Open bottom mesons in asymmetric nuclear matter in presence of strong magnetic fields

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Abstract

The modifications of the masses of the $B$ and $\bar{B}$ mesons in asymmetric nuclear matter in the presence of strong magnetic fields, are investigated using a chiral effective model. The medium modifications of these open bottom mesons arise due to their interactions with the scalar mesons and the nucleons. In the magnetized nuclear matter, the proton has contributions from the Landau levels. In the chiral effective model, the masses of the $B$ and $\bar{B}$ mesons are calculated from the leading term, namely the vectorial Weinberg Tomozawa term as well as from the next to leading order contributions, i.e., due to the scalar exchange and the range terms. Due to the Weinberg-Tomozawa term, the $\bar{B}$ mesons experience an attractive interaction in the symmetric nuclear matter, whereas the $B$ mesons have a repulsive interaction. Inclusion of the contributions from the scalar exchange and the range terms as well, leads to drop of the masses of both $B$ and $\bar{B}$ mesons. The effect of the isospin asymmetry breaks the mass degeneracy of the $B^+$ and $B^0$ (as well as of the $B^-$ and $\bar{B}^0$) mesons, and its effect is observed to be large at high densities. The effects of anomalous magnetic moments of the nucleons are taken into account in the present study of the masses of the open bottom mesons in magnetized nuclear matter.

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I. INTRODUCTION

The study of properties of hadrons in the hadronic medium is an intense field of research in high energy physics. This is due to the relevance of the subject in ultra-relativistic heavy-ion collision experiments, where the experimental observables are affected by the medium modifications of the hadrons. At high energy heavy ion collision experiments, huge magnetic fields are created \( eB \sim 6m^2 \) at RHIC, BNL, and \( eB \sim 15m^2 \) at LHC, CERN. It is thus important to study the effects of the magnetic fields on the hadrons in the strongly interacting matter resulting from the high energy nuclear collisions. The study of magnetic fields is also important for astrophysical objects like the neutron stars with magnetic fields/magnetars. The presence of magnetic fields also lead to novel effects like chiral magnetic effect and the inverse magnetic catalysis, which has made the study of effects of the magnetic fields in strongly interacting matter even more interesting. The magnetic fields created in non-central relativistic heavy ion collision experiments drop rapidly after the collision, creating induced currents, which tend to slow down the decay of the magnetic field. Also at later times, the matter effects are believed to play an important role in the time evolution of the magnetic field, with the effect of slowing down the decay of the magnetic field. The time evolution of the magnetic field in the heavy ion collision experiments is still an open problem, which needs careful investigation of the electrical conductivity of the medium as well as solutions of the magnetohydrodynamic equations. The fact that there are huge magnetic fields created in relativistic heavy ion collision experiments, has initiated number of studies of the hadrons, including those of heavy flavour mesons, e.g., open charm and open bottom mesons, as well as heavy quarkonium states, e.g, the charmonium states, in the presence of magnetic fields.

The masses of the heavy (charm and bottom)-light mesons, e.g. the \( D \) and \( B \) mesons, are modified appreciably in the hadronic medium. This is because, due to the presence of the light quark (antiquark) (along with the heavy charm (bottom) antiquark (quark)), these mesons interact with the light quark condensates, which are modified significantly in the medium. The effects of these medium modifications should show in observables, e.g. the production and propagation of these particles. There have been various approaches to study the in-medium properties of these heavy flavour mesons, e.g., QCD sum rule approach.
(where the in-medium masses are calculated from the medium modifications of the quark and gluon condensates), the Quark Meson Coupling (QMC) model \[25\], in which the quarks interact via scalar and vector mesons \[26\], the coupled channel approach \[27–30\], using effective hadronic models, e.g. chiral effective model \[31–34\], studies using heavy quark symmetry and $\bar{D}$ ($B$) interaction with nucleon using pion exchange \[35\], and, using a heavy meson effective theory with $O(1/M)$ corrections \[36\]. There has also been a study of the heavy flavour meson (heavy quark) embedded as an impurity in the nuclear matter (quark matter) by Yasui and Sudoh, which have very similar behaviour at large mass (of the meson/quark) limit \[37\]. The open charm and open bottom mesons studied using pion exchange for the heavy meson–nucleon interaction \[35\], give attractive interactions of the $\bar{D}$ and $B$ mesons in nuclear matter, which predict possibility of the bound states of these mesons with the nuclei \[25\]. Due to attractive interaction of $J/\psi$ in nuclear matter \[33, 34, 38, 39\], the bound states of the $J/\psi$ to nuclei have also been predicted within the QMC model \[40\].

