Collective modes in hadronic matter in a relativistic model with medium dependent coupling

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The propagation of the lightest mesons in hadronic matter is studied at finite temperature, by including Dirac sea effects. The meson-nucleon interaction is described in a field theory model, which exhibits a residual interaction beyond the ground state, with a medium dependent coupling constant. The effective meson masses, the low lying collective excitations of nuclear matter, and the giant isoscalar resonances of a system with finite particle number are studied and compared with other theoretical predictions. The effect of different regularization prescriptions is also considered.

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I. INTRODUCTION

Properties of hadrons are expected to vary considerably when immersed in a dense and/or hot medium. The study of the density dependence of meson masses can provide information about the underlying strong interaction, as has been suggested in [1]. In that paper a qualitative expression, the so-called Brown-Rho universal scaling law, is given for the behavior of the in-medium hadronic masses, valid near the chiral transition point. In later works [2,3] a relationship between the chiral picture and the hadronic language is proposed. According to the scaling law, all hadrons masses decrease approximately at the same rate as the system approaches the chiral phase transition (with the exception of the pseudoscalar meson masses). This fact could explain experimental results, for example on dilepton production. Therefore it would be desirable that hadronic models that work well far from the chiral limit, could reproduce these features asymptotically.

In this work we intend to compare meson properties under extreme conditions of density and temperature, and to study collective phenomena of nuclear matter. For this purpose we have selected a model of the quantum hadrodynamics [4], which in its former version (QHD-I) has no chiral invariance. However it has been argued [5] that the predictions of the model are intrinsically consistent with chiral symmetry. Subsequent developments have led to its present day interpretation as an effective field theory [6]. Specifically we have used the derivative scalar coupling model (DSCM) proposed by Zimanyi and Moszkowski [7] and profusely applied to describe nuclear properties [8–14].

We organize this paper by presenting the DSCM beyond the lowest order approximation in Sec. II, the evaluation of the meson propagators in the relativistic random phase approximation (RRPA) is reviewed in Sec. III, results and discussion are given in Sec. IV, and the conclusions are presented in Sec. V.

II. THE DSCM BEYOND THE MEAN FIELD APPROXIMATION

QHD was originally outlined as a renormalizable theory with a very restricted set of free parameters used to fit bulk nuclear matter properties. The ground state solution is obtained in the mean field approximation and assumed to become more and more precise as baryon density grows. However qualitative discrepancies have shown that the inclusion of vacuum contributions is essential even at the lowest order. Furthermore technical difficulties in the summation of higher order diagrams have spoiled the program of systematic corrections [15]. Despite these conceptual inconsistencies the successes obtained in the description of nuclear phenomena have encouraged its development, until its current interpretation in the framework of effective field and density functional theories [4,6].

In this work we have selected a quantum hadrodynamics model which uses the same degrees of freedom as the QHD-I, but the nucleon-scalar meson interaction is given in a nonpolynomic parametrization. It has been used to study many body effects in several applications [8–11], related to an effective quark description of hadronic properties [12], and extended to include tensor coupling [14]. The DSCM has two important features which distinguish it from the QHD-I. First it is nonrenormalizable *ab initio* and there is no immediate way to introduce vacuum corrections to the ground state. This state is obtained in the mean field approximation, and the main properties of nuclear matter are successfully described. A possible interpretation of this fact is that DSCM is more efficient than QHD-I for describing low-energy hadron physics, at this order of approximation. Secondly, a residual interaction can be extracted beyond the lowest order solution, whose strength monotonically decreases with baryon density [13]. This fact ensures the ground state predominance at high density as assumed in quantum hadrodynamics. The properties just enumerated motivated us to use DSCM to study meson propagation in extreme baryon density and temperature. In a first approach we look for a qualitative description and consider only nucleons, scalar meson (σ), and vector meson (ω_μ) fields. But resonances [10], hyperons [11], and heavier mesons can also be included in a more realistic treatment.

