

Predictions of masses for light hybrid baryons

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Within the method of parity-projected QCD sum rules, we study the mass spectra of light hybrid baryons with $I(J^P) = 1/2(1/2^\pm), 3/2(1/2^\pm), 1/2(3/2^\pm), 3/2(3/2^\pm)$ by constructing the local $qqqq$ interpolating currents. We calculate the correlation functions up to dimension eight condensates at the leading order of α_s . The stable QCD Lapalce sum rules can be established for the positive-parity $N_{1/2^+}, \Delta_{3/2^+}, \Delta_{1/2^+}$ and negative-parity $N_{1/2^-}, N_{3/2^-}, \Delta_{1/2^-}$ channels to extract their mass spectra. The lowest-lying hybrid baryons are predicted to be the positive-parity $N_{1/2^+}$ state around 2.01 GeV. These hybrid baryons mainly decay into conventional baryon plus meson final states. We propose to search for the light hybrid baryons through the $\Upsilon/\psi(3686)$ decays via the three-gluon emission mechanism in BESIII and BelleII experiments. Hopefully, our studies of the light hybrid baryons will be useful for understanding the excited baryon spectrum and the behavior of gluonic degrees of freedom in QCD.

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I. INTRODUCTION

As a distinctive prediction of QCD, the existence of hybrid states is a longstanding issue in both experimental and theoretical aspects in the past several decades. Different from the conventional $q\bar{q}$ meson and qqq baryon [1,2], hybrid states contain the gluonic degrees of freedom in the configurations, such as $\bar{q}gq$ hybrid mesons and $qqqg$ hybrid baryons. It is of crucial importance to identify the existence of hybrid states for understanding the non-perturbative properties of QCD [3–10].

In the past several decades, there have been numerous investigations on hybrid mesons in various theoretical

methods, such as the MIT bag model [11,12], Coulomb gauge QCD [13], Lattice QCD (LQCD) [14–22], the flux tube model [23–25], Bethe-Salpeter equation [26,27], QCD sum rules [28–43], and so on. The most intriguing candidates for light hybrid mesons are $\pi_1(1400)$ [44], $\pi_1(1600)$ [45], $\pi_1(2015)$ [46], and $\eta_1(1855)$ [47,48] with exotic $J^{PC} = 1^{-+}$. Comparing to hybrid mesons, a hybrid baryon is composed of a color octet triquark field and a color octet gluonic excitation. To date, the studies of hybrid baryons are much less appealing since they cannot carry exotic quantum numbers. In general, QCD allows arbitrarily strong mixing effect between $qqqq$ hybrid baryons and excited qqq states with the same quantum numbers.

In Refs. [49–51], the Roper resonance $N(1440)$ was studied as a hybrid baryon with $J^P = 1/2^+$. Such a possibility is excluded according to the meson-cloud theory from the constituent quark model [52]. In Ref. [49], the masses of light hybrid baryon states with spin-parity $J^P = 1/2^+, 3/2^+, 5/2^+$ were calculated in the MIT bag model. They predicted the masses of these hybrids in the range of 1.6–2.5 GeV, among which the lightest one is the $1/2^+$ state. In the flux-tube model, the lightest hybrid baryons are the mass-degenerate states with $J^P = 1/2^+$ and $3/2^+$ near 1.87 GeV [53,54]. Using the relativistic quark model, the

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authors of Ref. [55] studied the mass spectra of the nucleon and delta hybrids and suggested that $N(1710)$ and $\Delta(1600)$ could be hybrid baryons. In Refs. [56,57], the mass spectra of positive-parity hybrid baryons were calculated in LQCD, in which $N_{1/2^+}$ and $N_{3/2^+}$ were predicted to be the lowest-lying hybrid states around 2.3–2.6 GeV. The QCD sum rule method has also been applied to study the hybrid baryons [51,58–60]. In Ref. [58], a light hybrid baryon $N_{1/2^+}$ was studied and found to be approximately 2.1 GeV. Using the same interpolating current, the calculation was improved to give different results of the gluon, trigluon, and mixed condensates, while the hybrid mass was found to be around the observed mass of the Roper resonance [51]. In QCD sum rules, the $\Lambda(1600)$ and $\Lambda(1405)$ were suggested as strange hybrid baryons with $I(J^P) = 0(1/2^+)$ [59] and $I(J^P) = 0(1/2^-)$ [60], respectively. Recently, the triply heavy $cccg$ and $bbbg$ hybrids were studied in QCD sum rules [61] and a semirelativistic potential model [62].

In this work, we systematically study the mass spectra of nucleon and delta hybrid baryons with both positive and negative parities by using the parity-projected QCD sum rules. We construct the $[qqq]_{8c}[G]_{8c}$ local hybrid baryon

interpolating currents without the covariant derivative operators and introduce the parity projected sum rules to separate the hybrid baryons with different parities. We find stable mass sum rules and establish extensive mass spectra for the positive-parity $N_{1/2^+}, \Delta_{3/2^+}, \Delta_{1/2^+}$ and negative-parity $N_{1/2^-}, N_{3/2^-}, \Delta_{1/2^-}$ hybrid baryons. Our results show that the lowest-lying hybrid baryon is predicted to be the positive-parity $N_{1/2^+}$ state around 2.01 GeV. We propose searching for these light hybrid baryons through the χ_{cJ}/Υ decays via the three-gluon emission mechanism in BESIII and BelleII experiments.

