Hindawi Publishing Corporation
Advances in High Energy Physics
Volume 2015, Article ID 697514, 16 pages
http://dx.doi.org/10.1155/2015/697514

Research Article

D_S Mesons in Asymmetric Hot and Dense Hadronic Matter

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Received 8 June 2015; Revised 9 August 2015; Accepted 18 August 2015

Academic Editor: Juan José Sanz-Cillero

The in-medium properties of D_S mesons are investigated within the framework of an effective hadronic model, which is a generalization of a chiral $SU(3)$ model, to $SU(4)$, in order to study the interactions of the charmed hadrons. In the present work, the D_S mesons are observed to experience net attractive interactions in a dense hadronic medium, hence reducing the masses of the D_S^+ and D_S^- mesons from the vacuum values. While this conclusion holds in both nuclear and hyperonic media, the magnitude of the mass drop is observed to intensify with the inclusion of strangeness in the medium. Additionally, in hyperonic medium, the mass degeneracy of the D_S mesons is observed to be broken, due to opposite signs of the Weinberg-Tomozawa interaction term in the Lagrangian density. Along with the magnitude of the mass drops, the mass splitting between D_S^+ and D_S^- mesons is also observed to grow with an increase in baryonic density and strangeness content of the medium. However, all medium effects analyzed are found to be weakly dependent on isospin asymmetry and temperature. We discuss the possible implications emanating from this analysis, which are all expected to make a significant difference to observables in heavy ion collision experiments, especially the upcoming Compressed Baryonic Matter (CBM) experiment at the future Facility for Antiproton and Ion Research (FAIR), GSI, where matter at high baryonic densities is planned to be produced.

1. Introduction

An effective description of hadronic matter is fairly common in low-energy QCD [1–3]. Realizing that baryons and mesons constitute the effective degrees of freedom in this regime, it is quite sensible to treat QCD at low-energies as an effective theory of these quark bound states [1]. This approach has been vigorously pursued in various incarnations over the years, with the different adopted strategies representing merely different manifestations of the same underlying philosophy. The actual manifestations range from the quark-meson coupling model [4, 5], phenomenological, relativistic mean-field theories based on the Walecka model [6], along with their subsequent extensions, the method of QCD sum rules [7–9], as well as the coupled channel approach [10, 11] for treating dynamically generated resonances, which has further evolved into more specialized forms, namely, the local hidden gauge theory [12, 13], and formalisms based on incorporating heavy-quark spin symmetry (HQSS) [14, 15] into the coupled channel framework [16–20]. Additionally, the method of chiral-invariant Lagrangians [1] (which will also be embraced in this work) has developed over the years into a very successful strategy. This method uses an effective field theoretical model in which the specific form of hadronic interactions is dictated by symmetry principles and the physics governed predominantly by the dynamics of chiral symmetry—its spontaneous breakdown implying a nonvanishing scalar condensate $\langle \bar{q}q \rangle$ in vacuum. One naturally expects then that the hadrons composed of these quarks would also be modified in accordance with these condensates [2, 3]. But while all hadrons would be subject to medium modifications from this perspective, pseudoscalar mesons have a special role in this context. In accordance with Goldstone's theorem [1], spontaneous breaking of chiral symmetry leads to the occurrence of massless pseudoscalar modes, the so-called Goldstone bosons, which are generally identified with the spectrum of light pseudoscalar mesons, like the pions or kaons and antikaons [2, 3]. In a strict sense, however, none of these physical mesons is a true Goldstone mode, since they are all massive, while Goldstone modes are supposed to be massless [21]. The origin of the masses of these mesons is related to the nonzero masses of light quarks, as can be easily discerned from the Gell-Mann-Oakes-Renner (GOR) relations and, hence, from explicit symmetry breaking terms (or explicit mass terms) in the chiral effective framework [1]. In fact, if one considers the limiting situation of vanishing quark masses ($m_q \rightarrow 0$), the masses of these pseudoscalar mesons would also vanish, so that the perfect Goldstone modes are indeed recovered. For this reason, the physically observed light pseudoscalar mesons are dubbed pseudo-Goldstone bosons [21]. In purely this sense, therefore, there is an inherent similarity between all these classes of pseudoscalar mesons; masses are acquired through explicit quark mass terms, and only the magnitude of these mass terms differs between these cases, being small for the pions (since $m_u, m_d < 10\,\mathrm{MeV}$), comparatively larger for the kaons and antikaons (since $m_s \sim 150\,\mathrm{MeV}$), and appreciably larger for the charmed pseudoscalar mesons. Thus, as we advance from the pions to the strange pseudoscalar mesons, with increasing mass, these mesons depart more from the ideal Goldstone mode character pertaining to the theorem and there is considerable departure of the charmed pseudoscalar mesons from Goldstone mode behaviour due to the explicit chiral symmetry breaking arising from the large charm quark mass ($m_c \simeq 1.3\,\mathrm{GeV}$). An understanding of the in-medium properties of the pseudoscalar mesons has been an important topic of research, both theoretically and experimentally. Within the chiral effective approach, the pseudoscalar mesons are modified in the medium due to the modifications of the quark condensates in the hadronic medium. For pions, it is observed, however, that medium effects for them are weakened by the smallness of explicit symmetry breaking terms [2, 3]. Considerably detailed analysis of medium effects has been performed over the years, particularly in a chiral $SU(3)$ approach, for strange pseudoscalar mesons (kaons and antikaons) [22–25]. For studying the charmed mesons, one needs to generalize the $SU(3)$ model to $SU(4)$, in order to incorporate the interactions of the charmed mesons to the light hadrons. Such a generalization from $SU(3)$ to $SU(4)$ was initially done in [26], where the interaction Lagrangian was constructed for the pseudoscalar mesons for $SU(4)$ from a generalization of the lowest order chiral $SU(3)$ Lagrangian. Since the chiral symmetry is explicitly broken for the $SU(4)$ case due to the large mass of the charm quark ($m_c \simeq 1.3\,\mathrm{GeV}$), which is much larger than the masses of the light quarks, for the study of the charmed (D) pseudoscalar mesons [27–29], we adopt the philosophy of generalizing the chiral $SU(3)$ model to $SU(4)$ to derive the interactions of these mesons with the light hadrons but use the observed masses of these heavy hadrons as well as empirical/observed values of their decay constants [30]. With all these studies proving to be informative, the most natural direction of extension of this approach would be to analyze these medium effects for a *strange-charmed* system (the D_S mesons). Apart from pure theoretical interest, an understanding of the medium modifications of D_S mesons is important, since these can make a considerable difference to experimental observables in the (ongoing and future) relativistic heavy ion collision experiments, besides being significant in questions concerning their production and transport in such experimental situations. For instance, in a recent work, He et al. [31] have shown that the modifications of the D_S meson spectrum can serve as a useful probe for understanding key issues regarding hadronization in heavy ion collisions. It is suggested that, by comparing observables for D and D_S mesons, it is possible to constrain the hadronic transport coefficient. This comparison is useful since it allows for a clear distinction between hadronic and quark-gluon plasma behavior.

However, as far as the existing literature on this strange-charmed system of mesons is concerned, we observe that only the excited states of D_S mesons have received considerable attention, predominantly as dynamically generated resonances in various coupled channel frameworks [32, 33]. One must bear in mind that, in certain situations, a molecular interpretation of these excited states (resonances) is more appropriate [34] for an explanation of their observed, larger than expected lifetimes. From this perspective, a whole plethora of possibilities have been entertained for the excited D_S states, the standard quark-antiquark picture aside. These include their description as molecular states, borne out of two mesons, four-quark states, or the still further exotic possibilities—as two-diquark states and as a mixture of quark-antiquark and tetraquark states [32]. Prominent among these is a treatment of $D^*_{S0}(2317)$ as a DK bound state [35, 36], $D_{S1}(2460)$ as a dynamically generated D^*K resonance [37], and $D^*_{S2}(2573)$ being treated within the hidden local gauge formalism in coupled channel unitary approach [38, 39], as well as the vector D^*_S states and the $D^+_S(2632)$ resonance treated in a multichannel approach [40]. The former three have also been covered consistently under the four-quark picture [41]. So, while considerable attention has been paid towards dynamically generating the higher excited states of the D_S mesons, there is a conspicuous dearth [42] of available information about the medium modifications of the lightest pseudoscalar D_S mesons, $D_S(1968.5)$, the one that we know surely is well described within the quark-antiquark picture. In fact, to the best of our knowledge, the entire existing literature about the ($J_P = 0^-$)D_S mesons in a hadronic medium is limited to the assessment of their spectral distributions and medium effects on the dynamically generated resonances borne out of the interaction of these D_S mesons with other hadron species, in the coupled channel analyses of [42–45]. This situation is quite unlike their open-charm, non-strange counterparts, the D mesons, which boast of a sizeable amount of literature having been extensively investigated using a multitude of approaches over the years. In stark contrast, the available literature concerning the in-medium behavior of pseudoscalar D_S mesons can at best be described as scanty, and there is need for more work on this subject. If one considers this problem from the point of view of the (extended) chiral effective approach, this scantiness is most of all because of the lack of a proper framework where the relevant form of the interactions for the D_S mesons with the light hadrons (or in more generic terms, of meson-baryon interactions with the charm sector covered), based on arguments of symmetry and invariance, could be written down. Clearly, such interactions would have to be based on

