S-Wave Single Heavy Baryons with Spin-3/2 at Finite Temperature

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The thermal behavior of the spectroscopic parameters of the S-wave single heavy baryons $\Sigma_Q^0, \Xi_Q^0$ and $\Lambda_Q^0$ with spin-3/2 are investigated in QCD at finite temperature. We analyze the variations of the mass and residue of these baryons taking into consideration the contributions of QCD thermal condensates up to dimension eight in Wilson expansion. At finite temperature, due to the collapse of the Lorentz invariance by the choice of reference frame and appearance of the extra $O(3)$-symmetry, new four-dimensional operators come out in the form of the fermionic and gluonic parts of the energy momentum tensor that are taken into account in the calculations. Our analyses show that at lower temperatures, the parameters of baryons under consideration are not affected by the medium. These parameters, however, show rapid variations with respect to temperature at higher temperatures near to a critical temperature, after which the baryons are melted. The results of the masses and residues at $T \to 0$ limit are compatible with the available experimental data and other theoretical studies.

I. INTRODUCTION

With the rising number of experimental data on charmed and bottom baryons, the interest in the investigation of heavy baryons has increased, considerably. Before giving the details of the experimental studies on heavy baryons, it would be useful to give some theoretical information. The Quark Model is one of the most successful tools to classify the mesons and baryons. The traditional single heavy baryons ($Qqq$) consist of one heavy ($Q = b$ or $c$) and two light quarks ($q = u, d$ or $s$). The mass of heavy quark is very large compared to the light quark masses and the light degrees of freedom form a diquark $qq$, which orbits the nearly static heavy $Q$ quark. Therefore, infinitely heavy mass limit ($m_Q \to \infty$) for the heavy quark is utilized to classify the single heavy baryons $\Xi_c^0, \Xi_c^\pm$. In this case, for the two light quarks, the total flavor-spin wave function has to be symmetric because their color wave function is antisymmetric. Hence there are two different representations for the S-wave heavy baryons ($3 \otimes 3 = \bar{3} \oplus 6$): antisymmetric $\bar{3}$ or symmetric 6. The antitriplet ($\bar{3}$) of baryons contain only spin-1/2 states while the sextet (6) of baryons contain both spin-1/2 and spin-3/2 states. In this study, we investigate the thermal properties of the single heavy bottom/charmed spin-3/2 sextet states: The members for charmed baryons are shown in Figure 1.

Experimentally, the $\Sigma_c^*$ antitriplet ($\Lambda_c^+, \Xi_c^+, \Omega_c^+$) states, the $\Xi_c^+$ sextet ($\Omega_c^+, \Sigma_c, \Xi_c^0$) baryons and the $\Lambda_c^+$ sextet ($\Omega_c^+, \Sigma_c, \Xi_c^0$) resonances have been observed in the charmed sector while the only $\Lambda_b, \Sigma_b^{(*)}, \Xi_b^{(*)}$ and $\Omega_b$ have been discovered in the bottom picture. Some history of discoveries are in order: In 2006 the CDF collaboration reported observation of $\Lambda_b$ [3] and $\Omega_c^*$ discovered by the Babar collaboration [6]. The CDF collaboration reported the first observation of $\Sigma_b$ and $\Sigma_b^*$ baryons later [7]. The D0 collaboration declared the observation of $\Xi_b$ [8] and it was confirmed by CDF in a short time [9]. The observation of ground and excited states of $\Xi_c^0$ were proclaimed by Belle and BABAR collaborations [10, 11]. $\Xi_c^+$ observed by Belle in 2008 [12] and discovery of $\Sigma_c^*$ was reported by CMS and LHCb collaborations [13, 14].

On the other hand, various theoretical studies in vac-
uum have been utilized to investigate the spectroscopic parameters of single heavy baryons. In 1982 Shuryak primarily calculated the heavy baryon masses via QCD sum rule [15]. In Ref. [16] Capstick and Isgur examined the heavy baryon systems in a quark potential model. Bagan et al. investigated the heavy baryons by taking into account the separation of negative and positive parity contributions [17]. Grozin and Yakovlev evaluated the masses of $\Lambda_Q$ and $\Sigma_Q^{(*)}$ using the heavy quark effective theory (HQET) [18]. Charmed baryons were investigated in Chiral perturbation theory by Savage and also results were extended for b-baryons in the same study [19]. Roncaglia et al. in Ref. [20] estimated the heavy baryon masses in the framework of Feynman-Hellman theorem. In Ref. [21] Jenkins studied the masses of heavy baryons in the $1/m_Q$ and $1/N_c$ expansions. The $1/m$ corrections to heavy baryon masses were calculated by Dai et al. in the framework of the HQET [22]. QCD sum rule for heavy baryons at leading order in $1/m_Q$ and to the leading order in $\alpha_s$ were evaluated by Groote et al. in Ref. [23]. Wang et al. improved the analysis for the $\Lambda_Q$ and $\Sigma_Q$ baryon masses to order $\Lambda_{QCD}/m_Q$ from QCD sum rule [24]. Mathur et al. predicted the mass spectrum of charmed and bottom baryons from Lattice QCD [25]. Wang and Huang in Ref. [26] studied the mass, coupling constant, and Isgur-Wise function for ground-state heavy baryons within the framework of HQET by taking into account both the two and three-point correlation functions. Ebert et al. computed heavy baryon masses in the heavy-quark light-diquark approximation in the framework of constituent quark model [27]. Garcilazo et al. solved exactly the three quark problem via Faddeev method in momentum space [28]. In Ref. [29] Zang and Huang calculated the charm and bottom baryon masses up to operator dimension six in operator product expansion (OPE) by the help of the QCD sum rule approach. The mass and residue of $\Omega_c^*$ and $\Omega_b^*$ with spin parity $3/2^+$ were studied by Wang via QCD sum rule [30]. A quark model was applied to the spectrum of baryons containing one heavy baryon by Roberts and Pervin [31]. Bottom baryon spectra were investigated using Faddeev method in momentum space by Valcarce et al. [32]. Liu et al. performed a systematic study of the masses of bottom baryons up to $1/m_Q$ in HQET [33]. Using the coupled channel formalism, Gerasyuta and Matskevich calculated the S-wave bottom baryons masses [34]. In Ref. [35] Karliner et al. investigated the b-baryons in the quark model. Groote et al. computed the NLO perturbative corrections for the static properties of heavy baryons [36]. In Ref. [37] the mass of $\Lambda_Q$ and $\Sigma_Q^{(*)}$ baryons and in Ref. [38] charmed and bottom baryons were calculated by Zang and Huang via QCD sum rule taking into account operators up to dimension six. In two-point and light cone QCD sum rule methods Aliev et al. studied the mass and magnetic moments of single heavy baryons with spin-3/2 [39]. Lewis and Shyn predicted the bottom baryon masses based on a $2 + 1$ flavor dynamical lattice QCD simulation [40]. The spin-3/2 heavy and doubly heavy baryons were investigated by subtracting the contributions from the corresponding negative parity by Z. G. Wang. The mass spectra of heavy baryons were studied by the help of the motivated relativistic quark model by Ebert [41]. Kim et al. investigated the single heavy baryon mass based on the self-consistent Chiral quark soliton model in Ref. [42]. Finally, Azizi and Er studied the in-medium properties of spin-3/2 heavy baryons in nuclear matter using QCD sum rule in a dense medium [43].

