Doubly hidden-charm/bottom $QQ\bar{Q}\bar{Q}$ tetraquark states

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Abstract. We study the mass spectra for the $cc\bar{c}\bar{c}$ and $bb\bar{b}\bar{b}$ tetraquark states by developing a moment sum rule method. Our results show that the $bb\bar{b}\bar{b}$ tetraquarks lie below the threshold of $\eta_b(1S)\eta_b(1S)$. They are probably stable and very narrow. The masses for the doubly hidden-charm states $cc\bar{c}\bar{c}$ are higher than the spontaneous dissociation thresholds of two charmonium mesons. We suggest to search for such states in the $J/\psi J/\psi$ and $\eta_c(1S)\eta_c(1S)$ channels.

1 Introduction

The configurations of multiquark states were proposed by Gell-Mann [1] and Zweig [2] at the birth of quark model (QM). In the past fifty years, it has been an extremely intriguing research issue of searching for multiquark matter. The light tetraquark $qq\bar{q}\bar{q}$ state has been used to investigate the scalar mesons below 1 GeV [3]. Since 2003, plenty of charmoniumlike states have been observed and the hidden-charm $qc\bar{q}\bar{c}$ tetraquark fomalism is extensively discussed to explain the nature of these new XYZ states [4–11].

The doubly hidden-charm/bottom tetraquark $QQ\bar{Q}\bar{Q}$ is composed of four heavy quarks. Such tetraquark states did not receive much attention in both experimental and theoretical aspects [12–21]. Recently, there are some discussions about the masses and decays of the $QQ\bar{Q}\bar{Q}$ states [22–30]. The masses of these $QQ\bar{Q}\bar{Q}$ tetraquarks are far away from the mass regions of the conventional $Q\bar{Q}$ mesons and the XYZ states. It will be very easy to distinguish them from the XYZ and $Q\bar{Q}$ states in the spectroscopy. On the other hand, the $QQ\bar{Q}\bar{Q}$ states favor the compact tetraquark configuration than the loosely bound hadron molecular configuration, since the light mesons can not be exchanged

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between the two charmonium/bottomonium states. In this paper, we develop a moment QCD sum rule method to calculate the mass spectra for the doubly hidden-charm/bottom $cc\bar{c}\bar{c}$ and $bb\bar{b}\bar{b}$ tetraquark states.

2 QCD sum rules

In this section we briefly introduce the method of QCD sum rules [31–33]. Comparing to the traditional SVZ QCD sum rules, we use another version of QCD sum rules, the moment QCD sum rules in our analyses for the doubly hidden-charm/bottom $QQ\bar{Q}\bar{Q}$ tetraquark systems. The moment QCD sum rules have been very successfully used for studying the charmonium and bottomonium mass spectra [31, 32, 34–36] and determining the heavy quark masses and the strong coupling constant [37–39].

We start by considering the following two-point correlation functions

$$\begin{aligned}
\Pi(q) &= i\int d^4x e^{iq\cdot x}\langle 0|T[J(x)J^\dagger(0)]|0\rangle\,,\\
\Pi_{\mu\nu}(q) &= i\int d^4x e^{iq\cdot x}\langle 0|T[J_\mu(x)J_\nu^\dagger(0)]|0\rangle\,,\\
\Pi_{\mu\nu,\rho\sigma}(q) &= i\int d^4x e^{iq\cdot x}\langle 0|T[J_{\mu\nu}(x)J_{\rho\sigma}^\dagger(0)]|0\rangle\,,
\end{aligned}\tag{1}$$

in which the interpolating currents $J(x)$, $J_\mu(x)$ and $J_{\mu\nu}(x)$ couple to the scalar, vector and tensor states respectively.

