D mesons at finite temperature and density in the Polyakov–Nambu–Jona-Lasinio model

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We study D meson resonances in hot, dense quark matter within the Nambu–Jona-Lasinio model and its Polyakov-loop extension. We show that the mass splitting between D^+ and D^- mesons is moderate, not in excess of 100 MeV. When the decay channel into quasifree quarks opens (Mott effect) at densities above twice saturation density, the decay width reaches rapidly the value of 200 MeV which entails a spectral broadening sufficient to open J/ψ dissociation processes. Contrary to results from hadronic mean-field theories, the chiral quark model does not support the scenario of a dropping D meson masses so that scenarios for J/ψ dissociation by quark rearrangement built on the lowering of the threshold for this process in a hot and dense medium have to be reconsidered and should account for the spectral broadening.

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I. INTRODUCTION

The modification of the D meson properties (masses and widths) in hot, dense matter has consequences for scenarios of J/ψ suppression, e.g., by the processes of the type

$$J/\psi + \pi, \rho \rightarrow D^* + \bar{D}, \qquad J/\psi + N \rightarrow \Lambda_c + \bar{D}, \quad (1)$$

which couple hidden charm to open charm states and thus lead to a dissociation of charm in the medium [1], see Refs. [2–5] for early controversial estimates of the cross sections of such processes. The reverse process [6–8] of charmonium regeneration by open charm recombination should play a dominant role for J/ψ production at LHC [9,10] where charm is abundant in the medium.

Since either dropping masses [11,12] or increasing widths [13,14] of the D mesons in a hot and dense medium can lead to a lowering of the reaction threshold and thus to a strong increase of the rate for the processes (1), both effects may contribute to an explanation of the anomalous J/ψ suppression found in the NA50 experiment [15–17] and subsequently confirmed by NA60 [18,19] and PHENIX [20]. For a recent review, see [21].

In contrast to results from a relativistic mean-field model of D mesons in nuclear matter which predicts a strong downwards shift of D^+ meson masses due to the renormalization with a scalar mean field [22], the consideration of the quark substructure of these mesons leads to qualitatively different behavior. As we will show in this work on the basis of a chiral quark model of the Nambu-Jona-Lasinio (NJL) type and its Polyakov-loop extension (PNJL), the Pauli blocking effect in the Bethe-Salpeter equation for the D mesons largely compensates the dropping masses of their quark constituents. As a result, D meson masses do not drop but stay almost constant or rather increase with increasing density (and temperature) of the matter. Their decay width, however, increases rapidly and reaches values which allow for a subthreshold quark rearrangement dissociation of J/ψ. Therefore, scenarios for J/ψ suppression built on the quark rearrangement reaction (1) have to be reconsidered. As has been demonstrated in [23,24], a sufficient width of D mesonic correlations in the quark plasma is essential for understanding charm thermalization and diffusion in RHIC experiments, see [25] for a review.

It is interesting to note that a recent self-consistent coupled channel approach for D mesons in hot, dense nuclear matter [26,27] supports the picture of a spectral broadening with a negligible mass shift up to temperatures $T = 150$ MeV and densities $n = 2n_0$ with $n_0 = 0.16$ fm^{-3} being the nucleon density of nuclear matter at saturation. Therefore, it seems likely that a quark hadron duality similar to that discussed for low-mass dilepton production [28] can be observed also in the D meson channel at the deconfinement transition. The spectral

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broadening of D mesons rather than their mass shift has been suggested for an explanation of anomalous J/ψ suppression in [13,14]. This question, however, awaits a thorough investigation.

In the present note, we investigate the extension of previous exploratory calculations of D mesonic correlations in quark matter, based on the NJL model [14,29] to the domain of finite baryon densities which will become accessible in the CBM experiment at FAIR. Of particular interest will be the question whether the suggested mass splitting of D meson states [30] will be observable or rather washed out by spectral broadening. Furthermore, a strong isospin dependence of the D meson broadening could result in observable signatures, possibly relevant for quark-gluon plasma diagnostics.