$SU(4)$ symmetry and bear all these pseudoscalar mesons and baryons in 15-plet and 20-plet representations, respectively, with meson-baryon interaction terms still in accordance with the general framework for writing chiral-invariant structures, as well as bearing appropriate symmetry breaking terms obeying the requisite transformation behavior under chiral transformations [1], which is quite a nontrivial problem. Of late, such formalisms have been proposed in [43, 46], as an extension of the frameworks based on chiral Lagrangians, where $SU(4)$ symmetry forms the basis for writing down the relevant interaction terms. However, since the mass of the charm quark is approximately 1.3 GeV [47], which is considerably larger than that of the up, down, and strange quarks, the $SU(4)$ symmetry is explicitly broken by this large charm quark mass. Hence, this formalism only uses the symmetry to derive the form of the interactions, whereas an explicit symmetry breaking term accounts for the large quark mass through the introduction of mass terms of the relevant (D or D_S) mesons. Also, $SU(4)$ symmetry being badly broken implies that any symmetry and order in the masses and decay constants, as predicted on the basis of $SU(4)$ symmetry, would not hold in reality. The same is acknowledged in this approach [46] and, as has been already mentioned, one does not use the masses and decay constants as expected on the basis of $SU(4)$ symmetry but rather their observed Particle Data Group (PDG) [47] values. Overall, therefore, $SU(4)$ symmetry is treated (appropriately) as being broken in this approach. Also, it is quite well-known, through both model-independent [48] and model-dependent [29] calculations, that the light quark condensates ($\langle \bar{u}u\rangle$, $\langle \bar{d}d\rangle$) are modified significantly in a hadronic medium with medium parameters like density and temperature; the strange quark condensate $\langle \bar{s}s\rangle$ is comparatively stolid and its variation is significantly more subdued, while upon advancing to the charm sector the variation in the charmed quark condensate $\langle \bar{c}c\rangle$ is altogether negligible in the entire hadronic phase [48]. These observations form the basis for treating the charm degrees of freedom of open-charm pseudoscalar mesons as frozen in the medium, as was the case in the treatments of [28, 29, 46]. Thus, as we advance from pions and kaons to the charmed pseudoscalar mesons, the generalization is perfectly natural but with the aforementioned caveats. Provided all these aspects are taken into account, a generalization of this chiral effective framework to open-charm pseudoscalar mesons is quite reasonable and sane, and the predictions from such an extended chiral effective approach bode very well [28] with alternative calculations based on the QCD sum rule approach, quark-meson coupling model, coupled channel approach, and studies of quarkonium dissociation using heavy-quark potentials from lattice QCD at finite temperatures. Additionally, it is interesting to note that this approach, followed in [28, 29, 46] for the charmed pseudoscalar mesons, has recently been extended to the bottom sector and used to study the medium behavior of the open bottom pseudoscalar B, $\bar{B}$, and B_S mesons [49, 50]. The inherent philosophy beneath this extension continues to be the same; the dynamics of the heavy quark/antiquark is treated as frozen, and the interactions of the light quark (or antiquark) of the meson, with the particles constituting the medium, are responsible for the medium modifications. With this subsequent generalization as well, the physics of the medium behavior that follows from this approach is in agreement [49] with works based on alternative, independent approaches, like the heavy meson effective theory, quark-meson coupling model, and the QCD sum rule approach. Thus, these aforementioned, prior works based on the generalization of the original chiral effective approach to include heavy flavored mesons, are totally concordant with results from alternative approaches followed in the literature, which lends an aura of credibility to this strategy. Given this backdrop, it is clear that these formalisms wipe out the reason why such an investigation for the D_S mesons within the effective hadronic model, obtained by generalizing the chiral $SU(3)$ model to $SU(4)$, has not been undertaken till date and permit this attempt to fill the void.

We organize this paper as follows. In Section 2, we outline the chiral $SU(3)_L \times SU(3)_R$ model (and its generalization to the $SU(4)$ case) used in this investigation. In Section 3, the Lagrangian density for the D_S mesons, within this extended framework, is explicitly written down and is used to derive their in-medium dispersion relations. In Section 4, we describe and discuss our results for the in-medium properties of D_S mesons, first in the nuclear matter case and then in the hyperonic matter situation, following which we briefly discuss the possible implications of these medium modifications. Finally, we summarize the entire investigation in Section 5.

2. The Effective Hadronic Model

As mentioned previously, this study is based on a generalization of the chiral $SU(3)_L \times SU(3)_R$ model [51], to $SU(4)$. We summarize briefly the rudiments of the model, while referring the reader to [51, 52] for the details. This is an effective hadronic model of interacting baryons and mesons, based on a nonlinear realization of chiral symmetry [53, 54], where chiral invariance is used as a guiding principle, in deciding the form of the interactions [55–57]. Additionally, the model incorporates a scalar dilaton field, χ, to mimic the broken scale invariance of QCD [52]. Once these invariance arguments determine the form of the interaction terms, one resorts to a phenomenological fitting of the free parameters of the model, to arrive at the desired effective Lagrangian density for these hadron-hadron interactions. The general expression for the chiral model Lagrangian density reads

$$\mathscr{L} = \mathscr{L}_{\text{kin}} + \sum_W \mathscr{L}_{\text{BW}} + \mathscr{L}_{\text{vec}} + \mathscr{L}_0 + \mathscr{L}_{\text{scale break}} + \mathscr{L}_{\text{SB}}. \tag{1}$$

In (1), $\mathscr{L}_{\text{kin}}$ is the kinetic energy term, while $\mathscr{L}_{\text{BW}}$ denotes the baryon-meson interaction term. Here, baryon-pseudoscalar meson interactions generate the baryon masses. $\mathscr{L}_{\text{vec}}$ treats the dynamical mass generation of the vector mesons through couplings with scalar mesons. The self-interaction terms of these mesons are also included in this term. $\mathscr{L}_0$ contains the meson-meson interaction terms, which induce spontaneous breaking of chiral symmetry. $\mathscr{L}_{\text{scale break}}$ introduces scale invariance breaking, via a logarithmic potential term in the scalar dilaton field, χ. Finally, $\mathscr{L}_{\text{SB}}$ refers to the explicit symmetry breaking term. This approach has been employed extensively to study the in-medium properties of hadrons, particularly pseudoscalar mesons [22–25]. As was observed in Section 1 as well, this would be most naturally extended to the charmed (nonstrange and strange) pseudoscalar mesons. However, that calls for this chiral $SU(3)$ formalism to be generalized to $SU(4)$, which has been addressed in [43, 46]. For studying the in-medium behavior of pseudoscalar mesons, the following contributions need to be analyzed [28, 29, 46]:

$$\mathscr{L} = \mathscr{L}_{\text{WT}} + \mathscr{L}_{\text{1st Range}} + \mathscr{L}_{d_1} + \mathscr{L}_{d_2} + \mathscr{L}_{\text{SME}}. \quad (2)$$

In (2), $\mathscr{L}_{\text{WT}}$ denotes the Weinberg-Tomozawa term, given by the expression [46]

$$\mathscr{L}_{\text{WT}} = -\frac{1}{2}\left[\bar{B}_{ijk}\gamma^{\mu}\left(\left(\Gamma_{\mu}\right)_l^{\ k} B^{ijl} + 2\left(\Gamma_{\mu}\right)_l^{\ j} B^{ilk}\right)\right], \quad (3)$$

with repeated indices summed over. Baryons are represented by the tensor B^{ijk}, which is antisymmetric in its first two indices [43]. The indices i, j, and k run from 1 to 4, and one can directly read the quark content of a baryon state, with the following identification: $1 \leftrightarrow u$, $2 \leftrightarrow d$, $3 \leftrightarrow s$, $4 \leftrightarrow c$. However, the heavier, charmed baryons are discounted from this analysis. In (3), Γ_{μ} is defined as

$$\Gamma_{\mu} = -\frac{i}{4}\left(u^{\dagger}\left(\partial_{\mu}u\right) - \left(\partial_{\mu}u^{\dagger}\right)u + u\left(\partial_{\mu}u^{\dagger}\right) - \left(\partial_{\mu}u\right)u^{\dagger}\right), \quad (4)$$

where the unitary transformation operator, $u = \exp(iM\gamma_5/\sqrt{2}\sigma_0)$, is defined in terms of the matrix of pseudoscalar mesons, $M = (M^a\lambda_a/\sqrt{2})$, λ_a representing the generalized Gell-Mann matrices. Further, $\mathscr{L}_{\text{SME}}$ is the scalar meson exchange term, which is obtained from the explicit symmetry breaking term

$$\mathscr{L}_{\text{SB}} = -\frac{1}{2}\operatorname{Tr}\left(A_p\left(uXu + u^{\dagger}Xu^{\dagger}\right)\right), \quad (5)$$

where $A_p = (1/\sqrt{2})\,\text{diag}[m_\pi^2 f_\pi, m_\pi^2 f_\pi, (2m_K^2 f_K - m_\pi^2 f_\pi), (2m_D^2 f_D - m_\pi^2 f_\pi)]$ and X refers to the scalar meson multiplet [28]. Also, the first range term is obtained from the kinetic energy term of the pseudoscalar mesons and is given by the expression

$$\mathscr{L}_{\text{1st Range}} = \operatorname{Tr}\left(u_{\mu}Xu^{\mu}X + Xu_{\mu}u^{\mu}X\right), \quad (6)$$

where $u_{\mu} = -i((u^{\dagger}(\partial_{\mu}u) - (\partial_{\mu}u^{\dagger})u) - (u(\partial_{\mu}u^{\dagger}) - (\partial_{\mu}u)u^{\dagger}))/4$. Lastly, the d_1 and d_2 range terms are

$$\mathscr{L}_{d_1} = \frac{d_1}{4}\left(\bar{B}_{ijk}B^{ijk}\left(u_{\mu}\right)_l^{\ m}\left(u^{\mu}\right)_m^{\ l}\right), \quad (7)$$

$$\mathscr{L}_{d_2} = \frac{d_2}{2}\left[\bar{B}_{ijk}\left(u_{\mu}\right)_l^{\ m}\left(\left(u^{\mu}\right)_m^{\ k}B^{ijl} + 2\left(u^{\mu}\right)_m^{\ j}B^{ilk}\right)\right]. \quad (8)$$

Adopting the mean-field approximation [6, 52], the effective Lagrangian density for scalar and vector mesons simplifies; the same is used subsequently to derive the equations of motion for the nonstrange scalar-isoscalar meson σ, scalar-isovector meson δ, and strange scalar meson ζ and for the vector-isovector meson ρ, nonstrange vector meson ω, and strange vector meson ϕ, within this model.

The D_S meson interaction Lagrangian density and in-medium dispersion relations, as they follow from the above general formulation, are described next.