To study the doubly hidden-charm/bottom tetraquarks, we construct the $QQ\bar{Q}\bar{Q}$ interpolating currents with four heavy quarks in the compact diquark-antidiquark configuration. We use all diquark fields $Q_a^TCQ_b$, $Q_a^TC\gamma_5Q_b$, $Q_a^TC\gamma_\mu\gamma_5Q_b$, $Q_a^TC\gamma_\mu Q_b$, $Q_a^TC\sigma_{\mu\nu}Q_b$ and $Q_a^TC\sigma_{\mu\nu}\gamma_5Q_b$ and consider the Pauli principle to determine the color and flavor structures for the tetraquark operators. Following Refs. [22, 40], we obtain the $QQ\bar{Q}\bar{Q}$ tetraquark interpolating currents as the following. The interpolating currents with $J^{PC}=0^{++}$ are

$$\begin{aligned}
J_1 &= Q_a^TC\gamma_5Q_b\bar{Q}_a\gamma_5C\bar{Q}_b^T\,,\\
J_2 &= Q_a^TC\gamma_\mu\gamma_5Q_b\bar{Q}_a\gamma^\mu\gamma_5C\bar{Q}_b^T\,,\\
J_3 &= Q_a^TC\sigma_{\mu\nu}Q_b\bar{Q}_a\sigma^{\mu\nu}C\bar{Q}_b^T\,,\\
J_4 &= Q_a^TC\gamma_\mu Q_b\bar{Q}_a\gamma^\mu C\bar{Q}_b^T\,,\\
J_5 &= Q_a^TCQ_b\bar{Q}_aC\bar{Q}_b^T\,,
\end{aligned}\tag{2}$$

where J_1, J_2, J_5 belong to the symmetric $[\mathbf{6_c}]_{QQ}\otimes[\bar{\mathbf{6}}_\mathbf{c}]_{\bar{Q}\bar{Q}}$ color structure while J_3, J_4 belong to the antisymmetric $[\bar{\mathbf{3}}_\mathbf{c}]_{QQ}\otimes[\mathbf{3_c}]_{\bar{Q}\bar{Q}}$ color structure. The interpolating currents with $J^{PC}=0^{-+}$ and 0^{--} are

$$\begin{aligned}
J_1^\pm &= Q_a^TCQ_b\bar{Q}_a\gamma_5C\bar{Q}_b^T \pm Q_a^TC\gamma_5Q_b\bar{Q}_aC\bar{Q}_b^T\,,\\
J_2^+ &= Q_a^TC\sigma_{\mu\nu}Q_b\bar{Q}_a\sigma^{\mu\nu}\gamma_5C\bar{Q}_b^T\,,
\end{aligned}\tag{3}$$

in which J_1^+ and J_2^+ couple to the states with $J^{PC}=0^{-+}$, and J_1^- couples to the states with $J^{PC}=0^{--}$. The currents $J_1^\pm$ belong to the symmetric color structure while J_2^+ belongs to antisymmetric color structure. The interpolating currents with $J^{PC}=1^{++}$ and 1^{+-} are

$$\begin{aligned}
J_{1\mu}^\pm &= Q_a^TC\gamma_\mu\gamma_5Q_b\bar{Q}_aC\bar{Q}_b^T \pm Q_a^TCQ_b\bar{Q}_a\gamma_\mu\gamma_5C\bar{Q}_b^T\,,\\
J_{2\mu}^\pm &= Q_a^TC\sigma_{\mu\nu}\gamma_5Q_b\bar{Q}_a\gamma^\nu C\bar{Q}_b^T \pm Q_a^TC\gamma^\nu Q_b\bar{Q}_a\sigma_{\mu\nu}\gamma_5C\bar{Q}_b^T\,,
\end{aligned}\tag{4}$$

in which $J^+_{1\mu}$ and $J^+_{2\mu}$ couple to the states with $J^{PC} = 1^{++}$, and $J^-_{1\mu}$ and $J^-_{2\mu}$ couple to the states with $J^{PC} = 1^{+-}$. The currents $J^\pm_{1\mu}$ belong to the symmetric color structure while $J^\pm_{2\mu}$ belongs to antisymmetric color structure. The interpolating currents with $J^{PC} = 1^{-+}$ and 1^{--} are