II. HYBRID BARYON CURRENTS AND PARITY-PROJECTED QCD SUM RULES

A hybrid baryon current is composed of a color-octet triquark operator $[qqq]_{8c}$ and a color-octet excited gluon field. To investigate the low-lying states, we use the S-wave triquark fields introduced as “Ioffe currents” [63]. The gluon field is combined to compose the color-singlet hybrid baryon interpolating currents with $J^P = 1/2^+$ and $3/2^+$,

$$\begin{aligned}
J_1 &= \varepsilon^{abc} g_s [(u^{aT} C \gamma^\mu u^b) \gamma^\nu \gamma^5 (G_{\mu\nu} d)^c - (u^{aT} C \gamma^\mu d^b) \gamma^\nu \gamma^5 (G_{\mu\nu} u)^c], \\
J_2 &= \varepsilon^{abc} g_s [(u^{aT} C \gamma^\mu u^b) \gamma^\nu \gamma^5 (G_{\mu\nu} d)^c + 2(u^{aT} C \gamma^\mu d^b) \gamma^\nu \gamma^5 (G_{\mu\nu} u)^c], \\
J_3 &= \varepsilon^{abc} g_s [(u^{aT} C \sigma^{\mu\nu} u^b) \gamma^5 (G_{\mu\nu} d)^c - (u^{aT} C \sigma^{\mu\nu} d^b) \gamma^5 (G_{\mu\nu} u)^c], \\
J_4 &= \varepsilon^{abc} g_s [(u^{aT} C \sigma^{\mu\nu} u^b) \gamma^5 (G_{\mu\nu} d)^c + 2(u^{aT} C \sigma^{\mu\nu} d^b) \gamma^5 (G_{\mu\nu} u)^c], \\
J_5 &= \varepsilon^{abc} g_s [(u^{aT} C \gamma^\mu u^b) \gamma^\nu (\tilde{G}_{\mu\nu} d)^c - (u^{aT} C \gamma^\mu d^b) \gamma^\nu (\tilde{G}_{\mu\nu} u)^c], \\
J_6 &= \varepsilon^{abc} g_s [(u^{aT} C \gamma^\mu u^b) \gamma^\nu (\tilde{G}_{\mu\nu} d)^c + 2(u^{aT} C \gamma^\mu d^b) \gamma^\nu (\tilde{G}_{\mu\nu} u)^c], \\
J_7 &= \varepsilon^{abc} g_s [(u^{aT} C \sigma^{\mu\nu} u^b) (\tilde{G}_{\mu\nu} d)^c - (u^{aT} C \sigma^{\mu\nu} d^b) (\tilde{G}_{\mu\nu} u)^c], \\
J_8 &= \varepsilon^{abc} g_s [(u^{aT} C \sigma^{\mu\nu} u^b) (\tilde{G}_{\mu\nu} d)^c + 2(u^{aT} C \sigma^{\mu\nu} d^b) (\tilde{G}_{\mu\nu} u)^c], \\
J_{1\mu} &= \varepsilon^{abc} g_s [(u^{aT} C \gamma^\nu u^b) (G_{\mu\nu} d)^c - (u^{aT} C \gamma^\nu d^b) (G_{\mu\nu} u)^c], \\
J_{2\mu} &= \varepsilon^{abc} g_s [(u^{aT} C \gamma^\nu u^b) (G_{\mu\nu} d)^c + 2(u^{aT} C \gamma^\nu d^b) (G_{\mu\nu} u)^c], \\
J_{3\mu} &= \varepsilon^{abc} g_s [(u^{aT} C \gamma^\nu u^b) \gamma^5 (\tilde{G}_{\mu\nu} d)^c - (u^{aT} C \gamma^\nu d^b) \gamma^5 (\tilde{G}_{\mu\nu} u)^c], \\
J_{4\mu} &= \varepsilon^{abc} g_s [(u^{aT} C \gamma^\nu u^b) \gamma^5 (\tilde{G}_{\mu\nu} d)^c + 2(u^{aT} C \gamma^\nu d^b) \gamma^5 (\tilde{G}_{\mu\nu} u)^c],
\end{aligned} \tag{1}$$

where u/d is the up/down quark field, $G_{\mu\nu}$ is the gluon field strength, and $\tilde{G}_{\mu\nu} = \frac{1}{2} \varepsilon_{\mu\nu\alpha\beta} G^{\alpha\beta}$ the dual gluon field strength. The currents $J_{1,3,5,7,1\mu,3\mu}$ can couple to the nucleon hybrids $N_{1/2^\pm}, N_{3/2^\pm}$ with $I = 1/2$, while $J_{2,4,6,8,2\mu,4\mu}$ couple to the delta hybrids $\Delta_{1/2^\pm}, \Delta_{3/2^\pm}$ with $I = 3/2$. The couplings between currents and the corresponding positive-parity states can be written as [64–67]

$$\begin{aligned}
\langle 0 | J_i | H_+(p) \rangle &= f_+ u(p), \\
\langle 0 | J_{i\mu} | H_+(p) \rangle &= f'_+ u_\mu(p),
\end{aligned} \tag{2}$$

in which $u(p)$ and $u_\mu(p)$ are the Dirac spinor and Rarita-Schwinger vector, respectively, and f_+ is the coupling constant. These currents can also couple to the negative-parity states via different coupling relations

$$\begin{aligned}
\langle 0 | J_i | H_-(p) \rangle &= f_- \gamma_5 u(p), \\
\langle 0 | J_{i\mu} | H_-(p) \rangle &= f'_- \gamma_5 u_\mu(p).
\end{aligned} \tag{3}$$

