

Research Article

Search for the Charmed Baryonium and Dibaryon Structures via the QCD Sum Rules

Xiu-Wu Wang ^{1,2} and Zhi-Gang Wang ¹

¹School of Mathematics and Physics, North China Electric Power University, Baoding 071003, China

²School of Nuclear Science and Engineering, North China Electric Power University, Beijing 102206, China

Correspondence should be addressed to Xiu-Wu Wang; wangxiuwu2020@163.com

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In the present work, we construct eight six-quark currents to study the $\Sigma_c \bar{\Sigma}_c$ baryonium and $\Sigma_c \Sigma_c$ dibaryon states via the QCD sum rules. For either $\Sigma_c \bar{\Sigma}_c$ baryonium or $\Sigma_c \Sigma_c$ dibaryon states, we construct four currents with the $J^P = 1^-, 1^+, 0^+$ and 0^- . Except the current 1^+ of $\Sigma_c \bar{\Sigma}_c$, we find the Borel windows for the other seven. There are two possible $\Sigma_c \bar{\Sigma}_c$ molecular states or compact six-quark states with the $J^P = 1^-$ and 0^- . The results of the current 0^+ of $\Sigma_c \Sigma_c$ show a possible $\Sigma_c \Sigma_c$ molecular state with 20 MeV binding energy. We find two $\Sigma_c \bar{\Sigma}_c$ baryonium resonances states with the $J^P = 1^+, 0^+$ and two $\Sigma_c \Sigma_c$ dibaryon resonances states with the $J^P = 0^-, 1^-$, respectively. Pole residues of the seven states are also calculated.

1. Introduction

The dibaryon is composed of two baryons, and the baryonium is formed by a baryon and antibaryon. In 1980s, the research of the interaction between the two baryons caused great attention both theoretically and experimentally [1]. The first known dibaryon state discovered in 1932 has been the deuteron, the bound state of proton and neutron [2]. Except for the deuteron, another genuine dibaryon resonance observed in experiment is the $d^*(2380)$ [3–5]. The baryon consists of three quarks, six quarks for the dibaryon or baryonium, it is natural to extend the concept of hexaquarks for which the quarks sitting in a common quark bag. The data of the electromagnetic form factors of the state are possible to differentiate whether it is a molecular type or a compact six-quark entity. For the deuteron, its charge radius is 2.1 fm, the proton and the neutron inside the deuteron are on average 4 fm apart from each other and do not overlap [6, 7], these data depict a clear physical picture of the hadronic molecule for the deuteron. One could also make

the differentiation via the decay branching; in Ref. [8], the decay width of $d^*(2380)$ is considered via the chiral quark model; it is proposed that $d^*(2380)$ spends 2/3 of its time as a six-quark and the rest as a $\Delta\Delta$ molecule. Theoretically and experimentally, the ways which may find the difference are there; however, seldom work has been done yet for both the $\Sigma_c \Sigma_c$ dibaryon and $\Sigma_c \bar{\Sigma}_c$ baryonium in this aspect. At present, it is too hard for us to differentiate $\Sigma_c \Sigma_c$ dibaryon or $\Sigma_c \bar{\Sigma}_c$ baryonium from the compact six-quark state. We now could only say: it may be the molecular state or compact six-quark state if the mass is below the threshold of the two baryon constituents (bound state); as for the dibaryon resonance state, its mass is above the threshold of the two baryon constituents.

The oldest prediction of the existence of dibaryon with nearly 100 MeV binding energy based on SU(6) symmetry was done by Dyson and Xuong [9] which is in good agreement with the experimental observations of $d^*(2380)$ [3–5]. Although the insufficient experimental data for the solid observations of the dibaryon and baryonium states, lots

of theoretical calculations have been done in search of them. In recent years, several QCD-inspired models are applied to study the dibaryons and baryoniums, such as the quark delocalization color screening model [10, 11], solving the Bethe-Salpeter equation via effective Lagrangians [12, 13], constituent quark model [14, 15], lattice QCD [16, 17], and the QCD sum rules [18–21]. Especially, in Refs. [22, 23], the authors study the charmed hadrons to investigate the mutual interactions between them and search for poles by solving the Bethe-Salpeter equation. It is natural to ask: what are the predictions of the interactions of these charmed hadrons in the framework of the QCD sum rules? In Ref. [24], the scalar and axialvector $\Xi_{cc}\Sigma_c$ dibaryon states are studied in consideration of two-baryon scattering states. It is proposed that the two-baryon scattering states cannot saturate the QCD sum rules and their contributions can be neglected. We follow this idea and do not consider the two-baryon scattering states. The earlier study of the hidden-charm tetraquark molecular states has previously been published in a preprint [25]. In this article, we construct eight six-quark local interpolating currents to study both the $\Sigma_c\bar{\Sigma}_c$ baryoniums with the $J^P = 1^-, 1^+, 0^+, 0^-$ and the $\Sigma_c\Sigma_c$ dibaryons with the $J^P = 0^+, 0^-, 1^+, 1^-$, respectively.

The article is arranged as follows: we derive the QCD sum rules of the dibaryon and baryonium states in Section 2. We present the numerical results and discussions in Section 3. Section 4 is reserved for our conclusions.

2. The QCD Sum Rules for the Baryonium and Dibaryon States

Now, let us write down the corresponding two-point correlation functions from the eight constructed currents:

$$\begin{aligned}
\Pi_{1;\alpha\beta}(p) &= i \int d^4x e^{ip \cdot x} \langle 0 | J_{1,\alpha}(x) J_{1,\beta}^\dagger(0) | 0 \rangle, \\
\Pi_{2;\alpha\beta}(p) &= i \int d^4x e^{ip \cdot x} \langle 0 | J_{2,\alpha}(x) J_{2,\beta}^\dagger(0) | 0 \rangle, \\
\Pi_3(p) &= i \int d^4x e^{ip \cdot x} \langle 0 | J_3(x) J_3^\dagger(0) | 0 \rangle, \\
\Pi_4(p) &= i \int d^4x e^{ip \cdot x} \langle 0 | J_4(x) J_4^\dagger(0) | 0 \rangle, \\
\Pi_5(p) &= i \int d^4x e^{ip \cdot x} \langle 0 | J_5(x) J_5^\dagger(0) | 0 \rangle, \\
\Pi_6(p) &= i \int d^4x e^{ip \cdot x} \langle 0 | J_6(x) J_6^\dagger(0) | 0 \rangle, \\
\Pi_{7;\alpha\beta}(p) &= i \int d^4x e^{ip \cdot x} \langle 0 | J_{7,\alpha}(x) J_{7,\beta}^\dagger(0) | 0 \rangle, \\
\Pi_{8;\alpha\beta}(p) &= i \int d^4x e^{ip \cdot x} \langle 0 | J_{8,\alpha}(x) J_{8,\beta}^\dagger(0) | 0 \rangle,
\end{aligned} \tag{1}$$

