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Relation between the mass modification of heavy–light mesons and the chiral symmetry structure in dense matter

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We point out that the study of the density dependences of the masses of heavy–light mesons gives some clues to the chiral symmetry structure in nuclear matter. We include the omega meson effect as well as the sigma meson effect at mean-field level on the density dependence of the masses of heavy–light mesons with chiral partner structure. It is found that the omega meson affects the masses of the heavy–light mesons and their antiparticles in opposite ways, while it affects the masses of chiral partners in the same way. This is because the ω meson is sensitive to the baryon number of the light degrees included in the heavy–light mesons. We also show that the mass difference between chiral partners is proportional to the mean field of sigma, reflecting the partial restoration of chiral symmetry in the nuclear matter. In addition to the general illustration of the density dependence of the heavy–light meson masses, we consider two concrete models for nuclear matter, the parity doublet model and the skyrmion crystal model in the sense of the mean-field approximation.

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1. Introduction

Spontaneous chiral symmetry breaking is one of the most important properties of low-energy QCD. It is expected that the spontaneous chiral symmetry breaking characterized by non-zero value of the quark condensate generates part of the hadron masses and causes the mass splitting between chiral partners. Then, schematically, hadron masses can be expressed as the sum of the chiral invariant mass and the chiral non-invariant mass coming from the spontaneous chiral symmetry breaking. For example, for the nucleon mass, one has [1–7]

$$m_N = m_0 + \Delta(\langle\bar{q}q\rangle),$$

where m_0 is the chiral invariant mass and Δ is the part of the mass that vanishes in the chiral symmetric phase. Naturally, it is interesting to ask how much of the hadron mass is generated by chiral symmetry breaking. An ideal environment to estimate the magnitude of the hadron mass coming from spontaneous chiral symmetry breaking is QCD at extreme conditions in which the chiral symmetry is believed to be partially restored. In such an environment this can be accessed

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by studying the temperature and/or, as will be done in this work, the density dependence of hadron mass.

In the nucleon sector, by using an effective model with parity doublet structure of baryons, it has been found that $\sim$ 70% of nucleon mass comes from chiral symmetry breaking [1,2]. However, when the baryon as a topological soliton in the hidden local symmetry Lagrangian is immersed in the dense matter, which is treated as skyrmion matter, it was found that the chiral invariant mass comprises $\sim$ 60% of the nucleon mass [4,5], which is roughly the same as that obtained based on the renormalization group analysis of the hidden local symmetry Lagrangian with baryons [3]. In this paper, we study the medium-modified mass splitting of heavy–light mesons with chiral partner structure, in which it is widely accepted that the mass splitting arises from the spontaneous breaking of chiral symmetry [8,9]. This type of study can be tested in the planned experiments at J-PARC, FAIR, and so on.

Studying the properties of heavy–light mesons in medium is also expected to give clues for understanding the chiral symmetry structure (see, e.g., Ref. [10] for a review). The medium-modified heavy–light meson spectrum has been studied by several groups in the literature [11–17]. In Ref. [15], it was shown that the D meson ($J^P = 0^-$) is mixed with a D^* meson ($J^P = 1^-$) in the spin–isospin-correlated matter, in which the mixing strength reflects the strength of the correlation. In Refs. [11,16,17], by regarding the D_0^* ($J^P = 0^+$) and D_1 ($J^P = 1^+$) mesons as chiral partners to the D and D^* mesons, it was shown that the mass splitting of the chiral partner is reduced at high density and temperature. In particular, in Ref. [17], by replacing the chiral field for pions interacting with the heavy mesons with its mean-field value obtained in the nuclear matter created by the skyrmion crystal approach [4], it was shown that the masses of D and D^* increase with density while the masses of D_0^* and D_1 decrease, and that their masses approach the average value. In other words, the degenerated mass (actually, the difference between the degenerated mass and the heavy quark mass) agrees with the chiral invariant mass, which is given by the average at vacuum. However, in the analyses of Refs. [15,17], only the pion is included in the light hadron sector, and effects of other mesons are not included. In particular, the analysis in Ref. [12] shows that the ω meson increases the mass of the D meson, while it decreases that of the $\bar{D}$ meson.

In this paper, we study the effects of the ω meson as well as the σ meson on the density dependence of the effective masses of heavy–light mesons. We show that the effect of the σ meson increases the masses of the (D, D^*) heavy quark doublet, while it decreases the masses of the chiral partners, i.e., the (D_0^*, D_1) doublet; this is similar to the analysis in Refs. [16,17]. On the other hand, the effect of the ω meson increases the masses of both doublets. Nevertheless, the difference between the masses of chiral partners decreases proportionally to the mean-field value of the σ meson, which reflects the partial chiral symmetry restoration. As a result, the masses of the (D, D^*) and (D_0^*, D_1) doublets approach a certain degenerate value. Unlike the previous analysis, the degenerate value does not agree with the average value at vacuum, which is the chiral invariant mass of these doublets. In the following analysis, after general consideration, we consider two concrete models, the parity doublet model [18] and skyrmion crystal model based on hidden local symmetry [4], to give quantitative results.