3. D_S Mesons in Hadronic Matter

The Lagrangian density for the D_S mesons in isospin-asymmetric, strange, hadronic medium is given as

$$\mathscr{L}_{\text{total}} = \mathscr{L}_{\text{free}} + \mathscr{L}_{\text{int}}. \quad (9)$$

This $\mathscr{L}_{\text{free}}$ is the free Lagrangian density for a complex scalar field (which corresponds to the D_S mesons in this case) and reads

$$\mathscr{L}_{\text{free}} = \left(\partial^{\mu}D_S^{+}\right)\left(\partial_{\mu}D_S^{-}\right) - m_{D_S}^2\left(D_S^{+}D_S^{-}\right). \quad (10)$$

On the other hand, $\mathscr{L}_{\text{int}}$ is determined to be

$$\begin{aligned}\mathscr{L}_{\text{int}} = {} & -\frac{i}{4f_{D_S}^2}\Big[\Big(2\left(\bar{\Xi}^0\gamma^{\mu}\Xi^0 + \bar{\Xi}^-\gamma^{\mu}\Xi^-\right) + \bar{\Lambda}^0\gamma^{\mu}\Lambda^0 \\ & + \bar{\Sigma}^+\gamma^{\mu}\Sigma^+ + \bar{\Sigma}^0\gamma^{\mu}\Sigma^0 + \bar{\Sigma}^-\gamma^{\mu}\Sigma^-\Big)\Big(D_S^+\left(\partial_{\mu}D_S^-\right) \\ & - \left(\partial_{\mu}D_S^+\right)D_S^-\Big)\Big] + \frac{m_{D_S}^2}{\sqrt{2}f_{D_S}}\left[\left(\zeta' + \zeta_c'\right)\left(D_S^+D_S^-\right)\right] \\ & - \frac{\sqrt{2}}{f_{D_S}}\left[\left(\zeta' + \zeta_c'\right)\left(\left(\partial_{\mu}D_S^+\right)\left(\partial^{\mu}D_S^-\right)\right)\right] + \frac{d_1}{2f_{D_S}^2}\Big[\Big(\bar{p}p \\ & + \bar{n}n + \bar{\Lambda}^0\Lambda^0 + \bar{\Sigma}^+\Sigma^+ + \bar{\Sigma}^0\Sigma^0 + \bar{\Sigma}^-\Sigma^- + \bar{\Xi}^0\Xi^0 \\ & + \bar{\Xi}^-\Xi^-\Big)\left(\left(\partial_{\mu}D_S^+\right)\left(\partial^{\mu}D_S^-\right)\right)\Big] \\ & + \frac{d_2}{2f_{D_S}^2}\Big[\Big(2\left(\bar{\Xi}^0\Xi^0 + \bar{\Xi}^-\Xi^-\right) + \bar{\Lambda}^0\Lambda^0 + \bar{\Sigma}^+\Sigma^+ \\ & + \bar{\Sigma}^0\Sigma^0 + \bar{\Sigma}^-\Sigma^-\Big)\left(\left(\partial_{\mu}D_S^+\right)\left(\partial^{\mu}D_S^-\right)\right)\Big].\end{aligned} \quad (11)$$

In this expression, the first term (with coefficient $-i/4f_{D_S}^2$) is the Weinberg-Tomozawa term, obtained from (3), the second term (with coefficient $m_{D_S}^2/\sqrt{2}f_{D_S}$) is the scalar meson exchange term, obtained from the explicit symmetry breaking term of the Lagrangian (see (5)), the third term (with coefficient $-\sqrt{2}/f_{D_S}$) is the first range term (see (6)), and the fourth and fifth terms (with coefficients $(d_1/f_{D_S}^2)$ and $(d_2/f_{D_S}^2)$, resp.) are the d_1 and d_2 terms, calculated from (7) and (8), respectively. Also, $\zeta' = \zeta - \zeta_0$ is the fluctuation of the strange scalar field from its vacuum value. The mean-field approximation, mentioned earlier, is a useful, simplifying measure in this context, since it permits us to have the following replacements:

$$\begin{aligned}\overline{B}_i B_j &\longrightarrow \langle \overline{B}_i B_j \rangle \equiv \delta_{ij}\rho_i^s, \\ \overline{B}_i\gamma^\mu B_j &\longrightarrow \langle \overline{B}_i\gamma^\mu B_j \rangle = \delta_{ij}\left(\delta_\mu^0 \left(\overline{B}_i\gamma^\mu B_j\right)\right) \equiv \delta_{ij}\rho_i.\end{aligned} \tag{12}$$

Thus, the interaction Lagrangian density can be recast in terms of the baryonic number densities and scalar densities, given by the following expressions:

$$\rho_i = \frac{\gamma_s}{(2\pi)^3}\int d^3k\left(\frac{1}{\exp\left(\left(E_i^*(k) - \mu_i^*\right)/T\right) + 1} - \frac{1}{\exp\left(\left(E_i^*(k) + \mu_i^*\right)/T\right) + 1}\right), \tag{13}$$

$$\rho_i^s = \frac{\gamma_s}{(2\pi)^3}\int d^3k\frac{m_i^*}{E_i^*(k)}\left(\frac{1}{\exp\left(\left(E_i^*(k) - \mu_i^*\right)/T\right) + 1} + \frac{1}{\exp\left(\left(E_i^*(k) + \mu_i^*\right)/T\right) + 1}\right). \tag{14}$$

In the above, m_i^* and μ_i^* are the effective mass and effective chemical potential of the ith baryon, given as $m_i^* = -(g_{\sigma i}\sigma + g_{\zeta i}\zeta + g_{\delta i}\delta)$, $\mu_i^* = \mu_i - (g_{\rho i}\tau_3\rho + g_{\omega i}\omega + g_{\phi i}\phi)$, $E_i^*(k) = (k^2 + m_i^{*2})^{1/2}$, and $\gamma_s = 2$ is the spin degeneracy factor. One can find the equations of motion for the D_S^+ and D_S^- mesons, by the use of Euler-Lagrange equations on this Lagrangian density. The linearity of these equations follows from (11), which allows us to assume plane wave solutions ($\sim e^{i(\vec{k}\cdot\vec{r} - \omega t)}$), and hence, "Fourier transform" these equations, to arrive at the in-medium dispersion relations for the D_S mesons. These have the general form

$$-\omega^2 + \vec{k}^2 + m_{D_S}^2 - \Pi\left(\omega, \left|\vec{k}\right|\right) = 0, \tag{15}$$

where m_{D_S} is the vacuum mass of the D_S mesons and $\Pi(\omega, |\vec{k}|)$ is the self-energy of the D_S mesons in the medium. Explicitly, the latter reads

$$\begin{aligned}\Pi\left(\omega, \left|\vec{k}\right|\right) = {} & \left[\left(\frac{d_1}{2f_{D_S}^2}\left(\rho_p^s + \rho_n^s + \rho_\Lambda^s + \rho_{\Sigma^+}^s + \rho_{\Sigma^0}^s + \rho_{\Sigma^-}^s + \rho_{\Xi^0}^s + \rho_{\Xi^-}^s\right)\right) + \left(\frac{d_2}{2f_{D_S}^2}\left(2\left(\rho_{\Xi^0}^s + \rho_{\Xi^-}^s\right) + \rho_\Lambda^s + \rho_{\Sigma^+}^s + \rho_{\Sigma^0}^s + \rho_{\Sigma^-}^s\right)\right) - \left(\frac{\sqrt{2}}{f_{D_S}}\left(\zeta' + \zeta_c'\right)\right)\right]\left(\omega^2 - \vec{k}^2\right) \\ & \pm \left[\frac{1}{2f_{D_S}^2}\left(2\left(\rho_{\Xi^0} + \rho_{\Xi^-}\right) + \rho_\Lambda + \rho_{\Sigma^+} + \rho_{\Sigma^0} + \rho_{\Sigma^-}\right)\right]\omega \\ & + \left[\frac{m_{D_S}^2}{\sqrt{2}f_{D_S}}\left(\zeta' + \zeta_c'\right)\right],\end{aligned} \tag{16}$$

where the + and − signs in the coefficient of ω refer to D_S^+ and D_S^-, respectively, and we have used (12) to simplify the bilinears. In the rest frame of these mesons (i.e., setting $|\vec{k}| = 0$), these dispersion relations reduce to

$$-\omega^2 + m_{D_S}^2 - \Pi(\omega, 0) = 0, \tag{17}$$

which is a quadratic equation in ω, that is, of the form $A\omega^2 + B\omega + C = 0$, where the coefficients A, B, and C depend on various interaction terms in (15) and read

$$A = \left[1 + \left(\frac{d_1}{2f_{D_S}^2}\sum_{(N+H)}\rho_i^s\right) + \left(\frac{d_2}{2f_{D_S}^2}\left(2\sum_{\Xi}\rho_i^s + \sum_{(H-\Xi)}\rho_i^s\right)\right) - \left(\frac{\sqrt{2}}{f_{D_S}}\left(\zeta' + \zeta_c'\right)\right)\right], \tag{18}$$

$$B = \pm\left[\frac{1}{2f_{D_S}^2}\left(2\sum_{\Xi}\rho_i + \sum_{(H-\Xi)}\rho_i\right)\right], \tag{19}$$

$$C = \left[-m_{D_S}^2 + \frac{m_{D_S}^2}{\sqrt{2}f_{D_S}}\left(\zeta' + \zeta_c'\right)\right]. \tag{20}$$

As before, the "+" and "−" signs in B correspond to D_S^+ and D_S^-, respectively. In writing the above summations, we have used the following notation: H denotes the set of all hyperons, N is the set of nucleons, Ξ represents the Xi hyperons ($\Xi^{-,0}$), and $(H-\Xi)$ denotes all hyperons other than Xi hyperons, that is, the set of baryons ($\Lambda, \Sigma^{+,-,0}$) which all carry one strange quark. This form is particularly convenient for later analysis. Also, the optical potential of D_S mesons is defined as

$$U(k) = \omega(k) - \sqrt{k^2 + m_{D_S}^2}, \tag{21}$$

where k ($=|\vec{k}|$) refers to the momentum of the respective D_S meson, and $\omega(k)$ represents its momentum-dependent in-medium energy.

In the next section, we study the sensitivity of the D_S meson effective mass on various characteristic parameters of hadronic matter, namely, baryonic density (ρ_B), temperature (T), isospin asymmetry parameter $\eta = -\sum_i I_{3i}\rho_i/\rho_B$, and the strangeness fraction $f_s = \sum_i |S_i|\rho_i/\rho_B$, where S_i and I_{3i} denote the strangeness quantum number and the third component of isospin of the ith baryon, respectively.

4. Results and Discussion

Before describing the results of our analysis of D_S mesons in a hadronic medium, we first discuss our parameter choice and the various simplifying approximations employed in this investigation. The parameters of the effective hadronic model are fitted to the vacuum masses of baryons, nuclear saturation properties, and other vacuum characteristics in the mean-field approximation [51, 52]. In this investigation, we have used the same parameter set that has earlier been used to study charmed (D) mesons within this effective hadronic model [46]. In particular, we use the same values of the parameters d_1 and d_2 of the range terms ($d_1 = 2.56/m_K$ and $d_2 = 0.73/m_K$), fitted to empirical values of kaon-nucleon scattering lengths for the $I = 0$ and $I = 1$ channels, as employed in earlier treatments [22, 23, 28, 46]. For an extension to the strange-charmed system, the only extra parameter that needs to be fitted is the D_S meson decay constant, f_{D_S}, which is treated as follows. By extrapolating the results of [57], we arrive at the following expression for f_{D_S} in terms of the vacuum values of the strange and charmed scalar fields:

$$f_{D_S} = \frac{-(\zeta_0 + \zeta_{c0})}{\sqrt{2}}. \tag{22}$$

For fitting the value of f_{D_S}, we retain the same values of σ_0 and ζ_0 as in the earlier treatments of this chiral effective model [23, 51] and determine ζ_{c0} from the expression for f_D in terms of σ_0 and ζ_{c0} [28], using the Particle Data Group (PDG) value of $f_D = 206$ MeV. Substituting these in (22), our fitted value of f_{D_S} comes out to be 235 MeV, which is close to its PDG value of 260 MeV and, particularly, is of the same order as typical lattice QCD calculations for the same [47]. We therefore persist with this fitted value $f_{D_S} = 235$ MeV in this investigation. Also, we treat the charmed scalar field (ζ_c) as being arrested at its vacuum value (ζ_{c0}) in this investigation, as has been considered in other treatments of charmed mesons in a hadronic context [28, 46]. This neglect of charm dynamics appears natural from a physical viewpoint, owing to the large mass of a charm quark ($m_c \sim 1.3$ GeV) [47]. The same was verified in an explicit calculation in [48], where the charm condensate was observed to vary very weakly in the temperature range of interest to us in this regime [58–60]. As our last approximation, we point out that, in the current investigation, we work within the *"frozen glueball limit"* [52], where the scalar dilaton field (χ) is regarded as being frozen at its vacuum value (χ_0). This approximation was relaxed in a preceding work [29] within this effective hadronic model, where the in-medium modifications of this dilaton field were found to be quite meager. We conclude, therefore, that this weak dependence only serves to justify the validity of this assumption.

