

Open heavy flavor mesons in hot asymmetric strange hadronic matter: A QCD sum rule approach

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The in-medium masses of the pseudoscalar open charm (D , $\bar{D}$, D_s and $\bar{D}_s$) and open bottom (B , $\bar{B}$, B_s and $\bar{B}_s$) mesons in hot asymmetric strange hadronic matter are studied within a QCD sum rule approach. These are computed using the medium modifications of the light quark condensates ($\langle \bar{q}_i q_i \rangle$, with $(q_i, i = 1, 2, 3) \equiv (u, d, s)$), the scalar gluon condensate, $\langle \frac{\alpha_s}{\pi} G_{\mu\nu}^a G^{a\mu\nu} \rangle$, the twist-2 tensorial gluon operator, $\langle \frac{\alpha_s}{\pi} (G_{\mu\sigma}^a G_{\nu}^{\sigma} u^\mu u^\nu - \frac{1}{4} G_{\mu\sigma}^a G^{a\mu\sigma}) \rangle$ (where u^μ is the four-velocity of the medium) and other operators in the operator product expansion up to mass dimension 5. Within a chiral SU(3) model, the quark condensates are obtained from the medium modifications of the nonstrange and strange scalar-isoscalar fields (σ and ζ), and the scalar-isovector field, δ , whereas, the gluon (scalar and twist-2) condensates are computed from the in-medium value of the dilaton field, χ , which is incorporated within the chiral model to mimic the broken scale invariance of QCD. The mixed quark-gluon condensate $\langle \bar{q}_i g_s \sigma \cdot G q_i \rangle$ is calculated from the in-medium quark condensate $\langle \bar{q}_i q_i \rangle$. The splittings of the masses of the $D - \bar{D}$ ($B - \bar{B}$), as well as $D_s - \bar{D}_s$ ($B_s - \bar{B}_s$) in the hadronic medium are due to the odd part of the spectral function. For the strange-charm (strange-bottom) mesons ($(q_i, i = 3) \equiv s$), the value of $\langle q_i^\dagger q_i \rangle$ dominates over the contributions from the other operators, $\langle q_i^\dagger i D_0^2 q_i \rangle$ and $\langle q_i^\dagger g_s \sigma \cdot G q_i \rangle$, of the odd part of the spectral function. This is observed as opposite behavior of the particle-antiparticle mass splittings with the mass of the particle to be larger (smaller) than the antiparticle mass in nuclear (hyperonic) matter. The density and isospin asymmetry effects are observed to be the dominant medium effects which might have observable consequences on the production of charmonia and open charm mesons, and, on the production ratios, D^+/D^- and $D^0/\bar{D}^0$, in asymmetric heavy ion collisions in the Compressed Baryonic Matter experiment at FAIR at the future facility at Gesellschaft für Schwerionenforschung (GSI).

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I. INTRODUCTION

The study of the properties of hadrons under extreme conditions of temperature and/or density is an important subject of research in strong interaction physics. The topic is of direct relevance for the experimental observables of the relativistic heavy-ion collision experiments, as well as of relevance for study of the bulk properties of compact objects, e.g., the neutron stars. The light vector mesons have been extensively studied in the literature, as the medium modifications of these mesons can be observed in the dilepton spectra arising from heavy ion collision experiments [1]. In the recent years, a lot of studies on the in-medium properties of heavy flavor hadrons have been

initiated due to its relevance and possible observability in the relativistic heavy ion collision experiments [2]. The spectral modification of the charm hadrons will be investigated via $\bar{p}$ -A collisions at the Facility for Antiproton and Ion Research (FAIR) at the future facility at GSI, Darmstadt [3,4] and charm production using pion beams on protons in Japan Proton Accelerator Research Complex (J-PARC) [5–7]. The charm (bottom) meson-nucleon interaction can also lead to possibility of formation of charm (bottom) mesons bound in atomic nuclei.

The properties of heavy flavor mesons in the (hot) hadronic matter have been studied extensively in the literature. These have been investigated using the potential models [8], the coupled channel approach [9], the quark meson coupling model [10–12] heavy quark symmetry and interaction of these mesons with nucleons via pion exchange [13], heavy meson effective theory [14], studying the heavy flavor meson as an impurity in nuclear matter [15], and the QCD sum rule approach [16–33]. Using leading order QCD formula [34–36], the medium modifications of the masses of the charmonium states due to the medium change of the scalar gluon condensate were

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studied in a linear density approximation in Ref. [37], which showed much larger mass shifts for the excited states [$\psi(3686)$ and $\psi(3770)$], as compared to the mass shift of J/ψ . Within a chiral effective model [38–42], the in-medium masses of the charmonium (bottomonium) states [43–45] as well as of the open heavy flavor mesons [43,44,46–50] have been studied. Within the chiral effective model, the mass modifications of the quarkonium states arise due to the medium modification of a scalar dilaton field, which is introduced in the model to mimic the broken scale invariance of QCD. The open heavy flavor mesons have been studied due to their interactions with the baryons and the scalar mesons in the hadronic medium. The model has been used to study the partial decay widths of the heavy quarkonium states to the open heavy flavor mesons, in the (magnetized) hadronic medium [44,51] using a light quark pair creation model [52], namely the 3P_0 model [53–56] as well as using a field theoretical model for composite hadrons [57–60]. Within the chiral effective model, in the presence of an external magnetic field, the effects of magnetized Dirac sea have been investigated on the properties of heavy flavor mesons in the hadronic medium [61–63]. The magnetized Dirac sea, accounting for the anomalous magnetic moment of nucleons is observed to lead to the phenomenon of inverse magnetic catalysis at finite densities for symmetric as well as asymmetric nuclear matter in presence of strong magnetic fields. Recently, the effects of the Dirac sea as well as mixing of the pseudoscalar and vector mesons on the masses, decay widths, and the production cross-sections of the charmonium states in hot magnetized strange hadronic matter have been investigated in Ref. [64].

Within the chiral SU(3) model, using the medium changes of the light quark condensates ($\langle\bar{u}u\rangle$, $\langle\bar{d}d\rangle$ and $\langle\bar{s}s\rangle$), obtained from the nonstrange isoscalar (σ) isovector (δ) and strange (ζ) scalar fields and gluon condensates calculated from the scalar dilaton field, χ , the light vector mesons (ω , ρ and ϕ), in (magnetized) hadronic matter have been studied within the framework of the QCD sum rule approach [65,66]. The effects of the magnetized Dirac sea on the spectral properties of light vector and axial-vector mesons in hot magnetized nuclear matter [67] and of the ground state heavy quarkonia [68] have also been studied with the QCD sum rule approach. In the present work, we investigate the masses of the open charm (D , $\bar{D}$, D_s) and open bottom (B , $\bar{B}$, B_s) mesons in isospin asymmetric strange hadronic matter at finite temperature. In Sec. II, we briefly describe the QCD sum rule approach used to compute the masses of the open charm and open bottom mesons, retaining the operators up to dimension 5 in the operator product expansion of the spectral function. Unlike the case of a vacuum, in the hadronic medium, there is splitting between the masses of the particles and their antiparticles arising due to the odd part of the spectral function. In Sec. III, we describe the chiral SU(3) model used to compute the in-medium values of the quark (light

and strange), the gluon (scalar and twist-2), and the mixed quark-gluon operators, which are used for the calculation of the in-medium masses of the pseudoscalar open charm and open bottom mesons. Section IV discusses the results and in Sec. V, we summarize the findings of the present work.

II. QCD SUM RULE APPROACH

The QCD sum rule approach relates the hadronic properties, encoded in spectral functions, to the QCD condensates [69,70]. The basic ingredient is the two-point correlation function given as

$$\Pi(q) = i \int d^4x e^{-iq \cdot x} \langle \Omega | T(j(x)j(0)^\dagger) | \Omega \rangle, \quad (1)$$

where, $|\Omega\rangle$ is the physical ground state. Unlike the case of vacuum, the two-point correlation function in the hadronic medium has contributions, which are odd in q_0 , and can be written in terms of even and odd functions of q_0 as

$$\Pi(q_0, \vec{q}) = \Pi^e(q_0, \vec{q}) + q_0 \Pi^o(q_0, \vec{q}), \quad (2)$$

with $\Pi^e(q_0, \vec{q}) = \frac{1}{2}(\Pi(q_0, \vec{q}) + \Pi(-q_0, \vec{q}))$ and $\Pi^o(q_0, \vec{q}) = \frac{1}{2q_0}(\Pi(q_0, \vec{q}) - \Pi(-q_0, \vec{q}))$. It might be noted here that $\Pi^{e(o)}(q_0, \vec{q})$ depends on q_0^2 . The Borel transformed sum rules for the even and odd parts, for $\vec{q} = 0$, are given as [25]

$$\begin{aligned} \mathcal{B}[\Pi_{OPE}^e(\omega^2, \vec{q} = 0)](M^2) &= \frac{1}{\pi} \left[\int_{s_0^-}^{s_0^+} ds \text{Im} \Pi^{(ph)}(s) \right. \\ &\quad \left. + \left(\int_{-\infty}^{s_0^-} + \int_{s_0^+}^{\infty} \right) ds \text{Im} \Pi(s) \right] s e^{-\frac{s^2}{M^2}}, \end{aligned} \quad (3)$$

and

$$\begin{aligned} \mathcal{B}[\Pi_{OPE}^o(\omega^2, \vec{q} = 0)](M^2) &= \frac{1}{\pi} \left[\int_{s_0^-}^{s_0^+} ds \text{Im} \Pi^{(ph)}(s) \right. \\ &\quad \left. + \left(\int_{-\infty}^{s_0^-} + \int_{s_0^+}^{\infty} \right) ds \text{Im} \Pi(s) \right] e^{-\frac{s^2}{M^2}}. \end{aligned} \quad (4)$$

In the above equations, $\omega^2 = -q_0^2$, the subscript “OPE” denotes the operator product expansion $\langle \Omega | T(j(x)j(y)^\dagger) | \Omega \rangle = \sum_n C_n(x-y) \langle \Omega | O_n | \Omega \rangle$, where $\langle \Omega | O_n | \Omega \rangle$ are the QCD condensates for the operators O_n , $C_n(x-y)$ are the Wilson coefficients, $s_0^+(s_0^-)$ is the continuum threshold for the particle (antiparticle), and M is the Borel mass. The open charm and open bottom mesons have been studied in symmetric nuclear matter within the QCD sum rule approach keeping operators up to dimension 5 in the OPE side of the spectral function in Ref. [25]. In the present work, we study the in-medium masses of these pseudoscalar mesons in hot isospin asymmetric strange hadronic matter, again, retaining