2. Framework

To explain the main point explicitly, we work in the heavy quark limit and consider a simple chiral effective model for a heavy meson multiplet of charmed mesons with $J^P = 0^-$, 1^-, 0^+, and 1^+ based on the chiral doubling structure [8,9]. Let H and G denote the heavy quark doublets of heavy–light mesons with the expression

$$H = \frac{1+v^\mu\gamma_\mu}{2}\left[D^*_\mu\gamma^\mu + iD\gamma_5\right],$$
$$G = \frac{1+v^\mu\gamma_\mu}{2}\left[D^*_0 - i\gamma^\mu D'_{1\mu}\gamma_5\right], \tag{1}$$

where v^μ is the velocity of the heavy–light mesons, and D, D^*_μ, D^*_0, and $D'_{1\mu}$ are corresponding meson fields. We introduce the chiral fields $\mathcal{H}_{L,R}$ as

$$\mathcal{H}_R = \frac{1}{\sqrt{2}}\left[G + iH\gamma_5\right]\ ,\ \ \mathcal{H}_L = \frac{1}{\sqrt{2}}\left[G - iH\gamma_5\right], \tag{2}$$

which transform linearly under the chiral symmetry: $\mathcal{H}_{R,L} \to \mathcal{H}_{R,L} g^\dagger_{R,L}$ with $g_{R,L} \in \mathrm{SU}(2)_{R,L}$.

The relevant Lagrangian used in the present calculation is expressed as [17,24][1]

$$\begin{aligned}\mathcal{L} = {}& \mathrm{tr}\left[\mathcal{H}_L(iv\cdot\partial)\bar{\mathcal{H}}_L\right] + \mathrm{tr}[\mathcal{H}_R(iv\cdot\partial)\bar{\mathcal{H}}_R] \\ & - g_{\omega DD}\,\mathrm{Tr}\left[\mathcal{H}_L v^\mu\omega_\mu\bar{\mathcal{H}}_L + \mathcal{H}_R v^\mu\omega_\mu\bar{\mathcal{H}}_R\right] \\ & + \frac{\Delta_M}{2f_\pi}\mathrm{tr}\left[\mathcal{H}_L M\bar{\mathcal{H}}_R + \mathcal{H}_R M^\dagger\bar{\mathcal{H}}_L\right] \\ & - i\frac{g_A}{2f_\pi}\mathrm{tr}\left[\mathcal{H}_R\gamma_5\gamma^\mu\partial_\mu M^\dagger\bar{\mathcal{H}}_L - \mathcal{H}_L\gamma_5\gamma^\mu\partial_\mu M\bar{\mathcal{H}}_R\right],\end{aligned} \tag{3}$$

where Δ_M is the mass difference between G and H doublets, f_π is the pion decay constant, and g_A is a dimensionless real parameter. In the above Lagrangian, the omega meson field ω_μ is introduced as a chiral singlet and the field M is parametrized as $M = \sigma + i\sum_{a=1}^3 \pi_a\tau_a$ with the Pauli matrix τ_a, which transforms as $M \to g_L M g^\dagger_R$. We rewrite the effective Lagrangian (3) in terms of the H and G fields as

$$\begin{aligned}\mathcal{L} = {}& \mathrm{tr}\left[Gv^\mu\left(i\partial_\mu + g_{\omega DD}\omega_\mu\right)\bar{G} - Hv^\mu\left(i\partial_\mu + g_{\omega DD}\omega_\mu\right)\bar{H}\right] \\ & + \frac{\Delta_M}{4f_\pi}\mathrm{tr}\left[G\left(M+M^\dagger\right)\bar{G} + H\left(M+M^\dagger\right)\bar{H} - iG\left(M-M^\dagger\right)\gamma_5\bar{H} + iH\left(M-M^\dagger\right)\gamma_5\bar{G}\right] \\ & - \frac{ig_A}{4f_\pi}\mathrm{tr}\Big[G\gamma_5\left(\partial\!\!\!/ M^\dagger - \partial\!\!\!/ M\right)\bar{G} - H\gamma_5\left(\partial\!\!\!/ M^\dagger - \partial\!\!\!/ M\right)\bar{H} \\ & \qquad + iG\left(\partial\!\!\!/ M^\dagger + \partial\!\!\!/ M\right)\bar{H} - iH\left(\partial\!\!\!/ M^\dagger + \partial\!\!\!/ M\right)\bar{G}\Big].\end{aligned} \tag{4}$$