$$
\begin{aligned}
J^\pm_{1\mu} &= Q_a^T C\gamma_\mu\gamma_5 Q_b \bar{Q}_a\gamma_5 C\bar{Q}_b^T \pm Q_a^T C\gamma_5 Q_b \bar{Q}_a\gamma_\mu\gamma_5 C\bar{Q}_b^T \,, \\
J^\pm_{2\mu} &= Q_a^T C\sigma_{\mu\nu} Q_b \bar{Q}_a\gamma^\nu C\bar{Q}_b^T \pm Q_a^T C\gamma^\nu Q_b \bar{Q}_a\sigma_{\mu\nu} C\bar{Q}_b^T \,,
\end{aligned}
\tag{5}
$$

in which $J^+_{1\mu}$ and $J^+_{2\mu}$ couple to the states with $J^{PC} = 1^{-+}$, and $J^-_{1\mu}$ and $J^-_{2\mu}$ couple to the states with $J^{PC} = 1^{--}$. The currents $J^\pm_{1\mu}$ belong to the symmetric color structure while $J^\pm_{2\mu}$ belongs to antisymmetric color structure. The interpolating currents with $J^{PC} = 2^{++}$ are

$$
\begin{aligned}
J_{1\mu\nu} &= Q_a^T C\gamma_\mu Q_b \bar{Q}_a\gamma_\nu C\bar{Q}_b^T + Q_a^T C\gamma_\nu Q_b \bar{Q}_a\gamma_\mu C\bar{Q}_b^T \,, \\
J_{2\mu\nu} &= Q_a^T C\gamma_\mu\gamma_5 Q_b \bar{Q}_a\gamma_\nu\gamma_5 C\bar{Q}_b^T + Q_a^T C\gamma_\nu\gamma_5 Q_b \bar{Q}_a\gamma_\mu\gamma_5 C\bar{Q}_b^T \,,
\end{aligned}
\tag{6}
$$

where current $J_{1\mu\nu}$ belongs to the antisymmetric color structure while $J_{2\mu\nu}$ belongs to symmetric color structure.

At the hadronic level, the correlation functions in Eq.(1) can be described by the dispersion relation

$$
\Pi(q^2) = \frac{(q^2)^N}{\pi}\int_{M_H^2}^\infty \frac{\mathrm{Im}\Pi(s)}{s^N(s-q^2-i\epsilon)}ds + \sum_{n=0}^{N-1} b_n (q^2)^n \,,
\tag{7}
$$

where M_H is the hadron mass and b_n are unknown subtraction constants. A narrow resonance approximation is usually used to describe the spectral function

$$
\begin{aligned}
\rho(s) = \frac{1}{\pi}\mathrm{Im}\Pi(s) &= \sum_n \delta(s-m_n^2)\langle 0|J|n\rangle\langle n|J^\dagger|0\rangle + \cdots \\
&= f_X^2\delta(s-m_X^2) + \cdots \,,
\end{aligned}
\tag{8}
$$

where "$\cdots$" represents the excited higher states and continuum contributions and f_X is a coupling constant between the interpolating current and hadron state

$$
\begin{aligned}
\langle 0|J|X\rangle &= f_X \,, \\
\langle 0|J_\mu|X\rangle &= f_X\epsilon_\mu \,, \\
\langle 0|J_{\mu\nu}|X\rangle &= f_X\epsilon_{\mu\nu} \,,
\end{aligned}
\tag{9}
$$

in which ϵ_μ and $\epsilon_{\mu\nu}$ are the polarization vector and tensor, respectively. To pick out the contribution of the lowest lying resonance in Eq. (8), we define moments in Euclidean region $Q^2 = -q^2 > 0$ [32, 41]:

$$
M_n(Q_0^2) = \frac{1}{n!}\left(-\frac{d}{dQ^2}\right)^n \Pi(Q^2)|_{Q^2=Q_0^2} = \int_{16m_Q^2}^\infty \frac{\rho(s)}{(s+Q_0^2)^{n+1}}ds
\tag{10}
$$

$$
= \frac{f_X^2}{(m_X^2+Q_0^2)^{n+1}}\left[1+\delta_n(Q_0^2)\right],
\tag{11}
$$

in which $\delta_n(Q_0^2)$ contains the contributions of higher states and continuum. It tends to zero as n goes to infinity. We consider the following ratio to eliminate f_X in Eq. (11)

$$
r(n, Q_0^2) \equiv \frac{M_n(Q_0^2)}{M_{n+1}(Q_0^2)} = (m_X^2+Q_0^2)\frac{1+\delta_n(Q_0^2)}{1+\delta_{n+1}(Q_0^2)}.
\tag{12}
$$