QCD operators up to dimension 5 in the operator product expansion. The spectral functions of the open heavy flavor mesons are obtained using the current operators for the particles (p_i), $j_i = j_i^p \equiv i\bar{q}_i\gamma_5 Q$, with $(q_i, i = 1, 2, 3) \equiv (u, d, s)$, where the heavy quark $Q = c$ corresponds to D^0 , D^+ , and D_s^+ mesons, and $Q = b$ corresponds to B^- , $\bar{B}^0$, and $\bar{B}_s^0$ mesons, respectively. The current operator is $j_i^\dagger \equiv i\bar{Q}\gamma_5 q_i$, for the antiparticles ($\bar{D}^0$, D^- , D_s^- for charm sector and B^+ , B^0 and B_s^0 for the bottom sector). For the phenomenological hadronic spectral function, we take the ansatz [25]

$$\text{Im}\Pi^{(ph)}(s) = \pi F_+ \delta(s - m_+) - \pi F_- \delta(s + m_-), \quad (5)$$

where $m_\pm$ corresponds to the mass of the particle (antiparticle). In Eqs. (3) and (4), we change the integration variable from s to $\tilde{s} (= s^2)$ in the integrals corresponding to the continuum parts of the spectral functions, and choose the continuum thresholds for the particle and antiparticle to be $s_0^\pm = \pm\sqrt{\tilde{s}_0}$, similar to the case of vacuum. We also assume, according to semilocal duality hypothesis, that $\text{Im}\Pi(\tilde{s})$ corresponding to the particle and antiparticle continuum

parts of the spectral function on the right-hand side of the Eqs. (3) and (4) are the same as the perturbative contribution of the spectral function, $\text{Im}\Pi_{OPE}^{per}(s)$, as calculated in operator product expansion [26]. It might be noted here that for the odd part of the spectral function, the perturbative contributions from the particle and antiparticle cancel with each other for the choice of the continuum thresholds $s_0^\pm = \pm\sqrt{\tilde{s}_0}$, as in the present work. The Borel transformed sum rules for the even and odd parts of the spectral function are then obtained as [25]

$$\begin{aligned} \mathcal{B}[\Pi_{OPE}^e(\omega^2, \vec{q} = 0)](M^2) &= \frac{1}{\pi} \int_{m_Q^2}^{\infty} d\tilde{s} \text{Im}\Pi_{OPE}^{per}(\tilde{s}, \vec{q} = 0) e^{-\tilde{s}/M^2} + e^{-m_Q^2/M^2} P^e(M^2) \\ &= m_+ F_+ e^{-m_+^2/M^2} + m_- F_- e^{-m_-^2/M^2} \\ &\quad + \frac{1}{\pi} \int_{\tilde{s}_0}^{\infty} d\tilde{s} \text{Im}\Pi_{OPE}^{per}(\tilde{s}) e^{-\tilde{s}/M^2} \end{aligned} \quad (6)$$

with

$$\begin{aligned} P^e(M^2) &= -m_Q \langle \bar{q}_i q_i \rangle + \frac{1}{2} m_{q_i} \langle \bar{q}_i q_i \rangle \left(1 + \frac{m_Q^2}{M^2} \right) + \frac{m_Q}{2M^2} \left(\frac{m_Q^2}{2M^2} - 1 \right) \langle \bar{q}_i g_s \sigma \cdot G q_i \rangle \\ &\quad + \frac{1}{12} \left\langle \frac{\alpha_s}{\pi} G_{\mu\nu} G^{\mu\nu} \right\rangle + \left[\left(\frac{7}{18} + \frac{1}{3} \ln \left(\frac{\mu^2 m_Q^2}{M^4} \right) - \frac{2\gamma_E}{3} \right) \left(\frac{m_Q^2}{M^2} - 1 \right) - \frac{2}{3} \frac{m_Q^2}{M^2} \right] \\ &\quad \times \left\langle \frac{\alpha_s}{\pi} \left(G_{\mu\sigma}^a G_{\nu}^a \sigma^{\mu\nu} - \frac{1}{4} G_{\mu\sigma}^a G^{a\mu\sigma} \right) \right\rangle + 2 \left(\frac{m_Q^2}{M^2} - 1 \right) \langle q_i^\dagger i D_0 q_i \rangle \\ &\quad + 4 \left(\frac{m_Q^3}{2M^4} - \frac{m_Q}{M^2} \right) \left[\langle \bar{q}_i D_0^2 q_i \rangle - \frac{1}{8} \langle \bar{q}_i g_s \sigma \cdot G q_i \rangle \right], \end{aligned} \quad (7)$$

and,

$$\begin{aligned} \mathcal{B}[\Pi_{OPE}^o(\omega^2, \vec{q} = 0)](M^2) &= e^{-m_Q^2/M^2} P^o(M^2) \\ &= F_+ e^{-m_+^2/M^2} - F_- e^{-m_-^2/M^2} \end{aligned} \quad (8)$$

with

$$\begin{aligned} P^o(M^2) &= \langle q_i^\dagger q_i \rangle - 4 \left(\frac{m_Q^2}{2M^4} - \frac{1}{M^2} \right) \langle q_i^\dagger D_0^2 q_i \rangle \\ &\quad - \frac{1}{M^2} \langle q_i^\dagger g_s \sigma \cdot G q_i \rangle. \end{aligned} \quad (9)$$

The QCD condensates corresponding to the even and odd parts of the spectral function occur in the polynomials $P^e(M^2)$ and $P^o(M^2)$, given by Eqs. (7) and (9), respectively. For the polynomial $P^e(M^2)$, the first and second terms arise due to the quark-antiquark condensate

($\langle \bar{q}_i q_i \rangle$), retaining the term proportional to the light quark mass (m_{q_i}) [16] as given by the second term. The first term, due to the large mass of the heavy quark (m_Q) is the dominant contribution to the mass of the open charm (bottom) meson, arising from the dimension 3 quark-antiquark condensate in the operator product expansion. The strong coupling constant, α_s is given as $\alpha_s = 4\pi/[(11 - 2N_f/3) \ln(\mu^2/\Lambda_{\text{QCD}}^2)]$, with $\Lambda_{\text{QCD}} = 500$ MeV, and, $N_f = 4$ (5) for the study of the open charm (bottom) mesons. μ , the renormalization scale, is taken to be mass of the largest quark in the system, i.e., the mass of the charm (bottom) quark, for investigating the mass of the open charm (bottom) mesons [25], γ_E is the Euler constant. The other terms are the higher mass dimension operators, which include the dimension-4 scalar and twist-2 gluon condensates, the mixed quark-gluon operator (the abbreviation $\sigma \cdot G$ is used for $\sigma_{\mu\nu} G^{\mu\nu}$), and terms proportional to the operators $\langle q_i^\dagger i D_0 q_i \rangle$ and $\langle \bar{q}_i D_0^2 q_i \rangle$. In the term proportional to the twist-2 gluon condensate operator, u^μ

is the four-velocity of the hadronic medium. We study the spectral functions of the open heavy flavor mesons in the rest frame of the hadronic medium, for which $u^\mu = (1, \vec{u} = 0)$. In the odd part of the Borel transformed spectral function, the dominant contribution in Eq. (9) is from the number density of the quark ($\langle q_i^\dagger q_i \rangle$), along with the contributions from the dimension 5 operators, $\langle \bar{q}_i D_0^2 q_i \rangle$ and $\langle q_i^\dagger g_s \sigma \cdot G q_i \rangle$. The masses for the particles and antiparticle, $m_\pm$, corresponding to the open heavy flavor mesons, are solved by using the spectral function given by Eq. (5) in the Eqs. (3) and (4) [25]. The expressions for the masses of the particle and antiparticle are given as [26]

$$m_\pm = \left[\frac{1}{4} \left(\frac{gf' - fg'}{f^2 + gg'} \right)^2 - \left(\frac{ff' + g'g}{f^2 + gg'} \right) \right]^{1/2} \pm \frac{1}{2} \left(\frac{gf' - fg'}{f^2 + gg'} \right) \quad (10)$$

where, $f = \frac{1}{\pi} \int_{m_Q^2}^{\tilde{s}_0} d\tilde{s} \text{Im} \Pi_{OPE}^{per}(\tilde{s}, \vec{q} = 0) e^{-\tilde{s}/M^2} + e^{-m_Q^2/M^2} P^e(M^2)$, with $P^e(M^2)$ given by Eq. (7), $g = e^{-m_Q^2/M^2} P^o(M^2)$, the odd part of the spectral function as given by Eq. (8), and, f' and g' denote the derivatives of f and g with respect to $1/M^2$. As can be observed from Eq. (10), the mass difference between the particles (D^0 , D^+ and D_s^+) and their antiparticles ($\bar{D}^0$, D^- and D_s^-) for the charm sector, and between the particles (B^- , B^0 and $\bar{B}_s^0$) and their antiparticles (B^+ , B^0 and B_s^0) for the bottom sector, in the hadronic medium arise from the values of $\langle q_i^\dagger q_i \rangle$, $\langle q_i^\dagger i D_0^2 q_i \rangle$ and $\langle q_i^\dagger g_s \sigma \cdot G q_i \rangle$ of the odd part of the spectral function (g), hence due to the expectation values of the QCD operators $\langle q_i^\dagger q_i \rangle$, $\langle q_i^\dagger D_0^2 q_i \rangle$ and $\langle q_i^\dagger g_s \sigma \cdot G q_i \rangle$ in $P^o(M^2)$, given by Eq. (9). The in-medium QCD condensates are calculated within a chiral SU(3) model, as described in the following section.

III. CHIRAL SU(3) MODEL

In this section, we briefly describe the chiral SU(3) model [40–42,64], which is used to obtain the in-medium QCD condensates in the hot asymmetric strange hadronic matter. The expectation values of these QCD operators are central to the calculation of the in-medium masses of the open charm and open bottom mesons within the framework of the QCD sum rule approach, as described in the previous section. The chiral SU(3) model [40–42] is based on a nonlinear realization of chiral symmetry [71–73] and incorporates the broken scale invariance of QCD through the introduction of a scalar dilaton field, χ [38,39]. The Lagrangian density of chiral SU(3) model is written as [40]

$$\mathcal{L} = \mathcal{L}_{\text{kin}} + \sum_{W=X,Y,V,A,u} \mathcal{L}_{BW} + \mathcal{L}_{\text{vec}} + \mathcal{L}_0 + \mathcal{L}_\chi + \mathcal{L}_{SB}. \quad (11)$$

In Eq. (11), $\mathcal{L}_{\text{kin}}$ is the kinetic energy term, $\mathcal{L}_{BW}$ is the baryon-meson interaction term, where, $W = X, Y, V, A, u$

correspond to the interactions of the baryons with the scalar, pseudoscalar singlet, vector meson, axial-vector meson, and pseudoscalar octet mesons, respectively. The baryon-spin-0 interaction terms generate the baryon masses. $\mathcal{L}_{\text{vec}}$ describes the dynamical mass generation of the vector mesons via couplings to the scalar mesons, containing additionally quartic self-interactions of the vector fields. $\mathcal{L}_0$ contains the meson-meson interaction terms, $\mathcal{L}_\chi$ is the QCD scale breaking term in the Lagrangian density, including a logarithmic potential of the scalar dilaton field, χ , and $\mathcal{L}_{SB}$ describes the explicit chiral symmetry breaking. In the present study of isospin asymmetric strange hadronic matter at finite temperature, the mean field approximation is used, in which the meson fields are replaced by their expectation values. The meson fields that have nonzero expectation values are the scalar (σ, ζ, δ) and the vector fields, $\omega^\mu \rightarrow \omega \delta^{\mu 0}$, $\rho^{\mu a} \rightarrow \rho \delta^{\mu 0} \delta^{a3}$, and $\phi^\mu \rightarrow \phi \delta^{\mu 0}$, and the expectation values of the other mesons are zero. We also use the approximations $\bar{\psi}_i \psi_j \rightarrow \delta_{ij} \langle \bar{\psi}_i \psi_i \rangle \equiv \delta_{ij} \rho_i^s$ and $\bar{\psi}_i \gamma^\mu \psi_j \rightarrow \delta_{ij} \delta^{\mu 0} \langle \bar{\psi}_i \gamma^0 \psi_i \rangle \equiv \delta_{ij} \delta^{\mu 0} \rho_i$, where ρ_i^s and ρ_i are the scalar and number density of baryon of species, i . In the above approximation, the terms that contribute to the baryon-meson interaction, $\mathcal{L}_{BW}$ of Eq. (11) arise from the interactions due to the scalar (S) and the vector (V) mesons, and are given as

$$\mathcal{L}_{BS} + \mathcal{L}_{BV} = \sum_i \bar{\psi}^i (g_{\sigma i} \sigma + g_{\zeta i} \zeta + g_{\delta i} \delta - g_{\omega i} \gamma^0 \omega - g_{\rho i} \gamma^0 \rho - g_{\phi i} \gamma^0 \phi) \psi^i. \quad (12)$$

The explicit symmetry breaking term in the chiral model is given by

$$\mathcal{L}_{SB} = - \left(\frac{\chi}{\chi_0} \right)^2 \text{Tr} \left[\text{diag} \left(\frac{m_\pi^2 f_\pi}{2} (\sigma + \delta), \frac{m_\pi^2 f_\pi}{2} (\sigma - \delta), \left(\sqrt{2} m_K^2 f_K - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi \right) \zeta \right) \right], \quad (13)$$

which reduces to Eq. (8), and in QCD, this term is given as $\mathcal{L}_{SB}^{\text{QCD}} = -\text{Tr}[\text{diag}(m_u \bar{u}u, m_d \bar{d}d, m_s \bar{s}s)]$. In the above, the matrices, whose traces give the corresponding Lagrangian densities, have been explicitly written down, and a comparison of their matrix elements relates the quark condensates to the values of the scalar fields as [65]

$$\begin{aligned} m_u \langle \bar{u}u \rangle &= \left(\frac{\chi}{\chi_0} \right)^2 \frac{m_\pi^2 f_\pi}{2} (\sigma + \delta), \\ m_d \langle \bar{d}d \rangle &= \left(\frac{\chi}{\chi_0} \right)^2 \frac{m_\pi^2 f_\pi}{2} (\sigma - \delta), \\ m_s \langle \bar{s}s \rangle &= \left(\frac{\chi}{\chi_0} \right)^2 \left(\sqrt{2} m_K^2 f_K - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi \right) \zeta. \end{aligned} \quad (14)$$

The medium modifications of the quark condensates $\langle \bar{q}_i q_i \rangle$, with $q_i \equiv u, d, s$ for $i = 1, 2, 3$, are obtained from the medium modifications of the scalar fields (σ and ζ and δ).

