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Analysis of $\frac{1}{2}^+$ baryon states containing fourth-family quarks from QCD sum rules*

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Abstract: When the fourth generation of quarks have sufficiently small mixing with ordinary standard-model quarks, the hadrons made up from these quarks can be long-lived enough. We analyze the $\frac{1}{2}^+$ baryon states containing fourth-generation quarks and standard-model quarks, i.e. the charm or bottom quarks, in the QCD sum rules approach. Considering the perturbative and two gluon condensate contributions in the calculation, we give the numerical results of the masses and pole residues.

Key words: sum rules, the fourth generation of quarks, baryons

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1 Introduction

The number of lepton and quark families is one of the problems that cannot be addressed by the Standard Model (SM). The existence of the fourth generation of quarks was predicted about 20 years ago based on flavor democracy arguments [1–3]. The addition of the fourth generation of quarks, which is one of the simple extensions of the SM, has many important physical consequences, such as in flavor physics and CP violation [4–6].

Recent analysis of the Tevatron data implies $m_{\mathrm{d}_4} > 372$ GeV [7] and $m_{\mathrm{u}_4} > 358$ GeV [8], and this can form hadronic states due to the small mixing between these quarks and ordinary SM quarks [9–11]. The condition for the existence of new hadrons containing ultra-heavy quarks (Q) is [12]:

$$|V_{\mathrm{Qq}}| \leqslant \left(\frac{100}{m_{\mathrm{Q}}}\right)^{3/2}.$$

For top quarks, the above formula leads to $V_{\mathrm{tq}} < 0.44$, but the single top quark production experiment at the Tevatron gives $V_{\mathrm{tb}} > 0.74$ [13]. Thus, there will be no top flavor hadrons. As for the fourth-family quarks, when they have small enough mixing with SM quarks, they can form new hadronic states with the ordinary SM quarks. The parameterization obeying this condition has been proposed in Ref. [14].

If such new hadronic states exist, it will be possible to observe them at the Tevatron or LHC. Therefore, it is necessary to investigate their properties theoretically and phenomenologically. In order to calculate their hadronic parameters, such as their masses and pole residues, we need to consult some nonperturbative approaches. Among the nonperturbative methods, the QCD sum rules originally devised for low-energy hadronic physics, have been successfully extended to heavy quark (c or b) physics, and remain one of the most applicable and predictive approaches to hadronic physics[15–17]. Recently, V. Bashiry et al. calculated the masses and decay constants of mesons containing the fourth generation of quarks from QCD sum rules[18]. In this work, we analyze the $\frac{1}{2}^+$ baryon states containing the fourth-generation quarks and the charm or bottom quark.

The paper is organized as follows. In Section 2, we analytically calculate the masses and pole residues of the considered baryon states within the framework of QCD sum rules. The numerical results are presented in Section 3 and some discussions are included in Section 4.

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2 QCD sum rules for the new $\frac{1}{2}^+$ baryon states

In order to analyze these new baryon states within the framework of the QCD sum rules, we must write down their interpolating currents. As we know, Ioffe currents are more suitable for the analysis of the masses and other static properties of the baryons [19], and we use the Ioffe currents $J(x)$ for the analysis of the $\frac{1}{2}^+$ baryon states containing the fourth-generation quarks and the charm or bottom quark:

$$J(x)=\epsilon^{ijk}(Q_i^{\mathrm{T}}(x)C\gamma_\mu Q_j(x))\gamma_5\gamma^\mu q_k(x), \qquad (1)$$

where Q and q represent the fourth-family quark u_4 1) and the SM quark b(c), i, j and k are the color indexes, and C is the charge conjunction matrix.