Theoretical investigations of spectroscopic parameters of the single heavy baryons at finite temperature will help us better understand and analyze the results of heavy-ion collision experiments and gain valuable information on the internal structures of these baryons, behavior of these baryons near to the critical temperature, possible phase transition to quark gluon plasma (QGP) (adopted as a new phase of matter) as well as the perturbative and nonperturbative dynamics of QCD. At short distances, to describe the strong interaction quantum chromodynamics (QCD) is a suitable theory. However, the calculations of hadronic parameters including nonperturbative effects (occur in low energy scale) usually need some nonperturbative phenomenological models. Many phenomenological models are available in the literature: QCD sum rule is one of the powerful ones among them. This method firstly suggested by Shifman, Vainshtein ve Zakharov to investigate the vacuum properties of mesons [44] and then Ioffe [45] applied this method for baryons. The thermal version of the QCD sum rule was extended by Bochkarev and Shaposhnikov [46]. In addition to the vacuum expectation values of quark and gluon condensates, their thermal forms and some new operators appear in the thermal version.

In this study, we investigate the temperature effects on the spectroscopic parameters of the ground state sextet baryons including single heavy quark and with spin-3/2 via thermal QCD sum rule method. Taking into account
the additional operators coming from OPE due to breaking of the Lorentz invariance by the choice of the thermal rest frame, condensates up to dimension eight are considered. The article is arranged in the following form. In Sec. II, the in-medium sum rules for the mass and residues of the $\Sigma^Q_0$, $\Xi^*_0$, and $\Omega^Q_0$ single heavy baryons are obtained. In Sec. III the numerical analysis for the spectroscopic parameters under consideration is performed. The last section includes the summary and our concluding remarks.

II. CALCULATIONS

In this section, QCD sum rules for the spectroscopic parameters of the spin-$3/2$ $\Sigma^0_0$, $\Xi^*_0$ and $\Omega^0_0$ baryons are obtained at finite temperature. To this end, we start with the following two-point thermal correlation function:

$$\Pi_{\mu\nu}(q, T) = i \int d^4 x \ e^{iq \cdot x} \langle \Psi | T \{ J_\mu(x), \bar{J}_\nu(0) \} | \Psi \rangle,$$

where $q$ is the four-momentum of the chosen baryon, $\Psi$ is the ground state of the hot medium, $T$ denotes the time-ordering operator and $J_\mu(x)$ is the interpolating current of the single heavy baryon, $B_{SH}$.

As the standard procedures of the QCD sum rule, the correlation function given above can be calculated at different contexts. At large distances, it is evaluated in terms of the hadronic parameters such as the mass and residue of hadron. We call it the physical or hadronic representation of the correlator. The same correlator can be expressed in terms of the quark, gluon and mixed condensed by the help of the OPE at $q^2 << 0$ region. The computations in this way contain short distance effects. This representation, is generally called the OPE or QCD side of the correlation function. At the end of the day, we match the two windows and compare the coefficients of the same Lorentz structures from both sides. To remove the unwanted contributions coming from the higher states and continuum, the Borel transformation as well as continuum subtraction, supplied by the quark-hadron duality assumption at finite temperature, are performed. These procedures bring some auxiliary parameters, which we fix them before making any numerical estimations on the physical quantities.

To obtain the physical side of the correlator, a complete set of intermediate state with the same quantum numbers and quark content as the chosen current is inserted between the interpolating currents in correlation function. This is followed by the integral over four-$x$, which leads to

$$\Pi^{phys}_{\mu\nu}(q, T) = \langle \Psi | J_\mu(0) | B_{SH}(q, s) \rangle \langle B_{SH}(q, s) | J_\nu^{(2)}(0) | \Psi \rangle,$$

where $m_{B_{SH}}(T)$ is the temperature-dependent mass of the ground state of $B_{SH}$. The matrix element $\langle \Psi | J_\mu(0) | B_{SH}(q, s) \rangle$ is defined in terms of the temperature dependent residue, $\lambda_{B_{SH}}(T)$, as

$$\langle \Psi | J_\mu(0) | B_{SH}(q, s) \rangle = \lambda_{B_{SH}}(T) u_\mu(q, s),$$

where $u_\mu(q, s)$ is the Rarita-Schwinger spinor. The final form of the physical side can be obtained by inserting Eq. (3) into Eq. (2) and summing over the spins of the $B_{SH}$. The summation over Rarita-Schwinger spinors is performed using

$$\sum_s u_\mu(q, s) \bar{u}_\nu(q, s)$$

$$= - \left( \frac{q + m_{B_{SH}}}{3m_{B_{SH}}} \right) \left[ g_{\mu\nu} - \frac{1}{3} \gamma_{\mu} \gamma_{\nu} - \frac{2 q_\mu q_\nu}{3m_{B_{SH}}} - \frac{q_\mu \gamma_\nu - q_\nu \gamma_\mu}{3m_{B_{SH}}} \right].$$