where

$$\begin{aligned}
J_{1,\alpha} &= \bar{J}_\Sigma \gamma_\alpha J_\Sigma, \\
J_{2,\alpha} &= \bar{J}_\Sigma \gamma_5 \gamma_\alpha J_\Sigma, \\
J_3 &= \bar{J}_\Sigma J_\Sigma, \\
J_4 &= \bar{J}_\Sigma \gamma_5 J_\Sigma, \\
J_5 &= J_\Sigma^T C \gamma_5 J_\Sigma, \\
J_6 &= J_\Sigma^T C J_\Sigma, \\
J_{7,\alpha} &= J_\Sigma^T C \gamma_\alpha J_\Sigma, \\
J_{8,\alpha} &= J_\Sigma^T C \gamma_5 \gamma_\alpha J_\Sigma, \\
J_\Sigma &= \epsilon^{ijk} \left(u^{iT} C \gamma_\mu d^j \right) \gamma^\mu \gamma_5 c^k,
\end{aligned} \tag{2}$$

the superscripts i, j, k are color indices and C represents the charge conjugation matrix. The parities of these currents are shown as

$$\begin{aligned}
\hat{P} J_{1,\alpha}(x) \hat{P}^{-1} &= J_{1,\alpha}^\alpha(\tilde{x}), \\
\hat{P} J_{2,\alpha}(x) \hat{P}^{-1} &= -J_{2,\alpha}^\alpha(\tilde{x}), \\
\hat{P} J_3(x) \hat{P}^{-1} &= J_3(\tilde{x}), \\
\hat{P} J_4(x) \hat{P}^{-1} &= -J_4(\tilde{x}), \\
\hat{P} J_5(x) \hat{P}^{-1} &= J_5(\tilde{x}), \\
\hat{P} J_6(x) \hat{P}^{-1} &= -J_6(\tilde{x}), \\
\hat{P} J_{7,\alpha}(x) \hat{P}^{-1} &= -J_{7,\alpha}^\alpha(\tilde{x}), \\
\hat{P} J_{8,\alpha}(x) \hat{P}^{-1} &= J_{8,\alpha}^\alpha(\tilde{x}),
\end{aligned} \tag{3}$$

where \hat{P} is the parity operator, $x = (x^0, x^1, x^2, x^3)$ and $\tilde{x} = (x^0, -x^1, -x^2, -x^3)$. A complete set of intermediate hadronic states are inserted into each $J_i(x) J_i^\dagger(0)$ in the above correlation functions. The correlation functions at the hadron sides are simplified as

$$\begin{aligned}
\Pi_{1,2,7,8;\alpha\beta}(p) &= A_{1,2,7,8}(p) \left(-g_{\alpha\beta} + \frac{p_\alpha p_\beta}{p^2} \right) + \dots, \\
\Pi_{3,4,5,6}(p) &= \frac{\lambda_{3,4,5,6}^2}{M_{3,4,5,6}^2 - p^2} + \dots,
\end{aligned} \tag{4}$$

where

$$\begin{aligned}
A_{1,2,7,8}(p) &= \frac{\lambda_{1,2,7,8}^2}{M_{1,2,7,8}^2 - p^2} + \dots, \\
\langle 0 | J_{1,2,7,8;\alpha}(0) | Z_{1,2,7,8} \rangle &= \lambda_{1,2,7,8} \epsilon_\alpha, \\
\langle 0 | J_{3,4,5,6}(0) | Z_{3,4,5,6} \rangle &= \lambda_{3,4,5,6},
\end{aligned} \tag{5}$$

$|Z_i\rangle$ denote the ground states which have the same J^P as J_i , the ε_α are the polarization vectors and λ_i are the pole residues which show the coupling of the currents J_i to the states ($i = 1, 2, \dots, 8$).

For the QCD sides of the correlation functions, we contract the quark fields via the Wick theorem and express the correlation functions in terms of full quark propagators. After the detailed calculations, we find $\Pi_{7;\alpha\beta}(p) \equiv 0$ which means the net contribution of all the Feynman diagrams of this state including the leading order, perturbation terms and all kinds of condensates is zero. Except for this term, the others are given by

$$\Pi_{1;\alpha\beta}(p) = -\varepsilon^{ijk} \varepsilon^{lmn} \varepsilon^{i'j'k'} \varepsilon^{l'm'n'} i \int d^4x e^{ip \cdot x}$$

$$\left\{ \text{Tr} \left[\gamma_5 \gamma^\mu \gamma_\alpha \gamma^\tau \gamma_5 B^{nk'}(x) \gamma_5 \gamma^{\mu'} \gamma_\beta \gamma^{\tau'} \gamma_5 B^{n'k}(-x) \right] \cdot \text{Tr} \left[\gamma_{\tau'} Q^{m'j}(-x) \gamma_\mu C Q^{i'T}(-x) C \right] \cdot \text{Tr} \left[\gamma_\tau Q^{mj}(x) \gamma_{\mu'} C Q^{l'T}(x) C \right] \right\},$$

$$\Pi_{2;\alpha\beta}(p) = -\varepsilon^{ijk} \varepsilon^{lmn} \varepsilon^{i'j'k'} \varepsilon^{l'm'n'} i \int d^4x e^{ip \cdot x}$$

$$\left\{ \text{Tr} \left[\gamma_5 \gamma^\mu \gamma_\alpha \gamma^\tau \gamma_5 B^{nk'}(x) \gamma_5 \gamma^{\mu'} \gamma_\beta \gamma^{\tau'} \gamma_5 B^{n'k}(-x) \right] \cdot \text{Tr} \left[\gamma_{\tau'} Q^{m'j}(-x) \gamma_\mu C Q^{i'T}(-x) C \right] \cdot \text{Tr} \left[\gamma_\tau Q^{mj}(x) \gamma_{\mu'} C Q^{l'T}(x) C \right] \right\},$$

$$\Pi_3(p) = -\varepsilon^{ijk} \varepsilon^{lmn} \varepsilon^{i'j'k'} \varepsilon^{l'm'n'} i \int d^4x e^{ip \cdot x}$$

$$\left\{ \text{Tr} \left[\gamma_5 \gamma^\mu \gamma^\tau \gamma_5 B^{nk'}(x) \gamma_5 \gamma^{\mu'} \gamma^{\tau'} \gamma_5 B^{n'k}(-x) \right] \cdot \text{Tr} \left[\gamma_{\tau'} Q^{m'j}(-x) \gamma_\mu C Q^{i'T}(-x) C \right] \cdot \text{Tr} \left[\gamma_\tau Q^{mj}(x) \gamma_{\mu'} C Q^{l'T}(x) C \right] \right\},$$

$$\Pi_4(p) = -\varepsilon^{ijk} \varepsilon^{lmn} \varepsilon^{i'j'k'} \varepsilon^{l'm'n'} i \int d^4x e^{ip \cdot x}$$

$$\left\{ \text{Tr} \left[\gamma_5 \gamma^\mu \gamma_5 \gamma^\tau \gamma_5 B^{nk'}(x) \gamma_5 \gamma^{\mu'} \gamma_5 \gamma^{\tau'} \gamma_5 B^{n'k}(-x) \right] \cdot \text{Tr} \left[\gamma_{\tau'} Q^{m'j}(-x) \gamma_\mu C Q^{i'T}(-x) C \right] \cdot \text{Tr} \left[\gamma_\tau Q^{mj}(x) \gamma_{\mu'} C Q^{l'T}(x) C \right] \right\},$$

