D-mesons in asymmetric nuclear matter

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Abstract

We calculate the in-medium $D$ and $\bar{D}$-meson masses in isospin asymmetric nuclear matter in an effective chiral model. The $D$ and $\bar{D}$ - mass modifications arising from their interactions with the nucleons and the scalar mesons in the effective hadronic model are seen to be appreciable at high densities and have a strong isospin dependence. These mass modifications can open the channels of the decay of the charmonium states ($\Psi', \chi_c, J/\Psi$) to $D\bar{D}$ pairs in the dense hadronic matter. The isospin asymmetry in the doublet $D = (D^0, D^+)$ is seen to be particularly appreciable at high densities and should show in observables like their production and flow in asymmetric heavy ion collisions in the compressed baryonic matter experiments in the future facility of FAIR, GSI. The results of the present work are compared to calculations of the $D(\bar{D})$ in-medium masses in the literature using the QCD sum rule approach, quark meson coupling model, coupled channel approach as well as from the studies of quarkonium dissociation using heavy quark potentials from lattice QCD at finite temperatures.

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I. INTRODUCTION

The in-medium properties of hadrons is a topic of intense research in high energy physics since the last few decades. The ongoing and future relativistic heavy ion collision experiments, at the high energy accelerators SPS, CERN, Switzerland; SIS, GSI, Germany; RHIC, BNL, USA; LHC, CERN, Switzerland, and the compressed baryonic matter (CBM) experiments at the future facilities at GSI, Germany, are intended to probe matter at high temperatures and densities. The hadrons modified in the hot and dense hadronic medium resulting from heavy ion collision experiments, affect the experimental observables. E.g., the dilepton spectra observed from heavy ion collision experiments at the SPS [1, 2] are attributed to the medium modifications of vector mesons [3, 4, 5, 6, 7], and can not be explained by vacuum hadronic properties. The production of kaons and antikaons as well as their collective flow are attributed to the modifications of their spectral functions in the medium [4, 8, 9, 10, 11, 12, 13, 14]. Due to the presence of the light quark/antiquark, the medium modifications of the $D$ ($\bar{D}$) mesons can be appreciable in the dense hadronic medium [15, 16, 17, 18, 19, 20, 21]. The observation of open charm enhancement in nuclear collisions [22, 23] as well as $J/\Psi$ suppression as observed at the SPS [24, 25, 26], can be due to medium modifications of the $D$ ($\bar{D}$) in the medium. In higher energy heavy ion collision experiments at RHIC as well as at LHC, the $J/\Psi$ suppression can arise from formation of a quark-gluon plasma (QGP) [27, 28]. However, the effect of the hadron absorption of $J/\Psi$ is not negligible [29, 30, 31]. With the drop in the mass of $D\bar{D}$ pair in the medium, the excited states of $J/\Psi$, a major source of yield of $J/\Psi$ [32], can decay to the $D(\bar{D})$ final states [33], thereby leading to $J/\Psi$ suppression [34] in the hadronic medium. The medium modification of the masses of $J/\Psi$ and their excited states in the hadronic medium, due to the D-meson mass modifications have also been considered in the literature [35] and the excited states of $J\Psi$ are seen to have appreciable mass dop in the medium. The effects due to the mass modifications of the D-mesons as well as the charmonium bound states on experimental observables could be explored at the future accelerator facility at GSI [36]. It is thus important to study the modifications of the charmed mesons in the medium and hence one has to understand the charmed meson interactions in the hadronic phase.

In the QCD sum rule approach, due to the presence of the light quark/antiquark, the
mass modification of the $D$-mesons arises from the light quark condensate \[15, 16, 19\]. In the quark meson coupling (QMC) model, the contribution from the $m_c\langle \bar{q}q \rangle_N$ term is represented by a quark-$\sigma$ meson coupling. The QMC model predicts the mass shift of the $D$-meson to be of the order of 60 MeV at nuclear matter density \[21\], which is very similar to the value obtained in the QCD sum rule calculations of Ref. \[15, 19\]. Furthermore, lattice calculations for heavy quark potentials at finite temperature suggest a similar drop \[20, 37\].