One expects $\delta_n(Q_0^2) \cong \delta_{n+1}(Q_0^2)$ for sufficiently large n to suppress the contributions of higher states and continuum [32]. Then hadron mass of the lowest lying resonance m_X is then extracted as

$$m_X = \sqrt{r(n, Q_0^2) - Q_0^2}\,. \tag{13}$$

Using the operator production expansion (OPE) method, the two-point function can also be evaluated at the quark-gluonic level as a function of various QCD parameters. In the fully heavy tetraquark systems, we only need to calculate the perturbative term and the gluon condensate contributions to the correlation functions. One can find the results of the moments $M_n(Q_0^2)$ in Ref. [22].

3 Numerical results

We perform the numerical analyses by using the following values of parameters [42–45]

$$\begin{aligned} m_c(\overline{\mathrm{MS}}) &= 1.27 \pm 0.03 \text{ GeV}, \\ m_b(\overline{\mathrm{MS}}) &= 4.18 \pm 0.03 \text{ GeV}\,, \\ \langle g_s^2 GG\rangle &= (0.48 \pm 0.14) \text{ GeV}^4\,. \end{aligned} \tag{14}$$

To provide reliable moment sum rule analyses, one needs to find suitable working regions of the two parameters n and Q_0^2 in the ratio $r(n, Q_0^2)$. We define $\xi = Q_0^2/16m_c^2$ for $cc\bar{c}\bar{c}$ and Q_0^2/m_b^2 for $bb\bar{b}\bar{b}$ systems for convenience. These two parameters will affect the pole contribution and the OPE convergence. For small value of ξ, the high dimension condensates in OPE will give large contributions, and thus leading to bad OPE convergence [32, 36]. However, a larger value of ξ means slower convergence of $\delta_n(Q_0^2)$ in Eq. (11). Such behavior can be compensated by n: the OPE convergence becomes good for small n while $\delta_n(Q_0^2)$ tends to zero for sufficiently large n. One needs to find suitable working regions for (n, ξ) where the lowest lying resonance dominates the moments and the OPE converges well.

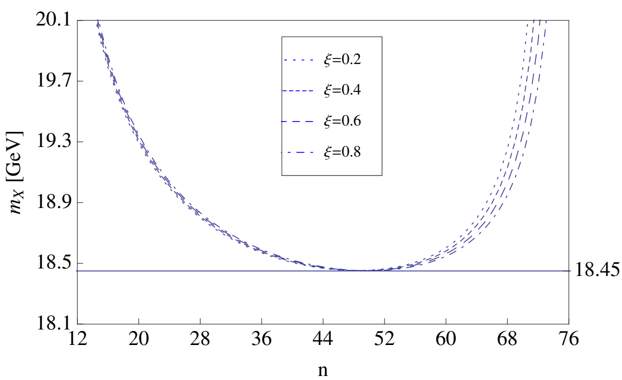

Figure 1. Hadron mass m_{X_b} for $J_1(bb\bar{b}\bar{b})$ with $J^{PC} = 0^{++}$, as a function of n for different value of ξ.

As an example, we use the interpolating current J_1 with $J^{PC} = 0^{++}$ in Eq. (2) to perform numerical analyses. Requiring the perturbative term to be larger than the gluon condensate term, we obtain upper