The two-point correlation functions can be written as

$$\begin{aligned}\Pi(p^2) &= i \int d^4x e^{ip \cdot x} \langle 0 | T[J(x) \bar{J}(0)] | 0 \rangle = \Pi_{1/2}(p^2), \\ \Pi_{\mu\nu}(p^2) &= i \int d^4x e^{ip \cdot x} \langle 0 | T[J_\mu(x) \bar{J}_\nu(0)] | 0 \rangle \\ &= -g_{\mu\nu} \Pi_{3/2}(p^2) + \dots,\end{aligned}\quad (4)$$

where $\Pi_{1/2}(p^2)$ and $\Pi_{3/2}(p^2)$ are the invariant functions for the hadron states with spin 1/2 and 3/2, respectively. They contain information of both positive- and negative-parity states

$$\begin{aligned}\Pi(p^2) &= \not{p} \Pi_{\not{p}} + \Pi_I \\ &= \frac{1}{2} \left(\frac{\not{p}}{\sqrt{p^2}} + 1 \right) \Pi_+(p^2) + \frac{1}{2} \left(\frac{\not{p}}{\sqrt{p^2}} - 1 \right) \Pi_-(p^2)\end{aligned}\quad (5)$$

in which $\not{p} = \gamma_\mu p^\mu$. Π_+ and Π_- are the correlation function for positive- and negative-parity states respectively, which satisfy the equation

$$\Pi_\pm(p^2) = \frac{f_\pm^2}{M_\pm^2 - p^2} + \dots, \quad (6)$$

where “...” denote the excited states and continuum contributions. At the hadronic level, we use the dispersion relation to describe the invariant function as [68]

$$\Pi(p^2) = \frac{(p^2)^N}{\pi} \int_0^\infty \frac{\text{Im} \Pi(s)}{s^N (s - p^2 - i\epsilon)} ds + \sum_{n=0}^{N-1} b_n (p^2)^n, \quad (7)$$

where N is the number of subtractions required to make the integral convergent. The unknown subtraction terms shall disappear after applying the Borel transform. The spectral function can be defined in the “narrow resonance” approximation

$$\begin{aligned}\rho(s) &\equiv \frac{\text{Im} \Pi(s)}{\pi} = \not{p} \rho^{\not{p}}(s) + \rho^I(s) \\ &= \frac{1}{2} \left(\frac{\not{p}}{\sqrt{s}} + 1 \right) f_+^2 \delta(s - M_+^2) \\ &\quad + \frac{1}{2} \left(\frac{\not{p}}{\sqrt{s}} - 1 \right) f_-^2 \delta(s - M_-^2) + \dots,\end{aligned}\quad (8)$$

where

$$\begin{aligned}\rho^{\not{p}}(s) &= \frac{1}{2} \frac{f_+^2}{\sqrt{s}} \delta(s - M_+^2) + \frac{1}{2} \frac{f_-^2}{\sqrt{s}} \delta(s - M_-^2) + \dots, \\ \rho^I(s) &= \frac{1}{2} f_+^2 \delta(s - M_+^2) - \frac{1}{2} f_-^2 \delta(s - M_-^2) + \dots.\end{aligned}\quad (9)$$

The negative- and positive-parity states can be extracted as

$$\rho_\pm(s) = \sqrt{s} \rho^{\not{p}}(s) \pm \rho^I(s). \quad (10)$$

Then, the QCD sum rules can be established as

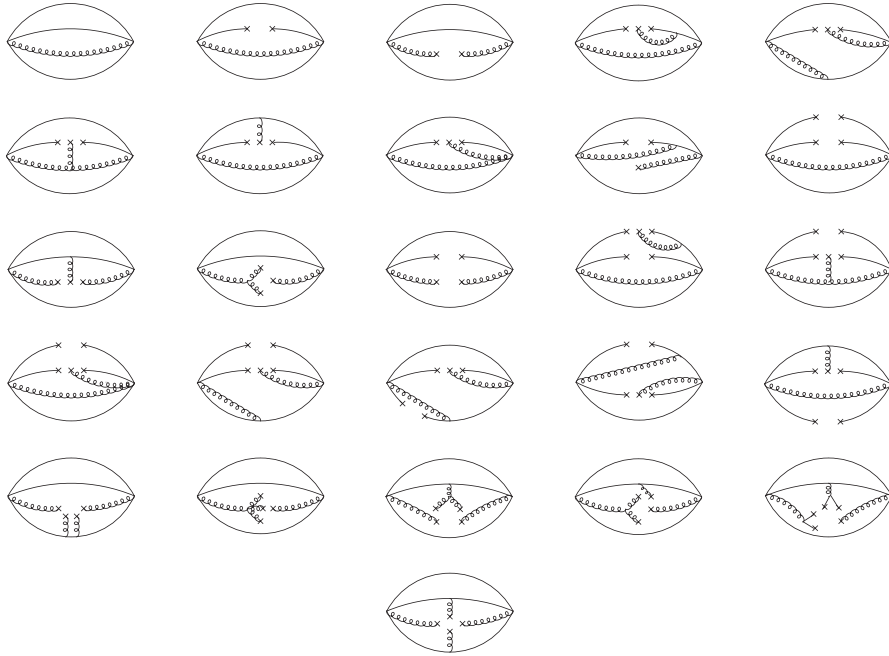


FIG. 1. Feynman diagrams involved in the calculations.

