QCD sum rules study of $QQ - u\bar{d}$ mesons

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We use QCD sum rules to study the possible existence of $QQ - u\bar{d}$ mesons, assumed to be a state with $J^P = 1^+$. For definiteness, we work with a current with an axial heavy diquark and a scalar light antidiquark, at leading order in $\alpha_s$. We consider the contributions of condensates up to dimension eight. For the $b$-quark, we predict $M_{T_{bb}} = (10.2 \pm 0.3) \text{ GeV}$, which is below the $BB^*$ threshold. For the $c$-quark, we predict $M_{T_{cc}} = (4.0 \pm 0.2) \text{ GeV}$, in agreement with quark model predictions.

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The general idea of possible stable heavy tetraquarks has been first suggested by Jaffe [1]. The case of a tetraquark $QQ\bar{u}\bar{d}$ with quantum numbers $I = 0$, $J = 1$ and $P = +1$ which, following ref. [2], we call $T_{QQ}$, is especially interesting. As already noted previously [2, 3], the $T_{bb}$ and $T_{cc}$ states cannot decay strongly or electromagnetically into two $B$ or two $D$ mesons in the $S$ wave due to angular momentum conservation nor in $P$ wave due to parity conservation. If their masses are below the $BB^*$ and $DD^*$ thresholds, these decays are also forbidden. Moreover, in the large $m_Q$ limit, the light degrees of freedom cannot resolve the closely bound $QQ$ system. This results in bound states similar to the $\Lambda_Q$ states, with $QQ$ playing the role of the heavy antiquark [4]. Therefore, the stability of $\Lambda_Q$ implies that $QQ\bar{u}\bar{d}$ is also safe from decaying through $QQ\bar{u}\bar{d} \to QQq + \bar{q}u\bar{d}$. As a result, $T_{QQ}$ is stable with respect to strong interactions and must decay weakly.

There are some predictions for the masses of the $T_{QQ}$ states. In ref. [3] the authors use a color-magnetic interaction, with flavor symmetry breaking corrections, to study heavy tetraquarks. They assume that the Belle resonance, $X(3872)$, is a $c\bar{c}q\bar{q}$ tetraquark, and use its mass as input to determine the mass of other tetraquark states. They get $M_{T_{cc}} = 3966$ MeV and $M_{T_{bb}} = 10372$ MeV. In ref. [2], the authors use one-gluon exchange potentials and two different spatial configurations to study the mesons $T_{cc}$ and $T_{bb}$. They get $M_{T_{cc}} = 3876 - 3905$ MeV and $M_{T_{bb}} = 10519 - 10651$ MeV. There are also calculations using expansion in the harmonic oscillator basis [6], and variational method [7].

In this work we use QCD sum rules (QCDSR) [8, 9, 10] to study the two-point functions of the state $T_{QQ}$. There are several reasons, why it is interesting to investigate this channel. First of all, having two heavy quarks, it is an explicit exotic state. The experimental observation would already prove the existence of the tetraquark state without any theoretical extrapolation. Moreover, from a technical point of view, this means that there are no contributions from the disconnected diagrams, which are technically very difficult to estimate in QCD sum rules or in lattice gauge theory calculation.

In previous calculations, the QCDSR approach was used to study the light scalar mesons [11, 12, 13, 14, 15] the $D^*_{J}(2317)$ meson [16, 17] and the $X(3872)$ meson [18], considered as four-quark states and a good agreement with the experimental masses was obtained. However, the tests were not decisive as the usual quark–antiquark assignments also provide predictions consistent with data [10, 12, 19, 20].

Considering $T_{QQ}$ as an axial diquark-antidiquark state, a possible current describing such state is given by:

$$j_\mu = i\{Q_a^T C\gamma_\mu Q_b\}[\bar{u}_a\gamma_5 C\bar{d}_b^T] ,$$

where $a$, $b$ are color indices, $C$ is the charge conjugation matrix and $Q$ denotes the heavy quark.