In this work we study the medium modification of the masses of open charm mesons ($D^{\pm}$) due to their interactions with the isoscalar-scalar mesons ($\sigma$ and $\zeta$ mesons), isovector-scalar $\delta$ mesons and the nucleons in asymmetric nuclear matter. The medium modification of the properties of kaons and antikaons in (isospin asymmetric) dense nuclear matter has been studied in a chiral $SU(3)$-flavor model in Ref. \[38, 39\] and have been extended to include the effects from hyperons in the asymmetric strange hadronic matter in ref. \[40\]. We generalise the model to $SU(4)$-flavor to derive the interactions of the $D$ mesons with the nucleons and scalar mesons and investigate their mass modifications in the asymmetric matter. The masses of the $D^{\pm}$ mesons in symmetric nuclear matter have been studied earlier in such an effective chiral model \[41\]. In a coupled channel approach for the study of $D$ mesons, using a separable potential, it was shown that the resonance $\Lambda_c(2593)$ is generated dynamically in the $I=0$ channel \[42\] analogous to $\Lambda(1405)$ in the coupled channel approach for the $\bar{K}N$ interaction \[43\]. The approach has been generalized to study the spectral density of the $D$-mesons at finite temperatures and densities \[44\], taking into account the modifications of the nucleons in the medium. The results of this investigation seem to indicate a dominant increase in the width of the $D$-meson whereas there is only a very small change in the $D$-meson mass in the medium \[44\]. However, these calculations \[42, 44\], assume the interaction to be $SU(3)$ symmetric in $u,d,c$ quarks and ignore channels with charmed hadrons with strangeness. A coupled channel approach for the study of $D$-mesons has been developed based on $SU(4)$ symmetry \[45\] to construct the effective interaction between pseudoscalar mesons in a 16-plet with baryons in 20-plet representation through exchange of vector mesons and with KSFR condition \[46\]. This model \[45\] has been modified in aspects like regularization method and has been used to study DN interactions in Ref. \[47\]. This reproduces the resonance $\Lambda_c(2593)$ in the $I=0$ channel and in addition generates
another resonance in the I=1 channel at around 2770 MeV. These calculations have been
generalized to finite temperatures \cite{48} accounting for the in-medium modifications of the
nucleons in a Walecka type $\sigma - \omega$ model, to study the $D$ and $\bar{D}$ properties \cite{49} in the hot
and dense hadronic matter.

Within the effective chiral model considered in the present investigation, the $D(\bar{D})$
energies are modified due to a vectorial Weinberg-Tomozawa, due to scalar exchange terms
($\sigma, \zeta, \delta$) as well as range terms \cite{39, 40}. The isospin asymmetric effects among $D^0$ and $D^+$
in the doublet, $D \equiv (D^0, D^+)$ as well as between $\bar{D}^0$ and $D^-$ in the doublet, $\bar{D} \equiv (\bar{D}^0, D^-)$
arise due to the scalar-isovector $\delta$ meson, due to asymmetric contributons in the Weinberg-
Tomozawa term, as well as in the range term \cite{38}.

We organize the paper as follows: We briefly recapitulate the $SU(3)$-flavor chiral model
adopted for the description of the asymmetric hadronic matter \cite{39, 40} in Section II. The
properties then are studied within this approach. These give rise to medium modifications for
the $D$-masses through their interactions with the nucleons and scalar mesons as presented in
Section III. Section IV discusses the results of the present investigation, while we summarise
our findings and discuss possible outlook in Section V.

II. THE HADRONIC CHIRAL $SU(3) \times SU(3)$ MODEL

The effective hadronic chiral Lagrangian used in the present work is given as

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

are discussed. Eq. (1) corresponds to a relativistic quantum field theoretical model of
baryons and mesons adopting a nonlinear realization of chiral symmetry \cite{50, 51, 52} and
broken scale invariance (for details see \cite{53, 54, 55}) as a description of the hadronic matter.
The model was used successfully to describe nuclear matter, finite nuclei, hypernuclei
and neutron stars. The Lagrangian contains the baryon octet, the spin-0 and spin-1 meson
multiplets as the elementary degrees of freedom. In Eq. (1), $\mathcal{L}_{\text{kin}}$ is the kinetic energy
term, $\mathcal{L}_{BW}$ contains the baryon-meson interactions in which the baryon-spin-0 meson interaction
terms generate the baryon masses. $\mathcal{L}_{\text{vec}}$ 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. $L_0$ contains the meson-meson interaction terms inducing the spontaneous breaking of chiral symmetry as well as a scale invariance breaking logarithmic potential. $L_{SB}$ describes the explicit chiral symmetry breaking.

The baryon-scalar meson interactions generate the baryon masses and the parameters corresponding to these interactions are adjusted so as to obtain the baryon masses as their experimentally measured vacuum values. For the baryon-vector meson interaction terms, there exist the $F$-type (antisymmetric) and $D$-type (symmetric) couplings. Here we will use the antisymmetric coupling \[39, 40, 53\] because, following the universality principle \[56\] and the vector meson dominance model, one can conclude that the symmetric coupling should be small. Additionally we choose the parameters \[38, 53\] so as to decouple the strange vector field $\phi_\mu \sim \bar{s}\gamma_\mu s$ from the nucleon, corresponding to an ideal mixing between $\omega$ and $\phi$. A small deviation of the mixing angle from the ideal mixing \[57, 58, 59\] has not been taken into account in the present investigation.