The DSCM model consists of nucleon and meson fields in interaction, the simplest version [7] has a Yukawa type N-ω coupling and a N-σ nonpolynomic term $\mathcal{L}_{N\sigma}$,

$$\mathcal{L}_{DSCM}=\bar{\psi}\left(i\not{\partial}-\frac{M}{1+g_s\sigma/M}-g_v\not{\omega}\right)\psi+\frac{1}{2}\partial^\mu\sigma\partial_\mu\sigma -\frac{1}{2}m_s^2\sigma^2-\frac{1}{4}F^{\mu\nu}F_{\mu\nu}+\frac{1}{2}m_v^2\omega^\mu\omega_\mu, \quad (1)$$

where $\psi(x)$, is the nucleon field, M is the nucleon mass and g_s, g_v are a dimensional coupling constants. As usual in quantum hadrodynamics the ground state for homogeneous infinite matter is described in a mean field approximation, considering the meson fields as classical quantities and assimilating them to effective nucleon properties. Thus we can separate c-number contributions

$$\sigma=\sigma_0+s, \tag{2}$$

$$\omega_\mu=\omega_0\delta_{\mu 0}+w_\mu, \tag{3}$$

where σ_0,ω_0 are the classical mean field values and s,w_μ the quantum fluctuations which are not included in the ground state. Expressions for σ_0 and ω_0 can be evaluated by taking statistical mean values of the Euler-Lagrange equations for the mesons and requiring self-consistency. In this way we obtain

$$\sigma_0=\frac{g_s/m_s^2}{(1+g_s\sigma_0/M)^2}\langle\overline{\psi}(x)\psi(x)\rangle, \tag{4}$$

$$\omega_0=\frac{g_v}{m_v^2}\langle\psi^\dagger(x)\psi(x)\rangle. \tag{5}$$

Here $\psi(x)$ denotes the ground state solution for the nucleon field which in turn depends on σ_0 through the effective nucleon mass $M^*=M/(1+g_s\sigma_0/M)$. In Eq. (5) the term between angular brackets represents the conserved baryon density.

A residual nucleon-σ interaction arises beyond the mean field approximation [13] by introducing Eq. (2) into Eq. (1), thus the N-σ coupling becomes

$$\mathcal{L}_{N\sigma}=-\overline{\psi}(x)M^*\psi(x)+\mathcal{L}_{res}, \tag{6}$$

$$\mathcal{L}_{res}=g_s^*\frac{\overline{\psi}(x)s\psi(x)}{1+m^*\dfrac{g_s s}{M}}, \tag{7}$$

where $m^*=M^*/M$ and $g_s^*=g_s m^{*2}$ is an effective coupling constant. Thus $\mathcal{L}_{N\sigma}$, besides the nucleon effective mass term, generates a residual interaction $\mathcal{L}_{res}$ with a medium dependent coupling constant. Variable couplings arise naturally in nuclear many-body calculations [16,17], and on the other hand this is an expected feature whenever hadron substructure becomes relevant [18–20]. The expression for $\mathcal{L}_{res}$ is nonpolynomic and cannot be directly used in a diagrammatic expansion, although restriction to the physical regime for which quantum fluctuations are negligible as compared to the mean value σ_0, enables us to approximate $\mathcal{L}_{res}\simeq g_s^*\overline{\psi}(x)s\psi(x)$. This linearized version can be used to study the bosonic excitations of the system as it approaches an eventual chiral phase transition.

III. THE MESON PROPAGATORS IN THE RRPA

Corrections to the free meson propagators can be introduced by means of the relativistic random phase approximation. This is a nonperturbative scheme of solution which uses the Dyson-Schwinger equation to sum up the lowest order nucleon ring diagrams. Due to σ-ω mixing it is more adequate to unify scalar and vector propagators in a five dimensional matrix notation:

$$\mathcal{D}^0_{ab}=\begin{pmatrix}\Delta^0 & 0\\ 0 & D^0_{\mu\nu}\end{pmatrix}^{ab}, \tag{8}$$

where $a,b=-1,4$ and $\mu,\nu=0,4$. The index -1 is assigned to scalar meson contribution and $\mu=0,4$, to the μ component of the vector field. The free meson propagators are $\Delta^0(p)=1/(\mathrm{p}^2-m_s^2+i\varepsilon)$, $D^0_{\mu\nu}(p)=\xi_{\mu\nu}/(\mathrm{p}^2-m_v^2+i\varepsilon)$, with $\mathrm{p}^2=p_\mu p^\mu$ and $\xi_{\mu\nu}=p_\mu p_\nu/m_v^2-g_{\mu\nu}$. The nucleon propagator at finite temperature T in the mean field approximation is given by