J^{PC}	Currents	m_{X_c}(GeV)	m_{X_b}(GeV)
0^{++}	J_1	6.44 ± 0.15	18.45 ± 0.15
	J_2	6.59 ± 0.17	18.59 ± 0.17
	J_3	6.47 ± 0.16	18.49 ± 0.16
	J_4	6.46 ± 0.16	18.46 ± 0.14
	J_5	6.82 ± 0.18	19.64 ± 0.14
0^{-+}	J_1^+	6.84 ± 0.18	18.77 ± 0.18
	J_2^+	6.85 ± 0.18	18.79 ± 0.18
0^{--}	J_1^-	6.84 ± 0.18	18.77 ± 0.18
1^{++}	$J_{1\mu}^+$	6.40 ± 0.19	18.33 ± 0.17
	$J_{2\mu}^+$	6.34 ± 0.19	18.32 ± 0.18
1^{+-}	$J_{1\mu}^-$	6.37 ± 0.18	18.32 ± 0.17
	$J_{2\mu}^+$	6.51 ± 0.15	18.54 ± 0.15
1^{-+}	$J_{1\mu}^+$	6.84 ± 0.18	18.80 ± 0.18
	$J_{2\mu}^+$	6.88 ± 0.18	18.83 ± 0.18
1^{--}	$J_{1\mu}^-$	6.84 ± 0.18	18.77 ± 0.18
	$J_{2\mu}^-$	6.83 ± 0.18	18.77 ± 0.16
2^{++}	$J_{1\mu\nu}$	6.51 ± 0.15	18.53 ± 0.15
	$J_{2\mu\nu}$	6.37 ± 0.19	18.32 ± 0.17

Table 1. Mass spectra for the $cc\bar{c}\bar{c}$ and $bb\bar{b}\bar{b}$ tetraquarks.

limits n_{max}, which increases with respect to the value of ξ. We show the hadron mass m_{X_b} as a function of n for $\xi = 0.2, 0.4, 0.6, 0.8$ in Fig. 1. One notes that the mass curves have plateaus which provide stable mass prediction

$$m_{X_b} = (18.45 \pm 0.15)\,\mathrm{GeV}\,, \tag{15}$$

in which the error comes from the uncertainties of ξ, the heavy quark mass and the gluon condensate in Eq. (14). Using the interpolating currents in Eqs. (2)–(6), we perform numerical analyses for all $cc\bar{c}\bar{c}$ and $bb\bar{b}\bar{b}$ systems with various quantum numbers. We collect the numerical results in Table 1. It is shown that the negative parity states ($J^{PC} = 0^{-+}, 0^{--}, 1^{-+}, 1^{--}$) are slightly heavier than the positive parity states ($J^{PC} = 0^{++}, 1^{++}, 1^{+-}, 2^{++}$).

It is interesting to compare the mass spectra with the corresponding two-meson mass thresholds. As shown in Fig. 2, the masses of $bb\bar{b}\bar{b}$ tetraquarks are below the $\eta_b(1S)\eta_b(1S)$ threshold while all $cc\bar{c}\bar{c}$ tetraquarks lie above the $\eta_c(1S)\eta_c(1S)$ threshold. The two bottomonium mesons decays for the $bb\bar{b}\bar{b}$ tetraquarks are thus forbidden by the kinematics. For the doubly hidden-charm $cc\bar{c}\bar{c}$ tetraquarks, they can decay via the spontaneous dissociation mechanism by considering the restrictions of the symmetries. In Table 2, we collect the possible S-wave and P-wave dissociation decay channels for the $cc\bar{c}\bar{c}$ states.

In principle, the $bb\bar{b}\bar{b}$ tetraquark can also decay into $B^{(*)}\bar{B}^{(*)}$ via a heavy quark pair annihilation and a light quark pair creation processes. The suppression by the annihilation of a heavy quark pair will be compensated by the large phase space factor. Such $B^{(*)}\bar{B}^{(*)}$ decay modes may dominate the total width of the doubly hidden-bottom $bb\bar{b}\bar{b}$ tetraquark state.