In the present work, the mass modifications of the $B$ and $\bar{B}$ mesons are studied in symmetric as well as asymmetric nuclear matter in the presence of strong magnetic fields, using a chiral effective model. The model is a generalization of a chiral SU(3) model to include the interactions of the charm and bottom mesons with the light hadronic sector. The chiral SU(3) model, based on a nonlinear realization of chiral symmetry, has been used to study nuclear matter, finite nuclei \[49\], hyperonic matter \[50\], vector mesons \[51\], kaons and antikaons \[52–55\], as well as to study the charge neutral matter present in the interior of the (proto) neutron stars \[56\]. The chiral SU(3) model, generalized to the charm and bottom sectors, and has been used to study the $D$ and $\bar{D}$ mesons \[32–34\], the strange charm mesons (the $D_s$ mesons) \[57\], the $B$, $\bar{B}$ mesons \[58\], and, the strange bottom mesons, i.e.,
the $B_s$ mesons \cite{59}, as arising from their interactions to the baryons and the scalar mesons. Within effective hadronic model, the broken scale symmetry of QCD \cite{49,50,60} has also been incorporated through a scalar dilaton field which mimicks the gluon condensates of QCD. The mass modifications of the heavy quarkonium mesons, e.g., charmonium \cite{33,34} and bottomonium states \cite{61}, which do not contain any light quark/antiquark, arise due to their interactions with the gluon condensate of QCD. These are calculated within the chiral effective model from the medium modifications of the dilaton field in the nuclear matter. The in-medium partial decay widths of the charmonium states ($J/\psi$ and the excited states of charmonium, e.g., $\psi(3686)$, $\psi(3770)$, as well as $\chi_c$) to $D\bar{D}$ have been studied in the literature, using a light quark pair creation model, namely $^3P_0$ model \cite{63,64}. These have been calculated by assuming mass modifications of the $D$ and $\bar{D}$ mesons \cite{62}, but without accounting for any mass modifications of the charmonium states. Using the $^3P_0$ model, from the medium modifications of the masses of the open charm mesons and the charmonium states, calculated within the chiral effective model, the partial decay widths of the charmonium states to $D\bar{D}$ in hadronic matter, have also been studied \cite{34}. Subsequently, the in-medium charmonium decay widths have been studied, using a field theoretic model of composite hadrons \cite{65}. The partial decay widths of the $\Upsilon$-states to $B\bar{B}$, in hadronic matter have later been studied using the field theoretic model of composite hadrons with quark constituents \cite{66}, using the in-medium masses of the $B$ and $\bar{B}$ \cite{58}, as well as the bottomonium states \cite{61}, as calculated within the chiral effective model. In the mean field approximation, the values of the scalar mesons, $\sigma$, $\zeta$ and $\delta$, as calculated from their equations of motion, are related to the light quark condensates, and the dilaton field is related to the gluon condensate in the medium, which have been used to study the properties of the light vector mesons ($\omega$, $\rho$ and $\phi$) \cite{67}, as well as the charmonium states, $J/\psi$ and $\eta_c$ \cite{68} using the QCD sum rule approach. The $D$ mesons in the presence of an external magnetic field have been studied in the literature, within the QCD sum rule approach, accounting for the mixing of the pseudoscalar mesons and the vector mesons, as well as the Landau quantization effects for the charged $D$ mesons \cite{18}. Using a semiclassical approach \cite{16}, the open charm and open bottom mesons have been studied in the presence of an external magnetic field, where the mass of the open charm (bottom) meson is due to the interaction of the magnetic field to the
spin of the quarks. Different spin orientations of the heavy-light quark-antiquark systems were observed to have different masses in the presence of the magnetic field. The $D$ and $B$ mesons were observed to have lowering of their masses (arising from the mass reductions of specific spin orientations) [16]. Further, M1 transitions become dominant in the presence of strong magnetic fields, leading to the mixing of the spin–0 states $D$ ($B$) to the spin–1 states $D^*$ ($B^*$). The masses of the $D$ and $\bar{D}$ mesons in nuclear matter in the presence of strong magnetic fields have been recently investigated [19]. These masses are calculated as arising due to their interactions with the nucleons and the scalar mesons in the magnetized nuclear matter. The proton has contributions from the Landau energy levels, in the presence of the external magnetic field. The effects of isospin asymmetry as well as anomalous magnetic moments of the nucleons [5–8, 69–72] are considered for the study of the masses of $D$ mesons in the magnetized nuclear matter. In the present work, we study the medium modifications of the masses of the $B$ and $\bar{B}$ mesons in nuclear matter in the presence of strong magnetic fields within the chiral effective model, including the effects from isospin asymmetry as well as anomalous magnetic moments of the proton and neutron.

We organize the paper as follows. In section II, we describe the chiral effective model used for the study of the modifications of the masses $B$ and $\bar{B}$ mesons in the strongly magnetized (asymmetric) nuclear matter in the present investigation. In Section III, we discuss the results of the effects of the magnetic field (including the effects of the anomalous magnetic moments of the nucleons) on the $B$ and $\bar{B}$ masses in the isospin asymmetric magnetized nuclear medium. Section IV summarizes the findings of the present work.