By using the above behast, we recast the physical side as

$$\Pi^{phys}_{\mu\nu}(q, T) = \frac{\lambda^2_{B_{SH}}(T)(q + m_{B_{SH}})}{q^2 - m^2_{B_{SH}}} \left[ g_{\mu\nu} - \frac{1}{3} \gamma_{\mu} \gamma_{\nu} - \frac{2 q_\mu q_\nu}{3m_{B_{SH}}} - \frac{q_\mu \gamma_\nu - q_\nu \gamma_\mu}{3m_{B_{SH}}} \right] + ..., \quad (5)$$

where $\lambda^2_{B_{SH}}(T) = \lambda_{B_{SH}}(T) \bar{\lambda}_{B_{SH}}(T)$. It should also be specified that the interpolating current $J_\mu(x)$ couples to both the spin-1/2 and spin-3/2 states. In this study, we only consider the contribution of spin-3/2 heavy baryons and we need to comb out the pollution of spin-1/2 state. These unwanted contributions can be eliminated in two different ways: 1) For spin-3/2 state, it should be introduced a projection operator which destroys the spin-1/2 contributions, 2) By a specific ordering of the Dirac matrices and remove the terms corresponding to the spin-1/2 particles (for more details see for instance [47]). The contribution of the spin-1/2 states can be traced using

$$\langle \Psi | J_\mu(0) | \frac{1}{2}(q) \rangle = [ \kappa_1 q_\mu + \kappa_2 \gamma^\mu ] u(q), \quad (6)$$

where $\kappa_1$ and $\kappa_2$ are some constants. By applying the condition $J_\mu \gamma^\mu = 0$ (for more details see [48]), we get $\kappa_1$
in terms of \( \kappa_2 \). Hence,

\[
\langle \Psi | J_\mu(0) | \frac{1}{\sqrt{2}}(q) \rangle = \kappa_2 \left( \gamma_\mu - \frac{4}{m_2} q_\mu \right) u(q). \tag{7}
\]

As is seen from Eq. (7), the pollution coming from spin-1/2 resonances are commensurate to either \( q_1 \) or \( \gamma_\mu \). To remove these contributions, the Dirac matrices are ordered as \( \gamma_\mu \) and terms proportional to \( q_1 \) or \( q_\mu \), also those beginning with \( \gamma_\mu \) or ending with \( \gamma_\mu \) are set to zero. Finally, the clean physical side of the correlator, in the Borel scheme, is obtained as

\[
B\Pi^{\text{phys}}_{\mu \nu}(q, T) = \lambda_B^2 e^{-m_B^2/T}/M^2 \delta_{\mu \nu} + \lambda_B e^{-m_B^2/T}/M^2 \gamma_{\mu \nu} + \cdots,
\tag{8}
\]

where \( M^2 \) is the Borel parameter and dots denote the contributions of other structures as well as the higher states and continuum.

The next step is to calculate the OPE side of the correlation function. In deep Euclidean region, the correlation function is evaluated in terms of the quark and gluon degrees of freedom by the help of Wilson expansion. To achieve this goal, the basic point it to choose a suitable interpolating current defining the particles under study. The general form of the interpolating current for spin-3/2 \( B_{SH} \) in a compact form can be written as \([49, 51]\)

\[
J_\mu(x) = A \epsilon_{abc} \left( q^{aT}(x) C \gamma_\mu q^b_T(x) \right) Q^c(x) + \left( q^{aT}(x) C \gamma_\mu q^b_T(x) \right) q^c(x) + \left( Q^{aT}(x) C \gamma_\mu q^b_T(x) \right) q^c(x),
\tag{9}
\]

where \( A \) is the normalization constant, \( \epsilon_{abc} \) is the antisymmetric Levi-Civita tensor, \( a, b, c \) are color indices, \( q_1(2) \) denotes the light quark \((u, d \text{ or } s)\), \( Q \) is the bottom \((b)\) or charm \((c)\) quark and \( C \) is the charge conjugation operator. The normalization constant \( A \) and the \( q_1(2) \) quark for the considered baryons are given in Table II.

By inserting the explicit form of the interpolating current into the correlator and contracting all heavy and light quark fields via Wick’s theorem, we get the correlation function in the case of \( q_1 \neq q_2 \) in terms of the thermal light/heavy quark propagators, \( S_{q}(Q) \), as

\[
\Pi^{\text{OPE}}_{\mu \nu}(q, T) = -\frac{2i}{3} \epsilon_{abc} \epsilon_{a'b'c'} \int d^4 x e^{iq \cdot x} \left\{ S^{cc'_c T}_{q} [S^{ba'_{b'}}_{q_2} \bar{S}^{ab}_{q_1} \gamma_\mu + S^{cc'_c T}_{q_1} [S^{ba'_{b'}}_{q_2} \bar{S}^{ab}_{q_1} \gamma_\mu] + S^{cc'_{c'c}}_{q} [S^{ba'_{b'}}_{q_2} \bar{S}^{ab}_{q_1} \gamma_\mu] + S^{cc'_{c'c}}_{q_1} [S^{ba'_{b'}}_{q_2} \bar{S}^{ab}_{q_1} \gamma_\mu] + S^{cc'_{c'c}}_{q_2} [S^{ba'_{b'}}_{q_2} \bar{S}^{ab}_{q_1} \gamma_\mu] \right\}.
\tag{10}
\]