The corresponding $\frac{1}{2}^-$ baryon states can be interpolated by the currents $J_-=\mathrm{i}\gamma_5 J_+$ [20], where J_+ denotes the currents $J(x)$. The correlation functions $\Pi_\pm(p)$ are defined by

$$\Pi_\pm(p)=\mathrm{i}\int d^4x\mathrm{e}^{\mathrm{i}p\cdot x}\langle 0|T\{J_\pm(x)\bar{J}_\pm(0)\}|0\rangle, \qquad (2)$$

and can be decomposed as

$$\Pi_\pm(p)=\not{p}\Pi_1(p^2)\pm\Pi_0(p^2), \qquad (3)$$

due to the Lorentz covariance. The currents J_+ couple not only to the positive-parity, but also to the negative-parity baryon states [21], $\langle 0|J_+|B^-\rangle\cdot\langle B^-|\bar{J}_+|0\rangle=-\gamma_5\langle 0|J_-|B^-\rangle\langle B^-|\bar{J}_-|0\rangle\gamma_5$, where B^- denotes the negative-parity baryon states. In order to obtain the hadronic representation, we insert a complete set of intermediate baryon states with the same quantum numbers as the current $J_\pm(x)$ into the correlation functions $\Pi_\pm(p)$. After isolating the pole terms of the lowest states of the new baryons, we obtain:

$$\Pi_+(p)=\lambda_+^2\frac{\not{p}+M_+}{M_+^2-p^2}+\lambda_-^2\frac{\not{p}-M_-}{M_-^2-p^2}+\cdots, \qquad (4)$$

where $M_\pm$ are the masses of the lowest states with parity $\pm$, respectively, and $\lambda_\pm$ are the corresponding pole residues. If we take $\vec{p}=0$, we have

$$\begin{aligned}\lim_{\epsilon\to 0}\frac{\mathrm{Im}\Pi_+(p_0+\mathrm{i}\epsilon)}{\pi} &= \lambda_+^2\frac{\gamma_0+1}{2}\delta(p_0-M_+)\\ &\quad+\lambda_-^2\frac{\gamma_0-1}{2}\delta(p_0-M_-)+\cdots\\ &=\gamma_0 A(p_0)+B(p_0)+\cdots, \qquad (5)\end{aligned}$$

where

$$\begin{aligned}A(p_0)&=\frac{1}{2}\left[\lambda_+^2\delta(p_0-M_+)+\lambda_-^2\delta(p_0-M_-)\right],\\ B(p_0)&=\frac{1}{2}\left[\lambda_+^2\delta(p_0-M_+)-\lambda_-^2\delta(p_0-M_-)\right],\end{aligned} \qquad (6)$$

the $A(p_0)+B(p_0)$ and $A(p_0)-B(p_0)$ contain the contributions from the positive-parity and negative-parity baryon states, respectively.

The correlation functions are calculated in the deep Euclidean region in the QCD side. When including perturbative α_s corrections, a factor $(1+\alpha_s/\pi)$ appears, while α_s/π is much smaller than 0.1 at the energy scale Tev, therefore we can safely ignore α_s corrections. On the other hand, as the heavy quark condensates are suppressed by the inverse powers of the heavy quark mass, we calculate the two-gluon condensate as the first nonperturbative contributions. We then use the dispersion relation to obtain the spectral densities $\rho^A(p_0)$ and $\rho^B(p_0)$ at the level of quark-gluon degrees of freedom after taking the limit $\vec{p}=0$, where $\rho^A(p_0)$ and $\rho^B(p_0)$ correspond to the tensor structures γ_0 and 1, respectively. Finally, we introduce the weight functions $\exp\left[-\frac{p_0^2}{T^2}\right]$, and obtain the following sum rule:

$$\lambda_+^2\exp\left[-\frac{M_+^2}{T^2}\right]=\int_\Delta^{\sqrt{s_0}}\mathrm{d}p_0[\rho^A(p_0)+\rho^B(p_0)]\exp\left[-\frac{p_0^2}{T^2}\right]. \qquad (7)$$