The Lagrangian densities corresponding to the interaction for the vector meson, $L_{vec}$, the meson-meson interaction $L_0$ and that corresponding to the explicit chiral symmetry breaking $L_{SB}$ have been described in detail in references \[38, 53\].

To investigate the hadronic properties in the medium, we write the Lagrangian density within the chiral SU(3) model in the mean field approximation and determine the expectation values of the meson fields by minimizing the thermodynamical potential \[54, 55\].

### III. $D$ AND $\bar{D}$ MESONS IN THE MEDIUM

We now examine the medium modifications for the $D$ and $\bar{D}$-meson masses in the asymmetric nuclear matter. The properties of nucleons and scalar mesons have been studied in the asymmetric hadronic matter within a chiral SU(3) model \[39\]. We assume that the additional effect of charmed particles in the medium leads to only marginal modifications \[60\] of these hadronic properties and do not need to be taken into account here. However, to investigate the medium modification of the $D$-meson mass, we need to know the interactions of the $D$-mesons with the light hadron sector.

The light quark condensate has been shown to play an important role for the shift in the $D$-meson mass in the QCD sum rule calculations \[15\]. In the present chiral model,
the interactions to the scalar fields (nonstrange, \(\sigma\) and strange, \(\zeta\)) as well as a vectorial Weinberg-Tomozawa interaction term modify the masses for \(D^\pm\) mesons in the medium. These interactions were considered within the SU(3) chiral model to investigate the modifications of K-mesons in the dense (asymmetric) hadronic medium [38, 39, 40].

To consider the medium effects on the \(D\) and \(\bar{D}\)-meson masses we generalize the chiral \(SU(3)\)-flavor model to include the charmed mesons. The scalar meson multiplet has now the expectation value

\[
\langle X \rangle = \begin{pmatrix}
\frac{(\sigma + \delta)}{\sqrt{2}} & 0 & 0 & 0 \\
0 & \frac{(\sigma - \delta)}{\sqrt{2}} & 0 & 0 \\
0 & 0 & \zeta & 0 \\
0 & 0 & 0 & \zeta_c
\end{pmatrix},
\]

with \(\zeta_c\) corresponding to the \(\bar{c}c\)--condensate. The pseudoscalar meson field \(P\) can be written, including the charmed mesons, as

\[
P = \begin{pmatrix}
\frac{\pi^0}{\sqrt{2}} & \pi^+ & \frac{2K^+_0}{1+w} & \frac{2D^+_0}{1+w_c} \\
\pi^- & -\frac{\pi^0}{\sqrt{2}} & \frac{2K^-_0}{1+w} & \frac{2D^-_0}{1+w_c} \\
\frac{2K^-}{1+w} & \frac{2K^+_0}{1+w} & 0 & 0 \\
\frac{2D^-}{1+w_c} & \frac{2D^+_0}{1+w_c} & 0 & 0
\end{pmatrix},
\]

where \(w = \sqrt{2}\zeta/\sigma\) and \(w_c = \sqrt{2}\zeta_c/\sigma\). From PCAC, one gets the decay constants for the pseudoscalar mesons as \(f_\pi = -\sigma\), \(f_K = - (\sigma + \sqrt{2}\zeta)/2\) and \(f_D = - (\sigma + \sqrt{2}\zeta_c)/2\). In the present calculations, the value for the D-decay constant will be taken to be 135 MeV [19]. We note that for the decay constant of \(D^+\), the Particle Data Group [61] quotes a value of \(f_D^+ \approx 200\,\text{MeV}\). Taking a similar value also for \(f_D\) would not affect our results qualitatively, however (see also [62]).

The interaction Lagrangian modifying the \(D\)-meson mass can be written as [39]