II. $B$ AND $\bar{B}$ MESONS IN MAGNETIZED NUCLEAR MATTER

In the present work, we use the effective chiral model for the study of the medium modifications of $B$ and $\bar{B}$ mesons in nuclear matter in the presence of magnetic field. The Lagrangian density for the model is given as [19]

$$\mathcal{L} = \mathcal{L}_{\text{kin}} + \sum_{W=X,Y,V,A,u} \mathcal{L}_{BW} + \mathcal{L}_{\text{vec}} + \mathcal{L}_0 + \mathcal{L}_{\text{scalebreak}} + \mathcal{L}_{\text{SB}} + \mathcal{L}_{\text{mag}},$$  

(1)

which is based on a nonlinear realization of SU(3) chiral symmetry [73–75] and broken scale invariance [49, 50, 60]. In Eq. (1), $\mathcal{L}_{\text{kin}}$ is the kinetic energy term, $\mathcal{L}_{BW}$ contains the baryon-
meson interactions. The baryon masses are generated by the baryon-scalar meson interaction terms in the Lagrangian density and the parameters corresponding to these interactions are adjusted so as to obtain the baryon masses as their experimentally measured vacuum values. $\mathcal{L}_{vec}$ describes the dynamical mass generation of the vector mesons via couplings to the scalar fields and contains additionally quartic self-interactions of the vector fields. $\mathcal{L}_0$ contains the meson-meson interaction terms inducing the spontaneous breaking of chiral symmetry, $\mathcal{L}_{scalebreak}$ is a scale invariance breaking logarithmic potential, $\mathcal{L}_{SB}$ describes the explicit chiral symmetry breaking. The last term, $\mathcal{L}_{mag}$ is the contribution from the magnetic field, given as

$$\mathcal{L}_{mag} = -\bar{\psi}_i q_i \gamma\mu A^\mu \psi_i - \frac{1}{4} k_i \mu_N \bar{\psi}_i \sigma^{\mu\nu} F_{\mu\nu} \psi_i - \frac{1}{4} F^{\mu\nu} F_{\mu\nu}. \quad (2)$$

In the above, $\psi_i$ corresponds to the $i$-th baryon. The second term in equation (2) corresponds to the tensorial interaction with the electromagnetic field and is related to the anomalous magnetic moment of the baryon (proton and neutron for nuclear matter as considered in the present investigation). In this term, $\mu_N$ is the Nuclear Bohr magneton, given as $\mu_N = \frac{e}{2m_N}$, where $m_N$ is the vacuum mass of the nucleon. $F^{\mu\nu}$ is the electromagnetic tensor given as $F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu$. We choose the magnetic field to be uniform and along the z-axis, i.e., $\vec{B} = (0, 0, B)$, and take the vector potential to be $A^\mu = (0, 0, Bx, 0)$.

For the study of $B$ and $\bar{B}$ mesons in the nuclear matter in the presence of magnetic fields, we write the Lagrangian density in the mean field approximation (where the meson fields are treated as classical fields), and determine these fields from their coupled equations of motion [50, 51]. We use the frozen glueball approximation, i.e., neglect the medium dependence of the dilaton field, which is small compared to those of the scalar fields. The equations of motion of the scalar fields are given as

$$k_0 \chi^2 \sigma - 4k_1 \left(\sigma^2 + \zeta^2 + \delta^2\right) \sigma - 2k_2 \left(\sigma^3 + 3\sigma \delta^2\right) - 2k_3 \chi \sigma \zeta - \frac{d}{3} \chi^4 \left(\frac{2\sigma}{\sigma^2 - \delta^2}\right) + \left(\frac{\chi}{\chi_0}\right)^2 m_\pi^2 f_\pi - g_{\sigma N} (\rho^p_\pi + \rho^n_\pi) = 0, \quad (3)$$

$$k_0 \chi^2 \zeta - 4k_1 \left(\sigma^2 + \zeta^2 + \delta^2\right) \zeta - 4k_2 \zeta^3 - k_3 \chi \left(\sigma^2 - \delta^2\right) - \frac{d}{3} \chi^4 \left[\sqrt{2} m_k^2 f_k \chi \left(\frac{1}{\sqrt{2}} m_k^2 f_k\right) - g_{\zeta N} (\rho^p_\zeta + \rho^n_\zeta) = 0, \quad (4)$$
\[ k_0 \chi^2 \delta - 4k_1 \left( \sigma^2 + \zeta^2 + \delta^2 \right) \delta - 2k_2 \left( \delta^3 + 3\sigma^2 \delta \right) + k_3 \chi \delta^3 \]

\[ + \frac{2}{3} d \left( \frac{\delta}{\sigma^2 - \delta^2} \right) - g_{\delta N} (\rho_p^s - \rho_n^s) = 0, \]  