$$\Pi_5(p) = -2\varepsilon^{ijk} \varepsilon^{lmn} \varepsilon^{i'j'k'} \varepsilon^{l'm'n'} i \int d^4x e^{ip \cdot x}$$

$$\left\{ \text{Tr} \left[\gamma_5 \gamma^\tau \gamma_5 \gamma^\mu \gamma_5 B^{kk'}(x) \gamma_5 \gamma^{\mu'} \gamma_5 \gamma^{\tau'} \gamma_5 C B^{nn'T}(x) C \right] \cdot \text{Tr} \left[\gamma_\mu Q^{ij}(x) \gamma_{\mu'} C Q^{i'T}(x) C \right] \text{Tr} \left[\gamma_\tau Q^{mmm'}(x) \gamma_{\tau'} C Q^{ll'T}(x) C \right] \right. \\ \left. + \text{Tr} \left[\gamma_5 \gamma^\tau \gamma_5 \gamma^\mu \gamma_5 B^{kn'}(x) \gamma_5 \gamma^{\mu'} \gamma_5 \gamma^{\tau'} \gamma_5 C B^{nk'T}(x) C \right] \cdot \text{Tr} \left[\gamma_\mu Q^{ij}(x) \gamma_{\tau'} C Q^{i'T}(x) C \right] \text{Tr} \left[\gamma_\tau Q^{mmm'}(x) \gamma_{\mu'} C Q^{ll'T}(x) C \right] \right. \\ \left. - 2\text{Tr} \left[\gamma_5 \gamma^\tau \gamma_5 \gamma^\mu \gamma_5 B^{kk'}(x) \gamma_5 \gamma^{\mu'} \gamma_5 \gamma^{\tau'} \gamma_5 C B^{nn'T}(x) C \right] \cdot \text{Tr} \left[\gamma_\mu Q^{ij}(x) \gamma_{\mu'} C Q^{i'T}(x) C \right] \gamma_\tau Q^{mmm'}(x) \gamma_{\tau'} C Q^{ll'T}(x) C \right\},$$

$$\Pi_6(p) = 2\varepsilon^{ijk} \varepsilon^{lmn} \varepsilon^{i'j'k'} \varepsilon^{l'm'n'} i \int d^4x e^{ip \cdot x}$$

$$\left\{ \text{Tr} \left[\gamma_5 \gamma^\tau \gamma^\mu \gamma_5 B^{kk'}(x) \gamma_5 \gamma^{\mu'} \gamma^{\tau'} \gamma_5 C B^{nn'T}(x) C \right] \cdot \text{Tr} \left[\gamma_\mu Q^{ij}(x) \gamma_{\mu'} C Q^{i'T}(x) C \right] \text{Tr} \left[\gamma_\tau Q^{mmm'}(x) \gamma_{\tau'} C Q^{ll'T}(x) C \right] \right. \\ \left. + \text{Tr} \left[\gamma_5 \gamma^\tau \gamma^\mu \gamma_5 B^{kn'}(x) \gamma_5 \gamma^{\mu'} \gamma^{\tau'} \gamma_5 C B^{nk'T}(x) C \right] \cdot \text{Tr} \left[\gamma_\mu Q^{ij}(x) \gamma_{\tau'} C Q^{i'T}(x) C \right] \text{Tr} \left[\gamma_\tau Q^{mmm'}(x) \gamma_{\mu'} C Q^{ll'T}(x) C \right] \right. \\ \left. - 2\text{Tr} \left[\gamma_5 \gamma^\tau \gamma^\mu \gamma_5 B^{kk'}(x) \gamma_5 \gamma^{\mu'} \gamma^{\tau'} \gamma_5 C B^{nn'T}(x) C \right] \cdot \text{Tr} \left[\gamma_\mu Q^{ij}(x) \gamma_{\mu'} C Q^{i'T}(x) C \right] \gamma_\tau Q^{mmm'}(x) \gamma_{\tau'} C Q^{ll'T}(x) C \right\},$$

$$\Pi_{8;\alpha\beta}(p) = 2\varepsilon^{ijk} \varepsilon^{lmn} \varepsilon^{i'j'k'} \varepsilon^{l'm'n'} i \int d^4x e^{ip \cdot x}$$

$$\left\{ \text{Tr} \left[\gamma_5 \gamma^\tau \gamma_5 \gamma^\mu \gamma_\alpha \gamma^\tau \gamma_5 B^{kk'}(x) \gamma_5 \gamma^{\mu'} \gamma_5 \gamma_\beta \gamma^{\tau'} \gamma_5 C B^{nn'T}(x) C \right] \cdot \text{Tr} \left[\gamma_\mu Q^{ij}(x) \gamma_{\mu'} C Q^{i'T}(x) C \right] \text{Tr} \left[\gamma_\tau Q^{mmm'}(x) \gamma_{\tau'} C Q^{ll'T}(x) C \right] \right. \\ \left. + \text{Tr} \left[\gamma_5 \gamma^\tau \gamma_5 \gamma^\mu \gamma_\alpha \gamma^\tau \gamma_5 B^{kn'}(x) \gamma_5 \gamma^{\mu'} \gamma_5 \gamma_\beta \gamma^{\tau'} \gamma_5 C B^{nk'T}(x) C \right] \cdot \text{Tr} \left[\gamma_\mu Q^{ij}(x) \gamma_{\tau'} C Q^{i'T}(x) C \right] \text{Tr} \left[\gamma_\tau Q^{mmm'}(x) \gamma_{\mu'} C Q^{ll'T}(x) C \right] \right. \\ \left. - 2\text{Tr} \left[\gamma_5 \gamma^\tau \gamma_5 \gamma^\mu \gamma_\alpha \gamma^\tau \gamma_5 B^{kk'}(x) \gamma_5 \gamma^{\mu'} \gamma_5 \gamma_\beta \gamma^{\tau'} \gamma_5 C B^{nn'T}(x) C \right] \cdot \text{Tr} \left[\gamma_\mu Q^{ij}(x) \gamma_{\mu'} C Q^{i'T}(x) C \right] \gamma_\tau Q^{mmm'}(x) \gamma_{\tau'} C Q^{ll'T}(x) C \right\}, \quad (6)$$

where $Q^{ab}(x)$ and $B^{ab}(x)$ are the full light and heavy quark propagators, respectively, under the operator product expansion at the quark level, they are expressed as

$$Q^{ab}(x) = \frac{ix\delta^{ab}}{2\pi^2 x^4} - \frac{\delta^{ab}}{12} \langle \bar{q}q \rangle - \frac{\delta^{ab} x^2}{192} \langle \bar{q}g_s \sigma Gq \rangle - \frac{i\delta^{ab} x^2 x l g_s^2 \langle \bar{q}q \rangle^2}{7776} \\ - (t^n)^{ab} (x/\sigma^{\alpha\beta} + \sigma^{\alpha\beta} x/l) \frac{i}{32\pi^2 x^2} g_s G_{\alpha\beta}^n - \frac{\delta^{ab} x^4 \langle \bar{q}q \rangle \langle GG \rangle}{27648} \\ - \frac{1}{8} \langle \bar{q}^b \sigma^{\alpha\beta} q^a \rangle \sigma_{\alpha\beta} - \frac{1}{4} \langle \bar{q}^b \gamma_\mu q^a \rangle \gamma^\mu + \dots, \\ B_{ab}(x) = \frac{i}{(2\pi)^4} \int d^4k e^{-ik \cdot x} \left\{ \frac{\delta_{ab}}{k-l-m_c} - \frac{g_s C_{\alpha\beta}^h t_{ab}^h}{4} \frac{\sigma^{\alpha\beta} (k+l+m_c) + (k+l+m_c) \sigma^{\alpha\beta}}{(k^2 - m_c^2)^2} \right. \\ \left. + \frac{g_s D_\alpha C_{\beta\lambda}^h t_{ab}^h (f^{\lambda\beta\alpha} + f^{\lambda\alpha\beta})}{3(k^2 - m_c^2)^4} + \dots \right\}, \\ f^{\lambda\alpha\beta} = (k+l+m_c) \gamma^\lambda (k+l+m_c) \gamma^\alpha (k+l+m_c) \gamma^\beta (k+l+m_c), \quad (7)$$

TABLE 1: The related numerical results extracted from the Borel windows of the seven currents.

