Excited heavy baryon masses to order $\Lambda_{\mathrm{QCD}}/m_Q$ from QCD sum rules

Dao-Wei Wang
Department of Applied Physics, National University of Defense Technology, Hunan 410073, China

Ming-Qiu Huang
CCAST (World Laboratory), P.O. Box 8730, Beijing 100080, China
and Department of Applied Physics, National University of Defense Technology, Hunan 410073, China

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The masses of the p-wave excited heavy baryons have been calculated to $\Lambda_{\mathrm{QCD}}/m_Q$ order using the QCD sum rule method within the framework of heavy quark effective theory. Numerical results for the kinetic energy λ_1 and chromomagnetic interaction λ_2 are presented. The splitting between spin 1/2 and 3/2 doublets derived from our calculation is given; the agreement with the current experiment is good.

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I. INTRODUCTION

In the past decade continuous progress has been made in the investigation of excited heavy baryons. The lowest lying orbitally excited charmed states $\Lambda_c(2593)$ and $\Lambda_c(2625)$ have been observed by several collaborations [1]; the excited states of Ξ_c have also been reported recently [2], and more data are expected in the near future. From the theoretical point of view those data need to be studied comprehensively. Furthermore, with the collection of more experimental data for excited heavy baryon states, it is useful to make some theoretical predictions about their spectroscopies.

Heavy baryons containing a single heavy quark can be described exactly by heavy quark effective theory (HQET) [3–5] in the heavy quark limit. This fact should contribute to the spin-flavor symmetry of the system comprised of infinitely heavy quarks. HQET has been applied successfully to learn about the properties of heavy mesons and baryons, including spectroscopy and weak decays. The mass formula for a spin symmetry doublet of heavy baryons up to order $1/m_Q$ corrections can be written as

$$M=m_Q+\bar{\Lambda}-\frac{1}{2m_Q}(\lambda_1+d_M\lambda_2), \tag{1}$$

where the parameter $\bar{\Lambda}$ is the effective mass of the light degrees of freedom in the $m_Q\to\infty$ limit, and λ_1 and λ_2 are related to the heavy quark kinetic energy and the chromomagnetic energy of HQET, respectively,

$$\lambda_1=\langle B(v)|\bar{h}_v(iD^{\perp})^2h_v|B(v)\rangle,$$

$$d_M\lambda_2=\left\langle B(v)\left|\bar{h}_v\sigma_{\mu\nu}\frac{g_s}{2}G^{\mu\nu}h_v\right|B(v)\right\rangle. \tag{2}$$

The constant d_M characterizes the spin-orbit interaction of the heavy quark and the gluon field; it is zero for the singlet, and $1,-\frac{1}{2}$ for the spin 1/2, spin 3/2 doublets, respectively. Thus the splitting of the spin 1/2 and 3/2 doublets is

$$M^2_{B^*_Q}-M^2_{B_Q}=\frac{3}{2}\lambda_2, \tag{3}$$

where B_Q and B^*_Q denote the spin 1/2 and 3/2 doublets, respectively.

The heavy baryon mass parameters λ_1 and λ_2 play a most significant role in our understanding of the spectroscopy [6] and inclusive decay rates [7]. They must be estimated in some nonperturbative approaches due to the asymptotic freedom property of QCD. A viable approach is the QCD sum rules [8] formulated in the framework of HQET [9]. This method allows us to relate the hadronic observables to QCD parameters via the operator product expansion (OPE) of the correlator. In the case of the ground state heavy baryon, predictions about the mass spectroscopy have been made to leading and next-to-leading order in α_s [10,11] and to order $1/m_Q$ [12,13] using the QCD sum rule method. As to excited baryon mass spectroscopy, only results to leading order in $1/m_Q$ expansion have been obtained from the QCD sum rules [14,15]. In [13] we calculated the baryonic parameters λ_1 and λ_2 for the ground state Λ_Q and Σ_Q baryons using the QCD sum rules in the HQET. Employing the baryonic currents from [14], we now derive these parameters for excited Λ- and Σ-type baryons following the same procedure.

The remainder of this paper is organized as follows. In Sec. II A we introduce the interpolating currents for excited state heavy baryons and briefly present the two-point sum rules. The direct Laplace sum rule analysis for the matrix elements is presented in Sec. II B. Section III is devoted to numerical results and our conclusions. Some comments are also available in Sec. III.

II. DERIVATION OF THE SUM RULES

A. Heavy baryonic currents and two-point sum rules

In this work we adopt the currents constructed from the Bethe-Salpeter equation in Ref. [14] as

$$j_{\Sigma_{Qk1}}=\epsilon_{abc}(q_1^{T\,a}\tau C\gamma_5 D_t^{\mu_1}q_2^b)\Gamma' h_v^c, \tag{4a}$$

$$j_{\Lambda_{Qk0}}=\epsilon_{abc}(q_1^{T\,a}\tau C\gamma_t^{\mu}D_{t\mu}q_2^b)\Gamma' h_v^c, \tag{4b}$$

$$j_{\Lambda_{Qk1}}=\epsilon_{abc}(q_1^{T\,a}\tau C\epsilon_{\mu_1\nu\sigma\rho}\gamma_t^{\nu}v^{\rho}D_t^{\sigma}q_2^b)\Gamma' h_v^c, \tag{4c}$$

