ANALYSIS OF THE $\frac{1}{2}^{\pm}$ FLAVOR ANTITRIPLET HEAVY BARYON STATES WITH QCD SUM RULES

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Abstract

In this article, we study the masses and pole residues of the $\frac{1}{2}^{\pm}$ flavor antitriplet heavy baryon states ($\Lambda^+_c$, $\Xi^+_c$, $\Xi^0_c$) and ($\Lambda^0_b$, $\Xi^0_b$, $\Xi^-_b$) by subtracting the contributions from the corresponding $\frac{1}{2}^{\pm}$ heavy baryon states with the QCD sum rules, and observe the masses are in good agreement with the experimental data and make reasonable predictions for the unobserved $\frac{1}{2}^{-}$ bottom baryon states. Once reasonable values of the pole residues $\lambda_{\Lambda}$ and $\lambda_{\Xi}$ are obtained, we can take them as basic parameters to study the relevant hadronic processes with the QCD sum rules.

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1 Introduction

In recent years, several new excited charmed baryon states have been observed by the BaBar, Belle and CLEO Collaborations, such as the $\Lambda_c(2765)^+$, $\Lambda_c^+(2880)$, $\Lambda_c^+(2940)$, $\Sigma_c^+(2800)$, $\Xi_c^+(2980)$, $\Xi_c^+(3077)$, $\Xi_c^0(2980)$, $\Xi_c^0(3077)$ [1-3], and re-vivified the interest in the charmed baryon spectrum. On the other hand, the QCD sum rules is a powerful theoretical tool in studying the ground state heavy baryon states [4, 5]. In the QCD sum rules, the operator product expansion is used to expand the time-ordered currents into a series of quark and gluon condensates which parameterize the long distance properties of the QCD vacuum. Based on the quark-hadron duality, we can obtain copious information about the hadronic parameters at the phenomenological side [4, 5]. There have been several works on the masses of the heavy baryon states with the full QCD sum rules and the QCD sum rules in the heavy quark effective theory [6-15, 16-21]. In Refs. [22, 23, 24], we study the $\frac{1}{2}^{+}$ heavy baryon states $\Omega_Q$, $\Xi'_Q$ and $\Sigma_Q$ and $\frac{3}{2}^{+}$ heavy baryon states $\Omega'^*_Q$, $\Xi'^*_Q$ and $\Sigma'^*_Q$ with the full QCD sum rules, and observe that the pole residues of the $\frac{3}{2}^{+}$ heavy baryon states from the sum rules with different tensor structures are consistent with each other, while the pole residues of the $\frac{1}{2}^{+}$ heavy baryon states from the sum rules with different tensor structures differ from each other greatly. Those pole residues are important parameters in studying the radiative decays $\Omega_Q \to \Omega_Q \gamma$, $\Xi'_Q \to \Xi'_Q \gamma$ and $\Sigma'^*_Q \to \Sigma'^*_Q \gamma$ [24, 25], we should refine those parameters to improve the predictive ability.

In Ref. [26], Jido et al introduce a novel approach based on the QCD sum rules to separate the contributions of the negative-parity light flavor baryon states from the positive-parity light flavor baryon states, as the interpolating currents may have non-vanishing couplings to both the negative- and positive-parity baryon states [27]. Before the work of...
Jido et al, Bagan et al take the infinite mass limit for the heavy quarks to separate the contributions of the positive and negative parity heavy baryon states unambiguously [6].

In Ref.[28], we follow Ref.[26] and re-study the masses and pole residues of the \( \frac{1}{2}^+ \) flavor sextet heavy baryon states \( \Omega_Q, \Xi_Q^0 \) and \( \Sigma_Q \) by subtracting the contributions of the negative parity heavy baryon states with the full QCD sum rules. In this article, we use the same approach to study the \( \frac{1}{2}^± \) flavor antitriplet heavy baryon states \( (\Lambda_0^+, \Xi_0^+, \Xi_c^0) \) and \( (\Lambda_0^0, \Xi_0^0, \Xi_0^−) \).

The article is arranged as follows: we derive the QCD sum rules for the masses and the pole residues of the heavy baryon states \( (\Lambda_0^+, \Xi_0^+, \Xi_c^0) \) and \( (\Lambda_0^0, \Xi_0^0, \Xi_0^−) \) in Sect.2; in Sect.3, we present the numerical results and discussions; and Sect.4 is reserved for our conclusions.