We next describe our analysis for the in-medium behavior of D_S mesons, beginning first with the nuclear matter ($f_s = 0$) situation and including the hyperonic degrees of freedom only later. This approach has the advantage that many features of the D_S meson in-medium behavior, common between these regimes, are discussed in detail in a more simplified context, and the effect of strangeness becomes a lot clearer.

In nuclear matter, the Weinberg-Tomozawa term and the d_2 range term vanish, since they depend on the number densities and scalar densities of the hyperons, and have no contribution from the nucleons. It follows from the self-energy expression that only the Weinberg-Tomozawa term differs between D_S^+ and D_S^-; all other interaction terms are absolutely identical for them. Thus, a direct consequence of the vanishing of Weinberg-Tomozawa contribution is that D_S^+ and D_S^- are degenerate in nuclear matter. As mentioned earlier, the D_S meson dispersion relations, with $|\vec{k}| = 0$, takes the quadratic equation form ($A\omega^2 + B\omega + C = 0$), where the coefficients (given by (18)–(20)) reduce to the following expressions in nuclear matter:

$$A_{(N)} = \left[1 + \left(\frac{d_1}{2f_{D_S}^2}\left(\rho_p^s + \rho_n^s\right)\right) - \left(\frac{\sqrt{2}}{f_{D_S}}\left(\zeta' + \zeta_c'\right)\right)\right], \tag{23}$$

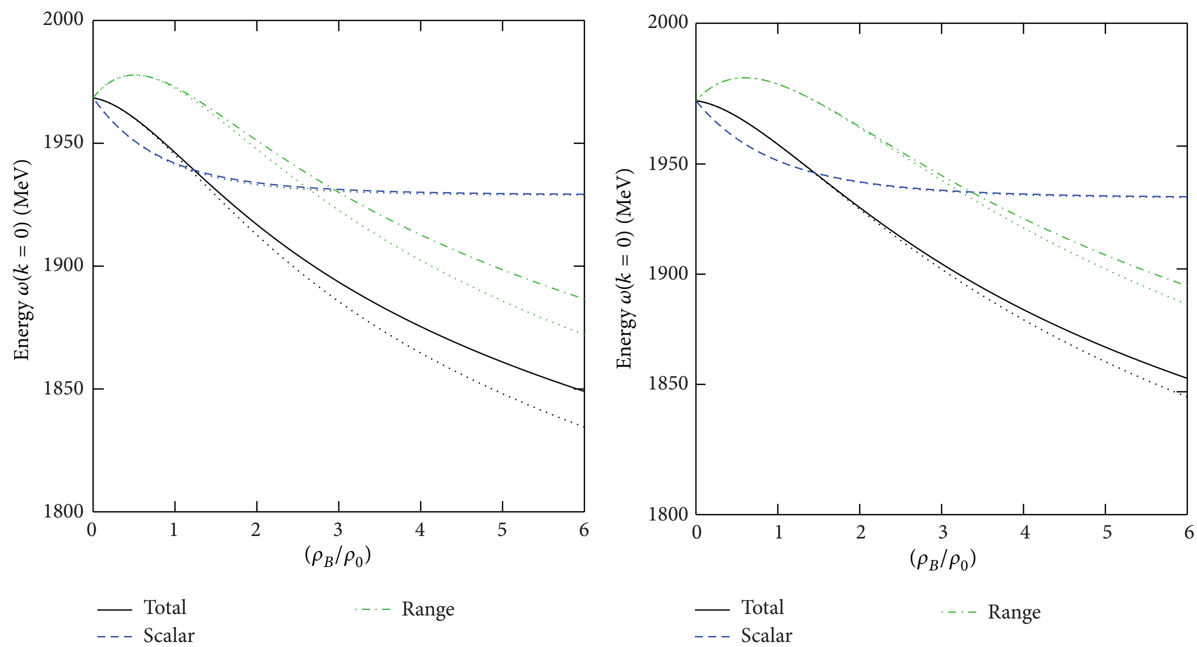
$$B_{(N)} = 0, \tag{24}$$

$$C_{(N)} = \left[-m_{D_S}^2 + \frac{m_{D_S}^2}{\sqrt{2}f_{D_S}}\left(\zeta' + \zeta_c'\right)\right], \tag{25}$$

where the subscript (N) emphasizes the nuclear matter context. For solving this quadratic equation, we require the values of ζ', as well as the scalar densities of protons and neutrons, which are obtained from a simultaneous solution of coupled equations of motion for the scalar fields, subject to constraints of fixed values of ρ_B and η. The behavior of the scalar fields, thus obtained, has been discussed in detail in [29]. Here, we build upon these scalar fields and proceed to discuss the behavior of solutions of the in-medium dispersion relations of D_S mesons, given by (23) to (25). The variation of the D_S meson in-medium mass, $\omega(\vec{k} = 0)$, with baryonic density in nuclear matter, along with the individual contributions to this net variation, is shown in Figure 1 for both zero and finite temperature value ($T = 100$ MeV). It is observed that the in-medium mass of D_S mesons decreases with density, while being weakly dependent on temperature and isospin asymmetry parameter. We can understand the observed behavior through the following analysis. From (23) to (25), we arrive at the following closed-form solution for the D_S meson effective mass in nuclear matter:

$$\omega = m_{D_S}\left[\frac{1 - \left(\zeta' + \zeta_c'\right)/\sqrt{2}f_{D_S}}{1 + \left(d_1/2f_{D_S}^2\right)\left(\rho_p^s + \rho_n^s\right) - \left(\sqrt{2}/f_{D_S}\right)\left(\zeta' + \zeta_c'\right)}\right]^{1/2}. \tag{26}$$

From this exact closed-form solution, several general conclusions regarding the in-medium behavior of D_S mesons in nuclear matter can be drawn. On the basis of this expression, the d_1 range term appearing in the denominator will lead to a decrease in the medium mass with increase in density. The in-medium mass is, additionally, also modified by the medium dependence of ζ (we assume the value of ζ_c to be frozen at its vacuum value and hence $\zeta_c' = 0$). However, the density dependence of ζ' is observed to be quite subdued in nuclear matter. The total change in the value of ζ from its vacuum value is observed; that is, $(\zeta')_{\text{max.}} \approx 15$ MeV. This value is acquired at a baryonic density of $\rho_B \approx 3\rho_0$ and thereafter it appears to saturate. This behavior of the strange scalar field, though ubiquitous in this chiral model [29, 51], is not just limited to it. Even in calculations employing the relativistic Hartree approximation to determine the variation of these scalar fields [61, 62], a behavior similar to this mean-field treatment is observed. We also point out that the saturation of the scalar meson exchange and first range term follows



(a) $T = 0$ (b) $T = 100$ MeV

FIGURE 1: The various contributions to the D_S meson energies, $\omega(\vec{k}=0)$, in nuclear matter, plotted as functions of baryonic density in units of the nuclear saturation density (ρ_B/ρ_0). In each case, the isospin asymmetric situation ($\eta = 0.5$), as described in the legend, is also compared against the symmetric situation ($\eta = 0$), represented by dotted lines.

as a direct consequence of the saturation of ζ', as is implied by their proportionality in (11) and (25). At larger densities, while the terms having ζ' saturate, the contribution from the d_1 range term continues to rise. With ζ' saturating at a value of around 15 MeV and the value of f_{D_S} chosen to be 235 MeV in the present investigation, the last term in the numerator as well as the denominator in (26) turns out to be much smaller as compared to unity. Moreover, the d_1 term in the denominator in the expression for the in-medium mass of D_S given by (26) also turns out to be much smaller as compared to unity, with $(d_1(\rho_p^s + \rho_n^s)/2f_{D_S}^2) \simeq 0.05$ at a density of $\rho_B = 6\rho_0$. Hence the smallness of this term as compared to unity is justified for the entire range of densities we are concerned with, in the present investigation. In order to read more into our solution, we expand the argument of the square root, in a binomial series (assuming $(\zeta' + \zeta_c')/\sqrt{2}f_{D_S} \ll 1$ and $(d_1(\rho_p^s + \rho_n^s)/2f_{D_S}^2) \ll 1$), and retain up to only the first-order terms. This gives, as the approximate solution,

$$\omega \approx m_{D_S}\left[1 + \frac{(\zeta' + \zeta_c')}{2\sqrt{2}f_{D_S}} - \frac{d_1}{4f_{D_S}^2}\left(\rho_p^s + \rho_n^s\right)\right]. \qquad (27)$$