In general, one should consider all possible combinations of different $1^+$ four-quark operators, as was done in [21] for the $0^{++}$ light mesons. However, the current in Eq. (1) well represents the most attractive

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configuration expected with two heavy quarks. This is so because the most attractive light antidiquark is expected to be the in the color triplet, flavor anti-symmetric and spin 0 channel. This is also expected quite naturally from the color magnetic interaction, which can be phenomenologically parameterized as,

\[ V_{ij} = -\frac{C}{m_im_j} \lambda_i \cdot \lambda_j \sigma_i \cdot \sigma_j. \]  

(2)

Here, \( m, \lambda, \sigma \) are the mass, color and spin of the constituent quark \( i, j \). Eq. (2) favors the anti-diquark to be in the color triplet and spin 0 channel. The flavor anti-symmetric condition then follows from requiring anti-symmetric wave function of the anti-diquark. Similarly, since the anti-diquark is in the color triplet state the remaining QQ should be in the color anti-triplet spin 1 state. Although the spin 1 configuration is repulsive, its strength is much smaller than that for the light diquark due to the heavy charm quark mass. Therefore a constituent quark picture for \( T_{QQ} \) would be a light anti-diquark in color triplet, flavor anti-symmetric and spin 0 \( (\epsilon_{abc}[u_b\gamma_5C\bar{d}_c^T]) \) combined with a heavy diquark of spin 1 \( (\epsilon_{aef}[Q^T_2C\gamma_\mu Q_f]) \).

The simplest choice for the current to have a non zero overlap with such a \( T_{QQ} \) configuration is given in Eq. (1). While a similar configuration \( T_{ss} \) is also possible, we believe that the repulsion in the strange diquark with spin 1 will be larger and hence energetically less favorable. As discussed above, since the quantum number is \( 1^+ \), the decay into \( DD \) or \( BB \) would be forbidden and the allowed decay into \( DD^* \) or \( BB^* \) would have a smaller phase space, and the tetraquark state might have a small width, or may even be bound.

The QCDSR is constructed from the two-point correlation function

\[ \Pi_{\mu\nu}(q) = i \int d^4x \, e^{iq\cdot x} \langle 0 | T [j_\mu(x)J_\nu^\dagger(0)] | 0 \rangle = -\Pi_1(q^2)(g_{\mu\nu} - \frac{q_\mu q_\nu}{q^2}) + \Pi_0(q^2) \frac{q_\mu q_\nu}{q^2}. \]  

(3)

Since the axial vector current is not conserved, the two functions, \( \Pi_1 \) and \( \Pi_0 \), appearing in Eq. (3) are independent and have respectively the quantum numbers of the spin 1 and 0 mesons.

The calculation of the phenomenological side proceeds by inserting intermediate states for the meson \( T_{QQ} \). Parametrizing the coupling of the axial vector meson \( 1^+ \), to the current, \( j_\mu \), in Eq. (1) in terms of the meson decay constant \( f_T \) and the meson mass \( M_T \) as:

\[ \langle 0 | j_\mu | T_{QQ} \rangle = \sqrt{2} f_T M_T^2 \epsilon_\mu, \]  

(4)

the phenomenological side of Eq. (3) can be written as

\[ \Pi_{\mu\nu}^{phen}(q^2) = \frac{2f_T^2M_T^8}{M_T^4 - q^2} \left( -g_{\mu\nu} + \frac{q_\mu q_\nu}{M_T^2} \right) + \cdots, \]  

(5)

where the Lorentz structure \( g_{\mu\nu} \) gets contributions only from the \( 1^+ \) state. The dots denote higher axial-vector resonance contributions that will be parametrized, as usual, through the introduction of a continuum threshold parameter \( s_0 \). [20]

On the OPE side, we work at leading order in \( \alpha_s \) and consider the contributions of condensates up to dimension eight. To keep the charm quark mass finite, we use the momentum-space expression for the charm quark propagator. We follow ref. [22] and calculate the light quark part of the correlation function in the coordinate-space, which is then Fourier transformed to the momentum space in \( D \) dimensions. The resulting light-quark part is combined with the charm quark part before it is dimensionally regularized at \( D = 4 \).