	J_i	J^P	T^2 (GeV ²)	$\sqrt{s_0}$ (GeV)	μ (GeV)	M (GeV)	PC	λ (10 ⁻³ GeV ⁸)
$\Sigma_c \bar{\Sigma}_c$	J_1	1^-	3.6 ~ 4.0	5.60 ± 0.1	3.2	4.88 ^{+0.09} _{-0.08}	(54 ~ 41)%	3.23 ^{+0.46} _{-0.43}
$\Sigma_c \bar{\Sigma}_c$	J_2	1^+	3.6 ~ 4.0	5.90 ± 0.1	3.8	5.31 ^{+0.07} _{-0.07}	(54 ~ 42)%	4.74 ^{+0.72} _{-0.64}
$\Sigma_c \bar{\Sigma}_c$	J_3	0^+	3.4 ~ 3.8	5.80 ± 0.1	3.7	5.23 ^{+0.07} _{-0.07}	(55 ~ 43)%	4.19 ^{+0.67} _{-0.59}
$\Sigma_c \bar{\Sigma}_c$	J_4	0^-	3.6 ~ 4.0	5.60 ± 0.1	3.2	4.88 ^{+0.08} _{-0.08}	(53 ~ 40)%	3.34 ^{+0.48} _{-0.45}
$\Sigma_c \Sigma_c$	J_5	0^+	3.3 ~ 3.7	5.55 ± 0.1	3.2	4.89 ^{+0.11} _{-0.10}	(56 ~ 42)%	4.09 ^{+0.61} _{-0.56}
$\Sigma_c \Sigma_c$	J_6	0^-	3.6 ~ 4.0	5.90 ± 0.1	3.9	5.37 ^{+0.08} _{-0.08}	(52 ~ 40)%	6.65 ^{+1.03} _{-0.93}
$\Sigma_c \Sigma_c$	J_8	1^-	3.3 ~ 3.7	5.85 ± 0.1	3.7	5.23 ^{+0.06} _{-0.06}	(57 ~ 43)%	4.66 ^{+0.71} _{-0.63}
$\Lambda_c \Lambda_c$	J_1	0^+	3.4 ~ 3.8	5.70 ± 0.1	3.6	5.12 ^{+0.15} _{-0.12}	(54 ~ 42)%	1.81 ^{+0.31} _{-0.27}
$\Lambda_c \bar{\Lambda}_c$	J_2	0^-	3.2 ~ 3.6	5.30 ± 0.1	2.8	4.66 ^{+0.10} _{-0.09}	(54 ~ 41)%	0.64 ^{+0.10} _{-0.09}
$\Lambda_c \bar{\Lambda}_c$	J_3	1^+	3.4 ~ 3.8	5.60 ± 0.1	3.4	5.00 ^{+0.10} _{-0.09}	(55 ~ 42)%	0.64 ^{+0.10} _{-0.09}
$\Lambda_c \bar{\Lambda}_c$	J_4	1^-	3.2 ~ 3.6	5.25 ± 0.1	2.7	4.57 ^{+0.09} _{-0.09}	(53 ~ 39)%	0.52 ^{+0.09} _{-0.08}

where $t^n = \lambda^n/2$, λ^n are the Gell-Mann matrices ($n = 1, 2, \dots, 8$) and $D_\alpha = \partial_\alpha - ig_s G_\alpha^h t^h$ [26–28].

Consider the Feynman graphes of the leading orders of these nonzero correlation functions, each one contains two heavy quark lines and four light quark lines. If each heavy quark line emits a gluon and each light quark line contributes quark-antiquark pair, they form the quark-gluon operator $g_s G_{\alpha\beta} g_s G_{\mu\nu} \bar{q}q \bar{q}q \bar{q}q \bar{q}q$ with dimension 16. This operator lead to vacuum condensates $\langle(\alpha_s/\pi)GG\rangle\langle\bar{q}q\rangle^4$ and $\langle\bar{q}g_s\sigma Gq\rangle^2\langle\bar{q}q\rangle^2$. So, the highest dimension of the considered condensates is up to 16 in the present work. As for the truncation of the order $\mathcal{O}(\alpha_s^k)$, the contributions of the related vacuum condensates are tiny while $k \geq 3/2$; it is accurate enough for us to calculate terms for $k \leq 1$ [29]. Detailed calculations show that the vacuum condensates $\langle\bar{q}q\rangle$ and $\langle\bar{q}g_s\sigma Gq\rangle$ have no contribution for all the considered states in this paper. Thus, the chosen vacuum condensates of the operator product expansions are $\langle(\alpha_s/\pi)GG\rangle$, $\langle\bar{q}q\rangle^2$, $\langle\bar{q}g_s\sigma Gq\rangle\langle\bar{q}q\rangle$, $\langle\bar{q}g_s\sigma Gq\rangle^2$, $\langle(\alpha_s/\pi)GG\rangle\langle\bar{q}q\rangle^2$, $\langle\bar{q}q\rangle^4$, $\langle\bar{q}g_s\sigma Gq\rangle\langle\bar{q}q\rangle^3$, $\langle\bar{q}g_s\sigma Gq\rangle^2\langle\bar{q}q\rangle^2$, and $\langle(\alpha_s/\pi)GG\rangle\langle\bar{q}q\rangle^4$.

After the calculations of the integrals in the coordinate space for the light quarks and momentum space for heavy quarks of the correlation functions, we perform the Borel transforms for both the hadronic sides and the QCD sides, we get the QCD sum rules,

$$\lambda_i^2 \exp\left(-\frac{M_i^2}{T^2}\right) = \int_{\Delta^2}^{s_0} ds \rho_i(s) \exp\left(-\frac{s}{T^2}\right), \quad (8)$$

where $\rho_i(s)$ are the QCD spectral densities of J_i , $\Delta^2 = 4m_c^2$. As for the continuum threshold parameters s_0 , we take the exponential values [24, 29, 30],

$$\sqrt{s_0} = M_Z + (0.5 \sim 0.7)\text{GeV}, \quad (9)$$

where M_Z are the masses of the ground states.