$$j_{\Lambda_{QK1}}=\epsilon_{abc}(q_1^{T\,a}\tau C\gamma_5 q_2^b)\Gamma' D_t^{\mu_1}h_v^c, \tag{4d}$$

$$j_{\Sigma_{QK0}}=\epsilon_{abc}(q_1^{T\,a}\tau C\gamma_t^{\mu} q_2^b)\Gamma' D_{t\mu}h_v^c, \tag{4e}$$

$$j_{\Sigma_{QK1}}=\epsilon_{abc}(q_1^{T\,a}\tau C\epsilon_{\mu_1\nu\sigma\rho}\gamma_t^{\nu}v^{\rho} q_2^b)\Gamma' D_t^{\sigma}h_v^c, \tag{4f}$$

in which C is the charge conjugation matrix, τ is the flavor matrix, which is antisymmetric for Λ-type baryons and symmetric for Σ-type baryons, Γ' are some gamma matrices, and a,b,c denote the color indices. Γ' can be chosen covariantly as

$$\Gamma'=\gamma_{t\mu_1}\gamma_5 \tag{5}$$

for the $\Sigma_{Qk1},\Lambda_{Qk1},\Lambda_{QK1},\Sigma_{QK1}$ doublet spin 1/2 baryon, and

$$\Gamma'=\Gamma_{\mu_1\rho_1}=-\frac{1}{3}(g_{t\mu_1\rho_1}+\gamma_{t\mu_1}\gamma_{t\rho_1}) \tag{6}$$

for the $\Sigma_{Qk1},\Lambda_{Qk1},\Lambda_{QK1},\Sigma_{QK1}$ doublet spin 3/2 baryon, in which $g_{t\mu_1\rho_1}$ and $\gamma_{t\mu_1}$ are perpendicular to the heavy quark velocity v, defined as $g_{t\mu\nu}=g_{\mu\nu}-v_\mu v_\nu$, $\gamma_{t\mu}=\gamma_\mu-\not{v}v_\mu$. For the singlets Λ_{Qk0} and Σ_{QK0}, the Γ' is simply the unit matrix I. The notation used here to describe excited state heavy baryons is the same as that used in Refs. [5,14]: k denotes $l_k=1$ and $l_K=0$ whereas K denotes $l_k=0$ and $l_K=1$, where l_k is the orbital angular momentum describing the relative motion of the two light quarks and the orbital angular momentum l_K describes the orbital motion of the center of mass of the two light quarks relative to the heavy quark. Q denotes the heavy quark and the number in subscript is the total angular momentum of the light diquark system. As for ground state baryons, the flavor configuration of Λ-type baryons is antisymmetric and that of Σ-type baryons is symmetric. Depending on the number of derivatives or the form of Γ', the interpolating current can have forms different from those listed above and may also be used in applications [14,15].

In the following analysis we will use those currents to interpolate excited heavy baryon states. At the leading order of the $1/m_Q$ expansion, they do not mix with each other even with the same quantum number, but to the next-to-leading order mixing of interpolating currents will appear. In our subsequent calculations we use only those currents and do not consider the effect resulting from the mixing of interpolating currents, because Refs. [10,16,17] have shown that the stability criterion for QCD sum rule applications excludes the existence of interpolating currents mixing; even if mixing does exist, the numerical result will not change drastically compared with the case without mixing.

The baryonic coupling constants in HQET are also needed in our calculation. They are defined as follows:

$$\langle 0|j|B(v)\rangle=Fu, \tag{7}$$

where $|B(v)\rangle$ denotes the excited baryon state and u can be the ordinary spinor u or the Rarita-Schwinger spinor u_α in HQET, corresponding to the spin 1/2 or spin 3/2 doublet, respectively. Irrespective of an irrelevant constant factor in the leading order, the coupling constants for spin 1/2 and spin 3/2 doublets are equivalent because of their identical spin-parity of the light degrees of freedom.

In order to determine the effective mass of the excited baryons, we analyze the two-point correlator defined as

$$i\int d^4x e^{ik\cdot x}\langle 0|T\{j(x)\overline{j}(0)\}|0\rangle=\frac{1+\not{v}}{2}\mathrm{Tr}[\tau\tau^+]\Pi(\omega), \tag{8}$$

where k is the residual momentum and $\omega=v\cdot k$. For large negative values of ω, $\Pi(\omega)$ can be expressed in terms of perturbative and nonperturbative contributions. The nonperturbative effects can be accounted for by including quark and gluon condensates ordered by increasing dimension, giving a series of power corrections in the "small" $1/\omega$ variable. The Borel transformation in the variable ω can help to improve the convergence of these nonperturbative series.