Some extra contractions arise because of the identical particles in the case of \( q_1 = q_2 = q \), and the correlator is obtained as

\[
\Pi^{\text{OPE}}_{\mu \nu}(q, T) = \frac{i}{3} \epsilon_{abc} \epsilon_{a'b'c'} \int d^4 x e^{iq \cdot x} \left\{ 2 S^{cc'_c T} [S^{ba'_{b'}}_{q_2} \bar{S}^{ab}_{q_1} \gamma_\mu] + 2 S^{cc'_c T} [S^{ba'_{b'}}_{q_2} \bar{S}^{ab}_{q_1} \gamma_\mu] + S^{cc'_{c'c}}_q [S^{ba'_{b'}}_{q_2} \bar{S}^{ab}_{q_1} \gamma_\mu] + S^{cc'_{c'c}}_q [S^{ba'_{b'}}_{q_2} \bar{S}^{ab}_{q_1} \gamma_\mu] + S^{cc'_{c'c}}_q [S^{ba'_{b'}}_{q_2} \bar{S}^{ab}_{q_1} \gamma_\mu] \right\},
\tag{11}
\]

where \( \bar{S}_{i}(Q) = C \bar{S}_{i}^{qT} C \). To go further in the calculations, the thermal light quark propagator in coordinate space is selected as (see also \([52, 53]\))

\[
S^{ij}(x) = \frac{-i}{2 \pi^2 x^2} \delta_{ij} - \frac{m_4}{4 \pi^2 x^2} \delta_{ij} - \frac{\langle \bar{q}q \rangle T}{12} \delta_{ij} - \frac{x^2}{192} \frac{m_4^2}{\langle \bar{q}q \rangle T} \left[ 1 - \frac{m_4}{6} \right] \delta_{ij} + \frac{i}{3} \left[ \frac{m_4}{16} \langle \bar{q}q \rangle T - \frac{1}{12} \langle u^\mu \Theta^I_{\mu \nu} u^\nu \rangle \right] \delta_{ij} + \frac{1}{3} \langle u \cdot x \bar{q} (u^\mu \Theta^I_{\mu \nu} u^\nu) \rangle \delta_{ij} - \frac{i q_A \lambda^{ij}}{32 \pi^2 x^2} G^A_{\mu \nu} \left( \delta_{\sigma \mu \nu} + \sigma^{\mu \nu} \cdot \right) - \frac{x^2 \langle \bar{q}q \rangle^2 T}{7776} \delta_{ij} - \frac{x^4 \langle \bar{q}q \rangle T \langle \bar{q}q \rangle G^2 T}{27648} + \cdots,
\tag{12}
\]
where includes the thermal quark and gluon condensates \((\langle q\bar{q}\rangle)_T \) and \((g_s^2 G^2)_T \), gluon fields in thermal bath as well as new operators containing the energy momentum tensor, \(\Theta_{\mu\nu} \). For the heavy quark, the following propagator including the thermal gluon condensate and gluon fields in hot medium is used \cite{54}:

\[
S^{ij}_Q(x) = i \int \frac{dk}{(2\pi)^4} e^{-ikx} \left( \frac{k + m_Q}{k^2 - m_Q^2} \delta_{ij} - \frac{g_s G^{ij}_0}{4} \sigma^{\alpha\beta} \left[ \frac{(k + m_Q)(\bar{G} + m_Q) + (k - m_Q)(\bar{G})}{(k^2 - m_Q^2)^2} \right] + \frac{m_Q k^2 + m_Q \bar{k}}{12} (k^2 - m_Q^2) + \cdots \right). \tag{13}
\]

In Eqs. \(12\) and \(13\), \(m_{q(Q)}\) denotes the light (heavy) quark mass.

The thermal quark condensate, \(\langle q\bar{q}\rangle_T\), which is valid up to a critical temperature \(T_c = 197\) MeV is parameterized in terms of the vacuum condensate, \(\langle 0|q\bar{q}|0\rangle\).

For this quantity, we use the following parametrization presented in \cite{54}, which is based on the lattice QCD predictions on the behavior of the quark condensate with respect to temperature:

\[
\langle \bar{q}q \rangle_T = \frac{\langle 0|q\bar{q}|0\rangle}{1 + e^{18.10(1.85)q_{min}\frac{T}{T_c} + 4.99q_{min}\frac{T}{T_c}(T - 1)}}. \tag{14}
\]

As we previously mentioned, because of the choice of the thermal rest frame in Wilson expansion, the Lorentz invariance is broken. To restore that the four-velocity vector of the medium \(u^\mu = (1, 0, 0, 0)\) is introduced, which implies \(u^2 = 1\) and \(q \cdot u = q_0\). In thermal version, as also mentioned above, new operators representing the fermionic and gluonic parts of the energy-momentum tensor arises in OPE. The fermionic part \(\Theta^{ij\mu}_Q\) appears explicitly in the light-quark propagator, while the gluonic part of the energy-momentum tensor \(\Theta^{ij\mu}_A\) appears in the expansion of the trace of two-gluon field strength tensor in heat bath \cite{56}:

\[
\langle T \gamma^c \gamma^a G_{\mu\nu} \rangle = \frac{1}{24} \langle g_{\alpha\mu} g_{\beta\nu} - g_{\alpha\nu} g_{\beta\mu} \rangle (G^2)^T + \frac{1}{6} \left[ g_{\alpha\mu} g_{\beta\nu} - g_{\alpha\nu} g_{\beta\mu} - 2 (u_{\alpha} u_{\mu} g_{\beta\nu} + u_{\beta} u_{\mu} g_{\alpha\nu}) \right] \times \langle u^\lambda \Theta^q_{\lambda\sigma} u^\gamma \rangle. \tag{15}
\]