(5)  

where, \( \rho_p^s \) and \( \rho_n^s \) are the scalar densities for the proton and neutron respectively. In the magnetized nuclear matter, the expressions for these scalar densities are given as \([7, 8]\)

\[ \rho_p^s = \frac{e B m_p^*}{2\pi^2} \left[ \sum_{\nu=0}^{\nu_{\text{max}}} \frac{\nu m_p^{*2} + 2e B \nu + \Delta_p}{\sqrt{m_p^{*2} + 2e B \nu}} \ln \left| \frac{k_{f,\nu,1}^{(p)} + E_f^{(p)}}{\sqrt{m_p^{*2} + 2e B \nu + \Delta_p}} \right| \right. \]

\[ + \left. \sum_{\nu=1}^{\nu_{\text{max}}-1} \frac{\nu m_p^{*2} + 2e B \nu - \Delta_p}{\sqrt{m_p^{*2} + 2e B \nu}} \ln \left| \frac{k_{f,\nu-1}^{(p)} + E_f^{(p)}}{\sqrt{m_p^{*2} + 2e B \nu - \Delta_p}} \right| \right], \]  

(6)  

and

\[ \rho_n^s = \frac{m_n^*}{4\pi^2} \sum_{s=\pm 1} \left[ k_{f,s}^{(n)} E_f^{(n)} - (m_n^* + s\Delta_n)^2 \ln \left| \frac{k_{f,s}^{(n)} + E_f^{(n)}}{m_n^* + s\Delta_n} \right| \right]. \]  

(7)  

In the expression for the scalar density for the proton as given by equation (6), one sees that there are contributions from the Landau energy levels due to the presence of the external magnetic field, \( B \) along z-axis. \( k_{f,\nu,\pm 1}^{(p)} \) are the Fermi momenta of protons for the Landau level, \( \nu \) for the spin index, \( s = \pm 1 \) (for the spin up and down projections), which are related to the Fermi energy of the proton as

\[ k_{f,\nu,s}^{(p)} = \sqrt{E_f^{(p)}^2 - \left( \sqrt{m_p^{*2} + 2e B \nu + s\Delta_p} \right)^2}. \]  

(8)  

In the scalar density for the neutron given by equation (7), the Fermi momenta for the neutron, for spin projection, \( s = \pm 1 \) (corresponding to the spin up (down) projection of the neutron), are given as

\[ k_{f,s}^{(n)} = \sqrt{E_f^{(n)}^2 - (m_n^* + s\Delta_n)^2}, \]  

(9)

where \( E_f^{(n)} \) is the Fermi energy of the neutron. The effects of the anomalous magnetic moments for the proton and neutron are encoded in the parameter \( \Delta_{p(n)} \), which is given as

\[ \Delta_{p(n)} = -\frac{1}{2} \kappa_{p(n)} \mu_N B, \]  

where, \( \kappa_i \) \( (i = p, n \text{ for nuclear matter}) \) is as defined in the electromagnetic tensor term in the Lagrangian density given by (2). The values of \( \kappa_p \) and \( \kappa_n \) of 3.5856 and −3.8263 are the values of the gyromagnetic ratio arising from the anomalous magnetic moments of the proton and neutron respectively.

In the present work, we study the in-medium masses of the open bottom mesons in the presence of a magnetic field, for given values isospin asymmetry parameter, \( \eta = (\rho_n - \rho_p)^s \).
\( \rho_p/(2\rho_B) \), where \( \rho_B \) is the baryon density, and the number densities for the proton and the neutron are given as

\[
\rho_p = \frac{eB}{4\pi^2} \left[ \sum_{\nu=0}^{(s+1)} k_{f,\nu}^{(p)} + \sum_{\nu=1}^{(\nu_{\max})} k_{f,\nu-1}^{(p)} \right]
\]

and

\[
\rho_n = \frac{1}{4\pi^2} \sum_{s=\pm 1} \left\{ \frac{2}{3} k_{f,s}^{(n)}, \Delta_n \left[ (m_n^s + s\Delta_n)k_{f,s}^{(n)} + E_f^{(n)} (\arcsin\left( \frac{m_n^s + s\Delta_n}{E_f^{(n)}} \right) - \frac{\pi}{2}) \right] \right\}
\]

respectively. As has already been mentioned, in the present work, the in-medium masses of the \( B \) and \( \bar{B} \) in the magnetized nuclear matter are studied in a chiral effective model, which is a generalization of a chiral SU(3) model, to the charm and bottom sectors, so as to derive the interactions of the charm mesons \([31–34, 57]\), and bottom mesons \([58, 59]\) with the light hadronic sector. The interaction Lagrangian modifying the \( B \) and \( \bar{B} \) mesons can be written as \([58]\)

\[
\mathcal{L}_{BN} = -\frac{i}{8f_B^2} \left[ 3(\bar{p}\gamma^\mu p + \bar{n}\gamma^\mu n) \left\{ ((\partial_\mu B^+) B^- - B^+(\partial_\mu B^-)) + ((\partial_\mu B^0) \bar{B}^0 - B^0(\partial_\mu \bar{B}^0)) \right\} \right.
\]