Now we replace the light meson fields by their mean-field values in medium. Here we consider symmetric matter only and assume no pion condensation, so that $\langle M\rangle = \langle\sigma\rangle$ and $\langle\partial_\mu M\rangle = 0$. Note that the mean-field value of σ at vacuum agrees with the pion decay constant, $\langle\sigma\rangle_0 = f_\pi$. From the above form, we obtain

$$\begin{aligned}\mathcal{L}_{\mathrm{eff}} = {}& \mathrm{tr}\left[G\left(i\partial_0 + g_{\omega DD}\langle\omega_0\rangle\right)\bar{G}\right] \\ & - \mathrm{tr}\left[H\left(i\partial_0 + g_{\omega DD}\langle\omega_0\rangle\right)\bar{H}\right] \\ & + \frac{\Delta_M}{2f_\pi}\langle\sigma\rangle\,\mathrm{tr}\left[G\bar{G} + H\bar{H}\right],\end{aligned} \tag{5}$$

[1] In Ref. [17], another term for the pionic interaction is included. In the present analysis we do not explicitly include the term, since one-pion interaction terms do not contribute to the following analysis.

where we have used $v^\mu = (1, \vec{0})$. This means that the effective masses of the H and G doublets are obtained as

$$m_H^{(\mathrm{eff})} = m - \frac{\Delta_M}{2f_\pi}\langle\sigma\rangle + g_{\omega DD}\langle\omega_0\rangle,$$
$$m_G^{(\mathrm{eff})} = m + \frac{\Delta_M}{2f_\pi}\langle\sigma\rangle + g_{\omega DD}\langle\omega_0\rangle, \tag{6}$$

where m is the average mass of the H and G doublets with $m = (m_H + m_G)/2$. The masses of the H and G doublets are determined by the spin average of the physical masses as

$$m_H = \frac{m_D + 3m_{D^*}}{4}, \quad m_G = \frac{m_{D_0^*} + 3m_{D_1}}{4}. \tag{7}$$

We should note that, for the anti-charmed mesons $\bar{D}$, $\bar{D}^*$, $\bar{D}_0^*$, and $\bar{D}_1$, the sign in front of the coupling to the omega meson is flipped, so that the effective masses are written as

$$m_{\bar{H}}^{(\mathrm{eff})} = m - \frac{\Delta_M}{2f_\pi}\langle\sigma\rangle - g_{\omega DD}\langle\omega_0\rangle,$$
$$m_{\bar{G}}^{(\mathrm{eff})} = m + \frac{\Delta_M}{2f_\pi}\langle\sigma\rangle - g_{\omega DD}\langle\omega_0\rangle. \tag{8}$$

Now let us study the density dependence of masses in Eqs. (6) and (8). For the mean-field value of ω, we simply take

$$\langle\omega_0\rangle = \frac{g_{\omega NN}}{m_\omega^2}\rho_B, \tag{9}$$

where $g_{\omega NN}$ is the omega meson coupling to the nucleon, m_ω is the mass of the omega meson, and ρ_B is the baryon number density. For the mean field of σ we adopt the linear density approximation as

$$\frac{\langle\sigma\rangle}{\langle\sigma\rangle_0} = 1 - \frac{\sigma_{\pi N}}{m_\pi^2 f_\pi^2}\rho_B, \tag{10}$$

where $\sigma_{\pi N}$ is the coefficient of the π–N sigma term.

To make a numerical estimation, we use $m_G = 2.40$ GeV, $m_H = 1.97$ GeV, and $\Delta_M = m_G - m_H = 430$ MeV for the masses, in addition to $m_\omega = 783$ MeV, $m_\pi = 137$ MeV, and $f_\pi = 92.1$ MeV. For the other parameters, we use $\sigma_{\pi N} = 45$ MeV, $|g_{\omega DD}| = 3.7$, estimated in Appendix A, and $|g_{\omega NN}| = 6.23$,[2] which lead to $|g_{\omega NN} g_{\omega DD}| = 23$, as reference values. We note that the D meson includes the anti-light quark and the $\bar{D}$ meson the light quark. Therefore, it is natural to consider that the $\bar{D}$ meson is affected by Pauli blocking in a dense medium, which is represented by the effect of the mean field of the ω meson. Then, for concreteness, we take $g_{\omega NN} g_{\omega DD} < 0$ below. When we study the case with $g_{\omega NN} g_{\omega DD} > 0$, we just exchange H with $\bar{H}$ and G with $\bar{G}$ in the following discussion.