We differentiate the Equation (8) with respect to $\tau = 1/T^2$ and eliminate the pole residues, the extracted masses are shown as

$$M_i^2 = \frac{-d/d\tau \int_{\Delta^2}^{s_0} ds \rho_i(s) \exp(-s\tau)}{\int_{\Delta^2}^{s_0} ds \rho_i(s) \exp(-s\tau)}. \quad (10)$$

3. Numerical Results and Discussions

We apply the standard values of the vacuum condensates $\langle\bar{q}q\rangle = -(0.24 \pm 0.01 \text{ GeV})^3$, $\langle\bar{q}g_s\sigma Gq\rangle = m_0^2\langle\bar{q}q\rangle \text{ GeV}^2$, $m_0^2 = (0.8 \pm 0.1) \text{ GeV}^2$, $\langle(\alpha_s/\pi)GG\rangle = (0.33 \text{ GeV})^4$ at the energy scale $\mu = 1 \text{ GeV}$ [26, 31], and choose the $\overline{\text{MS}}$ mass $m_c(m_c) = 1.275 \pm 0.025 \text{ GeV}$ [32]. Consider the energy-scale dependence for the above condensates,

$$\langle\bar{q}q\rangle(\mu) = \langle\bar{q}q\rangle(1 \text{ GeV}) \left[\frac{\alpha_s(1 \text{ GeV})}{\alpha_s(\mu)} \right]^{12/(33-2n_f)},$$

$$\langle\bar{q}g_s\sigma Gq\rangle(\mu) = \langle\bar{q}g_s\sigma Gq\rangle(1 \text{ GeV}) \left[\frac{\alpha_s(1 \text{ GeV})}{\alpha_s(\mu)} \right]^{12/(33-2n_f)},$$

$$m_c(\mu) = m_c(m_c) \left[\frac{\alpha_s(\mu)}{\alpha_s(m_c)} \right]^{12/(33-2n_f)},$$

$$\alpha_s(\mu) = \frac{1}{b_0 t} \left[1 - \frac{b_1 \log t}{b_0^2 t} + \frac{b_1^2 (\log^2 t - \log t - 1) + b_0 b_2}{b_0^4 t^2} \right], \quad (11)$$

where $t = \log(\mu^2/\Lambda_{\text{QCD}}^2)$, $b_0 = (33 - 2n_f)/12\pi$, $b_1 = (153 - 19n_f)/24\pi^2$, $b_2 = (2857 - (5033/9)5033/9n_f + (325/27)325/27n_f^2)/128\pi^3$, and $\Lambda_{\text{QCD}} = 213 \text{ MeV}$, 296 MeV , and 339 MeV for the flavors $n_f = 5, 4, 3$, respectively [32–34]. In this paper, we choose the flavor number $n_f = 4$ for all the states. Compared with the heavy quarks, the masses of the light quarks are too small to make effective difference. So, we ignore

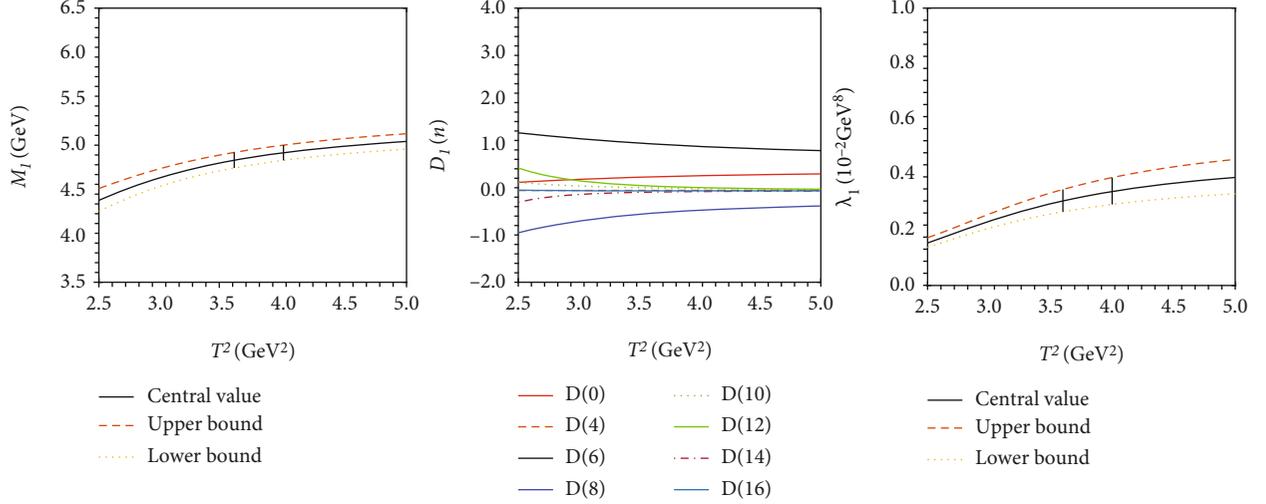


FIGURE 1: The numerical results of J_1 : $M_1 - T^2$ for (a), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values; $D_1(n) - T^2$ for (b); $\lambda_1 - T^2$ for (c), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values.

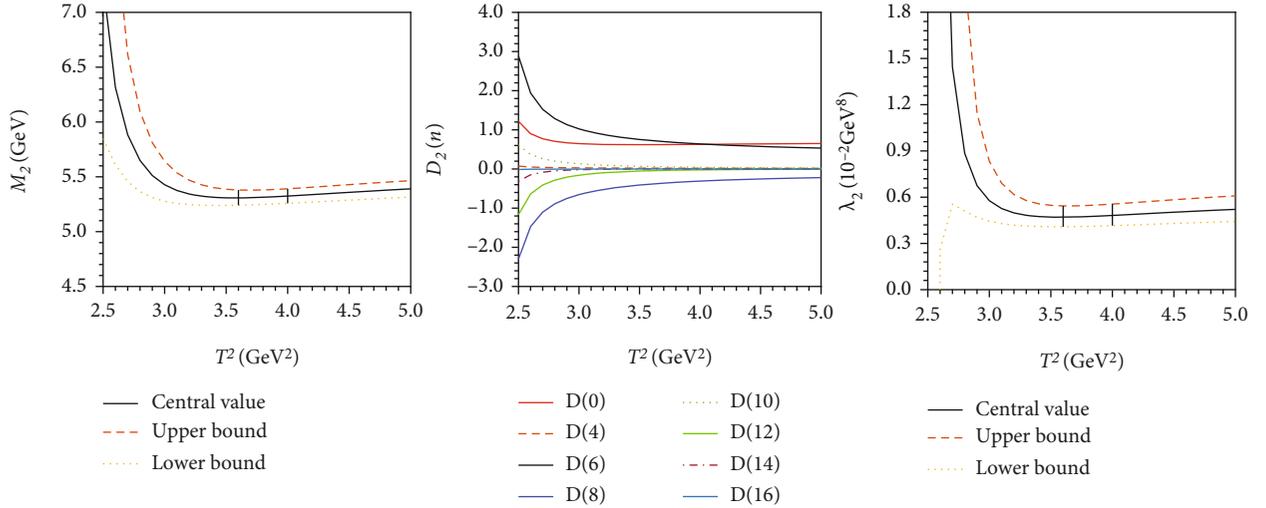


FIGURE 2: The numerical results of J_2 : $M_2 - T^2$ for (a), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values; $D_2(n) - T^2$ for (b); $\lambda_2 - T^2$ for (c), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values.

