Mass modification of D-meson in hot hadronic matter

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We evaluate the in-medium D and $\bar{D}$-meson masses in hot hadronic matter induced by interactions with the light hadron sector described in a chiral SU(3) model. The effective Lagrangian approach is generalized to SU(4) to include charmed mesons. We find that the D-mass drops substantially at finite temperatures and densities, which open the channels of the decay of the charmonium states (Ψ', χ_c, J/Ψ) to $D\bar{D}$ pairs in the thermal medium. The effects of vacuum polarisations from the baryon sector on the medium modification of the D-meson mass relative to those obtained in the mean field approximation are investigated. The results of the present work are compared to calculations based on the QCD sum-rule approach, the quark-meson coupling model, chiral perturbation theory, as well as to studies of quarkonium dissociation using heavy quark potential from lattice QCD.

1. INTRODUCTION

The in-medium properties of hadrons have been a field of active theoretical research in the last decade. From the experimental side in-medium effects have also been probed in relativistic heavy-ion collisions from SIS (1-2 A·GeV) to SPS energies (30-160 A·GeV). At the SPS the experimentally observed dilepton spectra [1, 2] – that offer a direct glance at the in medium spectral properties of the vector mesons – have been attributed to medium modifications of the spectral function especially of the ρ-meson [3, 4, 5, 6, 7] and can not be explained by vacuum hadronic properties.

Furthermore, experiments on $K^{\pm}$ production from nucleus-nucleus collisions at SIS energies (1– 2 A· GeV) have shown that in-medium properties of kaons are seen in the collective flow pattern of K^+ mesons both, in-plane and out-of-plane, as well as in the abundancy and spectra of antikaons [4, 8, 9, 10, 11, 12, 13, 14]. Since the strangeness sector ($K^{\pm}, \Lambda_s, \Sigma_s$) shows some analogy to the open charm sector ($D^{\pm}, \Lambda_c, \Sigma_c$), which results from exchanging strange/antistrange quarks by charm/anticharm quarks, the medium modifications for the D-mesons have become a subject of recent interest, too [15, 16, 17, 18, 19, 20]. It is expected that one might find open charm enhancement in nucleus-nucleus collisions [21] as well as J/Ψ suppression as experimentally observed at the SPS [22]. Indeed, the NA50 Collaboration has claimed to see an open charm enhancement by up to a factor of three in central $Pb + Pb$ collisions at 158 A·GeV [23].

Open charm mesons ($D, D^*, \bar{D}, \bar{D}^*$) can be produced abundantly in high energy heavy-ion collisions and might even dominate the high mass (M > 2 GeV) dilepton spectra [24]. The medium modification of D-mesons is worth investigating since they should modify the J/Ψ absorption cross section in the medium with baryons and mesons and could also provide a possible explanation for the observed J/Ψ suppression [22]. On the other hand, in high energy heavy-ion collisions at RHIC ($\sqrt{s} \sim 200$ GeV), an appreciable contribution of J/Ψ suppression is expected to be due to the formation of a quark-gluon Plasma (QGP) [25]. However, the effect of the hadron absorption of J/Ψ's is still not negligible [26, 27, 28]. It is thus important to understand the charmed meson interactions in the hadronic phase.

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The D-meson mass modifications have been studied using the QCD sum rule approach (QSR) [15]. It was found that due to the presence of a light quark in the D-meson the mass modification of D-mesons has a large contribution from the light quark condensates [18]. The J/Ψ - being a $c\bar{c}$ vector state - has a dominant contribution from the gluon condensates. Accordingly, the substantial medium modification of the light quark condensate – as compared to the gluon condensate – is attributed to a larger drop of the D-meson mass as compared to the J/Ψ mass modification. Alternatively, within a linear density approximation in the QSR, the DN interaction is seen to be more attractive than the J/Ψ-N interaction [15]. This leads to a larger decrease of the D-meson mass than that of the charmonium (about 10 times larger at nuclear matter density ρ_0). It is seen that the large mass shift of the D-meson originates from the contribution of the $m_c\langle\bar{q}q\rangle_N$ term in the operator product expansion. In the QMC model, the contribution from the $m_c\langle\bar{q}q\rangle_N$ term is represented by a quark-σ 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 [20], which is very similar to the value obtained in the QCD sum rule calculations of Ref. [15, 18]. Furthermore, lattice calculations for heavy quark potentials at finite temperature suggest a similar drop [19, 29].

A reduction of the D-mass has direct consequences for the production of open charm [21] as well as J/Ψ suppression [30]. Recently, the NA50 collaboration [31] has reported a strong (so called 'anomalous') J/Ψ suppression [32] in Pb-Pb collision at 158 AGeV. A possible explanation of the J/Ψ suppression [15] is the large difference in the mass shifts for D and J/Ψ. As suggested by the p-A collision data [33], a large part of the observed J/Ψ is produced from the excited states χ_c and Ψ'. Thus an appreciable drop of the D-meson mass could lead to the decay of these excited states (χ_c or Ψ') in the medium to $D\bar{D}$ final states [34] and might lead to a lower yield for J/Ψ accordingly. These effects could be explored at the future accelerator facility at GSI [35].

In this work we study the medium modification of the masses of charmed mesons ($D^{\pm}$) due to their interaction with the light hadron sector. The medium modification of hadronic properties in hot and dense hyperonic matter has been studied in a chiral $SU(3)$-flavor model in Ref. [36]. We generalise the model to $SU(4)$-flavor to include charm mesons. Thus knowing the interactions of the charmed pesudoscalar mesons, we investigate their mass modification in the hot and dense matter due to their interactions with light hadrons. Within the model, both $D^{\pm}$ experience a drop in the thermal medium due to an attractive interaction via exchange of scalar mesons. The D^{+} has a drop in the medium due to a vectorial (as well as vector ω exchange) term, whereas the D^{-} has positive contributions from the vector terms in the thermal medium. The effects on the D-meson mass arising from terms of the type $(\partial_\mu D^+)(\partial^\mu D^-)$ are taken into consideration, too. We compare the results obtained in the model with those obtained from the interactions arising from chiral perturbation theory, where a drop of the $D^{\pm}$ mass is found due to an attractive nucleon scalar interaction (the so called sigma term) while the vectorial interaction is responsible for a drop (rise) for the D^{+} (D^{-}) mesons. The repulsive scalar interaction ($\sim (\partial_\mu D^+)(\partial^\mu D^-)$) which is of the same order as the attractive scalar interaction in chiral model, is also taken into consideration to study the in–medium D-meson mass. In the present model, the effect is seen to be larger than that arising from chiral perturbation theory. The effect of taking into account the nucleon Dirac sea for the study of hadronic properties, and hence their effect on D-mesons due to their interaction with the nucleons and scalar and vector mesons, is studied additionally. It is seen to give rise to higher values for the masses as compared to the mean field approximation, since such vacuum polarizations effects decrease the strength of the scalar and vector fields associated with a softening of the equation of state.

We organize the paper as follows: We briefly recapitulate the $SU(3)$-flavor chiral model adopted for the description of the hot and dense hadronic matter in Section 2. The hadronic 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 and vector mesons as presented in Section 3. Section 4 discusses the results of the present investigation, while we summarise our findings and discuss open questions in Section 5.

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

In this section the various terms of the effective Hadronic Lagrangian

$$\mathcal{L} = \mathcal{L}_{kin} + \sum_{W=X,Y,V,\mathcal{A},u} \mathcal{L}_{BW} + \mathcal{L}_{VP} + \mathcal{L}_{vec} + \mathcal{L}_0 + \mathcal{L}_{SB} \qquad (1)$$

are discussed. Eq. (1) corresponds to a relativistic quantum field theoretical model of baryons and mesons built on chiral symmetry and broken scale invariance [36, 37, 38] to describe strongly interacting nuclear matter. We adopt a nonlinear realization of the chiral symmetry which allows for a simultaneous description of hyperon potentials and properties of finite nuclei [37]. This Lagrangian contains the baryon octet, the spin-0 and spin-1 meson multiplets as the elementary degrees of freedom. In Eq. (1), $\mathcal{L}_{kin}$ 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. $\mathcal{L}_{VP}$ describes the interactions of vector mesons with the pseudoscalar mesons (and with photons). $\mathcal{L}_{vec}$ describes the dynamical mass generation of the vector mesons via couplings to the scalar mesons 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 as well as a scale invariance breaking logarithmic potential. $\mathcal{L}_{SB}$ describes the explicit symmetry breaking of $U(1)_A$, $SU(3)_V$ and the chiral symmetry.