With those interpolating currents listed in Eq. (4) it is straightforward to obtain the two-point sum rules:

$$F^2_{\Lambda_{Qk0}}e^{-\bar{\Lambda}_{\Lambda_{Qk0}}/T}=\frac{18\,T^8}{\pi^4}\delta_7(\omega_c/T)+\frac{T^4}{4\,\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{m_0^2}{16}\langle\overline{q}q\rangle^2 e^{-m_0^2/8T^2}, \tag{9a}$$

$$F^2_{\Sigma_{Qk1}}e^{-\bar{\Lambda}_{\Sigma_{Qk1}}/T}=\frac{90\,T^8}{\pi^4}\delta_7(\omega_c/T)-\frac{3T^4}{2^5\,\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{m_0^2}{16}\langle\overline{q}q\rangle^2 e^{-m_0^2/8T^2}, \tag{9b}$$

$$F^2_{\Lambda_{Qk1}}e^{-\bar{\Lambda}_{\Lambda_{Qk1}}/T}=\frac{216\,T^8}{\pi^4}\delta_7(\omega_c/T)-\frac{T^4}{4\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{m_0^2}{8}\langle\overline{q}q\rangle^2 e^{-m_0^2/8T^2}, \tag{9c}$$

$$F^2_{\Lambda_{QK1}}e^{-\bar{\Lambda}_{\Lambda_{QK1}}/T}=\frac{216\,T^8}{\pi^4}\delta_7(\omega_c/T)-\frac{T^4}{8\,\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{m_0^2}{4}\langle\overline{q}q\rangle^2 e^{-m_0^2/8T^2}, \tag{9d}$$

$$F^2_{\Sigma_{QK1}}e^{-\bar{\Lambda}_{\Sigma_{QK1}}/T}=\frac{288\,T^8}{\pi^4}\delta_7(\omega_c/T)-\frac{3T^4}{2\,\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{m_0^2}{2}\langle\overline{q}q\rangle^2 e^{-m_0^2/8T^2}, \tag{9e}$$

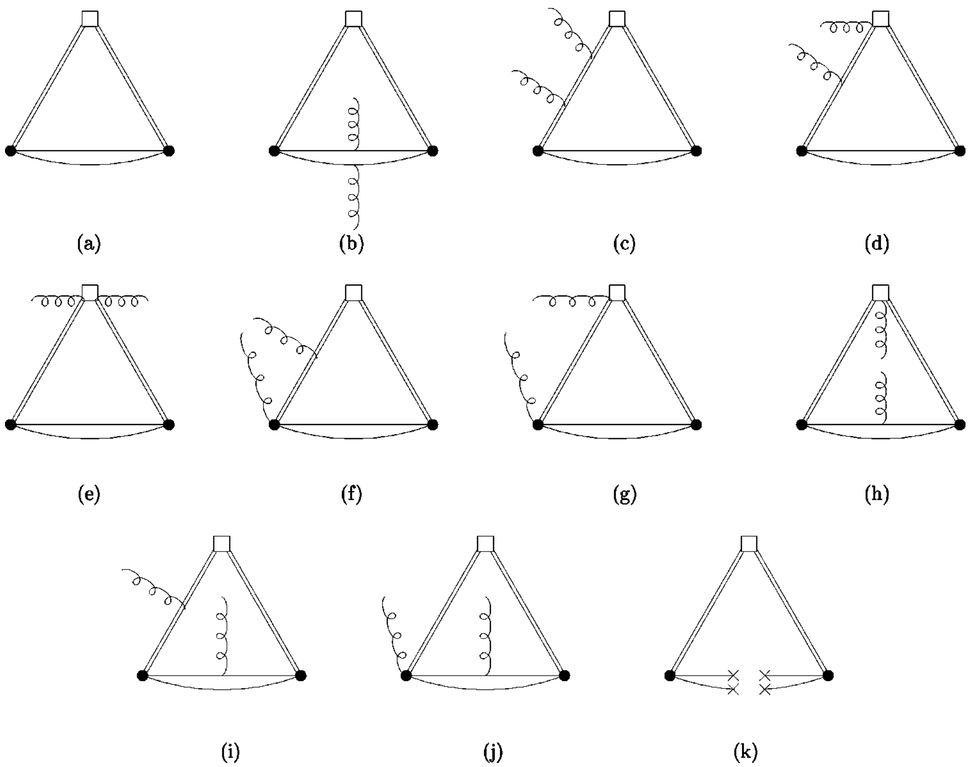

FIG. 1. Nonvanishing diagrams for the kinetic energy. The kinetic energy operator is denoted by a white square, the interpolating baryon currents by black circles. Heavy quark propagators are drawn as double lines.

$$F^2_{\Sigma_{QK0}} e^{-\bar{\Lambda}_{\Sigma_{QK0}}/T} = \frac{504\, T^8}{\pi^4}\, \delta_7(\omega_c/T) - \frac{T^4}{8\,\pi^2}\left\langle \frac{\alpha_s}{\pi} G^2 \right\rangle + \frac{m_0^2}{4}\langle \bar{q}q\rangle^2 e^{-m_0^2/8T^2}. \tag{9f}$$

In calculations we adopted the Gaussian ansatz for the nonlocal quark condensate to get the dimension 6 condensate contribution. Dimension $D>6$ condensates are not included. The functions $\delta_n(\omega_c/T)$ arise from the continuum subtraction and are defined in [13].