The scalar gluon condensate $\langle \frac{\alpha_s}{\pi} G_{\mu\nu}^a G^{a\mu\nu} \rangle$ and the expectation value of the twist-2 tensorial gluon operator, $\langle \frac{\alpha_s}{\pi} (G_{\mu\sigma}^a G_{\nu}^{\sigma} u^\mu u^\nu - \frac{1}{4} G_{\mu\sigma}^a G^{a\mu\sigma}) \rangle$, are obtained from the in-medium value of the dilaton field, χ [19], which is incorporated in the chiral SU(3) model to simulate the broken scale invariance of QCD. The Lagrangian density for the scalar dilaton field, χ , is given as [19]

$$\mathcal{L}_\chi = \frac{1}{2}(\partial_\mu \chi)(\partial^\mu \chi) - k_4 \chi^4 - \frac{1}{4} \chi^4 \ln\left(\frac{\chi^4}{\chi_0^4}\right) + \frac{d}{3} \chi^4 \ln\left(\left(\frac{(\sigma^2 - \delta^2)\zeta}{\sigma_0^2 \zeta_0}\right)\left(\frac{\chi}{\chi_0}\right)^3\right), \quad (15)$$

where the last two terms in the above Lagrangian density are the scale breaking terms. The energy momentum tensor obtained from the above Lagrangian density is given as

$$T_{\mu\nu} = (\partial_\mu \chi) \left(\frac{\partial \mathcal{L}_\chi}{\partial (\partial^\nu \chi)} \right) - g_{\mu\nu} \mathcal{L}_\chi \equiv (\partial_\mu \chi)(\partial_\nu \chi) - g_{\mu\nu} \mathcal{L}_\chi. \quad (16)$$

Accounting for the finite mass of the quarks, the energy momentum of the QCD can be written in terms of a symmetric traceless part and a trace part as

$$T_{\mu\nu} = -ST(G_{\mu\sigma}^a G_{\nu}^{a\sigma}) + \frac{g_{\mu\nu}}{4} \left(\frac{\beta_{\text{QCD}}}{2g} G_{\sigma\kappa}^a G^{a\sigma\kappa} + \sum_i m_i \bar{q}_i q_i \right), \quad (17)$$

writing

$$\left\langle \frac{\alpha_s}{\pi} G_{\mu\sigma}^a G_{\nu}^{a\sigma} \right\rangle = \left(u_\mu u_\nu - \frac{g_{\mu\nu}}{4} \right) G_2, \quad (18)$$

where u_μ is the four-velocity of the nuclear medium, taken as $u_\mu = (1, 0, 0, 0)$, we obtain the energy momentum tensor in QCD as

$$T_{\mu\nu} = -\left(\frac{\pi}{\alpha_s}\right) \left(u_\mu u_\nu - \frac{g_{\mu\nu}}{4} \right) G_2 + \frac{g_{\mu\nu}}{4} \left(\frac{\beta_{\text{QCD}}}{2g} G_{\sigma\kappa}^a G^{a\sigma\kappa} + \sum_i m_i \bar{q}_i q_i \right). \quad (19)$$

We calculate the expectation values of the scalar and twist-2 gluon operators by equating the energy momentum tensor for the dilaton field with that of QCD, multiplying both sides by $g^{\mu\nu}$ and $(u^\mu u^\nu - \frac{g^{\mu\nu}}{4})$ [$u^\mu = (1, 0, 0, 0)$ is the four-velocity of the hadronic medium at rest] respectively, and summing over the indices (μ and ν) [19]. We then obtain [19]

$$T_\mu{}^\mu = \sum_i m_i \bar{q}_i q_i + \left\langle \frac{\beta_{\text{QCD}}}{2g} G_{\mu\nu}^a G^{a\mu\nu} \right\rangle = -(1-d)\chi^4 \quad (20)$$

and,

$$\begin{aligned} & \left\langle \frac{\alpha_s}{\pi} G_{\mu\sigma}^a G_{\nu}^{a\sigma} \right\rangle \left(u^\mu u^\nu - \frac{g^{\mu\nu}}{4} \right) \\ & \equiv \frac{3}{4} G_2 = \frac{3\alpha_s}{4\pi} \left[-(1-d+4k_4)(\chi^4 - \chi_0^4) \right. \\ & \quad \left. - \chi^4 \ln\left(\frac{\chi^4}{\chi_0^4}\right) + \frac{4}{3} d \chi^4 \ln\left(\left(\frac{(\sigma^2 - \delta^2)\zeta}{\sigma_0^2 \zeta_0}\right)\left(\frac{\chi}{\chi_0}\right)^3\right) \right]. \quad (21) \end{aligned}$$

In Eq. (20), the first term of the trace of the energy-momentum tensor within the chiral SU(3) model is the negative of the explicit chiral symmetry breaking term, $\mathcal{L}_{SB}$ given by Eq. (13). It might be noted here that in deriving the relation of the twist-2 gluon operator with the dilaton field [given by Eq. (21)], we have assumed that the dilaton field, χ , is nonrelativistic, and hence have assumed that $\langle (\partial_i \chi)(\partial_j \chi) \rangle_{\text{medium}} \simeq \langle (\partial_i \chi)(\partial_j \chi) \rangle_{\text{vac}}$. Using the QCD β function at one loop level, for N_c colors and N_f flavors

$$\beta_{\text{QCD}}(g) = -\frac{11N_c g^3}{48\pi^2} \left(1 - \frac{2N_f}{11N_c} \right), \quad (22)$$

for $N_c = 3$ and $N_f = 3$, the scalar gluon condensate is obtained as

$$\begin{aligned} \left\langle \frac{\alpha_s}{\pi} G_{\mu\nu}^a G^{a\mu\nu} \right\rangle &= \frac{8}{9} \left[(1-d)\chi^4 + \left(\frac{\chi}{\chi_0}\right)^2 \left(m_\pi^2 f_\pi \sigma \right. \right. \\ & \quad \left. \left. + \left(\sqrt{2} m_k^2 f_k - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi \right) \zeta \right) \right]. \quad (23) \end{aligned}$$

The parameter d in Eq. (20) originates from the second logarithmic term of Eq. (15). The parameter d , along with the other parameters corresponding to the meson-meson interaction Lagrangian density, $\mathcal{L}_0$ of the general Lagrangian given by (11), are fitted so as to ensure extrema in the vacuum for the σ , ζ , and χ field equations, to reproduce the vacuum masses of the η and η' mesons, the mass of the σ meson around 500 MeV, and pressure, $p(\rho_0) = 0$, with ρ_0 as the nuclear matter saturation density [40,47]. The scalar fields (the nonstrange isoscalar field, σ , the strange isoscalar field, ζ , and the isovector field δ), the dilaton field, χ , and the vector fields (nonstrange isoscalar ω , strange isoscalar ϕ , and nonstrange isovector ρ) are obtained by minimizing the thermodynamic potential Ω for given values of temperature, baryonic density, ρ_B , the isospin asymmetry parameter, $\eta = -\frac{\sum_i I_{3i} \rho_i}{\rho_B}$, and the strangeness fraction, $f_s = \frac{\sum_i |s_i| \rho_i}{\rho_B}$, where, I_{3i} and $|s_i|$ denote the third component of the isospin quantum number and the number of strange quarks in the i^{th} baryon, respectively.

The effects of the scalar and twist-2 gluon condensates, as obtained from the medium modification of the dilaton field, χ , within the chiral SU(3) model, on the masses of the vector and pseudoscalar charmonium ground states (J/ψ and η_c) were studied in Ref. [19] and on the masses of the heavy quarkonium (charmonium and bottomonium) states in Ref. [68], using the QCD sum rule approach. In the present work, the effects of the medium modifications of the QCD operators, up to dimension 5, on the masses of the open charm and open bottom mesons are investigated within the QCD sum rule approach. The nonstrange ($\langle\bar{u}u\rangle$, $\langle\bar{d}d\rangle$) and strange ($\langle\bar{s}s\rangle$) quark condensates are obtained from the scalar fields using Eq. (14) and the scalar and twist-2 gluon condensates are calculated from the dilaton field, χ , using Eqs. (20) and (21) within the chiral SU(3) model. The dimension 5 mixed quark-gluon condensate $\langle\bar{q}_i g_s \sigma \cdot G q_i\rangle$ in the operator product expansion is obtained from the in-medium quark condensate $\langle\bar{q}_i q_i\rangle$, with $i = 1, 2, 3 \equiv u, d, s$, as computed within the chiral SU(3) model, by using the relation

$$\langle\bar{q}_i g_s \sigma \cdot G q_i\rangle = 2\lambda_q^2 \langle\bar{q}_i q_i\rangle. \quad (24)$$

The values of the parameter λ_q estimated using different methods are, however, observed to be widely different [69,74–79]. Using the leading order approximations for the parton distribution functions [80], the values of the higher order operators (or combinations of these operators) used in the present investigation are as follows [81]. The estimates for $q_i = u, d$ are

$$\langle q_i^\dagger i D_0 q_i \rangle = 0.18 \text{ GeV} \rho_B, \quad (25)$$

$$-\langle\bar{q}_i D_0^2 q_i\rangle + \frac{1}{8} \langle\bar{q}_i g_s \sigma \cdot G q_i\rangle = 0.3 \text{ GeV}^2 \rho_B, \quad (26)$$

and

$$-\langle q_i^\dagger D_0^2 q_i\rangle + \frac{1}{12} \langle q_i^\dagger g_s \sigma \cdot G q_i\rangle = 0.031 \text{ GeV}^2 \rho_B, \quad (27)$$

for the case of symmetric nuclear matter. For the strange quark ($q_i = s$), the estimate for the above operator is taken to be $\langle s^\dagger i D_0 s \rangle = 0.06 \text{ GeV} \rho_B$ [69]. For the consideration of the strange-heavy flavor mesons, the estimates of the combinations of the operators given by Eqs. (26) and (27) (for $q_i = s$) are given as [25]

$$\begin{aligned} & -\langle\bar{s} D_0^2 s\rangle + \frac{1}{8} \langle\bar{s} g_s \sigma \cdot G s\rangle \\ & = y \left(-\langle q_i^\dagger D_0^2 q_i\rangle + \frac{1}{12} \langle q_i^\dagger g_s \sigma \cdot G q_i\rangle \right) \end{aligned} \quad (28)$$

and

$$\begin{aligned} & -\langle s^\dagger D_0^2 s\rangle + \frac{1}{12} \langle s^\dagger g_s \sigma \cdot G s\rangle \\ & = y \left(-\langle q_i^\dagger D_0^2 q_i\rangle + \frac{1}{12} \langle q_i^\dagger g_s \sigma \cdot G q_i\rangle \right), \end{aligned} \quad (29)$$

where, $q_i = u$ or d in the rhs of the above equations, which have equal values for the operators for symmetric nuclear matter. In Eqs. (28) and (29), y is the strangeness content of the nucleon, given as

$$y = \langle\bar{s}s\rangle_{\text{med}} / \langle\bar{q}q\rangle_{\text{med}}, \quad (30)$$

with $\langle\bar{s}s\rangle_{\text{med}} = \langle\bar{s}s\rangle - \langle\bar{s}s\rangle_{\text{vac}}$, and $\langle\bar{q}q\rangle_{\text{med}} = (\langle\bar{u}u\rangle_{\text{med}} + \langle\bar{d}d\rangle_{\text{med}})/2$ is the average of the quark-antiquark condensates for u and d quarks after subtracting the vacuum contributions. For the isospin asymmetric nuclear (hyperonic) matter as considered in the present investigation, we generalize the above relations as follows. For the non-strange quarks, $q_i = u, d$, for $i = 1, 2$,

$$-\langle\bar{q}_i D_0^2 q_i\rangle + \frac{1}{8} \langle\bar{q}_i g_s \sigma \cdot G q_i\rangle = 0.3 \text{ GeV}^2 \left(\frac{\langle\bar{q}_i q_i\rangle_{\text{med}}}{\langle\bar{q}q\rangle_{\text{med}}} \right) \rho_B, \quad (31)$$

$$\begin{aligned} & -\langle q_i^\dagger D_0^2 q_i\rangle + \frac{1}{12} \langle q_i^\dagger g_s \sigma \cdot G q_i\rangle \\ & = 0.031 \text{ GeV}^2 \left(\frac{\langle\bar{q}_i q_i\rangle_{\text{med}}}{\langle\bar{q}q\rangle_{\text{med}}} \right) \rho_B. \end{aligned} \quad (32)$$