The temperature dependent gluon condensate \((G^2)_T\) is parameterized in terms of the vacuum gluon condensate \(\langle 0|G^2|0\rangle\) as \cite{57}:

\[
\langle G^2 \rangle_T = \langle 0|G^2|0\rangle \left[ 1 - 1.65 \left( \frac{T}{T_c} \right)^{8.735} + 0.04967 \left( \frac{T}{T_c} \right)^{0.7211} \right]. \tag{16}
\]

For the temperature-dependent strong coupling constant \cite{58,59} we utilize

\[
g_s^{-2}(T) = \frac{11}{8\pi^2} \ln \left( \frac{2\pi T}{\Lambda_{MS}} \right) + \frac{51}{8\pi^2} \ln \left[ 2 \ln \left( \frac{2\pi T}{\Lambda_{MS}} \right) \right], \tag{17}
\]

where, \(\Lambda_{MS} \simeq T_c/1.14\).

Alike to the physical part, the correlation function on the OPE side is expanded in terms of the Lorentz structures as

\[
\Pi_{\mu\nu}^{OPE}(q, T) = \Gamma_1^{OPE} \delta_{\mu\nu} + \Gamma_2^{OPE} g_{\mu\nu} + \text{other structures}, \tag{18}
\]

where \(\Gamma_1^{OPE}\) is the coefficient of the selected Lorentz structure. These functions can be expressed by the help of following dispersion integral:

\[
\Gamma_1^{OPE}(1) = \int_{s_{\min}}^\infty ds \rho_1^{OPE}(s, T) + \Gamma_{nonpert}^{1(2)}, \tag{19}
\]

where \(s_{\min} = (m_q + m_g + m_Q)^2\), \(\rho_1^{OPE}(s, T)\) is the spectral density obtained via the imaginary part of the perturbative correlation function (pert in the following equation stands for the perturbative contributions)

\[
\rho_1^{OPE}(s, T) = -\frac{1}{\pi} \text{Im} \Gamma_1^{OPE,pert}(1), \tag{20}
\]

and \(\Gamma_{nonpert}\) represents the contributions coming from all the nonperturbative effects. In this step, our main aim is to calculate the spectral densities, corresponding to the perturbative effects in the present study, as well as the nonperturbative contributions to the QCD side. To this end, the explicit forms of the heavy and light quark propagators are inserted into Eqs. \(10\) and \(11\). The next.
step is to perform the standard but lengthy calculations: These calculations contain Fourier integrals appearing in different forms, Borel transformation as well as continuum subtraction. By matching the coefficients of the selected structures from both the physical and OPE sides of the correlation function, we find the desired sum rules:

\[
\lambda^2_{B_{SH}}(T)e^{-m^2_{B_{SH}}(T)/M^2} = \hat{B}\Gamma_{1}^{OPE},
\]

and

\[
\lambda^2_{B_{SH}}(T)m_{B_{SH}}(T)e^{-m^2_{B_{SH}}(T)/M^2} = \hat{B}\Gamma_{2}^{OPE},
\]

where the functions \(\hat{B}\Gamma_{1(2)}^{OPE}\) denote the \(\Gamma_{1(2)}^{OPE}\) in Borel scheme and are given as

\[
\hat{B}\Gamma_{1(2)}^{OPE} = \int_{s_{\text{min}}}^{s_{0}(T)} ds \rho_{1(2)}^{OPE}(s, T)e^{-s/M^2} + \hat{B}\Gamma_{1(2)}^{\text{nonpert}},
\]

with \(s_{0}(T)\) being the temperature-dependent continuum threshold. We will use the above sum rules to extract the values of the mass and residue of the baryons under consideration as well as their thermal behavior in next section.

As examples, we would like to present the explicit forms of the \(\rho_{1}^{OPE}(s, T)\) and \(\hat{B}\Gamma_{1}^{\text{nonpert}}\) for the \(\Sigma_{b}\) baryon. They are obtained as

\[
\rho_{1}^{OPE}(s, T) = -\frac{1}{96\pi^4\beta} \int_{0}^{1} dz \left\{ z \left( m_b^2 + s\beta \right) \left[ z \left( 3m_b^2(z + 1) - 12m_b m_u + s\beta(7z + 3) \right) 
- 12m_d(m_b z - 2m_u\beta) \right] \right\} \Theta[L(s, z)],
\]