The correlation function, \( \Pi_1 \), in the OPE side can be written as a dispersion relation:

\[ \Pi_1^{OPF}(q^2) = \int_{4m_{Q}^2}^{\infty} ds \frac{\rho(s)}{s - q^2}, \]  

(6)

where the spectral density is given by the imaginary part of the correlation function: \( \pi\rho(s) = \text{Im}[\Pi_1^{OPE}(s)] \). After making a Borel transform of both sides, and transferring the continuum contribution to the OPE side, the sum rule for the axial vector meson \( T_{QQ} \) up to dimension-eight condensates can be written as:

\[ 2f_T^2M_T^8e^{-M_T^2/M^2} = \int_{4m_{Q}^2}^{s_0} ds \, e^{-s/M^2} \rho(s) + \Pi_{1}^{\text{mix}(\bar{q}q)}(M^2), \]  

(7)
where
\[\rho(s) = \rho^{\text{pert}}(s) + \rho^{(\bar{q}q)}(s) + \rho^{(G^2)}(s) + \rho^{\text{mix}}(s) + \rho^{(\bar{q}q)^2}(s) + \rho^{\text{mix}\{\bar{q}q\}}(s), \]

with
\[\rho^{\text{pert}}(s) = \frac{1}{2^{9/2} \pi^5} \int_{\alpha_{\text{min}}}^{\alpha_{\text{max}}} \frac{d\alpha}{\alpha^3} \int_{\beta_{\text{min}}}^{1-\alpha} \frac{d\beta}{\beta^3} (1 - \alpha - \beta) \left[ (\alpha + \beta) m_{Q}^2 - \alpha \beta s \right]^3 \times \left[ 1 + \frac{\alpha + \beta}{4} \left( (\alpha + \beta) m_{Q}^2 - \alpha \beta s - m_{Q}^2 (1 - \alpha - \beta) \right), \right.\]
\[\rho^{(\bar{q}q)}(s) = 0,\]
\[\rho^{(G^2)}(s) = -\frac{(g^2 G^2)}{2^{10} \pi^6} \left\{ -\frac{1}{4} \int_{\alpha_{\text{min}}}^{\alpha_{\text{max}}} \frac{d\alpha}{\alpha(1 - \alpha)} (m_{Q}^2 - \alpha (1 - \alpha) s)^2 \right.\]
\[+ \int_{\alpha_{\text{min}}}^{\alpha_{\text{max}}} \frac{d\alpha}{\alpha} \int_{\beta_{\text{min}}}^{1-\alpha} \frac{d\beta}{\beta} \left[ \frac{(\alpha + \beta) m_{Q}^2 - \alpha \beta s}{4 \beta} \left( (\alpha + \beta) m_{Q}^2 - \alpha \beta s + 2 m_{Q}^2 \right) \right.\]
\[+ m_{Q}^2 (1 - \alpha - \beta) \left[ m_{Q}^2 (1 - \alpha - \beta) + \left( (\alpha + \beta) m_{Q}^2 - \alpha \beta s \right) \left( -4 - \alpha - \beta + 3 \frac{\beta}{1 - \alpha} \right) \right.\]
\[+ \frac{1}{48 \alpha \beta^2} (1 - \alpha - \beta) \left( (\alpha + \beta) m_{Q}^2 - \alpha \beta s \right)^2 (5 - \alpha - \beta) \right\} \left. \right\} + \int_{\alpha_{\text{min}}}^{\alpha_{\text{max}}} \frac{d\alpha}{\alpha(1 - \alpha)} (m_{Q}^2 - \alpha (1 - \alpha) s)^2 \right.\]
\[\rho^{\text{mix}}(s) = 0,\]
\[\rho^{\text{mix}\{\bar{q}q\}}(s) = \frac{(\bar{q}q)^2}{24 \pi^2} s \sqrt{1 - 4 m_{Q}^2/s}. \]

where the integration limits are given by \(\alpha_{\text{min}} = (1 - \sqrt{1 - 4 m_{Q}^2/s})/2\), \(\alpha_{\text{max}} = (1 + \sqrt{1 - 4 m_{Q}^2/s})/2\) and \(\beta_{\text{min}} = \alpha m_{Q}^2/(s a - m_{Q}^2)\). The contribution of dimension-six condensates \(\langle g^3 G^3 \rangle\) is neglected, since it is assumed to be suppressed by the loop factor \(1/16\pi^2\). We have included, for completeness, a part of the dimension-8 condensate contributions. We should note that a complete evaluation of these contributions require more involved analysis including a non-trivial choice of the factorization assumption basis [28].