$$\Pi_{\pm}(M_B^2, s_0) = \int_{(3m_q)^2}^{s_0} \rho_{\pm}(s) e^{-s/M_B^2} ds = f_{\pm}^2 e^{-M_{\pm}^2/M_B^2}, \quad (11)$$

where $\rho_{\pm}(s) = \text{Im} \Pi_{\pm}(s)/\pi$, m_q is the light quark mass ($q = u/d$), and M_B and s_0 are the Borel mass and continuum threshold, respectively. It is clear that $\rho_{\pm}(s)$ contains only contributions from positive/-negative-parity baryon states. The lowest-lying hadron mass can be obtained as

$$M_{\pm}(M_B^2, s_0) = \sqrt{\frac{\frac{\partial}{\partial(-1/M_B^2)} \Pi_{\pm}(M_B^2, s_0)}{\Pi_{\pm}(M_B^2, s_0)}}. \quad (12)$$

At the QCD side, the correlation functions and the spectral functions can be evaluated via the operator product expansion (OPE). In this work, we calculate the correlation functions at the leading order of α_s up to dimension-8 condensates, including the Feynman diagrams in Fig. 1. As an example, we show the spectral function for the current J_7 as

$$\begin{aligned} \rho_7^{\not{p}}(s) = & \frac{3\alpha_s s^4}{20480\pi^5} + \frac{\alpha_s \langle \bar{q}q \rangle^2 s}{12\pi} - \frac{49\langle \alpha_s GG \rangle^2}{2304\pi^2} + \frac{11\langle g_s^3 G^3 \rangle M^2}{3072\pi^4} \\ & - \frac{47\alpha_s \langle g_s \bar{q}\sigma Gq \rangle \langle \bar{q}q \rangle}{576\pi} - \frac{93\alpha_s^2 \langle \bar{q}q \rangle^2 M^2}{256\pi^2} \\ & + \frac{19\alpha_s^2 \langle \bar{q}q \rangle^2 M^2 \log(\frac{s}{\mu^2})}{64\pi^2} + \frac{13\langle \alpha_s GG \rangle^2 \log(\frac{s}{\mu^2})}{1152\pi^2}, \end{aligned} \quad (13)$$

$$\rho_7^I(s) = -\frac{\langle \alpha_s GG \rangle \langle \bar{q}q \rangle s}{32\pi} + \frac{55\alpha_s \langle g_s \bar{q}\sigma Gq \rangle s^2}{768\pi^3} - \frac{5\alpha_s \langle \bar{q}q \rangle s^3}{192\pi^3}. \quad (14)$$

At the leading order of α_s , the dimension-4 and dimension-6 gluon condensates give no contribution for the spin-1/2 currents $J_1 - J_8$. We calculate the $\langle \bar{q}q \rangle \langle \bar{q}q \rangle$ and $\langle \bar{q}q \rangle \langle \alpha_s GG \rangle$ condensates by applying the factorization assumption, which is usually adopted in QCD sum rules to estimate the values of the high dimension condensates [69,70]. There are 11 independent gluon condensates for dimension-8 operators [71,72],

$$\begin{aligned} \phi_1 &\equiv \langle \text{Tr}(G_{\nu\mu} G^{\nu\mu}) \text{Tr}(G_{\tau\rho} G^{\tau\rho}) \rangle, \\ \phi_2 &\equiv \langle \text{Tr}(G_{\nu\mu} G^{\rho\mu}) \text{Tr}(G^{\nu}_{\tau} G^{\rho}_{\tau}) \rangle, \\ \phi_3 &\equiv \langle \text{Tr}(G_{\nu\mu} G^{\tau\rho}) \text{Tr}(G^{\nu\mu} G_{\tau\rho}) \rangle, \\ \phi_4 &\equiv \langle \text{Tr}(G_{\nu\mu} G^{\tau\rho}) \text{Tr}(G^{\nu}_{\tau} G^{\mu}_{\rho}) \rangle, \\ \phi_5 &\equiv \langle \text{Tr}(G_{\mu\nu} G_{\mu\nu} G_{\alpha\beta} G_{\alpha\beta}) \rangle, \\ \phi_6 &\equiv \langle \text{Tr}(G_{\mu\nu} G_{\alpha\beta} G_{\mu\nu} G_{\alpha\beta}) \rangle, \\ \phi_7 &\equiv \langle \text{Tr}(G_{\mu\nu} G_{\nu\alpha} G_{\alpha\beta} G_{\beta\mu}) \rangle, \\ \phi_8 &\equiv \langle \text{Tr}(G_{\mu\nu} G_{\alpha\beta} G_{\nu\alpha} G_{\beta\mu}) \rangle, \\ \phi_9 &\equiv \langle f_{abc} G_{\mu\nu}^a J_{\mu}^b J_{\nu}^c \rangle, \\ \phi_{10} &\equiv \langle J_{\mu}^a D^2 J_{\mu}^a \rangle, \\ \phi_{11} &\equiv \langle f^{abc} G_{\mu\nu}^a G_{\nu\lambda}^b D^2 G_{\lambda\mu}^c \rangle, \end{aligned} \quad (15)$$