and

\[
\hat{B}\Gamma_{1}^{\text{nonpert}}.
\]
\[ \hat{B}^{1}_{1, \text{nonpert}} = \frac{-1}{1152\pi^4} \frac{1}{M_{\text{in}}} \int_{s_{\text{min}}}^{s_{0}(T)} ds \int_{0}^{1} dz \left\{ -96\pi^2 \left\{ \langle \bar{u}u \rangle + 2\langle u\Theta^g u | z(2z + 23) + 15 \rangle + 256\pi^2 \beta \langle u\Theta f u | (5z - 1) \rangle \right\} + \Theta[L(s, z)] \right\} \nonumber \\
+ \frac{1}{13824\pi^2 M_{\beta}^3} \left( \langle \bar{d}d \rangle \left\{ \langle G^2 \rangle \left[ -16m_{d}^4 m_{d}z + 8m_{d}^3 \beta(m_{d}m_{u} + 2M^2z) - 4m_{d}^2 M^2 \beta \left( 5m_{d}(3z + 1) - 4m_{u} \right) \right] + 8m_{d} M^2 \beta \left( 2m_{d}m_{u} (3z - 2) + M^2(z(10z - 3) - 3) \right) + 8m_{d}^3 (3M^2 - 8q_{d}^2) \right\} + 16 \langle \bar{d}d \rangle \left\{ \langle G^2 \rangle \left[ -16m_{d}^4 m_{d}z + 8m_{d}^3 \beta(m_{d}m_{u} + 2M^2z) - 4m_{d}^2 M^2 \beta \left( 5m_{d}(3z + 1) - 4m_{u} \right) \right] + 8m_{d} M^2 \beta \left( 2m_{d}m_{u} (3z - 2) + M^2(z(10z - 3) - 3) \right) + 8m_{d}^3 (3M^2 - 8q_{d}^2) \right\} \right\} \Theta[L(s, z)] \right\} \nonumber \\
+ \frac{1}{972 M_{4}^4} \left( \langle \bar{u}u \rangle \left[ 27\pi^2 \langle \bar{d}d \rangle \left( 3m_{d}^2 m_{d}m_{u} - 8m_{b} M^2 (m_{d} + m_{u} ) + 2M^2 (3m_{d}m_{u} + 8M^2) \right) \right] + 4M_{4}^2 g_{s}^2 \langle \bar{u}u \rangle \left( m_{b} (m_{d} + m_{u} + M^2) \right) \right\} + \frac{1}{6912 M_{2}^2} \left( \langle \bar{d}d \rangle \left[ 13m_{d} M^2 \langle G^2 \rangle g_{s}^2 + 52m_{d} M^2 g_{s}^2 \langle u\Theta^g u \rangle \right] + 512\pi^2 \langle u\Theta f u \rangle \left( 3m_{d}^2 m_{d} - 4m_{b} M^2 - 4m_{d} (M^2 + 2q_{d}^2) \right) \right\} \right\} \nonumber \\
+ \frac{1}{162 M_{8}^4} \left( 3\pi^2 m_{d}^3 \langle \bar{u}u \rangle \left[ 3m_{d}^2 m_{d}m_{u} - 5m_{d}^2 M^2 (m_{d} + m_{u} ) + 3m_{b}^2 M^2 (4M^2 - m_{d}m_{u}) + 3M^4 (4M^2 - m_{d}m_{u}) \right] \right\} - M_{4}^2 \langle u\Theta f u \rangle \left[ \langle G^2 \rangle g_{s}^2 + 4M_{4}^2 g_{s}^2 \langle u\Theta^g u \rangle + 16\pi^2 \langle u\Theta f u \rangle \left( 3m_{d}^2 + 8M^2 + 16q_{d}^2 \right) \right] \right\}, \] 

(25)

where \( \Theta \) stands for the unit-step function, \( L(s, z) = s \ z(1 - z) - m_{d}^2 \ z \) and \( \beta = z - 1 \).

### III. NUMERICAL RESULTS

In this section, we analyze the obtained sum rules for the masses and residues. They includes some input.
parameters such as the heavy and light quark masses, quark, gluon and mixed condensates in vacuum and energy of the quasi-particle in medium, \( q_0 \). Their numerical values are presented in Table II.

| Parameter | \( \langle 0 \vert \frac{1}{2} \alpha G^2 \vert 0 \rangle \) | \( m_{q0} \) | \( m_{d0} \) | \( m_{c0} \) | \( m_{s0} \) | \( m_{g0} \) |
|-----------|-------------|----------|----------|----------|----------|----------|
| \( \langle 0 \vert \bar{d} d \vert 0 \rangle \) | \( \langle 0 \vert \bar{s} s \vert 0 \rangle \) | \( \langle 0 \vert \bar{u} u \vert 0 \rangle \) | \( \langle 0 \vert \bar{c} c \vert 0 \rangle \) |
| \( 6.08 \pm 0.40 \) GeV | \( 2.2^{+0.5}_{-0.4} \) MeV | \( 2.75^{+0.7}_{-0.3} \) MeV | \( 0.8 \pm 0.2 \) GeV |
| \( 5.94 \pm 0.31 \) MeV | \( 2.3 \) MeV | \( 6.0 \) MeV | \( 8 \) MeV |
| \( 6.0 \) MeV | \( 4.7 \) MeV | \( 1.27 \) MeV |
| \( 6.0 \pm 0.3 \) MeV | \( 0.8 \pm 0.2 \) MeV |

TABLE II: Input parameters used in calculations \(^{60,63}\).

In addition, we also need to have the gluonic and fermionic parts of the energy density. Based on the lattice QCD results on the thermal behavior of the energy-momentum tensor given in \(^{64}\), its parametrization was previously obtained in Ref. \(^{65}\) as

\[
\langle \Theta_{(0)}^G \rangle = \frac{1}{T^4} \left[ \frac{113.87 T^4}{4 \pi^2} \right] - 10.14 \frac{1}{GeV} T^5, \quad \text{(26)}
\]

which, we are going to use it in our calculations. The next problem is to obtain the parametrization of \( s_0(T) \) as a function of temperature. This function shall reduce to the vacuum threshold, \( s_0 \), at zero temperature. We parameterize it as

\[
s_0(T) = s_0 f(T), \quad \text{(27)}
\]

such that at \( T \to 0 \) limit, \( f(T) \to 1 \). Hence, we should first determine \( s_0 \) based on the standard prescriptions of the method, afterwards we will extract the function \( f(T) \) from the calculations.

