S-wave $D^{(*)}N$ molecular states: $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$?

Jian-Rong Zhang

Department of Physics, College of Science, National University of Defense Technology, Changsha 410073, Hunan, People’s Republic of China

Theoretically, some works have proposed the hadronic resonances $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ to be S-wave $DN$ and $D^*N$ molecular candidates, respectively. In the framework of QCD sum rules, we investigate that whether $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ could be explained as the S-wave $DN$ state with $J^P = \frac{1}{2}^-$ and the S-wave $D^*N$ state with $J^P = \frac{3}{2}^-$, respectively. Technically, contributions of operators up to dimension 12 are included in the operator product expansion (OPE). The final results are 3.64±0.33 GeV and 3.73±0.35 GeV for the S-wave $DN$ state of $J^P = \frac{1}{2}^-$ and the S-wave $D^*N$ state of $J^P = \frac{3}{2}^-$, respectively. They are somewhat bigger than the experimental data of $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$, respectively. In view of that corresponding molecular currents are constructed from local operators of hadrons, the possibility of $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ as molecular states can not be arbitrarily excluded merely from these disagreements between molecular masses using local currents and experimental data. But then these results imply that $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ could not be compact states. This may suggest a limitation of the QCD sum rule using the local current to determine whether some state is a molecular state or not. As byproducts, masses for their bottom partners are predicted to be 6.97±0.34 GeV for the S-wave $B\Lambda$ state of $J^P = \frac{1}{2}^-$ and 6.98±0.34 GeV for the S-wave $\bar{B}^*\Lambda$ state of $J^P = \frac{5}{2}^-$. 

PACS numbers: 11.55.Hx, 12.38.Lg, 12.39.Mk

I. INTRODUCTION

In the past several years, many new excited charmed baryonic states have been discovered experimentally. For example, Belle Collaboration observed an isotriplet of new states $\Sigma_c(2800)$ decaying into $\Lambda_c^0\pi$, and they tentatively identified the quantum numbers of these states as $J^P = \frac{3}{2}^-$. In particular, the neutral state $\Sigma_c(2800)^0$ was possibly confirmed in the $B^- \rightarrow \Sigma_c(2800)^0\bar{p}$ channel by Babar Collaboration [2]. However, the measured mass 2846±8±10 MeV for $\Sigma_c(2800)^0$ is 3σ higher (assuming Gaussian statistics) than the Belle’s measured value, and Babar indicated that there is weak evidence that the excited $\Sigma_c^0$ observed by them is $J = \frac{1}{2}$. Moreover, Babar collaboration reported the observation of a new charmed state $\Lambda_c(2940)^+$ decaying to $D^0\bar{p}$ with a mass of 2939.8±1.3±1.0 MeV and an intrinsic width of 17.5±5.2±5.9 MeV [2]. Subsequently, Belle Collaboration confirmed it in the $\Lambda_c(2940)^+ \rightarrow \Sigma_c(2455)^{0,+}\pi^+\pi^-$ decay and measured its mass and width to be $2938.0\pm1.3\pm0.0$ MeV and $13\pm5\pm7$ MeV, respectively [4].

The experimental observations have triggered theorists’ great interest in understanding the internal structures of $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$. One direct way of theoretical studies grounds on the assignments of them as conventional charmed baryons. From a relativized potential model prediction, masses of $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ are close to theoretical values of $\Sigma_c^*$ with $J^P = \frac{1}{2}^-$ or $\frac{3}{2}^-$ and $\Lambda_c^*$ with $J^P = \frac{1}{2}^-$ or $\frac{3}{2}^-$, respectively [2]. In the relativistic quark-diquark picture, Ebert et al. suggested $\Sigma_c(2800)$ as one of the orbital (1P) excitations of the $\Sigma_c$ with $J^P = \frac{1}{2}^-$, $\frac{3}{2}^-$, or $\frac{5}{2}^-$, and proposed $\Lambda_c(2940)^+$ as the first radial excitation of $\Sigma_c$ with $J^P = \frac{3}{2}^+$. Strong decays of $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ as charmed baryons were analyzed in heavy baryon chiral perturbation theory [2], with the $3P_0$ model [3], and using the chiral quark model [4]. In Ref. [10], Garcilazo et al. indicated that $\Sigma_c(2800)$ would correspond to an orbital excitation with $J^P = \frac{1}{2}^-$ or $\frac{3}{2}^-$ and $\Lambda_c(2940)^+$ may constitute the second orbital excitation of the $\Lambda_c$ by the Faddeev method. In a mass loaded flux tube model, Chen et al. suggested that $\Lambda_c(2940)^+$ could be the orbitally excited $\Lambda_c^*$ with $J^P = \frac{5}{2}^-$. He et al. evaluated the production rate of $\Lambda_c(2940)^+$ as a charmed baryon at PANDA [11].

Another different way bases on the assumption that $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ are some molecular can-
candidates. Lutz et al. interpreted Σc(2800) as a chiral molecule \[13\]. In Ref. \[14\], Σc(2800) was deciphered as a dynamically generated resonance with a dominant D\*N configuration. Dong et al. pursued a possible hadronic molecule interpretation of Σc(2800) as a bound state of the charmed D meson and the nucleon N, since isospectral states of Σc(2800) are very close to respective D\*N thresholds \[15\]. They chose several possible quantum number assignments of Σc(2800) as \(J^P = \frac{1}{2}^+\) and \(\frac{3}{2}^+\). Here \(J^P = \frac{1}{2}^+\) corresponds to a S-wave D\*N configuration, \(J^P = \frac{1}{2}^+\) and \(\frac{3}{2}^+\) represent a P-wave, and \(J^P = \frac{3}{2}^-\) has a relative D-wave in the D\*N system. They finally concluded that Σc → Λcπ decay widths are consistent with current data for \(J^P = \frac{1}{2}^+\) and \(J^P = \frac{3}{2}^-\) assignments. Coming down to \(\Lambda_c(2940)^+\), it was firstly proposed to be a S-wave D\*0p molecular state with \(J^P = \frac{1}{2}^-\) in Ref. \[16\], because its mass is just a few MeV below the D\*0p threshold. From an effective Lagrangian approach, the strong two-body decay of \(\Lambda_c(2940)^+\) was studied under the \(J^P = \frac{1}{2}^-\) and \(\frac{1}{2}^+\) D\*N molecular assignments, and it showed that \(J^P = \frac{1}{2}^-\) should be ruled out \[17\]. Later, the radiative and strong three-body decays of \(\Lambda_c(2940)^+\) were also researched in the D\*N molecule picture with \(J^P = \frac{1}{2}^+\) \[18, 19\]. He et al. systematically studied the interaction between \(D^*\) and N, and concluded that the \(D^*N\) systems may behave as \(J^P = \frac{1}{2}^\pm\) and \(\frac{3}{2}^\pm\) states \[20\]. In Ref. \[21\], García-Recio et al. found a possible molecular candidate for the \(\Lambda_c(2940)^+\) in the \(\frac{1}{2}^-\) channel. Ortega et al. studied \(\Lambda_c(2940)^+\) as a \(D^*N\) molecule with \(J^P = \frac{3}{2}^-\) in a constituent quark model, and claimed obtaining a mass which agrees with the experimental data \[22\].

