In-medium vector meson masses in a chiral SU(3) model

D. Zschiesche,[1,*] A. Mishra,[1,2,†] S. Schramm,[3] H. Stöcker,[1] and W. Greiner[1]
[1]*Institut für Theoretische Physik, Robert Mayer Strasse 8-10, D-60054 Frankfurt am Main, Germany*
[2]*Department of Physics, I.I.T. Delhi, Hauz Khas, New Delhi - 110 016, India*
[3]*Argonne National Laboratory, 9700 South Cass Avenue, Argonne, Illinois 60439, USA*
(Received 4 March 2003; revised manuscript received 1 June 2004; published 11 October 2004)

A significant drop of the vector meson masses in nuclear matter is observed in a chiral SU(3) model due to the effects of the baryon Dirac sea. This is taken into account through the summation of baryonic tadpole diagrams in the relativistic Hartree approximation. The appreciable decrease of the in-medium vector meson masses is due to the vacuum polarization effects from the nucleon sector and is not observed in the mean field approximation.

DOI: 10.1103/PhysRevC.70.045202 PACS number(s): 21.65.+f, 24.10.Jv

I. INTRODUCTION

The medium modifications of the vector mesons (ρ and ω) in hot and dense matter have recently been a topic of great interest in strong interaction physics research, both experimentally [1–5] and theoretically [6–12]. One of the explanations of the experimental observation of enhanced dilepton production [1–3] in the low invariant mass regime could be a reduction in the vector meson masses in the medium. It was first suggested by Brown and Rho that the vector meson masses drop in the medium according to a simple scaling law [6], given as $m_V^*/m_V = f_\pi^*/f_\pi$. f_π is the pion decay constant and the asterisk refers to in-medium quantities. There have also been QCD sum rule approaches extensively used in the literature [8–11] for consideration of the in-medium vector meson properties. In the framework of quantum hadrodynamics (QHD) [13] as a description of the hadronic matter, it is seen that the dropping of the vector meson masses has its dominant contribution arising from the vacuum polarization effects in the baryon sector [14–17]. This drop is not observed in the mean field approximation. The vector meson properties [18] and their effects on the low mass dilepton spectra [19] have been investigated recently including the quantum correction effects from the baryon as well as the scalar meson sectors in the Walecka model [20].

In the present investigation we use the SU(3) chiral model [21,22] for description of the hadronic matter. This model has been shown to successfully describe hadronic properties in the vacuum as well as nuclear matter, finite nuclei, and neutron star properties. Furthermore the model consistently includes the lowest lying baryon and meson multiplets, including the vector mesons. In the mean field approximation the vector meson masses do not show any significant drop, similar to results in the Walecka model. The effect of the Dirac sea is taken into account by summing over baryonic tadpole diagrams in the relativistic Hartree approximation (RHA). In an alternative approach to QHD renormalizability [23], a chiral effective model has been studied emphasizing the concept of naturalness as the guiding principle for constructing the model. While naturalness arguments sound appealing in themselves, they have clear limits in the study of nuclear matter properties with relativistic meson field theories, as—independent of the approximation scheme—the calculations always involve a rather "unnatural" fine tuned cancellation of the strong scalar attraction and vector repulsion inherent in those approaches. We therefore follow the direct field-theoretical path by including the full tadpole contributions from the lower continuum states. It is seen that an appreciable decrease of the vector meson masses arises from the nucleon Dirac sea. This shows the importance of taking into account these contributions.

We organize the paper as follows. In Sec. II we introduce the chiral SU(3) model used in the present investigation. Section III describes the mean field approximation for nuclear matter. In Sec. IV the nuclear matter properties are considered in the relativistic Hartree approximation. Section V gives the in-medium vector meson properties due to the contributions from the nucleon Dirac sea. The results are presented and discussed in Sec. VI. Finally, in Sec. VII we summarize the findings of the present work.

II. THE HADRONIC CHIRAL SU(3)×SU(3) MODEL

We consider a relativistic field-theoretical model of baryons and mesons built on chiral symmetry and broken scale invariance [21,22]. A nonlinear realization of chiral symmetry is adopted, that has been successful in a simultaneous description of finite nuclei and hyperon potentials [21]. The general form of the Lagrangian is as follows:

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

$\mathcal{L}_{kin}$ is the kinetic energy term; $\mathcal{L}_{BW}$ includes the interaction terms of the baryons with the spin-0 and spin-1 mesons, the former generating the baryon masses. $\mathcal{L}_{VP}$ contains the interaction terms of vector mesons with pseudoscalar mesons. $\mathcal{L}_{vec}$ generates the masses of the spin-1 mesons through interactions with spin-0 fields and contains quartic self-interactions of the vector fields. $\mathcal{L}_0$ gives the meson-meson interaction terms which induce the spontaneous breaking of

*Electronic address: ziesche@th.physik.uni-frankfurt.de
†Electronic address: mishra@th.physik.uni-frankfurt.de

chiral symmetry. It also includes a scale-invariance breaking logarithmic potential. Finally, $\mathcal{L}_{SB}$ introduces an explicit symmetry breaking of the $U(1)_A, SU(3)_V$, and chiral symmetry.

A. Kinetic terms

The kinetic energy terms are given as [21]

$$\mathcal{L}_{kin} = i\,\mathrm{Tr}\,\overline{B}\gamma_\mu D^\mu B + \tfrac{1}{2}\mathrm{Tr}\,D_\mu X D^\mu X + \mathrm{Tr}(u_\mu X u^\mu X + X u_\mu u^\mu X) + \tfrac{1}{2}\mathrm{Tr}\,D_\mu Y D^\mu Y + \tfrac{1}{2}D_\mu \chi D^\mu \chi - \tfrac{1}{4}\mathrm{Tr}(\tilde{V}_{\mu\nu}\tilde{V}^{\mu\nu}) - \tfrac{1}{4}\mathrm{Tr}(F_{\mu\nu}F^{\mu\nu}) - \tfrac{1}{4}\mathrm{Tr}(\mathcal{A}_{\mu\nu}\mathcal{A}^{\mu\nu}), \tag{2}$$

where B is the baryon octet, X is the scalar multiplet, Y is the pseudoscalar chiral singlet, $\tilde{V}^\mu$ is the vector meson multiplet with field tensor $\tilde{V}_{\mu\nu} = \partial^\nu \tilde{V}^\mu - \partial^\nu \tilde{V}^\mu$,[1] $A_{\mu\nu} = \partial^\nu A^\mu - \partial^\nu A^\mu$ is the axial-vector field tensor, $F_{\mu\nu}$ is the electromagnetic field tensor, and χ is the scalar, isoscalar glueball field. The kinetic energy term for the pseudoscalar mesons is given in terms of the axial vector $u_\mu = -(i/2)[u^\dagger \partial_\mu u - u \partial_\mu u^\dagger]$, where $u = \exp[(i/2\sigma_0)\pi^a \lambda_a \gamma_5]$ is the unitary transformation operator [21]. The pseudoscalar mesons are given as parameters of the symmetry transformation. Since the fields in the nonlinear realization of chiral symmetry contain the local unitary transformation operator, covariant derivatives $D_\mu = \partial_\mu + i[\Gamma_\mu, \ ]$, with $\Gamma_\mu = -(i/2)[u^\dagger \partial_\mu u + u \partial_\mu u^\dagger]$ have to be used to guarantee chiral invariance [21]. For example, for the baryons this yields

$$D_\mu B = \partial_\mu B + i[\Gamma_\mu, B]. \tag{3}$$

B. Baryon-meson interaction

The SU(3) structure of the the baryon-meson interaction terms are the same for all mesons, except for the difference in Lorentz space. For a general meson field W they read

$$\mathcal{L}_{BW} = -\sqrt{2}g_8^W\{\alpha_W[\overline{B}\mathcal{O}BW]_F + (1-\alpha_W)[\overline{B}\mathcal{O}BW]_D\} - g_1^W\frac{1}{\sqrt{3}}\mathrm{Tr}(\overline{B}\mathcal{O}B)\mathrm{Tr}\,W, \tag{4}$$

with $[\overline{B}\mathcal{O}BW]_F := \mathrm{Tr}(\overline{B}\mathcal{O}WB - \overline{B}\mathcal{O}BW)$ and $[\overline{B}\mathcal{O}BW]_D := \mathrm{Tr}(\overline{B}\mathcal{O}WB + \overline{B}\mathcal{O}BW) - \frac{2}{3}\mathrm{Tr}(\overline{B}\mathcal{O}B)\mathrm{Tr}\,W$. The different terms to be considered are those for the interaction of baryons with scalar mesons ($W=X$, $\mathcal{O}=1$), with vector mesons ($W=\tilde{V}_\mu$, $\mathcal{O}=\gamma_\mu$ for the vector and $W=\tilde{V}_{\mu\nu}$, $\mathcal{O}=\sigma^{\mu\nu}$ for the tensor interaction), with axial-vector mesons ($W=\mathcal{A}_\mu$, $\mathcal{O}=\gamma_\mu\gamma_5$), and with pseudoscalar mesons ($W=u_\mu$, $\mathcal{O}=\gamma_\mu\gamma_5$), respectively. In the following we discuss the relevant couplings for the current investigation.