Chiral dynamics has been applied successfully not only in the light quark sector but also especially for the investigation of heavy-light pseudoscalar meson properties. This has been most impressively demonstrated within the Dyson-Schwinger equation approach in Ref. [31] which reproduced a complete set of heavy meson observables. A particular feature of these systems is the heavy-quark symmetry which allows to separate the physics of the heavy and the light quark components from each other and absorb the chiral dynamics of the latter into the universal Isgur-Wise function [32]. For more details, see the reviews on heavy-quark effective theory [33–35]. It is interesting to note that already the chiral dynamics encoded in the rather schematic NJL model reproduces features like heavy-quark symmetry and Isgur-Wise function [36] when extended to the heavy-quark sector. These properties remain unaffected when confining properties are mimicked in the heavy-quark extended NJL model by an infrared cutoff procedure [37] or by a confining interquark potential in a relativistic potential model of heavy mesons [38].

D meson properties as reported in our study come from a simultaneous solution of two types of nonperturbative equations: the gap equations for the quark masses and the Bethe-Salpeter equations for the meson masses. The results for the dependence of pseudoscalar meson masses M_P on the current quark masses m_q of the heavier quark in the meson nicely reproduce the transition from a Gell-Mann–Oakes–Renner-like behavior for pion and kaon $(M_P \sim \sqrt{m_q})$ to the additive quark model like behavior $(M_P \sim m_q)$ for D and B mesons. These results are then readily generalized from the vacuum to finite temperature and chemical potential within the Matsubara formalism.

II. MODEL AND FORMALISM

We employ a four-flavor model with NJL-type interaction kernel as a straightforward generalization of recent work on the $SU_f(3)$ scalar and pseudoscalar meson spectrum [39–41] developed on the basis of Ref. [42] and its generalization by coupling to the Polyakov loop [43,44],

$$\mathcal{L} = \bar{q}(i\gamma^\mu D_\mu + \hat{m})q + G_S \sum_{a=0}^{15}[(\bar{q}\lambda^a q)^2 + (\bar{q}i\gamma_5\lambda^a q)^2]$$
$$- \mathcal{U}(\Phi[A], \bar{\Phi}[A]; T). \quad (2)$$

Here q denotes the quark field with four flavors, $N_f = 4$, $f = u, d, s, c$, and three colors, $N_c = 3$; λ^a are the flavor $SU_f(4)$ Gell-Mann matrices ($a = 0, 1, 2, \ldots, 15$), G_S is a coupling constant. The global symmetry of the Lagrangian (2) is explicitly broken by the current quark masses $\hat{m} = \text{diag}(m_u, m_d, m_s, m_c)$. The covariant derivative is defined as $D^\mu = \partial^\mu - iA^\mu$, with $A^\mu = \delta_0^\mu A^0$ (Polyakov gauge); in Euclidean notation $A^0 = -iA_4$. The strong coupling constant g_s is absorbed in the definition of $A^\mu(x) = g_s \mathcal{A}_a^\mu(x)\frac{\lambda_a}{2}$, where $\mathcal{A}_a^\mu$ is the ($SU_c(3)$) gauge field and λ_a are the (color) Gell-Mann matrices.

The Polyakov loop field Φ appearing in the potential term of (2) is related to the gauge field through the gauge-covariant average of the Polyakov line [45]

$$\Phi(\vec{x}) = \langle\langle l(\vec{x})\rangle\rangle = \frac{1}{N_c}\text{Tr}_c\langle\langle L(\vec{x})\rangle\rangle, \quad (3)$$

where

$$L(\vec{x}) = \mathcal{P}\exp\left[i\int_0^\beta d\tau A_4(\vec{x}, \tau)\right]. \quad (4)$$

Concerning the effective potential for the (complex) Φ field, we adopt the form and parametrization proposed in Ref. [46].

This effective chiral field theory has the same chiral symmetry of QCD, which is also shared by the quark interaction terms. The (P)NJL model is a primer for describing the dynamical breakdown of this symmetry in the vacuum and its partial restoration at high temperatures and chemical potentials. At the same time it provides a field-theoretic description of pseudoscalar meson properties which is in accordance with the low energy theorems (such as the Goldstone theorem) of QCD.

In the vacuum the PNJL model with the Lagrangian (2) goes over to the NJL one and the pseudoscalar meson properties are described in the standard way by analyzing the polarization operators

$$\Pi_{ij}(P) = iN_c\int\frac{d^4p}{(2\pi)^4}\text{tr}_D[S_i(p)(i\gamma_5)S_j(p+P)(i\gamma_5)], \quad (5)$$

where tr_D is the trace over Dirac matrices, $S_i(p)$ is the quark Green function with the dynamical quark mass M_i.