\[
\mathcal{L}_{DN} = -\frac{i}{8f_D^2} \left[ 3 \bar{p} \gamma^\mu p + \bar{n} \gamma^\mu n \right] \left[ (D^0(\partial_\mu \bar{D}^0) - (\partial_\mu D^0)\bar{D}^0) + (D^+(\partial_\mu \bar{D}^-) - (\partial_\mu D^+)\bar{D}^-) \\
+ \left( \bar{p} \gamma^\mu p - \bar{n} \gamma^\mu n \right) \left( D^0(\partial_\mu \bar{D}^-) - (\partial_\mu D^-)\bar{D}^0 \right) - \left( D^+(\partial_\mu D^-) - (\partial_\mu D^+)\bar{D}^- \right) \right] \\
+ \frac{m_D^2}{2f_D^2} \left[ (\sigma + \sqrt{2}\zeta_c)(D^0 D^0 + (D^- D^+)) + \delta(D^0 D^0) - (D^- D^+) \right] \\
- \frac{1}{f_D} \left[ (\sigma + \sqrt{2}\zeta_c)((\partial_\mu \bar{D}^0)(\partial^\mu D^0) + (\partial_\mu D^-)(\partial^\mu D^+)) \right]
\]
\[ + \delta \left( (\partial_\mu \bar{D}^0)(\partial^\mu D^0) - (\partial_\mu D^-)(\partial^\mu D^+) \right) \]
\[ + \frac{d_1}{2f_D^2} (\bar{p}p + \bar{n}n)((\partial_\mu D^-)(\partial^\mu D^+) + (\partial_\mu D^0)(\partial^\mu D^0)) \]
\[ + \frac{d_2}{4f_D^2} \left[ (\bar{p}p + \bar{n}n)((\partial_\mu \bar{D}^0)(\partial^\mu D^0) + (\partial_\mu D^-)(\partial^\mu D^+)) \right. \]
\[ + (\bar{p}p - \bar{n}n)((\partial_\mu \bar{D}^0)(\partial^\mu D^0) - (\partial_\mu D^-)(\partial^\mu D^+)) \] \hspace{1cm} (4)

In (4) the first term is the vectorial interaction term obtained from the kinetic term in (1). The second term, which gives an attractive interaction for the \( D \)-mesons, is obtained from the explicit symmetry breaking term in (1). The third term arises from kinetic term of the pseudoscalar mesons [39, 40]. The fourth and fifth terms have been written down for the DN interactions, in analogy with the \( d_1 \) and \( d_2 \) terms in chiral SU(3) model [39, 40]. The last three terms in (4) represent the range term in the chiral model. It might be noted here that the interaction of the pseudoscalar mesons to the vector mesons, in addition to the pseudoscalar meson–nucleon vectorial interaction leads to a double counting in the linear realization of the chiral effective theory [63]. Within the nonlinear realization of the chiral effective theories, such an interaction does not arise in the leading or sub-leading order, but only as a higher order contribution [63]. Hence the vector meson-pseudoscalar interaction will not be considered within the present investigation.

The Fourier transformations of the equations of motion for \( D \) and \( \bar{D} \) mesons yield the dispersion relations,
\[ \omega^2 + \vec{k}^2 + m_D^2 - \Pi(\omega, |\vec{k}|) = 0, \] \hspace{1cm} (5)
where \( \Pi \) denotes the self energy of the \( D \) (\( \bar{D} \)) meson in the medium.

Explicitly, the self energy \( \Pi(\omega, |\vec{k}|) \) for the \( D \) meson doublet, \( (D^0, D^+) \) arising from the interaction (4) is given as
\[ \Pi(\omega, |\vec{k}|) = \frac{1}{4f_D^2} \left[ 3(\rho_p + \rho_n) \pm (\rho_p - \rho_n) \right] \omega \]
\[ + \frac{m_D^2}{2f_D} (\sigma' + \sqrt{2}\zeta_c' \pm \delta') \]
\[ + \left[ - \frac{1}{f_D}(\sigma' + \sqrt{2}\zeta_c' \pm \delta') + \frac{d_1}{2f_D^2}(\rho^s_p + \rho^s_n) \right. \]
\[ + \left. \frac{d_2}{4f_D^2} \left( (\rho^s_p + \rho^s_n) \pm (\rho^s_p - \rho^s_n) \right) \right] (\omega^2 - \vec{k}^2), \] \hspace{1cm} (6)
where the ± signs refer to the $D^0$ and $D^+$ respectively. In the above, $\sigma' (= \sigma - \sigma_0)$, $\zeta_c' (= \zeta_c - \zeta_{c0})$ and $\delta' (= \delta - \delta_0)$ are the fluctuations of the scalar-isoscalar fields $\sigma$ and $\zeta_c$, and the third component of the scalar-isovector field, $\delta$, from their vacuum expectation values. The vacuum expectation value of $\delta$ is zero ($\delta_0 = 0$), since a nonzero value for it will break the isospin symmetry of the vacuum (we neglect here the small isospin breaking effect arising from the mass and charge difference of the up and down quarks). In the above, $\rho_p$ and $\rho_n$ are the number densities of proton and neutron and $\rho^s_p$ and $\rho^s_n$ are their scalar densities.

