$ below 2 GeV cannot be accommodated in one $q\bar{q}$ nonet; some are supposed to be glueballs, molecules and multiquark states. The more elusive things are the constituent structures of the mesons $f_0(980)$ and $a_0(980)$ with almost degenerate masses. In the naive quark model, we have $a_0 = (u\bar{u} - d\bar{d})/\sqrt{2}$ and $f_0 = s\bar{s}$, while in the framework of four-quark models, the mesons $f_0(980)$ and $a_0(980)$ could either be compact objects i.e. nucleon-like bound states of quarks with the symbolic quark structures $f_0 = s\bar{s}(u\bar{u} + d\bar{d})/\sqrt{2}$ and $a_0 = s\bar{s}(u\bar{u} - d\bar{d})/\sqrt{2}$ [2], or spatially extended objects i.e. deuteron-like bound states of hadrons; the $f_0(980)$ and $a_0(980)$ mesons are usually taken as $K\bar{K}$ molecules [3]. The hadronic dressing mechanism takes the point of view that the mesons $f_0(980)$ and $a_0(980)$ have small $q\bar{q}$ cores of typical $q\bar{q}$ meson size, and the strong couplings to the hadronic channels enrich the pure $q\bar{q}$ states with other components and spend part (or most part) of their lifetime as virtual $K\bar{K}$ states [4]. In the hybrid model, those mesons are four-quark states $(qq)_{\bar{3}}(\bar{q}\bar{q})_3$ in S-wave near the center, with some constituent $q\bar{q}$ in P-wave, but further out they rearrange themselves into $(q\bar{q})_1(q\bar{q})_1$ states and finally as meson–meson states [5]. All those interpretations have both outstanding advantages and obvious shortcomings in one way or the other.

There maybe exist two scalar nonets below 1.7 GeV. The attractive interactions of one gluon exchange favor the formation of diquarks in the color antitriplet $\overline{3}_c$, the flavor antitriplet $\overline{3}_f$ and the spin singlet 1_s. The strong attractions between the states $(qq)_{\overline{3}}$ and $(\bar{q}\bar{q})_3$ in S-wave may result in a nonet manifest below 1 GeV, while the conventional ${}^3P_0$ $\bar{q}q$ nonet would have masses of about 1.2–1.6 GeV. Taking the diquarks and antidiquarks as the basic constituents, and keeping the effects of the s quark mass at the first order, the two isoscalars $\bar{u}\bar{d}ud$ and $\bar{s}s\frac{\bar{u}u+\bar{d}d}{\sqrt{2}}$ mix ideally; the $\bar{s}s\frac{\bar{u}u+\bar{d}d}{\sqrt{2}}$ are degenerate with the isovectors $\bar{s}s\bar{d}u$, $\bar{s}s\frac{\bar{u}u-\bar{d}d}{\sqrt{2}}$ and $\bar{s}s\bar{u}d$ naturally. Comparing with the traditional $\bar{q}q$ nonet mesons, the mass spectrum is inverted. The lightest state is the non-strange isosinglet $(\bar{u}\bar{d}ud)$, the heaviest ones are the degenerate isosinglet and isovectors with hidden $\bar{s}s$ pairs, while the four strange states lie in between [5,6].

In this article, we take the point of view that the well confirmed $f_0(980)$ and $a_0(980)$ mesons are four-quark states $(qq)_{\bar{3}}(\bar{q}\bar{q})_3$ in the ideal mixing limit, and devote our attention to the determination of the values of their masses, m_{f_0} and m_{a_0}, in the framework of the QCD sum rule approach [7–9]. Detailed studies of the other scalar four-quark states (the $\kappa(800)$s have not been confirmed yet) and the mixing between the two isoscalars (the $f_0(980)$ meson and the broad $f_0(600)$ meson) will be our next work.

This article is arranged as follows: in Sect. 2, we obtain the QCD sum rules for the masses of the mesons

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$f_0(980)$ and $a_0(980)$; in Sect. 3, we present numerical results; Sect. 4 is reserved for our conclusions.