2.0.1. The kinetic energy terms

An important property of the nonlinear realization of chiral symmetry is that all terms of the model-Lagrangian only have to be invariant under $SU(3)_V$ transformations in order to ensure chiral symmetry. This vector transformation depends in general on the pseudoscalar mesons and thus is local. Covariant derivatives have to be introduced for the kinetic terms in order to preserve chiral invariance [37]. The covariant derivative used in this case, reads: $D_\mu = \partial_\mu + [\Gamma_\mu, \ ]$ with $\Gamma_\mu = -\frac{i}{2}[u^\dagger \partial_\mu u + u \partial_\mu u^\dagger]$, where $u = \exp\left[\frac{i}{\sigma_0}\pi^a \lambda^a \gamma_5\right]$ is the unitary transformation operator. The pseudoscalar mesons are given as parameters of the symmetry transformation.
The kinetic energy terms read

$$\begin{aligned} \mathcal{L}_{kin} &= i\mathrm{Tr}\overline{B}\gamma_\mu D^\mu B + \frac{1}{2}\mathrm{Tr}D_\mu X D^\mu X + \mathrm{Tr}(u_\mu X u^\mu X + X u_\mu u^\mu X) + \frac{1}{2}\mathrm{Tr}D_\mu Y D^\mu Y \\ &+ \frac{1}{2}D_\mu \chi D^\mu \chi - \frac{1}{4}\mathrm{Tr}\left(\tilde{V}_{\mu\nu}\tilde{V}^{\mu\nu}\right) - \frac{1}{4}\mathrm{Tr}\left(F_{\mu\nu}F^{\mu\nu}\right) - \frac{1}{4}\mathrm{Tr}\left(\mathcal{A}_{\mu\nu}\mathcal{A}^{\mu\nu}\right) . \end{aligned} \qquad (2)$$

In (2) B is the baryon octet, X the scalar meson multiplet, Y the pseudoscalar chiral singlet, $\tilde{V}^\mu$ ($\mathcal{A}^\mu$) the renormalised vector (axial vector) meson multiplet with the field strength tensor $\tilde{V}_{\mu\nu} = \partial_\mu \tilde{V}_\nu - \partial_\nu \tilde{V}_\mu$ ($\mathcal{A}_{\mu\nu} = \partial_\mu \mathcal{A}_\nu - \partial_\nu \mathcal{A}_\mu$). $F_{\mu\nu}$ is the field strength tensor of the photon and χ is the scalar, iso-scalar dilaton (glueball) -field.

2.0.2. Baryon-meson interaction

Except for the difference in Lorentz indices, the SU(3) structure of the baryon -meson interaction terms are the same for all mesons. This interaction for a general meson field W has the form

$$\mathcal{L}_{BW} = -\sqrt{2} g_8^W \left(\alpha_W [\overline{B}\mathcal{O}BW]_F + (1-\alpha_W)[\overline{B}\mathcal{O}BW]_D\right) - g_1^W \frac{1}{\sqrt{3}} \mathrm{Tr}(\overline{B}\mathcal{O}B)\mathrm{Tr}W \,, \tag{3}$$

with $[\overline{B}\mathcal{O}BW]_F := \mathrm{Tr}(\overline{B}\mathcal{O}WB - \overline{B}\mathcal{O}BW)$ and $[\overline{B}\mathcal{O}BW]_D := \mathrm{Tr}(\overline{B}\mathcal{O}WB + \overline{B}\mathcal{O}BW) - \frac{2}{3}\mathrm{Tr}(\overline{B}\mathcal{O}B)\mathrm{Tr}W$. The different terms – to be considered – are those for the interaction of baryons with scalar mesons ($W = X, \mathcal{O} = 1$), with vector mesons ($W = \tilde{V}_\mu, \mathcal{O} = \gamma_\mu$ for the vector and $W = \tilde{V}_{\mu\nu}, \mathcal{O} = \sigma^{\mu\nu}$ for the tensor interaction), with axial vector mesons ($W = \mathcal{A}_\mu, \mathcal{O} = \gamma_\mu\gamma_5$) and with pseudoscalar mesons ($W = u_\mu, \mathcal{O} = \gamma_\mu\gamma_5$), respectively. For the current investigation the following interactions are relevant: Baryon-scalar meson interactions generate the baryon masses through coupling of the baryons to the non-strange $\sigma(\sim \langle \bar{u}u + \bar{d}d\rangle)$ and the strange $\zeta(\sim \langle \bar{s}s\rangle)$ scalar quark condensate. After insertion of the scalar meson matrix X, one obtains the baryon masses as

$$\begin{aligned}
m_N &= m_0 - \frac{1}{3} g_8^S (4\alpha_S - 1)(\sqrt{2}\zeta - \sigma) \\
m_\Lambda &= m_0 - \frac{2}{3} g_8^S (\alpha_S - 1)(\sqrt{2}\zeta - \sigma) \\
m_\Sigma &= m_0 + \frac{2}{3} g_8^S (\alpha_S - 1)(\sqrt{2}\zeta - \sigma) \\
m_\Xi &= m_0 + \frac{1}{3} g_8^S (2\alpha_S + 1)(\sqrt{2}\zeta - \sigma)
\end{aligned} \tag{4}$$

with $m_0 = g_1^S(\sqrt{2}\sigma + \zeta)/\sqrt{3}$. The parameters g_1^S, g_8^S and α_S can be used to fix the baryon masses to their experimentally measured vacuum values. It should be emphasised that the nucleon mass also depends on the *strange condensate* ζ. Recently, the vector meson properties were investigated in nuclear matter for the special case of $\alpha_S = 1$ and $g_1^S = \sqrt{6} g_8^S$ [38]. Then the nucleon mass depends only on the non–strange quark condensate. In the present investigation, the general case will be used to study hot and strange hadronic matter [36] and takes into account the baryon couplings to both scalar fields (σ and ζ) while summing over the baryonic tadpole diagrams to investigate the effect from the baryonic Dirac sea in the relativistic Hartree approximation [36].

In analogy to the baryon-scalar meson coupling there exist two independent baryon-vector meson interaction terms corresponding to the F-type (antisymmetric) and D-type (symmetric) couplings. Here we will use the symmetric coupling because – from the universality principle [39] and the vector meson dominance model – one can conclude that the antisymmetric coupling should be small. We realize it by setting $\alpha_V = 1$ for all fits. Additionally we decouple the strange vector field $\phi_\mu \sim \bar{s}\gamma_\mu s$ from the nucleon by setting $g_1^V = \sqrt{6} g_8^V$. The remaining baryon-vector meson interaction reads

$$\mathcal{L}_{BV} = -\sqrt{2} g_8^V \left\{ [\bar{B}\gamma_\mu B V^\mu]_F + \mathrm{Tr}\big(\bar{B}\gamma_\mu B\big)\mathrm{Tr}V^\mu \right\} . \tag{5}$$

2.0.3. Meson-meson interactions

The Lagrangian describing the interaction for the scalar mesons, X, and pseudoscalar singlet, Y, is given as [37]

$$\mathcal{L}_0 = -\frac{1}{2}k_0\chi^2 I_2 + k_1(I_2)^2 + k_2 I_4 + 2k_3\chi I_3, \tag{6}$$

with $I_2 = \mathrm{Tr}(X+iY)^2$, $I_3 = \det(X+iY)$ and $I_4 = \mathrm{Tr}(X+iY)^4$. In the above, χ is the scalar color singlet gluon field. It is introduced in order to satisfy the QCD trace anomaly, i.e. the nonvanishing energy-momentum tensor $\Theta^\mu_\mu = (\beta_{QCD}/2g)\langle G^a_{\mu\nu}G^{a,\mu\nu}\rangle$, where $G^a_{\mu\nu}$ is the gluon field tensor.
A scale breaking potential

$$\mathcal{L}_{\text{scalebreak}} = -\frac{1}{4}\chi^4 \ln\frac{\chi^4}{\chi_0^4} + \frac{\delta}{3}\chi^4 \ln\frac{I_3}{\det\langle X\rangle_0} \tag{7}$$

is introduced additionally and yields

$$\theta^\mu_\mu = 4\mathcal{L} - \chi\frac{\partial\mathcal{L}}{\partial\chi} - 2\partial_\mu\chi\frac{\partial\mathcal{L}}{\partial(\partial_\mu\chi)} = \chi^4. \tag{8}$$

It allows for the identification of the χ field width the gluon condensate $\Theta^\mu_\mu = (1-\delta)\theta^\mu_\mu = (1-\delta)\chi^4$. Finally the term $\mathcal{L}_\chi = -k_4\chi^4$ generates a phenomenologically consistent finite vacuum expectation value. We shall use the frozen glueball approximation i.e. assume $\chi = \langle 0|\chi|0\rangle \equiv \chi_0$, since the variation of χ in the medium is rather small [37].