Similarly, for the $\bar{D}$ meson doublet, $(\bar{D}^0, D^-)$, the self-energy is calculated as

$$
\Pi(\omega, |\vec{k}|) = -\frac{1}{4f_D^2} \left[ 3(\rho_p + \rho_n) \pm (\rho_p - \rho_n) \right] \omega 
+ \frac{m_D^2}{2f_D}(\sigma' + \sqrt{2}\zeta_c' \pm \delta') 
+ \left[ -\frac{1}{f_D} (\sigma' + \sqrt{2} \zeta_c' \pm \delta') + \frac{d_1}{2f_D^2}(\rho^s_p + \rho^s_n) 
\right. 
+ \frac{d_2}{4f_D^2} \left( (\rho^s_p + \rho^s_n) \pm (\rho^s_p - \rho^s_n) \right) (\omega^2 - |\vec{k}|^2),
$$

where the ± signs refer to the $\bar{D}^0$ and $D^-$ respectively. The optical potentials are calculated from the energies of the $D$ and $\bar{D}$ mesons

$$
U(\omega, k) = \omega(k) - \sqrt{k^2 + m_D^2},
$$

where $m_D$ is the vacuum mass for the $D(\bar{D})$ meson and $\omega(k)$ is the momentum dependent energy of the $D(\bar{D})$ meson.

The parameters $d_1$ and $d_2$ are determined by a fit of the empirical values of the KN scattering lengths $[64, 65, 66]$ for $I=0$ and $I=1$ channels $[39, 40]$. In the next section, we shall discuss the results for the $D$-meson mass modification obtained in the present effective chiral model as compared to the results in the literature, obtained from other approaches.

IV. RESULTS AND DISCUSSIONS

To study the $D(\bar{D})$-meson masses in asymmetric nuclear medium due to its interactions with the light hadrons, we have generalized the chiral SU(3) model used for the study
of the dense hadronic matter to SU(4) for the meson sector. The contributions from the various terms of the interaction Lagrangian (4) to the masses of the $D \equiv (D^0, D^+)$ and $\bar{D} \equiv (\bar{D}^0, D^-)$ are shown in figures 1 and 2, as functions of density. These are illustrated for the isospin asymmetric case with the value of the asymmetry parameter, $\eta = (\rho_n - \rho_p)/(2\rho_B)$ as 0.5 and compared with the results obtained for the isospin symmetric case ($\eta = 0$). The isospin symmetric part ($\sim (\rho_n + \rho_p)$) of the first term of (4), called the Weinberg-Tomozawa term, is attractive for $D \equiv (D^0, D^+)$ mesons and leads to a drop of the masses of the $D^+$ and $D^0$ mesons, whereas it is repulsive for the $\bar{D}$ mesons in the nuclear medium, leads to an increase of the $\bar{D}^0$ and $D^-$ meson masses. In the isospin asymmetric nuclear medium, the Weinberg-Tomozawa term, leads to a mass splitting of the $D^0$ and $D^+$ mesons, giving a further drop in the mass of $D^+$, whereas the asymmetry reduces the drop of the mass in $D^0$. The D-meson self energy arising from the Weinberg-Tomozawa interaction, $\Pi_{WT}(\omega, |\vec{k}|)$ is given by the first term of equation (6), and at low densities, this turns out to be much smaller than $(\vec{k}^2 + m_D^2)$. One can then, as a first approximation replace $\Pi_{WT}(\omega, |\vec{k}|)$ by $\Pi_{WT}(m_D, |\vec{k}|)$ and solve for dispersion relation given by (5). Confining our attention to the Weinberg-Tomozawa interaction only, the energies of the $D^0$ and $D^+$ mesons, in the above approximation, are given by

$$\omega(|\vec{k}|) \simeq (|\vec{k}|^2 + m_D^2)^{1/2} - \frac{1}{8f_D^2}[3(\rho_p + \rho_n) \mp (\rho_p - \rho_n)]$$

(9)

One can note from the above equation (9) that at low densities, a given isospin asymmetry introduces equal increase (drop) for mass of $D^0(D^+)$ meson at low densities. However, at higher densities, there are deviations from the analytical expressions given by (9) as expected, though the qualitative feature of the $D^+ (D^0)$ experiencing an attractive (repulsive) interaction from the vectorial Weinberg-Tomozawa term of (6) still remains the same, as can be seen from figure 1. One sees, from figure 2, that there is an increase in the masses of the $D^-$ and $\bar{D}^0$ and their mass shifts are equal for the case of isospin symmetric matter ($\eta = 0$). However, in the presence of asymmetry, the $D^-$ mass is seen to have an increase whereas $\bar{D}^0$ mass drops in the asymmetric medium. These behaviours for $D^-$ and $\bar{D}^0$ can be understood by examining the asymmetric contributions of the Weinberg-Tomozawa term (the first term of (7)).