Besides the continuum threshold in vacuum the sum rules obtained in previous section include another auxiliary parameter, Borel parameter \( M^2 \), which should also be fixed. We need to determine the working regions of \( s_0 \) and \( M^2 \) such that the physical quantities under consideration show mild dependence on these parameters. The continuum threshold \( s_0 \) is not totally free but it is related to the energy of the first excited state in the same channel. Thanks to the experiments that have provided many new results not only on the ground states but also on the first determine \( M^2 \) and \( s_0 \), which should also be fixed. We need to determine the working regions of \( s_0 \) and \( M^2 \) such that the physical quantities under consideration show mild dependence on these parameters. The continuum threshold \( s_0 \) is not totally free but it is related to the energy of the first excited state in the same channel. Thanks to the experiments that have provided many new results not only on the ground states but also on the excited states of some single heavy baryons, recently \(^{63}\). In view of PDG, we see that the excited states generally have energies about 300 MeV above the ground states masses. In choosing the working window for the \( s_0 \), we also look after the pole dominance and OPE convergence in our sum rules. These considerations leads to the window:

\[
| m_{B_{SH}} + 0.3 |^2 \text{ GeV}^2 \leq s_0 \leq | m_{B_{SH}} + 0.5 |^2 \text{ GeV}^2. \quad \text{(28)}
\]

The upper and lower limits of the Borel parameter are fixed consider the criteria of the QCD sum rule method. To find the lower limit, we apply the criterion of the OPE convergence at the chosen window for the continuum threshold. To this end, we demand that the perturbative part exceeds the total nonperturbative contributions and the slogan of the higher the dimension of the nonperturbative operator the lower its contribution is satisfied. Our calculations show that the operators having eight dimensions, the higher dimension that we include into the analyses, constitute only one percent of the total contribution at lower value of \( M^2 \), i.e. \( \Gamma_{(1/2)}^8, O(P) = m_{B_{min}, s_0} / \Gamma_{(1/2)} = 0.01 \). Figure 2 shows the convergence of the OPE in our calculations. As it is clear, the perturbative contribution is about 60% of the total at \( M^2_{min} = 6 \) and the main contribution in nonperturbative part belongs to the quark condensate, \( \langle \overline{q} q \rangle \). Note that, Figure 2 shows the OPE contributions in an additive form as stated in the caption of the figure.

FIG. 2: The contributions of the OPE operators: the perturbative contribution (black-solid line), \( \langle \overline{q} q \rangle \) (blue-solid line), \( \langle G^2 \rangle \), \( \langle u \Theta^f u \rangle \), \( \langle u \Theta^g u \rangle \) (red-dashed line), \( \langle q G q \rangle \) (green-dashed line), \( \langle \overline{q} \overline{q} \rangle \) (purple-solid line), \( \langle \overline{q} q \rangle \) (yellow-dashed line), \( \langle G^2 \rangle \), \( \langle u \Theta^f u \rangle \), \( \langle u \Theta^g u \rangle \) (brown-solid line).

To obtain \( M^2_{max} \), we utilize the condition of the pole
dominance as
\[ PC = \frac{\Gamma_{1(2)}^{OPE}(M^2, s_0)}{\Gamma_{1(2)}^{OPE}(M^2, \infty)} \geq \frac{1}{2}. \]  
(29)

As a result, we get the working region of the Borel parameter as \( M^2 \in [6, 10] \). We plot, as an example, a 3D graphic of the mass of \( \Sigma^*_b \) baryon as functions of \( M^2 \) and \( s_0 \) at \( T = 0 \) in Figure 3. As is seen the mass shows good stability against the variations of the auxiliary parameters in the selected windows.

Now, we proceed to find the function \( f(T) \) and the temperature dependent mass \( m_{B_{SH}}(T) \) and residue \( \lambda_{B_{SH}}(T) \) of the single heavy spin-3/2 baryons. To this end, we use the two sum rules in Eqs. (21) and (22) and one extra equation obtained by applying the derivative with respect to \( \frac{d}{d\left(\frac{M^2}{T^2}\right)} \) to both sides of Eq. (21). Simultaneous solving of the resultant three equations with the aim of obtaining the three mentioned unknowns gives the function \( f(T) \) as
\[ f(T) = 1 - 0.99 \left(\frac{T}{T_c}\right)^{12}. \]  
(30)

In the following, we proceed to discuss the thermal behavior of the masses and residues under study as the main goal of the present work. In this context, as examples, we plot the \( m(T)/m(0) \) and \( \lambda(T)/\lambda(0) \) for the bottom members as functions of \( T/T_c \) and \( M^2 \) in Figure 4 at average value of the vacuum continuum threshold. This figure shows that the spectroscopic parameters of the \( \Sigma^*_b, \Xi^*_b \) and \( \Lambda^*_b \) baryons are stable against the changes in temperature until a certain temperature but after that, they start to decrease with increasing the temperature. Our analyses show that the charmed baryons present similar behavior, as well. The points that the stability starts to break down for mass and residue are \( T \cong 0.15 \text{ GeV} \) and \( T \cong 0.13 \text{ GeV} \), respectively. After these points the mass and residue starts to diminish. The mass and residue fall substantially near to the critical temperature. The amount of decrements at \( T_c \) are 87% (86%) for the mass of bottom (charmed) and 60% (80%) for the residue of bottom (charmed) baryons, respectively compared to their vacuum values. These behavior of baryons near to the deconfinement/critical temperature can be interpreted as melting of the heavy baryons and their conversion to a soup of quark-gluon as probable new phase of the matter.

At the end of this section, we would like to present our results for the masses of the single heavy spin-1/2 baryons at \( T \to 0 \) limit. This is done in table III. For comparison, we also present the existing theoretical predictions in the literature and experimental data in the same table. With a quick glance in this table, we see that our predictions, within the errors, are overall consistent with other theoretical predictions made using different methods and approaches. Our predictions are also well consistent with the existing experimental data for five members within the presented uncertainties. \( \Omega^*_b \) baryon is only missing member, which has not been discovered in the experiment. We hope that, our result together with other theoretical predictions will help experimental group in the course of search for this particle.