The Lagrangian for the vector meson interaction is written as

$$\mathcal{L}_{vec} = \frac{m_V^2}{2}\frac{\chi^2}{\chi_0^2}\mathrm{Tr}\big(\tilde{V}_\mu\tilde{V}^\mu\big) + \frac{\mu}{4}\mathrm{Tr}\big(\tilde{V}_{\mu\nu}\tilde{V}^{\mu\nu}X^2\big) + \frac{\lambda_V}{12}\Big(\mathrm{Tr}\big(\tilde{V}^{\mu\nu}\big)\Big)^2 + 2(\tilde{g}_4)^4\mathrm{Tr}\big(\tilde{V}_\mu\tilde{V}^\mu\big)^2. \tag{9}$$

The vector meson fields, $\tilde{V}_\mu$ are related to the renormalized fields by $V_\mu = Z_V^{1/2}\tilde{V}_\mu$, with $V = \omega, \rho, \phi$. The masses of ω, ρ and ϕ are fitted from m_V, μ and λ_V.

2.0.4. Explicit chiral symmetry breaking

The explicit symmetry breaking term is given as [37]

$$\mathcal{L}_{SB} = \mathrm{Tr}A_p\left(u(X+iY)u + u^\dagger(X-iY)u^\dagger\right) \tag{10}$$

with $A_p = 1/\sqrt{2}\mathrm{diag}(m_\pi^2 f_\pi, m_\pi^2 f_\pi, 2m_K^2 f_K - m_\pi^2 f_\pi)$ and $m_\pi = 139$ MeV, $m_K = 498$ MeV. This choice for A_p, together with the constraints $\sigma_0 = -f_\pi$, $\zeta_0 = -\frac{1}{\sqrt{2}}(2f_K - f_\pi)$ on the VEV on the scalar condensates assure that the PCAC-relations of the pion and kaon are fulfilled. With $f_\pi = 93.3$ MeV and $f_K = 122$ MeV we obtain $|\sigma_0| = 93.3$ MeV and $|\zeta_0| = 106.56$ MeV.

2.1. Mean field approximation

We next proceed to study the hadronic properties in the chiral SU(3) model. The Lagrangian density in the mean field approximation is given as

$$\mathcal{L}_{BX} + \mathcal{L}_{BV} = -\sum_i \overline{\psi_i}\left[g_{i\omega}\gamma_0\omega + g_{i\phi}\gamma_0\phi + m_i^*\right]\psi_i \tag{11}$$

$$\mathcal{L}_{vec} = \frac{1}{2}m_\omega^2\frac{\chi^2}{\chi_0^2}\omega^2 + g_4^4\omega^4 + \frac{1}{2}m_\phi^2\frac{\chi^2}{\chi_0^2}\phi^2 + g_4^4\left(\frac{Z_\phi}{Z_\omega}\right)^2\phi^4 \tag{12}$$

$$\begin{aligned}\mathcal{V}_0 &= \frac{1}{2}k_0\chi^2(\sigma^2+\zeta^2) - k_1(\sigma^2+\zeta^2)^2 - k_2(\frac{\sigma^4}{2}+\zeta^4) - k_3\chi\sigma^2\zeta \\ &+ k_4\chi^4 + \frac{1}{4}\chi^4\ln\frac{\chi^4}{\chi_0^4} - \frac{\delta}{3}\chi^4\ln\frac{\sigma^2\zeta}{\sigma_0^2\zeta_0}\end{aligned} \tag{13}$$

$$\mathcal{V}_{SB} = \left(\frac{\chi}{\chi_0}\right)^2\left[m_\pi^2 f_\pi\sigma + (\sqrt{2}m_K^2 f_K - \frac{1}{\sqrt{2}}m_\pi^2 f_\pi)\zeta\right], \tag{14}$$

where $m_i^* = -g_{\sigma i}\sigma - g_{\zeta i}\zeta$ is the effective mass of the baryon of type i ($i = N, \mathbf{\Sigma}, \Lambda, \Xi$). In the above, $g_4 = \sqrt{Z_\omega}\tilde{g}_4$ is the renormalised coupling for ω-field. The thermodynamical potential of the grand canonical ensemble, Ω, per unit volume V at given chemical potential μ and temperature T can be written as

$$\begin{aligned}\frac{\Omega}{V} &= -\mathcal{L}_{vec} - \mathcal{L}_0 - \mathcal{L}_{SB} - \mathcal{V}_{vac} + \sum_i \frac{\gamma_i}{(2\pi)^3}\int d^3k\, E_i^*(k)\Big(f_i(k) + \bar{f}_i(k)\Big) \\ &- \sum_i \frac{\gamma_i}{(2\pi)^3}\,\mu_i^* \int d^3k\,\Big(f_i(k) - \bar{f}_i(k)\Big).\end{aligned} \tag{15}$$

Here the vacuum energy (the potential at $\rho = 0$) has been subtracted in order to get a vanishing vacuum energy. In (16) γ_i are the spin-isospin degeneracy factors. The f_i and $\bar{f}_i$ are thermal distribution functions for the baryon of species, i given in terms of the effective single particle energy, E_i^*, and chemical potential, μ_i^*, as

$$f_i(k) = \frac{1}{\mathrm{e}^{\beta(E_i^*(k)-\mu_i^*)}+1} \quad , \quad \bar{f}_i(k) = \frac{1}{\mathrm{e}^{\beta(E_i^*(k)+\mu_i^*)}+1}, \tag{16}$$

with $E_i^*(k) = \sqrt{k_i^2 + {m_i^*}^2}$ and $\mu_i^* = \mu_i - g_{i\omega}\omega$. The mesonic field equations are determined by minimizing the thermodynamical potential [36, 38]. These are expressed in terms of the scalar and vector densities for the baryons at finite temperature

$$\rho_i^s = \gamma_i \int \frac{d^3k}{(2\pi)^3}\,\frac{m_i^*}{E_i^*}\,\Big(f_i(k) + \bar{f}_i(k)\Big)\ ;\ \ \rho_i = \gamma_i \int \frac{d^3k}{(2\pi)^3}\,\Big(f_i(k) - \bar{f}_i(k)\Big)\ . \tag{17}$$

The energy density and the pressure are given as, $\epsilon = \Omega/V + \mu_i\rho_i$+TS and $p = -\Omega/V$.

2.2. Relativistic Hartree approximation

The relativistic Hartree approximation takes into account the effects from the Dirac sea by summing over the baryonic tadpole diagrams and the interacting propagator for a baryon of type i has the form [40]

$$\begin{aligned}G_i^H(p) &= (\gamma^\mu\bar{p}_\mu + m_i^*)\Bigg[\frac{1}{\bar{p}^2 - {m_i^*}^2 + i\epsilon} \\ &+ \frac{\pi i}{E_i^*(p)}\left\{\frac{\delta(\bar{p}^0 - E_i^*(p))}{\mathrm{e}^{\beta(E_i^*(p)-\mu_i^*)}+1} + \frac{\delta(\bar{p}^0 + E_i^*(p))}{\mathrm{e}^{\beta(E_i^*(p)+\mu_i^*)}+1}\right\}\Bigg] \\ &\equiv G_i^F(p) + G_i^D(p),\end{aligned} \tag{18}$$

where $E_i^*(p) = \sqrt{\mathbf{p}^2 + {m_i^*}^2}$, $\bar{p} = p + \Sigma_i^V$ and $m_i^* = m_i + \Sigma_i^S$. Σ_i^V and Σ_i^S are the vector and scalar self energies of baryon, i respectively. In the present investigation (for the study of hot hyperonic matter) the baryons couple to both the non-strange (σ) and strange (ζ) scalar fields, so that we have

$$\Sigma_i^S = -(g_{\sigma i}\tilde{\sigma} + g_{\zeta i}\tilde{\zeta})\,, \tag{19}$$

where $\tilde{\sigma} = \sigma - \sigma_0$, $\tilde{\zeta} = \zeta - \zeta_0$. The scalar self-energy Σ_i^S can be written

$$\Sigma_i^S = i\left(\frac{g_{\sigma i}^2}{m_\sigma^2} + \frac{g_{\zeta i}^2}{m_\zeta^2}\right)\int \frac{\mathrm{d}^4\mathrm{p}}{(2\pi)^4}\mathrm{Tr}\big[G_i^F(p) + G_i^D(p)\big]e^{ip^0\eta} \equiv \left(\Sigma_i^S\right)^F + \left(\Sigma_i^S\right)^D. \tag{20}$$