The polarization operators can be presented in terms of two integrals which for mesons at rest in the medium are given by

$$\Pi^{ij}(P_0) = 4\{(I_1^i + I_1^j) - [P_0^2 - (M_i - M_j)^2]I_2^{ij}(P_0)\}, \quad (6)$$

where

$$I_1^i = iN_c \int \frac{d^4p}{(2\pi)^4} \frac{1}{p_0^2 - E_i^2} = \frac{N_c}{4\pi^2} \int_0^\Lambda \frac{p^2 dp}{E_i}, \quad (7)$$

$$I_2^{ij}(P_0) = iN_c \int \frac{d^4p}{(2\pi)^4} \frac{1}{(p_0^2 - E_i^2)((p_0 + P_0)^2 - E_j^2)}$$
$$= \frac{N_c}{4\pi^2} \int_0^\Lambda \frac{p^2 dp}{E_i E_j} \frac{E_i + E_j}{P_0^2 - (E_i + E_j)^2}, \quad (8)$$

where $E_{i,j} = \sqrt{p^2 + M_{i,j}^2}$ is the quark energy.

As the Lagrangian (2) defines a nonrenormalizable field theory, we introduce the three-momentum cutoff with the parameter Λ to regularize the integrals. When $P_0 > M_i + M_j$ it is necessary to take into account the imaginary part of the second integral. It may be found, with help of the $i\epsilon$–prescription $P_0^2 \to P_0^2 - i\epsilon$, that

$$I_2^{ij}(P_0) = \frac{N_c}{4\pi^2} \mathcal{P} \int_0^\Lambda \frac{p^2 dp}{E_i E_j} \frac{E_i + E_j}{P_0^2 - (E_i + E_j)^2}$$
$$+ i\frac{N_c}{16\pi} \frac{p^*}{(E_i^* + E_j^*)}, \quad (9)$$

where $p^* = \sqrt{(P_0^2 - (M_i - M_j)^2)(P_0^2 - (M_i + M_j)^2)}/2P_0$ is the momentum and $E_{i,j}^* = \sqrt{(p^*)^2 + M_{i,j}^2}$ the corresponding energy.

The quark mass M_i we find from the gap equation

$$M_i = m_i + 16 M_i G_S I_1^i. \quad (10)$$

The meson mass spectrum we obtain from the condition

$$1 - 2G_S \Pi^{ij}(P_0 = M_P, \mathbf{P} = 0) = 0. \quad (11)$$

The pseudoscalar meson-quark-antiquark coupling constants are defined as

$$g_{P\bar{q}_i q_j}^{-2} = -\frac{1}{2M_P} \frac{\partial}{\partial P_0} [\Pi^{ij}(P_0)]|_{P_0 = M_P}. \quad (12)$$

Note that when $P_0 > M_i + M_j$, then Eq. (11) has to be solved in their complex form in order to determine the mass of the resonance M_P and the respective decay width Γ_P. Thus, we assume that this equation can be written as a system of two coupled equations

$$-\left[M_P^2 - \frac{1}{4}\Gamma_P^2 - (M_i - M_j)^2\right]$$
$$= \frac{(8G_S)^{-1} - (I_1^i + I_1^j)}{|I_2(P_0 = M_P + i\epsilon)|^2} \mathrm{Re} I_2(P_0 = M_P + i\epsilon), \quad (13)$$

$$-[M_P \Gamma_P] = \frac{(8G_S)^{-1} - (I_1^i + I_1^j)}{|I_2(P_0 = M_P + i\epsilon)|^2} \mathrm{Im} I_2(P_0 = M_P + i\epsilon), \quad (14)$$

which have solutions of the form

TABLE I. Results for pseudoscalar meson properties in the light, strange, charm and bottom sectors for the corresponding current quark masses and dynamically generated quark masses of the model in the vacuum at $T = \mu = 0$.

flavor	P	M_P [MeV]	f_P [MeV]	m_f [MeV]	M_f [MeV]
u, d	π	135.0	92.4	5.5	368.0
s	K	497.7	95.4	140.7	587.4
c	D	1869.3	79.6	1279.9	1828.8
b	B	5279.4		4634.8	

$$P_0 = M_P - i\frac{1}{2}\Gamma_P. \quad (15)$$

As shown in [39], the model (2) successfully describes meson properties in the vacuum at $T = \mu = 0$. We use here the parametrization given in Table III of Ref. [47] for the case of the NJL model. The value of the coupling constant is $G_S\Lambda^2 = 2.32$, the three-momentum cutoff is at $\Lambda = 602.3$ MeV; see also [48] for an online tool for the three-flavor NJL model. The parametrization of the current quark masses for the heavy flavors (c, b) is performed with the above formulas for the masses of the corresponding heavy-light pseudoscalar mesons (D, B). The results are summarized in Table I. The dependence of the heavy-light meson mass on the current mass of the heavier quark is shown by the solid line in Fig. 1 and provides satisfactory agreement with the particle data group listings [49].