For the QCD operators involving the strange quark, we take

$$-\langle\bar{s} D_0^2 s\rangle + \frac{1}{8} \langle\bar{s} g_s \sigma \cdot G s\rangle = 0.3 \text{ GeV}^2 \rho_B y, \quad (33)$$

$$-\langle s^\dagger D_0^2 s\rangle + \frac{1}{12} \langle s^\dagger g_s \sigma \cdot G s\rangle = 0.031 \text{ GeV}^2 \rho_B y, \quad (34)$$

where y is the strangeness content of the nucleon defined by Eq. (30). The estimate for the operator $\langle q_i^\dagger g_s \sigma \cdot G q_i\rangle = 0.33 \text{ GeV}^2 \rho_B$, for $q_i = u$ or d for the symmetric nuclear matter [25] is generalized to the isospin asymmetric case in the present investigation as

$$\langle q_i^\dagger g_s \sigma \cdot G q_i\rangle = 0.33 \text{ GeV}^2 \left(\frac{\langle\bar{q}_i q_i\rangle_{\text{med}}}{\langle\bar{q}q\rangle_{\text{med}}} \right) \rho_B, \quad (35)$$

for $q_i = u$ and d , and, for the operator with strange quark given as,

$$\langle s^\dagger g_s \sigma \cdot G s\rangle = 0.33 \text{ GeV}^2 \rho_B y. \quad (36)$$

Using the estimates for the combinations of the operators given by Eqs. (32) and (34), one obtains the estimates for $\langle q_i^\dagger D_0^2 q_i\rangle$, for $(q_i, i = 1, 2, 3) \equiv (u, d, s)$. Using the values

for the QCD operators for the even and odd parts of the spectral function as obtained above, the masses of the open charm and open bottom mesons are computed within the QCD sum rule approach.

IV. RESULTS AND DISCUSSIONS

In this section, we discuss the results obtained for the in-medium masses of the open charm (D , $\bar{D}$, D_s and $\bar{D}_s$) and open bottom (B , $\bar{B}$, B_s and $\bar{B}_s$) mesons, and their anti-particles in asymmetric strange hadronic matter at finite temperatures. These are computed using the QCD sum rule approach retaining QCD operators up to dimension 5, with the medium dependence, e.g., the temperature, isospin asymmetry, strangeness, and density dependence of these operators calculated using a chiral SU(3) model. In the present investigation, we choose the quark mass parameters to be $m_u = m_d = 6$ MeV, $m_s = 185$ MeV, $m_c = 1.5$ GeV and $m_b = 4.7$ GeV. The light quark condensates ($\langle \bar{q}_i q_i \rangle$, ($i = 1, 2, 3 \equiv (u, d, s)$)) are obtained from the values of the scalar fields, σ , ζ , and δ in the chiral SU(3) model by using Eq. (14). For the chosen values of the masses of the quarks, the vacuum values of the light quark

condensates are obtained as $\langle \bar{q}q \rangle_{\text{vac}} = \langle \bar{u}u \rangle_{\text{vac}} = \langle \bar{d}d \rangle_{\text{vac}} = (-241.1 \text{ MeV})^3$ and $\langle \bar{s}s \rangle_{\text{vac}} = (-288.1 \text{ MeV})^3$. The non-strange quark condensate, $\langle \bar{q}q \rangle_{\text{vac}}$ may be compared with the value of $(-225 \text{ MeV})^3$ [81] and the average estimate of $(-272 \text{ MeV})^3$ of the flavor lattice averaging group [82]. As may be observed, for the strange quark condensate, the value of $(-288.1 \text{ MeV})^3$ in the present investigation has close agreement with the estimates of $\langle \bar{s}s \rangle_{\text{vac}} = (-290 \text{ MeV})^3$ [83] and $\langle \bar{s}s \rangle_{\text{vac}} = (-296 \text{ MeV})^3$ [84] from lattice calculations. The value of $\langle \bar{s}s \rangle_{\text{vac}}$ of $(-288.1 \text{ MeV})^3$ in the present work as well as from the lattice calculations [83,84] are much higher than the value $\langle \bar{s}s \rangle_{\text{vac}} = (-208 \text{ MeV})^3$ obtained using the parametrization $\langle \bar{s}s \rangle = 0.8 \langle \bar{q}q \rangle$ used in Ref. [81]. In the present work, the light quark condensates in vacuum are observed to be modified to the values $\langle \bar{u}u \rangle = \langle \bar{d}d \rangle = (-206.1 \text{ MeV})^3$ and $\langle \bar{s}s \rangle = (-277.3 \text{ MeV})^3$ in symmetric nuclear matter at nuclear matter saturation density (ρ_0). The strangeness content nucleon, y , as defined in Eq. (30) has a value of 0.49 for $\rho_B = \rho_0$ in the chiral model as used in the present investigation. The strangeness content, y , is related to the nucleon σ term as [69,85]

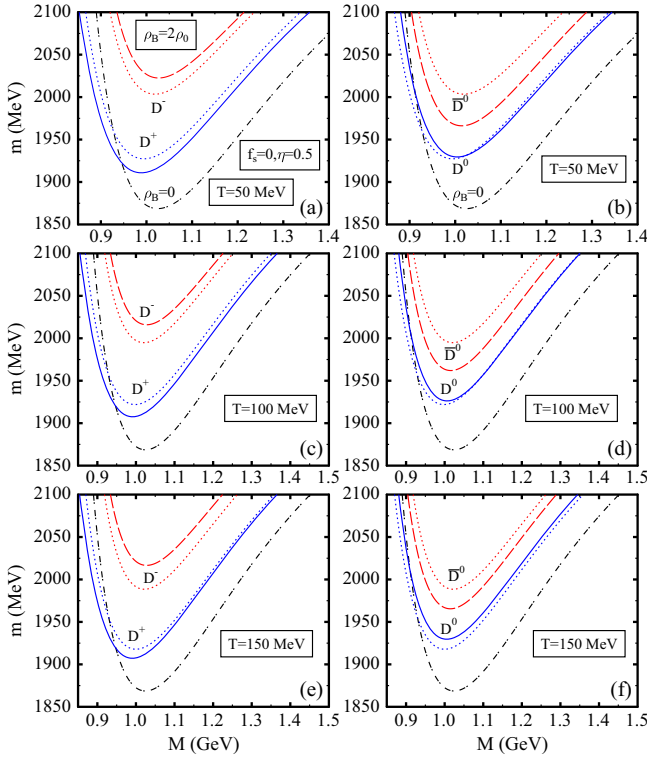


FIG. 1. Masses of $D^+(D^-)$ [in panels (a), (c) and (e)] and $D^0(\bar{D}^0)$ [in panels (b), (d) and (f)] mesons are plotted as functions of the Borel mass M for $T = 50, 100$, and 150 MeV in nuclear matter ($f_s = 0$) for baryon density $\rho_B = 2\rho_0$ and $\eta = 0.5$, as the solid (dashed) lines. These are compared to the case of $\eta = 0$ (shown as dotted lines). The Borel curve for $\rho_B = 0$ is also shown in the figure as the dot-dashed line.

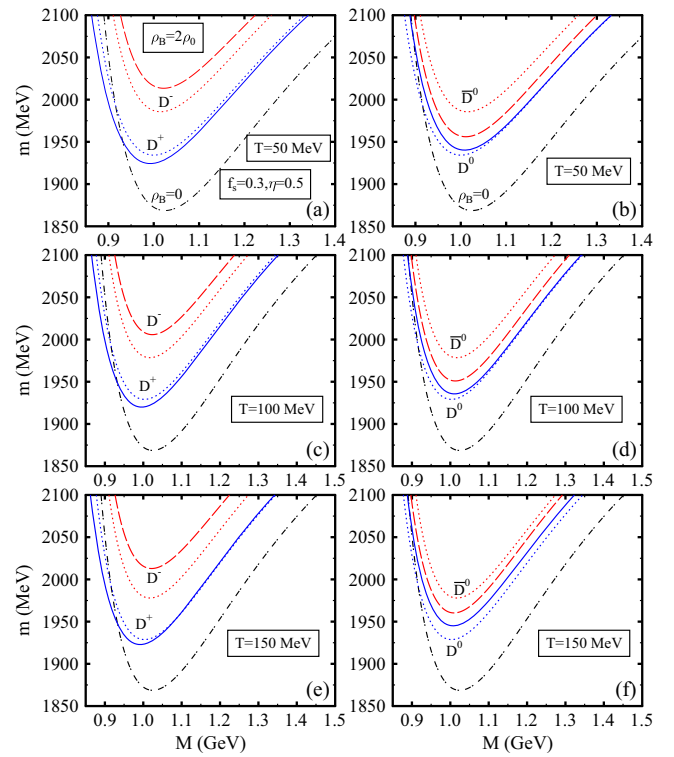


FIG. 2. Masses of $D^+(D^-)$ (in panels (a), (c) and (e)) and $D^0(\bar{D}^0)$ (in panels (b), (d) and (f)) mesons are plotted as functions of the Borel mass M for $T = 50, 100$ and 150 MeV in hyperonic matter (with $f_s = 0.3$) for baryon density $\rho_B = 2\rho_0$ and $\eta = 0.5$, as the solid (dashed) lines. These are compared to the case of $\eta = 0$ (shown as dotted lines). The Borel curve for $\rho_B = 0$ is also shown in the figure as the dot-dashed line.

$$y = \frac{\sigma_N^0}{\sigma_N} - 1, \quad (37)$$

with the nucleon σ term defined as

$$\sigma_N = (m_u + m_d) \frac{(\langle \bar{q}q \rangle_{\text{med}} - \langle \bar{q}q \rangle_{\text{vac}})}{\rho_B}. \quad (38)$$

In Eq. (37), σ_N^0 is the value of the nucleon σ term σ_N , when the strangeness content of the nucleon (y) is taken to be zero. Using the estimates for σ_N and σ_N^0 as 45 and 35 MeV respectively [69,85], the value of y is obtained to be 0.286 at the nuclear matter saturation density. The value for y estimated using lattice calculations in Ref. [86] is 0.36. The value of y is thus not well estimated, and is varied within the range 0–0.5 [87]. We might note here that the value of the nucleon σ term [using the Eq. (38)] for the obtained value of $y = 0.49$ in the present investigation, is calculated to be $\sigma_N = 54.8$ MeV. This value of the present work may be compared to the often used value of 45 MeV [69].