2 Masses of the $f_0(980)$ and $a_0(980)$ mesons with the QCD sum rules

In the four-quark models, the structures of the scalar mesons $f_0(980)$ and $a_0(980)$ in the ideal mixing limit can be symbolically taken as [2,5,6]

$$f_0(980) = \frac{us\bar{u}\bar{s} + ds\bar{d}\bar{s}}{\sqrt{2}}, \quad a_0(980) = \frac{us\bar{u}\bar{s} - ds\bar{d}\bar{s}}{\sqrt{2}}. \quad (1)$$

The four-quark configurations of the $J^{PC} = 0^{++}$ mesons can give a lot of satisfactory descriptions of the hadron phenomenon, for example, the mass degeneracy of the $f_0(980)$ and $a_0(980)$ mesons, the mass hierarchy pattern of the scalar nonet, the large radiative widths of the $f_0(980)$ and $a_0(980)$ mesons, and the $D_s^+(c\bar{s})$ to $\pi^+\pi^+\pi^-$ decay.

In the following, we write down the interpolating currents for the scalar mesons $f_0(980)$ and $a_0(980)$ based on the four-quark model [8,9]:

$$J^A_{f_0} = \frac{\epsilon_{abc}\epsilon_{ade}}{\sqrt{2}} \times \left[(u_b^{\mathrm{T}} C\gamma_5 s_c)(\bar{u}_d \gamma_5 C\bar{s}_e^{\mathrm{T}}) + (d_b^{\mathrm{T}} C\gamma_5 s_c)(\bar{d}_d \gamma_5 C\bar{s}_e^{\mathrm{T}})\right], \quad (2)$$

$$J^B_{f_0} = \frac{\epsilon_{abc}\epsilon_{ade}}{\sqrt{2}} \times \left[(u_b^{\mathrm{T}} C s_c)(\bar{u}_d C\bar{s}_e^{\mathrm{T}}) + (d_b^{\mathrm{T}} C s_c)(\bar{d}_d C\bar{s}_e^{\mathrm{T}})\right], \quad (3)$$

$$J^A_{a_0} = \frac{\epsilon_{abc}\epsilon_{ade}}{\sqrt{2}} \times \left[(u_b^{\mathrm{T}} C\gamma_5 s_c)(\bar{u}_d \gamma_5 C\bar{s}_e^{\mathrm{T}}) - (d_b^{\mathrm{T}} C\gamma_5 s_c)(\bar{d}_d \gamma_5 C\bar{s}_e^{\mathrm{T}})\right], \quad (4)$$

$$J^B_{a_0} = \frac{\epsilon_{abc}\epsilon_{ade}}{\sqrt{2}} \times \left[(u_b^{\mathrm{T}} C s_c)(\bar{u}_d C\bar{s}_e^{\mathrm{T}}) - (d_b^{\mathrm{T}} C s_c)(\bar{d}_d C\bar{s}_e^{\mathrm{T}})\right], \quad (5)$$

where a, b, c, ... are color indices and C is the charge conjugation matrix. The constituents $S^a(x) = \epsilon^{abc} u_b^{\mathrm{T}}(x) C\gamma_5 s_c(x)$ and $P^a(x) = \epsilon^{abc} u_b^{\mathrm{T}}(x) C s_c(x)$ represent the scalar $J^P = 0^+$ and the pseudoscalar $J^P = 0^-$ us diquarks respectively. They both belong to the antitriplet $\bar{3}$ representation of the color $SU(3)$ group and can cluster together to form S^a–$\bar{S}^a$ type and P^a–$\bar{P}^a$ type diquarks pairs to give the correct spin and parity for the scalar mesons $J^P = 0^+$. The scalar diquarks correspond to the $^1S_0$ states of us and ds diquark systems. The one gluon exchange force and the instanton induced force can lead to significant attractions between the quarks in the 0^+ channels [10]. The pseudoscalar diquarks do not have a non-relativistic limit, and can be taken as the $^3P_0$ states.