$(\Sigma_i^S)^D$ is the density dependent part and is identical to the mean field contribution

$$\left(\Sigma_i^S\right)^D = -\left(\frac{g_{\sigma i}^2}{m_\sigma^2} + \frac{g_{\zeta i}^2}{m_\zeta^2}\right)\rho_i^s, \tag{21}$$

with ρ_i^s as defined in (17). The Feynman part $(\Sigma_i^S)^F$ of the scalar part of the self-energy is divergent. We carry out a dimensional regularization to extract the convergent part. Adding the counter terms [36]

$$\left(\Sigma_i^S\right)_{CTC} = -\left(\frac{g_{\sigma i}^2}{m_\sigma^2} + \frac{g_{\zeta i}^2}{m_\zeta^2}\right)\sum_{n=0}^{3}\frac{1}{n!}(g_{\sigma i}\tilde{\sigma} + g_{\zeta i}\tilde{\zeta})^n\beta_{n+1}^i\,, \tag{22}$$

yields the additional contribution from the Dirac sea to the baryon self energy [36]. The field equations for the scalar meson fields are then modified to

$$\left.\frac{\partial(\Omega/V)}{\partial\Phi}\right|_{RHA} = \left.\frac{\partial(\Omega/V)}{\partial\Phi}\right|_{MFT} + \sum_i \frac{\partial m_i^*}{\partial\Phi}\Delta\rho_i^s = 0 \quad \text{with} \quad \Phi = \sigma, \zeta\,, \tag{23}$$

where the additional contribution to the nucleon scalar density is given as [36]

$$\Delta\rho_i^s = -\frac{\gamma_i}{4\pi^2}\left[{m_i^*}^3\ln\left(\frac{m_i^*}{m_i}\right) + m_i^2(m_i - m_i^*) - \frac{5}{2}m_i(m_i - m_i^*)^2 + \frac{11}{6}(m_i - m_i^*)^3\right]. \tag{24}$$

3. *D*-MESON MASS MODIFICATION IN THE MEDIUM

We now examine the medium modification for the D-meson mass in the hot and dense hadronic matter. In the last section, the SU(3) chiral model was used to study the hadronic properties in the medium within the relativistic Hartree approximation. We assume that the additional effect of charmed particles in the medium leads to only marginal modifications [41] 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, σ and strange, ζ) as well as a vectorial interaction and a ω- exchange 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 thermal medium [44]. The scalar meson exchange gives an attractive interaction leading to a drop of the D (K) -meson masses similar to a scalar sigma term in the chiral perturbation theory [45]. In fact, the sigma term corresponding to the scalar kaon-nucleon [44] (as well as πN sigma term) and D-nucleon attractive interaction are predicted in our approach automatically by using SU(3) and SU(4) symmetry, respectively. The pion-nucleon and kaon-nucleon sigma terms as calculated from the scalar meson exchange interaction of our Lagrangian are around 28 MeV and 463 MeV respectively. The value for KN sigma term calculated in our model is close to the value of Σ_{KN}=450 MeV found by lattice gauge calculations of [49]. The value for the DN sigma term within our chiral model turns out to be $\Sigma_{DN} = 7366$ MeV ignoring the contribution of the charm condensate to the nuclear scalar density. In the present investigation we also consider the effect of repulsive scalar contributions ($\sim (\partial_\mu D^+)(\partial^\mu D^-)$) which contribute in the same order as the attractive sigma term in chiral perturbation theory.

To consider the medium effect on the 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

$$\langle X \rangle = \begin{pmatrix} \sigma/\sqrt{2} & 0 & 0 & 0 \\ 0 & \sigma/\sqrt{2} & 0 & 0 \\ 0 & 0 & \zeta & 0 \\ 0 & 0 & 0 & \zeta_c \end{pmatrix}, \tag{25}$$

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

$$P = \begin{pmatrix} \pi^0/\sqrt{2} & \pi^+ & \frac{2K^+}{1+w} & 0 \\ \pi^- & -\pi^0/\sqrt{2} & 0 & \frac{2D^-}{1+w_c} \\ \frac{2K^-}{1+w} & 0 & 0 & 0 \\ 0 & \frac{2D^+}{1+w_c} & 0 & 0 \end{pmatrix}, \tag{26}$$

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 [18]. We note that for the decay constant of D_s^+, the Particle Data Group [42] quotes a value of $f_{D_s^+} \simeq 200 MeV$. Taking a similar value also for f_D would not affect our results qualitatively, however (see also [43]). The vector meson interaction with the pseudoscalar mesons, which modifies the masses of the K(D) mesons, is given as [44]

$$\mathcal{L}_{VP} = -\frac{m_V^2}{2g_V}\mathrm{Tr}(\Gamma_\mu V^\mu) + h.c. \tag{27}$$

The vector meson multiplet is given as $V = \mathrm{diag}\big((\omega+\rho_0)/\sqrt{2},\ (\omega-\rho_0)/\sqrt{2},\ \phi,\ J/\Psi\big)$. The non-diagonal components in the multiplet have not been written down as they are not relevant for the present investigation. With the interaction (27), the coupling of the D-meson to the ω-meson is related to the pion-rho coupling as $g_{\omega D}/g_{\rho\pi\pi} = f_\pi^2/(2f_D^2)$.

The scalar meson exchange interaction term, which is attractive for the D-mesons, is given from the explicit symmetry breaking term by equation (10), where $A_p = 1/\sqrt{2}$ diag $(m_\pi^2 f_\pi,\ m_\pi^2 f_\pi,\ 2\ m_K^2 f_K - m_\pi^2 f_\pi,\ 2\ m_D^2 f_D - m_\pi^2 f_\pi)$.

The interaction Lagrangian modifing the D-meson mass can be written as [44]

$$\begin{aligned} \mathcal{L}_{\mathcal{D}} &= -\frac{3i}{8f_D^2}\bar{N}\gamma^\mu N(D^+\partial_\mu D^- - \partial_\mu D^+ D^-) \\ &+ \frac{m_D^2}{2f_D}(\sigma + \sqrt{2}\zeta_c)D^+D^- - ig_{\omega D}(D^+\partial_\mu D^- - \partial_\mu D^+ D^-)\omega^\mu \\ &- \frac{1}{f_D}(\sigma + \sqrt{2}\zeta_c)(\partial_\mu D^+)(\partial^\mu D^-) + \frac{d_1}{2f_D^2}(\bar{N}N)(\partial_\mu D^+)(\partial^\mu D^-). \end{aligned} \tag{28}$$

In (28) the first term is the vectorial interaction term obtained from the first term in (2). The second term, which gives an attractive interaction for the D-mesons, is obtained from the explicit symmetry breaking term (10). The third term, referring to the interaction in terms of ω-meson exchange, is attractive for the D^+ and repulsive for D^-. The fourth term arises within the present chiral model from the kinetic term of the pseudoscalar mesons given by the third term in equation (2), when the scalar fields in one of the meson multiplets, X are replaced by their vacuum expectation values. The fifth term in (28) has been written down for the DN interactions, analogous to a term of the type

$$\mathcal{L}_{\tilde{D}}^{BM} = d_1 Tr(u_\mu u^\mu \bar{B} B), \tag{29}$$

in the SU(3) chiral model. The last two terms in (28) represent the range term in the chiral model.