The scalar meson exchange contribution to the self energy is given by the second term of
equation (6) for $D$ mesons and by the second term of (7) for $\bar{D}$ mesons. Its interaction is attractive and is identical for both the $D$ and $\bar{D}$ doublets for the isospin symmetric nuclear matter (a negligible difference in the energies is due to the difference in their vacuum masses, $m_{D^+} = m_{D^-} = 1869$ MeV and $m_{D^0} = m_{\bar{D}^0} = 1864.5$ MeV). One might notice from the self energy terms due to the scalar meson exchange that a nonzero value of the scalar isovector $\delta$-meson arising due to isospin asymmetry in nuclear matter, gives a drop in the masses for the $D^+$ and $D^-$, whereas the $\delta$ contribution is repulsive for $D^0$ and $\bar{D}^0$ ($\sigma' = \sigma - \sigma_0 > 0$ and $\delta' = \delta < 0$). This gives the $D^0$ mass to be higher than $D^+$ mass for asymmetric nuclear matter as seen in figure 1 and the mass of $\bar{D}^0$ mass (identical to $D^0$) to be higher than $D^-$ mass (identical to $D^+$ mass) as plotted in figure 2. However, one might notice that the shifts in the $D^+(D^-)$ and $D^0(\bar{D}^0)$ masses about the isospin symmetric case ($\eta=0$) case are not equal and opposite. The reason for the seen asymmetry in the mass splittings is the following. For low baryon densities, one has $\sigma' \sim \rho_s \approx \rho_B$, and $\delta' \sim (\rho_p^s - \rho_n^s) \approx (\rho_p - \rho_n)$, so that one would expect the splittings of the masses of $D^+$ and $D^0$ to be symmetrical (about the isospin symmetric matter) for a given baryon density $\rho_B$ and isospin asymmetry parameter, $\eta$. However, these no longer hold good for higher densities and the mean field $\sigma$ calculated in the isospin symmetric situation ($\eta=0$) turns out to be different for the asymmetric situation when the coupled equations of motions are solved for the scalar mean fields due to the presence of the scalar isovector $\delta$ field, as compared to when these equations are solved for symmetric nuclear matter (for $\eta=0$) in the absence of the $\delta$ meson.

The contributions to the $D$ and $\bar{D}$ self energies due to range terms are given by the last three terms of the right hand side of equations (6) and (7) respectively. The first term of the range terms given by the third term in (6) is repulsive whereas the second and third range terms have attractive contributions, when the isospin asymmetry is not taken into account. However, the isospin asymmetry, due to a nonzero value of the $\delta$ field, leads to an increase in the masses of the $D^+$ and $D^-$ mesons and a drop in the masses of $D^0$ and $\bar{D}^0$ from the isospin symmetric case. The second of the range terms (the $d_1$ term) is attractive and gives identical mass drops for $D^+$ and $D^0$ in the $D$ doublet as well as for $D^-$ and $\bar{D}^0$ in the $\bar{D}$ doublet. This term is proportional to $(\rho_p^s + \rho_n^s)$, which turns out to be different for the isospin asymmetric case as compared to the isospin symmetric nuclear matter, due
to the presence of the $\delta$ meson. This is because the equations of motion for the scalar fields for the two situations (with/without $\delta$ mesons) give different values for the mean field, $\sigma \sim (\rho_s^p + \rho_s^n)$. The last term of the range term (the $d_2$ term) has a negative contribution for the energies of $D^+$ and $D^0$ mesons as well as for $D^-$ and $\bar{D}^0$ mesons for the isospin symmetric matter. The isospin asymmetric part arising from the $((\rho_s^n - \rho_s^p))$ term of the $d_2$ term has a further drop in the masses for $D^\pm$ mesons, whereas it increases the masses of the $D^0$ and $\bar{D}^0$ mesons from their isospin symmetric values. One sees from figures 1 and 2 that for both $D^+$ and $D^-$ mesons, the modification in the masses arising from the range terms, for the asymmetric matter ($\eta=0.5$), as compared to the isospin symmetric matter is negligible due to the increase from the first two range terms almost cancelling with the drop due to the $d_2$ term. For the $D^0$ and $\bar{D}^0$ mesons, the mass is seen to increase as compared to the isospin symmetric matter, due to the increase due to the second and third range terms dominating over the drop due to first range term. The values of the parameters $d_1$ and $d_2$ are fitted from the kaon-nucleon scattering lengths [38, 39] to be $2.56/m_K$ and $0.73/m_K$ respectively [40] and it is seen that the $d_2$ term has a smaller contribution as compared to the $d_1$ term. At high densities, these attractive $d_1$ and $d_2$ terms dominate over the first range term (repulsive) and this leads to a decrease of the masses of the $D$ and $\bar{D}$ mesons. There is seen to be a substantial drop of D-meson masses at high densities due to the inclusion of this range term.