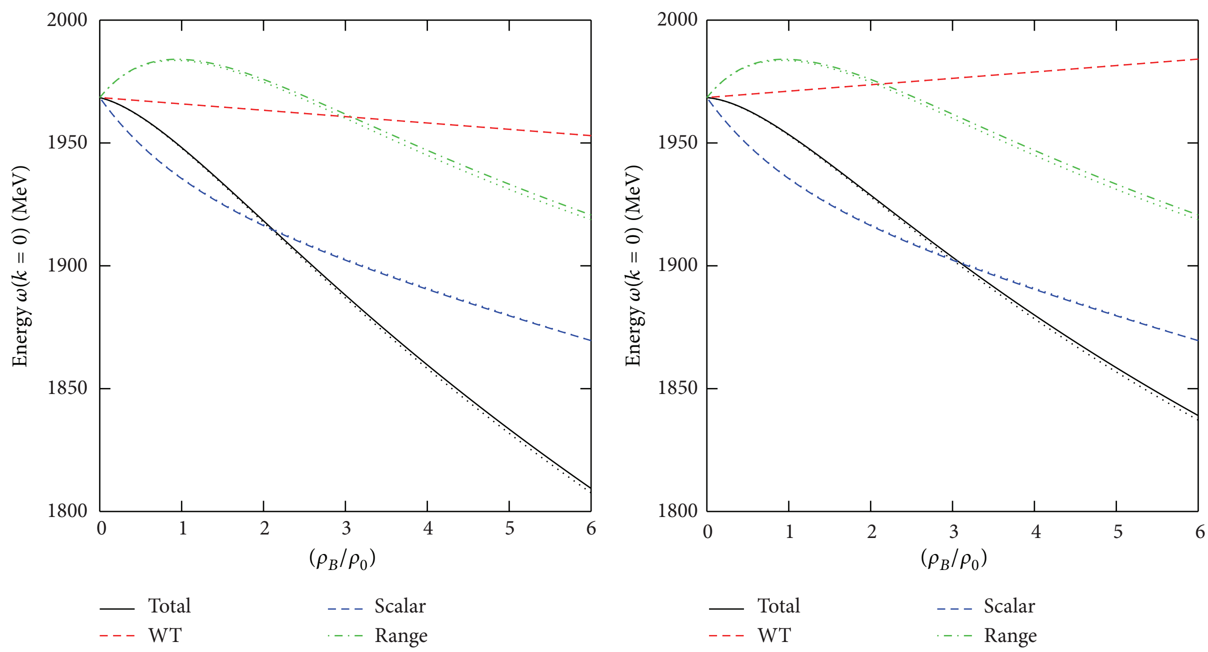
Moreover, since $\rho_i^s \approx \rho_i$ at small densities, the d_1 range term contribution in the above equation approximately equals $(d_1/4f_{D_S}^2)\rho_B$ at low densities. The first of the terms in the approximate expression given by (27) is an increasing function while the second one is a decreasing function of density. Their interplay generates the observed curve shape, the repulsive contribution being responsible for the small hump in an otherwise linear fall, at small densities ($0 < \rho_B < 0.5\rho_0$). While the attractive d_1 term would have produced a linear decrease right away, the role of repulsive contribution is to impede this decrease, hence producing the hump. At moderately higher densities, however, the contribution from the second term outweighs the first, which is why we see a linear drop with density. This is observed to be the case, till around $\rho_B \approx 2\rho_0$. At still larger densities, the approximation ($\rho_i^s \approx \rho_i$) breaks down, though our binomial expansion is still valid. Since scalar density falls slower than number density, the term $[-(d_1/4f_{D_S}^2)(\rho_p^s + \rho_n^s)]$ will fall faster than $[-(d_1/4f_{D_S}^2)\rho_B]$. This is responsible for the change in slope of the curve at intermediate densities, where a linear fall with density is no longer obeyed. So, at intermediate and large densities, the manner of the variation is dictated by the scalar densities, in the d_1 range term. From (18) to (20), one would expect the mass modifications of D_S mesons to be insensitive to isospin asymmetry, since, for example, in this nuclear matter case, the dispersion relation bears isospin symmetric terms like $(\rho_p^s + \rho_n^s)$. (This is in stark contrast with earlier treatments of kaons and antikaons [22, 23] as well as the D mesons [28] within this effective model, where the dispersion relations had terms like $(\rho_p^s - \rho_n^s)$ or $(\rho_p - \rho_n)$, which contributed to asymmetry.) However, the D_S meson effective mass is observed to depend on asymmetry in Figure 1, though the dependence is weak. For example, the values of $\omega(k = 0)$, for the isospin symmetric situation ($\eta = 0$), are 1913, 1865, and 1834 MeV, respectively, at $\rho_B = 2\rho_0$, $4\rho_0$, and $6\rho_0$, while the corresponding numbers for the (completely) asymmetric ($\eta = 0.5$) situation are 1917, 1875, and 1849 MeV, respectively, at $T = 0$. Since the η-dependence of $\mathscr{L}_{\mathrm{SME}}$ and

$\mathscr{L}_{\text{1st Range}}$ contributions must be identical, one can reason from Figure 1 that this isospin dependence of D_S meson effective mass is almost entirely due to the d_1 range term ($\sim(\rho_p^s + \rho_n^s)$), which was expected to be isospin symmetric. This apparently counterintuitive behavior has been observed earlier in [29], in the context of D mesons. This is because the value of $(\rho_p^s + \rho_n^s)$ turns out to be different for symmetric and asymmetric situations, contrary to naive expectations. Since the scalar-isovector δ meson ($\delta \sim \langle \bar{u}u - \bar{d}d \rangle$) is responsible for introducing isospin asymmetry in this effective hadronic model [51], owing to the equations of motion of the scalar fields being coupled, the values of the other scalar fields turn out to be different in the symmetric ($\delta = 0$) and asymmetric ($\delta \neq 0$) cases. The same is also reflected in the values of the scalar densities calculated from these scalar fields, which leads to the observed behavior.

Additionally, we observe from a comparison of Figures 1(a) and 1(b) that the magnitude of the D_S meson mass drop decreases with an increase in temperature from $T = 0$ to $T = 100$ MeV. For example, in the symmetric ($\eta = 0$) situation, at $T = 0$, the D_S meson mass values, at $\rho_B = 2\rho_0$, $4\rho_0$, and $6\rho_0$, are 1913, 1865, and 1834 MeV, respectively, which grow to 1925, 1879, and 1848 MeV, respectively, at $T = 100$ MeV. Likewise, with $\eta = 0.5$, the corresponding values read 1917, 1875, and 1849 MeV at $T = 0$, while the same numbers, at $T = 100$ MeV, are 1925, 1883, and 1855 MeV. Thus, though small, there is a definite reduction in the magnitude of the mass drops, as we go from $T = 0$ to $T = 100$ MeV, in each of these cases. This behavior can be understood, from the point of view of the temperature variation of scalar condensates, in the following manner. It is observed that the scalar fields decrease with an increase in temperature from $T = 0$ to 100 MeV [27, 29]. In [27], the same effect was understood as an increase in the nucleon mass with temperature. The equity of the two arguments can be seen by invoking the expression relating the baryon mass to the scalar fields' magnitude, $m_i^* = -(g_{\sigma i}\sigma + g_{\zeta i}\zeta + g_{\delta i}\delta)$ [51]. The temperature dependence of the scalar fields (σ, ζ, δ) has been studied within the model in [29], which are observed to be different for the zero and finite baryon densities. At zero baryon density, the magnitudes of the scalar fields σ and ζ are observed to remain almost constant up to a temperature of about 125 MeV, above which these are observed to decrease with temperature. This behaviour can be understood from the expression of ρ_s^i given by (14) for the situation of zero density, that is, for $\mu_i^* = 0$, which decreases with increase in temperature. The temperature dependence of the scalar density in turn determines the behaviour of the scalar fields. The scalar fields which are solved from their equations of motion behave in a similar manner as the scalar density. At finite densities, that is, for nonzero values of the effective chemical potential, μ_i^*, however, the temperature dependence of the scalar density is quite different from the zero density situation. For finite baryon density, with increase in temperature, there are contributions also from higher momenta, thereby increasing the denominator of the integrand on the right-hand side of the baryon scalar density given by (14). This leads to a decrease in the scalar density. At finite baryon densities, the competing effects of the thermal distribution functions and the contributions from the higher momentum states give rise to the temperature dependence of the scalar density, which in turn determine the behaviour of the σ and ζ fields with temperature. These scalar fields are observed to have an initial increase in their magnitudes up to a temperature of around 125–150 MeV, followed by a decrease as the temperature is further raised [29]. This kind of behavior of the scalar σ field on temperature at finite densities has also been observed in the Walecka model in [63]. In fact, we point out that a decrease in the scalar condensates with an increase in temperature, though small in the hadronic regime (<170 MeV), is well-known in general model-independent terms [64] and was also observed to be a consistent feature of all $U(N_f)_L \times U(N_f)_R$ linear sigma models in the model-independent work of Röder et al. [48]. Since these scalar fields serve as an input in calculating the scalar densities [51], a decrease in the magnitude of the latter accompanies a decrease in the former, at larger temperatures. From the point of view of the dispersion relations, this results in a decrease in the coefficient A and, owing to the inverse dependence of ω on A, increases the value of $\omega(k = 0)$ in the finite temperature case, as compared to the $T = 0$ situation. Or stating it differently, the difference of this $\omega(k = 0)$, from the vacuum value, that is, the mass drop, decreases. Thus, from a physical viewpoint, since the origin of these mass drops is the attractive in-medium interactions, one can say that the reduction in the mass drop magnitudes is due to a weakening of the attractive strength of these in-medium interactions, represented in these models by a reduction in the quark condensates with increasing temperatures. Additionally, it is also observed from Figure 1 that isospin dependence of the D_S meson mass, feeble anyways, weakens further with temperature. For example, as mentioned earlier, mass of the D_S mesons at $\rho_B = 6\rho_0$, for the $\eta = 0$ and $\eta = 0.5$ cases, is 1834 and 1849 MeV, respectively (a difference of 15 MeV), when $T = 0$, which changes to 1848 and 1855 MeV at $T = 100$ MeV (a 7 MeV difference). This is, once again, due to a decrease in the magnitude of δ, with temperature, at any fixed value of the parameters η and ρ_B. In particular, the difference between the values of δ for the $\eta = 0$ and $\eta = 0.5$ cases is observed to decrease with temperature (as was shown explicitly in [29]). Since asymmetry is introduced through δ, a decrease in the difference of the values of ω, between symmetric and asymmetric situations, with temperature, follows naturally.

Next, we generalize our analysis by including hyperonic degrees of freedom as well, in the medium. However, as mentioned previously, in the ensuing discussion of D_S mesons in hyperonic matter, we focus predominantly on the new physics arising via the introduction of strangeness in the medium, since, for example, the weak dependence on isospin asymmetry, or a weak reduction in the mass drop magnitudes at higher temperatures, has, in principle, the same explanation that stood in the nuclear matter context.

Figure 2 shows the variation of the in-medium mass of the D_S mesons, along with the various individual contributions, in hyperonic medium as a function of baryonic density, for typical values of temperature, isospin asymmetry parameter, and strangeness fraction. The most drastic consequence of the inclusion of hyperons in the medium is that the mass



(a) D_S^+

(b) D_S^-

FIGURE 2: The various contributions to the effective mass of D_S^+ and D_S^- mesons, as a function of baryonic density, for typical values of temperature ($T = 100$ MeV) and strangeness fraction ($f_s = 0.5$). In each case, the asymmetric situation ($\eta = 0.5$), as described in the legend, is also compared against the symmetric situation ($\eta = 0$), represented by dotted lines.

degeneracy of D_S^+ and D_S^- now stands broken. For example, the values of $(m_{D_S^+}, m_{D_S^-})$ at $\rho_B = \rho_0$, $2\rho_0$, $4\rho_0$, and $6\rho_0$ are observed to be (1948, 1953), (1918, 1928), (1859, 1879), and (1808, 1838) MeV, respectively, as can be seen from the figure. Thus, except at vacuum ($\rho_B = 0$), mass difference of D_S^+ and D_S^- is nonzero at finite ρ_B, growing in magnitude with density. This mass degeneracy breaking is a direct consequence of nonzero contributions from the Weinberg-Tomozawa term, which follows from (11), (16), and (19). The same may be reconciled with Figure 2, from which it follows that, except for this Weinberg-Tomozawa term acquiring equal and opposite values for these mesons, all other terms are absolutely identical for them. Once again, on the basis of the following analysis of the D_S meson dispersion relations at zero momentum, we insist that this observed behavior is perfectly consistent with expectations.