the masses of the light quarks. The masses of the baryon and anti-baryon $M_{\Sigma_c} = M_{\bar{\Sigma}_c} = 2.455$ GeV are from the particle data group [32]. We abandon the expressions of the lengthy spectral densities of these currents which are dozens of pages, one can contact us via the email. Note that Equation (10) depends on the energy scale μ [19, 35], we apply the energy scale formula to determine the best energy scales of the QCD spectral densities [27],

$$\mu = \sqrt{M_{X|YZ}^2 - 4\mathbb{M}_c^2}, \quad (12)$$

where \mathbb{M}_c is the effective charm quark mass, we choose the value $\mathbb{M}_c = 1.84 \pm 0.01$ GeV [36]. Interestingly, in Ref. [29],

for the vector current $J_\mu = \bar{J}_\Lambda \gamma_\mu J_\Lambda$, if we set $\mathbb{M}_c = 1.84$ GeV, the central value of the mass of $\Lambda_c \bar{\Lambda}_c$ with the $J^P = 1^-$ is 4.57 GeV with $\sqrt{s_0} = 2m_\Lambda + 0.66$ GeV and PC = 39% – 53% which indicates a possible bound $\Lambda_c \bar{\Lambda}_c$ baryonium state, where $m_\Lambda = 2.29$ GeV is the mass of Λ_c [32]. As part of the systematic study for the dibaryons formed by the charmed baryons, the detailed results of Ref. [29] are also shown in Table 1. From Equation (10), we find that M_Z relies on the threshold parameter s_0 and Borel parameter T^2 . Of course, at first, we do not know the suitable energy scale, so we try and adjust the value of μ via trial and error. From the energy scale formula Equation (12), one can determine that M_Z increases with the parameter μ ; however, numerical results show that mass derived from Equation (10) decreases with the increment of

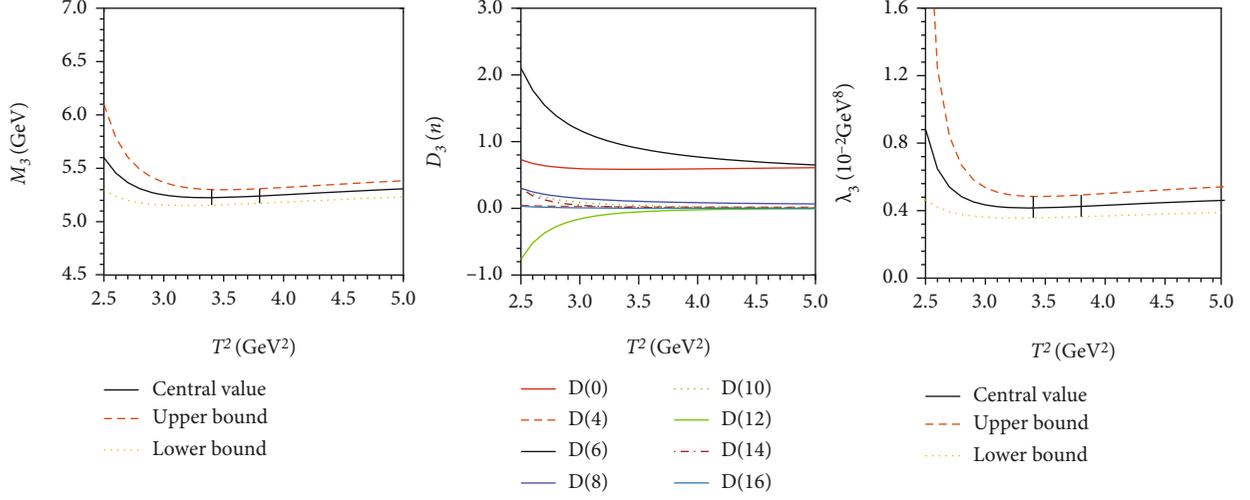


FIGURE 3: The numerical results of J_3 : $M_3 - T^2$ for (a), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values; $D_3(n) - T^2$ for (b); $\lambda_3 - T^2$ for (c), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values.

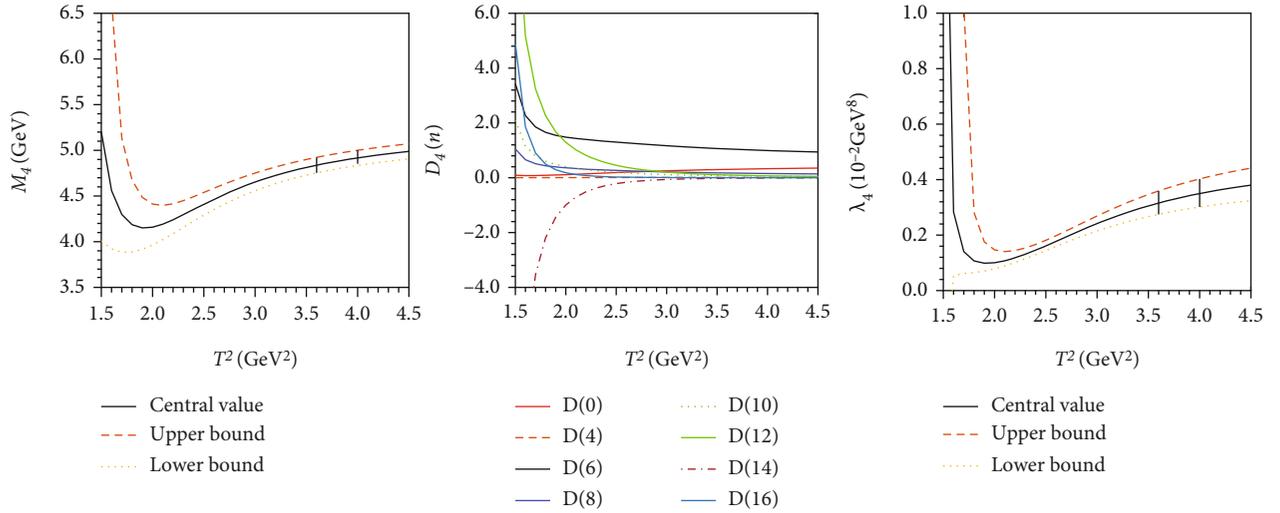


FIGURE 4: The numerical results of J_4 : $M_4 - T^2$ for (a), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values; $D_4(n) - T^2$ for (b); $\lambda_4 - T^2$ for (c), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values.

the parameter μ , so we can always find the best energy scale and determine the mass which satisfies the energy scale formula. As for the values of s_0 and T^2 , we search for the best Borel window which satisfies the following two principles, firstly, the pole dominance criterion; here, we define the pole contributions (PC) of the QCD sum rules as

$$\text{PC} = \frac{\int_{4m_c^2}^{s_0} ds \rho_{\text{QCD}}(s) \exp(-s/T^2)}{\int_{4m_c^2}^{\infty} ds \rho_{\text{QCD}}(s) \exp(-s/T^2)}, \quad (13)$$

secondly, the convergence of the contribution of operator product expansion, for which, we define the contributions of the vacuum condensates of dimension n as

$$D(n) = \frac{\int_{4m_c^2}^{s_0} ds \rho_{\text{QCD},n}(s) \exp(-s/T^2)}{\int_{4m_c^2}^{s_0} ds \rho_{\text{QCD}}(s) \exp(-s/T^2)}. \quad (14)$$