B. Sum rules for λ_1 and λ_2

For the evaluation of the matrix elements λ_1 and λ_2 we consider the three-point correlators as follows:

$$i^2 \int d^4x \int d^4y e^{ik\cdot x - ik'\cdot y} \langle 0 | T\{ j(x) \bar{h}_v (iD^{\perp})^2 h_v(0) \bar{j}(y)\} | 0\rangle = \frac{1+\not{v}}{2} \mathrm{Tr}[\tau\tau^+] T_K(\omega,\omega'),$$

$$i^2 \int d^4x \int d^4y e^{ik\cdot x - ik'\cdot y} \left\langle 0 \left| T\left\{ j(x) \bar{h}_v \sigma_{\mu\nu} \frac{g_s}{2} G^{\mu\nu} h_v(0) \bar{j}(y)\right\} \right| 0\right\rangle = d_M \frac{1+\not{v}}{2} \mathrm{Tr}[\tau\tau^+] T_S(\omega,\omega'), \tag{10}$$

where the coefficients $T_K(\omega,\omega')$ and $T_S(\omega,\omega')$ are analytic functions in the "off-shell energies" $\omega = v\cdot k$ and $\omega' = v\cdot k'$ with discontinuities for positive values of these variables. By saturating the three-point functions with the complete set of baryon states, one can isolate the part of interest, the contribution of the lowest lying baryon states associated with the heavy-light currents, as one having poles in both the variables ω and ω' at the value $\omega = \omega' = \bar{\Lambda}$.

Confining ourselves to taking into account these leading contributions to the perturbation and operators with dimension $D \leqslant 6$ in the OPE, the relevant diagrams in our theoretical calculations for the kinetic energy are shown in Fig. 1. The relevant diagrams for the chromomagnetic interaction do not differ from those for the ground state baryons [13], so we do not show them here. Using the dispersion relations, $T_K(\omega,\omega')$ and $T_S(\omega,\omega')$ can be cast into the form of integrals of the double spectral densities. Following Refs. [18–20], introducing new variables $\omega_+ = \frac{1}{2}(\omega+\omega')$ and $\omega_- = \omega - \omega'$, performing the integral over ω_-, assuming quark-hadron duality in ω_+, and employing the Borel transformation B^{ω}_{τ}, $B^{\omega'}_{\tau'}$ to suppress the continuum contributions and subtractions, we then obtain the desired sum rules. Considering the symmetry of the correlator, it is natural to set the parameters τ, τ' to be the same and equal to $2T$, where T is the Borel parameter of the two-point functions. We end up with the set of sum rules

$$\lambda_2^{\Sigma_{QK1}}F^2e^{-\bar{\Lambda}/T}=\frac{2^{11}3}{\pi^4}\frac{\alpha_s}{\pi}T^{10}\delta_9(\omega_c/T)-\frac{2T^6}{\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{16m_0^2T^2\alpha_s}{3\pi}\langle\bar{q}q\rangle^2e^{-m_0^2/16T^2}, \tag{11a}$$

$$\lambda_2^{\Lambda_{QK1}}F^2e^{-\bar{\Lambda}/T}=\frac{2^{10}3}{\pi^4}\frac{\alpha_s}{\pi}T^{10}\delta_9(\omega_c/T)-\frac{3T^6}{\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{4m_0^2T^2\alpha_s}{3\pi}\langle\bar{q}q\rangle^2e^{-m_0^2/16T^2}, \tag{11b}$$

$$\lambda_2^{\Sigma_{Qk1}}F^2e^{-\bar{\Lambda}/T}=\frac{2^8 35}{\pi^4}\frac{\alpha_s}{\pi}T^{10}\delta_9(\omega_c/T)-\frac{T^6}{2\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{m_0^2T^2\alpha_s}{3\pi}\langle\bar{q}q\rangle^2e^{-m_0^2/16T^2}, \tag{11c}$$

$$\lambda_2^{\Lambda_{Qk1}}F^2e^{-\bar{\Lambda}/T}=\frac{2^9 35}{\pi^4}\frac{\alpha_s}{\pi}T^{10}\delta_9(\omega_c/T)-\frac{3T^6}{\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{4m_0^2T^2\alpha_s}{3\pi}\langle\bar{q}q\rangle^2e^{-m_0^2/16T^2}, \tag{11d}$$

$$-\lambda_1^{\Lambda_{QK1}}F^2e^{-\bar{\Lambda}/T}=\frac{2^6 3^3 5T^{10}}{\pi^4}\delta_9(\omega_c/T)+\frac{19T^6}{2\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{5m_0^4}{16}\langle\bar{q}q\rangle^2e^{-m_0^2/8T^2}, \tag{11e}$$

$$-\lambda_1^{\Sigma_{QK0}}F^2e^{-\bar{\Lambda}/T}=\frac{2^6 3^4 5T^{10}}{\pi^4}\delta_9(\omega_c/T)+\frac{113T^6}{4\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{5m_0^4}{16}\langle\bar{q}q\rangle^2e^{-m_0^2/8T^2}, \tag{11f}$$

$$-\lambda_1^{\Sigma_{QK1}}F^2e^{-\bar{\Lambda}/T}=\frac{2^8 3^2 5T^{10}}{\pi^4}\delta_9(\omega_c/T)-\frac{3T^6}{\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{5m_0^4}{8}\langle\bar{q}q\rangle^2e^{-m_0^2/8T^2}, \tag{11g}$$

$$-\lambda_1^{\Sigma_{Qk1}}F^2e^{-\bar{\Lambda}/T}=\frac{2^4 3^2 17T^{10}}{\pi^4}\delta_9(\omega_c/T)+\frac{21T^6}{8\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{3m_0^4}{64}\langle\bar{q}q\rangle^2e^{-m_0^2/8T^2}, \tag{11h}$$