$$G^0_{\mu\nu}(p)=(\not{p}+M^*)_{\mu\nu}\left[\frac{1}{\mathrm{p}^2-M^{*2}+i\varepsilon}+2\pi i\delta(\mathrm{p}^2-M^{*2})n_F(p_0)\right], \tag{9}$$

the first term between square brackets is usually labeled the Feynman component $G_F(p)$ and the second one the density dependent part $G_D(p)$, although $G_F(p)$ also depends on the medium properties through M^*. Here $n_F(p_0)$ is the Fermi occupation number $n_F(p_0)=\Theta(p_0)n_+(p_0)+\Theta(-p_0)n_-(p_0)$, $n_\pm(p_0)=(1+e^{\beta(|p_0|\pm\mu)})^{-1}$ with μ the chemical potential which takes into account the conservation of the baryon density and $\beta=1/T$.

The Schwinger-Dyson equation for the full meson propagator $\mathcal{D}(p)$ reads $\mathcal{D}_{ab}=\mathcal{D}^0_{ab}+\mathcal{D}^0_{ac}\Pi^{cd}\mathcal{D}_{db}$, with Π_{ab} the proper polarization insertion. This equation can formally be solved to give

$$\mathcal{D}_{ab}=\mathcal{U}^{-1}_{ac}\mathcal{D}^{0c}{}_b, \tag{10}$$

where $\mathcal{U}_{ab}=g_{ab}-\mathcal{D}^0_{ac}\Pi^c{}_b$. In the RRPA approximation only the nucleon ring diagrams are included in the polarization Π_{ab}. Due to baryon current conservation its components are linked by the relation $p_\mu\Pi^{\mu\nu}=0$. Furthermore if we select the external momentum $p_\mu=(p_0,p,0,0)$, then isotropy avoids the mixing of transversal components among them and with scalar, temporal, and longitudinal components. There are only four independent elements which we select as $\Pi_v=\Pi_{00}$, $\Pi_M=\Pi_{-10}$, $\Pi_s=\Pi_{-1-1}$, and $\Pi_T=\Pi_{22}=\Pi_{33}$. Within this choice we find that det $\mathcal{U}$ is factorized into two terms det $\mathcal{U}=\epsilon_L\epsilon_T^2$,

$$\epsilon_T=1+\Pi_T D^0_{00}, \tag{11}$$

$$\epsilon_L=(1-\Pi_s\Delta^0)(1+\mathrm{p}^2\Pi_v D^0_{00}/p_0^2)+\mathrm{p}^2\Pi_M^2\Delta^0 D^0_{00}/p_0^2. \tag{12}$$

For practical purposes the term $p_\mu p_\nu/m_v^2$ is not considered in the vector meson propagator in Eqs. (11) and (12), since it gives null contribution due to baryon current conservation.

The formulas for the polarization components are

$$\Pi_v(p) = -ig_v^2\delta\int\frac{d^4q}{(2\pi)^4}\mathrm{Tr}[G^0(q)\gamma_0 G^0(p+q)\gamma_0],$$

$$\Pi_M(p) = ig_v g_s^*\delta\int\frac{d^4q}{(2\pi)^4}\mathrm{Tr}[G^0(q)\gamma_0 G^0(p+q)],$$

$$\Pi_s(p) = -ig_s^{*2}\delta\int\frac{d^4q}{(2\pi)^4}\mathrm{Tr}[G^0(q)G^0(p+q)],$$

$$\Pi_T(p) = -ig_v^2\delta\int\frac{d^4q}{(2\pi)^4}\mathrm{Tr}[G^0(q)\gamma_2 G^0(p+q)\gamma_2],$$