Due to the uncertainties of the input parameters, the uncertainties of the extracted masses and pole residues are determined by the formula,

$$\delta = \sqrt{\sum_i \left(\frac{\partial f}{\partial x_i} \right)^2 \Big|_{x_i = \bar{x}_i} (x_i - \bar{x}_i)^2}, \quad (15)$$

where f represent the masses or the pole residues, x_i are the input parameters $\langle \bar{q}q \rangle$, $\langle \bar{q}g\sigma Gq \rangle$, \dots , \bar{x}_i denote the central values of these parameters. Since it is difficult to find the

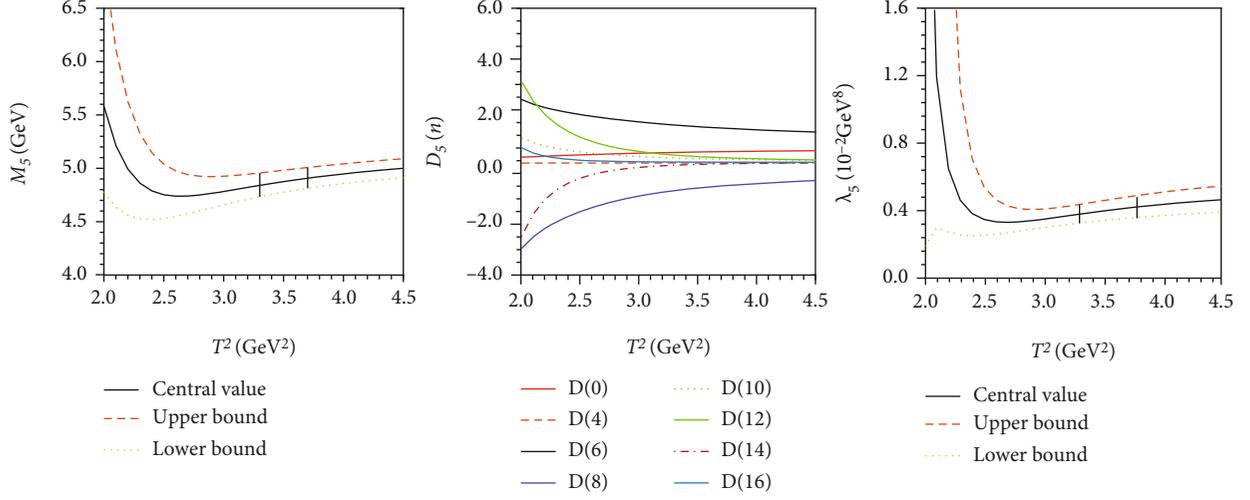


FIGURE 5: The numerical results of J_5 : $M_5 - T^2$ for (a), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values; $D_5(n) - T^2$ for (b); $\lambda_5 - T^2$ for (c), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values.

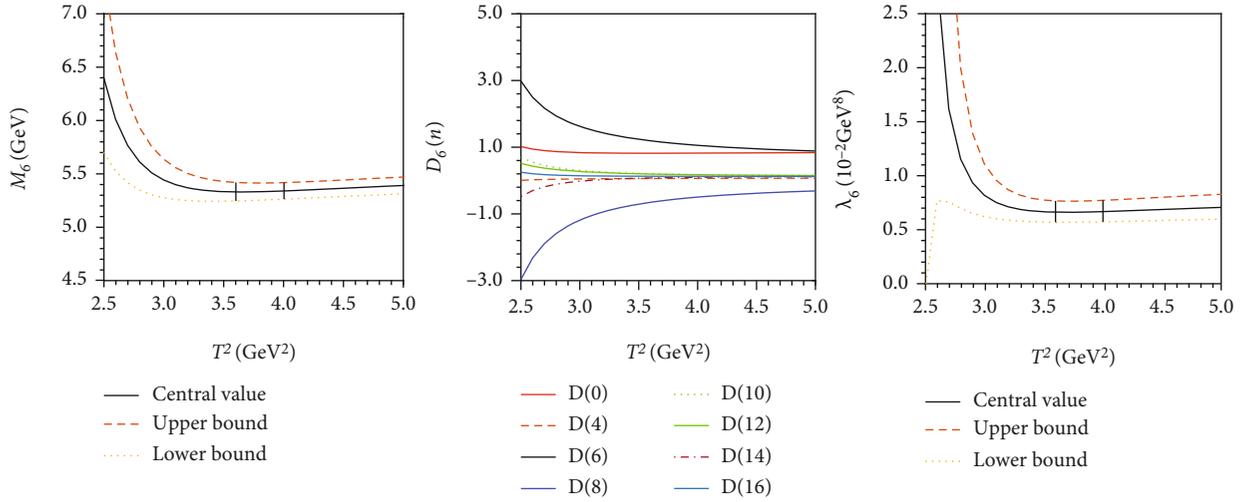


FIGURE 6: The numerical results of J_6 : $M_6 - T^2$ for (a), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values; $D_6(n) - T^2$ for (b); $\lambda_6 - T^2$ for (c), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values.

analytic expressions of the partial derivatives $\partial f/\partial x_i$, we carry out the approximate numerical calculations for the uncertainties via,

$$\left(\frac{\partial f}{\partial x_i}\right)^2 (x_i - \bar{x}_i)^2 \approx [f(\bar{x}_i \pm \Delta x_i) - f(\bar{x}_i)]^2. \quad (16)$$

Except J_7 , we find the Borel windows for the other seven currents. Their numerical results are shown in Figures 1–7, respectively. In each figure, we show the masses, dimensional contributions of vacuum condensates and pole residues of each current on $M - T^2$, $D(n) - T^2$, and $\lambda - T^2$ graphs, respectively. Currents J_1 , J_2 , J_3 , and J_4 are related to the $\Sigma_c \bar{\Sigma}_c$ baryonium structures. For the vector current J_1 , the extracted mass of the corresponding $\Sigma_c \bar{\Sigma}_c$ state with the J^P

$= 1^-$ is 4.88 GeV with the upper uncertainty being 0.09 GeV. Since the threshold of the two Σ_c and $\bar{\Sigma}_c$ baryon constituents is 4.91 GeV, so, we could not say for sure whether it is a molecular state formed by the two baryons or not. On the safe side, it is a possible molecular or compact six-quark state. In the same way, $\Sigma_c \bar{\Sigma}_c$ with the $J^P = 0^-$ is also a possible molecular or compact six-quark state. As for axial-vector current J_2 and scalar current J_3 , their extracted masses of the corresponding states with the $J^P = 1^+$ and 0^+ are 5.31 GeV and 5.23 GeV, respectively. They are all above the thresholds of the two Σ_c and $\bar{\Sigma}_c$ baryon constituents even if we take into the account of the lower bounds of uncertainties of the extracted masses. Both these two are the baryonium resonance states.