[1]As described in Sec. II C, the vector mesons need to be renormalized. The physical fields will be denoted as V_μ and $\rho_\mu, \omega_\mu, \phi_\mu$, respectively, and the unrenormalized, mathematical fields as $\tilde{V}_\mu$ and $\tilde{\rho}_\mu, \tilde{\omega}_\mu, \tilde{\phi}$.

1. Baryon–scalar-meson interaction (baryon masses)

The baryons and the scalar mesons transform equally in the left and right subspaces. Therefore, in contrast to the linear realization of chiral symmetry, an f-type coupling is allowed for the baryon-meson interaction. In addition, it is possible to construct mass terms for baryons and to couple them to chiral singlets. After insertion of the vacuum expectation value (VEV) for the scalar multiplet matrix $\langle X\rangle_0$, one obtains the baryon masses as generated by the VEV of the nonstrange $\sigma \sim \langle \overline{u}u + \overline{d}d\rangle$ and the strange $\zeta \sim \langle \overline{s}s\rangle$ scalar fields [21]. Here we will consider the limit $\alpha_S = 1$ and $g_1^S = \sqrt{6}g_8^S$. In this case the nucleon mass does depend only on the nonstrange condensate σ. Furthermore, the coupling constants between the baryons and the two scalar condensates are related to the additive quark model. This leaves only one coupling constant free that is adjusted to give the correct nucleon mass [21]. For a fine-tuning of the remaining masses, it is necessary to introduce an explicit symmetry breaking term, which breaks the SU(3) symmetry along the hypercharge direction (for details see [21]). Therefore the resulting baryon octet masses for the current investigation read

$$m_N = -g_{N\sigma}\sigma_0,$$

$$m_\Lambda = -g_{N\sigma}\left(\frac{2}{3}\sigma_0 - \frac{1}{3}\sqrt{2}\zeta_0\right) + \frac{m_1 + 2m_2}{3},$$

$$m_\Sigma = -g_{N\sigma}\left(\frac{2}{3}\sigma_0 - \frac{1}{3}\sqrt{2}\zeta_0\right) + m_1,$$

$$m_\Xi = -g_{N\sigma}\left(\frac{1}{3}\sigma_0 - \frac{2}{3}\sqrt{2}\zeta_0\right) + m_1 + m_2. \tag{5}$$

Alternative ways of mass generation have also been considered earlier [21].

2. Baryon–vector-meson interaction

Two independent interaction terms of baryons with spin-1 mesons can be constructed in analogy with the baryon–spin-0-meson interaction. They correspond to the antisymmetric (f-type) and symmetric (d-type) couplings, respectively. The general couplings are shown in [21]. From the universality principle [24] and the vector meson dominance model one may conclude that the d-type coupling should be small. Here we will use pure f-type coupling, i.e., $\alpha_V = 1$ for all fits, even though a small admixture of d-type coupling allows for some fine-tuning of the single particle energy levels of nucleons in nuclei (see [21]). As for the case with scalar mesons, we furthermore set $g_1^V = \sqrt{6}g_8^V$, so that the strange vector field $\tilde{\phi}_\mu \sim \overline{s}\gamma_\mu s$ does not couple to the nucleon. The resulting Lagrangian reads

$$\mathcal{L}_{BV} = -\sqrt{2}g_8^V([\overline{B}\gamma_\mu B\tilde{V}^\mu]_F + \mathrm{Tr}(\overline{B}\gamma_\mu B)\mathrm{Tr}\,\tilde{V}^\mu), \tag{6}$$

or explicitly written out for the nuclear matter case,

$$\mathcal{L}_{BV}^{N} = 3g_8^V\tilde{\omega}_\mu\overline{\psi_N}\gamma_\mu\psi_N + g_8^V\tilde{\rho}_\mu\overline{\psi_N}\gamma_\mu\tau_3\psi_N. \quad (7)$$

Note that in this limit all coupling constants are fixed once g_8^V is specified [21]. This is done by fitting the nucleon-ω coupling to the energy density at nuclear matter saturation ($E/A=-16$ MeV). Since we consider nuclear matter, the couplings of the vector mesons to the hyperons will not be discussed here.

C. Meson-meson interactions

1. Vector mesons

The vector meson-meson interactions contain the mass terms of the vector mesons and higher order vector meson self-interactions. The simplest scale invariant mass term is

$$\mathcal{L}_{vec}^{(1)} = \frac{1}{2}m_V^2\frac{\chi^2}{\chi_0^2}\mathrm{Tr}\,\tilde{V}_\mu\tilde{V}^\mu. \quad (8)$$

It implies a mass degeneracy for the vector meson nonet. The scale invariance is assured by the square of the glueball field χ (see Sec. II C 2 for details). To split the masses, one can add the chiral invariants [25,26]

$$\mathcal{L}_{vec}^{(2)} = \frac{1}{4}\mu\,\mathrm{Tr}[\tilde{V}_{\mu\nu}\tilde{V}^{\mu\nu}X^2] \quad (9)$$

and

$$\mathcal{L}_{vec}^{(3)} = \frac{1}{12}\lambda_V(\mathrm{Tr}[\tilde{V}_{\mu\nu}])^2. \quad (10)$$

Note that in Eq. (10) we replace the scalar multiplet X by its vacuum expectation value. Combining the contributions (9) and (10) with the kinetic energy term (2), one obtains the following terms for the vector mesons in the vacuum:

$$-\frac{1}{4}Z_\rho^{-1}(\tilde{V}_{\tilde{\rho}}^{\mu\nu})^2 - \frac{1}{4}Z_\omega^{-1}(\tilde{V}_{\tilde{\omega}}^{\mu\nu})^2 - \frac{1}{4}Z_\phi^{-1}(\tilde{V}_{\tilde{\phi}}^{\mu\nu})^2, \quad (11)$$

with, e.g., $\tilde{V}_{\tilde{\rho}}^{\mu\nu}=\partial^\mu\tilde{\rho}^\nu-\partial^\nu\tilde{\rho}^\mu$. With the renormalization constants the new vector meson fields are defined as $\rho=Z_\rho^{-1/2}\tilde{\rho}$, $\omega=Z_\omega^{-1/2}\tilde{\omega}$, $\phi=Z_\phi^{-1/2}\tilde{\phi}$. Explicitly the renormalization constants are given as

$$Z_\rho^{-1} = \left(1-\mu\frac{\sigma_0^2}{2}\right),$$

$$Z_{\omega,\phi}^{-1} = \left[\left(1-\frac{\mu(\sigma_0^2+2\xi_0^2)-2\lambda_V}{4}\right)\pm\frac{1}{2}D^{1/2}\right], \quad (12)$$

where

$$D = \frac{\mu^2}{4}(\sigma_0^2-2\xi_0^2)^2 + \lambda_V^2 - \frac{\lambda_V}{3}\mu(\sigma_0^2-2\xi_0^2). \quad (13)$$

Then the Lagrangian for the new fields in the vacuum reads

$$\mathcal{L}_{vec}^{\mathrm{vac}} = -\frac{1}{4}[(V_\rho^{\mu\nu})^2+(V_\omega^{\mu\nu})^2+(V_\phi^{\mu\nu})^2] + \frac{1}{2}\frac{\chi^2}{\chi_0^2}(m_\rho^2\rho^2+m_\omega^2\omega^2+m_\phi^2\phi^2) \quad (14)$$

where

$$m_\rho^2 = Z_\rho m_V^2,\quad m_\omega^2 = Z_\omega m_V^2,\quad m_\phi^2 = Z_\phi m_V^2 \quad (15)$$

denote the vector meson masses in the vacuum. Using $m_V=687.33$ MeV, $\mu\sigma_0^2=0.41$, and $\lambda_V=-0.041$, the correct ω, ρ, and ϕ masses are obtained. The vector meson self-interactions [27,28] read

$$\mathcal{L}_{vec}^{(4)} = 2(\tilde{g}_4)^4\mathrm{Tr}(\tilde{V}_\mu\tilde{V}^\mu)^2. \quad (16)$$

The coupling of this self-interaction term is also modified by the redefinition of the fields. The redefined coupling corresponding to the quartic interaction for the ω field can be expressed in terms of the coupling $\tilde{g}_4$ of the term (16). This term gives a contribution to the vector meson masses in the medium, i.e., for finite values of the ω or ρ fields. The resulting expressions for the vector meson masses in the medium (isospin symmetric) are

$$(m_\omega^{eff})^2 = m_\omega^2 + 12g_4^4\omega^2, \quad (17)$$

$$(m_\rho^{eff})^2 = m_\rho^2 + 12g_4^4\frac{Z_\rho}{Z_\omega}\omega^2, \quad (18)$$

$$(m_\phi^{eff})^2 = m_\phi^2 + 24g_4^4\frac{Z_\phi^2}{Z_\omega^2}\phi^2, \quad (19)$$

with $g_4=\sqrt{Z_\omega}\tilde{g}_4$ as the renormalized coupling. Since the quartic self-interaction contributes only in the medium, the coupling g_4 cannot be unambigiously fixed. It is fitted, so that the compressibility is in the desired region between 200–300 MeV in the mean field approximation. Note that the N-ω as well as the N-ρ couplings are also affected by the redefinition of the fields with the corresponding renormalized coupling constants as $g_{N\omega}\equiv 3g_V^8\sqrt{Z_\omega}$ and $g_{N\rho}\equiv g_V^8\sqrt{Z_\rho}$.