where the trace runs over the Lorentz indices. The isospin degeneration is taken into account by the factor $\delta = 1,2$, for neutron or nuclear matter, respectively. The evaluation of these expressions using Eq. (9), splits them into several terms containing the products $G_F(q)G_F(p+q)$, $G_F(q)G_D(p+q)$, and $G_D(q)G_D(p+q)$, which represents particle-antiparticle, Pauli blocking, and particle-hole effects, respectively. From these, only the first one diverges and requires some regularization procedure. We restrict ourselves to extract divergences by a well defined method and eventually additional free parameters could be fixed using experimental data. In the framework of dimensional regularization [21] it can be demonstrated that the finite part of Feynman contribution to Π_M is identically zero, while for the remaining three we impose null contribution at zero density and at the meson mass shell, $\mathsf{p}^2 = m_{s,v}^2$ for the regularization point. Thus we obtain the following results for the finite parts:

$$\Pi_v^{fin}(p) = \frac{g_v^2\delta}{2\pi^2}p^2\int_0^1 dz\, z(1-z)\ln\left[\frac{M^{*2}-z(1-z)\mathsf{p}^2}{M^2-z(1-z)m_v^2}\right],$$

$$\Pi_T^{fin}(p) = \frac{\mathsf{p}^2}{p^2}\Pi_v^{fin}(p),$$

$$\Pi_s^{fin}(p) = \lambda\frac{g_s^{*2}\delta}{8\pi^2}(m^{*2}m_s^2 - \mathsf{p}^2) - \frac{3g_s^{*2}\delta}{4\pi^2}\int_0^1 dz$$

$$\times[M^{*2}-z(1-z)\mathsf{p}^2]\ln\left[\frac{M^{*2}-z(1-z)\mathsf{p}^2}{M^2-z(1-z)m_s^2}\right].$$

In $\Pi_s^{fin}(p)$ the constant λ is not uniquely determined and we take it as a free parameter, in particular $\lambda = 1$ ensures the condition $\partial\Pi_s/\partial\mathsf{p}^2 = 0$ at zero density and at the scalar meson mass shell.

The formalism just described can be used to find the collective modes. These are self-sustained density fluctuations of the whole system, whose dispersion relation can be found by the condition $\det\mathcal{U} = 0$. That is, it corresponds to the states which break down the relation (10). In this work we neglect the imaginary part of the polarization insertions, this is an exact result in the range $0 < \mathsf{p}^2 < 4M^{*2}$ where imaginary parts are rigorously null.

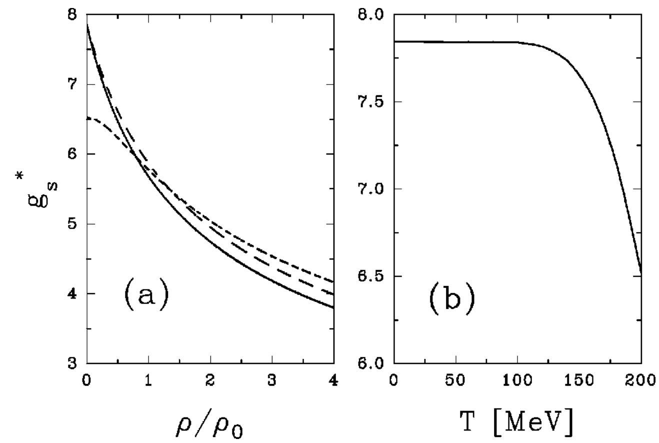

FIG. 1. The effective coupling g_s^* as a function of the medium properties. (a) The relative baryon density dependence for several temperatures; the curves with full, large-dashed, and short-dashed lines correspond to $T = 0$, 100, and 200 MeV, respectively. (b) The temperature dependence at zero baryon density.

IV. RESULTS AND DISCUSSION

In order to obtain numerical results we have taken $M = 939$ MeV, $m_s = 550$ MeV, and $m_v = 783$ MeV for the in-vacuum nucleon, σ-meson, and ω-meson masses. The coupling constants g_s, g_v are fixed in order to obtain the binding energy $\epsilon_B = -16$ MeV at the saturation baryon density $\rho_0 = 0.15$ fm^{-3} in the mean field approximation. In such a way we get $g_s = 7.845$, $g_v = 6.671$ and we obtain the effective nucleon mass $M^* = 0.85M$ and the compressibility $\kappa = 225$ MeV at the normal density ρ_0.