The calculation of the $a_0(980)$ meson as a four-quark state in the QCD sum rule approach was done originally for the decay constant and the hadronic coupling constants with the interpolating currents $J^1_{a_0}$ and $J^2_{a_0}$ [11, 12],

$$J^1_{f_0(a_0)} = \Sigma_{\Gamma=1,\pm\gamma_5} \bar{s}\Gamma s \frac{\bar{u}\Gamma u \pm \bar{d}\Gamma d}{\sqrt{2}},$$

$$J^2_{f_0(a_0)} = \Sigma_{\Gamma=1,\pm\gamma_5} \bar{s}\Gamma \frac{\lambda^a}{2} s \frac{\bar{u}\Gamma \frac{\lambda^a}{2} u \pm \bar{d}\Gamma \frac{\lambda^a}{2} d}{\sqrt{2}}, \quad (6)$$

where λ^a is a $SU(3)$ Gell-Mann matrix. Performing a Fierz transformation both in the Dirac spinor and color space, for example, we can obtain

$$\begin{aligned} J^2_{f_0} \propto\ & C_A J^A_{f_0} + C_B J^B_{f_0} \\ & + C_C \frac{\epsilon^{abc}\epsilon^{ade}}{\sqrt{2}} \\ & \times \left[(u_b^{\mathrm{T}} C\gamma_\mu s_c)(\bar{u}_d \gamma^\mu C\bar{s}_e^{\mathrm{T}}) + (d_b^{\mathrm{T}} C\gamma_\mu s_c)(\bar{d}_d \gamma^\mu C\bar{s}_e^{\mathrm{T}})\right] \\ & + C_D \frac{\epsilon^{abc}\epsilon^{ade}}{\sqrt{2}} \\ & \times \left[(u_b^{\mathrm{T}} C\gamma_\mu\gamma_5 s_c)(\bar{u}_d \gamma^\mu\gamma_5 C\bar{s}_e^{\mathrm{T}}) \right. \\ & \qquad \left. + (d_b^{\mathrm{T}} C\gamma_\mu\gamma_5 s_c)(\bar{d}_d \gamma^\mu\gamma_5 C\bar{s}_e^{\mathrm{T}})\right] \cdots, \\ J^2_{a_0} \propto\ & C_A J^A_{a_0} + C_B J^B_{a_0} \\ & + C_C \frac{\epsilon^{abc}\epsilon^{ade}}{\sqrt{2}} \\ & \times \left[(u_b^{\mathrm{T}} C\gamma_\mu s_c)(\bar{u}_d \gamma^\mu C\bar{s}_e^{\mathrm{T}}) - (d_b^{\mathrm{T}} C\gamma_\mu s_c)(\bar{d}_d \gamma^\mu C\bar{s}_e^{\mathrm{T}})\right] \\ & + C_D \frac{\epsilon^{abc}\epsilon^{ade}}{\sqrt{2}} \\ & \times \left[(u_b^{\mathrm{T}} C\gamma_\mu\gamma_5 s_c)(\bar{u}_d \gamma^\mu\gamma_5 C\bar{s}_e^{\mathrm{T}}) \right. \\ & \qquad \left. -(d_b^{\mathrm{T}} C\gamma_\mu\gamma_5 s_c)(\bar{d}_d \gamma^\mu\gamma_5 C\bar{s}_e^{\mathrm{T}})\right] \cdots \end{aligned} \quad (7)$$

Here C_A, C_B, C_C and C_D are coefficients which are not shown explicitly for simplicity. In the color superconductivity theory, the one gluon exchange induced Nambu–Jona–Lasinio-like models will also lead to the S^a–$\bar{S}^a$ type and P^a–$\bar{P}^a$ type diquark pairs [13],

$$G\bar{q}\gamma^\mu \frac{\lambda^a}{2} q \bar{q}\gamma_\mu \frac{\lambda^a}{2} q \propto C_A S^a \bar{S}^a + C_B P^a \bar{P}^a + \cdots \quad (8)$$

So we can take the point of view that the lowest lying scalar mesons are S-wave bound states of diquark–antidiquark pairs of S^a–$\bar{S}^a$ type and P^a–$\bar{P}^a$ type.