In the present work, the value for the scalar gluon condensate, $\langle \frac{\alpha_s}{\pi} G_{\mu\nu}^a G^{a\mu\nu} \rangle$, obtained in vacuum is $(373.1 \text{ MeV})^4 = 0.019 \text{ GeV}^4$, which is modified to

$(370.6 \text{ MeV})^4 = 0.0189 \text{ GeV}^4$ at $\rho_B = \rho_0$ in symmetric nuclear matter. The value of 0.019 GeV^4 in vacuum may be compared with the lattice results of 0.028 GeV^4 [88] and 0.077 GeV^4 [89]. The estimate of the scalar gluon condensate is thus observed to be larger in the lattice calculations as well as present work, as compared to the value of 0.012 GeV^4 obtained in the QCD sum rule calculations in Ref. [90]. As has already been mentioned, the parameter λ_q^2 relating the mixed quark-gluon condensate to the quark condensate [given by Eq. (24)] is not well estimated [69]. The estimates for λ_q^2 in vacuum are 0.4 GeV^2 using the standard QCD sum rule approach [74], 1.25 GeV^2 [75] and $0.5\text{--}0.55 \text{ GeV}^2$ [76,77] estimated using lattice QCD, and $\lambda_q^2 \sim 1.2 \text{ GeV}^2$ in a QCD sum rule analysis of the pion wave function using nonlocal condensates [78] as well as within an instanton liquid model [79]. In the present work, we take the value of this parameter to be $\lambda_q^2 = 1.2 \text{ GeV}^2$. For the charm sector, the vacuum masses for the $D(D^0, D^+)$ and $\bar{D}(\bar{D}^0, D^-)$ mesons are obtained to be 1869 MeV for $\tilde{s}_0 = 7.5 \text{ GeV}^2$, and the vacuum mass of $D_s^\pm$ is obtained to be around 1916 MeV for $\tilde{s}_0 = 15 \text{ GeV}^2$. For the bottom sector, the

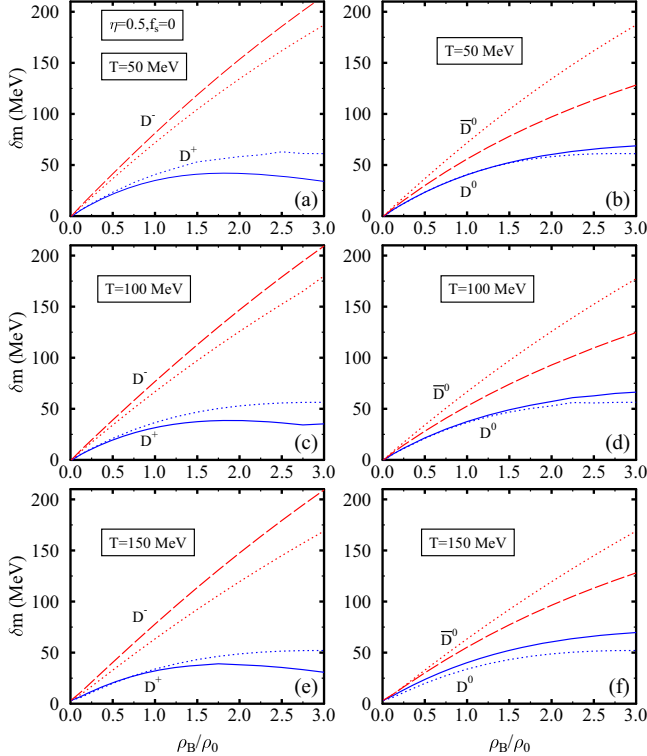


FIG. 3. Mass shifts of $D^+(D^-)$ [in panels (a), (c), and (e)] and $D^0(\bar{D}^0)$ [in panels (b), (d), and (f)] mesons are plotted as functions of the baryon density in units of nuclear matter saturation density for $T = 50, 100$, and 150 MeV in nuclear matter ($f_s = 0$) for $\eta = 0.5$, as the solid (dashed) lines, and are compared to the case of $\eta = 0$ (shown as dotted lines).

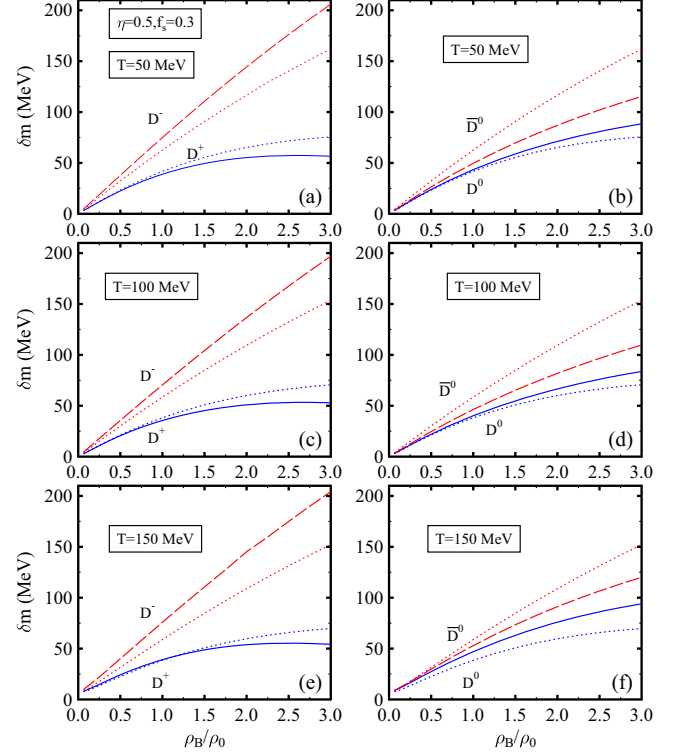


FIG. 4. Mass shifts of $D^+(D^-)$ (in panels (a), (c) and (e)) and $D^0(\bar{D}^0)$ (in panels (b), (d) and (f)) mesons are plotted as functions of the baryon density in units of nuclear matter saturation density for $T = 50, 100$ and 150 MeV in hyperonic matter (with $f_s = 0.3$) for $\eta = 0.5$, as the solid (dashed) lines, and, are compared to the case of $\eta = 0$ (shown as dotted lines).

vacuum masses for the $B(B^+, B^0)$ and $\bar{B}(B^-, \bar{B}^0)$ mesons are obtained to be 5277.8 MeV for $\tilde{s}_0 = 37.8 \text{ GeV}^2$, and the mass of $B_s^0(\bar{B}_s^0)$ is obtained to be around 5366.455 MeV for $\tilde{s}_0 = 42.3 \text{ GeV}^2$. There is a splitting between the masses of the open heavy flavor particles ($D^0, D^+, D_s^+, B^-, \bar{B}^0$ and $\bar{B}_s^0$) and their antiparticles ($\bar{D}^0, D^-, D_s^-, B^+, B^0$ and B_s^0), which arises from the contributions from the number density of the i -th quark ($\langle q_i^\dagger q_i \rangle$), along with the contributions from the dimension 5 operators, $\langle \bar{q}_i D_0^2 q_i \rangle$ and $\langle q_i^\dagger g_s \sigma \cdot G q_i \rangle$ of the odd part of the spectral function.

In Figs. 1 and 2, we plot the masses of the charged and neutral open charm mesons as functions of the Borel mass, M , for the cases of nuclear matter ($f_s = 0$) and strange hadronic matter (with $f_s = 0.3$) for $\rho_B = 2\rho_0$ and $T = 50, 100$, and 150 MeV. These are shown for the isospin asymmetry parameter, $\eta = 0.5$, and are compared to the case of isospin symmetric case ($\eta = 0$). The vacuum case is shown as the dot-dashed line. At finite densities, shown for $\rho_B = 2\rho_0$ in these figures, the mass of the antiparticle $D^-(\bar{D}^0)$ is observed to be higher than $D^+(D^0)$ both in the hot nuclear matter and strange hadronic matter. As can be observed from these figures, the effect of isospin

asymmetry is to reduce (increase) the mass of $D^+(D^0)$ within the iso-doublet (D^0, D^+), and of $\bar{D}^0(D^-)$ within the iso-doublet ($D^-, \bar{D}^0$) for both nuclear matter as well as in strange hadronic matter. The effect from isospin asymmetry is observed to be much smaller for the neutral particle D^0 as compared to the effects on its antiparticle $\bar{D}^0$.

The shifts in the masses of the charged ($D^\pm$) and neutral ($D^0, \bar{D}^0$) from the vacuum values are plotted as functions of the baryon density in units of nuclear matter saturation density, in Figs. 3 and 4 for nuclear matter ($f_s = 0$) and strange hadronic matter (with $f_s = 0.3$), respectively. The effect from isospin asymmetry is observed to be appreciable for higher values of the baryon density, especially for $\bar{D}^0$ meson. In isospin symmetric ($\eta = 0$) nuclear (hyperonic) matter, the masses of open charm mesons within the $D(D^0, D^+)$ and $\bar{D}(D^-, \bar{D}^0)$ doublets are degenerate. At $\rho_B = \rho_0$, the difference in the mass of the particle $D^0(D^+)$ ($\sim 1912 \text{ MeV}$) and the mass of its antiparticle $\bar{D}^0(D^-)$ ($\sim 1944 \text{ MeV}$) is around 32 MeV, which may be compared to the mass difference of around 30 MeV in Ref. [25] for the isospin symmetric ($\eta = 0$) nuclear matter. In the presence

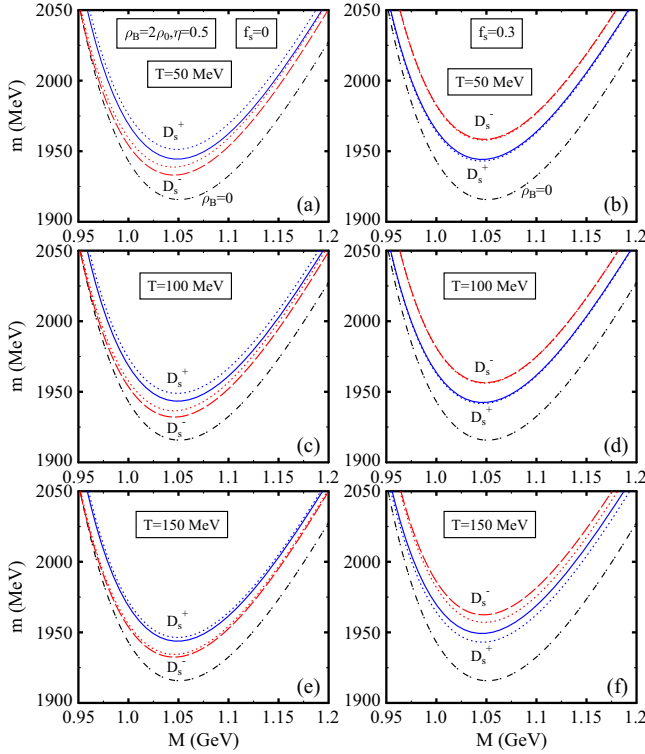


FIG. 5. Masses of D_s^+ and D_s^- mesons are plotted as functions of the Borel mass M for $T = 50, 100$, and 150 MeV in panels (a), (c), and (e) and in panels (b), (d), and (f), in nuclear ($f_s = 0$) and hyperonic (with $f_s = 0.3$) matter for baryon density $\rho_B = 2\rho_0$ and $\eta = 0.5$, as the solid (dashed) lines, and are compared to the case of $\eta = 0$ (shown as dotted lines). The Borel curve for $\rho_B = 0$ is shown as the dot-dashed line.

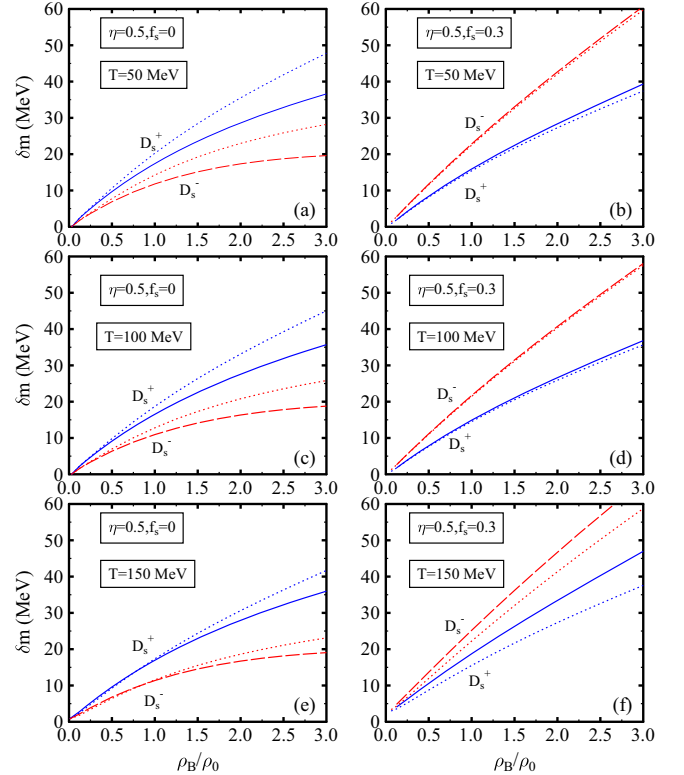


FIG. 6. Mass shifts of D_s^+ and D_s^- mesons are plotted as functions of the baryon density in units of nuclear matter saturation density for $\eta = 0.5$ [in panels (a), (c), and (e) for nuclear matter ($f_s = 0$) and in panels (b), (d), and (f) for strange hadronic matter with $f_s = 0.3$] for $T = 50, 100$, and 150 MeV as the solid (dashed) lines, and are compared to the case of $\eta = 0$ (shown as dotted lines).