Mean field results can be obtained by solving the self-consistent condition (4) for given values of the temperature T and the chemical potential μ, which is related to the baryon density ρ through

$$\rho = \delta\int\frac{d^3p}{(2\pi)^3}[n_+(E^*) - n_-(E^*)], \qquad (13)$$

with $E^* = \sqrt{p^2+M^{*2}}$. Once the mean value σ_0 is known, the effective coupling g_s^* can be evaluated for subsequent application to the RRPA. In Fig. 1(a) the effective coupling is drawn in terms of the baryon density for several temperatures. It can be seen that it is a monotonously decreasing function of the density for all the range of temperatures $0 < T < 200$ MeV, a shift down of approximately 40% is observed at the density $\rho = 4\rho_0$ relative to its vacuum value. In Fig. 1(b) the zero density results for g_s^* are examined in terms of the temperature, it has an approximately constant behavior until $T \simeq 150$ MeV is reached, from here on a fast decrease is observed.

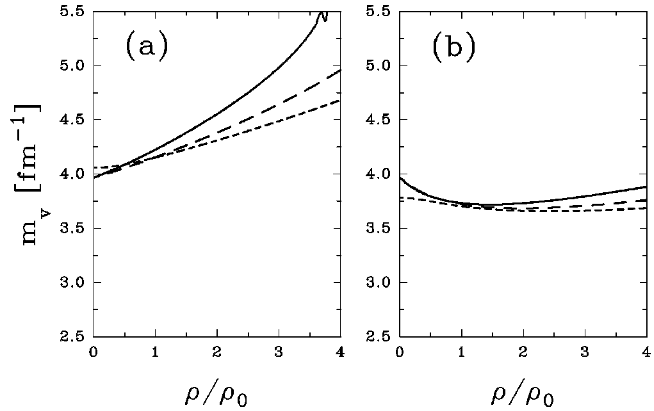

FIG. 2. The vector-meson effective mass as a function of the density for the temperatures $T=0$, 100, and 200 MeV. Line conventions are the same as in Fig. 1(a). The results corresponding to neglecting the Dirac sea contributions [case (a)] and those obtained by regularizing on the vector-meson mass shell [case (b)] are included.

The meson propagators evaluated in the RRPA can be used to extract the effective masses, which are defined as the roots of the inverse meson propagator at zero vector momentum, i.e., the p_0 solutions of

$$\mathcal{D}_{aa}^{-1}(p_0,p=0)=D_{aa}^{0-1}(p_0,p=0)-\Pi_{aa}(p_0,p=0)=0,$$

for $a=-1,1$.

The behavior of the lightest mesons masses (σ,ω,ρ) has been studied in the framework of the QHD-I model, since the early work of Chin [22] and under different approximations [23–27].

The density dependence of the vector meson effective mass is shown in Fig. 2, where we compare results obtained by completely neglecting divergent vacuum contributions [Fig. 2(a)] with those given by the on shell regularized polarization [Fig. 2(b)]. From Fig. 2(a) it can be seen that at zero temperature the vector meson mass increases with the density, at higher temperatures the growth becomes weaker. The inclusion of the vacuum inverts this trend for densities below a certain value, which changes with the temperature. For higher densities m_v turns increasing, but for all the range of densities and temperatures studied it remains below its vacuum value. Similar results have been found for the QHD-I model calculations, see for example [25–27]. The same comparison for the scalar meson is done in Fig. 3. Since the regularized scalar polarization depends on the free parameter λ, we have included the cases $\lambda=1$, 10, and 50.

For the scalar meson the situation is not so clear, different values of the parameter λ produce dissimilar density variations. In Fig. 3 the three cases $\lambda=1,10,50$ are shown. With these values we have obtained in the first case $(\lambda=1)$ a nonmonotonic density variation, for $\lambda=10$ an almost constant function for both density and temperature, finally for $\lambda=50$, a monotonously decreasing function is the outcome. The results without Dirac sea contribution can be compared for example with those of [9], where different versions of the DSCM are used at zero temperature (note that a different medium dependent coefficient is assigned to the effective coupling there). On the other hand, by choosing $\lambda=50$ in our calculations, an asymptotically decreasing σ-meson mass is obtained. This contrasts with the QHD-I model results of [25,27].