In this article, we investigate the masses of the scalar mesons $f_0(980)$ and $a_0(980)$ with two interpolating currents respectively and choose the following two-point correlation functions:

$$\Pi^i_S(p) = \mathrm{i} \int \mathrm{d}^4x\ \mathrm{e}^{\mathrm{i}p.x} \left\langle 0|T[J^i_S(x) J^{i\,\dagger}_S(0)]|0\right\rangle. \quad (9)$$

Here the current J^i_S denotes $J^A_{f_0}$, $J^B_{f_0}$, $J^A_{a_0}$ and $J^B_{a_0}$. According to the basic assumption of current–hadron duality in the QCD sum rule approach [7], we insert a complete series of intermediate states satisfying the unitarity principle with the same quantum numbers as the current operator $J^i_S(x)$ into the correlation functions in (9) to obtain the hadronic representation. Isolating the ground state contributions from the pole terms of the mesons $f_0(980)$ and $a_0(980)$, we get the result

$$\Pi^i_S(p) = \frac{2 f^{i2}_S m^{i8}_s}{m^{i2}_S - p^2} + \cdots, \quad (10)$$

where the following definitions have been used:

$$\langle 0|J_S^i|S\rangle = \sqrt{2} f_S^i m_S^{i4} \,. \quad (11)$$

We have not shown the contributions from the higher resonances and continuum states explicitly for simplicity.

The calculation of the operator product expansion in the deep Euclidean space-time region is straightforward and tedious; technical details are neglected for simplicity. In this article, we consider the vacuum condensates up to dimension six. Once the analytical results are obtained, we can take the current–hadron dualities below the thresholds s_0 and perform the Borel transformation with respect to the variable $P^2 = -p^2$, and finally we obtain the following sum rules:

$$2f^{A2}_{f_0(a_0)} m^{A8}_{f_0(a_0)} \mathrm{e}^{-\frac{m^{A2}_{f_0(a_0)}}{M^2}} = AA, \quad (12)$$

$$2f^{B2}_{f_0(a_0)} m^{B8}_{f_0(a_0)} \mathrm{e}^{-\frac{m^{B2}_{f_0(a_0)}}{M^2}} = BB, \quad (13)$$

$$AA = \int_{4m_s^2}^{s_0} \mathrm{d}s \mathrm{e}^{-\frac{s}{M^2}} \times \left\{ \frac{s^4}{2^9 5! \pi^6} + \frac{\langle\bar{s}s\rangle\langle\bar{q}q\rangle s}{12\pi^2} + \frac{3\langle\bar{q}g_s\sigma G q\rangle - \langle\bar{s}g_s\sigma G s\rangle}{2^6 3\pi^4} m_s s - \frac{2\langle\bar{q}q\rangle - \langle\bar{s}s\rangle}{2^6 3\pi^4} m_s s^2 + \frac{s^2}{2^9 3\pi^4}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle \right\},$$

$$BB = \int_{4m_s^2}^{s_0} \mathrm{d}s \mathrm{e}^{-\frac{s}{M^2}} \times \left\{ -\frac{s^4}{2^9 5! \pi^6} + \frac{\langle\bar{s}s\rangle\langle\bar{q}q\rangle s}{12\pi^2} + \frac{3\langle\bar{q}g_s\sigma G q\rangle + \langle\bar{s}g_s\sigma G s\rangle}{2^6 3\pi^4} m_s s - \frac{2\langle\bar{q}q\rangle + \langle\bar{s}s\rangle}{2^6 3\pi^4} m_s s^2 - \frac{s^2}{2^9 3\pi^4}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle \right\}.$$