2 QCD sum rules for the \( \Lambda_Q \) and \( \Xi_Q \)

The \( \frac{1}{2}^± \) flavor antitriplet heavy baryon states \( (\Lambda_0^+, \Xi_0^+, \Xi_c^0) \) and \( (\Lambda_0^0, \Xi_0^0, \Xi_0^−) \) can be interpolated by the following currents \( J_\Lambda(x) \) and \( J_\Xi(x) \) respectively,

\[
J_\Lambda(x) = e^{ijk} u_i^T(x) C \gamma_5 d_j(x) Q_k(x), \\
J_\Xi(x) = e^{ijk} \bar{u}_i^T(x) C \gamma_5 s_j(x) Q_k(x),
\]

(1)

where the \( Q \) represents the heavy quarks \( c \) and \( b \), the \( i, j \) and \( k \) are color indexes, and the \( C \) is the charge conjunction matrix.

The corresponding negative-parity heavy baryon states can be interpolated by the currents \( J_− = i\gamma_5 J_+ \) because multiplying \( i\gamma_5 \) to the \( J_+ \) changes the parity of the \( J_+ \) [26], where the \( J_+ \) denotes the currents \( J_\Lambda(x) \) and \( J_\Xi(x) \). The correlation functions are defined by

\[
\Pi_\pm(p) = i \int d^4xe^{ip·x} \langle 0 | T \{ J_\pm(x) J_\pm(0) \} | 0 \rangle,
\]

(2)

and can be decomposed as

\[
\Pi_\pm(p) = p\Pi_1(p) ± \Pi_0(p),
\]

(3)

due to Lorentz covariance. The currents \( J_\pm \) couple to both the positive- and negative-parity baryon states [27], i.e. \( \langle 0 | J_\pm | B^- \rangle \langle B^- | J_\pm | 0 \rangle = −\gamma_5 \langle 0 | J_\mp | B^- \rangle \langle B^- | J_\mp | 0 \rangle \gamma_5 \), where the \( B^- \) denote the negative parity baryon states.

We insert a complete set of intermediate baryon states with the same quantum numbers as the current operators \( J_+(x) \) and \( J_-(x) \) into the correlation functions \( \Pi_\pm(p) \) to obtain the hadronic representation [4][5]. After isolating the pole terms of the lowest states, we obtain the following result [26]:

\[
\Pi_\pm(p) = \lambda^- \frac{p_\perp^2 + M_+}{M_+^2 - p^2} + \lambda^+ \frac{p_\perp^2 - M_-}{M_-^2 - p^2} + \cdots,
\]

(4)

where the \( M_\pm \) are the masses of the lowest states with parity \( \pm \) respectively, and the \( \lambda_\pm \) are the corresponding pole residues (or couplings). If we take \( p_\perp = 0 \), then

\[
\text{limit}_{\epsilon \to 0} \frac{\text{Im}\Pi_\pm(p_0 + i\epsilon)}{\pi} = \lambda^+ \frac{\gamma_0 + 1}{2} \delta(p_0 - M_+) + \lambda^- \frac{\gamma_0 - 1}{2} \delta(p_0 - M_-) + \cdots
\]

\[
= \gamma_0 A(p_0) ± B(p_0) + \cdots,
\]

(5)
where

\[
A(p_0) = \frac{1}{2} \left[ \lambda_+^2 \delta(p_0 - M_+) + \lambda_-^2 \delta(p_0 - M_-) \right],
\]

\[
B(p_0) = \pm \frac{1}{2} \left[ \lambda_+^2 \delta(p_0 - M_+) - \lambda_-^2 \delta(p_0 - M_-) \right],
\]

(6)

de the contribution \(A(p_0) + B(p_0)\) \((A(p_0) - B(p_0))\) contains contributions from the positive parity (negative parity) states only for the \(\Pi^+\) \((\Pi^-)\) in space and use the momentum space expression for the heavy quark propagators, then

\[
\exp \left[ -\frac{p_0^2}{T^2} \right], \quad p_0 \exp \left[ -\frac{p_0^2}{T^2} \right],
\]

and obtain the following sum rules,

\[
\lambda_{\pm}^2 \exp \left[ -\frac{M_{\pm}^2}{T^2} \right] = \int_\Delta^{\sqrt{s_0}} dp_0 \left[ \rho^A(p_0) + \rho^B(p_0) \right] \exp \left[ -\frac{p_0^2}{T^2} \right],
\]