The density dependence of the masses of the $D$ mesons at specific values of the isospin asymmetric parameter, $\eta$, are shown in figure 3. The isospin asymmetry is seen to give a rise (drop) in the masses of the $D^0$ ($D^+$) as compared to their masses in the symmetric matter. For the isospin symmetric nuclear matter ($\eta=0$), the drop in the mass of the $D^+$ at $\rho_B = \rho_0$ is about 81 MeV from its vacuum value of 1869 MeV and $D^-$ meson mass also has a drop of about 30 MeV, giving a mass splitting between the $D^+$ and $D^-$ mesons as about 51 MeV. A similar drop of the D-meson mass is also predicted by the QMC model [21] as well as for the isospin averaged $D(\bar{D})$ mass in the QCD sum rule approach in [15]. In the recent QCD sum rule calculations [16], the mass drop for $D^+D^-$ is about 50 MeV and the mass splitting is about 90 MeV at $\rho = \rho_0$, whereas for the isospin symmetric nuclear matter, we obtain these values as 110 MeV and 50 MeV respectively. The coupled channel calculations
show a dominant modification of the D-meson width and only a small change in the $D$ meson mass in the medium. When the DN interaction is taken to be a Tomozawa Weinberg interaction supplemented by a scalar-isoscalar interaction \[48\], the mass modification of the $D$ mesons is seen to be around 10 MeV at $\rho_0$ and about 50 MeV for a density of $2\rho_0$. For the $D^-$ meson, there is an increase in the mass by about 10 MeV and 30 MeV at densities of $\rho_0$ and $2\rho_0$ respectively.

In the present investigation, the isospin dependence is seen to be quite prominent for $D^0$ as compared to $D^+$. The isospin asymmetry in the medium is seen to give an increase in the $D^0$ mass, whereas it gives a drop for the $D^+$ mass as compared to the $\eta=0$ case. For the isospin asymmetry parameter, $\eta=0.5$, the $D^0$ mass is seen to rise by about 23 MeV and 105 MeV for $\rho_B$ as $\rho_0$ and $5\rho_0$ respectively from their values of 1783 MeV and 1441 MeV in isospin symmetric matter. For the same value of asymmetry parameter, $\eta$, the $D^+$ is seen to have a mass drop of about 18 MeV and 39 MeV for $\rho_B = \rho_0$ and $5\rho_0$ respectively, from its $\eta=0$ values. This strong isospin dependence of the $D$ mesons should show up in observables like their production as well as flow in asymmetric heavy ion collisions planned at the future facility at FAIR, GSI. Figure 4 shows the masses of the $D^-$ and $\bar{D}^0$. It is seen that at high densities, both the $D^-$ and $\bar{D}^0$ are seen to have an increase in their masses in the asymmetric nuclear matter, as compared to the isospin symmetric case and this rise in the masses are seen to be similar for both $D^-$ and $\bar{D}^0$. For example, at nuclear matter density, the masses of $D^-$ and $\bar{D}^0$ are 1838 MeV and 1840 MeV respectively for $\eta=0.5$, which are very similar to the values of 1839 MeV and 1834.5 MeV for the isospin symmetric matter. For $\rho_B = 5\rho_0$, the masses of $D^-$ and $\bar{D}^0$ are 1690 MeV and 1698 MeV respectively, which are higher by about 31 MeV and 43 MeV from the isospin symmetric case. One sees that the values of $D^-$ and $\bar{D}^0$ masses remain very similar at a given isospin asymmetry. But it is seen that the density effects on these masses are quite appreciable (a drop of about 30 MeV for nuclear matter density and of about 180 MeV for a density about five times nuclear matter density for $\eta=0.5$).

Figures 5 and 6 show the optical potentials for the $D$ and $\bar{D}$ doublets as functions of the momentum. These are shown for densities $\rho_0$ and $5\rho_0$. The isospin dependence of the optical potentials is seen to be quite significant for high densities for the doublet $(D^+, D^0)$,
as has also been seen in the case of their masses. But, we see the optical potentials for the $D$ have very similar values for $D^-$ and $\bar{D}^0$ for a fixed value of the isospin asymmetric parameter. But, as already has been seen for the case of their masses, we see the optical potentials for both $D^-$ and $\bar{D}^0$ are quite different from the symmetric nuclear matter case at high densities.