[2] There are several values listed in the literature. Here, we use a value obtained in an analysis of nuclear matter based on the parity doublet model in Ref. [18], in which the saturation density, the binding energy, and the incompressibility are reproduced. Table I in that paper includes some errors, and the value $|g_{\omega NN}| = 6.23$ is the corrected value obtained for the chiral invariant mass $m_0 = 700$ MeV.

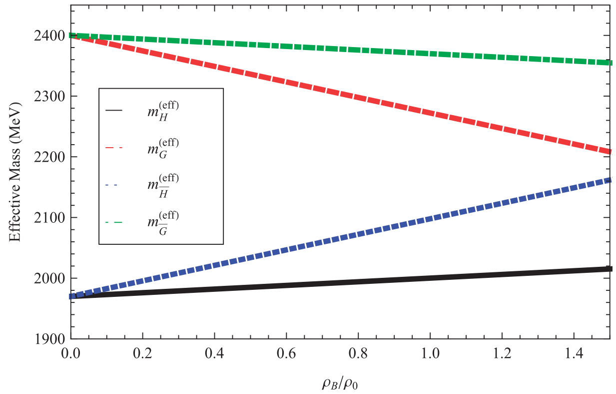

Fig. 1. Density dependence of the effective masses of the H doublet (black curve), $\bar{H}$ doublet (blue curve), G doublet (red curve), and $\bar{G}$ doublet (green curve) with $\sigma_{\pi N} = 45$ MeV and $g_{\omega DD}g_{\omega NN} = -23$.

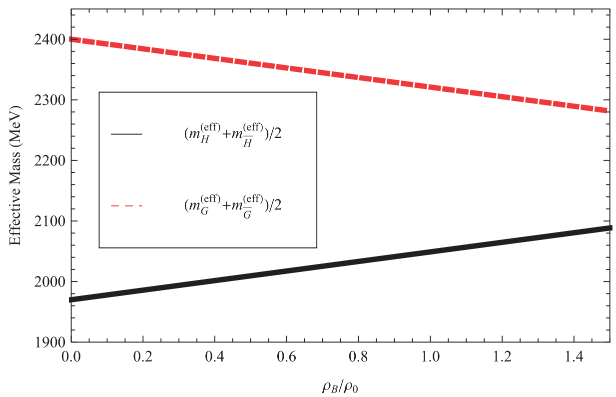

Fig. 2. Density dependence of the effective masses of charmed mesons. The dashed red and solid black curves show the sums of $m_G^{(\mathrm{eff})} + \overline{m}_G^{(\mathrm{eff})}$ and $m_H^{(\mathrm{eff})} + \overline{m}_H^{(\mathrm{eff})}$ divided by two, respectively, which do not depend on the sign of $g_{\omega DD}g_{\omega NN}$.

We plot the density dependence of the masses in Fig. 1. This shows that the masses of the H and $\bar{H}$ doublets as well as those of the G and $\bar{G}$ doublets are split by the ω contribution.

From the ω contribution combined with the σ contribution, the mass of the G doublet (indicated by the dashed red curve) decreases with increasing density, and the $\bar{H}$ mass (the dotted blue curve) increases. On the other hand, the H mass (the solid black curve) and the $\bar{G}$ mass (the dot-dashed green curve) are rather stable. As a result, the G mass tends to degenerate with the mass of the H doublet at certain high densities. If one measures the mass of H only, one might think that the chiral invariant mass would be almost the same as the mass of the H doublet. However, the actual chiral invariant mass is larger than the H mass at vacuum, which can be obtained by averaging the masses of the particles (G and H) and antiparticles ($\bar{H}$ and $\bar{G}$), as shown in Fig. 2. We should note that the sums of the particle and antiparticle masses are actually independent of the sign of $g_{\omega DD}g_{\omega NN}$.

The mass difference between the chiral partners, i.e., the H doublet and the G doublet, is caused by spontaneous chiral symmetry breaking. This structure is seen by subtracting the mass of the H doublet from that of the G doublet with Eq. (6) as

$$m_G^{(\mathrm{eff})} - m_H^{(\mathrm{eff})} = \frac{\Delta_M}{f_\pi}\langle\sigma\rangle. \tag{11}$$

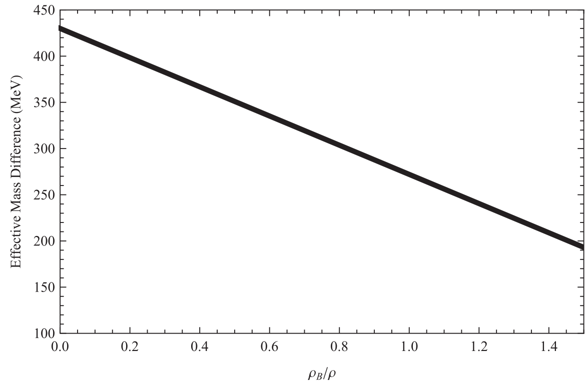

Fig. 3. Density dependence of the difference of the effective masses of charmed mesons defined by Eq. (11).