(7)

\[
\lambda_{\pm}^2 M_{\pm}^2 \exp \left[ -\frac{M_{\pm}^2}{T^2} \right] = \int_\Delta^{\sqrt{s_0}} dp_0 \left[ \rho^A(p_0) + \rho^B(p_0) \right] p_0^2 \exp \left[ -\frac{p_0^2}{T^2} \right],
\]

(8)

where

\[
\rho^A_{\pi}(p_0) = \frac{3p_0}{128\pi^4} \int_{t_i}^1 dt (1-t)^2 (p_0^2 - \bar{m}_Q^2)^2 + \frac{p_0 m_s (\bar{s}s - 2\bar{q}q)}{16\pi^2} \int_{t_i}^1 dt t
\]

\[
+ \frac{p_0}{128\pi^2} \left[ \frac{\alpha_s GG}{\pi} \right] \int_{t_i}^1 dt - \frac{m_Q^2}{768\pi^2} \left[ \frac{\alpha_s GG}{\pi} \right] \int_{t_i}^1 dt \left( \frac{1-t}{t^2} \right)^2 \delta(p_0 - \bar{m}_Q)
\]

\[
+ \frac{m_s [3\bar{q}g_s \sigma Gq] - [\bar{s}g_s \sigma Gs]}{192\pi^2} \delta(p_0 - m_Q) + \frac{\langle \bar{q}q \rangle \langle \bar{s}s \rangle}{12} \delta(p_0 - m_Q),
\]

(9)

\[
\rho^B_{\pi}(p_0) = \frac{3m_Q}{128\pi^4} \int_{t_i}^1 dt (1-t)^2 (p_0^2 - \bar{m}_Q^2)^2 + \frac{m_Q m_s (\bar{s}s - 2\bar{q}q)}{16\pi^2} \int_{t_i}^1 dt t
\]

\[
+ \frac{m_Q}{128\pi^2} \left[ \frac{\alpha_s GG}{\pi} \right] \int_{t_i}^1 dt + \frac{m_Q}{768\pi^2} \left[ \frac{\alpha_s GG}{\pi} \right] \int_{t_i}^1 dt \left( \frac{1-t}{t^2} \right)^3
\]

\[
- \frac{m_Q}{768\pi^2} \left[ \frac{\alpha_s GG}{\pi} \right] \int_{t_i}^1 dt \left( \frac{1-t}{t} \right)^2 \bar{m}_Q \delta(p_0 - \bar{m}_Q)
\]

\[
+ \frac{m_s [3\bar{q}g_s \sigma Gq] - [\bar{s}g_s \sigma Gs]}{192\pi^2} \delta(p_0 - m_Q) + \frac{\langle \bar{q}q \rangle \langle \bar{s}s \rangle}{12} \delta(p_0 - m_Q),
\]

(10)

\[
\rho^A_{Q}(p_0) = \frac{3p_0}{128\pi^4} \int_{t_i}^1 dt (1-t)^2 (p_0^2 - \bar{m}_Q^2)^2 + \frac{\langle \bar{q}q \rangle^2}{12} \delta(p_0 - m_Q)
\]

\[
+ \frac{p_0}{128\pi^2} \left[ \frac{\alpha_s GG}{\pi} \right] \int_{t_i}^1 dt - \frac{m_Q^2}{768\pi^2} \left[ \frac{\alpha_s GG}{\pi} \right] \int_{t_i}^1 dt \left( \frac{1-t}{t^2} \right)^2 \delta(p_0 - \bar{m}_Q),
\]