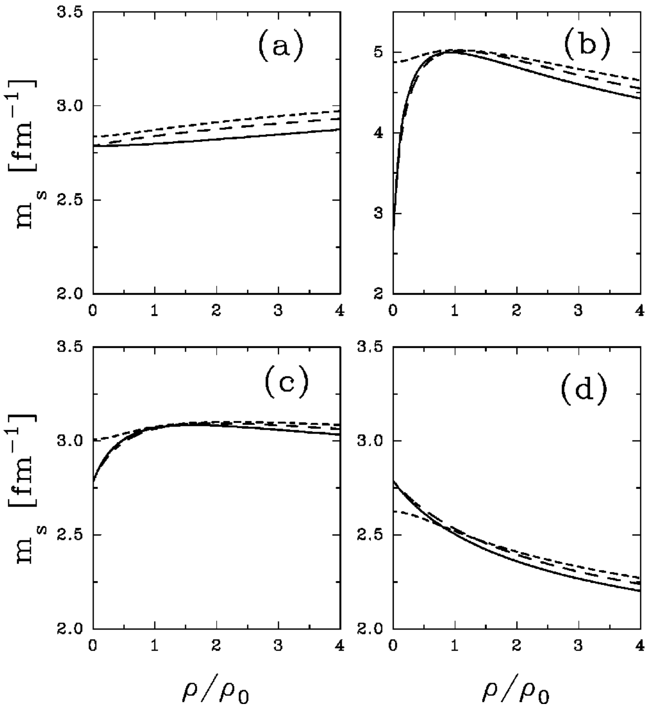

FIG. 3. The scalar-meson effective mass as a function of the density for the temperatures $T=0$, 100, and 200 MeV. Line conventions are the same as in Fig. 1(a). The results corresponding to neglecting the Dirac sea contributions [case (a)] and regularizing on the σ-meson mass shell for values of the regularizing parameter $\lambda=1$ [case (b)], $\lambda=10$ [case (c)], and $\lambda=50$ [case (d)] are included.

Thermal effects diminish the rate of variation with density. The significance of the σ meson in hadronic models is unclear at present. It has been interpreted as a correlated exchange of many pions [28], this hypothesis could explain the success of models like QHD-I which do not include pions explicitly. In other contexts this meson is regarded as the chiral partner of the pion. In any case the scalar meson seems to be linked to the manifestation of chiral symmetry in hadronic physics, and a dropping of the σ mass with density is generally accepted.

Collective excitations in nuclear matter have been previously investigated within the relativistic QHD-I model [22,23,25,27,29]. A very rich structure of the dispersion relations has been classified in [23] as zero sound, meson-branch modes, and instabilities. We have found similar results at zero temperature, but at higher temperatures only the so-called meson-branch modes survive. We have studied the range $0<p_0$, $p<2$ GeV of the momentum.

A brief digression on the so-called Landau ghost poles is opportune here. The presence of spacelike singularities in meson propagators was reported long ago [30]. Several recipes to extract spurious contributions have been stated [31,32]. However the imaginary part of the inverse meson propagators is nonzero for $\mathsf{p}^2<0$, and the existence of Lan-

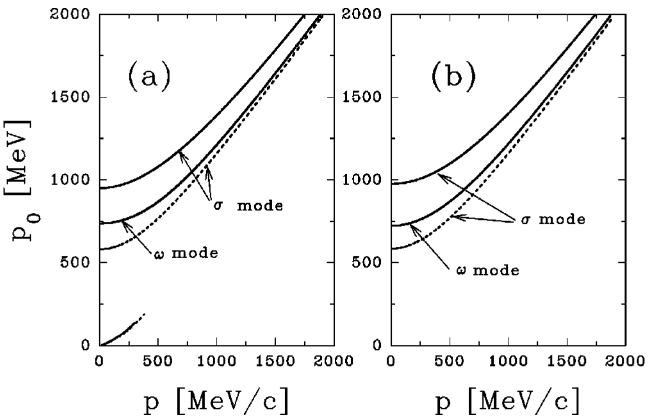

FIG. 4. The dispersion relations for longitudinal collective modes at $T=0$ (a) and $T=200$ MeV (b). In both figures the cases $\lambda=1$ and $\lambda=10$ are represented by solid and dotted lines, respectively. Here the presence of the so-called σ and ω excitation modes, as explained in the text, can be appreciated.

dau ghost at low densities could be a signal of nuclear matter instability against the gas-liquid phase transition [23]. At high densities these poles occur at several times the free nucleon mass and possibly its presence sets up the limit of applicability of the model. Notwithstanding, the numerical calculations in [31] show a negligible influence on the lower momentum applications.