Differentiating the above sum rules with respect to the variable $\frac{1}{M^2}$, eliminating the quantities $f^A_{f_0(a_0)}$ and $f^B_{f_0(a_0)}$, we obtain

$$m^{A2}_{f_0(a_0)} = \int_{4m_s^2}^{s_0} \mathrm{d}s \mathrm{e}^{-\frac{s}{M^2}} \times \left\{ \frac{s^5}{2^9 5! \pi^6} + \frac{\langle\bar{s}s\rangle\langle\bar{q}q\rangle s^2}{12\pi^2} + \frac{3\langle\bar{q}g_s\sigma G q\rangle - \langle\bar{s}g_s\sigma G s\rangle}{2^6 3\pi^4} m_s s^2 - \frac{2\langle\bar{q}q\rangle - \langle\bar{s}s\rangle}{2^6 3\pi^4} m_s s^3 + \frac{s^3}{2^9 3\pi^4}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle \right\} / AA, \quad (14)$$

$$m^{B2}_{f_0(a_0)} = \int_{4m_s^2}^{s_0} \mathrm{d}s \mathrm{e}^{-\frac{s}{M^2}} \times \left\{ -\frac{s^5}{2^9 5! \pi^6} + \frac{\langle\bar{s}s\rangle\langle\bar{q}q\rangle s^2}{12\pi^2} + \frac{3\langle\bar{q}g_s\sigma G q\rangle + \langle\bar{s}g_s\sigma G s\rangle}{2^6 3\pi^4} m_s s^2 - \frac{2\langle\bar{q}q\rangle + \langle\bar{s}s\rangle}{2^6 3\pi^4} m_s s^3 - \frac{s^3}{2^9 3\pi^4}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle \right\} / BB. \quad (15)$$

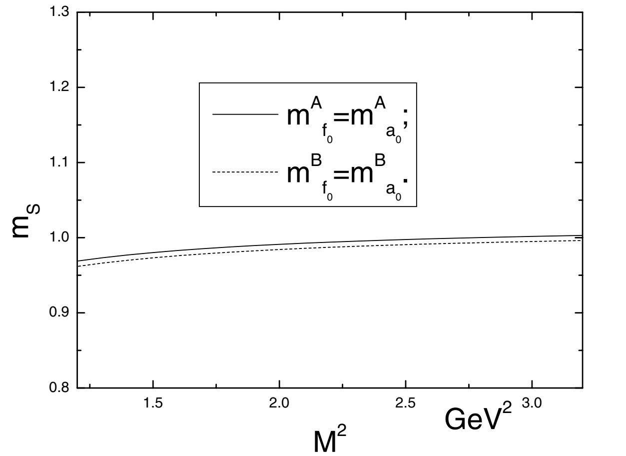

Fig. 1. m^S (in GeV) as a function of the Borel parameter M^2 for $s_0 = 1.50\,\mathrm{GeV}^2$

It is easy to perform the s integral in (12)–(15); we prefer this form for simplicity.

3 Numerical results

The parameters are taken as $\langle\bar{s}s\rangle = 0.8\langle\bar{u}u\rangle$, $\langle\bar{s}g_s\sigma G s\rangle = 0.8\langle\bar{s}s\rangle$, $\langle\bar{q}g_s\sigma G q\rangle = 0.8\langle\bar{q}q\rangle$, $\langle\bar{u}u\rangle = \langle\bar{d}d\rangle = \langle\bar{q}q\rangle = (-219\,\mathrm{MeV})^3$, $\langle\frac{\alpha_s GG}{\pi}\rangle = (0.33\,\mathrm{GeV})^4$, $m_u = m_d = 0$ and $m_s = 150\,\mathrm{MeV}$. The main contributions to the sum rules come from the quark condensates terms; here we have taken the standard values and neglected the uncertainties. Small variations of those condensates will not lead to larger changes about the numerical values. The threshold parameter s_0 is chosen to vary between $(1.4–1.6)\,\mathrm{GeV}^2$ to avoid possible pollutions from higher resonances and continuum states. In the region $M^2 = (1.2–3.2)\,\mathrm{GeV}^2$, the sum rules for $m^A_{f_0} = m^A_{a_0}$ and $m^B_{f_0} = m^B_{a_0}$ are almost independent of the Borel parameter M^2 and are plotted in Fig. 1 for $s_0 = 1.5\,\mathrm{GeV}^2$ as an example.