(11)
\[ p_{\Delta Q}^B(p_0) = \frac{3m_Q}{128\pi^4} \int_{t_i}^1 dt (1-t)^2 (p_0^2 - \tilde{m}_Q^2)^2 + \frac{\langle \bar{q}q \rangle^2}{12} \delta(p_0 - m_Q) \\
+ \frac{m_Q}{128\pi^2} \left( \frac{\alpha_s GG}{\pi} \right) \int_{t_i}^1 dt + \frac{m_Q}{192\pi^2} \left( \frac{\alpha_s GG}{\pi} \right) \int_{t_i}^1 dt \frac{(1-t)^3}{t^2} \\
- \frac{m_Q}{768\pi^2} \left( \frac{\alpha_s GG}{\pi} \right) \int_{t_i}^1 dt \frac{(1-t)^2}{t} \tilde{m}_Q \delta(p_0 - \tilde{m}_Q), \]  

where \( \tilde{m}_Q^2 = \frac{m_Q^2}{t_i^2} \), \( t_i = \frac{m_Q^2}{p_0^2} \), the \( s_0 \) are the threshold parameters, \( T^2 \) is the Borel parameter, \( \Delta = m_Q + m_s \) and \( m_Q \) in the channels \( \Xi_Q \) and \( \Lambda_Q \) respectively.

3 Numerical results and discussions

The input parameters are taken to be the standard values \( \langle \bar{q}q \rangle = -(0.24 \pm 0.01 \text{GeV})^3 \), \( \langle \bar{s}s \rangle = (0.8 \pm 0.2) \langle \bar{q}q \rangle \), \( \langle \bar{q}g_s\sigma Gq \rangle = m_0^2 \langle \bar{q}q \rangle \), \( \langle \bar{s}g_s\sigma Gs \rangle = m_0^2 \langle \bar{s}s \rangle \), \( m_0^2 = (0.8 \pm 0.2) \text{GeV}^2 \) \( \{29, 30\}, \) \( \langle \alpha_s GG \rangle = (0.012 \pm 0.004) \text{GeV}^4 \) \( \{30\} \), \( m_s = (0.14 \pm 0.01) \text{GeV} \), \( m_c = (1.35 \pm 0.10) \text{GeV} \) and \( m_b = (4.7 \pm 0.1) \text{GeV} \) \( \{31\} \) at the energy scale \( \mu = 1 \text{GeV} \).

The value of the gluon condensate \( \langle \alpha_s GG \rangle \) has been updated from time to time, and changes greatly \( \{32\} \). At the present case, the gluon condensate makes tiny contribution, the updated value \( \langle \alpha_s GG \rangle = (0.023 \pm 0.003) \text{GeV}^4 \) \( \{32\} \) and the standard value \( \langle \alpha_s GG \rangle = (0.012 \pm 0.004) \text{GeV}^4 \) \( \{30\} \) lead to a difference less than 15 MeV for the masses.

In the conventional QCD sum rules \( \{31, 35\} \), there are two criteria (pole dominance and convergence of the operator product expansion) for choosing the Borel parameter \( T^2 \) and threshold parameter \( s_0 \). We impose the two criteria on the heavy baryon states to choose the Borel parameter \( T^2 \) and threshold parameter \( s_0 \), the values are shown in Table 1. From Table 1, we can see that the contribution from the perturbative term is dominant, the operator product expansion is convergent certainly. In this article, we take the contribution from the pole term is larger than 45% (49%) for the positive (negative) parity baryon states, the uncertainty of the threshold parameter is 0.1 GeV, and the Borel window is 1 GeV$^2$.

In calculation, we neglect the contributions from the perturbative corrections. Those perturbative corrections can be taken into account in the leading logarithmic approximations through anomalous dimension factors. After the Borel transform, the effects of those corrections are to multiply each term on the operator product expansion side by the factor, \( \left[ \frac{\alpha_s(T\mu)}{\alpha_s(\mu')} \right]^{2T_j - \Gamma_{\mathcal{O}_n}} \), where the \( \Gamma_j \) is the anomalous dimension of the interpolating current \( J(x) \) and the \( \Gamma_{\mathcal{O}_n} \) is the anomalous dimension of the local operator \( \mathcal{O}_n(0) \). We carry out the operator product expansion at a special energy scale \( \mu^2 = 1 \text{GeV}^2 \), and set the factor \( \left[ \frac{\alpha_s(T\mu)}{\alpha_s(\mu')} \right]^{2T_j - \Gamma_{\mathcal{O}_n}} \approx 1 \), such an approximation maybe result in some scale dependence and weaken the prediction ability. In this article, we study the flavor antitriplet \( \frac{1}{2}^\pm \) heavy baryon states systemically, the predictions are still robust as we take the analogous criteria in those sum rules.