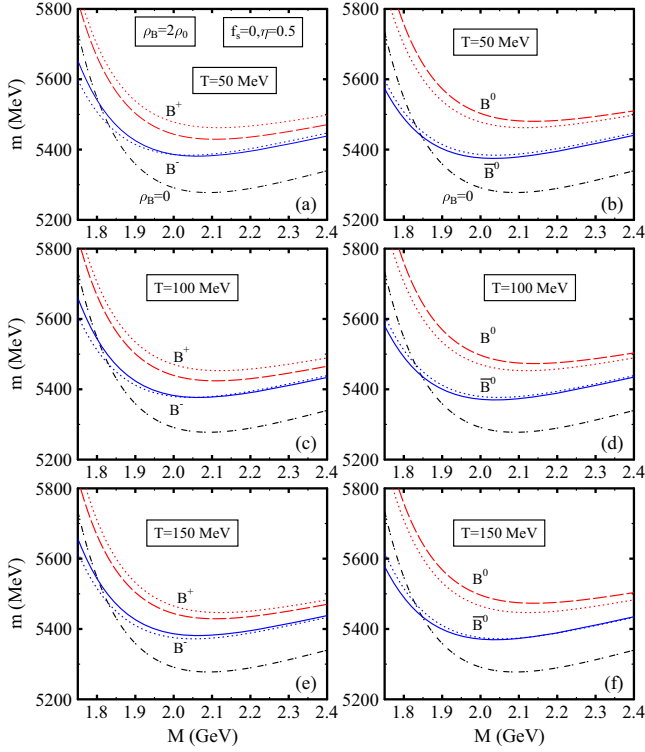


FIG. 7. Masses of $B^-(B^+)$ [in panels (a), (c), and (e)] and $\bar{B}^0(B^0)$ [in panels (b), (d), and (f)] mesons are plotted as functions of the Borel mass M for $T = 50, 100$ and 150 MeV in nuclear matter ($f_s = 0$) for baryon density $\rho_B = 2\rho_0$ and $\eta = 0.5$, as the solid (dashed) lines, and are compared to the case of $\eta = 0$ (shown as dotted lines). The Borel curve for $\rho_B = 0$ is shown as the dot-dashed line.

of strangeness in the medium (with $f_s = 0.3$), as plotted in Fig. 4, for isospin symmetric matter ($\eta = 0$), the mass difference of $D^0(D^+)$ (~ 1914.4 MeV) and $\bar{D}^0(D^-)$ (~ 1935.9 MeV) is observed to be modified to 21.5 MeV. In the presence of isospin asymmetry, the mass splitting between the charged (neutral) particle and antiparticle, $D^+(D^0)$ and $D^-(\bar{D}^0)$, is observed to increase (drop) as the density is raised. However, the effect from the isospin asymmetry is lessened in the presence of strangeness in the medium, as can be seen from Figs. 3 and 4.

In Fig. 5, the masses of the strange-charm mesons, $D_s^\pm$ are plotted as functions of the Borel mass, M for the nuclear and strange hadronic matter for $\rho_B = 2\rho_0$. The Borel curves for these masses are shown for the vacuum (shown as a dot-dashed line). In nuclear matter, unlike the observation of the mass of the particle, $D^+(D^0)$ to be smaller than the mass of the antiparticle, $D^-(\bar{D}^0)$ (as seen in Figs. 1 and 3), the opposite trend of the mass of particle (D_s^+) to be higher than the mass of antiparticle (D_s^-) is observed in nuclear matter ($f_s = 0$). On the other hand, the nonzero value of $\langle s^\dagger s \rangle$ in the strange hadronic matter leads to the antiparticle (D_s^-) mass to be higher than that of D_s^+ in hyperonic medium.

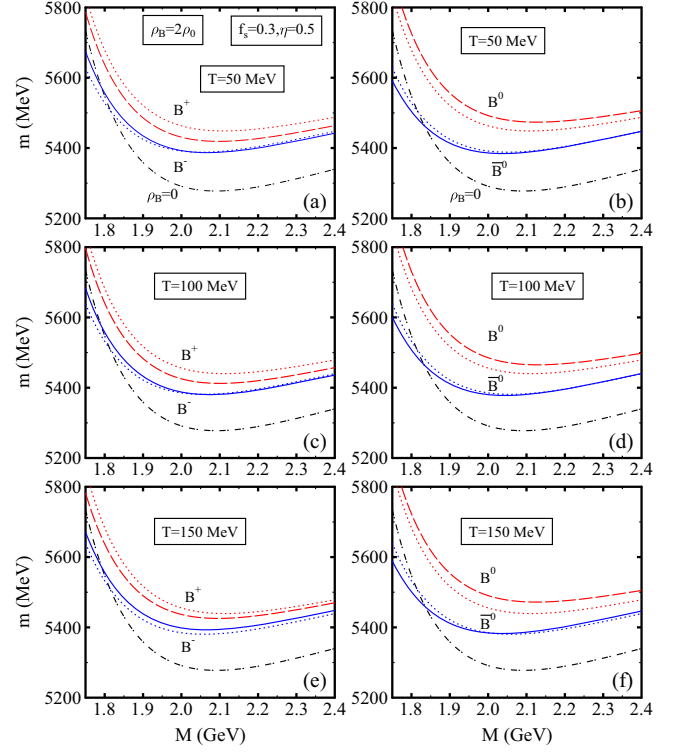


FIG. 8. Masses of $B^-(B^+)$ (in panels (a), (c) and (e)) and $\bar{B}^0(B^0)$ (in panels (b), (d) and (f)) mesons are plotted as functions of the Borel mass M for $T = 50, 100$ and 150 MeV in hyperonic matter (with $f_s = 0.3$) for baryon density $\rho_B = 2\rho_0$ and $\eta = 0.5$, as the solid (dashed) lines, and, are compared to the case of $\eta = 0$ (shown as dotted lines). The Borel curve for $\rho_B = 0$ is shown as the dot-dashed line.

Note that $\langle q_i^\dagger q_i \rangle$ for the $D(\bar{D})$ mesons corresponds to the number density of the nonstrange quarks (u or d), whereas, for $D_s^\pm$, $\langle q_i^\dagger q_i \rangle \equiv \langle s^\dagger s \rangle$, the number density of the s quark, which is 0 for nuclear matter. The mass splitting between the particle and antiparticle is thus mainly driven by the term $\langle q_i^\dagger q_i \rangle$, which dominates over the contributions from the dimension 5 operators $\langle q_i^\dagger i D_0^2 q_i \rangle$ and $\langle q_i^\dagger g_s \sigma \cdot G q_i \rangle$. In Fig. 6, the mass shifts of the strange-charm mesons are plotted as functions of the baryon density in units of nuclear matter saturation density. With isospin asymmetry, the mass shifts of both $D_s^\pm$ mesons are observed to be smaller as compared to the isospin symmetric case for nuclear matter. The trend is opposite in the presence of strangeness in the medium. The effect due to isospin asymmetry is observed to be marginal in strange hadronic matter for $T = 50$ and 100 MeV, whereas, at the higher temperature, $T = 150$ MeV, the effect is much larger.

The masses of the B and $\bar{B}$ mesons are plotted as functions of the Borel mass in Figs. 7 and 8 for $\rho_B = 2\rho_0$ for nuclear matter ($f_s = 0$) and strange hadronic matter (with $f_s = 0.3$), respectively, for values of the temperature, $T = 50, 100$, and 150 MeV. The density dependence of the mass shifts of B and $\bar{B}$ mesons are shown in Figs. 9 and 10

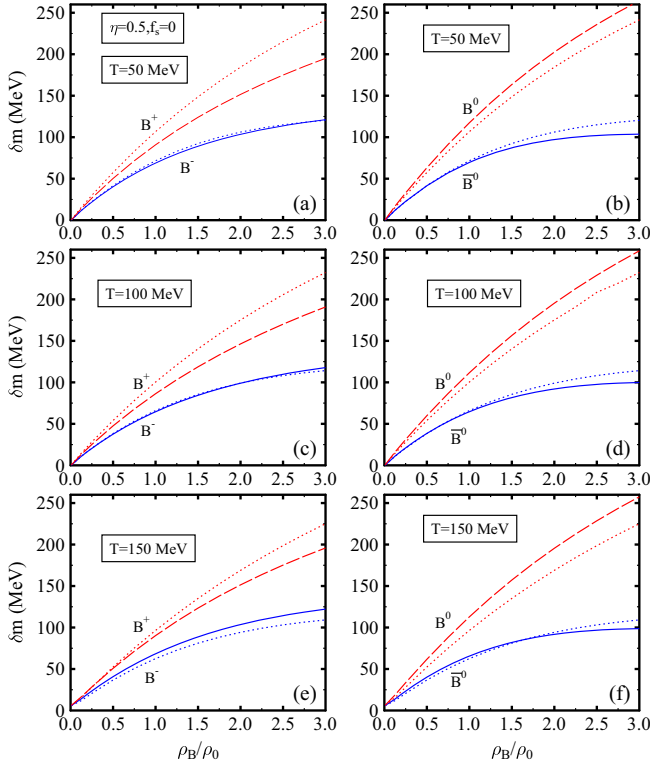


FIG. 9. Mass shifts of B^- (B^+) [in panels (a), (c), and (e)] and $\bar{B}^0$ (B^0) [in panels (b), (d), and (f)] mesons are plotted as functions of the baryon density in units of nuclear matter saturation density for $T = 50, 100$, and 150 MeV in nuclear matter ($f_s = 0$) for $\eta = 0.5$, as the solid (dashed) lines, and, are compared to the case of $\eta = 0$ (shown as dotted lines).

for $f_s = 0$ and $f_s = 0.3$, respectively. Similar to the cases of D and $\bar{D}$ mesons (see Figs. 1–4), the masses of the antiparticles are observed to be larger than the masses of the particles in nuclear matter as well as in strange hadronic matter. For isospin symmetric nuclear (hyperonic) matter, the masses of the mesons within the $B \equiv (B^+, B^0)$ and $\bar{B} \equiv (B^-, \bar{B}^0)$ doublets are degenerate. The masses of B ($\bar{B}$) for $T = 50, 100$, and 150 MeV are obtained as 5384.5 (5349.44), 5378.12 (5343.77), and 5374.27 (5340.35) for $\rho_B = \rho_0$ and 5462.26 (5384), 5453 (5376.87), and 5446.69 (5372) for $\rho_B = 2\rho_0$ in symmetric nuclear matter. For $\rho_B = \rho_0$, the mass shift of B ($\bar{B}$) in isospin symmetric nuclear matter from the vacuum value of 5277.8 MeV is thus observed to be 106.7 (71.64), 100.32 (66), and 96.47 (62.55), corresponding to $B - \bar{B}$ mass splittings as 35.06, 34.32 and 33.92, respectively. These values are similar to the value of the $D - \bar{D}$ mass splitting of around 32 MeV in isospin symmetric nuclear matter at $\rho_B = \rho_0$. The mass shifts (in MeV) from the vacuum values for the B and $\bar{B}$ mesons (~ 107 and 72) are observed to be much larger than the values of around 75 and 43 for the $\bar{D}$ and D mesons in isospin symmetric nuclear matter. The effect of isospin