2. Spin-0 potential

In the nonlinear realization of chiral symmetry the couplings of scalar mesons X and the pseudoscalar singlet Y with each other are governed only by $SU(3)_V$ symmetry. In this work we will use the same form of the potential as in the linear σ model with $U(1)_A$ breaking, as described in [21]. It reads

$$\mathcal{L}_0 = -\frac{1}{2}k_0\chi^2I_2 + k_1(I_2)^2 + k_2I_4 + 2k_3\chi I_3, \quad (20)$$

with $I_2=\mathrm{Tr}(X+iY)^2$, $I_3=\det(X+iY)$, and $I_4=\mathrm{Tr}(X+iY)^4$. Furthermore χ denotes a scalar color-singlet gluon field. It is introduced to construct the model to satisfy the QCD trace anomaly, i.e., the nonvanishing of the trace of the energy-momentum tensor $\theta_\mu^\mu=(\beta_{QCD}/2g)\mathcal{G}_{\mu\nu}^a\mathcal{G}^{\mu\nu a}$. Here, $\mathcal{G}_{\mu\nu}^a$ is the gluon field strength tensor of QCD.

All the terms in the Lagrangian are multiplied by appropriate powers of the glueball-field to obtain a dimension $(\mathrm{mass})^4$ in the fields. Then all coupling constants are dimensionless and therefore the model is scale invariant [29]. Then, a scale breaking potential

$$\mathcal{L}_{scale\ break} = -\frac{1}{4}\chi^4 \ln\frac{\chi^4}{\chi_0^4} + \frac{\delta}{3}\chi^4 \ln\frac{I_3}{\det\langle X\rangle_0} \tag{21}$$

is introduced. This yields $\theta_\mu^\mu = (1-\delta)\chi^4$. By identifying the χ field with the gluon condensate and the choice $\delta = 6/33$ for three flavors and three colors with β_{QCD} as given by the one-loop level, the correct trace anomaly is obtained. The first term in Eq. (21) corresponds to the contribution of the gluons and the second term describes the contribution from the quarks to the trace anomaly. Finally the term

$$\mathcal{L}_\chi = -k_4\chi^4 \tag{22}$$

generates a phenomenologically consistent finite vacuum expectation value.

The parameters k_0, k_2, and k_4 are used to ensure an extremum in the vacuum for the σ-, ζ-, and χ-field equations, respectively. As for the remaining constants, k_3 is constrained by the η and η' masses, which take the values $m_\eta = 520$ MeV and $m_{\eta'} = 999$ MeV in all parameter sets. k_1 is fixed in the mean field fit with quartic vector meson interaction such that the effective nucleon mass at saturation density is around $0.65 m_N$ and the σ mass is of the order of 500 MeV. Then it is kept constant in all the other fits, since a change in k_1 yields quite a strong modification of the other coupling constants in the self-consistency calculation. Since we want to focus on the influence of the Hartree terms, we try to keep everything else as little modified as possible.

Since the shift in the χ in the medium is rather small [21], we will in good approximation set $\chi = \chi_0$. We will refer to this case as the *frozen glueball limit*. The VEV of the gluon condensate, χ_0, is fixed to fit the pressure $p = 0$ at the saturation density $\rho_0 = 0.15\ \mathrm{fm}^{-3}$.

D. Explicitly broken chiral symmetry

In order to eliminate the Goldstone modes from a chiral effective theory, explicit symmetry breaking terms have to be introduced. Here, we again take the corresponding term of the linear σ model

$$\mathcal{L}_{SB} = \frac{1}{2}\mathrm{Tr}\, A_p(M + M^\dagger) = \mathrm{Tr}\, A_p[u(X + iY)u + u^\dagger(X - iY)u^\dagger] \tag{23}$$

with $A_p = 1/\sqrt{2}\,\mathrm{diag}(m_\pi^2 f_\pi, m_\pi^2 f_\pi, 2m_K^2 f_K - m_\pi^2 f_\pi)$ and $m_\pi = 139$ MeV, $m_K = 498$ MeV. This choice for A_p together with the constraints

$$\sigma_0 = -f_\pi, \quad \zeta_0 = -\frac{1}{\sqrt{2}}(2f_K - f_\pi) \tag{24}$$

on the VEV on the scalar condensates assure that the partial conservation of axial-vector coupling (PCAC) relations of the pion and kaon are fulfilled. With $f_\pi = 93.3$ MeV and $f_K = 122$ MeV we obtain $\sigma_0 = 93.3$ MeV and $\zeta_0 = 106.56$ MeV.

III. MEAN FIELD APPROXIMATION

The hadronic matter properties at finite density and temperature are studied in the mean field approximation [30]. Then the Lagrangian (1) becomes

$$\mathcal{L}_{BX} + \mathcal{L}_{BV} = -\overline{\psi_N}[g_{N\omega}\gamma_0\omega + m_N^*]\psi_N, \tag{25}$$

$$\mathcal{L}_{vec} = \frac{1}{2}m_\omega^2\frac{\chi^2}{\chi_0^2}\omega^2 + g_4^4\omega^4, \tag{26}$$

$$\mathcal{V}_0 = \frac{1}{2}k_0\chi^2(\sigma^2+\zeta^2) - k_1(\sigma^2+\zeta^2)^2 - k_2\left(\frac{\sigma^4}{2}+\zeta^4\right) - k_3\chi\sigma^2\zeta + k_4\chi^4 + \frac{1}{4}\chi^4\ln\frac{\chi^4}{\chi_0^4} - \frac{\delta}{3}\chi^4\ln\frac{\sigma^2\zeta}{\sigma_0^2\zeta_0}, \tag{27}$$

$$\mathcal{V}_{SB} = \left(\frac{\chi}{\chi_0}\right)^2\left[m_\pi^2 f_\pi\sigma + \left(\sqrt{2}m_K^2 f_K - \frac{1}{\sqrt{2}}m_\pi^2 f_\pi\right)\zeta\right], \tag{28}$$

where m_N^* is the effective mass of the nucleon. Only the scalar $(\mathcal{L}_{BX})$ and the vector meson terms $(\mathcal{L}_{BV})$ contribute to the baryon-meson interaction. For all other mesons, the expectation value vanishes in the mean field approximation. Now it is straightforward to write down the expression for the thermodynamical potential of the grand canonical ensemble, Ω, per volume V at a given chemical potential μ and at zero temperature:

$$\frac{\Omega}{V} = -\mathcal{L}_{vec} - \mathcal{L}_0 - \mathcal{L}_{SB} - \mathcal{V}_{\mathrm{vac}} + \frac{\gamma_N}{(2\pi)^3}\int_0^{\sqrt{\mu_N^{*\,2} - m_N^{*\,2}}} d^3k[E_N^*(k) - \mu_N^*] \tag{29}$$

The vacuum energy $\mathcal{V}_{\mathrm{vac}}$ (the potential at $\rho = 0$) has been subtracted in order to get a vanishing vacuum energy. The factor γ_N denotes the fermionic spin-isospin degeneracy factor, and $\gamma_N = 4$ for symmetric nuclear matter. The single particle energy is $E_N^*(k) = \sqrt{k_N^2 + m_N^{*\,2}}$ and the effective chemical potential reads $\mu_N^* = \mu_N - g_{N\omega}\omega$.

The mesonic fields are determined by extremizing the thermodynamic potential. Since we use the frozen glueball approximation (i.e., $\chi = \chi_0$), we have coupled equations only for the fields σ, ζ, and ω in the self-consistent calculation given as

$$\frac{\partial(\Omega/V)}{\partial\sigma} = k_0\chi^2\sigma - 4k_1(\sigma^2+\zeta^2)\sigma - 2k_2\sigma^3 - 2k_3\chi\sigma\zeta - 2\frac{\delta\chi^4}{3\sigma} + m_\pi^2 f_\pi + \frac{\partial m_N^*}{\partial\sigma}\rho_N^s = 0, \tag{30}$$

$$\frac{\partial(\Omega/V)}{\partial\zeta} = k_0\chi^2\zeta - 4k_1(\sigma^2+\zeta^2)\zeta - 4k_2\zeta^3 - k_3\chi\sigma^2 - \frac{\delta\chi^4}{3\zeta} + \left[\sqrt{2}m_K^2 f_K - \frac{1}{\sqrt{2}}m_\pi^2 f_\pi\right] = 0, \tag{31}$$

TABLE I. Parameters of the model, the corresponding terms in the Lagrangian, and constraints for fixing them.