Although various interpretations to Σc(2800) and \(\Lambda_c(2940)^+\) were put forward, at present their underlying structures are still unclear, which means that it is interesting and significative to make more theoretical efforts to reveal their properties. Therefore, we devote to studying that whether Σc(2800) and \(\Lambda_c(2940)^+\) could be the S-wave D\*N state with \(J^P = \frac{1}{2}^-\) and the S-wave D\*N state with \(J^P = \frac{3}{2}^-\), respectively. QCD is widely believed nowadays to be a true theory of strong interactions. At high energy, the effective coupling constant of the quark-gluon interaction becomes small because of asymptotic freedom and the interaction can be treated perturbatively. On the other hand, quark interaction within hadrons is strong since it binds quarks into unseparable pairs. Thus, low energy QCD involves a regime where it is futile to attempt perturbative calculations, and the strong interaction dynamics of hadronic systems is governed by nonperturbative QCD effects completely. There are still many questions on nonperturbative QCD remain unanswered or realized only at a qualitative level since one’s absence of knowledge on QCD confinement effects. Therefore, it is quite difficult to calculate the hadron spectrum from QCD first principles. The method of QCD sum rules \[23\], developed by Shifman, Vainshtein, and Zakharov, represents an attempt to bridge the gap between the perturbative and nonperturbative sectors by employing the language of dispersion relations. It is well known for the advantages of this method: instead of a model-dependent treatment in terms of constituent quarks, hadrons are represented by their interpolating quark currents and the interactions of quark-gluon currents with QCD vacuum fields critically depend on the quantum numbers (spin-parity, flavor content) of these currents. The QCD sum rule method is a nonperturbative formulation firmly based on the first principle of QCD, which has become a widely used working tool in hadron phenomenology. The mere fact that the seminal paper on QCD sum rules have already been cited more than 4000 times reflects its vigorously. It has been successfully applied to conventional mesons and baryons (for reviews see \[24–27\] and references therein) and multiquark states (e.g. see \[28\]). In particular, many theoretical practitioners began to study light pentaquark states in Refs. \[29, 30\] and heavy pentaquark systems in Ref. \[31\].

In this work, we intend to investigate that whether Σc(2800) and \(\Lambda_c(2940)^+\) could be the S-wave D\*N state with \(J^P = \frac{1}{2}^-\) and the S-wave D\*N state with \(J^P = \frac{3}{2}^-\) respectively from QCD sum rules. The rest of the paper is organized as follow. We derive QCD sum rules for molecular states in Sec. \[III\] with similar techniques as our previous works on heavy baryons \[32\] and molecular states \[33\]. The numerical analysis and discussions are presented in Sec. \[IV\] and masses of \(D^{(*)}N\) and \(B^{(*)}N\) molecular states are extracted out. Sec. \[IV\] contributes to the conclusions.
II. QCD SUM RULES FOR MESON-NUCLEON MOLECULAR STATES

A. constructions of interpolating currents

One basic point of QCD sum rules is to construct a proper interpolating current to represent the studied state. In the real world, one hadron in particular a molecular state can not be an ideal point particle in a rigorous manner because each constituent quark of a hadronic system is separated in the space. Without doubt, it would be best if one could describe a real hadron using some nonlocal current in QCD sum rules. However, the practitioners can find that it would become quite difficult or even unfeasible for QCD sum rule calculations when a hadron’s current is constructed nonlocal. Thus, interpolating currents used in QCD sum rules are commonly built local to characterize real hadrons, which is in fact a limitation inherent in the QCD sum rule method disposal of hadrons. The simplification has been widely proved feasible and the QCD sum rule method has been successfully applied to plenty of hadrons, involving a number of works on molecular states since the observations of so-called “X”, “Y”, and “Z” new hadrons in recent years (e.g. see [28] and references therein). Following the usual treatment, in this work we will construct molecular currents from local operators of hadrons. At present, molecular currents are built up with the color-singlet currents of composed hadrons to form hadron-hadron configurations of fields. Although molecular currents can be related to pentaquark currents by Fierz rearrangement, the former are different from currents of pentaquark states constructed by diquark-diquark-antiquark configurations of fields.

Therefore, currents for $S$-wave $D^{(*)}N$ or $B^{(*)}N$ molecular states can be built up with the color-singlet currents of $D^{(*)}$ or $B^{(*)}$ mesons and $N$ nucleons to form meson-nucleon configurations of fields. In full theory, interpolating currents for $D^{(*)}$ and $B^{(*)}$ mesons can be found in Ref. [34], and currents for nucleons have been listed in Ref. [35]. Therefore, we build following forms of currents:

$$j = (q^{c'} i\gamma_5 Q^{c'}) (\varepsilon_{abc} q^{T_a}_1 C^{\gamma_5} q^{b'}_2 \gamma_5 \gamma^\mu q^{d}_1),$$

for the $S$-wave $DN$ or $\bar{B}N$ molecular state with $J^P = \frac{1}{2}^-$, and

$$j^{p} = (q^{c'} \gamma^\rho Q^{c'}) (\varepsilon_{abc} q^{T_a}_1 C^{\gamma_5} q^{b'}_2 \gamma_5 \gamma^\mu q^{d}_1),$$

for the $S$-wave $D^*N$ or $B^*N$ molecular state with $J^P = \frac{3}{2}^-$. Here $Q$ is heavy quark $c$ or $b$, and $q_1$, $q_2$, as well as $q_3$ denote light quarks $u$ and/or $d$. The index $T$ means matrix transposition, $C$ is the charge conjugation matrix, with $a$, $b$, $c$ and $c'$ as color indices. Beside $J^P = \frac{3}{2}^-$, one may have noted that the quantum number for a $S$-wave $D^*N$ molecule could also be $\frac{1}{2}^-$. However, it is not straightforward to construct the interpolating current for the $S$-wave $D^*N$ molecule with $J^P = \frac{1}{2}^-$ from meson-nucleon configurations of fields. That’s the main reason why the $S$-wave $D^*N$ molecule with $J^P = \frac{1}{2}^-$ has not been involved here. In addition, it showed that the molecular assignment of $\Lambda_c(2940)^+$ with $J^P = \frac{1}{2}^-$ should be ruled out from an effective Lagrangian approach [17].