The zeros of det $\mathcal{U}$ can be classified as zeros of Eq. (11) (transverse mode excitations) or as zeros of Eq. (12) (longitudinal mode excitations). In the last case a two branch pattern is found, recognized as the σ- and ω-meson modes.

As an example of high density results, the dispersion relation for the longitudinal collective excitations at twice the normal density are shown in Fig. 4. The cases of $T=0$ and $T=200$ MeV correspond to Figs. 4(a) and 4(b), respectively. Two different values of the regularizing constant ($\lambda=1,10$) have been used. The negligible magnitude of Π_M as compared to Π_s and Π_v, produces a too small σ-ω mixing. As a consequence the longitudinal ω-meson mode is almost independent of the λ parameter. This fact is evident in Fig. 4, where the curves for the ω mode corresponding to $\lambda=1$ and $\lambda=10$ are indistinguishable. The same fact can be observed for the zero mode in the left-bottom corner of Fig. 4(a), which is nearly unaffected by the value of λ.

On the other hand, a lower-energy regime starting at $p_0 \simeq 600$ MeV for $\lambda=10$ replaces the scalar-meson curve starting at $p_0 \simeq 950$ MeV for $\lambda=1$. Finite temperature effects are almost negligible for the meson branches, although zero modes are drastically suppressed with temperature. In Fig. 5 the curves for the transverse mode dispersion relation are displayed for $T=0$ and 200 MeV. This component does not depend on λ. As in the previous case, temperature effects are small and are more apparent for lower momenta. These results can be compared with previous calculations using the Walecka model [25–27]. A qualitative agreement can be observed, but a global enhancement of all modes is revealed in the DSCM. The choice $\lambda=10$ in our calculations improves the coincidence between both models.

A qualitative description of the giant resonances of a system with a finite particle number A can be given in terms of the previous results. For this purpose we have based our results on the approach developed in Refs. [33,34], which is founded on the identification of the Landau-Migdal parameters in the expression of the RRPA dielectric function. This is a suitable procedure for sufficiently high A. The expression obtained by taking the limit $p \to 0$ of Eq. (12) at $T=0$ can be compared with the longitudinal dielectric function of the Landau liquid model

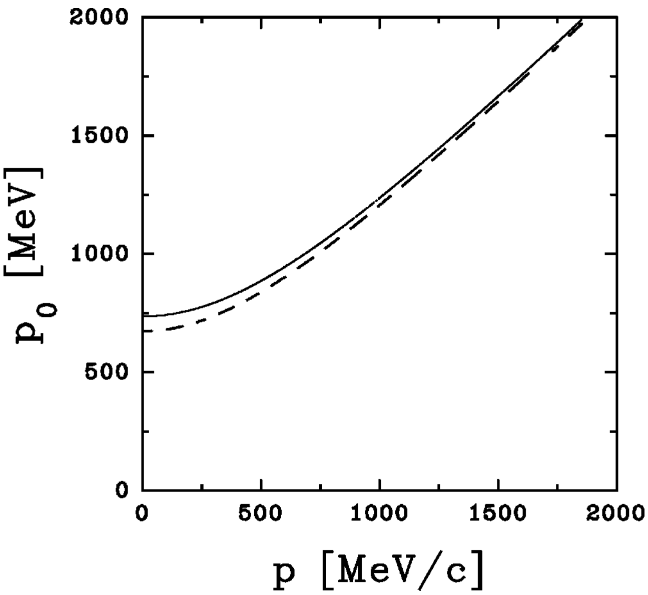

FIG. 5. The transverse mode dispersion relation for collective excitations in symmetric nuclear matter for $T=0$ (solid line) and $T=200$ MeV (dashed line).