3.1. Fitting to KN scattering data

The term (29) reduces for the KN interactions to

$$\mathcal{L}_{\tilde{D}}^{KN} = d_1 \frac{1}{2f_K^2} (\bar{N}N)(\partial_\mu K^+)(\partial^\mu K^-). \tag{30}$$

The coefficient, d_1 in the above shall be determined as consistent with the KN scattering data [46, 47, 48]. The isospin averaged KN scattering length

$$\bar{a}_{KN} = \frac{1}{4}(3a_{KN}^{I=1} + a_{KN}^{I=0}) \tag{31}$$

can be calculated to be

$$\begin{aligned} \bar{a}_{KN} &= \frac{m_K}{4\pi(1+m_K/m_N)} \Big[-\big(\frac{m_K}{2f_K}\big) \cdot \frac{g_{\sigma N}}{m_\sigma^2} - \big(\frac{\sqrt{2}m_K}{2f_K}\big) \cdot \frac{g_{\zeta N}}{m_\zeta^2} \\ &- \frac{2 g_{\omega K} g_{\omega N}}{m_\omega^2} - \frac{3}{4f_K^2} + \frac{d_1 m_K}{2f_K^2} \Big]. \end{aligned} \tag{32}$$

Choosing the empirical value of the isospin averaged scattering length [46, 48],

$$\bar{a}_{KN} \approx -0.255 \text{ fm} \tag{33}$$

determines the value for the coefficient d_1. The present calculations corresponds to the values $g_{\sigma N} = 10.618$, and $g_{\zeta N} = -0.4836$ as consistent with the vacuum baryon masses, and the other parameters as fitted from the nuclear matter saturation properties as listed in Ref. [36]. We consider the case when a quartic vector interaction is present. The coefficient d_1 is evaluated in the mean field and RHA cases as $5.63/m_K$ and $4.33/m_K$ respectively. The contribution from this term is thus seen to be attractive, contrary to the other term proportional to $(\partial_\mu D^+)(\partial^\mu D^-)$ in (28) which is repulsive.

3.2. Chiral perturbation theory

The effective Lagrangian obtained from chiral perturbation theory [45] has been used extensively in the literature for the study of kaons in dense matter. This has a vector interaction (called the Tomozawa-Weinberg term) as the leading term. At sub-leading order there are the attractive scalar nucleon interaction term (the sigma term) [45] as well as the repulsive scalar contribution (proportional to the kinetic term of the pesudoscalar meson). We generalize such an interaction to SU(4) to write down the interaction of a D-meson with a nucleon as

$$\begin{aligned}\mathcal{L}_{DN} &= -\frac{3i}{8f_D^2}\bar{N}\gamma^\mu N(D^+\partial_\mu D^- - \partial_\mu D^+ D^-) + \frac{\Sigma_{DN}}{f_D^2}(\bar{N}N)D^+D^- \\ &+ \frac{\tilde{D}}{f_D^2}(\bar{N}N)(\partial_\mu D^+)(\partial^\mu D^-).\end{aligned} \tag{34}$$

where $\Sigma_{DN} = \frac{m_d+m_c}{2}\langle N|(\bar{d}d+\bar{c}c)|N\rangle$ in analogy to the definition of $\Sigma_{KN} = \frac{\bar{m}+m_s}{2}\langle N|(\bar{u}u+\bar{s}s)|N\rangle$ [49]. Neglecting the charm condensate inside the nucleon, this is directly related to the pion-nucleon sigma term given as $\Sigma_{\pi N} = \bar{m}\langle N|(\bar{u}u+\bar{d}d)|N\rangle$, with $\langle N|\bar{u}u|N\rangle = \langle N|\bar{d}d|N\rangle$. In our calculations, we take $m_c = 1.3\ GeV$, $\bar{m} = (m_u + m_d)/2 = 7MeV$, and $\Sigma_{\pi N} = 45MeV$, which gives the value for $\Sigma_{DN} = 2089\ MeV$.

The last term of the Lagrangian (34) is repulsive and is of the same order as the attractive sigma term. This, to a large extent, compensates the scalar attraction due to the scalar Σ- term. We fix the coefficient $\tilde{D}$ from the KN scattering data [46]. This involves choosing a value for Σ_{KN}, which depends on the strange condensate content of the nucleon. Its value has, however, a large uncertainty. We consider the two extreme choices: $\Sigma_{KN} = 2m_\pi$ and $\Sigma_{KN} = 450$ MeV. The coefficient, $\tilde{D}$ as fitted to the empirical value of the KN scattering length (33) is:

$$\tilde{D} \approx 0.33/m_K - \Sigma_{KN}/m_K^2. \tag{35}$$

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 that using the interaction Lagrangian of chiral perturbation theory as well as from other approaches.

4. RESULTS AND DISCUSSIONS

To study the D-meson masses in hot and dense hadronic medium due to its interactions with the light hadrons, we have generalized the chiral SU(3) model used for the study of the hot and strange hadronic matter to SU(4) for the meson sector. The contributions from the various terms of the interaction Lagrangian (28) are shown in Fig. 1 in mean field approximation. The vector interaction (A) as well as the ω exchange (C) terms (given by the first and the third terms of equation (28), respectively) lead to a drop for the D^+ mass, whereas they are repulsive for the D^-. The scalar meson exchange term (B) is attractive for both D^+ and D^-. The first term of the range term (referred to as (D)) of eq. (28) is repulsive whereas the second term has an attractive contribution. This results in a turn over of the D-mass at around 0.4 ρ_0 above which the last term in (28) (attractive) dominates. The dominant contributions arise from the scalar exchange (B) and the term (D) (dominated by d_1 term at higher densities), which lead to a substantial drop of D meson mass in the medium. The vector terms (A) and (C) lead to a further drop of D^+ mass, whereas for D^- they compete with the contributions from the other two contributions. The effect from the nucleon Dirac sea on the mass modification of the D-mesons is shown in Fig. 2. This gives rise to smaller modifications as compared to the mean field calculations though qualitative features remain the same.

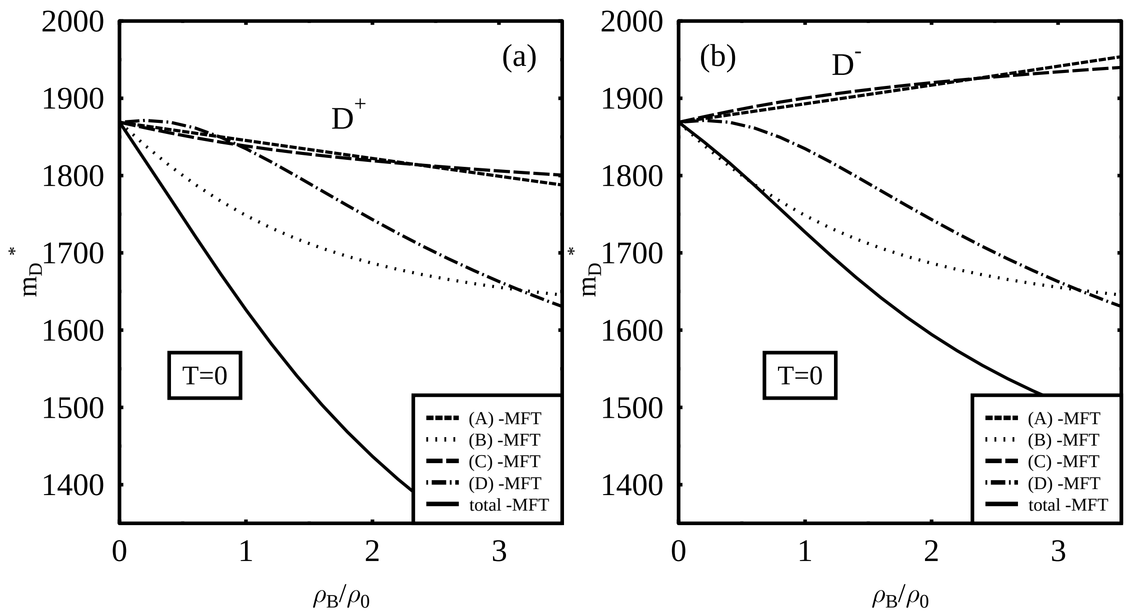

FIG. 1: Contributions to the masses of D$^{\pm}$ mesons due to the various interactions in the present chiral model in the mean field approximation. The curves refer to individual contributions from (A) the vectorial interaction, (B) scalar exchange, (C) ω exchange, (D) the last two terms of the equation (28). The solid line refers to the total contribution.