B. QCD sum rules for meson-nucleon molecular states

QCD sum rules for $DN$ and $\bar{B}N$ molecular states with $J^P = \frac{1}{2}^-$ are constructed from the two-point correlator

$$\Pi(q^2) = i \int d^4x e^{iq\cdot x} \langle 0 | T[j(x)\bar{j}(0)] | 0 \rangle.$$  

(3)

Lorentz covariance hints that the correlator has the form

$$\Pi(q^2) = \#\Pi_1(q^2) + \Pi_2(q^2).$$

(4)
Phenomenologically, the correlator can be expressed as

\[ \Pi(q^2) = \lambda_H^2 \frac{q - M_H}{M_H - q^2} + \frac{1}{\pi} \int_{s_0}^\infty ds \frac{\text{Im} \Pi_1^{\text{Phen}}(s) + \text{Im} \Pi_2^{\text{Phen}}(s)}{s - q^2} + \ldots, \]  

(5)

where \( M_H \) is the mass of the hadronic resonance, \( s_0 \) is the threshold parameter, and \( \lambda_H \) gives the coupling of the current to the hadron \( \langle 0 | j | H \rangle = \lambda_H v(q, s) \). In the OPE side, the correlator can be written as

\[ \Pi(q^2) = \int_{m_H^2}^\infty ds \rho_1(s) \frac{1}{s - q^2} + \Pi_1^{\text{cond}}(q^2) = \int_{m_H^2}^\infty ds \rho_2(s) \frac{1}{s - q^2} + \Pi_2^{\text{cond}}(q^2), \]

(6)

After equating the two sides for \( \Pi(q^2) \), assuming quark-hadron duality, making a Borel transform, and transferring the continuum contribution to the OPE side, the sum rules can be written as

\[ \lambda_H^2 e^{-M_H^2/M^2} = \int_{m_H^2}^{s_0} ds \rho_1(s) e^{-s/M^2} + \tilde{B} \Pi_1^{\text{cond}}, \]

(7)

\[ - \lambda_H^2 M_H e^{-M_H^2/M^2} = \int_{m_H^2}^{s_0} ds \rho_2(s) e^{-s/M^2} + \tilde{B} \Pi_2^{\text{cond}}, \]

(8)

where \( M^2 \) indicates the Borel parameter.

There is some difference while deriving mass sum rules for \( D^*N \) and \( B^*N \) molecular states with \( J^P = \frac{3}{2}^- \).

One can start from the two-point correlator

\[ \Pi^{\rho\tau}(q^2) = i \int d^4x e^{iq\cdot x} \langle 0 | T[j^\rho(x) j^{\rho\tau}(0)] | 0 \rangle. \]

(9)

Lorentz covariance implies that the two-point correlator in Eq. (9) has the form

\[ \Pi^{\rho\tau}(q^2) = -g^{\rho\tau}\l[ \Pi_1(q^2) + \Pi_2(q^2) \r] + \ldots, \]

(10)

where the ellipse denotes other Lorentz structures which acquire contributions from both \( J = \frac{1}{2} \) and \( J = \frac{3}{2} \). The tensor structures \( g^{\rho\tau} \partial \) and \( g^{\rho\tau} \) are contributed only by the \( J = \frac{3}{2} \) hadrons. The phenomenological side of \( \Pi^{\rho\tau}(q^2) \) can be expressed as

\[ \Pi^{\rho\tau}(q^2) = -g^{\rho\tau}\l\{ \lambda_H^2 \frac{q - M_H}{M_H - q^2} + \frac{1}{\pi} \int_{s_0}^\infty ds \frac{\text{Im} \Pi_1^{\text{Phen}}(s) + \text{Im} \Pi_2^{\text{Phen}}(s)}{s - q^2} \r\} + \ldots, \]

(11)

where \( \lambda_H \) gives the coupling of the hadronic state to the current \( j^\rho \) as \( \langle 0 | j^\rho | H \rangle = \lambda_H v^\rho(q, s) \). Here, \( v^\rho(q, s) \) is the Rarita-Schwinger spinor for \( J^P = \frac{3}{2}^- \). In the OPE side, one can write the correlator as

\[ \Pi^{\rho\tau}(q^2) = -g^{\rho\tau}\l\{ \frac{q}{M_H - q^2} \l[ \int_{m_H^2}^\infty ds \rho_1(s) + \Pi_1^{\text{cond}}(q^2) \r] + \int_{m_H^2}^\infty ds \frac{\rho_2(s)}{s - q^2} + \Pi_2^{\text{cond}}(q^2) \r\}, \]

(12)

Equating the two sides for \( \Pi^{\rho\tau}(q^2) \), assuming quark-hadron duality, and making a Borel transform, the sum rules can be expressed as

\[ \lambda_H^2 e^{-M_H^2/M^2} = \int_{m_H^2}^{s_0} ds \rho_1(s) e^{-s/M^2} + \tilde{B} \Pi_1^{\text{cond}}, \]

(13)

\[ - \lambda_H^2 M_H e^{-M_H^2/M^2} = \int_{m_H^2}^{s_0} ds \rho_2(s) e^{-s/M^2} + \tilde{B} \Pi_2^{\text{cond}}. \]