Taking into account all uncertainties of the relevant parameters, we obtain the values of the masses and pole residues of the flavor antitriplet \( \frac{1}{2}^\pm \) heavy baryon states \( (\Lambda_c^+, \Xi_c^+, \Xi_c^0) \).
and \((\Lambda_b^0, \Xi_b^0, \bar{\Xi}_b^-)\), which are shown in Figs.1-2 and Table 2. In this article, we calculate the uncertainties \(\delta\) with the formula

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

(13)

where the \(f\) denote the hadron mass \(M\) and the pole residue \(\lambda\), the \(x_i\) denote the input QCD parameters \(m_c, m_b, \langle \bar{q}q \rangle, \langle \bar{s}s \rangle, \ldots\), and the threshold parameter \(s_0\) and Borel parameter \(T^2\). As the partial derivatives \(\frac{\partial f}{\partial x_i}\) are difficult to carry out analytically, we take the approximation \(\left( \frac{\partial f}{\partial x_i} \right)^2 (x_i - \bar{x}_i)^2 \approx [f(\bar{x}_i + \Delta x_i) - f(\bar{x}_i)]^2\) in the numerical calculations.

In Ref.[7], Bagan et al study the masses and couplings (pole residues) of the heavy baryon states \(\Sigma_b^*, \Sigma_c^*, \Lambda_b\) and \(\Lambda_c\) with the full QCD sum rules, and observe that the \(\Sigma_b^*(c)\) and \(\Sigma_c^*(c)\) have degenerate masses within the uncertainties, while the \(\Lambda_b(c)\) is lighter than the \(\Sigma_b^*(c)\). In Refs.[10] [11], Zhang et al perform a systematic study of the masses of charmed and bottom baryon states with the full QCD sum rules, where the vacuum

| \(\Lambda_c\left(\frac{1}{2}^+\right)\) | \(T^2\) (GeV\(^2\)) | \(\sqrt{s_0}\) (GeV) | \(M\) (GeV) | \(\lambda\) (GeV\(^3\)) | \(M\) (GeV) [exp] |
|---|---|---|---|---|---|
| \(\Lambda_c\left(\frac{1}{2}^+\right)\) | 1.7 - 2.7 | 3.1 | 2.26 ± 0.27 | 0.022 ± 0.008 | 2.28646 |
| \(\Xi_c\left(\frac{1}{2}^+\right)\) | 1.9 - 2.9 | 3.2 | 2.44 ± 0.23 | 0.027 ± 0.008 | 2.4678 / 2.47088 |
| \(\Lambda_b\left(\frac{1}{2}^+\right)\) | 4.3 - 5.3 | 6.5 | 5.65 ± 0.20 | 0.030 ± 0.009 | 5.6202 |
| \(\Xi_b\left(\frac{1}{2}^+\right)\) | 4.4 - 5.4 | 6.5 | 5.73 ± 0.18 | 0.032 ± 0.009 | 5.7924 |
| \(\Lambda_c\left(\frac{1}{2}^-\right)\) | 2.2 - 3.2 | 3.4 | 2.61 ± 0.21 | 0.035 ± 0.009 | 2.5954 |
| \(\Xi_c\left(\frac{1}{2}^-\right)\) | 2.4 - 3.4 | 3.5 | 2.76 ± 0.18 | 0.042 ± 0.009 | 2.7891 / 2.7918 |
| \(\Lambda_b\left(\frac{1}{2}^-\right)\) | 4.7 - 5.7 | 6.7 | 5.85 ± 0.18 | 0.042 ± 0.012 | ? |
| \(\Xi_b\left(\frac{1}{2}^-\right)\) | 5.0 - 6.0 | 6.8 | 6.01 ± 0.16 | 0.051 ± 0.012 | ? |
Figure 1: The masses $M$ of the heavy baryon states, the $A$, $B$, $C$ and $D$ correspond to the $\frac{1}{2}^+$ heavy baryon channels $\Lambda_c$, $\Xi_c$, $\Lambda_b$ and $\Xi_b$ respectively, while the $E$, $F$, $G$ and $H$ correspond to the $\frac{1}{2}^-$ heavy baryon channels $\Lambda_c$, $\Xi_c$, $\Lambda_b$ and $\Xi_b$ respectively.
Figure 2: The pole residues $\lambda$ of the heavy baryon states, the $A$, $B$, $C$ and $D$ correspond to the $\frac{1}{2}^+$ heavy baryon channels $\Lambda_c$, $\Xi_c$, $\Lambda_b$ and $\Xi_b$ respectively, while the $E$, $F$, $G$ and $H$ correspond to the $\frac{1}{2}^-$ heavy baryon channels $\Lambda_c$, $\Xi_c$, $\Lambda_b$ and $\Xi_b$ respectively.
This work that the pole residues from the sum rules with different tensor structures (the negative parity heavy baryon states, and find that the predictions for the masses and energy \( \bar{\Lambda} \) and the study only the masses of the heavy baryon states \( \Xi_{c} \). In Ref.\[21\], Liu et al perform a systematic study of the masses of the bottom baryon states up to \( \frac{1}{m_Q} \) in the heavy quark effective field theory using the QCD sum rules.