\[\rho^{\text{mix}\{\bar{q}q\}}(s) = -\frac{(\bar{q}q \sigma G q) \langle \bar{q}q \rangle}{2^{9/2} \pi^6} \sqrt{1 - 4 m_{Q}^2/s},\]
\[\Pi^{\text{mix}\{\bar{q}q\}}_1(M^2) = -\frac{m_{Q}^2 (\bar{q}q \sigma G q) \langle \bar{q}q \rangle}{2^{5/2} 3 \pi^2} \int_0^{1} d\alpha \left[ 4 - \frac{m_{Q}^2}{\alpha(1 - \alpha) M^2} \right] \exp \left[ -\frac{m_{Q}^2}{\alpha(1 - \alpha) M^2} \right]. \]

In order to extract the mass \(M_T\) without worrying about the value of the decay constant \(f_T\), we take the derivative of Eq. (7) with respect to \(1/M^2\), divide the result by Eq. (7) and obtain:
\[M_T^2 = \frac{\int_{4 m_{Q}^2}^{s_0} ds \ e^{-s/M^2} \ s \ \rho(s)}{\int_{4 m_{Q}^2}^{s_0} ds \ e^{-s/M^2} \ \rho(s)}. \]

This quantity has the advantage to be less sensitive to the perturbative radiative corrections than the individual moments. Therefore, we expect that our results obtained to leading order in \(\alpha_s\) will be quite accurate.

In the numerical analysis of the sum rules, the values used for the quark masses and condensates are (see e.g. [10, 29]): \(m_u(m_u) = 1.23 \pm 0.05\) GeV, \(m_d(m_u) = 4.24 \pm 0.06\) GeV, \(\langle \bar{q}q \rangle = -(0.23 \pm 0.03)^3\) GeV\(^3\), \(\langle \bar{q}q \sigma G q \rangle = m_{Q}^2 \langle \bar{q}q \rangle\) with \(m_{Q}^2 = 0.8\) GeV\(^2\), \(g^2 G^2 = 0.88\) GeV\(^4\).

We start with the double charmed meson \(T_{cc}\). We evaluate the sum rules in the range \(2.0 \leq M^2 \leq 4\) GeV\(^2\) for \(s_0\) in the range: \(4.6 \leq \sqrt{s_0} \leq 5.0\) GeV.

Comparing the relative contribution of each term in Eqs. (9) to (11), to the right hand side of Eq. (7) we obtain a quite good OPE convergence (the perturbative contribution is at least 50% of the total).
FIG. 1: The relative OPE convergence in the region $2.0 \leq M^2 \leq 4.0$ GeV$^2$ for $\sqrt{s_0} = 4.8$ GeV. We start with the perturbative contribution divided by the total (long-dashed line) and each subsequent line represents the addition of one extra condensate dimension in the expansion: $+(g_5^2 G^2)$ (dot-dashed line), $+(\bar{q}q)^2$ (dotted-line), $+m_0^2(\bar{q}q)^2$ (solid line).

for $M^2 > 2.5$ GeV$^2$, as can be seen in Fig. 1. This analysis allows us to determine the lower limit constraint for $M^2$ in the sum rules window. This figure also shows that, although there is a change of sign between dimension-six and dimension-eight condensates contributions, the contribution of the latter is very small, where, we have assumed, in Fig. 1 to Fig. 4, the validity of the vacuum saturation for these condensates. The relatively small contribution of the dimension-eight condensates may justify the validity of our approximation, unlike in the case of the 5-quark current correlator, as noticed in [30].

FIG. 2: The solid line shows the relative pole contribution (the pole contribution divided by the total, pole plus continuum, contribution) and the dashed line shows the relative continuum contribution for $\sqrt{s_0} = 4.8$ GeV.

We get an upper limit constraint for $M^2$ by imposing the rigorous constraint that the QCD continuum contribution should be smaller than the pole contribution. The maximum value of $M^2$ for which this constraint is satisfied depends on the value of $s_0$. The comparison between pole and continuum contributions for $\sqrt{s_0} = 4.8$ GeV is shown in Fig. 2. The same analysis for the other values of the continuum threshold gives $M^2 \leq 3.1$ GeV$^2$ for $\sqrt{s_0} = 4.6$ GeV and $M^2 \leq 3.6$ GeV$^2$ for $\sqrt{s_0} = 5.0$ GeV.