(14)
To eliminate the hadron coupling constant $\lambda_H$ in sum rules (7), (8), (13), and (14), one can take the derivatives of sum rules with respect to $1/M^2$, divide the results by themselves to get

$$M_H^2 = \left\{ \int_{m_Q^2}^{\infty} ds \rho_1(s) e^{-s/M^2} + d/d(-\frac{1}{M^2}) \tilde{B}\Pi_1^{\text{cond}} \right\} / \left\{ \int_{m_Q^2}^{\infty} ds \rho_1(s) e^{-s/M^2} + \tilde{B}\Pi_1^{\text{cond}} \right\},$$

$$M_H^2 = \left\{ \int_{m_Q^2}^{\infty} ds \rho_2(s) e^{-s/M^2} + d/d(-\frac{1}{M^2}) \tilde{B}\Pi_2^{\text{cond}} \right\} / \left\{ \int_{m_Q^2}^{\infty} ds \rho_2(s) e^{-s/M^2} + \tilde{B}\Pi_2^{\text{cond}} \right\}. \quad (15)$$

$$M_H^2 = \left\{ \int_{m_Q^2}^{\infty} ds \rho_3(s) e^{-s/M^2} + d/d(-\frac{1}{M^2}) \tilde{B}\Pi_3^{\text{cond}} \right\} / \left\{ \int_{m_Q^2}^{\infty} ds \rho_3(s) e^{-s/M^2} + \tilde{B}\Pi_3^{\text{cond}} \right\}. \quad (16)$$

### C. Spectral Densities

The spectral density is given by the correlator's imaginary part

$$\rho_i(s) = \frac{1}{\pi} \text{Im} \Pi_i^{\text{OPE}}(s), \quad i = 1, 2.$$ \quad (17)

In the concrete OPE calculation, one works at leading order in $\alpha_s$ and considers condensates up to dimension 12. Note that $O(\alpha_s)$ corrections may be important in the QCD sum rule calculations. Meanwhile, one could expect the calculations of $O(\alpha_s)$ corrections especially for multiquark systems are complicated and tedious as one has to deal with many multi-loop massive propagator diagrams. Actually, a lot of hard calculations already need to be done even if one merely works at leading order since there include many Feynman diagrams up to dimension 12. However, it is expected that the $O(\alpha_s)$ corrections might be under control since a partial cancelation occurs in the ratio obtaining the mass sum rules (14) and (16). This has been proved to be true in the analysis for heavy mesons (the value of $f_D$ increases by 12% after the inclusion of the $O(\alpha_s)$ correction) and singly heavy baryons (the corrections increase the calculated baryon masses by about 10%). Furthermore, in order to improve on the accuracy of the QCD sum rule analysis for molecular states, one could take into account the $O(\alpha_s)$ corrections in the further work after fulfilling a burdensome task. To keep the heavy-quark mass finite, one can use the momentum-space expression for the heavy-quark propagator

$$S_Q(p) = \frac{i}{\not{p} - m_Q} - \frac{i}{4} g t A G_{A\alpha}(0) \frac{1}{(p^2 - m_Q^2)^2} \left[ \sigma_{\alpha\lambda} (\not{p} + m_Q) + (\not{p} + m_Q) \sigma_{\lambda\alpha} \right]$$

$$- \frac{i}{4} g^2 t A t B G^{A\alpha\beta}(0) G^B_{\mu\nu}(0) \frac{\not{p} + m_Q}{(p^2 - m_Q^2)^2} \gamma_\mu \gamma_\nu$$

$$+ \gamma_\alpha (\not{p} + m_Q) \gamma_\mu (\not{p} + m_Q) \gamma_\nu + \gamma_\alpha (\not{p} + m_Q) \gamma_\mu (\not{p} + m_Q) \gamma_\nu (\not{p} + m_Q) \gamma_\beta (\not{p} + m_Q)$$

$$+ \frac{i}{48} g^3 f_{ABC} G^{A\alpha}_{\mu} G^{B\beta}_{\nu} G^{C}_{\epsilon\gamma} \frac{1}{(p^2 - m_Q^2)^2} (\not{p} + m_Q) [\not{q} (p^2 - 3m_Q^2) + 2m_Q (2p^2 - m_Q^2)] (\not{p} + m_Q). \quad (18)$$

The light-quark part of the correlator can be calculated in the coordinate space employing the light-quark propagator

$$S_{ab}(x) = \frac{i \delta_{ab}}{2\pi^2 x^2} - \frac{m_q \delta_{ab}}{4\pi^2 x^2} t A A_{\mu\nu}(p^2 + \sigma_{\mu\nu}) - \frac{\delta_{ab}}{12} \not{q} \not{q} - \frac{i \delta_{ab}}{48} m_q \not{q}$$

$$- \frac{x^2 \delta_{ab}}{3 \cdot 2^6} (\not{q} \not{q} \not{q} - G_q) + \frac{i x^2 \delta_{ab}}{2^7 \cdot 3^2} m_q (\not{q} \not{q} \not{q} - G_q) - \frac{x^4 \delta_{ab}}{2^10 \cdot 3^4} (\not{q} \not{q} \not{q} \not{q}) (g^2 Q^2), \quad (19)$$

which is then Fourier-transformed to the momentum space in $D$ dimension. The resulting light-quark part is combined with the heavy-quark part before it is dimensionally regularized at $D = 4$.

The spectral densities can be listed as

$$\rho_i(s) = \rho_i^{\text{pert}}(s) + \rho_i^{(\bar{q}q)}(s) + \rho_i^{(\bar{q}q)^2}(s) + \rho_i^{(q\bar{q}g\sigma G_q)}(s) + \rho_i^{(qg^2G^2)}(s) + \rho_i^{(g^3G^2)}(s) + \rho_i^{(g^4)}(s) + \rho_i^{(\bar{q}q)(g\bar{q}g - G_q)}(s)$$

$$+ \rho_i^{(\bar{q}q) Q_g(q g^2 G^2)}(s) + \rho_i^{(\bar{q}q)(g^3 G^3)}(s) + \rho_i^{(\bar{q}q)(g^4 G^4)}(s) + \rho_i^{(g^5 G^5)}(s), \quad i = 1, 2. \quad (20)$$
Because many terms of $\rho_2(s)$ are proportional to light quarks' masses and approximate to zero, we merely present spectral densities resulted from $\Pi_1(q^2)$. Concretely, they read

$$
\rho_1^{(\bar{q}q)}(s) = -\frac{m_Q(q\bar{q})}{3 \cdot 211_{\pi}} \int \frac{1}{\alpha} \frac{1}{\alpha^2} (\alpha s - m_Q^2) (\alpha s + m_Q^2),
$$

$$
\rho_1^{(\bar{q}q)}(s) = \frac{m_Q(q\bar{q} \cdot Gq)}{211_{\pi}} \int \frac{1}{\alpha} \frac{1}{\alpha^2} (\alpha s - m_Q^2)^2,
$$