In the heavy quark limit, \( M_Q = m_Q + \bar{\Lambda} + O\left(\frac{1}{m_Q}\right) \). We can calculate the bound energy \( \bar{\Lambda} \) and the \( \frac{1}{m_Q} \) corrections of the lowest heavy baryon states \( \Lambda_Q \) with the QCD sum rules, compare them with the experimental data \( M_Q \), and determine the heavy quark masses \( m_Q \), which always suffer from large uncertainties, then use the \( m_Q \) as basic input parameters to calculate other heavy baryon masses, one can consult Refs.\[15, 16, 17\] for example. In Table 3, we also present the predictions from the full QCD sum rules and the QCD sum rules in the heavy quark effective theory (where systematic studies are preformed) \[9, 11, 21\].

From Tables 2-3, we can see that the present and other theoretical predictions are all in good agreement with the experimental data for the positive parity baryon states. The negative parity baryon states \( \Lambda_b \) and \( \Xi_b \) are not observed yet, we make reasonable predictions for their masses to confront with experimental data in the future at the LHCb \[33\].

The fractions

\[
R = \frac{\int_{\Delta}^{s_0} dp_0 \left[ \rho^A(p_0) - \rho^B(p_0) \right] \exp \left[ -\frac{m^2}{\bar{\Lambda}^2} \right]}{\int_{\Delta}^{s_0} dp_0 \left[ \rho^A(p_0) + \rho^B(p_0) \right] \exp \left[ -\frac{m^2}{\bar{\Lambda}^2} \right]}
\]

are less than 4.5% (6.0%) and 0.6% (0.8%) in the positive (negative) parity charmed baryon and bottom baryon channels, respectively. So the contaminations from the negative (or positive) parity baryon states are very small. In Refs.\[23, 24\], we study the \( \frac{1}{2}^+ \) flavor sextet heavy baryon states \( \Omega_Q, \Xi'_Q \) and \( \Sigma_Q \) with the full QCD sum rules, and observe that the pole residues from the sum rules with different tensor structures (\( \rho \) and 1) differ from each other greatly. In Ref.\[28\], we re-study the masses and pole residues of the \( \frac{1}{2}^+ \) flavor sextet heavy baryon states \( \Omega_Q, \Xi'_Q \) and \( \Sigma_Q \) by subtracting the contributions of the negative parity heavy baryon states, and find that the predictions for the masses and

| \( \Lambda_c(\frac{1}{2}^+) \) | [9] | [11] | [21] | This work |
|-----------------|-----|-----|-----|---------|
| \( \Xi_c(\frac{1}{2}^+) \) | 2.31 ± 0.19 | 2.271 ± 0.067 | 2.26 ± 0.27 |
| \( \Delta(\frac{1}{2}^+) \) | 2.5 ± 0.2 | 2.48 ± 0.21 | 2.44 ± 0.23 |
| \( \Xi_b(\frac{1}{2}^+) \) | 5.69 ± 0.13 | 5.637 ± 0.068 | 5.65 ± 0.20 |
| \( \Delta_b(\frac{1}{2}^+) \) | 5.75 ± 0.25 | 5.75 ± 0.13 | 5.73 ± 0.18 |
| \( \Lambda_c(\frac{1}{2}^-) \) | 2.53 ± 0.22 | 2.61 ± 0.21 |
| \( \Xi_c(\frac{1}{2}^-) \) | 2.65 ± 0.27 | 2.76 ± 0.18 |
| \( \Delta_b(\frac{1}{2}^-) \) | 5.84 ± 0.15 | 5.85 ± 0.18 |
| \( \Xi_b(\frac{1}{2}^-) \) | 5.95 ± 0.16 | 6.01 ± 0.16 |

Table 3: The masses \( M(\text{GeV}) \) of the heavy baryon states from the QCD sum rules.
pole residues are improved considerably. In the present case, we can choose the tensor structures \( \hat{\rho} \) or 1 or \( \gamma_0 + 1 \) freely to study the masses and pole residues.