$$-\lambda_1^{\Lambda_{Qk0}}F^2e^{-\bar{\Lambda}/T}=\frac{2^4 3^2 7T^{10}}{\pi^4}\delta_9(\omega_c/T)-\frac{17T^6}{8\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{3m_0^4}{64}\langle\bar{q}q\rangle^2e^{-m_0^2/8T^2}, \tag{11i}$$

$$-\lambda_1^{\Lambda_{Qk1}}F^2e^{-\bar{\Lambda}/T}=\frac{2^6 3^2 11T^{10}}{\pi^4}\delta_9(\omega_c/T)-\frac{5T^6}{2\pi^2}\left\langle\frac{\alpha_s}{\pi}G^2\right\rangle+\frac{3m_0^4}{32}\langle\bar{q}q\rangle^2e^{-m_0^2/8T^2}. \tag{11j}$$

The unitary normalization of the flavor matrix $\mathrm{Tr}[\tau\tau^+]=1$ has been applied to get these sum rules, as was done for the two-point sum rules.

III. NUMERICAL RESULTS AND CONCLUSIONS

In the following analysis, the standard values for the condensates are adopted [8]. From the two-point sum rules the effective mass can be obtained via a derivative of the Borel parameter as $\bar{\Lambda}=T^2 d\ln E/dT$, where E denotes the right hand side of the obtained two-point sum rule. Then, complying with the standard procedure of sum rule analysis, we change the continuum threshold ω_c and Borel parameter T to find the optimal stability window, and the numerical value of the effective mass is determined within this window. For the sum rules obtained above, we found that the typical value for the continuum threshold is $\omega_c\sim 1.6$ GeV and the typical interval for the Borel parameter is $\Delta T\sim 0.4$ GeV, which is narrower than the window for the ground state baryon. The only exception to this assertion is Λ_{Qk0}, for which we found a much lower continuum threshold $\omega_c\sim 1.2$ GeV and a narrower interval $\Delta T\sim 0.3$ GeV. Also it is worth noting that for Λ_{Qk0} the main contribution does not come from the perturbative part. The condensate contributions play an important role in the determination of the stability window. Indeed, if one keeps only the gluon condensate contribution, there will be no stability window at all, just as in the case in Ref. [14]. But if we assume the perturbative dominance and omit condensate contributions, we only end up with a better stability, and the numerical value is almost exactly the same. For the other sum rules for the effective mass the case is different. The dominant contribution to those sum rules comes from the perturbative part, and the dimension 4 and dimension 6 operators in the OPE only amount to 20% and 10% within the stability windows, respectively. The numerical results are presented in Fig. 2. For clarity we give our numerical average for the effective masses in Table I.

In order to get the numerical results for the two $1/m_Q$ order parameters, we divide our three-point sum rules by two-point functions to obtain λ_1 and λ_2 as functions of the continuum threshold ω_c and the Borel parameter T. This procedure can eliminate the systematic uncertainties and cancel the parameter $\bar{\Lambda}$. From the experience of QCD sum rule applications in the field of heavy quark physics, it is well known that three-point sum rule receives heavier contamina-

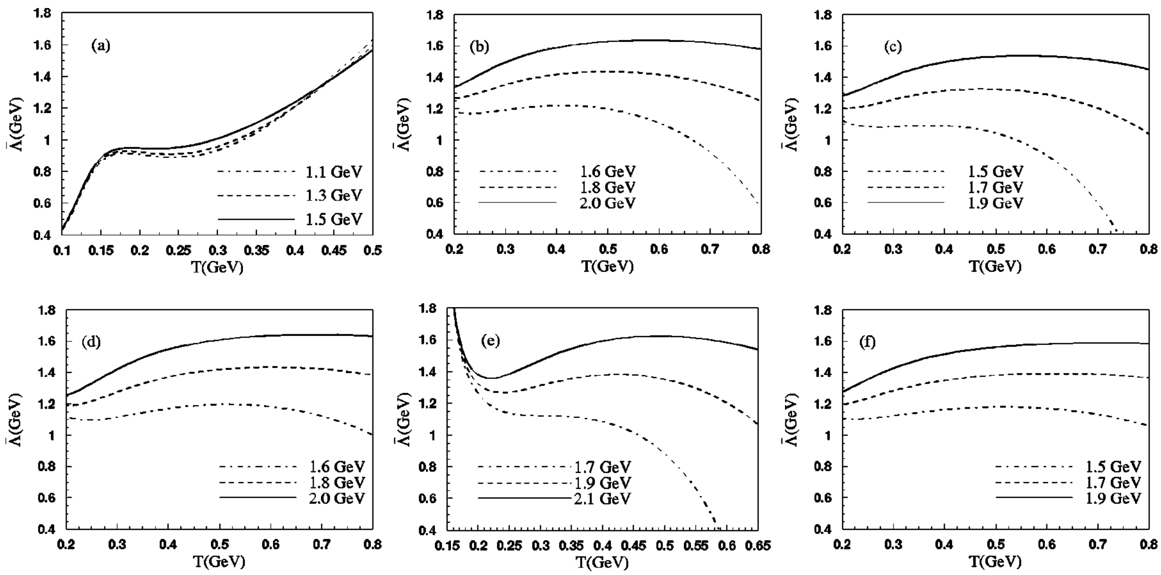

FIG. 2. Sum rules of the effective mass $\overline{\Lambda}$ for (a) Λ_{Qk0}, (b) Λ_{Qk1}, (c) Σ_{Qk1}, (d) Λ_{QK1}, (e) Σ_{QK1}, and (f) Σ_{QK0} baryons. The different choices of the continuum threshold ω_c corresponding to different curves are designated in the individual figures. Curves are plotted against the Borel parameter T.