$$
\rho_1^{(\bar{q}q)}(s) = \frac{m_Q^2(q\bar{q} \cdot Gq)}{5 \cdot 211_{\pi}} \int \frac{1}{\alpha} \frac{1}{\alpha^2} (\alpha s - m_Q^2)^2 (\alpha s - 2m_Q^2),
$$

$$
\rho_1^{(\bar{q}q)}(s) = \frac{m_Q^3(q\bar{q} \cdot Gq)}{5 \cdot 211_{\pi}} \int \frac{1}{\alpha} \frac{1}{\alpha^2} (\alpha s - m_Q^2)^2 [((\alpha s)^2 - 9\alpha s m_Q + 10m_Q^4],
$$

$$
\rho_1^{(\bar{q}q)}(s) = \frac{m_Q^4(q\bar{q} \cdot Gq)}{5 \cdot 211_{\pi}} \int \frac{1}{\alpha} \frac{1}{\alpha^2} (\alpha s - m_Q^2)^2 [((\alpha s)^2 - 9\alpha s m_Q + 10m_Q^4],
$$

for the $S$-wave $DN$ or $BN$ state with $J^P = \frac{1}{2}^-$, and

$$
\rho_1^{pert}(s) = \frac{1}{3 \cdot 5^2 \cdot 211_{\pi}} \int \frac{1}{\alpha} \frac{1}{\alpha^2} (\alpha s - m_Q^2)^4 [\alpha s + 4m_Q^2 - \alpha (\alpha s - m_Q^2)],
$$
\[
\rho_1^{(\bar{q}q)}(s) = -\frac{m_Q}{3 \cdot 2^{11} \pi^6} \int_\Lambda^1 d\alpha \left(1 - \alpha\right)^4 \frac{(as - m_Q^2)^3}{\alpha^3}, \\
\rho_1^{(\bar{q}q)^2}(s) = \frac{(\bar{q}q)^2}{3 \cdot 2^6 \pi^6} \int_\Lambda^1 d\alpha \left(1 - \alpha\right)^3 \frac{(as - m_Q^2)((as + m_Q^2) - \alpha(as - m_Q^2))}{\alpha^2}, \\
\rho_1^{(g\bar{q}Gq)}(s) = \frac{m_Q(g\bar{q}Gq)}{2^{11} \pi^6} \int_\Lambda^1 d\alpha \left(1 - \alpha\right)^3 \frac{(as - m_Q^2)^2}{\alpha^2}, \\
\rho_1^{(g^2G^2)}(s) = \frac{m_Q^2(g^2G^2)}{5 \cdot 2^{11} \pi^6} \int_\Lambda^1 d\alpha \frac{(as - m_Q^2)((as + m_Q^2) + 2(as - m_Q^2))}{\alpha^3}, \\
\rho_1^{(g^3G^3)}(s) = \frac{m_Q^3(g^3G^3)}{5 \cdot 2^{15} \pi^6} \int_\Lambda^1 d\alpha \frac{(as - m_Q^2)((as - 5m_Q^2) + 2((as)^2 - 9asm_Q^2 + 10m_Q^4))}{\alpha^5}, \\
\rho_1^{(\bar{q}q)^3}(s) = -\frac{m_Q}{3 \cdot 2^6 \pi^6} \int_\Lambda^1 d\alpha (1 - \alpha), \\
\rho_1^{(\bar{q}q)(g\bar{q}Gq)}(s) = \frac{(\bar{q}q)(g\bar{q}Gq)}{2^{10} \pi^4} \int_\Lambda^1 d\alpha (1 - \alpha)^3 \frac{(as - m_Q^2)}{\alpha}, \\
\rho_1^{(g\bar{q}Gq)(g\bar{q}Gq)}(s) = \frac{(g\bar{q}Gq)}{2^{10} \pi^4} \int_\Lambda^1 d\alpha (1 - \alpha)^2 \frac{(as - m_Q^2)}{\alpha^2}, \\
\rho_1^{(g^2G^2)(g\bar{q}Gq)}(s) = \frac{(g^2G^2)(g\bar{q}Gq)}{3 \cdot 2^{13} \pi^6} \int_\Lambda^1 d\alpha \frac{(1 - \alpha)^3}{\alpha^2}, \\
\hat{B}_{\text{H1 cond}} = \frac{m_Q\bar{q}q^2(g\bar{q}Gq)}{2^{10} \pi^4} \int_0^1 d\alpha e^{-m_Q^2/(\alpha M^2)} + \frac{m_Q^2(g\bar{q}Gq)^2}{2^{10} \pi^4} \int_0^1 d\alpha \frac{1 - \alpha}{\alpha} e^{-m_Q^2/(\alpha M^2)} \\
- \frac{m_Q^2(g\bar{q}Gq)}{3 \cdot 2^{15} \pi^6} \int_\Lambda^1 d\alpha \left(1 - \alpha\right)^2 \frac{(3 - m_Q^2)}{\alpha^2 M^2} e^{-m_Q^2/(\alpha M^2)} \\
- \frac{m_Q^3(g\bar{q}Gq)}{3 \cdot 2^{14} \pi^6} \int_\Lambda^1 d\alpha \left(1 - \alpha\right)^3 \frac{(- (a + 2) + m_Q^2)}{\alpha^4 M^2} e^{-m_Q^2/(\alpha M^2)} \\
- \frac{m_Q^2(g^2G^2)^2}{3 \cdot 2^{10} \pi^4} \int_\Lambda^1 d\alpha \left(1 - \alpha\right)^3 \frac{(- (a + 2) + m_Q^2)}{\alpha^3 M^2} e^{-m_Q^2/(\alpha M^2)} \\
- \frac{m_Q^3(g^3G^3)}{3 \cdot 2^{11} \pi^4} \int_\Lambda^1 d\alpha \left(1 - \alpha\right)^3 \frac{(- (a + 2) + m_Q^2)}{\alpha^4 M^2} e^{-m_Q^2/(\alpha M^2)} \\
- \frac{m_Q^3(g^2G^2)(g\bar{q}Gq)}{3 \cdot 2^{11} \pi^4} \int_\Lambda^1 d\alpha \left(1 - \alpha\right)^3 \frac{(- (a + 2) + m_Q^2)}{\alpha^3 M^2} e^{-m_Q^2/(\alpha M^2)} \\
- \frac{m_Q^3(g^3G^3)(g\bar{q}Gq)}{3 \cdot 2^{15} \pi^6} \int_\Lambda^1 d\alpha \left(1 - \alpha\right)^3 \frac{(- (a + 2) + m_Q^2)}{\alpha^3 M^2} e^{-m_Q^2/(\alpha M^2)},
\]