Taking the derivative with respect to $1/T^2$, we obtain another sum rule:

$$\lambda_+^2M_+^2\exp\left[-\frac{M_+^2}{T^2}\right]=\int_\Delta^{\sqrt{s_0}}\mathrm{d}p_0[\rho^A(p_0)+\rho^B(p_0)]p_0^2\exp\left[-\frac{p_0^2}{T^2}\right], \qquad (8)$$

where

$$\begin{aligned}\rho^A(p_0) = {}&\frac{3m_Q^2p_0}{8\pi^4}\int_{\alpha_i}^{\alpha_f}\mathrm{d}\alpha\int_{\beta_i}^{\beta_f}\mathrm{d}\beta(1-\alpha-\beta)(p_0^2-\tilde{m}^2)+\frac{3p_0}{8\pi^4}\int_{\alpha_i}^{\alpha_f}\mathrm{d}\alpha\int_{\beta_i}^{\beta_f}\mathrm{d}\beta\alpha\beta(1-\alpha-\beta)(p_0^2-\tilde{m}^2)(5p_0^2-3\tilde{m}^2)\\ &+\frac{m_Q^2}{192\pi^2}\left\langle\frac{\alpha_sGG}{\pi}\right\rangle\int_{\alpha_i}^{\alpha_f}\mathrm{d}\alpha\int_{\beta_i}^{\beta_f}\mathrm{d}\beta\delta(p_0-\tilde{m})\left[4\frac{\beta(1-3\alpha-\beta)+2\alpha(1-\alpha)}{\beta^2}\left(1-\frac{p_0^2}{T^2}\right)\right.\end{aligned}$$

1) As the mass difference between d_4 and u_4 is small, we will refer to both of the two members of the fourth family by u_4.

$$+4\frac{6\alpha^2-3\beta(1-\beta)+(1-\beta)(2-7\alpha)}{\beta^2}+\frac{3}{\beta}\Bigg]+\frac{p_0}{96\pi^2}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle\int_{\alpha_i}^{\alpha_f}d\alpha\int_{\beta_i}^{\beta_f}d\beta\left(\frac{1-\alpha-\beta}{1-\beta}+(9\alpha+\beta)\right)$$

$$+\frac{p_0^2}{64\pi^2}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle\int_{\alpha_i}^{\alpha_f}d\alpha\int_{\beta_i}^{\beta_f}d\beta\alpha\delta(p_0-\tilde{m})+\frac{m_q^2}{192\pi^2}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle\int_{\alpha_i}^{\alpha_f}d\alpha\int_{\beta_i}^{\beta_f}d\beta\delta(p_0-\tilde{m})$$

$$\cdot\left[\frac{\alpha(12\beta^2-10\beta-1)}{(1-\alpha-\beta)^2(1-\beta)}+\frac{2\alpha\beta}{(1-\alpha-\beta)^1}\left(1-\frac{p_0^2}{T^2}\right)\right]$$

$$-\frac{m_q^2 m_Q^2}{96\pi^2 T^2}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle\int_{\alpha_i}^{\alpha_f}d\alpha\int_{\beta_i}^{\beta_f}d\beta\frac{1}{(1-\alpha-\beta)^2}\delta(p_0-\tilde{m}), \quad (9)$$

$$\rho^B(p_0) = \frac{3m_q}{8\pi^4}\int_{\alpha_i}^{\alpha_f}d\alpha\int_{\beta_i}^{\beta_f}d\beta\alpha\beta(p_0^2-\tilde{m}^2)(2p_0^2-\tilde{m}^2)+\frac{3m_q m_Q^2}{4\pi^4}\int_{\alpha_i}^{\alpha_f}d\alpha\int_{\beta_i}^{\beta_f}d\beta(p_0^2-\tilde{m}^2)$$

$$+\frac{m_q m_Q^2}{96\pi^2 p_0}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle\int_{\alpha_i}^{\alpha_f}d\alpha\int_{\beta_i}^{\beta_f}d\beta\delta(p_0-\tilde{m})\left[\left(1-\frac{2p_0^2}{T^2}\right)\left(\frac{\alpha+2\beta}{\beta^2}+\frac{1}{1-\alpha-\beta}\right)\right.$$

$$\left.+\frac{4(2-\alpha-2\beta)}{\beta^2}+\frac{3}{\beta(1-\alpha-\beta)}+\frac{4(\alpha+\beta)}{(1-\alpha-\beta)^2}\right]+\frac{m_q p_0}{192\pi^2}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle\int_{\alpha_i}^{\alpha_f}d\alpha\int_{\beta_i}^{\beta_f}d\beta\delta(p_0-\tilde{m})$$