tion from the continuum modes than the two-point one, and the stability is not as good as that for the two-point sum rule [10,13,17,21]. Our results for the λ_1 sum rules of these excited heavy baryons are shown in Fig. 3, and in Fig. 4 the results for the λ_2 sum rules are presented. The typical value of the continuum threshold is $\omega_c \sim 1.8$ GeV and $\omega_c \sim 2.4$ GeV for the λ_1 and λ_2 sum rules, respectively. The typical interval of the stability window is also $\Delta T \sim 0.3$ GeV for both. The λ_1 sum rule for the Σ_{QK0} baryon is an exception to the above statement, for which the typical value of the continuum threshold is $\omega_c \sim 1.4$ GeV and the interval is $\Delta T \sim 0.2$ GeV. As for the convergence of the OPE within the stability windows for these sum rules, we would like to state some facts. For the λ_1 sum rules, except the case of Λ_{Qk1}, in which the dimension 4 and 6 operators give rise to contributions almost equal to that of the perturbative part, all the other sum rules behave well; in then the contributions of the dimension 4 and 6 operators amount to 50% and 20% of that of the perturbative part. On the other hand, for the λ_2 sum rules, the dimension 4 operator still contributes at almost 50% compared to the perturbative part, but the dimension 6 operator gives only a negligible contribution, typically less than 10%. The numerical results are also listed in Table I.

With these values and the heavy quark masses given in [13], $m_c = 1.41$ GeV and $m_b = 4.77$ GeV, the masses of the excited heavy baryons to order $1/m_Q$ can be obtained immediately. We give the masses in Table II. Our results are comparable to the predictions for the excited heavy baryon masses obtained by using the quark potential model [22].

The splitting between the spin 1/2 and spin 3/2 doublets can be obtained by multiplying λ_2 by a factor of 3/2 [cf. Eq. (3)]. It is 0.13 ± 0.02 GeV^2, 0.16 ± 0.05 GeV^2, 0.20 ± 0.02 GeV^2, and 0.19 ± 0.04 GeV^2 for Λ_{QK1}, Λ_{Qk1}, Σ_{QK1} and Σ_{Qk1} doublets, respectively. The approximately equal values for Λ_{QK1} and Λ_{Qk1}, Σ_{QK1} and Σ_{Qk1} may be interpreted as signal which implies that the current mixing effect cannot be large. If we take the middle value as the theoretical prediction for the physical state, then the splittings for the excited baryon states are

$$\Lambda_{Q1}^{*2} - \Lambda_{Q1}^{2} = 0.15 \pm 0.03 \ \text{GeV}^2,$$

$$\Sigma_{Q1}^{*2} - \Sigma_{Q1}^{2} = 0.20 \pm 0.03 \ \text{GeV}^2. \qquad (12)$$

Based on the current experimental data [23], the splitting of the excited Λ_c doublet is 0.17 GeV^2, which is in agreement with our theoretical result. When it is scaled up to the bottom quark mass scale there will be a factor of approximately 0.8 due to the renormalization group improvement.

TABLE I. Effective mass in GeV, and kinetic energy and chromomagnetic interaction energy in GeV^2 for excited heavy baryons. Errors quoted are due to the variation of the Borel parameter T and continuum threshold ω_c.

	Λ_{Qk0}	Σ_{Qk1}	Λ_{Qk1}	Λ_{QK1}	Σ_{QK1}	Σ_{QK0}
$\overline{\Lambda}$	1.04 ± 0.17	1.12 ± 0.22	1.36 ± 0.18	1.30 ± 0.13	1.27 ± 0.11	1.21 ± 0.09
$-\lambda_1$	1.20 ± 0.26	1.02 ± 0.25	0.84 ± 0.15	1.16 ± 0.11	1.48 ± 0.18	1.42 ± 0.25
λ_2		0.13 ± 0.03	0.11 ± 0.03	0.09 ± 0.01	0.13 ± 0.01	