for the S-wave $D^*N$ or $B^*N$ state with $J^P = \frac{1}{2}^-$. The lower limit of integration is given by $\Lambda = m_Q^2/s$. In the deriving of above results, we have applied the factorization hypothesis of the four quark condensate $\langle \bar{q}q\bar{q}q \rangle = \kappa \langle \bar{q}q \rangle \langle \bar{q}q \rangle$ and have set $\kappa = 1$ following the usual treatment. Numerically, the factor $\kappa$ may have some other value such as 2 or 3.
III. NUMERICAL ANALYSIS AND DISCUSSIONS

We perform numerical analysis of the sum rule \(15\) to extract mass values of studied states. One could take input parameters as \(m_c = 1.23 \pm 0.05\) GeV, \(m_b = 4.24 \pm 0.06\) GeV, \(\langle \bar{q}q \rangle = -(0.23 \pm 0.03)^3\) GeV\(^3\), \(\langle g \bar{q} q \cdot Gq \rangle = m_{q}^3 \langle \bar{q}q \rangle, \langle g^2G^2 \rangle = 0.88\) GeV\(^4\), and \(\langle g^3G^3 \rangle = 0.045\) GeV\(^6\). To choose proper work windows for the threshold \(s_0\) and the Borel parameter \(M^2\), one could consider two rules in the standard QCD approach: on one hand, the perturbative contribution should be larger than condensate contributions to have a good convergence in the OPE side; on the other hand, the pole contribution should be larger than the continuum state contributions in the phenomenological side. Besides the above two restrictions, the threshold parameter \(\sqrt{s_0}\) should not be taken arbitrarily. It is known that the first excitation of studied state defines the size of \(\sqrt{s_0}\), and \(\sqrt{s_0}\) should be higher than the extracted value \(M_H\) of studied state around 0.5 GeV for many hadrons. Taking the case of \(DN\) state as an example, if \(\sqrt{s_0}\) were taken as \(3.2 \sim 3.4\) GeV, one could obtain the mass \(M_H = 3.52 \pm 0.36\) GeV. However, the value of \(\sqrt{s_0} - M_H\) is too small or even minus, which means that the values of continuum threshold \(\sqrt{s_0}\) are taken a bit small and should be increased correspondingly.

However, it may have some difficulty to find a conventional work window critically satisfying all the above rules in this work, which has been discussed in detail for some other cases, e.g. Refs. 38, 41. The main reason is that some condensate contributions are very large, making the standard OPE convergence (i.e. perturbative contribution at least larger than each condensate contribution) to happen only at very large values of \(M^2\). For the case of \(S\)-wave \(DN\) state with \(J^P = \frac{1}{2}^-\) as an example, the comparison between pole and continuum contributions from the sum rule \(15\) for \(\sqrt{s_0} = 4.0\) GeV is shown in the left panel of Fig. 1, and its OPE convergence is shown by comparing the perturbative with other condensate contributions in the right panel. One can see that there are four main condensates (i.e., \(\langle \bar{q}q \rangle, \langle g \bar{q} q \cdot Gq \rangle, \langle \bar{q}q \rangle^2\), and \(\langle \bar{q}q \rangle \langle g \bar{q} q \cdot Gq \rangle\)), and they could cancel each other out to some extent as they have different signs. Besides, most of the other condensates calculated are very small and almost negligible. Thus, one could try releasing the rigid OPE convergence criterion (i.e., that the perturbative contribution should be larger than each condensate contribution) and restrict the ratio of the perturbative to the "total OPE contribution" (the sum of the perturbative and other condensates calculated) to be at least larger than one half or more. What is also very important that we have found that condensates higher than dimension 12 are quite small, and they could not radically influence the character of OPE convergence here. All these factors bring that the OPE convergence is still under control at relatively low values of \(M^2\).

The dependence on Borel parameter \(M^2\) for masses of \(S\)-wave \(DN\) and \(\bar{B}N\) states with \(J^P = \frac{1}{2}^-\) are shown in Fig. 2, for which continuum thresholds are taken as \(\sqrt{s_0} = 3.9 \sim 4.1\) GeV and \(\sqrt{s_0} = 7.4 \sim 7.6\) GeV, respectively. From the Borel curves, one can visually see that there indeed exist stable plateaus. Thus, we choose some transition range \(M^2 = 2.0 \sim 3.0\) GeV\(^2\) as a compromise Borel window and arrive at 3.75 \pm 0.14 GeV for \(DN\) state. Considering the uncertainty rooting in the variation of quark masses and condensates, we gain 3.75 \pm 0.14 \pm 0.08 GeV (the first error reflects the uncertainty due to variation of \(\sqrt{s_0}\) and \(M^2\), and the second error resulted from the variation of QCD parameters) for the \(S\)-wave \(DN\) state with \(J^P = \frac{1}{2}^-\). To investigate the effect of different factorization, we take \(\kappa = 2\) and arrive at 3.56 \pm 0.10 \pm 0.07 GeV from the similar analysis process. Similarly, the result is 3.45 \pm 0.07 \pm 0.07 GeV while \(\kappa = 3\). Finally, we average three results for \(\kappa = 1 \sim 3\) and arrive at the mass value 3.64 \pm 0.33 GeV concisely for the \(S\)-wave \(DN\) state with \(J^P = \frac{1}{2}^-\), which is somewhat higher than the experimental value of \(\Sigma_c(2800)\) even considering the uncertainty of result. For the \(S\)-wave \(\bar{B}N\) state with \(J^P = \frac{1}{2}^-\), we choose some transition range \(M^2 = 4.0 \sim 5.0\) GeV\(^2\) as a compromise Borel window and arrive at 7.06 \pm 0.13 GeV. Considering the uncertainty rooting in the variation of quark masses and condensates, we gain 7.06 \pm 0.13 \pm 0.12 GeV for the \(S\)-wave \(\bar{B}N\) state with \(J^P = \frac{1}{2}^-\). The respective results are 6.91 \pm 0.10 \pm 0.10 GeV and 6.82 \pm 0.09 \pm 0.09 GeV for \(\kappa = 2\) and 3. Averaging three results for \(\kappa = 1 \sim 3\), one can arrive at the final mass value 6.97 \pm 0.34 GeV in a nutshell for the \(S\)-wave \(\bar{B}N\) state with \(J^P = \frac{1}{2}^-\).
for \sqrt{s_0} = 4.0 \text{ GeV} for the S-wave $DN$ state with $J^P = \frac{1}{2}^-$.