\[
+ (\bar{p}\gamma^\mu p - \bar{n}\gamma^\mu n) \left\{ ((\partial_\mu B^0) B^- - B^0(\partial_\mu B^-)) - ((\partial_\mu B^0) \bar{B}^0 - B^0(\partial_\mu \bar{B}^0)) \right\} \]

\[
+ \frac{m_B^2}{2f_B} \left[ (\sigma + \sqrt{2}\zeta_b)((B^+ B^-) + B^0 \bar{B}^0) + \delta((B^+ B^-) - B^0 \bar{B}^0) \right]
\]

\[
- \frac{1}{f_B} \left[ (\sigma + \sqrt{2}\zeta_b)((\partial_\mu B^+)(\partial^\mu B^-) + (\partial_\mu B^0)(\partial^\mu \bar{B}^0)) \right.
\]

\[
+ \delta((\partial_\mu B^+)(\partial^\mu B^-) - (\partial_\mu B^0)(\partial^\mu \bar{B}^0)) \]

\[
+ \frac{d_1}{2f_B^2} (\bar{p}p + \bar{n}n)((\partial_\mu B^+)(\partial^\mu B^-) + (\partial_\mu B^0)(\partial^\mu \bar{B}^0))
\]

\[
+ \frac{d_2}{4f_B^2} \left[ 3(\bar{p}p + \bar{n}n)((\partial_\mu B^+)(\partial^\mu B^-) + (\partial_\mu B^0)(\partial^\mu \bar{B}^0)) \right.
\]

\[
+ (\bar{p}p - \bar{n}n)((\partial_\mu B^+)(\partial^\mu B^-) - (\partial_\mu B^0)(\partial^\mu \bar{B}^0)) \right\} \right]
\]

In \([12]\), the first term is the vectorial Weinberg Tomozawa interaction term, which is attractive for \( \bar{B} \) mesons, but repulsive for the \( B \) mesons. The second term is the scalar meson exchange term, which is attractive for both \( B \) and \( \bar{B} \) mesons. The third, fourth and fifth terms comprise the range term in the chiral model. The parameters \( d_1 \) and \( d_2 \) in the last two terms of the interaction Lagrangian given by \([12]\) are determined by fitting to the empirical values of the KN scattering lengths \([76–78]\) for \( I=0 \) and \( I=1 \) channels \([34, 55]\),

The dispersion relations for the $B$ and $\bar{B}$ mesons are obtained from the Fourier transformations of the equations of motion of these mesons. These are given as

$$-\omega^2 + \vec{k}^2 + m^2_{B(\bar{B})} - \Pi_{B(\bar{B})}(\omega, |\vec{k}|) = 0,$$

where $\Pi_{B(\bar{B})}$ denotes the self energy of the $B$ ($\bar{B}$) meson in the medium. For the $B$ meson doublet ($B^+, B^0$), and $\bar{B}$ meson doublet ($\bar{B}^-, \bar{B}^0$), the self energies are given by

$$\Pi_B(\omega, |\vec{k}|) = -\frac{1}{4f_B^2} \left[ 3(\rho_p + \rho_n) \pm (\rho_p - \rho_n) \right] \omega$$

$$+ \frac{m_B^2}{2f_B} (\sigma' + \sqrt{2}\zeta_b' \pm \delta')$$

$$+ \left[ -\frac{1}{f_B} (\sigma' + \sqrt{2}\zeta_b' \pm \delta') + \frac{d_1}{2f_B} (\rho_p^s + \rho_n^s) 
+ \frac{d_2}{4f_B^2} (3(\rho_p^s + \rho_n^s) \pm (\rho_p - \rho_n)) \right] \omega^2 - \vec{k}^2,$$

(13)

and

$$\Pi_{\bar{B}}(\omega, |\vec{k}|) = \frac{1}{4f_B^2} \left[ 3(\rho_p + \rho_n) \pm (\rho_p - \rho_n) \right] \omega$$

$$+ \frac{m_{\bar{B}}^2}{2f_B} (\sigma' + \sqrt{2}\zeta_b' \pm \delta')$$

$$+ \left[ -\frac{1}{f_B} (\sigma' + \sqrt{2}\zeta_b' \pm \delta') + \frac{d_1}{2f_B} (\rho_p^s + \rho_n^s) 
+ \frac{d_2}{4f_B^2} (3(\rho_p^s + \rho_n^s) \pm (\rho_p - \rho_n)) \right] \omega^2 - \vec{k}^2,$$

(14)

where the $\pm$ signs refer to the $B^+$ and $B^0$ respectively in equation (14) and to the $B^-$ and $\bar{B}^0$ mesons respectively in equation (15). In equations (14) and (15), $\sigma' (= \sigma - \sigma_0)$, $\zeta_b' (= \zeta_b - \zeta_{b0})$ and $\delta' (= \delta - \delta_0)$ are the fluctuations of $\sigma$, $\zeta_b$, and $\delta$, from their vacuum expectation values.