The general solution of the equivalent quadratic equation (17) is

$$\omega = \frac{\left(-B + \sqrt{B^2 - 4AC}\right)}{2A} \approx -\frac{B}{2A} + \sqrt{\frac{C_1}{A}} + \frac{B^2}{8A\sqrt{AC_1}} + \cdots, \quad (28)$$

where we have disregarded the negative root and have performed a binomial expansion of $(1 + B^2/4AC_1)^{1/2}$, with $C_1 = -C$. Upon feeding numerical values, we observe that the expansion parameter $(B^2/4AC_1)$ is much smaller than unity for our entire density variation, which justifies the validity of this expansion for our analysis. The same exercise also allows us to safely disregard higher-order terms and simply write

$$\omega_{\text{hyp}} \approx \sqrt{\frac{C_1}{A}} - \frac{B}{2A}. \quad (29)$$

Further, with the same justification as in the nuclear matter case, both $(C_1/A)^{1/2}$ and $(1/A)$ can also be expanded binomially. For example, for the second term, this gives

$$\frac{|B|}{2A} = \frac{|B|}{2} + \mathcal{O}\left(\frac{\rho_i \rho_i^s}{f_{D_S}^4}\right), \quad (30)$$

where the contribution from these higher-order terms is smaller owing to the large denominator, prompting us to retain only the first-order terms. Here, we point out that since $B = \pm|B|$ (+ sign for D_S^+ and − sign for D_S^-), this term, which represents the Weinberg-Tomozawa contribution to the dispersion relations, breaks the degeneracy of the D_S mesons. On the other hand, the first term in (29) is common between D_S^+ and D_S^- mesons. Thus, the general solution of the D_S meson dispersion relations in the hyperonic matter context can be written as

$$\omega_{\text{hyp}} = \omega_{\text{com}} \mp \omega_{\text{brk}}, \quad (31)$$

where

$$\omega_{\text{com}} = \sqrt{\frac{C_1}{A}} \approx m_{D_S}\left[1 - \left(\frac{d_1}{4f_{D_S}^2}\sum_{(N+H)}\rho_i^s\right) - \left(\frac{d_2}{4f_{D_S}^2}\left(2\sum_{\Xi}\rho_i^s + \sum_{(H-\Xi)}\rho_i^s\right)\right) + \left(\frac{(\zeta' + \zeta_c')}{2\sqrt{2}f_{D_S}}\right)\right], \quad (32)$$

$$\omega_{\text{brk}} = \frac{|B|}{2A} \approx \frac{|B|}{2} = \left[\frac{1}{4f_{D_S}^2}\left(2\sum_{\Xi}\rho_i + \sum_{(H-\Xi)}\rho_i\right)\right]. \quad (33)$$

Thus, $\omega_{D_S^+} = \omega_{\text{com}} - \omega_{\text{brk}}$ and $\omega_{D_S^-} = \omega_{\text{com}} + \omega_{\text{brk}}$ (where ω_{brk} is necessarily positive), which readily explains why the mass of D_S^+ drops more than that of D_S^-, with density. Also, this formulation readily accounts for symmetric fall of the medium masses of D_S^+ and D_S^- mesons, about ω_{com} in Figure 3. Trivially, it may also be deduced that ω_{brk} is extinguished in nuclear matter, so that the mass degeneracy of D_S^+ and D_S^- is recovered from these equations. However, we observe that though the curve corresponding to ω_{com} is exactly bisecting the masses of D_S^+ and D_S^- at small densities, this bisection is no longer perfect at high densities. This can be understood in the following manner. In essence, we are comparing the density dependence of a function $f(\rho_i^s)$, with the functions $f(\rho_i^s) \pm g(\rho_i)$. Since scalar densities fall sublinearly with the number density, it follows that the fall can not be absolutely symmetric at any arbitrary density. In fact, since a decreasing function of scalar density will fall faster than that of a number density, one expects ω_{com} to lean away from the curve for D_S^- meson and towards the curve corresponding to D_S^+, which is exactly what is observed in Figure 3.

Since this disparity between D_S^+ and D_S^- originates from the Weinberg-Tomozawa term, whose magnitude is directly proportional to the hyperonic number densities, one expects this disparity to grow with an increase in the strangeness content of the medium. The same also follows from the above formulation, since the mass splitting between the two, $\Delta m = (m_{D_S^-} - m_{D_S^+}) \equiv (\omega_{D_S^-} - \omega_{D_S^+}) = -2\omega_{\text{brk}}$, should grow with both f_s at fixed ρ_B (i.e., a larger proportion of hyperons) and ρ_B at fixed (nonzero) f_s (i.e., a larger hyperonic density), in accordance with (33). Naively, one expects this f_s dependence to be shared by the two D_S mesons; however, closer analysis reveals that this is not the case, as shown in Figure 4 where we consider the f_s dependence of the D_S meson medium mass. It is observed that while the f_s dependence for D_S^+ meson is quite pronounced, the same for the D_S^- is conspicuously subdued. Counterintuitive as it may apparently be, this observed behavior follows from the above formulation. Since $\omega_{D_S^-} = \omega_{\text{com}} + \omega_{\text{brk}}$, at any fixed density, the first of these is a decreasing function of f_s, while the second increases with f_s. These opposite tendencies are responsible for the weak f_s dependence of D_S^- meson. On the other hand, because $\omega_{D_S^+} = \omega_{\text{com}} - \omega_{\text{brk}}$, the two effects add up to produce a heightened overall decrease with f_s for the D_S^+ meson, as we observe in the figure.

A comparison of these hyperonic matter ($f_s \neq 0$) solutions with the nuclear matter ($f_s = 0$) solutions is shown in Figure 5 for typical values of the parameters. It is observed

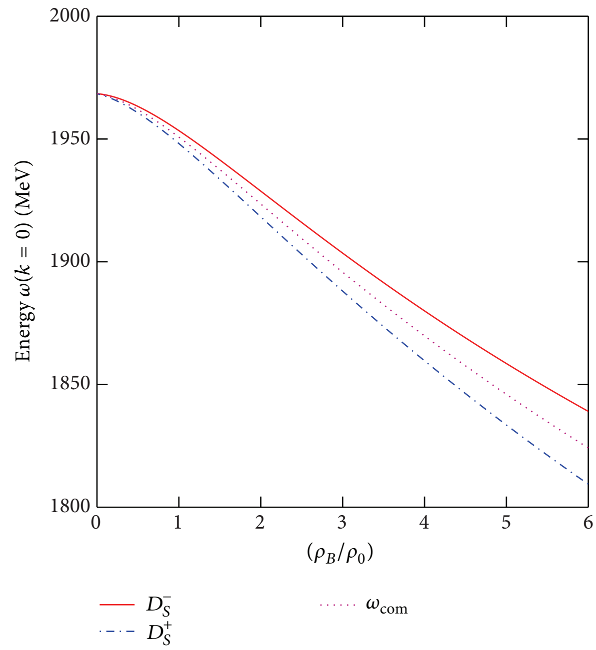


Figure 3: D_S meson mass degeneracy breaking in hyperonic matter, as a function of baryonic density, for typical values of the other parameters ($T = 100$ MeV, $f_s = 0.5$, and $\eta = 0.5$). The medium masses of D_S^+ and D_S^- are observed to fall symmetrically about ω_{com} (*see text*).

from Figure 5(a) that, at small densities, the nuclear matter curve, which represents both D_S^+ and D_S^-, bisects the mass degeneracy breaking curve, akin to ω_{com} in Figure 3. However, at larger densities, both of these $f_s \neq 0$ curves drop further than the $f_s = 0$ curve. From a casual comparison of expressions, it appears that the nuclear matter solution, given by (27), also enters the hyperonic matter solution, (32), albeit as a subset of ω_{com}. It is tempting to rearrange the latter then, such that this nuclear matter part is separated out, generating an expression of the type $\omega_{\text{hyp}} = \omega_{\text{nucl}} + f(\rho_i, \rho_i^s)|_{i=H}$, which would also entail the segregation of the entire f_s dependence. However, careful analysis reveals that it is impossible to achieve this kind of a relation, since, in spite of an identical expression, in hyperonic matter, the RHS of (27) does not give the nuclear matter solution. This is because the nuclear matter solution involves the scalar fields (and scalar densities computed using these scalar fields) obtained as a solution of coupled equations in the $f_s = 0$ situation, which are radically different from the solutions obtained in the $f_s \neq 0$ case. (This has been observed to be the case, in almost every treatment of hyperonic matter within this effective model but perhaps most explicitly in [46].) Thus, one can not retrieve the nuclear matter solution from the scalar fields obtained as solutions in the $f_s \neq 0$ case; moreover, due to the complexity of the concerned equations, it is impossible to achieve a closed-form relation between the values of the scalar fields (and hence also the scalar densities) obtained in the two cases. In fact, the low-density similarity between Figures 3 and 5(a) might lead one to presume that ω_{com} reduces to ω_{nucl} at small densities. This expectation is fueled further by their comparison in