The generalization of the model (2) to the case of non-zero temperature and density (details in [39]) is done within the imaginary time formalism by introducing the
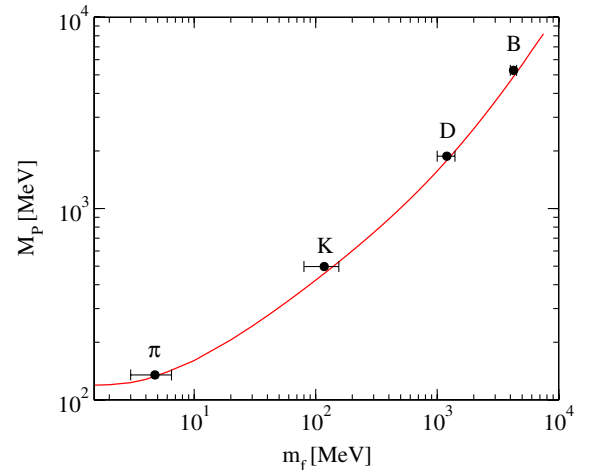

FIG. 1 (color online). The systematics of heavy-light pseudoscalar meson masses M_P in dependence on the mass m_f of the heavy flavor in the bound state with a light antiquark within the present model (red solid line) is compared to the masses of heavy-light pseudoscalar mesons (π, K, D, B) and the limits for the corresponding current quark masses (symbols with error bars) according to the particle data group [49].

Matsubara frequencies $\omega_n = (2n+1)\pi T$, $n = 0, \pm1, \pm2, \ldots$, so that $p_0 \to i\omega_n + \mu$ with the chemical potential μ and the temperature T. Instead of the integration over p_0, we have now to perform a sum over Matsubara frequencies. In the result we obtain the integrals

$$I_1^i = -\frac{N_c}{4\pi^2} \int \frac{\mathbb{p}^2 d\mathbb{p}}{E_i} (n_i^+ - n_i^-), \tag{16}$$

$$
\begin{aligned}
I_2^{ij}(P_0, T, \mu) = -N_c \int \frac{d^3\mathbf{p}}{(2\pi)^3} & \left[\frac{1}{2E_i} \frac{1}{(E_i + P_0)^2 - E_j^2} n_i^+ \right. \\
& - \frac{1}{2E_i} \frac{1}{(E_i - P_0)^2 - E_j^2} n_i^- \\
& + \frac{1}{2E_j} \frac{1}{(E_j - P_0)^2 - E_i^2} n_j^+ \\
& \left. - \frac{1}{2E_j} \frac{1}{(E_j + P_0)^2 - E_i^2} n_j^- \right],
\end{aligned}
\tag{17}
$$

where $n_i^\pm = f_\Phi(\pm E_i)$ are the generalized fermion distribution functions [44] for quarks of flavor i with positive (negative) energies in the presence of the Polyakov-loop values Φ and $\bar{\Phi}$

$$f_\Phi(E) = \frac{\bar{\Phi}e^{-\beta(E-\mu)} + 2\Phi e^{-2\beta(E-\mu)} + e^{-3\beta(E-\mu)}}{1 + 3(\bar{\Phi} + \Phi e^{-\beta(E-\mu)})e^{-\beta(E-\mu)} + e^{-3\beta(E-\mu)}}, \tag{18}$$

which go over to the ordinary Fermi functions in the case of the NJL model, where $\Phi = \bar{\Phi} = 1$

$$f_1(\pm E) = \frac{1}{1 + e^{\beta(\pm E - \mu)}}. \tag{19}$$

Note that we put the chemical potential for charm quarks to zero in the calculations discussed below. A small finite value might be necessary to ensure exactly vanishing net charm at high baryon density when, as we demonstrate below, the symmetry between masses of D mesons with charm and those with anticharm is broken by medium effects.