where $J_{\mu}^a = g_s \sum_{u,d,s} \bar{q} \gamma_{\mu} (\lambda^a/2) q$. We also apply the factorization assumption for these operators to simplify the calculations, as discussed in Refs. [71–73],

$$\begin{aligned} g_s^4 \phi_1 &= \frac{1}{4} \langle \alpha_s GG \rangle^2, \\ g_s^4 \phi_3 &= -\frac{1}{16} \langle \alpha_s GG \rangle^2 + 2g_s^4 \phi_2, \\ g_s^4 \phi_4 &= -\frac{1}{32} \langle \alpha_s GG \rangle^2 + g_s^4 \phi_2, \\ g_s^4 \phi_5 &= \frac{16\pi^2}{12} \langle \alpha_s GG \rangle^2, \\ g_s^4 \phi_6 &= -\frac{5 \times 16\pi^2}{48} \langle \alpha_s GG \rangle^2 + 2g_s^4 \phi_2, \\ g_s^4 \phi_7 &= g_s^4 \phi_8 = \frac{16\pi^2}{192} \langle \alpha_s GG \rangle^2 + \frac{1}{2} g_s^4 \phi_2, \\ g_s^3 \phi_9 &= -\frac{3}{2} g_s^4 \langle \bar{q}q \rangle^2 M^2, \\ g_s^2 \phi_{10} &= -\frac{4}{3} g_s^4 \langle \bar{q}q \rangle^2 M^2, \\ g_s^3 \phi_{11} &= \langle g_s^3 G^3 \rangle M^2, \end{aligned} \quad (16)$$

where $g_s^4 \phi_2$ is chosen as $\frac{7 \times 16\pi^2}{96} \langle \alpha_s GG \rangle^2$ [71] and $M^2 = 0.3 \text{ GeV}^2$ [74].

In Eq. (12), one should find reasonable working regions of s_0 and M_B^2 to determine the hadron mass. We require a good OPE convergence to obtain the lower bound of M_B^2 ,

$$R_{D>6} = \left| \frac{\Pi_{\pm}^{D>6}(M_B^2, \infty)}{\Pi_{\pm}^{\text{tot}}(M_B^2, \infty)} \right|, \quad (17)$$

where $\Pi_{\pm}^{D>6}(M_B^2, \infty)$ represents the contributions from $D > 6$ condensates. The upper bound of M_B^2 and the

optimal value of s_0 are found by ensuring a sufficiently large pole contribution

$$PC = \frac{\Pi_{\pm}(M_B^2, s_0)}{\Pi_{\pm}(M_B^2, \infty)}. \quad (18)$$

III. NUMERICAL ANALYSES

To perform numerical analyses, we use the following values for various QCD parameters at the renormalization scale $\mu = 1$ GeV and $\Lambda_{\text{QCD}} = 300$ MeV [72,75–77]:

$$\begin{aligned} m_u &= m_d = m_q = 0, \\ \alpha_s(\mu) &= \frac{4\pi}{29/3 \ln(\mu^2/\Lambda_{\text{QCD}}^2)}, \\ \langle \bar{q}q \rangle &= -(0.24 \pm 0.01)^3 \text{ GeV}^3, \\ \langle g_s \bar{q} \sigma G q \rangle &= (0.8 \pm 0.2) \times \langle \bar{q}q \rangle \text{ GeV}^2, \\ \langle \alpha_s GG \rangle &= (6.35 \pm 0.35) \times 10^{-2} \text{ GeV}^4, \\ \langle g_s^3 G^3 \rangle &= (8.2 \pm 1.0) \times \langle \alpha_s GG \rangle \text{ GeV}^2. \end{aligned} \quad (19)$$

As an example, we show our numerical analysis for the current J_7 , which can couple to both negative- and positive-parity hybrid baryons with $I(J) = 1/2(1/2)$. We first investigate the hybrid baryon with $J^P = 1/2^+$. We require $R_{D>6} \leq 10\%$ to guarantee the good OPE convergence, leading to the lower bound of the Borel mass $M_B^2 \geq 1.14 \text{ GeV}^2$. By using this value and requiring that the pole contribution be larger than 30%, the threshold value s_0 can be determined around 5.3 GeV^2 . We set 5% uncertainty of s_0 in our analyses. In Fig. 2(a), we show the variations of the hybrid baryon mass m_X with s_0 for various Borel mass M_B^2 , from which the optimal working region is $5.03 \leq s_0 \leq 5.57 \text{ GeV}^2$. Using the central value of s_0 , the upper bound of M_B^2 can be obtained by requiring the pole

contribution be larger than 30%. Finally, the Borel parameter can be determined to be $1.14 \leq M_B^2 \leq 1.46 \text{ GeV}^2$.

As shown in Fig. 2(b), the Borel curves are very stable in the working parameter regions so that the mass predictions are reliable. The mass of the hybrid baryon with $I(J^P) = 1/2(1/2)^+$ is predicted as

$$M_+ = 2.01_{-0.08}^{+0.12} \text{ GeV}, \quad (20)$$

where the errors come from s_0 and various condensates in Eq. (19).