In the present calculations, we neglect the fluctuation of $\zeta_b \sim \langle \bar{b}b \rangle$, due to the reason that the fluctuation of the heavy quark condensates are small in the medium [58].

The masses of the charged open bottom mesons, $B^\pm$ have an additional positive mass shift due to the presence of the magnetic field, which retaining only the lowest Landau level, are given as

$$m_{B^\pm}^{eff} = \sqrt{m_{B^\pm}^*^2 + |eB|},$$

(16)

whereas for the neutral $B(\bar{B})$ mesons, namely, $B^0(\bar{B}^0)$, the effective masses are given as $m_{B^0(\bar{B}^0)}^{eff} = m_{B^0(\bar{B}^0)}^*$, which are the solutions for $\omega$ at $|\vec{k}| = 0$, of the dispersion relations given by equation (13), using the self-energies for $B^0$ and $\bar{B}^0$ as given by (14) and (15).
In the next section, we shall discuss the results for the $B(B)$-meson mass modifications in symmetric as well as asymmetric nuclear matter in the presence of an external magnetic field, $B$, as obtained in the present chiral effective model. The in-medium masses of the $B$ and $\bar{B}$ mesons are studied accounting for the effects of the anomalous magnetic moments of the nucleons.

III. RESULTS AND DISCUSSIONS

We study the medium modifications of the masses of open bottom mesons in nuclear matter in presence of an external magnetic field using a chiral effective model. The modifications of the masses of the $B$ and $\bar{B}$ mesons arise from their interactions with the protons and neutrons, as well as the scalar mesons (isoscalar non-strange $\sigma$ meson, isoscalar strange $\zeta$ meson and isovector $\delta$ meson). The calculations are carried out within the frozen glue-ball approximation, where the medium modifications of the scalar dilaton field, $\chi$, which mimicks the gluon condensates of QCD, are not taken into account. We study the mass modifications in $B$ and $\bar{B}$ mesons accounting for the anomalous magnetic moments (AMM) of the nucleons and compared to the case when AMM effects are not taken into consideration. The presence of magnetic field gives rise to Landau energy levels for charged baryon in the medium i.e. proton. The in-medium masses of the $B$ mesons are calculated from the dispersion relation (13), with their self-energies being given by (14). As already stated, the charged $B$ mesons undergo further mass modification in presence of magnetic field which is given by (16), retaining contribution due to lowest Landau energy level.