$$\cdot\left[3+\frac{3\alpha}{1-\alpha-\beta}+\frac{4\alpha\beta(3\alpha+3\beta-2)}{(1-\alpha-\beta)^2}+\frac{\alpha\beta}{1-\alpha-\beta}\left(3-\frac{2p_0^2}{T^2}\right)\right]$$

$$+\frac{m_q}{192\pi^2}\left\langle\frac{\alpha_s GG}{\pi}\right\rangle\int_{\alpha_i}^{\alpha_f}d\alpha\int_{\beta_i}^{\beta_f}d\beta\left[12+\frac{12\alpha}{1-\alpha-\beta}+\frac{8\alpha\beta(3\alpha+3\beta-1)}{(1-\alpha-\beta)^2}\right], \quad (10)$$

s_0 is the threshold parameter, T^2 is the Borel parameter,

$$\alpha_i = \frac{-8a^2+2a+2ab+2a\sqrt{(b-4a-1)^2-16a}}{-b-4a+\sqrt{(b-4a-1)^2-16a}+1},$$

$$\alpha_f = \frac{-8a^2+2a+2ab-2a\sqrt{(b-4a-1)^2-16a}}{-b-4a-\sqrt{(b-4a-1)^2-16a}+1},$$

$$\beta_i = \frac{\alpha-\alpha^2-b\alpha+2a\alpha-a-\sqrt{(b\alpha-2a\alpha+a+\alpha^2-\alpha)^2-4(a\alpha-a\alpha^2)(\alpha-a)}}{2(\alpha-a)},$$

$$\beta_f = \frac{1+4a-b+\sqrt{(1+4a-b)^2-16a}}{2}-\alpha,$$

$$\tilde{m}^2 = \frac{m_q^2}{1-\alpha-\beta}+\frac{m_Q^2(\alpha+\beta)}{\alpha\beta},$$

$$a = m_Q^2/p_0^2,\quad b=m_q^2/p_0^2,\quad \Delta=2m_Q+m_q.$$

3 Numerical results

To obtain the numerical values for the masses and pole residues of the considered bound states from the above QCD sum rules, we take the mass of u_4 in the interval $m_{u_4} = 450-550$ GeV, $m_b = 4.7$ GeV, $m_c = 1.35$ GeV and

$$\left\langle\frac{\alpha_s GG}{\pi}\right\rangle = 0.012\ \text{GeV}^4.$$

In the standard QCD sum rules, there are two criteria for choosing the Borel parameter T^2 and threshold parameter s_0, i.e. the pole dominance and convergence of the operator product expansion [15–17]. We impose these two criteria on the considered baryon states to choose the Borel parameter T^2 and threshold parameter s_0.

We first give the masses with the variation in the threshold parameter s_0 in Fig. 1, and find that the masses obtained are not sensitive to the threshold parameters, although they increase with the threshold parameters. Since we have no information about the spectrum of the baryon states containing the fourth-family quarks, the dispersion relations used to obtain the spectral densities would be questionable at first sight, but this insensitivity indicates the reasonableness of the dispersion relations. In the calculation, we take the threshold parameters in the interval 1100 GeV $\leqslant \sqrt{s_0} \leqslant$ 1200 GeV for the $u_4u_4b(c)$ channel, and 550 GeV $\leqslant \sqrt{s_0} \leqslant$ 610 GeV for the $bb(cc)u_4$ channel.