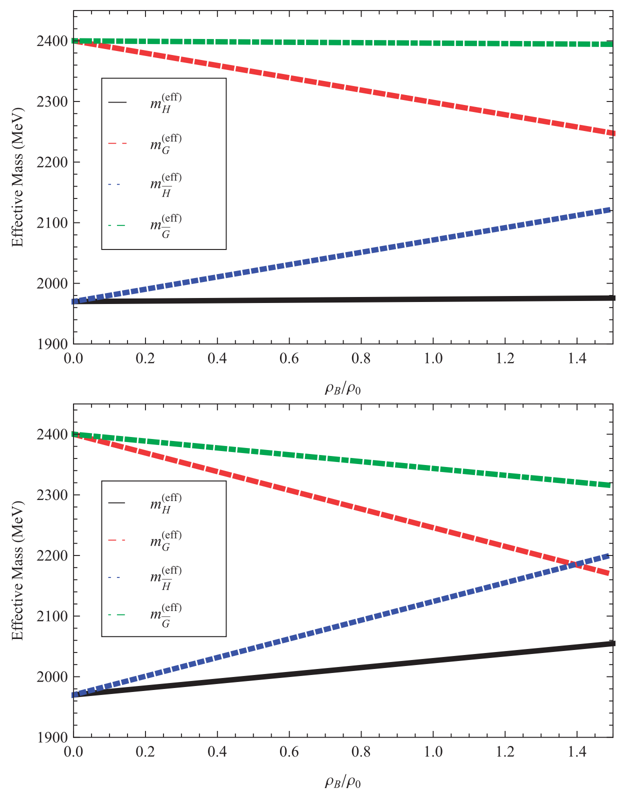

Fig. 4. Density dependence of the effective masses of charmed mesons with $\sigma_{\pi N} = 30$ MeV (upper panel) and $\sigma_{\pi N} = 60$ MeV (lower panel). Notations are the same as in Fig. 1.

So the mass difference is expected to give a clue to the chiral condensate. In the mean-field approximation, it is actually proportional to the mean field $\langle\sigma\rangle$, as shown in Fig. 3. This figure clearly shows that, with increasing nuclear matter density, the chiral symmetry is (partially) restored.

To check the π–N sigma term dependence of the effective masses, we vary the value of $\sigma_{\pi N}$ as 30 and 60 MeV, which are plotted in Fig. 4. This shows that the difference between the masses of H and G as well as that between $\bar{H}$ and $\bar{G}$ decreases more rapidly for larger values of $\sigma_{\pi N}$. As a result, the chiral symmetry is restored more rapidly for larger $\sigma_{\pi N}$.

We next check the dependence on the value of $|g_{\omega DD} g_{\omega NN}|$ in Fig. 5, by taking 30% deviation from the estimated value. This shows that the masses change more rapidly for larger values of $|g_{\omega DD} g_{\omega NN}|$.

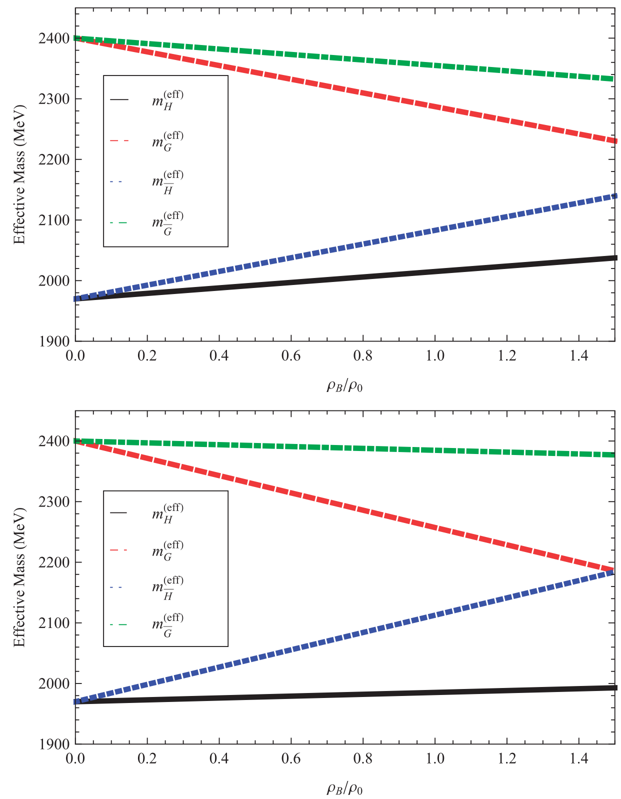

Fig. 5. Density dependence of the effective masses of charmed mesons for $\sigma_{\pi N} = 45$ MeV with $g_{\omega DD} g_{\omega NN} = -16$ (upper panel) and $g_{\omega DD} g_{\omega NN} = -30$ (lower panel). Notations are the same as in Fig. 1.