Due to the special quark constituents and Dirac structures of the interpolating currents, the scalar mesons $f_0(980)$ and $a_0(980)$ have degenerate masses. For the S^a–$\bar{S}^a$ type interpolating currents $J^A_{f_0}$ and $J^A_{a_0}$, the values for the masses are about $m^A_{f_0} = m^A_{a_0} = (0.96–1.02)\,\mathrm{GeV}$, while for the P^a–$\bar{P}^a$ type interpolating currents $J^B_{f_0}$ and $J^B_{a_0}$, the values for the masses are about $m^B_{f_0} = m^B_{a_0} = (0.95–1.01)\,\mathrm{GeV}$. Although the values for the masses $m^A_{f_0} = m^A_{a_0}$ lie a little above the masses $m^B_{f_0} = m^B_{a_0}$, we cannot get to the conclusion that the scalar mesons $f_0(980)$ and $a_0(980)$ prefer the S^a–$\bar{S}^a$ type interpolating currents $J^A_{f_0}$ and $J^A_{a_0}$ to the P^a–$\bar{P}^a$ type interpolating currents $J^B_{f_0}$ and $J^B_{a_0}$. A precise determination of what type interpolating currents we should choose calls for original theoretical approaches; the contributions from the direct instantons may do. In our recent work, we observed that the contributions from the direct instantons are considerable for the pentaquark state $\Theta^+(1540)$ [14]; furthermore, the contributions from the direct instantons can improve the QCD sum rule greatly in some channels, for example, the nonperturbative contributions from the direct instantons to the conventional operator product expansion can significantly improve the stability of chirally odd nucleon sum rules [15,16]. Despite whatever the interpolating currents may be, we observe that they both give the correct degenerate masses for the scalar mesons $f_0(980)$ and $a_0(980)$, and there must be some four-quark constituents in those mesons.

4 Conclusions

In this article, we take the point of view that the $f_0(980)$ and $a_0(980)$ mesons are the four-quark states $(qq)_{\bar{3}}(\bar{q}\bar{q})_3$ in the ideal mixing limit, and devote our attention to the determination of the values of their masses m_{f_0} and m_{a_0} in the framework of the QCD sum rule approach. Due to the special quark constituents and Dirac structures of the interpolating currents, the scalar mesons $f_0(980)$ and $a_0(980)$ have degenerate masses. For the S^a–$\bar{S}^a$ type interpolating currents $J^A_{f_0}$ and $J^A_{a_0}$, the values for the masses are about $m^A_{f_0} = m^A_{a_0} = (0.96\text{–}1.02)\,\text{GeV}$, while for the P^a–$\bar{P}^a$ type interpolating currents $J^B_{f_0}$ and $J^B_{a_0}$, the values for the masses are about $m^B_{f_0} = m^B_{a_0} = (0.95\text{–}1.01)\,\text{GeV}$. Although the values for the masses $m^A_{f_0} = m^A_{a_0}$ lie a little above the masses $m^B_{f_0} = m^B_{a_0}$, we cannot get to the conclusion that the scalar mesons $f_0(980)$ and $a_0(980)$ prefer the S^a–$\bar{S}^a$ type interpolating currents $J^A_{f_0}$ and $J^A_{a_0}$ to the P^a–$\bar{P}^a$ type interpolating currents $J^B_{f_0}$ and $J^B_{a_0}$. Despite whatever the interpolating currents may be, we observe that they both give the correct degenerate masses for the scalar mesons $f_0(980)$ and $a_0(980)$, and there must be some four-quark constituents in those mesons. Our results support the four-quark model and the hybrid model. In the hybrid model, those mesons are four-quark states, $(qq)_{\bar{3}}(\bar{q}\bar{q})_3$ in S-wave near the center, with some constituent $q\bar{q}$ in P-wave, but further out they rearrange themselves into $(q\bar{q})_1(q\bar{q})_1$ states and finally as meson–meson states [5]. A precise determination of what type of interpolating currents we should choose calls for original theoretical approaches. The contributions from the direct instantons may do.

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