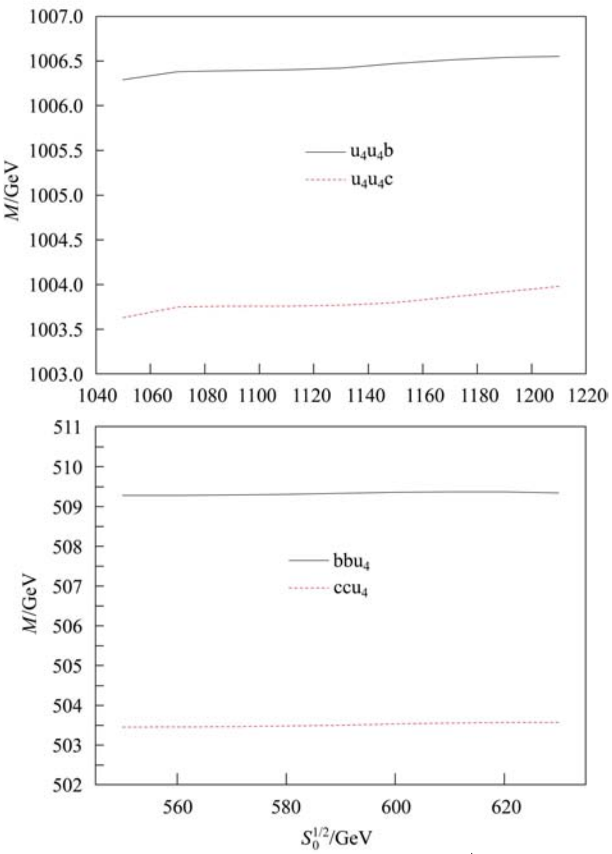

Fig. 1. The masses, M, of the new $\frac{1}{2}^+$ baryon states with variations in the threshold parameters, s_0. The Borel parameters, T^2, are taken to be the central values, and the masses of the fourth-generation quark u_4 are taken to be the central values, 500 GeV.

Next, we look for the working region of the Borel parameter T^2. The upper limit of T^2 is obtained by the requirement that the contributions of the higher states and continuum are only a few percent of the total dispersion integral. The lower bound of T^2 is obtained by the requirement that the operator product expansion (OPE) is convergent. We plot the masses with variations of the Borel parameters T^2 in Fig. 2, and the working region for the obtained Borel parameters is 1380 GeV^2 $\leqslant T^2 \leqslant$ 1580 GeV^2 for the $u_4u_4b(c)$ channel, and 360 GeV^2 $\leqslant T^2 \leqslant$ 450 GeV^2 for the $bb(cc)u_4$ channel. Taking into account the uncertainties of the relevant parameters, we obtain reasonable values of the masses of the new $\frac{1}{2}$ baryon states, as shown in Table 1.

We then investigate the pole residues with variations of the Borel parameters T^2, as shown in Fig. 3. Taking into account the uncertainties of the relevant parameters, we obtain the values of the pole residues of the new $\frac{1}{2}$ baryon states, as shown in Table 2. In Ref. [22], we obtain the pole residues of the $\frac{1}{2}^+$ doubly heavy baryon states using QCD sum rules, i.e.

$$\begin{aligned} \lambda_+(\varXi_{cc}) &= 0.115 \pm 0.027, \\ \lambda_+(\varOmega_{cc}) &= 0.138 \pm 0.030, \\ \lambda_+(\varXi_{bb}) &= 0.252 \pm 0.064, \\ \lambda_+(\varOmega_{bb}) &= 0.311 \pm 0.077. \end{aligned} \tag{11}$$

These doubly heavy baryon states are similar to the new $\frac{1}{2}^+$ baryon states in the quark structure, so by analogy, we see that the predicted pole residues are in a reasonable region.

Table 1. The masses of the new $\frac{1}{2}^+$ baryon states with different masses of the fourth-generation quark.

mass/GeV	m_{u_4}=450 GeV	m_{u_4}=500 GeV	m_{u_4}=550 GeV
u_4u_4b	906.77±0.13	1006.48±0.24	1107.42±0.41
u_4u_4c	904.16±0.28	1003.84±0.31	1105.01±0.47
bbu_4	459.48±0.22	509.33±0.29	559.53±0.46
ccu_4	453.67±0.18	503.50±0.27	553.94±0.39