3. Model analysis

After the above general discussion, let us study the density dependences of the effective masses based on some specific models. Here we use the nuclear matter described by the parity doublet model [18] and by the skyrmion crystal model based on hidden local symmetry [4].

3.1. *Parity doublet model*

In Ref. [18], the parity doublet model based on the linear σ model [1,2], in which an excited nucleon with negative parity, $N^*(1535)$, is regarded as the chiral partner to the ordinary nucleon, was extended by including a six-point interaction for the σ field and interactions to the ω and ρ mesons based on hidden local symmetry, to study the nuclear matter. It was shown that, for a wide range of chiral invariant masses for the nucleon, the model reproduces the saturation density, binding energy, incompressibility, and symmetry energy. In Refs. [19] and [20], it is shown that the ratio of the mean field $\langle\sigma\rangle$ at normal nuclear matter density to that at vacuum obtained for the chiral invariant mass of a nucleon $m_0 = 500$ MeV is consistent with the experimental value of that for the pion decay constant [21,22]. Here we use the density dependences of $\langle\sigma\rangle$ and $\langle\omega\rangle$ obtained from the model with $m_0 = 500$ MeV.

We show the resultant density dependence of the masses in Fig. 6. Here we use $|g_{\omega DD}| = 3.7$, estimated in Appendix A, as a typical value. This shows that the density dependence of all masses in the very-low-density region $\rho_B/\rho_0 \lesssim 0.3$ is similar to that in Fig. 1, reflecting that both $\langle\sigma\rangle$ and $\langle\omega\rangle$ in the parity doublet model are consistent with those obtained in the linear density approximation, as can be seen in Ref. [19]. However, around the density region $\rho_B/\rho_0 \sim 0.3$, the mass of the H

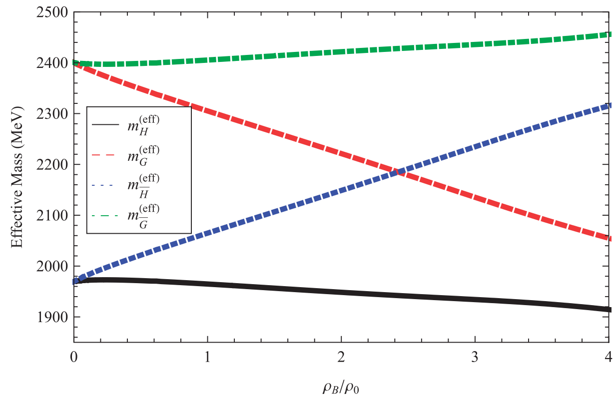

Fig. 6. Density dependence of the effective masses of charmed mesons in the parity doublet model. Notations are the same as in Fig. 1.

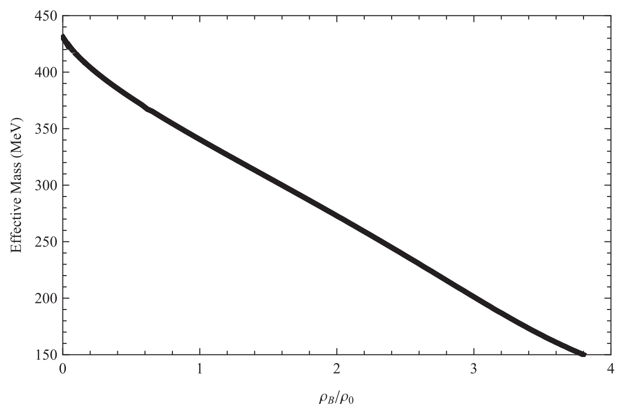

Fig. 7. Density dependence of the difference of the effective masses of charmed mesons defined by Eq. (11) in the parity doublet model.

doublet (black curve) starts to decrease and that of the G doublet (green curve) to increase, unlike the linear density approximation. In this model, the mean field $\langle\omega\rangle$ is proportional to the density, like the linear density approximation in Eq. (9). Then, the different density dependence of the masses originates from $\langle\sigma\rangle$. In Fig. 7, we plot the difference in the two masses of the H and G doublets. This shows that the difference decreases more slowly than that in the linear approximation shown in Fig. 3.