IV. SUMMARY AND CONCLUDING REMARKS

In this study we have performed two-point thermal QCD sum rule analyses for \( \Sigma^*_Q, \Xi^*_Q \) and \( \Omega^*_Q \) single heavy baryons including bottom/charm quark which are the members of the spin-3/2 sextet family. In the OPE, op-
operators up to dimension eight were taken into account which lead to a good OPE convergence as well as pole dominance. We included the thermal effects by two ways: We replaced the vacuum condensates by their thermal counterparts and considered the extra operators, appearing in the forms of the fermionic and gluonic parts of the energy momentum tensor due to the restoration of the Lorentz invariance. We fixed the auxiliary parameters entering the calculations by the standard prescriptions of the method. By simultaneous solving of the two sum rules obtained together with an extra equation derived from one of the sum rules, we found three unknowns: Thermal continuum threshold, temperature-dependent mass and temperature-dependent residue. We discussed the thermal behavior of the mass and residue for all the bottom and charmed baryon members having the spin-3/2. We observed that the spectroscopic parameters remain unchanged up to a certain temperature, after which they start to diminish considerably near to the critical temperature. The decrements order in the mass and residue of the considered baryons near to the deconfinement temperature are obtained as 86% and (60 – 80)%, respectively. The melting of the baryons under study at higher temperatures may be considered as the QCD phase transition to QGP as possible new phase of the matter. Our results can be used in the analyses of the in-medium and heavy ion collision experiments results with the aim of investigation of the behavior of the heavy baryons under extreme conditions.

We extracted the values of the masses for both the bottom and charmed baryons at $T \rightarrow 0$ limit and compared with the predictions of other phenomenological models and experimental data. The obtained results are well consistent with existing experimental data. Our result on the mass of $\Omega_b^+$ baryon as the only undiscovered member together with other predictions may help the experimental group to hunt this particle and measure its parameters.

|      | $m_{\Omega^+}$ | $m_{\Omega^+}$ | $m_{\Sigma^+}$ | $m_{\Sigma^+}$ | $m_{\Sigma^+}$ | $m_{\Sigma^+}$ |
|------|----------------|----------------|----------------|----------------|----------------|----------------|
| present work | $6.08^{+0.10}_{-0.15}$ | $2.75^{+0.08}_{-0.26}$ | $5.88^{+0.11}_{-0.11}$ | $2.56^{+0.08}_{-0.07}$ | $5.95^{+0.12}_{-0.13}$ | $2.65^{+0.08}_{-0.07}$ |
| 16   | -              | -              | 5.805          | 2.495          | -              | -              |
| 17   | -              | -              | 5.4 ~ 6.2      | 2.15 ~ 2.92    | -              | -              |
| 19   | -              | 2.768          | -              | 2.518          | -              | -              |
| 20   | $6.090 \pm 0.050$ | $2.770 \pm 0.030$ | $5.850 \pm 0.040$ | $2.520 \pm 0.020$ | $5.980 \pm 0.040$ | $2.650 \pm 0.020$ |
| 21   | 6.083          | 2.760          | 5.840          | -              | 5.966          | -              |
| 22   | -              | -              | 5.84 \pm 0.09  | 2.55 \pm 0.08  | -              | -              |
| 24   | -              | -              | 5.82 \pm 0.13  | 2.59 \pm 0.20  | -              | -              |
| 25   | 6.060          | 2.752          | 5.871          | 2.538          | 5.959          | 2.680          |
| 27   | 6.088          | 2.768          | 5.834          | 2.518          | 5.963          | 2.654          |
| 29, 37 | 6.00 \pm 0.16   | 2.74 \pm 0.23  | 5.81 \pm 0.19  | 2.56 \pm 0.24  | 5.94 \pm 0.17  | 2.64 \pm 0.22  |
| 30   | 6.06 \pm 0.13  | 2.76 \pm 0.10  | -              | -              | -              | -              |
| 32   | 6.079          | 2.767          | 5.829          | 2.502          | 5.961          | 2.642          |
| 33   | $6.063^{+0.082}_{-0.109}$ | $2.790^{+0.109}_{-0.105}$ | $5.835^{+0.082}_{-0.077}$ | $2.534^{+0.096}_{-0.081}$ | $5.929^{+0.083}_{-0.079}$ | $2.634^{+0.102}_{-0.094}$ |
| 34   | -              | -              | 5.829          | -              | -              | -              |
| 35   | 6.082          | -              | -              | 5.959 \pm 0.004 | -              | -              |
| 38   | 6.08 \pm 0.40  | 2.72 \pm 0.20  | 5.85 \pm 0.35  | 2.51 \pm 0.15  | 5.97 \pm 0.40  | 2.66 \pm 0.18  |
| 39   | 6.044 \pm 0.018 | -              | 5.842 \pm 0.026 | -              | 5.950 \pm 0.021 | -              |
| 40   | 6.17 \pm 0.15  | 2.79 \pm 0.19  | 5.85 \pm 0.20  | 2.48 \pm 0.25  | 6.02 \pm 0.17  | 2.65 \pm 0.20  |
| 41   | 6.088          | 2.768          | 5.834          | 2.519          | 5.963          | 2.649          |
| 42   | 6.073          | -              | 5.834          | -              | 5.954          | -              |
| Exp[3] | -              | $2.7659 \pm 0.0020$ | $5.83032 \pm 0.00027$ | $2.51848 \pm 0.00020$ | $5.9523 \pm 0.0009$ | $2.64638 \pm 0.00021$ |

TABLE III: The vacuum mass comparison of the single heavy spin-3/2 baryons in GeV with existing theoretical predictions and experimental data (Exp[3]).
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FIG. 4: The mass (right) and residue (left) of the bottom baryons as functions of $M^2$ and $T/T_c$. 