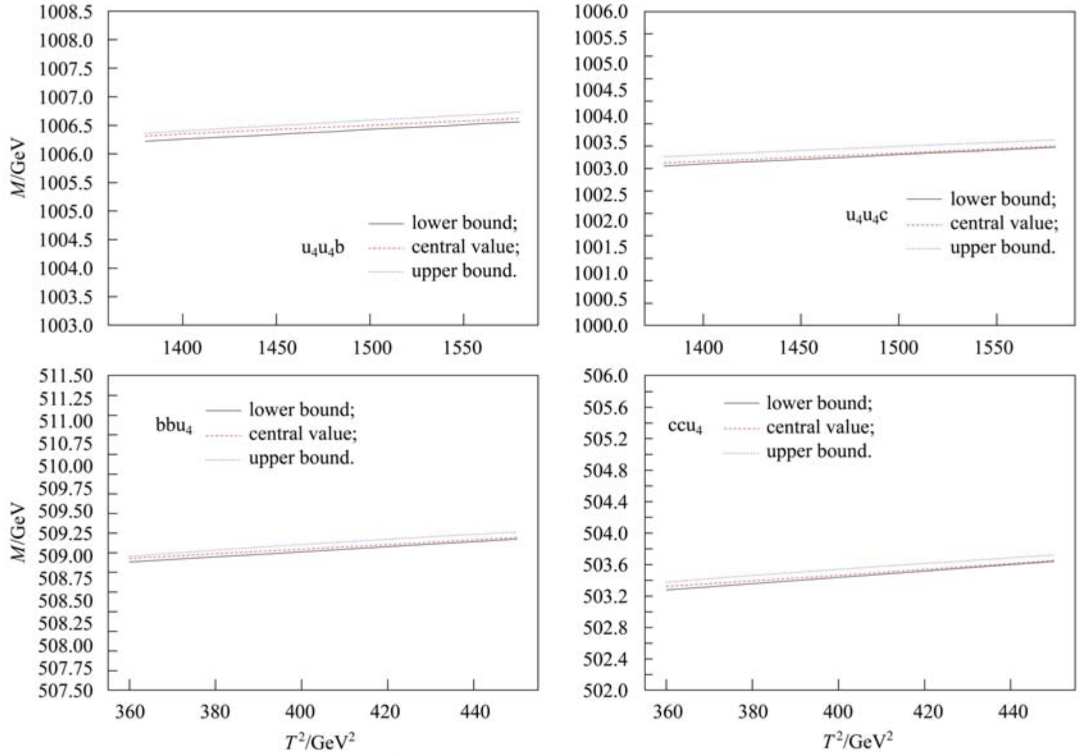

Fig. 2. The masses, M, of the new $\frac{1}{2}^{+}$ baryon states with variations in the Borel parameters, T^2. The threshold parameters, s_0, are taken to be the central values, and the masses of the fourth-generation quark u_4 are taken to be the central values, 500 GeV.