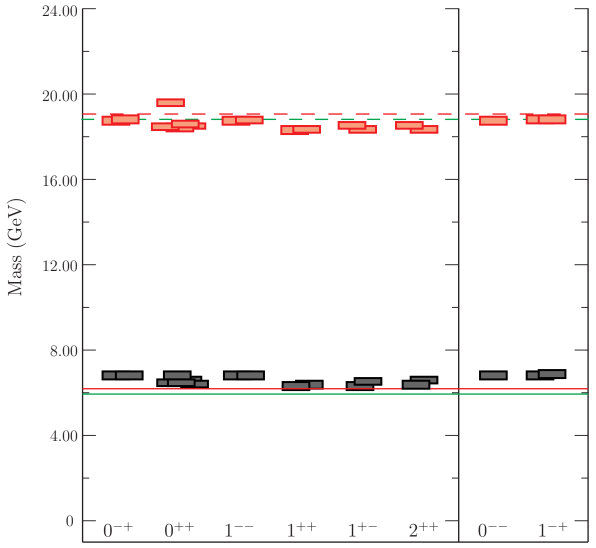

Figure 2. Summary of the doubly hidden-charm/bottom tetraquark spectra labelled by J^{PC}. The green and red solid (dashed) lines indicate the $\eta_c(1S)\eta_c(1S)$ ($\eta_b(1S)\eta_b(1S)$) and $J/\psi J/\psi$ ($\Upsilon(1S)\Upsilon(1S)$) thresholds, respectively.

J^{PC}	S-wave	P-wave
0^{++}	$\eta_c(1S)\eta_c(1S)$, $J/\psi J/\psi$	$\eta_c(1S)\chi_{c1}(1P)$, $J/\psi h_c(1P)$
0^{-+}	$\eta_c(1S)\chi_{c0}(1P)$, $J/\psi h_c(1P)$	$J/\psi J/\psi$
0^{--}	$J/\psi\chi_{c1}(1P)$	$J/\psi\eta_c(1S)$
1^{++}	–	$J/\psi h_c(1P)$, $\eta_c(1S)\chi_{c1}(1P)$, $\eta_c(1S)\chi_{c0}(1P)$
1^{+-}	$J/\psi\eta_c(1S)$	$J/\psi\chi_{c0}(1P)$, $J/\psi\chi_{c1}(1P)$, $\eta_c(1S)h_c(1P)$
1^{-+}	$J/\psi h_c(1P)$, $\eta_c(1S)\chi_{c1}(1P)$	–
1^{--}	$J/\psi\chi_{c0}(1P)$, $J/\psi\chi_{c1}(1P)$, $\eta_c(1S)h_c(1P)$	$J/\psi\eta_c(1S)$

Table 2. Possible decay modes of the $cc\bar{c}\bar{c}$ states by spontaneous dissociation into two charmonium mesons.

4 Summary

In this paper, we have calculated the mass spectra for the doubly hidden-charm/bottom $cc\bar{c}\bar{c}$ and $bb\bar{b}\bar{b}$ tetraquark states by using the moment QCD sum rule method. Our results show that the $cc\bar{c}\bar{c}$ tetraquarks lie above the two charmonium spontaneous dissociation thresholds and thus can mainly decay into two charmonium mesons. We suggest to search for these doubly hidden-charm $cc\bar{c}\bar{c}$ states in the $J/\psi J/\psi$ and $\eta_c(1S)\eta_c(1S)$ channels. For the $bb\bar{b}\bar{b}$ tetraquarks, their masses are lower than the $\eta_b(1S)\eta_b(1S)$ threshold so that the two bottomonium mesons decays are kinematical forbidden. These $bb\bar{b}\bar{b}$ tetraquark, if exist, may be very narrow and stable. In the near future, these doubly hidden-charm/bottom $cc\bar{c}\bar{c}$ and $bb\bar{b}\bar{b}$ tetraquark states can be searched for at facilities such as LHCb, CMS, RHIC and the forthcoming BelleII.

Acknowledgments

This project is supported by the Chinese National Youth Thousand Talents Program; the Natural Sciences and Engineering Research Council of Canada (NSERC); the National Natural Science Foundation of China under Grants No. 11475015, 373 No. 11375024, No. 11222547, No. 11175073, 374 No. 11575008, and No. 11621131001; the 973 program; 375 the Ministry of Education of China (SRFDP under Grant 376 No. 20120211110002 and the Fundamental Research 377 Funds for the Central Universities); and the National 378 Program for Support of Top-Notch Youth Professionals.

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