III. RESULTS

We have performed self-consistent solutions of the gap equations for the dynamically generated light quark masses at finite temperatures and chemical potentials for the PNJL model and its NJL model limit within the standard setting as summarized above. In Fig. 2 we display the results for the restoration of the approximate chiral symmetry in the u, d quark sector along trajectories in the QCD phase diagram with a constant ratio $r = T/\mu$, where $r = 0, 1/3, 1/2, 1$. The (pseudo)critical temperatures for the NJL model and for the PNJL model are shown in the phase diagram in Fig. 3 together with the trajectories along which we investigate the D meson properties.
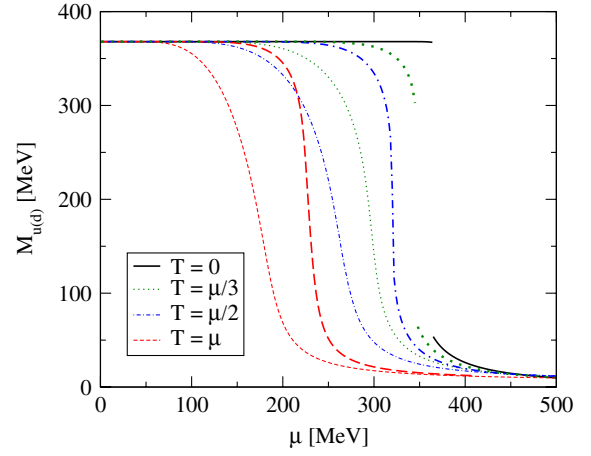

FIG. 2 (color online). Dynamically generated light quark mass $M_{u(d)}$ as a function of the quark chemical potential along trajectories in the quark matter phase diagram for $T = 0$ (black solid line), $T = \mu/3$ (green dotted line), $T = \mu/2$ (blue dash-dotted line) and $T = \mu$ (red dashed line). Results of the NJL model (thin lines) are compared to those of the PNJL model. They coincide for $T = 0$.

When lowering the ratio $r \to 0$, the phase transition turns from crossover to first order. The chiral restoration is a result of the phase space occupation (Pauli blocking) which effectively reduces the interaction strength in the gap equation.

Naïvely, one would expect that heavy-light mesons such as D mesons, should suffer a mass reduction when embedded in a hot and dense medium, as a result of the strong reduction of the light quark constituent mass (the charm quark mass is approximately unaffected). The solution of
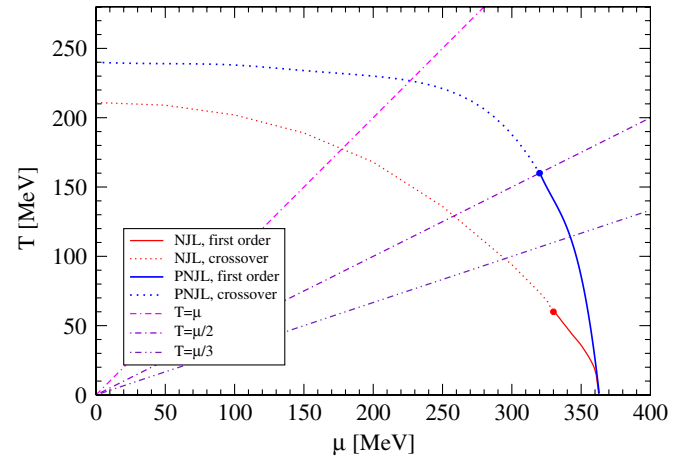