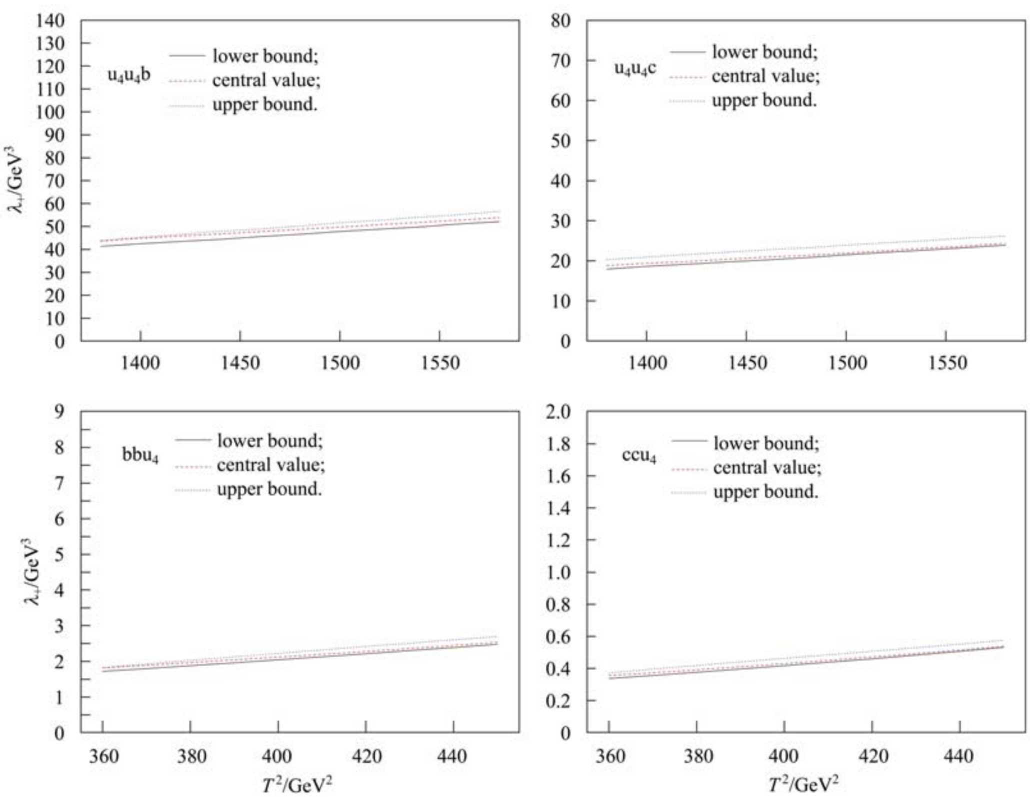

Fig. 3. The pole residues, λ_+, of the new $\frac{1}{2}^{+}$ baryon states with variations in the Borel parameters, T^2. The threshold parameters, s_0, are taken to be the central values, and the masses of the fourth-generation quark u_4 are taken to be the central values, 500 GeV.

Table 2. The pole residues of the new $\frac{1}{2}^+$ baryon states with different masses of the fourth-generation quark.

λ_+/GeV^3	m_{u_4}=450 GeV	m_{u_4}=500 GeV	m_{u_4}=550 GeV
$\mathrm{u_4u_4b}$	44.173±5.832	48.921±7.639	55.458±9.221
$\mathrm{u_4u_4c}$	19.752±3.172	22.049±4.008	28.384±4.898
$\mathrm{bbu_4}$	1.838±0.368	2.205±0.487	2.735±0.604
$\mathrm{ccu_4}$	0.323±0.088	0.456±0.149	0.589±0.217

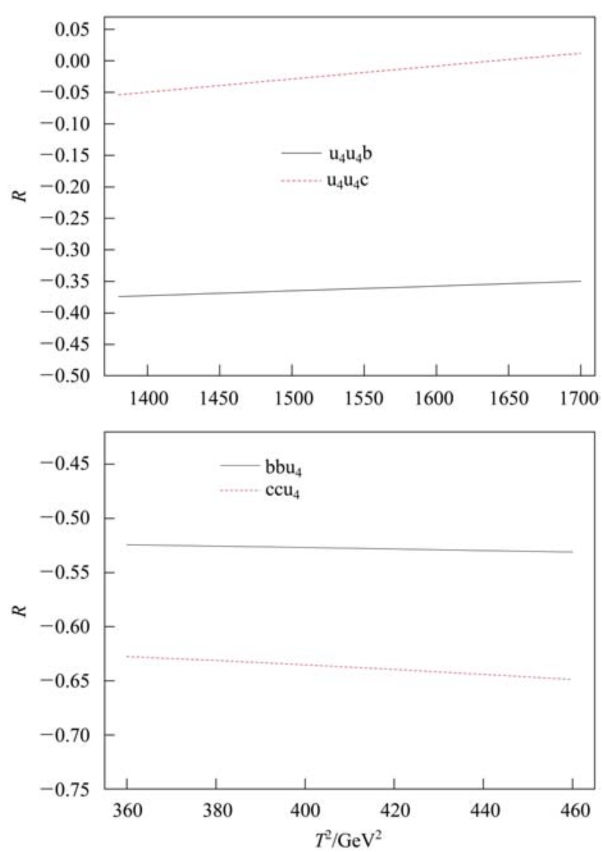

Fig. 4. The ratios between the contributions from the negative parity and positive parity baryon states with variations in the Borel parameters, T^2. The threshold parameters, s_0, are taken to be the central values, and the masses of the fourth-generation quark, u_4, are taken to be the central values, 500 GeV.