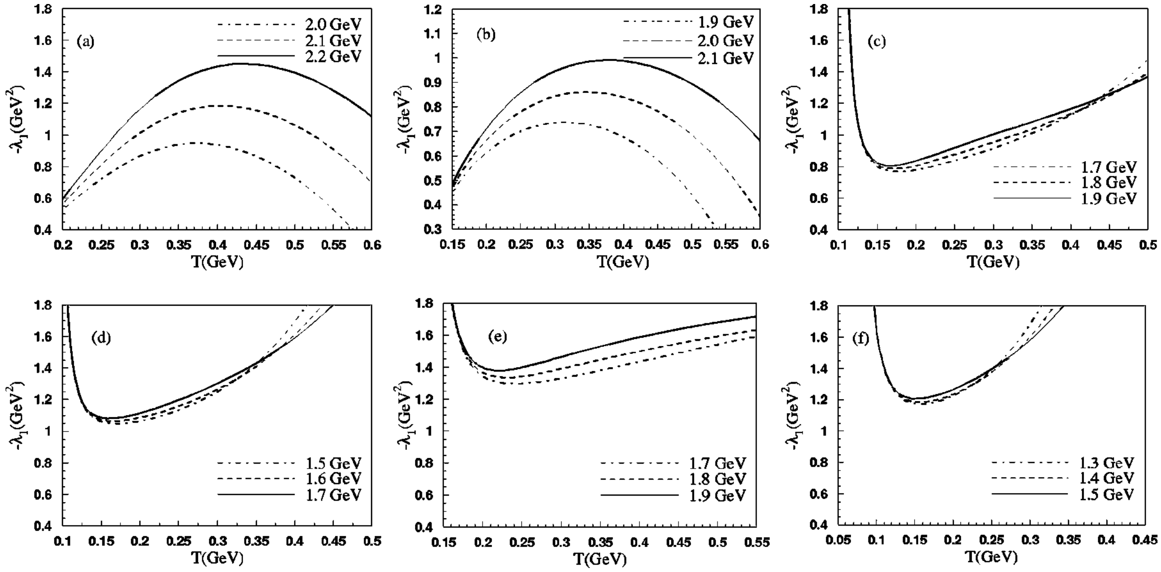

FIG. 3. Sum rules of the kinetic energy for (a) Λ_{Qk0}, (b) Λ_{Qk1}, (c) Σ_{Qk1}, (d) Λ_{QK1}, (e) Σ_{QK1}, and (f) Σ_{QK0} baryons. Other details are the same as in Fig. 2.

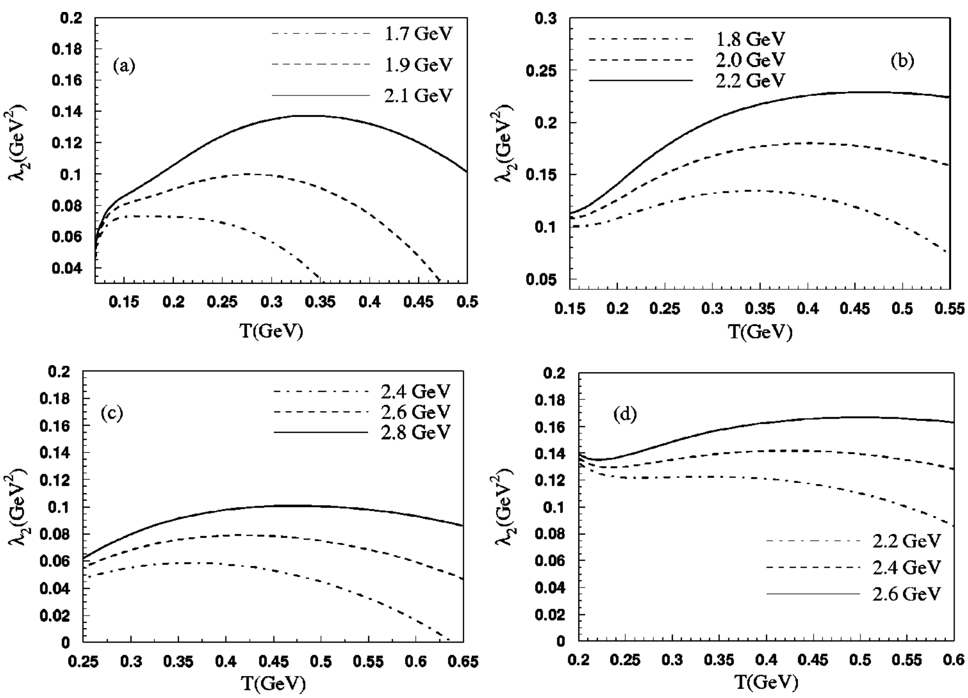

FIG. 4. Sum rules of the chromomagnetic interaction for (a) Λ_{Qk1}, (b) Σ_{Qk1}, (c) Λ_{QK1}, and (d) Σ_{QK1} baryons. Other details are the same as those in Fig. 2.

TABLE II. Masses in GeV for excited heavy baryons.

	Λ_{Qk0}	Λ_{Qk1}	Λ^*_{Qk1}	Σ_{Qk1}	Σ^*_{Qk1}	Λ_{QK1}	Λ^*_{QK1}	Σ_{QK1}	Σ^*_{QK1}	Σ_{QK0}
$Q=c$	2.863	3.019	3.076	2.831	2.900	3.083	3.129	3.148	3.219	3.113
$Q=b$	5.934	6.205	6.222	5.977	5.998	6.185	6.199	6.182	6.203	6.127

To conclude, we have calculated the $1/m_Q$ order corrections to the excited heavy baryon masses from the QCD sum rules within the framework of the HQET. From the spectrum thus obtained for the c quark case, we found that the Λ_{Qk0} and Σ_{Qk1} baryons lie $\sim$600 MeV above the ground state baryon Λ_c, while Λ_{Qk1}, Λ_{QK1}, and Σ_{QK1}, lie $\sim$800 MeV above Λ_c and for the Σ_{QK1} baryon this value is $\sim$900 MeV, typically with an error $\sim$300 MeV. When it comes to the b quark case, the result is that the Λ_{Qk0} and Σ_{Qk1} baryons lie $\sim$300 MeV above the ground state baryon Λ_b, while Λ_{Qk1}, Λ_{QK1}, Σ_{QK1}, and Σ_{QK1} lie $\sim$550 MeV above Λ_b, for which the typical error is $\sim$200 MeV. For the c quark case, the $1/m_Q$ order corrections are $\sim$400 MeV, which is not small because of the large value of the kinetic energy. For the b quark case, these corrections will be suppressed by the still larger b quark mass. Our theoretical predictions for the doublet splitting are in agreement with the current experimental data.