In the figures 1, 2, 3, 4, we show the masses of the $B$ and $\bar{B}$ mesons, plotted as functions of the baryon density in units of nuclear saturation density, $\rho_0$, for the values of the $eB$ as $2m_\pi^2$, $4m_\pi^2$, $6m_\pi^2$, $8m_\pi^2$, respectively. As can be seen from the expression for the self energy of the $B$ mesons given by (14), the isospin symmetric part of the Weinberg-Tomozawa term, is repulsive for the $B$ mesons (similar to the case of the $K$ ($\bar{D}$) mesons), whereas this term is attractive for the $\bar{B}$ mesons (similar to the $\bar{K}$ ($D$) mesons) in nuclear matter [31–34, 53–55]. The isospin symmetric part of Weinberg-Tomozawa term thus leads to an increase in the masses of both $B^+$ and $B^0$ mesons, whereas, the isospin asymmetric part of the Weinberg-Tomozawa term leads to a positive (negative) shift in the mass of the $B^0$ ($B^+$) meson, in
FIG. 1: The in-medium masses of the $B$ and $\bar{B}$ mesons plotted as functions of the baryon density (in units of nuclear matter saturation density) for various values of isospin asymmetry parameter, $\eta$, for value of the magnetic field, $eB = 2m_\pi^2$, accounting for the effects of anomalous magnetic moment. These results are compared to the case when the effects of the anomalous magnetic moment are not taken into account (shown as dotted line).
FIG. 2: The in-medium masses of the $B$ and $\bar{B}$ mesons plotted as functions of the baryon density (in units of nuclear matter saturation density) for various values of isospin asymmetry parameter, $\eta$, for value of the magnetic field, $eB = 4m_\pi^2$, accounting for the effects of anomalous magnetic moment. These results are compared to the case when the effects of the anomalous magnetic moment are not taken into account (shown as dotted line).
FIG. 3: The in-medium masses of the $B$ and $\bar{B}$ mesons plotted as functions of the baryon density (in units of nuclear matter saturation density) for various values of isospin asymmetry parameter, $\eta$, for value of the magnetic field, $eB = 6m_\pi^2$, accounting for the effects of anomalous magnetic moment. These results are compared to the case when the effects of the anomalous magnetic moment are not taken into account (shown as dotted line).
FIG. 4: The in-medium masses of the $B$ and $\bar{B}$ mesons plotted as functions of the baryon density (in units of nuclear matter saturation density) for various values of isospin asymmetry parameter, $\eta$, for value of the magnetic field, $eB = 8m^2_\pi$, accounting for the effects of anomalous magnetic moment. These results are compared to the case when the effects of the anomalous magnetic moment are not taken into account (shown as dotted line).
FIG. 5: The effective masses of $B$ ($B^+$, $B^0$) and $\bar{B}$ ($\bar{B}^-$, $\bar{B}^0$) mesons in MeV plotted as functions of $eB/m^2_\pi$, for baryon density, $\rho_B = \rho_0$, with different values of isospin asymmetry parameter, $\eta$, accounting for the effects of the anomalous magnetic moments (AMM) for the nucleons. The results are compared with the case of not accounting for the anomalous magnetic moments (shown as dotted lines).
FIG. 6: The effective masses of $B$ ($B^+ , B^0$) and $\bar{B}$ ($B^-, \bar{B}^0$) mesons in MeV plotted as functions of $eB/m^2_\pi$, for baryon density, $\rho_B = 3\rho_0$, with different values of isospin asymmetry parameter, $\eta$, accounting for the effects of the anomalous magnetic moments (AMM) for the nucleons. The results are compared with the case of not accounting for the anomalous magnetic moments (shown as dotted lines).
FIG. 7: The effective masses of $B$ ($B^+, B^0$) and $\bar{B}$ ($B^-, \bar{B}^0$) mesons in MeV plotted as functions of $eB/m^2$, for baryon density, $\rho_B = 5\rho_0$, with different values of isospin asymmetry parameter, $\eta$, accounting for the effects of the anomalous magnetic moments (AMM) for the nucleons. The results are compared with the case of not accounting for the anomalous magnetic moments (shown as dotted lines).
the isospin asymmetric nuclear medium. The scalar exchange term in the isospin symmetric nuclear matter is attractive for both the $B$ and $\bar{B}$ mesons. The last three terms of the self energies for the $B$ and $\bar{B}$ mesons, correspond to the range terms, which are repulsive for the first range term and attractive for the other two terms (the $d_1$ and $d_2$ terms) for isospin symmetric nuclear matter. Accounting for the contributions from all the terms (Weinberg-Tomozawa, scalar exchange and the range terms), as well as for a mass shift for the charged $B^\pm$ mesons due to Landau quantization, there is observed to be drop in the masses of the $B$ and $\bar{B}$ mesons with increase in density. There is observed to be increase in the in-medium masses of the $B$ as well as $\bar{B}$ mesons, when the effects of the anomalous magnetic moments are taken into account, as compared to when these are not taken into consideration. These effects are observed to be larger for higher values of the magnetic field.

For the $\bar{B}$ mesons, the isospin symmetric part of the Weinberg-Tomozawa term is attractive and leads to drop in their masses. With inclusion of isospin asymmetry in the medium, the mass of $\bar{B}^0$ is further decreased while that of $B^-$ has a positive contribution from isospin asymmetry. The behaviour of the $B^-$ meson mass is very similar to the mass of $B^+$ with isospin asymmetry, but the drop in the mass of $B^-$ is much larger than that of $B^+$ meson, due to the attractive Weinberg-Tomozawa contribution.

We next discuss, specifically, the effects of the magnetic field on the masses of the $B$ and $\bar{B}$ mesons for specific values of density and isospin asymmetry of the nuclear matter. These masses are plotted for densities $\rho_0$, $3\rho_0$ and $5\rho_0$ in figures 5, 6 and 7 respectively. For baryon density, $\rho_0$, as shown in figure 5 for symmetric nuclear matter, while accounting for the effects of anomalous magnetic moments (AMM) of the nucleons, the masses of $B^0$ as well as $\bar{B}^0$ are observed to show an initial increase with increase in magnetic field (upto $eB \sim 2 - 3m^2_\pi$), followed by negligible change with further increase in the magnetic field. In the absence of the AMM effects, in symmetric nuclear matter, the masses of these neutral $B$ and $\bar{B}$ mesons show a drop at small values of the magnetic field and marginal modifications as the magnetic field is further increased. However, the magnetic field is observed to have substantial influence on the masses of charged mesons $B^+$ and $B^-$ mesons (of around 15 MeV) in symmetric nuclear matter at $\rho_B = \rho_0$, due to the Landau quantization effects in the presence of magnetic field. The effects of the AMM are observed to lead to higher
masses of the $B$ as well as $\bar{B}$ mesons as compared to when these effects are not taken into account. For asymmetric nuclear matter, with AMM, the masses of $B^0$ and $\bar{B}^0$ undergo a slight increase ($\sim 5 - 8$ MeV) as magnetic field increases from zero to $8 m_{\pi}^2$, for $\rho_B = \rho_0$. The masses of charged mesons, $B^+$ and $B^-$ are observed to increase appreciably (by about $15 - 20$ MeV at $\rho_B = \rho_0$), for a similar increase in the magnetic field, due to the Landau quantization effects. There is no modification of the neutral $B^0$ and $\bar{B}^0$ meson masses for $\eta=0.5$ (when there are only neutrons in the system), due to the magnetic field, when the AMM effects are not taken into account, as the magnetic field effects for the neutrons are only due to the anomalous magnetic moment of the neutron. At higher densities of $\rho_B$ as $3 \rho_0$ and $5 \rho_0$, shown in figures 6 and 7, accounting for the effects of AMM is observed to increase the masses of $B$ and $\bar{B}$ mesons and these are observed to be larger at higher isospin asymmetry in the nuclear matter. In asymmetric nuclear matter whenever isospin asymmetry parameter, $\eta$ is different from 0.5 and if we neglect the effects of AMM, then the masses of these mesons are observed to drop as we increase the magnetic field. With AMM effects, at higher densities these masses are observed to rise with magnetic field, and, this increase is seen to be more appreciable for asymmetric nuclear matter. The difference between the masses, with and without considering the effects of AMM, is observed to be larger for higher values of magnetic field.