Parameter	Interaction	Lagrange term	Observable/constraint
g_8^S	$\mathcal{L}_{\mathcal{BM}}$	$g_8^S\sqrt{2}\{\mathrm{Tr}(\overline{B}B)\mathrm{Tr}\,X+[\overline{B}BX]_F\}$	$m_N=-g_{N\sigma}\sigma_0$
g_8^V	$\mathcal{L}_{\mathcal{BV}}$	$g_8^V\sqrt{2}\{\mathrm{Tr}(\overline{B}\gamma_\mu B)\mathrm{Tr}\,V^\mu+[\overline{B}\gamma_\mu BV^\mu]_F\}$	$E/A(\rho_0)=-16$ MeV, $g_{N\omega}\equiv 3g_8^V Z_\omega$
m_V	$\mathcal{L}_{vec}^{(1)}$	$\frac{1}{2}m_V^2\frac{\chi^2}{\chi_0^2}\mathrm{Tr}\,V_\mu V^\mu$	
μ	$\mathcal{L}_{vec}^{(2)}$	$\frac{1}{4}\mu\,\mathrm{Tr}[V_{\mu\nu}V^{\mu\nu}X^2]$	m_ω, m_ρ, m_ϕ
λ_V	$\mathcal{L}_{vec}^{(3)}$	$\frac{1}{4}\lambda_V(\mathrm{Tr}[V_{\mu\nu}])^2$	
g_4	$\mathcal{L}_{vec}^{(4)}$	$2g_4^4\mathrm{Tr}\,V_\mu V^\mu$	$K\approx 200$–300 MeV
k_0		$-\frac{1}{2}k_0\chi^2 I_2$	$\frac{\partial\Omega}{\partial\sigma}\big\vert_{\mathrm{vac}}=0$
k_1		$k_1(I_2)^2$	$m_N^*/m_N, m_\sigma$
k_2		$k_2 I_4$	$\frac{\partial\Omega}{\partial\zeta}\big\vert_{\mathrm{vac}}=0$
k_3	Scalar potential	$2k_3\chi I_3$	η,η' masses
k_4		$-k_4\chi^4$	$\frac{\partial\Omega}{\partial\chi}\big\vert_{\mathrm{vac}}=0$
δ		$\frac{\delta}{3}\chi^4\ln\frac{I_3}{\det X}$	β_{QCD}
χ_0			$p(\rho_0)=0$
m_π, m_K	$\mathcal{L}_{esb}$	$-\frac{1}{2}\mathrm{Tr}\,A_p[u(X+iY)u+u^\dagger(X+iY)u^\dagger]$	PCAC
σ_0,ζ_0			

$$\frac{\partial(\Omega/V)}{\partial\omega}=-m_\omega^2\omega-4g_4^4\omega^3+g_{N\omega}\rho_N=0. \qquad (32)$$

In the above, ρ_N^s and ρ_N are the scalar and vector densities for the nucleons, which can be calculated analytically for the case of $T=0$, yielding

$$\rho_N^s=\gamma_N\int\frac{d^3k}{(2\pi)^3}\frac{m_N^*}{E_N^*}=\frac{\gamma_N m_N^*}{4\pi^2}\left[k_{FN}E_{FN}^*-m_N^{*2}\ln\left(\frac{k_{FN}+E_{FN}^*}{m_N^*}\right)\right], \qquad (33)$$

$$\rho_N=\gamma_N\int_0^{k_{FN}}\frac{d^3k}{(2\pi)^3}=\frac{\gamma_N k_{FN}^3}{6\pi^2}. \qquad (34)$$

The parameters of the model are constrained by symmetry relations, characteristics of the vacuum or nuclear matter properties. Table I summarizes the various constraints for the parameters within the mean field approach.

IV. RELATIVISTIC HARTREE APPROXIMATION

If we go from the mean field to the Hartree approximation, additional terms in the grand canonical potential appear. These influence the energy, the pressure, and the meson field equations. In the present work, we use the version of the

TABLE II. Parameters for the mean field and the Hartree fits.

Parameter	Mean field		Hartree	
g_4	2.7	0	2.7	0
k_1	1.4	1.4	1.4	1.4
$g_{N\omega}$	12.83	10.52	10.51	9.37
$g_{N\rho}$	4.27	3.51	3.50	3.12
χ_0	402.7	430.1	430.2	446
k_3	−2.64	−2.07	−2.07	−1.73
k_0	2.37	2.07	2.07	1.93
k_2	−5.55	−5.55	−5.55	−5.55
k_4	−0.23	−0.23	−0.23	−0.24
$m_N^*/m_N(\rho_0)$	0.64	0.71	0.73	0.76
m_σ	475.6	560.2	560.4	610.7
K	266.1	359.5	304	377.8
a_4	29.0	27.4	23.9	24.1

chiral model with $g_{N\zeta}=0$, i.e., no coupling of the strange condensate to the nucleon. Hence additional terms will only appear due to summing over baryonic tadpole diagrams due to interaction with the scalar field σ, as in the Walecka model. The additional contribution to the energy density is given as

$$\Delta\epsilon = -\frac{\gamma_N}{16\pi^2}\Bigg[m_N^{*4}\ln\left(\frac{m_N^*}{m_N}\right) + m_N^3(m_N - m_N^*) - \frac{7}{2}m_N^2(m_N - m_N^*)^2 + \frac{13}{3}m_N(m_N - m_N^*)^3 - \frac{25}{12}(m_N - m_N^*)^4\Bigg], \tag{35}$$

where $m_N^* = -g_{N\sigma}\sigma$ and m_N is the nucleon mass in vacuum. This will subsequently modify the pressure and the σ field equations. With inclusion of the relativistic Hartree contributions, the field equation for σ as given by Eq. (30) gets modified to

$$k_0\chi^2\sigma - 4k_1(\sigma^2+\zeta^2)\sigma - 2k_2\sigma^3 - 2k_3\chi\sigma\zeta - 2\frac{\delta\chi^4}{3\sigma} + m_\pi^2 f_\pi + \frac{\partial m_N^*}{\partial\sigma}(\rho_N^s + \Delta\rho_N^s) = 0 \tag{36}$$

where the additional contribution to the nucleon scalar density is given as

$$\Delta\rho_N^s = -\frac{\gamma_N}{4\pi^2}\Bigg[m_N^{*3}\ln\left(\frac{m_N^*}{m_N}\right) + m_N^2(m_N - m_N^*) - \frac{5}{2}m_N(m_N - m_N^*)^2 + \frac{11}{6}(m_N - m_N^*)^3\Bigg]. \tag{37}$$

These make a refitting of some of the parameters necessary. First we have to account for the change in the energy and the pressure, i.e., $g_{N\omega}$ and χ_0 have to be refitted. Due to a change in χ_0 the parameters k_0, k_2, and k_4 must be adapted to ensure that the vacuum equations for σ,ζ, and χ have minima at the vacuum expectation values of the fields. Table II shows the parameters corresponding to the the mean field and the Hartree approximations.

V. VECTOR MESON PROPERTIES IN THE MEDIUM

A. In-medium vector meson masses

We now examine how the Dirac sea effects discussed in Sec. IV modify the masses of the vector mesons. Rewriting the expression for the vector interaction of these mesons given in Eq. (6) in terms of the renormalized couplings $g_{N\omega}$ and $g_{N\rho}$ yields

$$\mathcal{L}_{BV}^N = g_{N\omega}\omega_\mu\overline{\psi_N}\gamma_\mu\psi_N + g_{N\rho}\vec{\rho}_\mu\overline{\psi_N}\gamma_\mu\vec{\tau}\psi_N. \tag{38}$$

Furthermore a tensor coupling is introduced:

$$\mathcal{L}_{tensor} = -\frac{g_{NV}\kappa_V}{2m_N}[\overline{\psi_N}\sigma_{\mu\nu}\tau^a\psi_N\partial^\nu V_a^\mu] \tag{39}$$

where $(g_{NV},\kappa_V)=(g_{N\omega},\kappa_\omega)$ or $(g_{N\rho},\kappa_\rho)$ for $V_a^\mu=\omega^\mu$ or ρ_a^μ, $\tau_a=1$ or $\vec{\tau}$, $\vec{\tau}$ being the Pauli matrices. The vector meson self-energy is given as

$$\Pi_V^{\mu\nu}(k) = -\gamma_I g_{NV}^2\frac{i}{(2\pi)^4}\int d^4p\,\mathrm{Tr}[\Gamma_V^\mu(k)G(p) \times\Gamma_V^\nu(-k)G(p+k)], \tag{40}$$

where $\gamma_I=2$ is the isospin degeneracy factor for nuclear matter, and $\Gamma_V^\mu(k)=\gamma^\mu\tau_a-(\kappa_V/2m_N)\sigma^{\mu\nu}\tau_a$ represents the meson-nucleon vertex function. In the above, $G(k)$ is the interacting nucleon propagator resulting from summing over baryonic tadpole diagrams in the Hartree approximation. This is expressed, in terms of the Feynman and density dependent parts, as

$$G(k) = (\gamma_\mu\overline{k}^\mu + m_N^*)\Bigg[\frac{1}{\overline{k}^2 - m_N^{*2} + i\epsilon} + \frac{i\pi}{E_N^*(k)} \times\delta(\overline{k}^0 - E_N^*(k))\,\theta(k_F - |\vec{\overline{k}}|)\Bigg] \equiv G_F(k) + G_D(k). \tag{41}$$