Currents J_5 , J_6 , J_7 , and J_8 are related to the $\Sigma_c \Sigma_c$ dibaryon states. For the current $J_{7,\alpha} = J_\Sigma^T C \gamma_\alpha J_\Sigma$, its related

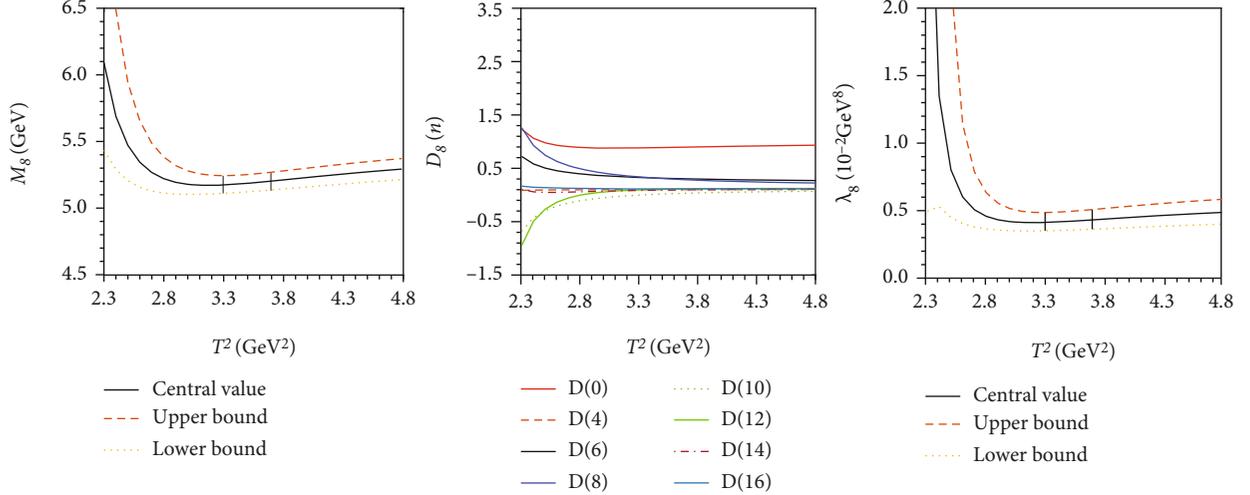


FIGURE 7: The numerical results of J_8 : $M_8 - T^2$ for (a), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values; $D_8(n) - T^2$ for (b); $\lambda_8 - T^2$ for (c), dashed curve is drawn from the upper values of the input parameters and dotted curve is from the lower values.

correlation function $\Pi_{7;\alpha\beta}(p) \equiv 0$. It indicates that there is no such $\Sigma_c \Sigma_c$ dibaryon state with the $J^P = 1^+$. The extracted mass from J_5 is 4.89 GeV which is 20 MeV below the threshold of the two Σ_c baryon constituents. Due to the uncertainties of input parameters, the related uncertainty of the mass of this state is up to 0.1 GeV, it is also a possible $\Sigma_c \Sigma_c$ molecular or a compact six-quark state with the $J^P = 0^+$. We also get the masses of J_6 and J_8 , they are all above the thresholds of the two baryon constituents. The detailed results of each state are listed in Table 1.

For all these seven currents, it is clear to be seen from the $D_i(n) - T^2$ curves in each titled figures, the most important items are the leading order, $\langle \bar{q}q \rangle^2$ and $\langle \bar{q}g_s \sigma Gq \rangle \langle \bar{q}q \rangle$. We find that the higher dimensional vacuum condensates ($n \geq 10$) play the less important roles. So, the convergence of the operator product expansion holds very well. It is worth mentioning that condensate $\langle \bar{q}q \rangle^2$ plays a dominant role for the current J_4 , what is more, the leading order of J_8 contributes the most important part for both the mass and pole residue of the corresponding state.

In 2017, the LHCb has observed the doubly-charmed baryon Ξ_{cc}^{++} with mass of 3621.40 ± 0.78 MeV [37]. It is proposed that there is large binding energy $B(cc) = 129$ MeV of the two heavy c -quarks in a baryon [38]. In Ref. [39], the quark-level analogue of nuclear fusion for the two Λ_c baryons was discussed. In a similar way, we also consider the quark-rearrangement reaction, the possible decays of $\Sigma_c \Sigma_c$ dibaryon states are given by

$$\begin{aligned} \Sigma_c \Sigma_c &\longrightarrow \Xi_{cc}^{++} n, \\ cud\ cud &\longrightarrow ccuddu. \end{aligned} \quad (17)$$

Since the J^P of Ξ_{cc}^{++} is unknown yet, thus, it is also impossible for us to determine the J^P of the produced neutron after the decay. We can only estimate the upper values of the exo-

thermic energy Q_{upper} of each $\Sigma_c \Sigma_c$ quark-rearrangement reaction by setting the minimum mass of the neutron as $m_n = 1440$ MeV. For $\Sigma_c \Sigma_c$ with the $J^P = 0^+$, it is impossible to have such kind of decay since its mass is less than the sum of $m_{\Xi_{cc}^{++}}$ and m_n . As for 0^- and 1^- , the corresponding Q_{upper} are 310 MeV and 170 MeV, respectively.

4. Conclusions

In this article, we construct the eight interpolating currents to study the $\Sigma_c \bar{\Sigma}_c$ baryonium and $\Sigma_c \Sigma_c$ dibaryon states via the QCD sum rules. The correlation function of J_7 is zero which indicates that there is no such $\Sigma_c \Sigma_c$ dibaryon state with the $J^P = 1^+$. We find two possible $\Sigma_c \bar{\Sigma}_c$ molecular or compact six-quark states with the $J^P = 1^-$ and 0^- . Their masses are $4.88_{-0.08}^{+0.09}$ GeV and $4.88_{-0.08}^{+0.08}$ GeV, respectively. For the currents J_2, J_3, J_6 , and J_8 , the extracted masses of these four states are $5.31_{-0.07}^{+0.07}$ GeV, $5.23_{-0.07}^{+0.07}$ GeV, $5.37_{-0.08}^{+0.08}$ GeV, and $5.23_{-0.06}^{+0.06}$ GeV, respectively. Both the four are above the thresholds of the two baryon constituents. We conclude that $\Sigma_c \bar{\Sigma}_c$ with the $J^P = 1^+$ and 0^+ are baryonium resonances states and $\Sigma_c \Sigma_c$ with the $J^P = 0^-$ and 1^- are dibaryon resonances states. We find the central value of the mass of the state with the $J^P = 0^+$ is just 20 MeV below the threshold of the two Σ_c baryon constituents. Considering the uncertainty, $\Sigma_c \Sigma_c$ dibaryon with the $J^P = 0^+$ is also a possible molecular state. For all the bound states related to $\Sigma_c \bar{\Sigma}_c$ and $\Sigma_c \Sigma_c$, we need more data, such as the electromagnetic form factors, to differentiate whether the bound dibaryon or baryonium is a molecular state or a compact six-quark entity. This is also the task of our further work. The possible decays, the quark-rearrangement reactions of each $\Sigma_c \Sigma_c$ states are discussed. The upper values of the exothermic energy are calculated for these decays which may be experimentally verified in the future.

Data Availability

All data included in this manuscript are available upon request by contacting the corresponding author.

Conflicts of Interest

The authors declare that they have no conflicts of interest.

Acknowledgments

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