The rations

$$R=\frac{\int_{\Delta}^{\sqrt{s_0}}\mathrm{d}p_0[\rho^A(p_0)-\rho^B(p_0)]\exp\left[-\frac{p_0^2}{T^2}\right]}{\int_{\Delta}^{\sqrt{s_0}}\mathrm{d}p_0[\rho^A(p_0)+\rho^B(p_0)]\exp\left[-\frac{p_0^2}{T^2}\right]} \tag{12}$$

between the contributions from the negative parity and positive parity baryon states are shown explicitly in Fig. 4, with the mass of the fourth-family quark, 500 GeV. At the T^2 = (1380–1700) GeV^2 value, $R=-34\%$–(-37%) in the $\mathrm{u_4u_4b}$ channel and $R=-5\%$–1% in the $\mathrm{u_4u_4c}$ channel. At the T^2 =(360–460) GeV^2 value, $R=-52\%$–(-53%) in the $\mathrm{bbu_4}$ channel and $R=-62\%$–(-64%) in the $\mathrm{ccu_4}$ channel. Thus, we can see that without separating the contributions of the positive-parity baryon states from the negative-parity baryon states explicitly, the two criteria for choosing the Borel parameter T^2 and threshold parameter s_0 in the standard QCD sum rules do not work efficiently; we may choose the Borel windows where the contaminations from the negative-parity baryon states are large. If we choose the tensor structure γ_0+1, the contaminations from the negative-parity baryon states are excluded explicitly.

The authors in Ref. [23] discussed the two-loop renormalization group equations for the Higgs quartic and Yukawa couplings in the Standard Model with the fourth-generation quarks, and showed a quasi fixed point structure. This quasi fixed point behavior indicates a possible restoration of scale symmetry above some physical cut-off scale Λ_{FP}, situated in the range of a few TeV to the order of 100 TeV, around which the authors found that strong Yukawa couplings make it possible for the fourth-generation quarks to form bound states. But their discussions are based on the conjecture that there exists a true fixed point which is reached at a similar energy scale. Our bound states are from QCD, not the electroweak sector, and their masses are small enough than Λ_{FP}. It would be possible that bound states exist at these two energy regions. Meanwhile, we should stress that our results are within QCD and do not include the contributions from Higgs couplings to the fourth-family quarks. In Ref. [9], the authors calculated the contributions to the binding energy in the ultraheavy meson sector from the Higgs couplings to the fourth-family quarks, and found that they are proportional to the product of two quark masses. The binding energy obtained in Ref. [18] is very small in comparison with the Higgs corrections in Ref. [9]. But on the other hand, the corresponding analyses in the ultraheavy baryon sector are insufficient and deserve further study.

The $\frac{1}{2}^{+}$ baryon states containing the quarks of the fourth generation and the SM charm or bottom quark are simple extensions of the ordinary $\frac{1}{2}^{+}$ doubly heavy baryon states. Hopefully, all the obtained results of the masses and pole residues will be tested by future Tevatron or LHC experiments. In fact, a process has been proposed to analyze the search potential to discover the hadrons containing the fourth-family quarks, such as search $p\bar{p} \rightarrow u_4\bar{u}_4 \rightarrow \gamma qg\bar{q}$ [5]. As for the $\frac{1}{2}^{+}$ baryon sector, we propose the process $p\bar{p} \rightarrow BX$ for the experimental search.

4 Conclusions

In this work, we studied the the $\frac{1}{2}^{+}$ baryon states containing fourth-generation quarks and the SM charm or bottom quark by separating the contributions from the corresponding $\frac{1}{2}^{-}$ baryon states within the framework of the QCD sum rules, and obtained their masses and pole residues. These new baryon states are similar to the ordinary $\frac{1}{2}^{+}$ double heavy baryon states in the quark structure, and our predicted results may be tested in future Tevatron or LHC experiments.

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