$$\epsilon_L = c\left[1+\Phi(x)\left(F_0+\frac{x^2F_1}{1+F_1/3}\right)\right], \tag{14}$$

where $\Phi(x)$ is the Linhardt function of argument $x=p_0/pv_F$ and c,F_0,F_1 are the Landau-Migdal parameters. The Fermi velocity v_F is defined in terms of the Fermi momentum (p_F) and energy ($E_F=\sqrt{p_F^2+M^{*2}}$) through $v_F=p_F/E_F$. By these means we obtain

$$c=1+F_s\left[\frac{3}{2}-v_F^2-3\frac{1-v_F^2}{4v_F}\ln\left(\frac{1+v_F}{1-v_F}\right)+A_{reg}\right],$$

$$F_0=\left[F_v-F_s\frac{1-v_F^2}{c}\right],$$

$$F_1=-\frac{F_v}{1+\frac{N_F}{3}\left(\frac{g_v v_F}{m_v}\right)^2},$$

with $N_F=\delta p_F E_F/\pi^2$ the density of states at the Fermi surface, $F_s=N_F(g_s/m_s)^2$, $F_v=N_F(g_v/m_v)^2$, and A_{reg} comes from the regularized part of Π_s

TABLE I. Comparison of the coefficients α_M and α_Q. The first column is the experimental accepted values, see for example [36], the second column corresponds to theoretical estimations of [33] without Dirac sea contribution, the third column is obtained from the numerical values in [34] which includes vacuum corrections, and the last three columns are our results for regularizing parameter $\lambda=1$, 10, and 50, respectively.

	Expt.	KS1	KS2	$\lambda=1$	$\lambda=10$	$\lambda=50$
α_M [MeV]	80	160	115	149	161	167.7
α_Q [MeV]	63	87.8	56.2	52.7	52.7	52.7

$$A_{reg}=\frac{1-v_F^2}{4v_F}\left[\frac{\lambda}{2}\left(\frac{m_s}{M}\right)^2+6\ \ln(M/M^*)-6+\frac{6}{m_s}\sqrt{4M^2-m_s^2}\ \arctan\left(\frac{m_s}{\sqrt{4M^2-m_s^2}}\right)\right]. \quad (15)$$

The giant monopole E_M and quadrupole E_Q excitation energies can be approximately described in terms of the parameters F_0 and F_1 previously obtained by

$$E_M=\frac{v_F}{1.2A^{1/3}}\sqrt{(1+F_0)(3+F_1)}=\alpha_M A^{-1/3},$$

$$E_Q=\frac{v_F}{1.2A^{1/3}}\sqrt{2(3+F_1)/5}=\alpha_Q A^{-1/3}.$$

The coefficients of $A^{-1/3}$ can be extracted from the experimental data. In order to estimate α_M and α_Q we have evaluated the expressions above at the saturation density. As can be seen in Table I, our results are of the same magnitude of those given by [34]. The result obtained for α_M is roughly twice the expected value, whereas α_Q is more well described. In Ref. [35] it has recently been reported that the inclusion of nonlocal pionic contributions improves the value of α_Q.

V. CONCLUSIONS

We have investigated the collective properties of hadronic matter in a relativistic model with density dependent nucleon-meson interaction. The effective meson masses as functions of the density and temperature of the hadronic medium have been evaluated by means of the proper polarization insertion. These polarizations have been constructed in the RRPA, including vacuum contributions. The effective ω-meson mass at zero temperature decreases for densities below 1.5 ρ_0, the behavior of the σ meson is strongly affected by the value of the parameter λ used in the regularization. The contribution of the Dirac sea to the RRPA polarization is significative, as can be seen by comparing with similar calculations [9]. The dispersion relations for collective excitations are qualitatively similar to those obtained with the QHD-I model, although the absolute magnitude of the scalar mode of the longitudinal component depends on the regularizing parameter used. The appearance of ghost poles at relative high momenta $p\simeq 3-3.5$ GeV occurs far from the range of the present work. Isoscalar giant resonances evaluated in terms of the relativistic Landau-Migdal parameters, are qualitatively described. However, a more realistic model could give a better quantitative agreement.

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