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[1] ARGUS Collaboration, H. Albrecht *et al.*, Phys. Lett. B **317**, 227 (1993); **402**, 207 (1997); E687 Collaboration, P.L. Frabetti *et al.*, Phys. Rev. Lett. **72**, 961 (1994); Phys. Lett. B **365**, 461 (1996); CLEO Collaboration, K.W. Edwards *et al.*, Phys. Rev. Lett. **74**, 3331 (1995).

[2] CLEO Collaboration, J.P. Alexander *et al.*, Phys. Rev. Lett. **83**, 3390 (1999).

[3] B. Grinstein, Nucl. Phys. **B339**, 253 (1990); E. Eichten and B. Hill, Phys. Lett. B **234**, 511 (1990); A.F. Falk, H. Georgi, B. Grinstein, and M.B. Wise, Nucl. Phys. **B343**, 1 (1990); F. Hussain, J.G. Körner, K. Schilcher, G. Thompson, and Y.L. Wu, Phys. Lett. B **249**, 295 (1990); J.G. Körner and G. Thompson, *ibid.* **264**, 185 (1991).

[4] M. Neubert, Phys. Rep. **245**, 259 (1994).

[5] J.G. Körner, M. Krämer, and D. Pirjol, Prog. Part. Nucl. Phys. **33**, 787 (1994).

[6] A.F. Falk and M. Neubert, Phys. Rev. D **47**, 2965 (1993); **47**, 2982 (1993).

[7] I.I. Bigi, N.G. Uraltsev, and A.I. Vainshtein, Phys. Lett. B **293**, 430 (1992); I.I. Bigi, M. Shifman, N.G. Uraltsev, and A.I. Vainshtein, Phys. Rev. Lett. **71**, 496 (1993); A.V. Manohar and M.B. Wise, Phys. Rev. D **49**, 1310 (1994); M. Luke and M.J. Savage, Phys. Lett. B **321**, 88 (1994); M. Neubert and C.T. Sachrajda, Nucl. Phys. **B483**, 339 (1997).

[8] M.A. Shifman, A.I. Vainshtein, and V.I. Zakharov, Nucl. Phys. **B147**, 385 (1979); **B147**, 448 (1979); V.A. Novikov, M.A. Shifman, A.I. Vainshtein, and V.I. Zakharov, Fortschr. Phys. **32**, 11 (1984).

[9] For a recent review of the QCD sum rule method, see P. Colangelo and A. Khodjamirian, in *At the Frontier of Particle Physics/Handbook of QCD*, edited by M. Shifman (World Scientific, Singapore, 2001), pp. 1495–1576.

[10] A.G. Grozin and O.I. Yakovlev, Phys. Lett. B **285**, 254 (1992); **291**, 441 (1992).

[11] S. Groote, J.G. Körner, and O.I. Yakovlev, Phys. Rev. D **55**, 3016 (1997).

[12] Y.-B. Dai, C.-S. Huang, C. Liu, and C.-D Lü, Phys. Lett. B **371**, 99 (1996).

[13] D.-W. Wang, M.-Q. Huang, and C.-Z. Li, Phys. Rev. D **65**, 094036 (2002).

[14] C.-S. Huang, A.-L. Zhang, and S.-L. Zhu, Phys. Lett. B **492**, 288 (2000).

[15] S.-L. Zhu, Phys. Rev. D **61**, 114019 (2000); J.P. Lee, C. Liu, and H.S. Song, Phys. Lett. B **476**, 303 (2000); M.-Q. Huang, J.P. Lee, C. Liu, and H.S. Song, *ibid.* **502**, 133 (2001).

[16] M.A. Ivanov, J.G. Körner, V.E. Lyubovitskij, M.A. Pisarev, and A.G. Rusetsky, Phys. Rev. D **61**, 114010 (2000).

[17] D.-W. Wang and M.-Q. Huang, Phys. Rev. D **67**, 074025 (2003).

[18] M. Neubert, Phys. Rev. D **45**, 2451 (1992); **46**, 3914 (1992).

[19] Y.-B. Dai, C.-S. Huang, M.-Q. Huang, and C. Liu, Phys. Lett. B **387**, 379 (1996).

[20] M. Neubert, Phys. Rev. D **46**, 1076 (1992).

[21] P. Ball and V.M. Braun, Phys. Rev. D **49**, 2472 (1994); M. Neubert, Phys. Lett. B **389**, 727 (1996).

[22] L.A. Copley, N. Isgur, and G. Karl, Phys. Rev. D **20**, 768 (1979); S. Capstick and N. Isgur, *ibid.* **34**, 2809 (1986).

[23] Particle Data Group, K. Hagiwara *et al.*, Phys. Rev. D **66**, 010001 (2002).