The vector meson self-energy can then be written as the sum of two parts

$$\Pi^{\mu\nu} = \Pi_F^{\mu\nu} + \Pi_D^{\mu\nu}. \tag{42}$$

In the above, $\Pi_F^{\mu\nu}$ is the contribution arising from the vacuum fluctuation effects, described by the coupling to the $N\overline{N}$ excitations, and $\Pi_D^{\mu\nu}$ is the density dependent contribution to the vector self-energy. For the ω meson, the tensor coupling is generally small as compared to the vector coupling to the nucleons [15]. This is neglected in the present calculations. The Feynman part of the self-energy, $\Pi_F^{\mu\nu}$, is divergent and needs renormalization. We use dimensional regularization to separate the divergent parts. For the ρ meson with tensor interactions, a phenomenological subtraction procedure [14,15] is adopted. After renormalization, the contributions to the meson self-energies from the Feynman part are given as follows. For the ω meson, one arrives at the expression

$$\Pi_F^{\omega}(k^2) \equiv \frac{1}{3}\mathrm{Re}(\Pi_F^{ren})_\mu^\mu = -\frac{g_{N\omega}^2}{\pi^2}k^2\int_0^1 dz\, z(1-z)\ln\left[\frac{m_N^{*2}-k^2z(1-z)}{m_N^2-k^2z(1-z)}\right], \quad (43)$$

and for the ρ meson,

$$\Pi_F^{\rho}(k^2) = -\frac{g_{N\rho}^2}{\pi^2}k^2\left[I_1 + m_N^*\frac{\kappa_\rho}{2m_N}I_2 + \frac{1}{2}\left(\frac{\kappa_\rho}{2m_N}\right)^2(k^2 I_1 + m_N^{*2}I_2)\right] \quad (44)$$

where

$$I_1 = \int_0^1 dz\, z(1-z)\ln\left[\frac{m_N^{*2}-k^2z(1-z)}{m_N^2-k^2z(1-z)}\right],$$

$$I_2 = \int_0^1 dz\,\ln\left[\frac{m_N^{*2}-k^2z(1-z)}{m_N^2-k^2z(1-z)}\right]. \quad (45)$$

The density dependent part for the self-energy is given as

$$\Pi^D(k_0,\boldsymbol{k}\to 0) = -\frac{4g_{NV}^2}{\pi^2}\int p^2dp\, F(|\boldsymbol{p}|,m_N^*)[f_{FD}(\mu^*,T) + \overline{f}_{FD}(\mu^*,T)] \quad (46)$$

with

$$F(|\boldsymbol{p}|,m_N^*) = \frac{1}{\epsilon^*(p)[4\epsilon^*(p)^2-k_0^2]}\left[\frac{2}{3}(2|\boldsymbol{p}|^2+3m_N^{*2}) + k_0^2\left\{2m_N^*\left(\frac{\kappa_V}{2m_N}\right) + \frac{2}{3}\left(\frac{\kappa_V}{2m_N}\right)^2(|\boldsymbol{p}|^2+3m_N^{*2})\right\}\right] \quad (47)$$

where $\epsilon^*(p)=(\boldsymbol{p}^2+m_N^{*2})^{1/2}$ is the effective energy for the nucleon. The effective mass of the vector meson is then obtained by solving the equation, with $\Pi=\Pi_F+\Pi_D$,

$$k_0^2 - (m_V^{eff})^2 + \mathrm{Re}\,\Pi(k_0,\boldsymbol{k}=0) = 0. \quad (48)$$

In the above, the vector meson masses m_V^{eff}, are obtained by inserting the classical expectation values of the meson fields in Eqs. (17), (18), and (19). These correspond to considering only the Fermi part of the baryon propagator and the tree level contribution to the vector meson mass. We compare the vector meson masses including the Dirac polarization to the tree level vector meson mass m_V^{eff}, which neglects the Fermi polarization effects corresponding to nucleon-hole excitations, as well as to the situation when such effects are taken into account.

B. Meson decay properties

We next proceed to study the vector meson decay widths as modified due to the effect of vacuum polarization effects through the RHA. The decay width for the process $\rho\to\pi\pi$ is calculated from the imaginary part of the self-energy and in the rest frame of the ρ meson, it becomes

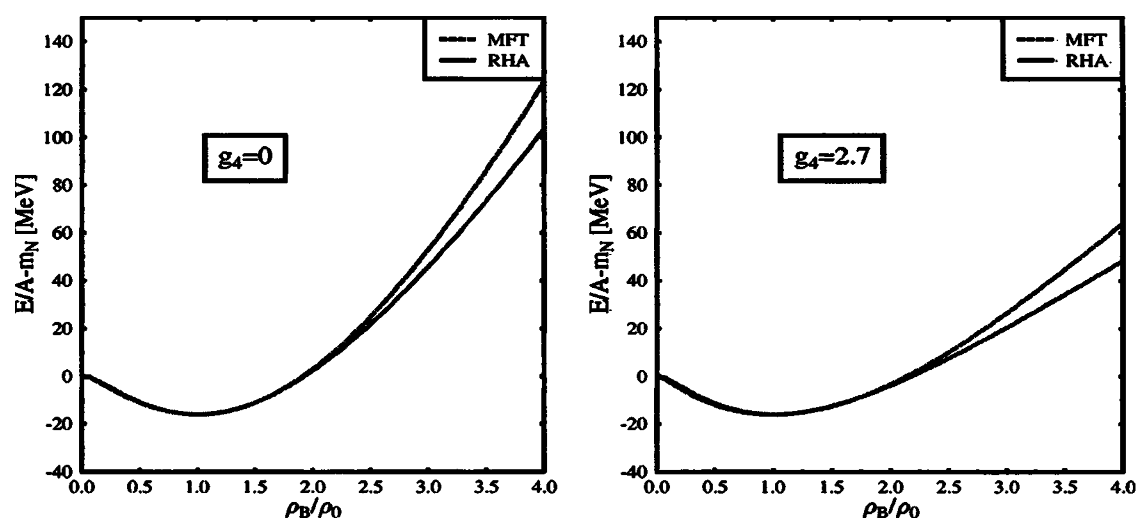

FIG. 1. Binding energy per particle as a function of density in the mean field and in the Hartree approximations.

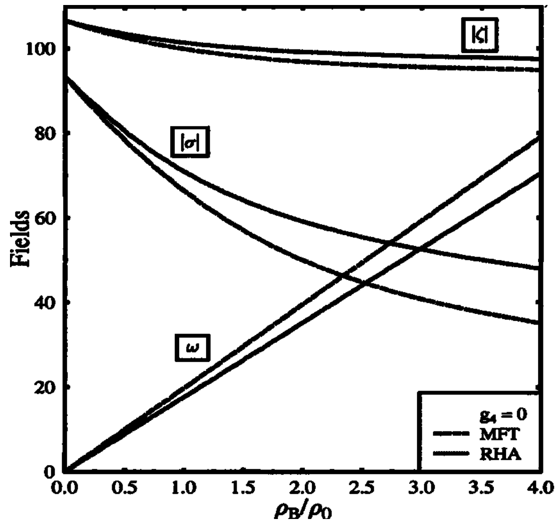

FIG. 2. Scalar fields σ, ζ and vector field ω as a function of density in the mean field and in the Hartree approximations.

$$\Gamma_\rho(k_0) = \frac{g_{\rho\pi\pi}^2}{48\pi}\frac{(k_0^2 - 4m_\pi^2)^{3/2}}{k_0^2}\left\{\left[1 + f\left(\frac{k_0}{2}\right)\right]\left[1 + f\left(\frac{k_0}{2}\right)\right] - f\left(\frac{k_0}{2}\right)f\left(\frac{k_0}{2}\right)\right\} \tag{49}$$

where $f(x) = [e^{\beta x} - 1]^{-1}$ is the Bose-Einstein distribution function. The first and the second terms in the above equation represent the decay and the formation of the resonance ρ. The medium effects have been shown to play a very important role for the ρ-meson decay width. In the calculation for the ρ decay width, the pion has been treated as free, i.e., any modification of the pion propagator due to effects like delta–nucleon-hole excitation [31] have been neglected. The coupling $g_{\rho\pi\pi}$ is fixed from the decay width of the ρ meson in vacuum ($\Gamma_\rho = 151$ MeV) decaying into two pions.