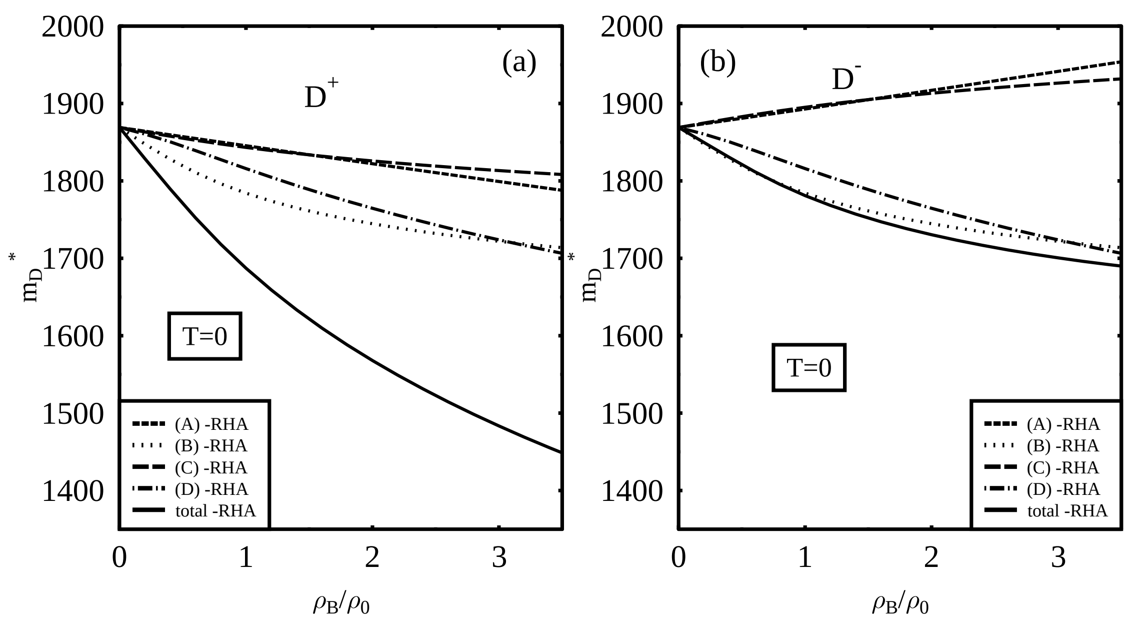

FIG. 2: Same as in figure (1), but in the relativistic Hartree approximation. The contributions to the masses of $D^{\pm}$ mesons due to the various interactions are seen to be smaller when the Dirac sea effects are taken into account.

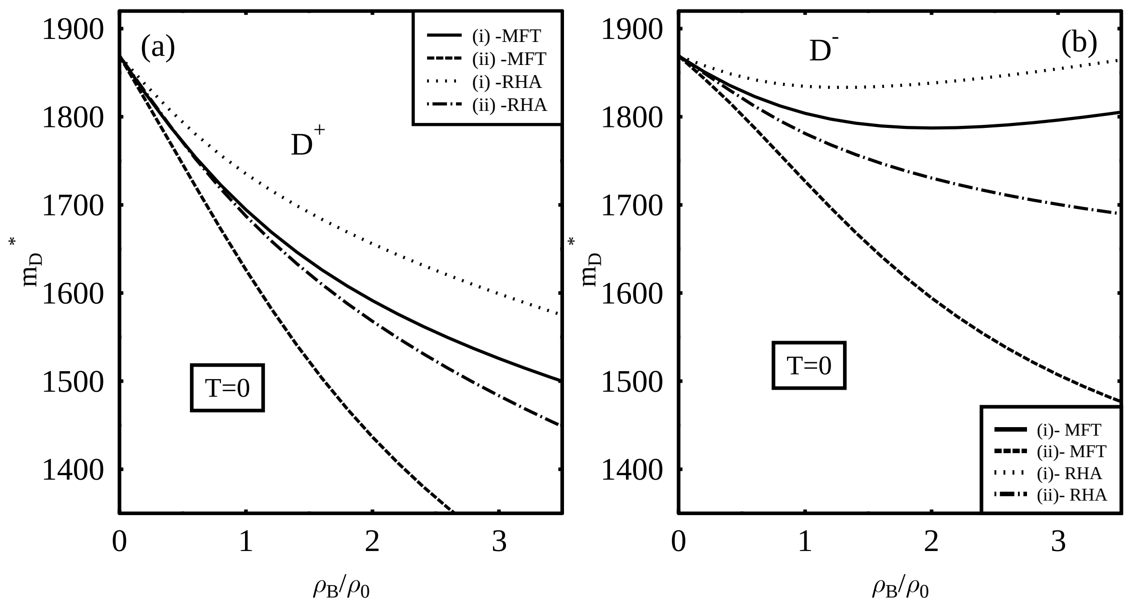

FIG. 3: Masses of $D^{\pm}$ mesons due to the interactions in the chiral effective model at $T = 0$.

In Fig. 3 the masses of the D-mesons are plotted for $T = 0$ in the present chiral model. We first consider the situation (case (i)) when the Weinberg-Tomozawa term is supplemented by the scalar and vector meson exchange interactions [44]. The other case (ii) corresponds to the inclusion of the last two terms in (28) with the parameter d_1 determined from the empirical value of the scattering length (33). There is seen to be a substantial drop of D-meson masses due to the inclusion of the range term.

In Fig. 4, the masses are plotted for $T = 150$ MeV. One might note here, that the drop is smaller as compared to at $T = 0$ at finite densities. This is due to the fact that the nucleon mass increases with temperature at finite densities [7, 36]. Such a behaviour of the nucleon mass with temperature was also observed earlier within the Walecka model by Ko and Li [7] in a mean field calculation. This subtle behaviour of the baryon self energy, given by (21) in the mean field approximation, can be understood in the following manner: The scalar self energy (21) increases due to the thermal distribution functions at finite temperatures, whereas at higher temperatures there are also contributions from higher momenta which lead to lower values of self energy. These competing effects give rise to the observed behaviour of the effective baryon masses with temperature at finite densities. This increase in the nucleon mass with temperature is also reflected in the vector meson (ω, ρ and ϕ) masses in the medium [36]. However at zero density, due to effects arising only from the thermal distribution functions, the masses are seen to drop with temperature.

We compare the results obtained in the present model to those using the interaction Lagrangian of chiral perturbation theory from Ref. [45]. The masses obtained from chiral perturbation theory are plotted in Figs. 5 and 6 for $T = 0$ and 150 MeV. The $D^{\pm}$ masses are plotted for the cases: (I) in the absence of the $(\partial_\mu D^+)(\partial^\mu D^-)$, (II) $\tilde{D}$

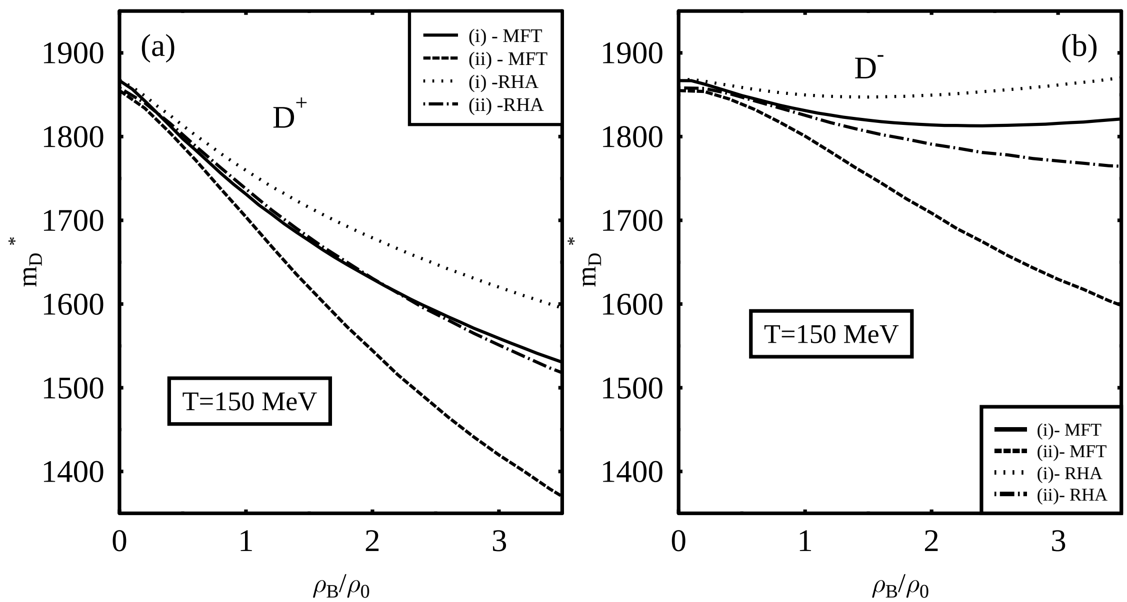

FIG. 4: Masses of $D^{\pm}$ mesons in the chiral effective model for $T = 150$ MeV.