In the right panel, the OPE convergence is shown by comparing the perturbative with other condensate contributions from the sum rule [Eq. (15)] for \sqrt{s_0} = 4.0 \text{ GeV} for the S-wave $DN$ state with $J^P = \frac{1}{2}^-$. 

For another example, the comparison between pole and continuum contributions from the sum rule [Eq. (15)] for \sqrt{s_0} = 4.1 \text{ GeV} for the S-wave $D^*N$ state with $J^P = \frac{3}{2}^-$ is shown in the left panel of Fig. 3, and its OPE convergence is shown by comparing the perturbative with other condensate contributions in the right panel. Masses of S-wave $D^*N$ and $\bar{B}^*N$ states with $J^P = \frac{3}{2}^-$ as a function of $M^2$ from sum rule [Eq. (15)] are shown in FIG. 4. Graphically, one can see that there are stable plateaus for the Borel curves. Similarly, we choose some transition range $M^2 = 2.0 \sim 3.0 \text{ GeV}^2$ as a compromise Borel window for $D^*N$ state, and arrive at $3.83 \pm 0.16 \text{ GeV}$. Considering the uncertainty rooting in the variation of quark masses and condensates, we obtain $3.83 \pm 0.16 \pm 0.09 \text{ GeV}$ for the S-wave $D^*N$ state with $J^P = \frac{3}{2}^-$. Taking $\kappa = 2$ and 3, the results are $3.62 \pm 0.09 \pm 0.07 \text{ GeV}$ and $3.52 \pm 0.07 \pm 0.07 \text{ GeV}$, respectively. Averaging three results for $\kappa = 1 \sim 3$, the final result is $3.73 \pm 0.35 \text{ GeV}$ in a concise form for the S-wave $D^*N$ state with $J^P = \frac{3}{2}^-$, which is bigger than the experimental data of $\Lambda_c(2940)^+$ even taking into account the uncertainty. For the S-wave $\bar{B}^*N$ state with $J^P = \frac{3}{2}^-$, we choose a compromise Borel window $M^2 = 4.0 \sim 5.0 \text{ GeV}^2$ and take $\sqrt{s_0} = 7.4 \sim 7.6 \text{ GeV}$. In the work windows, we obtain $7.07 \pm 0.12 \text{ GeV}$ for $\bar{B}^*N$ state. Varying
FIG. 3: In the left panel, the solid line shows the relative pole contribution (the pole contribution divided by the total, pole plus continuum contribution), and the dashed line shows the relative continuum contribution from the sum rule [Eq. (13)] for $\sqrt{s_0} = 4.1$ GeV for the $S$-wave $D^*N$ state with $J^P = \frac{3}{2}^-$. In the right panel, the OPE convergence is shown by comparing the perturbative with other condensate contributions from the sum rule [Eq. (13)] for $\sqrt{s_0} = 4.1$ GeV for the $S$-wave $D^*N$ state with $J^P = \frac{3}{2}^-$. 

FIG. 4: Masses of $S$-wave $D^*N$ and $\bar{B}^*N$ states with $J^P = \frac{3}{2}^-$ as a function of $M^2$ from the sum rule [Eq. (15)] are shown. The continuum thresholds are taken as $\sqrt{s_0} = 4.0 \sim 4.2$ GeV and $\sqrt{s_0} = 7.4 \sim 7.6$ GeV, respectively.

input values of quark masses and condensates, we attain $7.07 \pm 0.12 \pm 0.12$ GeV for the $S$-wave $\bar{B}^*N$ state with $J^P = \frac{3}{2}^-$. The results are $6.92 \pm 0.11 \pm 0.10$ GeV and $6.83 \pm 0.10 \pm 0.09$ GeV for $\kappa = 2$ and $\kappa = 3$, respectively. Making the average of three results for $\kappa = 1 \sim 3$, one could gain $6.98 \pm 0.34$ GeV concisely for the $S$-wave $\bar{B}^*N$ state with $J^P = \frac{3}{2}^-$. 

IV. CONCLUSIONS

In some theoretical approaches, the hadronic resonances $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ have been suggested to be $S$-wave $DN$ and $D^*N$ molecular states, respectively. From QCD sum rules, we investigate that whether $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ could be the $S$-wave $DN$ state with $J^P = \frac{1}{2}^-$ and the $S$-wave $D^*N$ state with $J^P = \frac{3}{2}^-$, respectively. In the OPE calculation, contributions of operators up to dimension 12 are included and one could find that its convergence is still under control. The final result for the $S$-wave $DN$ state of $J^P = \frac{1}{2}^-$ is $3.64 \pm 0.33$ GeV, which is somewhat bigger than the experimental value.
of $\Sigma_c(2800)$ even considering the uncertainty of result. The numerical result for the $S$-wave $D^*N$ state of $J^P = \frac{3}{2}^-$ is $3.73 \pm 0.35$ GeV, which is a bit higher than the experimental data of $\Lambda_c(2940)^+$ even taking into account the uncertainty. Considering that corresponding molecular currents are constructed from local operators, one can not arbitrarily exclude the possibility that $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ are molecular states just from these disagreements. However, one can infer that $\Sigma_c(2800)$ and $\Lambda_c(2940)^+$ could not be compact states from these results. This may suggest a limitation of the QCD sum rule using the local current to determine whether some state is a molecular state or not. By the way, we also study the corresponding bottom counterparts and predict their masses to be $6.97 \pm 0.34$ GeV for the $S$-wave $\bar{B}N$ state with $J^P = \frac{1}{2}^-$ and $6.98 \pm 0.34$ GeV for the $S$-wave $\bar{B}^*N$ state with $J^P = \frac{3}{2}^-$, which could be searched in future experiments.

Acknowledgments

The author would like to thank Prof. C. García-Recio and Prof. L. L. Salcedo for communications and discussions. This work was supported by the National Natural Science Foundation of China under Contracts No. 11105223, No. 11275268, and the project in NUDT for excellent youth talents.