The $D$ and $\bar{D}$ mesons have been investigated in asymmetric nuclear matter in the presence of strong magnetic fields [19] within the chiral effective model used in the present investigation. We compare the effects of the magnetic field on the masses of the open bottom mesons studied in the present work with the effects on the masses of the $D$ and $\bar{D}$ mesons. The $\bar{B}$ ($D$) mesons have larger drop in their masses in the nuclear medium as compared to the $B$ ($\bar{D}$), arising mainly due to the Weinberg-Tomozawa interaction which is attractive for the former and repulsive for the latter. In the absence of magnetic field, in the asymmetric nuclear matter, the masses of the $B^+$ and $B^0$ mesons behave like the $\bar{D}^0$ and $D^-$ masses respectively, and, the masses of the $B^-$ and $\bar{B}^0$ mesons behave like $D^0$ and $D^+$ masses. In the presence of magnetic fields, there is additional (positive) mass shifts for the charged $D^\pm$ as well as $B^\pm$ mesons. However, due to the much larger mass of the $B$ and $\bar{B}$ mesons as compared to the $D$ and $\bar{D}$ mesons, the mass shift is much more appreciable for the charged
open charm mesons as compared to the charged open bottom mesons. At low densities \( (\rho_B \sim \rho_0) \), the changes in the masses of neutral \( B \) and \( D \) mesons are marginal as magnetic field is increased (with modifications upto around 5 MeV for the \( D^0 \) and \( \bar{D}^0 \) mesons and upto around 8 MeV for the \( B^0 \) and \( \bar{B}^0 \) mesons). At higher densities, accounting the effects of anomalous magnetic moment, the masses of all open bottom and open charm mesons show steady increase as magnetic field increases. In the absence of AMM effects, the masses of the \( D \) and \( \bar{D} \), as well as \( B \) and \( \bar{B} \) mesons are observed to be smaller than the case when the AMM effects are taken into account and the difference in the masses in the two situations is observed to be larger at higher densities.

The effects of the isospin asymmetry as well as magnetic fields on the masses of the \( B \) and \( \bar{B} \) mesons are observed to be large at high densities. The effects of the anomalous magnetic moments (AMM) are observed to be large at high values of the magnetic fields leading to an increase in the masses of the \( B \) and \( \bar{B} \) mesons as compared to when these effects are not taken into account. The AMM effects are observed to be more prominent in the isospin symmetric system, and becomes smaller (but still appreciable) with increase in isospin asymmetry in the nuclear medium.

IV. SUMMARY

To summarize, we have studied the medium modifications of the masses of the open bottom mesons in isospin asymmetric nuclear matter in presence of strong magnetic fields with emphasis on the behaviors when AMM of the nucleons are taken into consideration. We have used a chiral effective model which is generalized from SU(3) to SU(5) in order to describe the interactions of the open bottom mesons with the hadrons in the medium. Due to their interactions with the nucleons and the scalar fields, their masses are observed to be modified. In presence of external magnetic fields the number densities and scalar densities of protons have contributions from the Landau energy levels. The magnetic fields also show effects through the anomalous magnetic moments of the nucleons which we have studied in the present work. The effects of the isospin asymmetry as well as magnetic fields are observed to be large for both the \( B \) and \( \bar{B} \) meson doublets, especially at high densities. The isospin asymmetric effects being large at high densities should have observable effects
in the ratios of $B^+/B^0$ and $B^-/\bar{B}^0$, as well as in the decay widths of bottomonium states to $B\bar{B}$, in asymmetric heavy ion collisions planned at Compressed baryonic matter (CBM) experiments at FAIR at the future facility at GSI.

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