FIG. 3 (color online). (Pseudo)critical temperatures versus quark chemical potential for the NJL model and for the PNJL model. The first-order transition lines are shown as solid lines, the crossover is indicated by dashed lines. Critical end points are denoted as full dots. The straight lines stand for the trajectories $T = r\mu$ along which we investigate the D meson properties: $r = 1, 1/2, 1/3$.
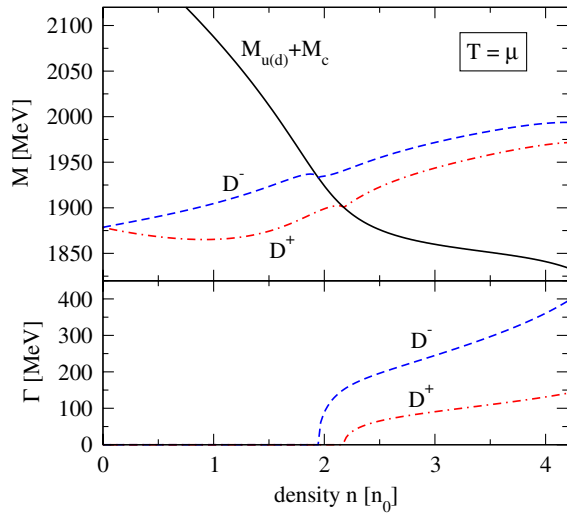
FIG. 4 (color online). Density dependence of *D* meson masses (upper panel) and widths (lower panel) along the trajectory $T = \mu$ in the QCD phase diagram for the NJL model case. While the *D* meson bound state masses are almost constant or moderately rise with increasing density the continuum threshold $(M_{u(d)} + M_c)$ gets lowered dramatically due to the chiral symmetry restoration transition. At about twice nuclear density, the decay channel into their quark-antiquark constituents opens and the bound states become resonant scattering states in the quark-antiquark continuum (Mott effect).

the in-medium D^- meson Bethe-Salpeter equation (BSE), however, shows a different result, displayed in Figs. 4 and 5 for the NJL case. This can be understood from inspecting the kernel of the BSE: The mass shift of the light quarks, which affects the energy denominators and would lead to a lowering of the meson masses, when evaluated in free space, gets compensated or even overcompensated by the Pauli blocking factor in the numerator. It is clear that *D*
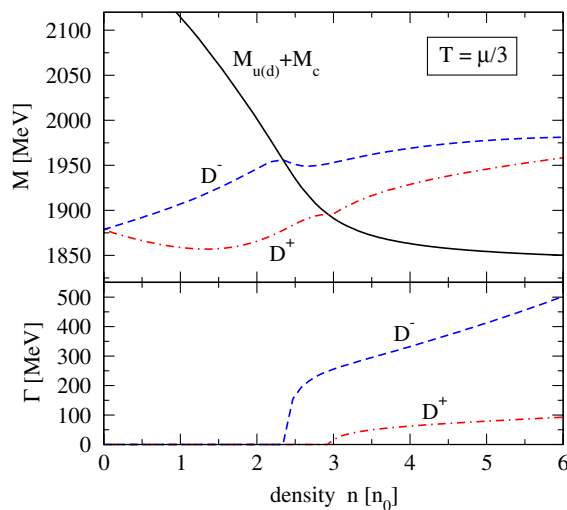
mesons containing a light quark, as the dominant quark species in a dense medium, feel a stronger Pauli blockig than those containing light antiquarks. Therefore, the D^- mass is shifted upwards while the D^+ mass stays approximately constant.

There is an important consequence of this fact that the heavy-light continuum threshold is lowered while the bound state masses are not: at a critical density the bound states merge the continuum (Mott effect) and become unstable against decay into their quark constituents, as signalled by the nonvanishing decay widths shown in the lower panels of Figs. 4 and 5. The stronger Pauli blocking for the D^- mesons leads not only to a lower Mott density, but also to a larger decay width as compared to the D^+. From these figures one can also read off the critical densities where the phenomenon of *D* meson Mott effect could be expected for the NJL model.

In Figs. 6 and 7 we present results as a function of the quark chemical potential in isospin-symmetric matter for the NJL model and its extension with coupling to the Polyakov loop. The results are rather similar, except for the smaller width of the transition region in the PNJL case and the different position of the critical point in the phase diagram. The higher temperature of the critical endpoint entails that along the trajectory for $r = 1/3$ the medium undergoes a first-order transition and the *D* meson properties change discontinuously.