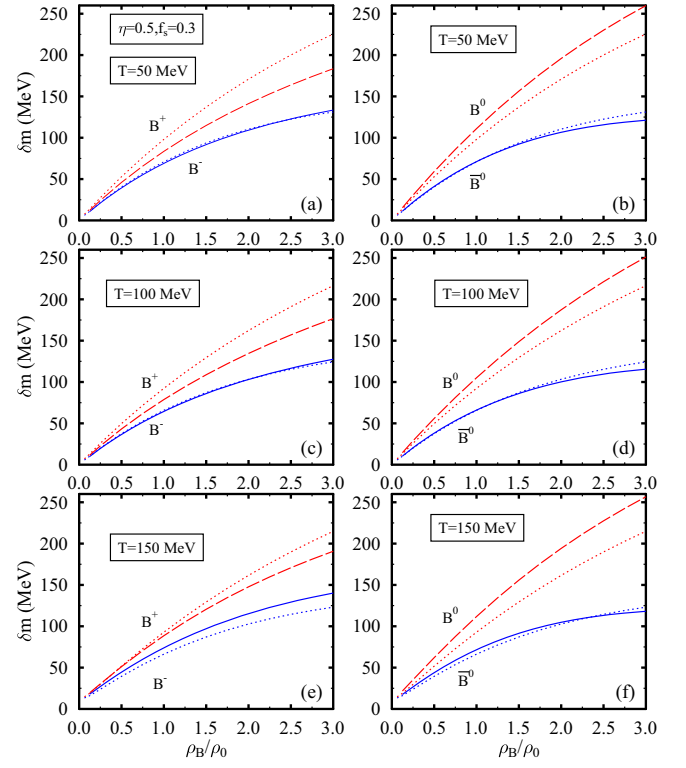


FIG. 10. Mass shifts of B^- (B^+) (in panels (a), (c) and (e)) and $\bar{B}^0$ (B^0) (in panels (b), (d) and (f)) mesons are plotted as functions of the baryon density in units of nuclear matter saturation density for $T = 50, 100$ and 150 MeV in hyperonic matter (with $f_s = 0.3$) for $\eta = 0.5$, as the solid (dashed) lines, and, are compared to the case of $\eta = 0$ (shown as dotted lines).

asymmetry is observed to be larger for the antiparticles (B^+ and B^0) as compared to the particles (B^- and $\bar{B}^0$). There is observed to be a drop (rise) in the mass of B^+ (B^0) within the isospin doublet (B^+, B^0) and in the mass of $\bar{B}^0$ (B^-) within the isospin doublet ($\bar{B}^0, B^-$), for isospin asymmetric (with $\eta = 0.5$) nuclear and hyperonic matter as compared to the isospin symmetric cases. However, the isospin asymmetry effect is observed to be marginal for the mass of $\bar{B}^0$.

In Fig. 11, the Borel curves for the strange-bottom mesons (B_s^0 and $\bar{B}_s^0$) are shown for density $\rho_B = 2\rho_0$ and Fig. 12 shows the density dependence of the mass shifts of these mesons for the nuclear and the strange hadronic matter. The mass of the antiparticle ($\bar{B}_s^0$) is observed to have lower mass as compared to the mass of the particle (B_s^0) in nuclear matter ($f_s = 0$) for which $\langle q_i^\dagger q_i \rangle \equiv \langle s^\dagger s \rangle$ is zero, whereas, the opposite trend is observed when there is nonzero strangeness in the medium. The lower (higher) mass of the particle, $\bar{B}_s^0$ as compared to the antiparticle, B_s^0 in the nuclear (hyperonic) matter shows that the term $\langle s^\dagger s \rangle$ dominates over the other two terms in the odd part of the spectral function, which is responsible

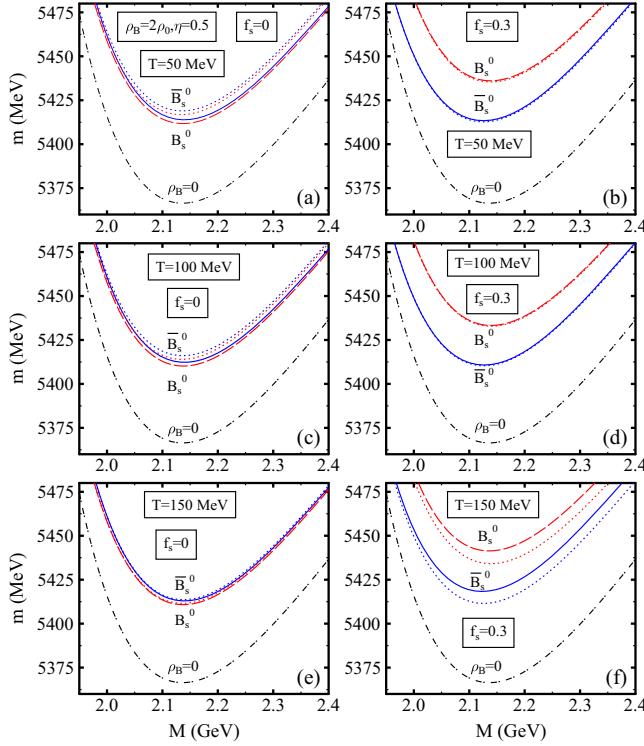


FIG. 11. Masses of $\bar{B}_s^0$ and B_s^0 mesons are plotted as function of the Borel mass M for $T = 50, 100$, and 150 MeV for nuclear matter ($f_s = 0$) in panels (a), (c), and (e) and for hyperonic (with $f_s = 0.3$) in panels (b), (d), and (f), for baryon density $\rho_B = 2\rho_0$ and $\eta = 0.5$ as the solid (dashed) line, and compared to the case of $\eta = 0$ (shown as dotted line). The Borel curve for $\rho_B = 0$ is shown as the dot-dashed line.

for the splitting of the masses between the particle and its antiparticle. The behavior of the mass splitting at finite densities is similar to the case of strange-charm mesons (as shown in Figs. 5 and 6). Similar to the strange-charm sector, where the isospin asymmetry leads to a drop (rise) of the masses of both D_s^+ and D_s^- mesons in nuclear (hyperonic) matter, the masses of both $\bar{B}_s^0$ and B_s^0 mesons are also observed to have a downward (upward) shift in nuclear (hyperonic) medium (see Figs. 11 and 12). However, the mass modifications due to isospin asymmetry are observed to be marginal for $T = 50$ and 100 MeV in hyperonic matter.

For the sake of comparison, we show the density dependence of the mass shifts of the mesons studied in the present work, in the hot nuclear and hyperonic matter in Figs. 13 and 14 for the charm sector and in Figs. 15 and 16 for the bottom sector. The masses are plotted for the isospin symmetric ($\eta = 0$) as well as for asymmetric matter (with $\eta = 0.5$) for values of the temperature, $T = 50, 100$, and 150 MeV. The mass shifts of the open heavy flavor mesons with the nonstrange quark (antiquark) constituents are observed to be larger than the strange-charm (strange-bottom) mesons. As has already been mentioned, the

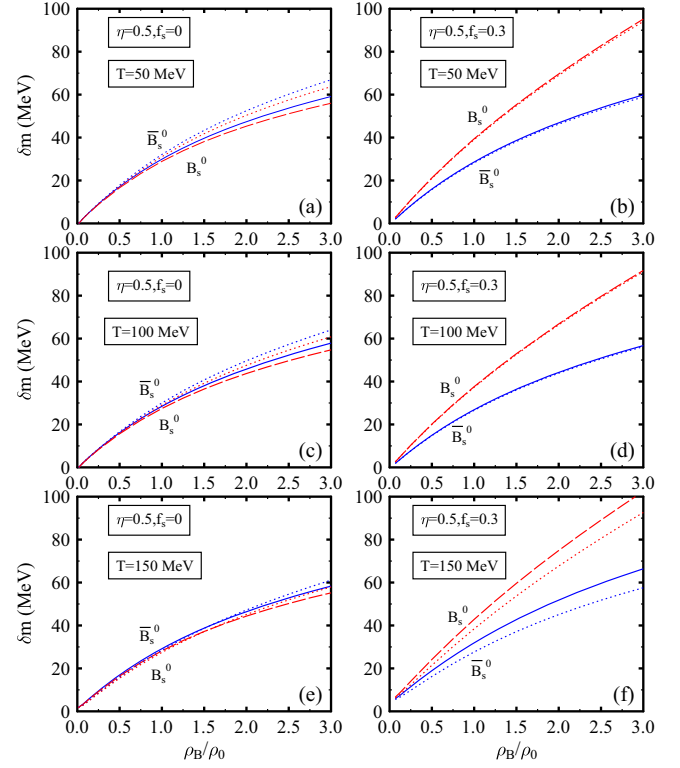


FIG. 12. Mass shifts of $\bar{B}_s^0$ and B_s^0 mesons are plotted as function of the baryon density in units of nuclear matter saturation density for $\eta = 0.5$ [in panels (a), (c), and (e) for nuclear matter ($f_s = 0$) and in panels (b), (d), and (f) for strange hadronic matter with $f_s = 0.3$] for $T = 50, 100$, and 150 MeV as the solid (dashed) line, and compared to the case of $\eta = 0$ (shown as dotted line).

antiparticles (D^- and $\bar{D}^0$ for the charm sector and B^+ and B^0 for the bottom sector) have larger mass shifts, as compared to the particles (D^+ and D^0 for the charm sector and B^- and $\bar{B}^0$ for the bottom sector), both in nuclear and hyperonic matter. The same trend of the antiparticle having a larger mass shift as compared to the particle is observed in hyperonic matter, for both strange-charm and strange-bottom mesons, as might be observed in Figs. 14 and 16, whereas the opposite trend is observed when the strangeness fraction in the medium is zero, i.e., in nuclear medium (see Figs. 13 and 15). This opposite behavior of the strange-heavy flavor mesons in nuclear and hyperonic media, as has already been seen from Figs. 5, 6, 11, and 12, arises due to the term proportional to the strange quark number density in the odd part of the spectral function for hyperonic matter, which is zero in nuclear matter. The isospin asymmetry in the medium leads to a splitting in the masses between the mesons of the iso-doublets [(D^0, D^+) , $(\bar{D}^0, D^-)$ for the open charm mesons and $(B^-, \bar{B}^0)$, (B^+, B^0) for the open bottom mesons]. As may be observed from Figs. 13 and 14, due to isospin asymmetry in the medium (shown for isospin asymmetry parameter, $\eta = 0.5$), there is an upward (downward) shift in the mass

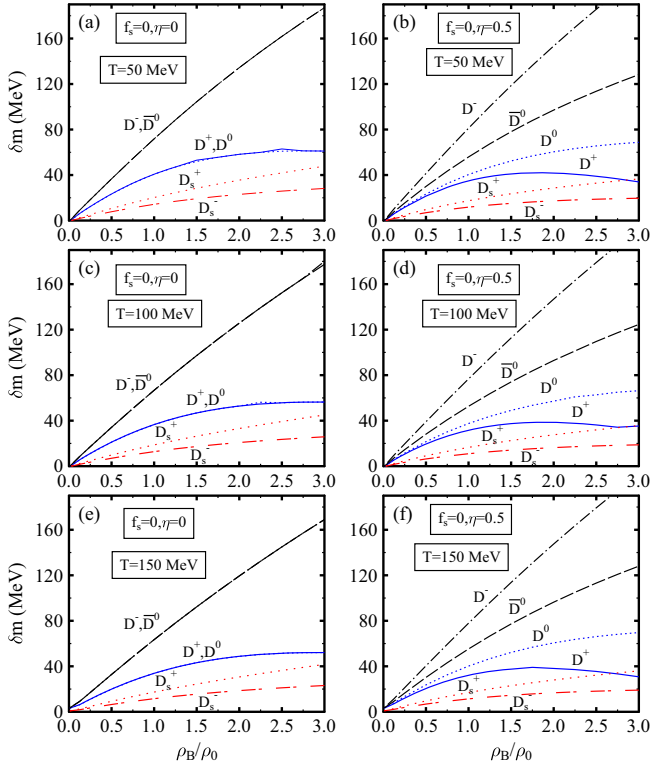


FIG. 13. Mass shifts of the open charm mesons are plotted as a functions of baryon density in units of nuclear matter saturation density for the isospin symmetric ($\eta=0$) nuclear matter ($f_s=0$) in panels (a), (c), and (e), and for the isospin asymmetric (with $\eta=0.5$) nuclear matter ($f_s=0$) in panels (b), (d), and (f), for $T=50, 100$, and 150 MeV, respectively. These mass shifts are shown as the short-dash-dotted, dashed, closely spaced dotted, solid, widely spaced dotted, and long-dash-dotted for the D^- , $\bar{D}^0$, D^0 , D^+ , D_s^+ and D_s^- mesons, respectively.