For all interpolating currents in Eq. (1), the numerical analyses show that the OPE convergences and pole contributions may vary significantly for different channels. For instance, the pole contributions for the positive-parity channels from currents J_6 and $J_{2\mu}$ and negative-parity channels from currents J_8 are significant so that we set $PC \geq 50\%$ in the analyses. However, for the negative-parity channels from J_4 and J_5 , the pole contributions are relatively small so that we set them as $PC \geq 20\%$. Besides, the OPE convergence behavior for some channels, such as the positive-parity channel from $J_{2\mu}$, is much better than other ones. These channels can actually provide more reliable mass predictions.

In Table I, we collect the numerical results for the hybrid baryons by using all interpolating currents in Eq. (1). It is shown that there are reliable mass predictions for the positive-parity $N_{1/2^+}, \Delta_{1/2^+}, \Delta_{3/2^+}$ and negative-parity $N_{1/2^-}, N_{3/2^-}, \Delta_{1/2^-}$ hybrid baryons. For some currents, the QCD sum rules are unstable, so the mass predictions are unreliable for these channels, e.g., $N_{3/2^+}$ and $\Delta_{3/2^-}$, for which we denote as “...” in Table I. As depicted in Fig. 3, the lightest hybrid baryon is predicted to appear in the positive-parity $N_{1/2^+}$ channel, the lightest negative-parity $N_{1/2^-}$ is about 2.43 GeV, and the lightest $\Delta_{1/2^-}$ state is about 400 MeV higher. Such a mass splitting is almost the same with the mass difference $m_{\Delta} - m_N$ for the

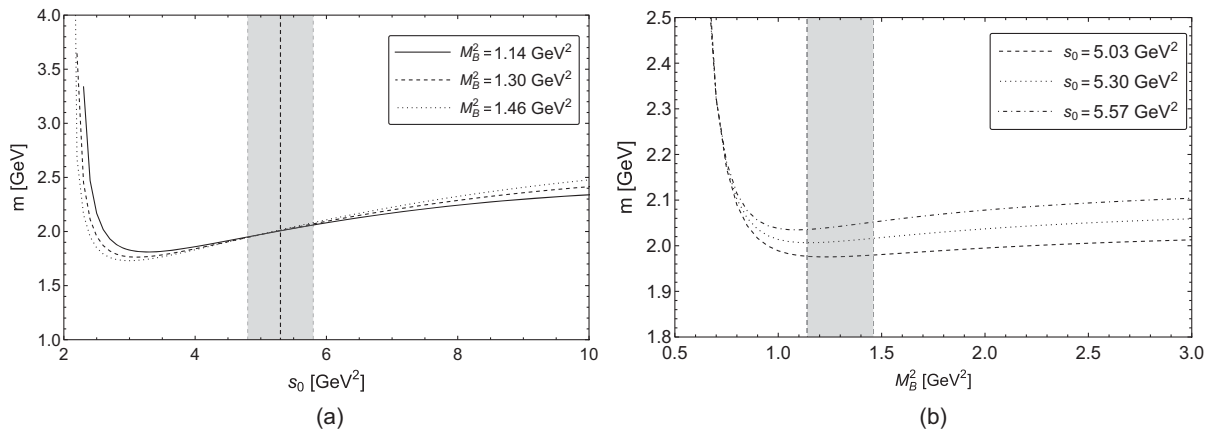


FIG. 2. Variations of hadron mass with (a) continuum threshold s_0 and (b) Borel parameter M_B^2 for the hybrid baryon with $I(J^P) = 1/2(1/2)^+$ from J_7 .

TABLE I. Mass spectra for all hybrid baryons. In the last column, the hadron masses are extracted by omitting the contributions of dimension-8 condensates. The notation “...” means no reliable mass prediction in this channel. The last column gives the results without considering the contributions of dimension-8 condensates.