In Fig. 3 we show the $T_{cc}$ meson mass obtained from Eq. (11), in the relevant sum rules window, with
the upper and lower validity limits indicated. From Fig. 3 we see that the results are reasonably stable as a function of $M^2$. In our numerical analysis, we shall then consider the range of $M^2$ values from 2.5 GeV$^2$ until the one allowed by the sum rule window criteria as can be deduced from Fig. 3 for each value of $s_0$.

We found that our results are not very sensitive to the value of the charm quark mass, neither to the value of the condensates. The most important source of uncertainty is the value of the continuum threshold and the Borel interval. Using the QCD parameters given above, the QCDSR predictions for the $T_{cc}$ mesons mass is:

$$M_{T_{cc}} = (4.0 \pm 0.2) \text{ GeV} ,$$

in a very good agreement with the predictions in refs. [2] and [5].

One can also evaluate the decay constant, defined in Eq. (4), to leading order in $\alpha_s$:

$$f_{T_{cc}} = (5.95 \pm 0.65) \times 10^{-5} \text{ GeV} ,$$

which can be more affected by radiative corrections than $M_{T_{cc}}$.

In the case of the double-beauty meson $T_{bb}$, using consistently the perturbative $\overline{MS}$-mass $m_b(m_b) = (4.24\pm0.6) \text{ GeV}$, and the continuum threshold in the range $11.3 \leq \sqrt{s_0} \leq 11.7 \text{ GeV}$, we find a good OPE convergence for $M^2 > 7.5 \text{ GeV}^2$. We also find that the pole contribution is bigger than the continuum contribution for $M^2 < 9.6 \text{ GeV}^2$ for $\sqrt{s_0} < 11.3 \text{ GeV}$, and for $M^2 < 11.2 \text{ GeV}^2$ for $\sqrt{s_0} < 11.7 \text{ GeV}$.

In Fig. 4 we show the $T_{bb}$ meson mass obtained from Eq. (11), in the relevant sum rules window, with the upper and lower validity limits indicated. From Fig. 4 we see that the results are very stable as a function of $M^2$ in the allowed region. Taking into account the variation of $M^2$ and varying $s_0$ and $m_b$ in the regions indicated above, we arrive at the prediction:

$$M_{T_{bb}} = (10.2 \pm 0.3) \text{ GeV} ,$$

also in a very good agreement with the results in refs. [2], [5] and [7]. For completeness, we predict the corresponding value of the decay constant to leading order in $\alpha_s$:

$$f_{T_{bb}} = (10.4 \pm 2.8) \times 10^{-6} \text{ GeV} .$$

We have presented a QCDSR analysis of the two-point functions of the double heavy-quark axial meson, $T_{QQ}$, considered as a four quark state. We find that the sum rules results for the masses of $T_{cc}$ and $T_{bb}$
FIG. 4: The $T_{bb}$ meson mass as a function of the sum rule parameter ($M^2$) for different values of the continuum threshold: $\sqrt{s_0} = 11.3$ GeV (dashed line), $\sqrt{s_0} = 11.7$ GeV (solid line). The bars delimit the region allowed for the sum rules.

are compatible with the results in refs. [2] and [3]. An improvement of this result needs an accurate determination of running masses $m_c$ and $m_b$ of the $\overline{MS}$-scheme and the inclusion of radiative corrections.

Our results show that while the $T_{cc}$ mass is bigger than the $D^*D$ threshold at about 3.875 GeV, the $T_{bb}$ mass is appreciably below the $\bar{B}^*\bar{B}$ threshold at about 10.6 GeV. Therefore, our results indicate that the $T_{bb}$ meson should be stable with respect to strong interactions and must decay weakly. Our result also confirms the naive expectation that the exotic states with heavy quarks tend to be more stable than the corresponding light states [31].

We present in Eqs. (13), and (15) predictions for the decay constants of the $T_{cc}$ and $T_{bb}$.

Different choices of the four-quark operators have been systematically presented for the $0^{++}$ light mesons in [21]. Though some combinations can provide a faster convergence of the OPE, we do not expect that the choice of the operators will affect much our results, where, in our analysis, the OPE has a good convergence.

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