The decay widths of the charmonium states can be modified by the level crossings between the excited states of $J/\Psi$ (i.e., $\Psi', \chi_c$) and the threshold for $D\bar{D}$ creation due to the medium modifications of the $D$-meson masses [18]. In the vacuum, the resonances above the $D\bar{D}$ threshold, for example the $\Psi''$ state, has a width of 25 MeV due to the strong open charm channel. On the other hand, the resonances below the threshold have a narrow width of a few hundreds of KeV, only. With the medium modification of the $D(\bar{D})$-meson masses, the channels for the excited states of $J/\Psi$, like $\chi_c$, $\Psi'$ decaying to $D^+D^-$ or $D^0\bar{D}^0$ pairs can open up in the dense hadronic medium. This can increase the decay widths of $\Psi'$ and $\chi_c$ states at high densities. In figure 7, we show the density dependence of the mass of the $D^+D^-$ as well as $D^0\bar{D}^0$ pair, calculated in the present investigation. We see that the decay to $D^+D^-$ channel is almost insensitive to isospin asymmetry. On the other hand, the decay channel to the $D^0\bar{D}^0$ is seen to have a strong isospin asymmetry dependence, with the asymmetry shifting the onset of the decay to higher densities. The decay of charmonium states to $D\bar{D}$ has been studied in Ref. [18, 33]. It is seen to depend sensitively on the relative momentum in the final state. These excited states might become narrow [18] though the $D$-meson mass is decreased appreciably at high densities. It may even vanish at certain momenta corresponding to nodes in the wavefunction [18]. Though the decay widths for these excited states can be modified by their wave functions, the partial decay width of $\chi_{c2}$, due to absence of any nodes, can increase monotonically with the drop of the $D^+D^-$ pair mass in the medium [18]. This can give rise to depletion in the $J/\Psi$ yield in heavy ion collisions. The dissociation of the quarkonium states ($\Psi', \chi_c, J/\Psi$) into $D\bar{D}$ pairs have also been studied [20, 67] by comparing their binding energies with lattice results on temperature dependence of heavy quark effective potential [37].
V. SUMMARY

To summarize we have investigated in a chiral model the in-medium masses of the $D$, $\bar{D}$-mesons in asymmetric nuclear matter, arising due to their interactions with the nucleons and scalar mesons. The properties of the light hadrons – as studied in a $SU(3)$ chiral model – modify the $D(\bar{D})$-meson properties in the dense hadronic medium. The $SU(3)$ model with parameters fixed from the properties of hadron masses in the vacuum, and low energy KN scattering data, is extended to $SU(4)$ to derive the interactions of $D(\bar{D})$-mesons with the light hadron sector. The mass modifications for the $D$ mesons are seen to be similar to earlier finite density calculations of QCD sum rules [16, 19] as well as to the quark meson coupling (QMC) model [21], in contrast to the small mass modifications in the coupled channel approach [44, 48]. In our calculations, the presence of the repulsive range term (the fourth term of (4) is compensated by the attractive $d_1$ and $d_2$ terms given by the last two terms in (4) and the latter (dominated by the $d_1$ term) have an effect of reducing the masses of both $D$ and $\bar{D}$ mesons at high densities.

The medium modification of the $D$-masses can lead to a suppression in the $J/\Psi$- yield in heavy-ion collisions, since the excited states of $J/\Psi$ and at a much higher density ($\simeq 5\rho_0$), $J/\Psi$ can decay to $D\bar{D}$ pairs in the dense hadronic medium. The decay to $D^+D^-$ pair seems to be insensitive to isospin dependence, whereas due to increase in the mass of $D^0\bar{D}^0$ in the asymmetric medium, isospin asymmetry is seen to disfavour the decay of the charmonium states to $D^0\bar{D}^0$ pair. The isospin dependence of $D^+$ and $D^0$ masses is seen to be dominant medium effect at high densities which might show in their production ($D^+/D^0$), whereas for the $D^-$ and $\bar{D}^0$, one does see that even though these have a strong density dependence, their in-medium masses remain similar at a given value for the isospin asymmetry parameter $\eta$. The strong density dependence as well as the isospin dependence of the $D(\bar{D})$ meson optical potentials in the asymmetric nuclear matter can be tested in the asymmetric heavy ion collision experiments at the future GSI facility [36].
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FIG. 2: (Color online) The various contributions to $\bar{D}$ meson energies at zero momentum (for $D^-$ in (a) and for $D^0$ in (b)) in MeV plotted as functions of the baryon density in units of nuclear matter saturation density, $\rho_B/\rho_0$ are shown for the isospin asymmetry parameter, $\eta=0.5$ and compared with the case of $\eta=0$ (dotted line).
FIG. 3: (Color online) The D meson energies for zero momentum \((k=0)\) (for \(D^+\) in (a) and for \(D^0\) in (b)) in MeV plotted as functions of the baryon density in units of \(\rho_0\), \(\rho_B/\rho_0\) for different values of the isospin asymmetry parameter, \(\eta\).

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