J_i	N/Δ	s_0 (GeV ²)	M_B^2 (GeV ²)	M_X (GeV)	PC (%)	$\tilde{M}_X$ (GeV)
J_1	$N_{1/2^-}$	$12.4(\pm 5\%)$	$2.75\text{--}2.82$	$3.15^{+0.06}_{-0.06}$	30.9	$3.30^{+0.06}_{-0.06}$
	$N_{1/2^+}$	...	...	...	...	...
J_2	$\Delta_{1/2^-}$	$16.5(\pm 5\%)$	$2.59\text{--}3.60$	$3.64^{+0.09}_{-0.08}$	41.6	$3.63^{+0.09}_{-0.08}$
	$\Delta_{1/2^+}$	...	...	...	...	...
J_3	$N_{1/2^-}$	$7.6(\pm 5\%)$	$1.66\text{--}2.05$	$2.43^{+0.06}_{-0.06}$	37.0	$2.14^{+0.06}_{-0.05}$
	$N_{1/2^+}$	$14.7(\pm 5\%)$	$2.29\text{--}2.93$	$3.50^{+0.17}_{-0.11}$	38.3	$3.29^{+0.19}_{-0.13}$
J_4	$\Delta_{1/2^-}$	$9.5(\pm 5\%)$	$2.43\text{--}2.54$	$2.82^{+0.20}_{-0.12}$	21.1	$3.00^{+0.19}_{-0.13}$
	$\Delta_{1/2^+}$	...	...	...	...	...
J_5	$N_{1/2^-}$	$10.5(\pm 5\%)$	$2.78\text{--}3.05$	$2.92^{+0.05}_{-0.05}$	22.7	$2.77^{+0.05}_{-0.05}$
	$N_{1/2^+}$	...	...	...	...	...
J_6	$\Delta_{1/2^-}$	...	...	...	...	...
	$\Delta_{1/2^+}$	$22.0(\pm 5\%)$	$3.35\text{--}3.70$	$4.18^{+0.10}_{-0.09}$	54.9	$4.31^{+0.10}_{-0.09}$
J_7	$N_{1/2^-}$	$15.0(\pm 5\%)$	$2.56\text{--}3.07$	$3.52^{+0.13}_{-0.10}$	36.1	$3.48^{+0.14}_{-0.11}$
	$N_{1/2^+}$	$5.3(\pm 5\%)$	$1.14\text{--}1.46$	$2.01^{+0.12}_{-0.08}$	38.7	...
J_8	$\Delta_{1/2^-}$	$13.8(\pm 5\%)$	$2.06\text{--}2.63$	$3.24^{+0.07}_{-0.06}$	59.5	$3.36^{+0.07}_{-0.07}$
	$\Delta_{1/2^+}$	...	...	...	...	...
$J_{1\mu}$	$N_{3/2^-}$	$17.6(\pm 5\%)$	$3.24\text{--}3.94$	$3.75^{+0.11}_{-0.09}$	36.9	$3.75^{+0.11}_{-0.09}$
	$N_{3/2^+}$	...	...	...	...	...
$J_{2\mu}$	$\Delta_{3/2^-}$	...	...	...	...	...
	$\Delta_{3/2^+}$	$15.5(\pm 5\%)$	$2.52\text{--}2.87$	$3.45^{+0.06}_{-0.06}$	55.5	$3.58^{+0.07}_{-0.07}$
$J_{3\mu}$	$N_{3/2^-}$	$13.1(\pm 5\%)$	$2.96\text{--}3.21$	$3.23^{+0.08}_{-0.07}$	32.8	$3.11^{+0.08}_{-0.07}$
	$N_{3/2^+}$	...	...	...	...	...
$J_{4\mu}$	$\Delta_{3/2^-}$	...	...	...	...	$2.68^{+0.14}_{-0.11}$
	$\Delta_{3/2^+}$	...	...	...	...	$3.49^{+0.10}_{-0.09}$

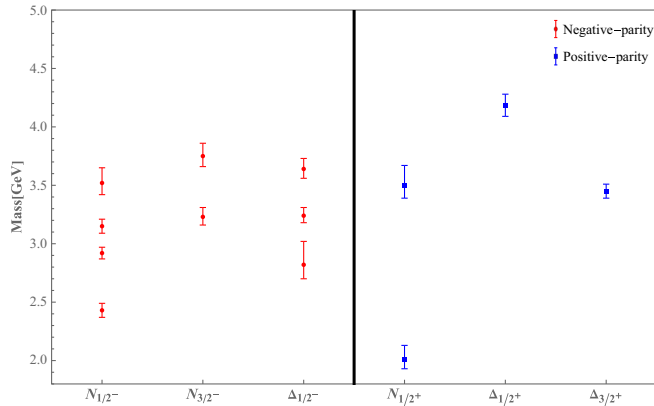


FIG. 3. Mass spectra for both negative-parity and positive-parity hybrid baryons.

conventional baryons. For the other channels, the hybrid baryons are predicted to be above 3 GeV. Considering the renormalization scale at $\mu = 2$ GeV, the predicted hybrid masses will be slightly increased without changing the overall spectrum structure. One may wonder that the traditional baryon states with the same quantum numbers also contribute to the correlation functions. Following the analysis of Ref. [34], we note that mixing with (lighter) conventional baryons that are weakly coupled to the hybrid currents can only raise the hybrid mass predictions of Table I.

The accuracy of the factorization hypothesis for high-dimensional condensates in QCD sum rules remains unestablished. To quantify this uncertainty, we reanalyzed the mass sum rules omitting dimension-8 condensates and summarized the results in the last column of

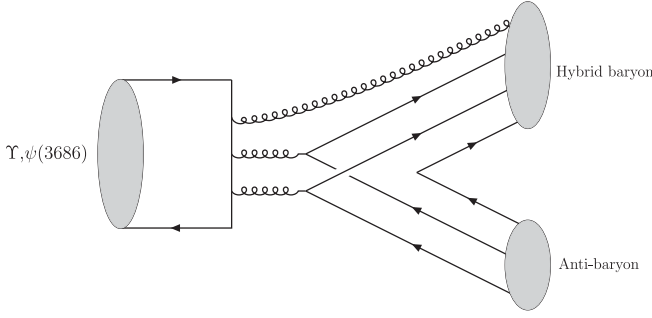


FIG. 4. Possible production mechanisms of light hybrid baryons from $\Upsilon/\psi(3686)$ decays.

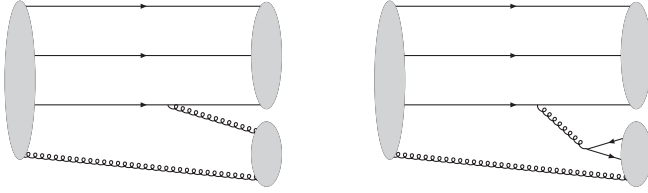


FIG. 5. Possible decay mechanisms of light hybrid baryon to gluon-rich channels.