For the nucleon-rho couplings, the vector and tensor couplings as obtained from the *N*-*N* forward dispersion relation [15,17,32] are used. With the couplings as described above, we consider the modification of ω- and ρ-meson properties in nuclear matter due to quantum correction effects.

To calculate the decay width for the ω meson, we consider the following interaction Lagrangian for the ω meson [33–35]:

$$\mathcal{L}_\omega = \frac{g_{\omega\pi\rho}}{m_\pi}\epsilon_{\mu\nu\alpha\beta}\partial^\mu\omega^\nu\partial^\alpha\rho_i^\beta\pi_i + \frac{g_{\omega 3\pi}}{m_\pi^3}\epsilon_{\mu\nu\alpha\beta}\epsilon_{ijk}\omega^\mu\partial^\nu\pi^i\partial^\alpha\pi^j\partial^\beta\pi^k. \tag{50}$$

The decay width of the ω meson in vacuum is dominated by the channel $\omega \to 3\pi$. In the medium, the decay width for $\omega \to 3\pi$ is given as

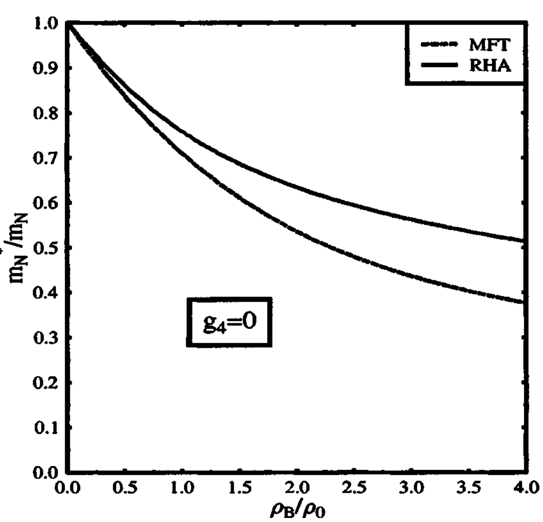

FIG. 3. Effective nucleon mass as a function of density in the mean field and in the Hartree approximations.

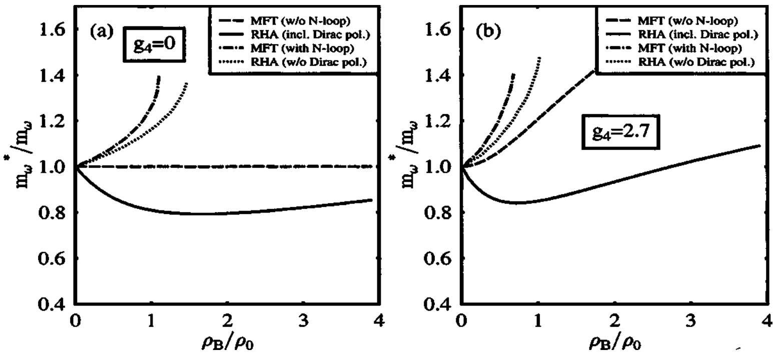

FIG. 4. Effective ω-meson mass in the mean field approximation and including the Hartree contributions. Left: no quartic vector-meson interaction. Right: including ω^4 interaction. There is a significant drop of the vector meson mass due to the Dirac sea effect, which is not seen in the mean field approximation.

$$\Gamma_{\omega\to 3\pi} = \frac{(2\pi)^4}{2k_0}\int d^3\tilde{p}_1 d^3\tilde{p}_2 d^3\tilde{p}_3 \delta^{(4)}(P - p_1 - p_2 - p_3) \times |M_{fi}|^2\{[1 + f(E_1)][1 + f(E_2)][1 + f(E_3)] - f(E_1)f(E_2)f(E_3)\}, \quad (51)$$

where $d^3\tilde{p}_i = d^3p_i/(2\pi)^3 2E_i$, p_i and the E_i's are four-momenta and energies for the pions, and the $f(E_i)$'s are their thermal distributions. The matrix element M_{fi} has contributions from the channels $\omega\to\rho\pi\to 3\pi$ [described by the first term in Eq. (50)] and the direct decay $\omega\to 3\pi$ resulting from the contact interaction [second term in Eq. (50)] [35–37]. For the $\omega\rho\pi$ coupling we take the value $g_{\omega\rho\pi}=2$ which is compatible with the vacuum decay width $\omega\to\pi\gamma$ [17]. We fix the point interaction coupling $g_{\omega 3\pi}$ by fitting the partial decay width $\omega\to 3\pi$ in vacuum (7.49 MeV) to be 0.24. The contribution arising from the direct decay turns out to be marginal, being of the order of up to 5% of the total decay width for $\omega\to 3\pi$.

With the modifications of the vector meson masses in the hot and dense medium, a new channel becomes accessible, the decay mode $\omega\to\rho\pi$ for $m_\omega^* > m_\rho^* + m_\pi$. This has been taken into account in the present investigation.

VI. RESULTS AND DISCUSSION

We shall now discuss the results of the present investigation: the nucleon properties as modified due to the Dirac sea contributions through the relativistic Hartree approximation and their effects on vector meson properties in the dense hadronic matter. Figure 1 shows the equation of state in the mean field and in the Hartree approximation with and without quartic self-interaction for the ω field. In both cases we observe that the additional terms resulting from the Hartree approximation lead to a softening of the equation of state at higher densities. However, the compressibility in the relativistic Hartree approximation is higher than the mean field value, as shown in Table II. Furthermore, the influence of the finite value for the quartic ω coupling g_4 is clearly visible. In this case the compressibility at nuclear saturation is strongly reduced (Table II). Also, the resulting equation of state is much softer in particular at higher densities. The reason for this can be seen from Fig. 2. The vector field ω, which causes the repulsion in the system, rises much more steeply as a function of density for $g_4=0$ than for the case of $g_4=2.7$, because the quartic self-interaction attenuates the ω field.

The effective nucleon mass for the different cases is depicted in Fig. 3. Here the RHA predicts higher nucleon masses than the mean field (MF) case. At higher densities these contributions become increasingly important. This is also reflected in the density dependence of the nonstrange σ field, showing a considerable increase due to the Hartree contributions (Fig. 2). In contrast, the strange condensate ζ, which does not couple to the nucleons, takes only slightly lower values in the MF case.

The in-medium properties of the vector mesons are modified due to the vacuum polarization effects. The nucleon-ω vector coupling $g_{N\omega}$ is calculated from the nuclear matter saturation properties. As already stated, the N-ω tensor coupling is neglected. Figure 4 shows the resulting modification of the ω-meson mass in the Hartree approximation as compared to the mean field case. For $g_4=0$, the tree level ω mass has no density dependence, because of the frozen glueball approximation. However, including the effect of the nucleon loop, the fluctuation of the Fermi sea, corresponding to the particle-hole excitations, leads to an increase in the ω mass.

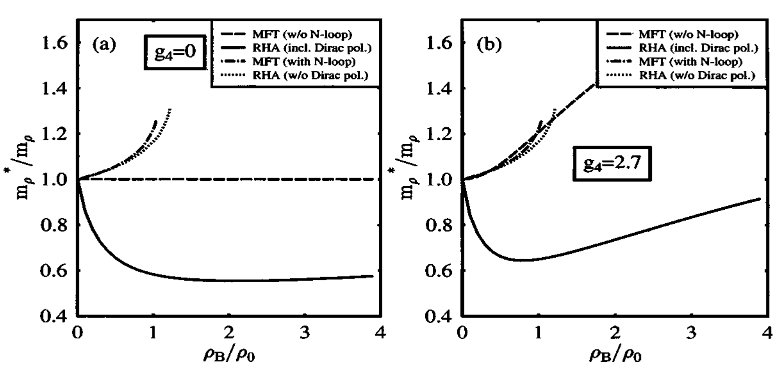

FIG. 5. Effective ρ-meson mass without and with the Hartree contributions, with the nucleon-rho vector and tensor couplings, as fitted from the NN scattering data ($g_{N\rho}=2.63$, $\kappa_\rho=6.1$). The Hartree approximation gives rise to the decrease of the ρ mass in the medium.

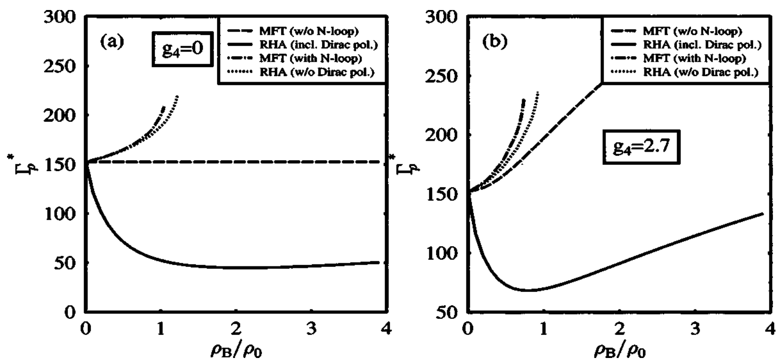

FIG. 6. Decay width of ρ meson in the absence and presence of the Dirac sea effect with the couplings fitted from the NN scattering data.