3.2. *Skyrme model*

In Refs. [4,5], the skyrmion crystal model is used to study the qualitative structure of nuclear matter by regarding the skyrmion matter as nuclear matter in the sense of the large-N_c limit of QCD. A robust conclusion drawn in the skyrmion crystal approach is that, when the density of the nuclear matter is increased, the skyrmion matter undergoes a topological phase transition to matter made of half-skyrmions, in which the space average of the chiral condensate vanishes, although it is locally non-zero and the chiral symmetry is still broken [23]. Since the half-skyrmion phase is not observed in nature, the critical density should be higher than the normal nuclear density. Recently, the description of nuclear matter from the skyrmion crystal and the implication of the topological phase transition in the equation of state of neutron stars have seen great progress (see, e.g., Ref. [7] for a recent review).

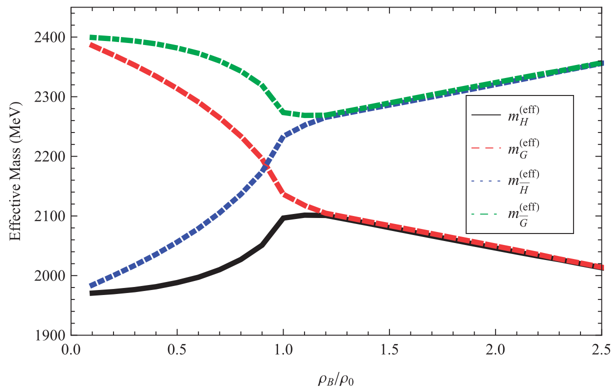

Fig. 8. Density dependence of the effective mass of charmed mesons in the skyrmion crystal model with $g_{\omega DD} = -3.7$. Notations are the same as in Fig. 1.

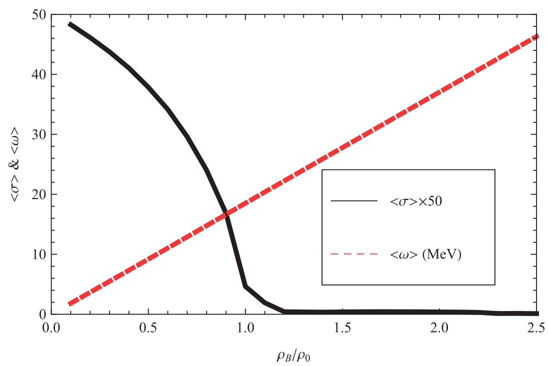

Fig. 9. Density dependence of $\langle\sigma\rangle$ and $\langle\omega\rangle$ calculated in the skyrmion crystal model.

In the present analysis, we calculate the mean fields $\langle\sigma\rangle$ and $\langle\omega\rangle$ in the skyrmion crystal model, and substitute the values into Eqs. (6) and (8) to obtain the density dependence of the charmed meson masses. We plot in Fig. 8 the density dependence of the effective masses of charmed mesons by using $\langle\sigma\rangle$ and $\langle\omega\rangle$ with $g_{\omega DD} = -3.7$ calculated by the skyrmion crystal model based on hidden local symmetry [4]; other parameters are the same as those used in the plot of Fig. 1. In this model, we find that both G and $\bar{G}$ masses decrease with density while both H and $\bar{H}$ increase with density. Because the density dependence of the $\bar{H}$ mass and the G mass is stronger than that of the $\bar{G}$ mass and the H mass, H and G as well as $\bar{H}$ and $\bar{G}$ become degenerate at density $\sim 1.2\rho_0$, at which the skyrmion phase transits to the half-skyrmion phase. This is because the mass difference between H and G as well as $\bar{H}$ and $\bar{G}$ is proportional to $\langle\sigma\rangle$; this vanishes in the half-skyrmion phase. Moreover, we find that the degenerated masses of H and G and those of $\bar{H}$ and $\bar{G}$ linearly depend on density in the half-skyrmion phase. The reason for this is that $\langle\omega\rangle$ is a linear function of density and this linear dependence agrees with Eq. (9). We plot in Fig. 9 the density dependence of $\langle\sigma\rangle$ and $\langle\omega\rangle$. Note that, as stressed above, the analysis shows just a qualitative structure, and the degeneracy of chiral partners does not imply chiral restoration but is due to the vanishing of the space average of the chiral condensate in half-skyrmion matter.

4. Summary and discussion

In this work, by regarding the (D_0^*, D_1) heavy quark doublet as the chiral partner of the (D, D^*) doublet, we have explicitly shown that the effect of the ω meson decreases the masses of both doublets, while the ($\bar{D}_0^*$, $\bar{D}_1$) and ($\bar{D}$, $\bar{D}^*$) meson masses are increased. We explicitly point out that the ω meson effect is significant for understanding the density dependence of effective hadron masses in medium. Even though the qualitative dependence is model dependent, the tendency that the masses of the heavy–light mesons and their antiparticles are split due to the ω meson effect is robust. We hope that these medium-modified masses of the heavy–light mesons can be detected in future experiments at J-PARC and FAIR through the strong and weak channels, such as $\psi(3770) \to D\bar{D}, J/\psi \to \bar{D}e^+\nu_e$, and so on (see, e.g., Ref. [25]).