Table I. For most channels, the $D = 8$ contributions shift the predicted hadron mass by less than 10%, consistent with findings in hybrid meson studies [73]. Significant deviations occur, however, in the $N_{1/2^+}$ channel of J_7 and $\Delta_{3/2^+}$ channels of $J_{4\mu}$, where the stability of the sum rules is qualitatively altered by the inclusion of these condensates.

IV. PRODUCTION AND DECAY MECHANISMS

The decay of charmonium and bottomonium mesons can provide an ideal platform for producing gluon-rich hadron states. For instance, the well known 1^{-+} hybrid candidate $\eta_1(1855)$ was observed in the radiative decays of J/ψ

[47,78]. Similarly, the charmonium/bottomonium decay processes may produce light hybrid baryons, if kinematically allowed. According to our results, all predicted hybrid baryons can be generated through the bottomonium decays via the three-gluon emission processes, such as $\Upsilon \rightarrow \bar{N}N_{hyb}/\bar{\Delta}\Delta_{hyb} + c.c.$ Besides, the lightest $N_{1/2^+}$ hybrid state can also be generated through the charmonium decays $\psi(3686) \rightarrow \bar{N}N_{1/2^+} + c.c.$ We propose such hybrid baryon production mechanisms in Fig. 4.

Carrying the same quantum numbers, the hybrid baryons could always mix with traditional baryon states, so their decay channels to the conventional baryon plus meson are similar to the latter ones, as shown in Table II. Moreover, a hybrid baryon may decay into final states containing gluon-rich hadrons such as hybrid mesons or glueballs, as depicted in Fig. 5. Such characteristic decays may be useful to distinguish hybrid baryons from conventional states and pentaquark states.

V. SUMMARY

We calculate the masses of the light hybrid baryons with quantum numbers $I(J^P) = 1/2(1/2^\pm), 3/2(1/2^\pm), 1/2(3/2^\pm)$, and $3/2(3/2^\pm)$ in parity-projected QCD sum rule method by constructing local interpolating currents. The correlation functions and spectral densities are evaluated up to dimension-8 condensates. We establish stable mass sum rules for the positive-parity $N_{1/2^+}, \Delta_{3/2^+}, \Delta_{1/2^+}$ and negative-parity $N_{1/2^-}, N_{3/2^-}, \Delta_{1/2^-}$ hybrid baryons, so the mass predictions in these channels are reliable. The lightest hybrid baryon is predicted as the positive-parity states $N_{1/2^+}$ around 2.01 GeV, while the negative-parity states are heavier, which is consistent with the LQCD's prediction [56]. Such mass spectra may suggest the same transverse electric gluon excitation with $J_g^{P_g} = 1^+$ in both hybrid baryon and meson systems [56]. All other hybrids in our results are predicted to

TABLE II. Some possible decay modes into conventional baryon-meson final states for light hybrid baryons.

Current	Hybrid	Decay modes
J_1	$N_{1/2^-}$	$N\pi/\eta^{(\prime)}/\rho/\omega, N(1440)\pi/\eta'/\rho/\omega, \Delta\rho, \Delta(1600)\rho, N(1535)f_0(980)/a_0(980)$
J_2	$\Delta_{1/2^-}$	$N\pi/\rho, N(1440)\pi/\rho, \Delta(1232)\rho, N(1535)a_0(980)$
J_3	$N_{1/2^-}$	$N\pi/\eta^{(\prime)}/\rho/\omega, N(1440)\pi/\eta^{(\prime)}, \Delta(1232)\rho$
	$N_{1/2^+}$	$Nf_0(980)/a_0(980), \Delta(1232)\pi/\rho, N(1440)f_0(980)/a_0(980)$
J_4	$\Delta_{1/2^-}$	$N\pi/\rho, N(1440)\pi/\rho, \Delta(1232)\rho, N(1535)a_0(980)$
J_5	$N_{1/2^-}$	$N\pi/\eta^{(\prime)}/\rho/\omega, N(1440)\pi/\eta^{(\prime)}/\rho/\omega, \Delta(1232)\rho, \Delta(1600)\rho, N(1535)f_0(980)/a_0(980)$
J_6	$\Delta_{1/2^+}$	$Na_0(980), N(1440)a_0(980), \Delta(1232)a_0(980)f_0(980)$
J_7	$N_{1/2^-}$	$N\pi/\eta^{(\prime)}/\rho/\omega, N(1440)\pi/\eta'/\rho/\omega, \Delta(1232)\rho, \Delta(1600)\rho, N(1535)f_0(980)/a_0(980)$
	$N_{1/2^+}$	$Nf_0(980)/a_0(980), \Delta(1232)\pi/\rho$
J_8	$\Delta_{1/2^-}$	$N\pi/\rho, N(1440)\pi/\rho, \Delta(1232)\rho, N(1535)a_0(980)$
$J_{1\mu}$	$N_{3/2^-}$	$N\rho/\omega, N(1400)\rho/\omega, \Delta(1232)\rho$
$J_{2\mu}$	$\Delta_{3/2^+}$	$Nb_1/\rho, N(1440)b_1/\rho, \Delta(1232)b_1/h_1/\rho$
$J_{3\mu}$	$N_{3/2^-}$	$N\rho/\omega, N(1400)\rho/\omega, \Delta(1232)\rho$

be above 3 GeV. These hybrid baryons may mainly decay into the conventional baryon-meson final states. We suggest hunting for such hybrids through the gluon-rich processes in BESIII and BelleII experiments.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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