As can be seen from Eq. (46), the difference in the contribution due to the Fermi polarization in the MF theory (MFT) and RHA at zero temperature, is purely due to difference in the effective nucleon mass and the N-ω coupling parameters. At low densities, these contributions are seen to be very similar in both the approaches. In the low density approximation, the Fermi part of the vector self-energy (assuming the tensor coupling to be zero) reduces to

$$\Pi_D = -\frac{4g_{VN}^2 M_N}{4M_N^2 - k_0^2}\rho_B. \tag{52}$$

Making the further assumption that the effective omega mass $k_0 \ll 2M_N$, which may be a valid assumption at low densities, while considering the Fermi polarization effects, one arrives at the expression for the density dependent vector self-energy as

$$\Pi_D = -\frac{g_{VN}^2}{M_N}\rho_B, \tag{53}$$

which is identical to the vector self-energy due to Fermi polarization, as obtained earlier [16]. At low densities, the ω mass increases with density. However, it is seen that in the mean field approximation and in the absence of a quartic coupling, no real solution for k_0 exists for the self-consistency equation (48) for the mass of the ω meson, for a density above $1.1\rho_0$. This feature remains also in the RHA as well as when $g_4 \neq 0$. This can be undertstood qualitatively in the following manner. At higher densitites, the assumption $k_0 \ll 2M_N$ is no longer valid, and one should solve the self-consistency equation

$$k_0^2 - (m_V^{eff})^2 - \frac{4g_{VN}^2 M_N}{4M_N^2 - k_0^2}\rho_B = 0, \tag{54}$$

which does not have a real solution for k_0^2 (and hence for k_0) when $[(m_V^{eff})^2 - 4M_N^2]^2 < 16 M_N g_{VN}^2 \rho_B$. In other words, there a critical density $(\rho_B)_{crit} = [(m_V^{eff})^2 - 4M_N^2]^2/(16 g_{VN}^2 M_N)$, above which no real solution for the medium modified vector meson mass k_0 exists. Solving the self-consistency equation as given by Eq. (48) and including only the Fermi polarization part as given by Eq. (46) retains the same qualitative behavior. In contrast, a strong reduction due to the Dirac sea polarization is found for densities up to around normal nuclear matter density. At higher densities, the Fermi polarization part of the ω self-energy starts to become important, leading to an increase in the mass. A similar behavior has been observed in the Walecka model [14–16]. The quartic term in the ω field considerably enhances the ω mass with increasing density. Thus in the mean field case, the mass rises monotonically. For the Hartree approximation a decrease of the ω mass for small densities can still be found. But at higher densities the contribution from the quartic term becomes more important and leads to an increase of the in-medium mass.

In Fig. 5, we illustrate the medium modification for the ρ-meson mass with the vector and tensor couplings to the nucleons being fixed from the NN forward dispersion relation [15,17,32]. The values for these couplings are given as $g_{N\rho}^2/4\pi = 0.55$ and $\kappa_\rho = 6.1$. We notice that the decrease in the ρ meson with increasing density is much sharper than that of the ω meson. Such a behavior of the ρ meson undergoing a much larger medium modification was also observed earlier [17] within the relativistic Hartree approximation in the Walecka model. This indicates that the tensor coupling, which is

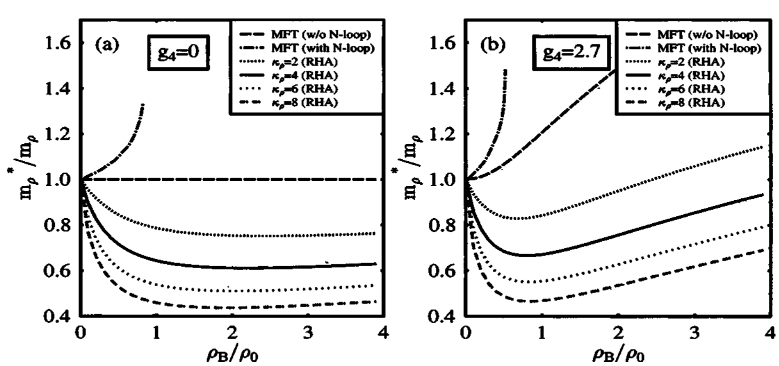

FIG. 7. Effective ρ-meson mass without and with the Hartree contributions, with the nucleon-rho vector coupling $g_{N\rho}$ as from the chiral model, which is compatible with the symmetry energy. Since we do not know the medium dependent tensor coupling κ_ρ, it is taken as a parameter. The Hartree approximation gives rise to the decrease of the ρ mass in the medium, which is seen to be quite sensitive to the nucleon-rho tensor coupling.

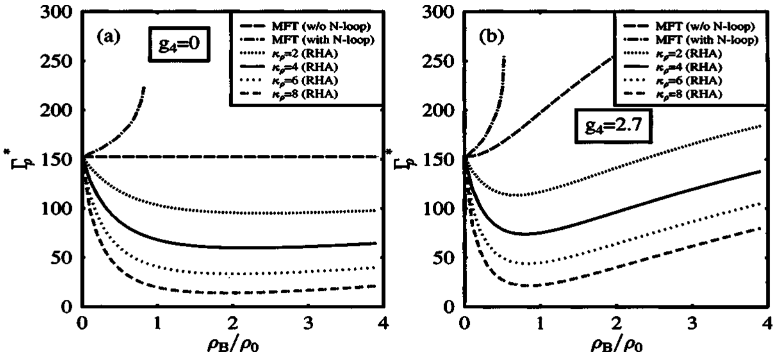

FIG. 8. Decay width of ρ meson in the absence and presence of the Dirac sea effect with the nucleon-rho vector coupling $g_{\rho N}$ as from the chiral model, which is compatible with the symmetry energy. The tensor coupling κ_ρ is taken as a parameter.

negligible for the ω meson, plays a significant role for the ρ meson. The Fermi polarization effects are seen to give a rise in the ρ mass with density, similar to what was seen for the ω meson mass.

The in-medium decay width for $\rho \rightarrow \pi\pi$, Γ^*_ρ reflects the behavior of the in-medium ρ mass. This is because in the present work only the case $T=0$ is considered, and so there is no Bose-enhancement effect. Therefore in the absence of the quartic vector-meson interaction, the significant drop of the mass of the ρ meson in the medium leads to a decrease of the ρ decay width in the relativistic Hartree approximation. However, only the Fermi polarization effects in the RHA or MFT lead to an increase in ρ decay width, which is a reflection of the increase in the mass of the ρ meson in the medium. These are shown in Fig. 6. Since the quartic self-interaction yields an increase in the mass at higher densities, it also leads to an increase of the ρ decay width.

In the previous calculations, the ρ-N coupling strengths were used as determined from the NN forward scattering data [32]. Now we consider the mass modification for the ρ meson, with the nucleon ρ coupling $g_{N\rho}$ as determined from the symmetry relations (Table II). The symmetry energy coefficient a_{sym} is given as [38]

$$a_{sym} = \frac{1}{2}\left[\frac{\partial^2}{\partial t^2}\left(\frac{\epsilon}{\rho}\right)\right]_{t=0}, \tag{55}$$

where $t=(\rho_n-\rho_p)/\rho_0$. The resulting values for the symmetry energy for the different cases are shown in Table II. They are compatible with the experiment. We take the tensor coupling as a parameter in our calculations since this coupling cannot be fixed from infinite nuclear matter properties. However, it influences the properties of finite nuclei. In a recent study, the importance of the tensor couplings (of ω and ρ) for the description of finite nuclei in the RHA has been discussed [39]. It is seen that the spectra of the shell model states are improved considerably with inclusion of these interactions, which are otherwise not well described in the RHA due to softening of the equation of state. The resulting in-medium mass of the ρ meson is plotted in Fig. 7 as a function of baryon density ρ_B/ρ_0. It is observed that the ρ-meson mass has a strong dependence on the tensor coupling. In the Hartree approximation the ρ-nucleon vector coupling does not differ too much in the two cases, i.e., depending on whether it is obtained from NN scattering data or from the symmetry relations. Thus we find a similar behavior for the ρ mass, if in the latter case we choose $\kappa_\rho=6$, i.e., close to the value from scattering data. For the same value of κ_ρ and the ρ-N vector coupling parameter as fitted from symmetry energy, the ρ mass is plotted in the mean field approximation. As already seen, the Fermi polarization gives rise to an increase in the ρ-meson mass.

Figure 8 shows the decay width for the ρ meson when we take the $N\rho$ vector coupling as determined from symmetry relations and the tensor coupling is taken as a parameter. In the mean field approximation, κ_ρ is chosen as 6, a value close to that from scattering data.

The decay width of the ω meson is plotted as a function of density in Fig. 9. In the vacuum the process $\omega \rightarrow 3\pi$ is the dominant decay mode. However, in the medium the channel $\omega \rightarrow \rho\pi$ also opens up, since the ρ meson has a stronger drop in the medium as compared to the ω-meson mass.