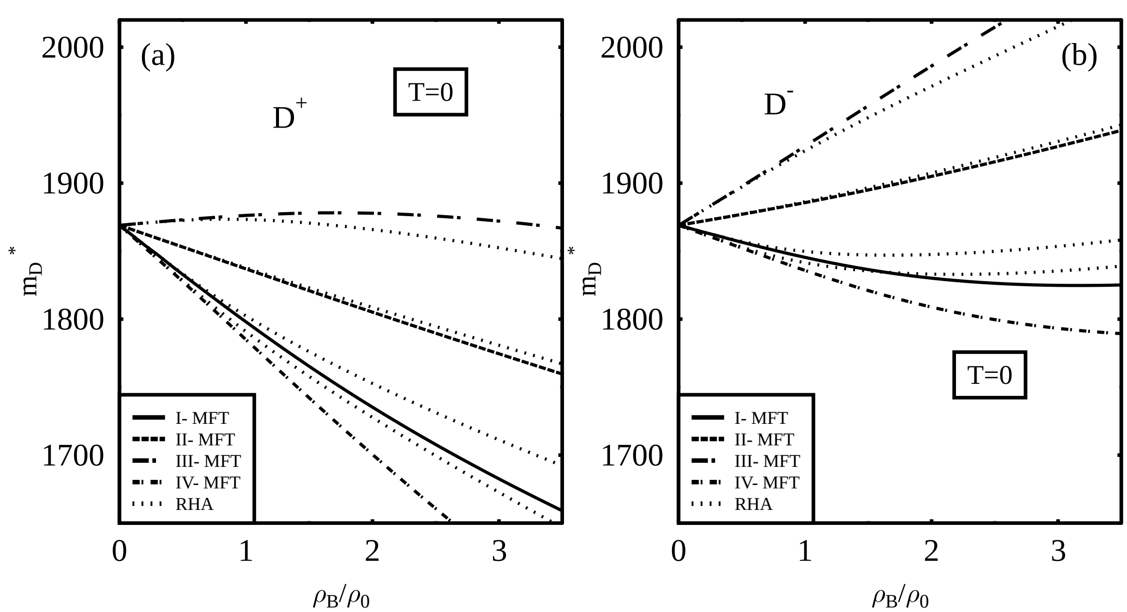

FIG. 5: Masses of $D^{\pm}$ mesons at $T = 0$ due to the interactions of chiral perturbation theory (see text for details).

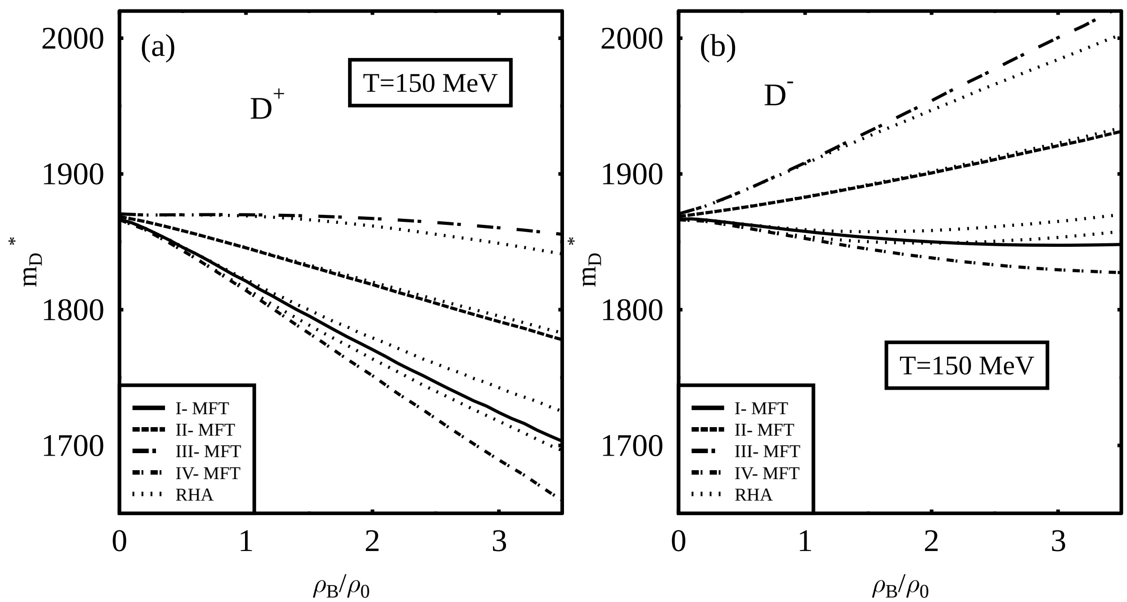

FIG. 6: Masses of $D^{\pm}$ mesons due to the interactions of chiral perturbation theory for $T = 150$ MeV (see text for details).

corresponding to $\Sigma_{KN} = 2m_\pi$, that is when the strangeness content of the nucleon is zero, (III) $\tilde{D}$ corresponding to $\Sigma_{KN} = 450$ MeV, (IV) $\tilde{D}$ corresponding to $\Sigma_{KN} = 450$ MeV and Σ_{DN} =7366 MeV, as calculated from the effective chiral model (28). The case (I) shows stronger drop of the D^+ mass in the medium as compared to (II) and (III) due to the exclusion of the scalar repulsion term. For D^- however there are cancelling effects from the sigma term and the Weinberg-Tomozawa interactions leading to only moderate mass modification. The repulsive term for (III) corresponding to the larger value of $\tilde{D}$ has a higher contribution as compared to (II) as expected. However, using the value for Σ_{DN} as predicted form the chiral effective model, leads to a larger drop for the D-meson masses from the sigma term as shown in curve (IV). The relativistic Hartree approximation gives rise to smaller mass modification as compared to the mean field case. In the chiral effective model (28) however the D^+-meson mass has a stronger drop as compared to the chiral perturbation theory. The value of d_1 in (28) as fitted from low energy KN scattering data has a higher value to overcome the repulsive ω- exchange term, which is absent in chiral perturbation theory. This leads to the range term in the present chiral model to be attractive of the same in chiral perturbation theory. As a result, the mass of the D^+D^- pair experiences a larger drop in the medium which as we shall see leads to the decay of charmonium states to such a pair.

When the vectorial interaction – supplemented by the scalar sigma term – is considered (case I), i.e., ignoring the repulsive terms proportional to $(\partial_\mu D)(\partial^\mu D)$, one obtains mass drops for $D^{\pm}$ to be around 67 MeV and 19 MeV at nuclear matter saturation density. These values are similar to those obtained in the QCD sum rule calculations of Ref. [18]. A similar drop of the D-meson mass is also predicted by the QMC model [20]. We stress that the present