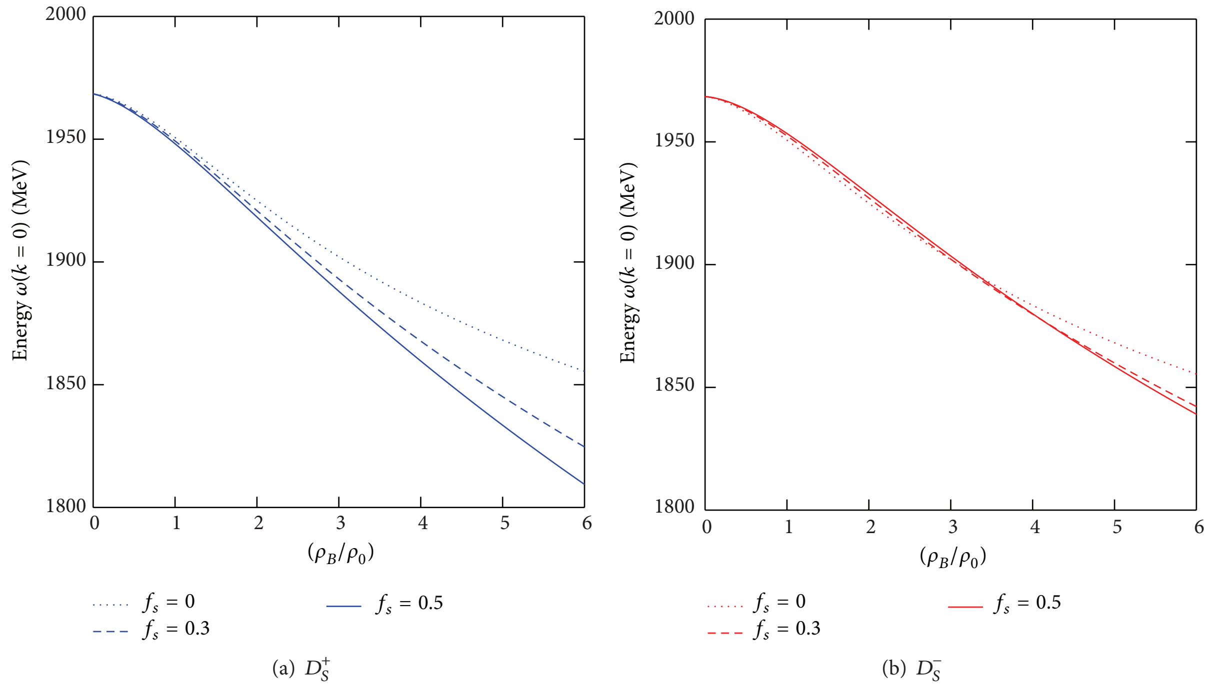


FIGURE 4: The sensitivity of the medium mass of (a) D_S^+ and (b) D_S^- mesons, to the strangeness fraction f_s, at typical values of temperature ($T = 100$ MeV) and isospin asymmetry parameter ($\eta = 0.5$).

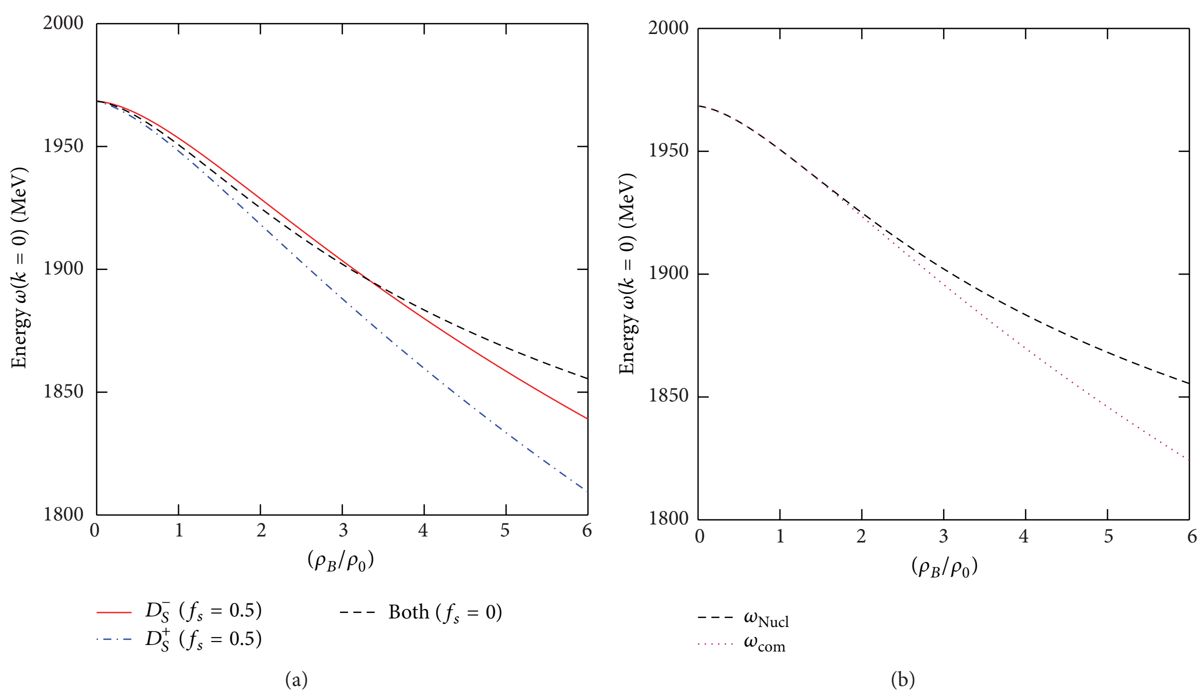
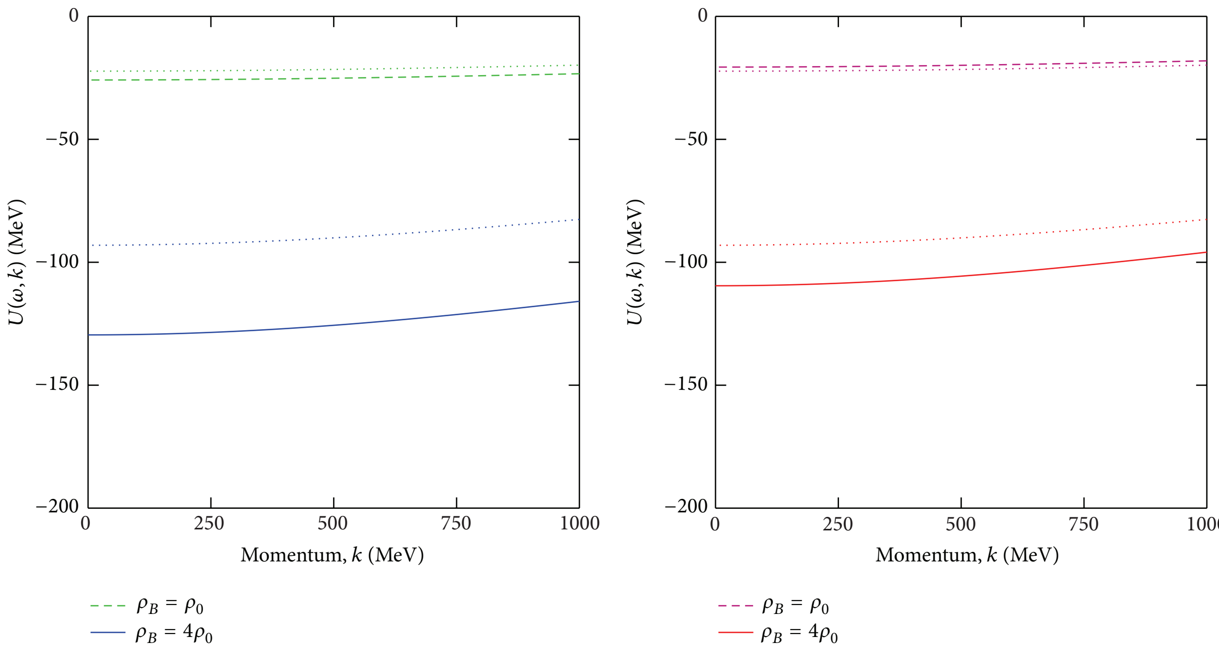


FIGURE 5: (a) A comparison of medium mass of D_S mesons, in nuclear and hyperonic matter, maintaining the same values of parameters as before. (b) Comparison of the nuclear matter solution, ω_{nucl}, with the ω_{com} in hyperonic matter, as observed in Figure 3.



(a) D_S^+

(b) D_S^-

FIGURE 6: The optical potentials of the (a) D_S^+ and (b) D_S^- mesons, as a function of momentum k $(\equiv|\vec{k}|)$, in cold ($T = 0$), asymmetric ($\eta = 0.5$) matter, at different values of ρ_B. In each case, the hyperonic matter situation ($f_s = 0.5$), as described in the legend, is also compared against the nuclear matter ($f_s = 0$) situation, represented by dotted lines.

Figure 5(b), which shows clearly that they coincide at small densities. However, in light of the above argument, it follows that they are absolutely unrelated. Their coincidence at small densities can be explained by invoking the relation $\rho_i^s \approx \rho_i$ in this regime, using which (32) reduces to

$$\omega_{\text{com}} \approx m_{D_S}\left(1 + \frac{(\zeta' + \zeta_c')}{2\sqrt{2} f_{D_S}} - \frac{\kappa_1}{4 f_{D_S}^2}\rho_B\right) - m_{D_S}\left(\frac{\kappa_2 \Delta}{4 f_{D_S}^2}\right), \tag{34}$$

with $\kappa_1 = (d_1 + d_2) \approx 1.28 d_1$, with our choice of parameters, κ_2 $(\equiv d_2) = 0.22\kappa_1$, and with the difference term, $\Delta = (\rho_{\Xi^0}^s + \rho_{\Xi^-}^s - \rho_p^s - \rho_n^s)$. The contribution from this second term in (34) is significantly smaller than the first term at small densities, owing to the smaller coefficient, as well as the fact that this depends on the difference of densities rather than on their sum (like the first term). This first part of (34) can be likened to the approximate nuclear matter solution at small ρ_B, as we concluded from (27). The marginal increase of κ above d_1 is compensated by a marginal increase in ζ' for the $f_s \neq 0$ case, as compared to the $f_s = 0$ situation, and hence the two curves look approximately the same at small densities in Figure 5(b). At larger densities, however, this simple picture breaks down, and the fact that they are unrelated becomes evident. Nevertheless, we may conclude in general from this comparison that the inclusion of hyperonic degrees of freedom in the hadronic medium makes it more attractive, regarding D_S mesons' in-medium interactions, especially at large ρ_B. From the point of view of the D_S meson dispersion relations, this is conclusively because of the contributions from the extra, hyperonic terms, which result in an overall increase of the coefficient A_{hyp} significantly above A_{nucl}, particularly at large f_s.