The pole residues of the \( \frac{1}{2}^\pm \) and \( \frac{3}{2}^+ \) heavy baryon sextets \( B_6 \) and \( B_6^* \) have been calculated in our previous works \[22, 24, 28, 36\]. Once reasonable values of the pole residues \( \lambda_\Lambda \) and \( \lambda_\Xi \) of the \( \frac{1}{2}^+ \) heavy baryon antitriplet \( \Lambda_\bar{3} \) are obtained, we can take them as basic input parameters and study the strong decays \( \Sigma_0^* \rightarrow \Lambda_Q \pi \), \( \Sigma_0 \rightarrow \Lambda_Q \pi \) and \( \Xi_0^* \rightarrow \Xi_Q \pi \) and the radiative decays \( B_6^* \rightarrow B_3 \gamma \) and \( B_6 \rightarrow B_3 \gamma \) in a systematic ways with the lightcone QCD sum rules or the QCD sum rules in external field, and confront the predictions with the experimental data in the future at the BESIII, PANDA and LHCb \[33, 34, 35\].

The strong decays \( \Sigma_c^*(2520) \rightarrow \Lambda_c \pi \), \( \Sigma_c(2455) \rightarrow \Lambda_c \pi \) saturate approximately the widths of the \( \Sigma_c^*(2520) \) and \( \Sigma_c(2455) \) respectively, while the strong decays \( \Xi_c^*(2645) \rightarrow \Xi_c \pi \) are seen \[31\]. From our previous works \[28, 36\], we can see that the corresponding strong decays \( \Sigma_b^* \rightarrow \Lambda_b \pi \), \( \Sigma_b \rightarrow \Lambda_b \pi \) and \( \Xi_b^* \rightarrow \Xi_b \pi \) are kinematically allowed; although the bottom baryon states \( \Xi_b^* \) have not been observed experimentally yet.

In Refs.\[24, 25\], we perform systematic studies for the radiative decays \( B_6^* \rightarrow B_0 \gamma \) with the light-cone QCD sum rules as the strong decays \( B_6^* \rightarrow B_0 \pi \) are forbidden due to the unavailable phase space, while the radiative channels are not phase space suppressed and become relevant; although the electromagnetic strength is weaker than that of the strong interaction. The radiative decays \( B_6^* \rightarrow B_3 \gamma \) and \( B_6 \rightarrow B_3 \gamma \) are important processes in testing the heavy quark symmetry and the chiral symmetry, for example, the \( \Xi_c^* \) and \( \Omega_c^* \) are governed by the radiative decays.

4 Conclusion

In this article, we study the \( \frac{1}{2}^\pm \) flavor antitriplet heavy baryon states \( (\Lambda_\bar{3}^+, \Xi_c^+, \Xi_\bar{b}^0) \) and \( (\Lambda_\bar{3}^0, \Xi_b^0, \Xi_\bar{b}^-) \) by subtracting the contributions from the corresponding \( \frac{1}{2}^\pm \) heavy baryon states with the QCD sum rules, obtain the masses which are in good agreement with the experimental data and make reasonable predictions for the unobserved \( \frac{1}{2}^\pm \) bottom baryon states. In calculation, we observe that the contaminations from the negative (or positive) parity baryon states are very small, one can choose the tensor structures \( \hat{\rho} \) or 1 or \( \gamma_0 + 1 \) freely to study the masses and pole residues. Once reasonable values of the pole residues \( \lambda_\Lambda \) and \( \lambda_\Xi \) are obtained, we can take them as basic input parameters and study the strong decays and radiative decays in a systematic ways with the light-cone QCD sum rules or the QCD sum rules in external field, and confront the predictions with the experimental data in the future at the BESIII, PANDA and LHCb.

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