[1] R. Mizuk et al., (Belle Collaboration), Phys. Rev. Lett. 94, 122002 (2005).
[2] B. Aubert et al., (BABAR Collaboration), Phys. Rev. D 78, 112003 (2008).
[3] B. Aubert et al., (BABAR Collaboration), Phys. Rev. Lett. 98, 012001 (2007).
[4] R. Mizuk et al., (Belle Collaboration), Phys. Rev. Lett. 98, 262001 (2007).
[5] S. Capstick and N. Isgur, Phys. Rev. D 34, 2809 (1986).
[6] D. Ebert, R. N. Faustova, and V. O. Galkin, Phys. Lett. B 659, 612 (2008).
[7] H. Y. Cheng and C. K. Chua, Phys. Rev. D 75, 014006 (2007).
[8] C. Chen, X. L. Chen, X. Liu, W. Z. Deng, and S. L. Zhu, Phys. Rev. D 75, 094017 (2007).
[9] X. H. Zhong and Q. Zhao, Phys. Rev. D 77, 074008 (2008).
[10] H. Garcilazo, J. Vijande, and A. Valcarce, J. Phys. G: Nucl. Part. Phys. 34, 961 (2007).
[11] B. Chen, D. X. Wang, and A. Zhang, Chin. Phys. C 33, 1327 (2009).
[12] J. He, Z. Ouyang, X. Liu, and X. Q. Li, Phys. Rev. D 84, 114010 (2011).
[13] M. F. M. Lutz and E. E. Kolomeitsev, Nucl. Phys. A 755, 29c (2005).
[14] C. E. Jiménez-Tejero, A. Ramos, and I. Vidaña, Phys. Rev. C 80, 035206 (2009).
[15] Y. Dong, A. Faessler, T. Gutsche, and V. E. Lyubovitskij, Phys. Rev. D 81, 074011 (2010).
[16] X. G. He, X. Q. Li, X. Liu, and X. Q. Zeng, Eur. Phys. J. C 51, 883 (2007).
[17] Y. Dong, A. Faessler, T. Gutsche, and V. E. Lyubovitskij, Phys. Rev. D 81, 014006 (2010).
[18] Y. Dong, A. Faessler, T. Gutsche, S. Kumano, and V. E. Lyubovitskij, Phys. Rev. D 82, 034035 (2010).
[19] Y. Dong, A. Faessler, T. Gutsche, S. Kumano, and V. E. Lyubovitskij, Phys. Rev. D 83, 094005 (2011).
[20] J. He and X. Liu, Phys. Rev. D 82, 114029 (2010).
[21] C. García-Recio, V. K. Magas, T. Mizutani, J. Nieves, A. Ramos, L. L. Salcedo, and L. Tolos, Phys. Rev. D 79, 054004 (2009).
[22] P. G. Ortega, D. R. Entem, and F. Fernández, [arXiv:1210.2633] [hep-ph].
[23] 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, 585 (1984).
[24] B. L. Ioffe, in The Spin Structure of The Nucleon, edited by B. Fröis, V. W. Hughes, and N. de Groot (World Scientific, Singapore, 1997).
[25] S. Narison, Camb. Monogr. Part. Phys. Nucl. Phys. Cosmol. 17, 1 (2002).
[26] P. Colangelo and A. Khodjamirian, At The Frontier of Particle Physics: Handbook of QCD, edited by M. Shifman, Boris Ioffe Festschrift Vol. 3 (World Scientific, Singapore, 2001), pp. 1495-1576; A. Khodjamirian, Continuous Advances in QCD 2002/ARKADYFEST, arXiv:0209166 [hep-ph].
[27] M. Neubert, Phys. Rev. D 45, 2451 (1992); M. Neubert, Phys. Rep. 245, 259 (1994).
[28] M. Nielsen, F. S. Navarra, and S. H. Lee, Phys. Rep. 497, 41 (2010).
[29] S. L. Zhu, Phys. Rev. Lett 91, 232002 (2003); W. Wei, P. Z. Huang, H. X. Chen, and S. L. Zhu, JHEP 07, 015 (2005).
[30] R. D. Matheus, F. S. Navarra, M. Nielsen, R. Rodrigues da Silva, and S. H. Lee, Phys. Lett. B 578, 323 (2004); J. Sugiyama, T. Doi, and M. Oka, Phys. Lett. B 581, 167 (2004); M. Eidemuller, Phys. Lett. B 597, 314 (2004); Z. G. Wang, W. M. Yang, and S. L. Wan, Phys. Rev. D 72, 034012 (2005).
[31] H. Kim, S. H. Lee, and Y. Oh, Phys. Lett. B 595, 293 (2004).
[32] J. R. Zhang and M. Q. Huang, Phys. Rev. D 77, 094002 (2008); Phys. Rev. D 78, 094007 (2008); Phys. Lett. B 674, 28 (2009).
[33] J. R. Zhang and M. Q. Huang, Phys. Rev. D 80, 056004 (2009); J. Phys. G: Nucl. Part. Phys. 37, 025005 (2010); Commun. Theor. Phys. 54, 1075 (2010); J. R. Zhang, M. Zhong, and M. Q. Huang, Phys. Lett. B 704, 312 (2011).
[34] L. J. Reinders, H. R. Rubinstein, and S. Yazaki, Phys. Rep. 127, 1 (1985).
[35] B. L. Ioffe, Nucl. Phys. B188, 317 (1981); E. V. Shuryak, Nucl. Phys. B198, 83 (1982).
[36] S. Narison, Phys. Lett. B 605, 319 (2005).
[37] S. Grewe, J. G. Körner, and O. I. Yakovlev, Phys. Rev. D 55, 3016 (1997).
[38] H. X. Chen, A. Hosaka, and S. L. Zhu, Phys. Lett. B 650, 369 (2007).
[39] Z. G. Wang, Nucl. Phys. A 791, 106 (2007).
[40] R. D. Matheus, F. S. Navarra, M. Nielsen, and R. Rodrigues da Silva, Phys. Rev. D 76, 056005 (2007).
[41] J. R. Zhang, L. F. Gan, and M. Q. Huang, Phys. Rev. D 85, 116007 (2012); J. R. Zhang and G. F. Chen, Phys. Rev. D 86, 116006 (2012); J. R. Zhang, Phys. Rev. D 87, 076008 (2013).