of the D^- ($\bar{D}^0$) for the antiparticle doublet, whereas there is a drop (rise) in the mass of D^+ (D^0) for the particle isodoublet, in both nuclear and hyperonic matter. The isospin asymmetry is observed to lead to a much larger particle–antiparticle mass difference for the charged sector (D^+ and D^-), whereas, there is observed to be a reduction in the particle–antiparticle mass difference for the neutral (D^0 and $\bar{D}^0$) mesons. Hence, the effects of the isospin asymmetry might modify the particle ratios, D^+/D^- and $D^0/\bar{D}^0$, as well as the dilepton yield from the $D\bar{D}$ in asymmetric heavy ion collision experiments. It might be noted here that the results discussed as well as shown in the figures in the present work correspond to the maximum possible isospin asymmetry ($\eta=0.5$), which is, however, not realized in the heavy ion collision experiments. For a smaller value of isospin asymmetry parameter, $\eta=0.3$ and $T=100$ MeV, one observes the particle–antiparticle mass difference (in MeV) to be modified to around 39 and 157 for densities ρ_0 and $3\rho_0$, respectively, for the charged open charm mesons, and 21 and 83 for the D^0 and $\bar{D}^0$ mesons, from the values of 30 and 121 for the charged as well as neutral open charm

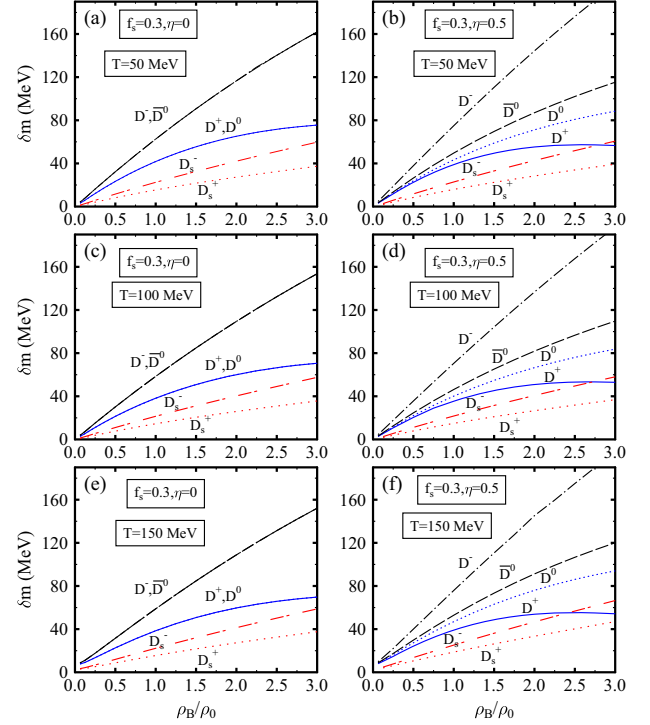


FIG. 14. Mass shifts of the open charm mesons are plotted as a functions of baryon density in units of nuclear matter saturation density for the isospin symmetric ($\eta=0$) hyperonic matter (with $f_s=0.3$) in panels (a), (c) and (e), and, for the isospin asymmetric (with $\eta=0.5$) hyperonic matter (with $f_s=0.3$) in panels (b), (d) and (f), for $T=50, 100$ and 150 MeV respectively. These mass shifts are shown as the short-dash-dotted, dashed, closely spaced dotted, solid, widely spaced dotted, long-dash-dotted for the D^- , $\bar{D}^0$, D^0 , D^+ , D_s^+ and D_s^- mesons respectively.

mesons in isospin symmetric ($\eta=0$) nuclear matter. The effect due to isospin asymmetry is thus observed to be appreciable, especially at high densities.

The effect of the baryon density is observed to be the dominant medium effect for the modifications of the masses of the open charm and open bottom mesons considered in the present work. There are significant contributions due to the isospin asymmetry of the medium, leading to appreciable increase (decrease) in the particle–antiparticle mass difference in the charged (neutral) mesons at high densities, for both charm (D , $\bar{D}$) and bottom (B , $\bar{B}$) sectors. With accessible energy for study of charm mesons at the CBM Experiment at FAIR at the future facility of GSI [91,92], which will be designed to investigate matter at high densities and moderate temperatures, the medium modifications of the masses of the D and $\bar{D}$ mesons could have observable consequences in the particle ratios, e.g., D^+/D^- , $D^0/\bar{D}^0$ in asymmetric nuclear collisions at CBM, FAIR. The medium modifications of the D and $\bar{D}$ mesons can modify the spectral properties of the charmonium states, which can subsequently affect the production of the open and hidden charmed mesons. Also, the dilepton

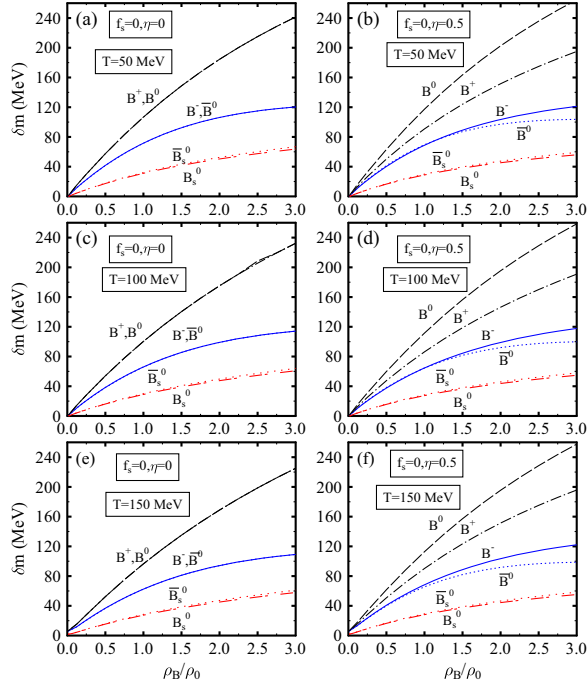


FIG. 15. Mass shifts of the open bottom mesons are plotted as a functions of baryon density in units of nuclear matter saturation density for the isospin symmetric ($\eta = 0$) nuclear matter ($f_s = 0$) in panels (a), (c), and (e), and for the isospin asymmetric (with $\eta = 0.5$) nuclear matter ($f_s = 0.5$) in panels (b), (d), and (f), for $T = 50, 100$, and 150 MeV, respectively. These mass shifts are shown as the short-dash-dotted, dashed, closely spaced dotted, solid, widely spaced dotted, and long-dash-dotted for the B^0 , B^+ , $\bar{B}^0$, B^- , $\bar{B}_s^0$, and B_s^0 mesons, respectively.

yield arising semileptonically from the $D\bar{D}$ correlated pairs can be modified in the isospin asymmetric hadronic matter. The measurement of the invariant dilepton spectra by an electron pair spectrometer at the J-PARC 50-GeV proton spectrometer is being carried out to explore the chiral symmetry of QCD at around nuclear density through study of the medium modifications of the vector mesons produced in pA collisions and decaying to e^+e^- pairs [93], with the E16 experiment at J-PARC specifically designed to study the mass of the ϕ mesons, due to its narrow width. The K^* and K_1 mesons can be produced at the J-PARC facility using the kaon beam, and partial chiral symmetry restoration in the hadronic medium can lead to reduction in the mass difference of these chiral partners. However, due to hadronic rescattering, there is expected to be enhancement in the production ratio of K^+/K^* , which can be a probe for chiral symmetry restoration [94]. In the future, with upgrade of the spectrometer at the J-PARC facility, the measurement of charm hadrons could be possible. The measurements of various hadrons can give a better understanding of the medium modifications of the hadron masses as well as a better insight to chiral symmetry restoration at high densities [95].

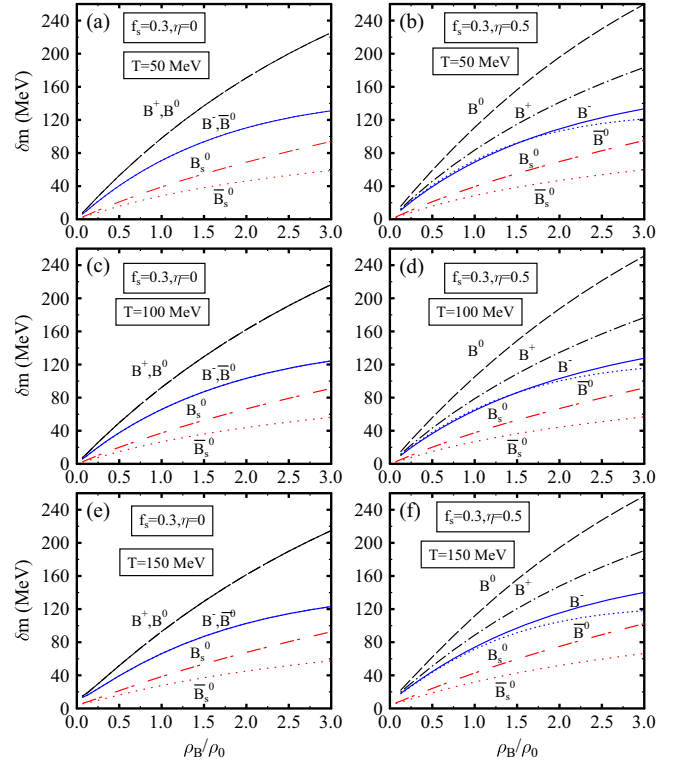


FIG. 16. Mass shifts of the open bottom mesons are plotted as a functions of baryon density in units of nuclear matter saturation density for the isospin symmetric ($\eta=0$) hyperonic matter (with $f_s=0.3$) in panels (a), (c) and (e), and, for the isospin asymmetric (with $\eta=0.5$) hyperonic matter (with $f_s=0.5$) in panels (b), (d) and (f), for $T = 50, 100$ and 150 MeV respectively. These mass shifts are shown as the short-dash-dotted, dashed, closely spaced dotted, solid, widely spaced dotted, long-dash-dotted for the B^0 , B^+ , $\bar{B}^0$, B^- , $\bar{B}_s^0$ and B_s^0 mesons respectively.

V. SUMMARY

In the present work, we have studied the medium modifications of the masses of the open charm (D , $\bar{D}$, D_s , and $\bar{D}_s$) and open bottom (B , $\bar{B}$, B_s , and $\bar{B}_s$) in isospin asymmetric strange hadronic matter at finite temperature. These are computed using a QCD sum rule approach retaining the QCD operators up to dimension 5 in the operator product expansion. The values of the QCD operators are calculated within a chiral SU(3) model. The mass splitting of the particles (D^0 , D^+ , D_s^+ , B^- , $\bar{B}^0$, $\bar{B}_s^0$) and their antiparticles ($\bar{D}^0$, D^- , D_s^- , B^+ , B^0 , B_s^0) arises due to the odd part of the spectral function. In nuclear matter, the masses of the particles $D^+(D^0)$ and $B^-(\bar{B}^0)$ mesons are observed to be smaller than the masses of the antiparticles $D^-(\bar{D}^0)$ and $B^+(B^0)$. For the strange-charm as well as strange-bottom mesons, the opposite behavior of the masses of particles (D_s^+ and $\bar{B}_s^0$) to be higher than the masses of antiparticles (D_s^- and B_s^0) is observed in nuclear matter, whereas, in the strange hadronic matter, the nonzero value of $\langle q_i^\dagger q_i \rangle \equiv \langle s^\dagger s \rangle$ leads the masses of the antiparticles

(D_s^- and B_s^0) to be higher than that of D_s^+ and $\bar{B}_s^0$ in hyperonic medium. The sign of the difference in the masses of the particle and antiparticle ($m_+ - m_-$) is observed to be opposite in nuclear and hyperonic matter, both for the strange-charm as well as strange-bottom mesons. In the mass splitting between the particle and antiparticle arising due to the odd part of the spectral function, the term proportional to $\langle q_i^\dagger q_i \rangle \equiv \langle s^\dagger s \rangle$ for the strange-charm (strange-bottom) mesons thus dominates over the other QCD operators $\langle q_i^\dagger i D_0^2 q_i \rangle$ and $\langle q_i^\dagger g_s \sigma \cdot G q_i \rangle$. The effects from density and isospin asymmetry, particularly, on the

masses of the D , $\bar{D}$, B , and $\bar{B}$ mesons are observed to be appreciable. These can modify, e.g., the production of the charmonia and open charm mesons, and the production ratios D^+/D^- , $D^0/\bar{D}^0$ in asymmetric relativistic heavy ion collision experiments at the CBM Experiment at the future facility at GSI.

DATA AVAILABILITY

The data created or analyzed in this study are not publicly available. The data are available from the authors upon reasonable request.

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