Finally, we consider momentum-dependent effects by investigating the behavior of the in-medium optical potentials for the D_S mesons. These are shown in Figure 6, where we consider them in the context of both asymmetric nuclear ($f_s = 0$) and hyperonic matter ($f_s = 0.5$), as a function of momentum k $(=|\vec{k}|)$ at fixed values of the parameters ρ_B, η, and T. As with the rest of the investigation, the dependence of the in-medium optical potentials on the parameters T and η is quite weak; hence, we neglect their variation in this context. In order to appreciate the observed behavior of these optical potentials, we observe that as per its definition, (21), at zero momentum, optical potential is just the negative of the mass drop of the respective meson (i.e., $U(k = 0) = -\Delta m(k = 0) \equiv \Delta m(\rho_B, T, \eta, f_s)$). It follows then that, at $k = 0$, the two D_S mesons are degenerate in nuclear matter; moreover, it is observed that this degeneracy extends to the finite momentum regime. This is because, from (15) and (16), the finite momentum contribution is also common for D_S^+ and D_S^- in nuclear matter. In hyperonic matter, however, nonzero contribution from these terms breaks the degeneracy in the $k = 0$ limit (as was already discussed before), and finite momentum preserves this nondegeneracy. Consequently, at any fixed density, the curves for different values of f_s differ in terms of their y-intercept but otherwise appear to run parallel. Further, as we had reasoned earlier for the $k = 0$

case, the effect of increasing f_s on the mass drops of D_S^+ and D_S^- is equal and opposite about the nuclear matter situation at small densities, which readily explains the behavior of optical potentials for these mesons at $\rho_B = \rho_0$. At larger densities, for example, the $\rho_B = 4\rho_0$ case shown in Figure 6, the lower values of optical potentials for both D_S^+ and D_S^-, as compared to the $f_s = 0$ case, can be immediately reconciled with the behavior observed in Figure 5. In fact, the large difference between the D_S^+ optical potentials for the $f_s = 0$ and $f_s = 0.5$ cases, as compared to that for D_S^-, is a by-product of their zero-momentum behavior, preserved at nonzero momenta as a consequence of (15) and (16).

We next discuss the possible implications of these medium effects. In the present investigation, we have observed a reduction in the mass of the D_S mesons, with an increase in baryonic density. This decrease in D_S meson mass can result in the opening up of extra reaction channels of the type $A \to D_S^+ D_S^-$ in the hadronic medium. In fact, a comparison with the spectrum of known charmonium states [47, 65] sheds light on the possible decay channels, as shown in Figure 7. As a starting approximation, we have neglected the variation of energy levels of these charmonia with density, since we intuitively expect medium effects to be larger for $s\bar{c}$ (or $c\bar{s}$) system, as compared to $c\bar{c}$ system. However, this approximation can be conveniently relaxed, for example, as was done in [46, 66] for charmonia and bottomonia states, respectively. It follows from Figure 7 that above certain threshold density, which is given by the intersection of the D_S meson curve with the respective energy level, the mass of the $D_S^+ D_S^-$ pair is smaller than that of the excited charmonium state. Hence, this decay channel opens up above this threshold density. The most immediate experimental consequence of this effect would be a decrease in the production yields for these excited charmonia in collision experiments. Additionally, since an extra decay mode will lower the lifetime of the state, one expects the decay widths of these excited charmonia to be modified as a result of these extra channels. Though it can be reasoned right away that reduction in lifetime would cause a broadening of the resonance curve, more exotic effects can also result from these level crossings. For example, it has been suggested in [67] that if the internal structure of the hadrons is also accounted for, there is even a possibility for these in-medium widths to become narrow. This counterintuitive result originates from the node structure of the charmonium wavefunction, which can even create a possibility for the decay widths to vanish completely at certain momenta of the daughter particles. Also, if the medium modifications of the charmonium states are accounted for, the decay widths will be modified accordingly (as was observed in the context of the D mesons, in [46]).

Apart from this, one expects signatures of these medium effects to be reflected in the particle ratio (D_S^+/D_S^-). A production asymmetry between D_S^+ and D_S^- mesons might be anticipated owing to their unequal medium masses, as observed in the current investigation. Also, due to the semileptonic and leptonic decay modes of the D_S mesons [47], one also expects that, accompanying an increased production of D_S mesons due to a lowered mass, there would be an enhancement

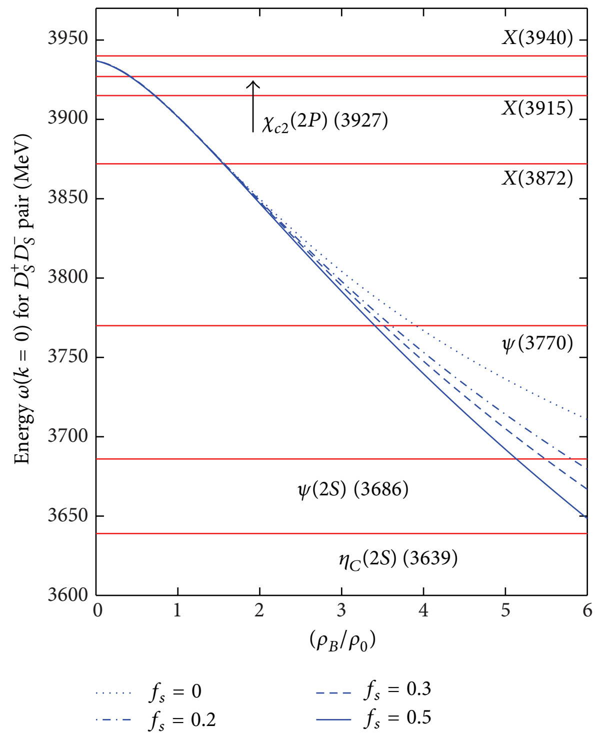


FIGURE 7: Mass of the $D_S^+D_S^-$ pair, compared against the (vacuum) masses of the excited charmonia states, at typical values of temperature ($T = 100$ MeV) and asymmetry ($\eta = 0.5$), for both nuclear matter ($f_s = 0$) and hyperonic matter ($f_s = 0.2$, 0.3, and 0.5) situations. The respective threshold density (*see text*) is given by the point of intersection of the two concerned curves.

in dilepton production as well. Further, we expect these medium modifications of the D_S mesons to be reflected in the observed dilepton spectra, due to the well-known fact that dileptons, with their small interaction cross section in hadronic matter, can serve as probes of medium effects in collision experiments [68, 69]. The attractive interaction of the D_S mesons with the hadronic medium might also lead to the possibility of formation of D_S-nucleus bound state, which can be explored at the CBM experiment at FAIR in the future facility at GSI [70].

We now discuss how the results of the present investigation compare with the available literature on the medium effects for pseudoscalar $D_S(1968.5)$ mesons. As has already been mentioned, [42–45] have treated the medium behavior of D_S mesons, using the coupled channel framework. The broad perspective that emerges from all these analyses is that of an attractive interaction in the medium [44, 45] between D_S mesons and baryon species like the nucleons or the Λ, Σ, Ξ hyperons [42–45], which is consistent with what we have observed in this work. Additionally, one observes in these approaches an attractive medium interaction even with the charmed baryons [43]. Some of these interaction channels also feature resonances, which is reflected in the relevant scattering amplitudes picking up imaginary parts as well, for example, the $N_{\bar{c}s}(2892)$ state in the D_S^+N channel [42, 43]. Treating such resonances comes naturally to the coupled channel framework, since, by default, the strategy is aimed at considering the scattering of various hadron species off one another and assessing scattering amplitudes and cross sections, and so forth [10, 11]. In the present work, we have investigated the in-medium masses of the D_S mesons arising due to their interactions with the baryons and scalar mesons in the nuclear (hyperonic) medium and have not studied the decays of these mesons. As has already been mentioned the mass modifications of the D_S mesons can lead to opening up of the new channels for the charmonium states decaying to $D_S^+ D_S^-$ at certain densities as can be seen from Figure 7. In fact, even if one contemplates working along the lines of [46] for calculating these decay widths, the situation is a bit more complicated for these D_S mesons, since the states $X(3872)$, $X(3915)$, and $X(3940)$, which are highly relevant in this context (as one can see from Figure 7), are not classified as having clear, definite quantum numbers in the spectrum of excited charmonium states [47]. This makes the assessment of medium effects for these states more difficult in comparison, since the identification of the relevant charmonium states with definite quantum numbers in the charmonium spectrum, such as $1S$ for J/ψ, $2S$ for $\psi(3686)$, and $1D$ for $\psi(3770)$, was a crucial requirement for evaluating their mass shifts in the medium in the QCD second-order Stark effect study in [46]. For some of these unconventional states, for example, $X(3872)$, this absence of definite identification with states in the spectrum has led to alternative possibilities being explored for these states. A prominent example is the possibility of a molecular structure for this $X(3872)$ state [65], borne out of contributions from the $(D^+ D^{*-} + D^- D^{*+})$ and $(D^0 \overline{D}^{*0} + \overline{D}^0 D^{*0})$ components in the s-wave.

Lastly, we point out that the medium effects described in this paper, and the possible experimental consequences entailed by these, are especially interesting in wake of the upcoming CBM [70] experiment at FAIR, GSI, where high baryonic densities are expected to be reached. Due to the strong density dependence of these medium effects, we expect that each of these mentioned experimental consequences would intensify at higher densities and should be palpable in the aforementioned future experiment.

5. Summary

To summarize, we have explored the properties of D_S mesons in a hot and dense hadronic environment, within the framework of the chiral $SU(3)_L \times SU(3)_R$ model, generalized to $SU(4)$. The generalization of chiral $SU(3)$ model to $SU(4)$ is done in order to derive the interactions of the charmed mesons with the light hadrons, needed for the study of the in-medium properties of the D_S mesons in the present work. However, realizing that the chiral symmetry is badly broken for the $SU(4)$ case due to the large mass of the charm quark, we use the interactions derived from $SU(4)$ for the D_S meson but use the observed masses of these heavy pseudoscalar mesons as well as empirical/observed values of their decay constants [30]. The D_S mesons have been considered in both (symmetric and asymmetric) nuclear and hyperonic matter, at finite densities that extend slightly beyond what can be achieved with the existing and known future facilities and at both zero and finite temperatures. Due to net attractive interactions in the medium, D_S mesons are observed to undergo a drop in their effective mass. These mass drops are found to intensify with an increase in the baryonic density of the medium, while being largely insensitive to changes in temperature as well as the isospin asymmetry parameter. However, upon adding hyperonic degrees of freedom, the mass degeneracy of D_S^+ and D_S^- is observed to be broken. The mass splitting between D_S^+ and D_S^- is found to grow significantly with an increase in baryonic density as well as the strangeness content of the medium. Through a detailed analysis of the in-medium dispersion relations for the D_S mesons, we have shown that the observed behavior follows precisely from the interplay of contributions from various interaction terms in the Lagrangian density. We have briefly discussed the possible experimental consequences of these medium effects, for example, in the D_S^+/D_S^- ratio, dilepton spectra, possibility of formation of exotic D_S-nucleus bound states, and modifications of the decays of charmonium states $D_S^+ D_S^-$ pair in the hadronic medium. The medium modifications of the D_S mesons are expected to be considerably enhanced at large densities and hence the experimental consequences may be accessible in the upcoming CBM experiment, at the future facilities of FAIR, GSI [70].

Conflict of Interests

The authors declare that there is no conflict of interests regarding the publication of this paper.

Acknowledgments

Divakar Pathak acknowledges financial support from University Grants Commission, India (Sr. no. 2121051124, Ref. no. 19-12/2010(i)EU-IV). Amruta Mishra would like to thank Department of Science and Technology, Government of India (Project no. SR/S2/HEP-031/2010), for financial support.

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