The modification of *D* meson properties in hot and dense nuclear matter as reported in the present work is essentially different from the one suggested in Ref. [11], where a lowering of the *D* meson mass had been conjectured with consequences for charmonium dissociation in heavy-ion collision experiments. In this reference the Pauli blocking effect was neglected. It is interesting to note that the Pauli blocking effect occurs not only on the quark level but
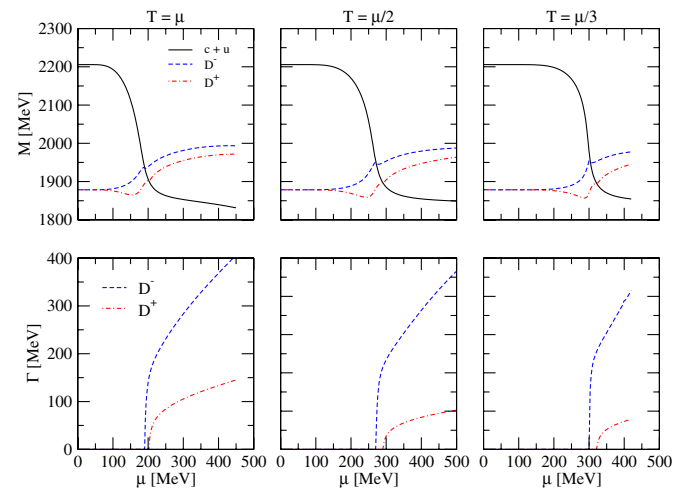

FIG. 5 (color online). Same as Fig. 4 along the trajectory $T = \mu/3$.

FIG. 6 (color online). *D* meson masses and widths along the trajectories $T = \mu$, $T = \mu/2$ and $T = \mu/3$ in the QCD phase diagram for the NJL model.
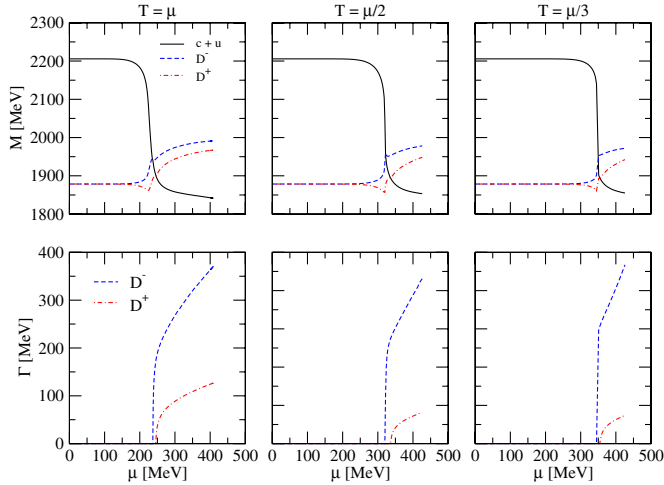
FIG. 7 (color online). Same as Fig. 6 for the PNJL model.

also on the hadronic level of description, when coupled channel equations for D mesons in nuclear matter are solved self-consistently as, e.g., in Ref. [26]. Also in this approach the D meson mass remains almost unshifted while a considerable spectral broadening is obtained under similar conditions of density and temperature as considered in the present work.

IV. CONCLUSIONS

In the isospin-symmetric quark matter case holds $M_{D^+} = M_{D^0}$ and $M_{D^-} = M_{\bar{D}^0}$.

The D mesons containing light quarks ($D^- = \bar{c}d$, $\bar{D}^0 = \bar{c}u$) suffer a positive mass shift due to the effective reduction of the coupling by Pauli blocking, since the phase space is occupied by light quarks abundant in the medium. The D mesons composed of light antiquarks ($D^+ = \bar{d}c$, $D^0 = \bar{u}c$) suffer no Pauli shift since there are

no antiquarks in the medium. Their mass at the Mott transition density is approximately the same as in vacuum.

It is interesting to compare the decay widths of the D mesons into their quark constituents. Because of the repulsive Pauli shift, the D^- and $\bar{D}^0$ are clearly above the threshold already at the first-order chiral phase transition and have a non-negligible decay width, whereas the D^+ and D^0 mesons are still very good resonances with negligible decay width at the transition.

The effect of coupling the chiral quark model to the Polyakov loop in the PNJL model is an effective suppression of the quark distribution functions in the BSE (13) and (14) as long as $\Phi \ll 1$ which leads to a narrowing [50] of the chiral transition region where medium effects on the D meson properties occur in the present model. In this region the dissociation of D mesons (Mott effect) occurs and is signalled by an increase in the spectral width of these states.

Summarizing the results of this model, we find no support for dropping D meson masses in the vicinity of the chiral/deconfinement phase transition in hot and dense matter but rather a strong spectral broadening. Therefore, scenarios of J/ψ suppression in dense matter via dissociation processes like (1), which are built on increasing widths [13,14] for the D mesons should be conceptually preferable over those built on dropping masses.

ACKNOWLEDGMENTS

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