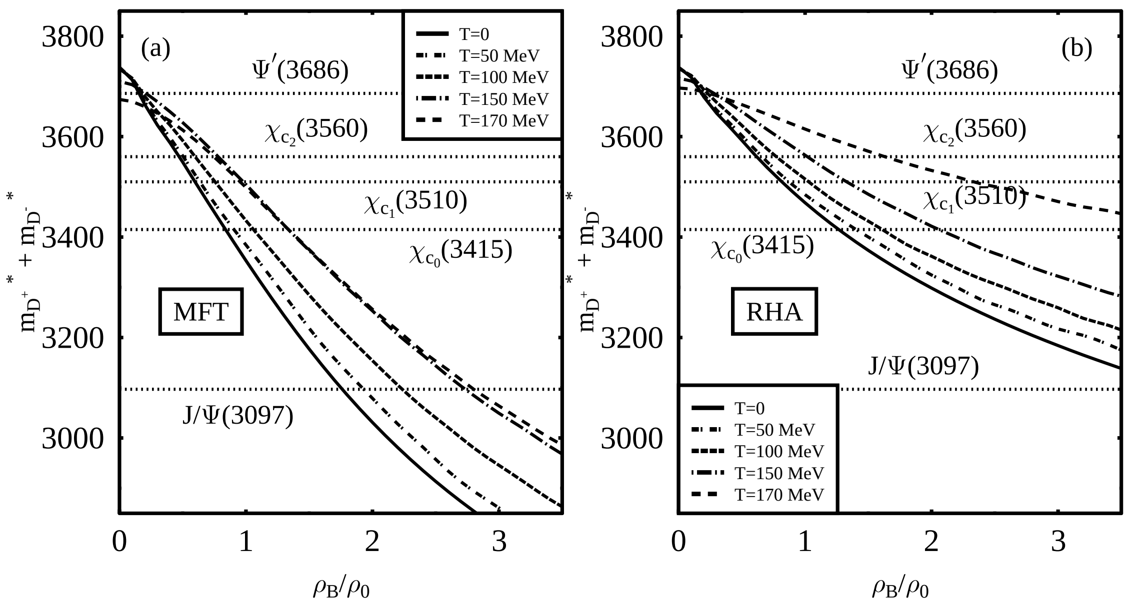

FIG. 7: The sum of the masses for $D^{\pm}$-mesons plotted versus density for different temperatures. The masses for the Ψ', χ_c and J/Ψ are also shown to indicate the threshold conditions for the decay of these quarkonium states into $D\bar{D}$ pairs. The relativistic Hartree approximation for a given temperature is seen to give rise to a higher threshold for the density as compared to the mean field calculations.

model gives stronger modifications for the D-meson masses than chiral perturbation theory.

It is interesting to compare the behaviour of D and K meson masses in a medium. The masses of K^- as well as D^+ drop in the medium. For the kaons, the vector interaction in the chiral perturbation theory is the leading contribution giving rise to a drop (increase) of the mass of K^- (K^+). The subleading contributions arise from the sigma and range terms with their coefficients as fitted to the KN scattering data [47]. Ignoring the charm condensate contribution in the nucleon in the chiral perturbation theory, the D^- mass also increases in the medium as in the cases II and III (see figures 5 and 6) similar to the behaviour of K^+. However, choosing the value for Σ_{DN} as calculated in the present chiral effective model, the mass of D^- drops in the medium (case IV). In the chiral effective model, the scalar exchange term as well as the range term (which turns attractive for densities above $0.4\rho_0$) lead to the drop of both D^+ and D^- masses in the medium. We note here that a similar behaviour is also obtained for the K^+. Firstly, it increases up to around a density of $0.8\rho_0$ and then drops due to the range term becoming attractive at higher densitites. However, though the qualitative features remain the same, the medium modification for K^+ is much less pronounced as compared to that of the D^+ in the medium. The medium modifications for the kaons and D-mesons in either model are obtained as consistent with the low energy KN scattering data.

The decay widths of the charmonium states can be modified by the level crossings between the excited states of J/Ψ (i.e., Ψ',χ_c) and the threshold for $D\bar{D}$ creation due to the medium modifications of the D-meson masses [17]. In the vacuum, the resonances above the $D\bar{D}$ threshold, for example the Ψ'' 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 $D^{\pm}$-meson masses, the channels for the excited states of J/Ψ, like χ_c, Ψ' decaying to D^+D^- pairs can open up at finite temperatures and densities. We show the temperature and density dependence of the mass of a D^+D^- pair in Fig. (7) for the chiral effective Lagrangian. The masses of the J/Ψ as well as its excited states are also shown to indicate the threshold values of temperature and density when their decay to D^+D^- pairs becomes accessible. Indeed we find that due to the substantial medium modification of the D-meson (especially of D^+ meson) that charmonium states will dissociate already at quite moderate densities. There is a strong drop of the mass of D^+D^- pair in the mean field calculations, which can lead to J/Ψ decaying to a pair of D-mesons at around 2-3 times the nuclear matter density. The relativistic Hartree approximation is seen to give rise to higher threshold values in density. At zero density, one sees that the Ψ' decaying to $D\bar{D}$ becomes accessible at a temperature of around 160 MeV which is around the chiral phase transition. Studies of charmonium dissociation using heavy quark potential inferred from lattice data, however, predict smaller values for the dissociation of Ψ' of around 0.1-0.2 T_c in [19] and 0.5 T_c in [50]. In our present model, χ_c remains stable at zero density, even up to a temperature of 170 MeV as compared to dissociation temperatures of χ_c of around 0.74 T_c in ref. [19] and around 0.9 T_c in [50]. The density modifications of the D-meson masses are seen to be large whereas the mass modifications are seen to be rather insensitive to temperature.

After the level crossings one would naively expect that the decay widths of Ψ' and χ_c states will increase drastically with density. The decay of charmonium states to $D\bar{D}$ has been studied in Ref. [17, 34]. It is seen to depend sensitively on the relative momentum in the final state. These excited states might become narrow [17] though the D-meson mass is decreased appreciably at high temperatures and densities. It may even vanish at certain momenta corresponding to nodes in the wavefunction [17]. Though the decay widths for these excited states can be modified by their wave functions, the partial decay width of χ_{c2}, due to absence of any nodes, can increase monotonically with the drop of the D^+D^- pair mass in the medium [17]. This can give rise to depletion in the J/Ψ yield in heavy ion collisions. The dissociation of the quarkonium states (Ψ', χ_c, J/Ψ) into $D\bar{D}$ pairs have also been studied [19, 50] by comparing their binding energies with lattice results on temperature dependence of heavy quark effective potential [29]. The dissociation occurs since the open charm mass drops faster with temperature than the mass of the excited charmonium [19]. The medium effects on the charmonium masses have recently been studied in a perturbative QCD approach [51] which shows an appreciable drop of Ψ' in nuclear matter. Accounting for the mass modifications of the charmonia will change the threshold conditions for the decay of these states to $D\bar{D}$ pairs, and, in turn, modify the J/Ψ yield in the heavy ion collision experiments. It can also have observable consequences in the dilepton spectra in the $\bar{p}A$ annihilation experiments [51] in the future GSI facility [35].

5. SUMMARY

To summarize we have investigated in a chiral model the temperature and density dependence of the $D, \bar{D}$-meson masses arising from the interactions with the nucleons and scalar and vector mesons. The properties of the light hadrons – as studied in a $SU(3)$ chiral model – modify the D-meson properties in the hot and dense hadronic medium. The SU(3) model with paramaters fixed from the properties of hadron masses, nuclei and hypernuclei and KN scattering data, is extended in a controlled fashion to SU(4) taking into account all terms up to the next to leading order arising in chiral perturbative expansion to derive the interactions of D-mesons with the light hadron sector. The important advantage of the present approach is that the DN, KN as well as πN sigma terms are calculated within the model. The model predictions for the πN and KN sigma terms are reasonable, the value for Σ_{KN} from the model being in agreement with lattice gauge calculations. Using the Lagrangian from chiral perturbation theory with a vectorial Tomozawa-Weinberg interaction, supplemented by an attractive scalar interaction (the sigma term) for the DN interactions, the results obtained are seen to be similar to earlier finite density calculations of QCD sum rules [18] as well as to the quark meson coupling (QMC) model [20]. However, the presence of the repulsive range term, given by the last term in (34) reduces the drop in mass. The chiral effective model, which is adjusted to describe nuclear properties, dominantly due to the scalar exchange and the attractive range term gives a larger drop of the D-meson masses at finite density as compared to chiral perturbation theory, The effect of the baryon Dirac sea for the hot hyperonic matter using the chiral model gives a higher value for the D-meson masses as compared to the mean field calculations. The medium modification of the D-masses can lead to a suppression in the J/Ψ- yield in heavy-ion collisions. In MFT, we find that J/Ψ dissociates into $D\bar{D}$ pairs already at 2-3 ρ_0. The relativistic Hartree approximation gives rise to somewhat higher values for the threshold densities for the quarkonium decaying to D^+D^- pairs. The density dependence of D-mass is seen to be the dominant medium effect as compared to the temperature dependence. The strong density dependence of the D-meson optical potential can be tested with the existing data on J/Ψ production from SPS experiments as well as by the future GSI facility [35] which is currently under investigation.

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