We would like to note that the result of the omega meson effect on the D and $\bar{D}$ mesons is consistent with the result obtained in Ref. [12]. In our analysis, we further introduce p-wave excited D and $\bar{D}$ mesons by using the chiral doubling model, and we have found that the difference between the masses of chiral partners decreases in proportion to the mean-field value of the σ meson, which reflects the partial chiral symmetry restoration even if the ω meson contribution is present.

In our calculation, we simply take the mean-field approach. An extension of the present work to include some loop contributions is reported in Ref. [26]. In the present work, we have only discussed medium-modified charmed mesons. The results presented here are intact for their bottom cousins, except that the average mass m should be taken as the value of bottom mesons.

Appendix. Estimation of $g_{\omega DD}$

In the present analysis, we estimate a reference value of $g_{\omega DD}$ defined by Eq. (3) in the heavy hadron limit by using the following naive scaling property:

$$\left|\frac{\tilde{g}_{\omega D\bar{D}}}{\tilde{g}_{\omega K\bar{K}}}\right| = \left|\frac{\tilde{g}_{D^{*+}D^0\pi^-}}{\tilde{g}_{K^{*0}K^-\pi^+}}\right|, \tag{A.1}$$

where the coupling constants are defined in the relativistic form of the interaction Lagrangian among a vector meson V and two pseudoscalar mesons P and P' expressed as

$$\mathcal{L}_{VPP'} = i\tilde{g}_{VPP'}V^\mu\left(\partial_\mu PP' - \partial_\mu P'P\right). \tag{A.2}$$

The coupling $\tilde{g}_{D^{*+}D^0\pi^+}$ appears in the decay width of $D^{*+} \to D^0\pi^+$ as

$$\Gamma(D^{*+} \to D^0\pi^+) = \frac{\tilde{g}_{D^*D\pi}^2|\vec{p}|^3}{24\pi m_H^2} = 56.5\ \text{keV}, \tag{A.3}$$

which with $|\vec{p}| = 39.4\,\text{MeV}$ and $m_H = 1.97\,\text{GeV}$ leads to

$$|\tilde{g}_{D^*D\pi}| = 16.5. \tag{A.4}$$

In a class of three-flavor chiral models for vector mesons, $g_{\omega K\bar{K}}$ and $g_{K^{*0}K^-\pi^+}$ are related to the vector meson masses as [27]

$$g_{\omega K^+K^-} = g_{\omega K^0\bar{K}^0} = g_{\omega K\bar{K}} = \frac{1}{4}\frac{m_\omega^2}{gf_K^2},$$
$$g_{K^{*0}K^-\pi^+} = \frac{1}{2\sqrt{2}}\frac{m_{K^*}^2}{gf_Kf_\pi}, \tag{A.5}$$

where m_ω and m_{K^*} are the masses of the ω and K^* mesons, f_π and f_K are the pion and kaon decay constants, and g is the gauge coupling constant of the hidden local symmetry (see, e.g., Ref. [28]).

Using $m_\omega = 783$ MeV, $m_{K^{*0}} = 896$ MeV, $f_\pi = 92.1$ MeV, and $f_K = 110$ MeV, the ratio of the two couplings in Eq. (A.5) is estimated as

$$\left|\frac{g_{\omega K\bar{K}}}{g_{K^{*0}K^-\pi^+}}\right| = 0.452, \tag{A.6}$$

which with Eq. (A.4) leads to

$$\left|\tilde{g}_{\omega D\bar{D}}\right| = \left|\frac{\tilde{g}_{\omega K\bar{K}}}{\tilde{g}_{K^{*0}K^-\pi^+}}\tilde{g}_{D^{*+}D^0\pi^-}\right| = 7.4. \tag{A.7}$$

From the heavy quark Lagrangian in Eq. (4), the ω–D–$\bar{D}$ interaction is written as

$$\mathcal{L}_{\omega DD} = 2g_{\omega DD} D\,\omega_\mu v^\mu \bar{D}, \tag{A.8}$$

with a scaling factor of the mass of the heavy meson M_H. Comparing this with Eq. (A.1), we estimate $g_{\omega DD}$ as

$$|g_{\omega DD}| = \frac{1}{2}\left|\tilde{g}_{\omega D\bar{D}}\right| = 3.7, \tag{A.9}$$

which is the value used in the present work.

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