The mean field approximation, does not have a contribution from the latter decay channel, whereas the inclusion of the relativistic Hartree approximation permits both processes in the medium. In the presence of the quartic vector meson interaction, the channel $\omega \rightarrow \rho\pi$, which opens up at around $0.3\rho_0$, no longer remains kinematically accessible at higher

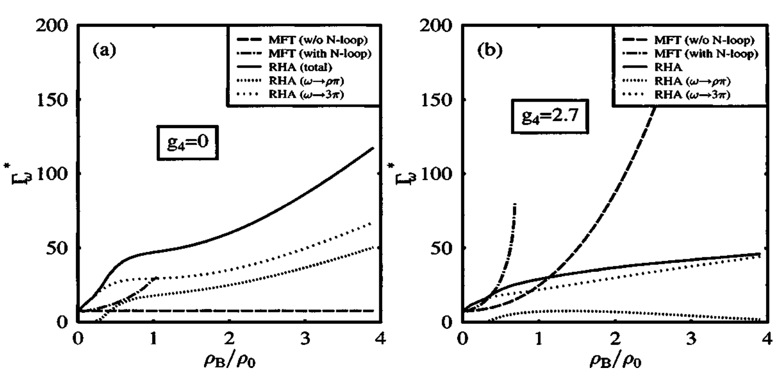

FIG. 9. Effective decay width of ω meson without and with the Hartree contributions. The decay width has contributions from $\omega \rightarrow 3\pi$ as well as $\omega \rightarrow \rho\pi$. The latter becomes accessible due to stronger medium modification of the ρ-meson mass as compared to the ω mass. The MFT has no contribution from the process $\omega \rightarrow \rho\pi$.

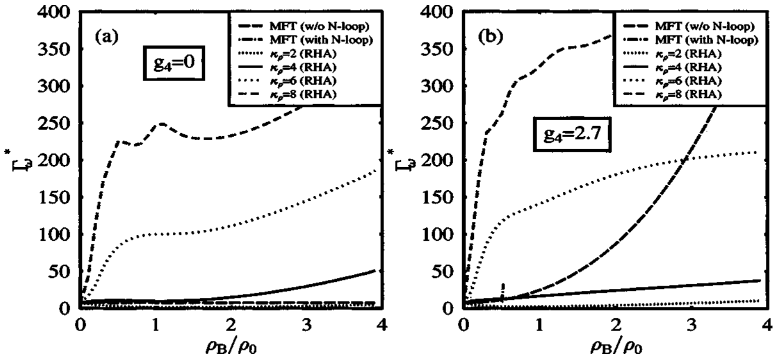

FIG. 10. Decay width of ω meson in the absence and presence of the Dirac sea effect. For the channel $\omega \rightarrow \rho\pi$, the ρ-meson properties are determined with the nucleon-rho vector coupling $g_{N\rho}$ as from the chiral model, which is compatible with the symmetry energy, and the tensor coupling κ_ρ is taken as a parameter.

densities. This is due to the increased importance of ω^4 contributions to the vector meson masses at high densities.

The strong enhancement of the ω-meson mass in the presence of a quartic self-interaction term for the ω field makes also the decay channel $\omega \rightarrow N\bar{N}$ kinematically accessible in the mean field approximation. In the present investigation the vector meson properties are considered at rest. Vector mesons with a finite three-momentum can also have additional decay channels to particle-hole pairs. These decay modes, e.g., have significant contributions, to the Δ decay width [40]. Additional channels that open up in the mean field approximation in the presence of a quartic term for ω, however, have not been taken into consideration in the present work. Here the emphasis is on the effect due to the relativistic Hartree approximation on the in-medium vector meson properties.

Figure 10 illustrates the decay width of the ω meson when the medium dependence of the ρN vector coupling is taken into account and the tensor coupling κ_ρ is taken as a parameter. In the mean field approximation, the κ_ρ chosen is 6, which is close to the value obtained from NN scattering data. The strong dependence of the ρ-meson properties on the tensor coupling in the relativistic Hartree approximation are reflected in the ω decay width through the channel $\omega \rightarrow \rho\pi$.

VII. SUMMARY

To summarize, in the present paper we have considered the modification of the vector meson properties due to vacuum polarization effects arising from the Dirac sea in nuclear matter in the chiral SU(3) model. The baryonic properties as modified due to such effects determine the vector meson masses in dense hadronic matter. A significant reduction of these masses in the medium is found, where the Dirac sea contribution dominates over the Fermi sea part. This shows the importance of the vacuum polarization effects for the vector meson properties, as has been emphasized earlier within the framework of quantum hadrodynamics [15,16].

The main aim of the present paper is to study the effcts of vacuum polarizations from the nucleon sector on the vector meson properties. The drastically different behavior of the vector meson masses in the two approaches (RHA and MFT) is due to the Dirac sea polarization effect leading to the large drop of the mass of the vector meson in the RHA, and is absent in the MFT. This effect is genuine to the RHA approach and does not reflect a mere rescaling of the MFA coupling strengths of the rho meson. The Fermi sea polarization is seen to be dominated by the Dirac sea polarization in the RHA, leading to a drop in the vector meson mass in the medium, whereas the particle-hole excitations in the MFT corresponding to the Fermi sea polarization give a rise in the vector meson mass [13].

The ρ-meson mass is seen to have a sharper drop as compared to the ω-meson mass in the medium. This reflects the fact that the vector-meson–nucleon tensor coupling, which is absent for the ω meson, plays an important role for the ρ mass. The decay width of $\rho \rightarrow \pi\pi$ is modified appreciably due to the modification of the ρ mass. At finite baryon densities, the scattering due to nucleons is seen to lead to a large increase in the ρ-meson decay width [7,11,42]. This, however, has not been taken into account in the present investigation.

The effects discussed above influence observables in finite nuclei, stellar objects, and relativistic heavy ion collisions. For example, the modified vector meson properties in a medium play an important role in the dilepton emission rates in relativistic heavy ion collisions [41]. This is reflected by the shift and broadening of the peaks in the low invariant mass regime in the dilepton spectra. Therefore, it will be important to investigate how the dilepton rates are modified by the in-medium vector meson properties in hot and dense hadronic matter. Generalization to finite temperatures [43] to study the spectral properties of the vector mesons and their effects on the dilepton spectra from the hot hadronic matter resulting from nuclear collisions are currently being studied [44]. This apart, the study of Hartree contributions in the analysis of the particle ratios from the relativistic heavy ion collision experiments [45] and related problems are also under investigation.

ACKNOWLEDGMENTS

One of the authors (A.M.) is grateful to J. Reinhardt for fruitful discussions and Institut fuer Theoretische Physik for warm hospitality. This work is supported by Deutsche Forschungsgemeinschaft (DFG), Gesellschaft für Schwerionenforschung (GSI), Bundesministerium für Bildung und Forschung (BMBF), the Graduiertenkolleg Theoretische und Experimentelle Schwerionenphysik, and the U.S. Department of Energy, Nuclear Physics Division (Contract No. W-31-109-Eng-38).

[1] HELIOS-3 Collaboration, N. Masera, Nucl. Phys. **A590**, 93c (1995).
[2] CERES Collaboration, G. Agakichiev *et al.*, Phys. Rev. Lett. **75**, 1272 (1995); Phys. Lett. B **422**, 405 (1998); Nucl. Phys. **A661**, 23c (1999).
[3] DLS Collaboration, R. J. Porter *et al.*, Phys. Rev. Lett. **79**, 1229 (1997); DLS Collaboration, W. K. Wilson *et al.*, Phys. Rev. C **57**, 1865 (1998).
[4] PHENIX Collaboration, D. P. Morrison, Nucl. Phys. **A638**, 565c (1998).
[5] HADES Collaboration, J. Stroth, Advances in Nuclear Dynamics **5**, 311 (1999).
[6] G. E. Brown and M. Rho, Phys. Rev. Lett. **66**, 2720 (1991).
[7] R. Rapp and J. Wambach, Adv. Nucl. Phys. **25**, 1 (2000); R. Rapp, G. Chanfray, and J. Wambach, Nucl. Phys. **A617**, 472 (1997).
[8] T. Hatsuda and Su H. Lee, Phys. Rev. C **46**, R34 (1992); T. Hatsuda, S. H. Lee, and H. Shiomi, *ibid.* **52**, 3364 (1995).
[9] X. Jin and D. B. Leinweber, Phys. Rev. C **52**, 3344 (1995); T. D. Cohen, R. D. Furnstahl, D. K. Griegel, and X. Jin, Prog. Part. Nucl. Phys. **35**, 221 (1995); R. Hofmann, Th. Gutsche, and A. Faessler, Eur. Phys. J. C **17**, 651 (2000); S. Mallik and K. Mukherjee, Phys. Rev. D **58**, 096011 (1998).
[10] S. Mallik and A. Nyffeler, Phys. Rev. C **63**, 065204 (2001).
[11] F. Klingl, N. Kaiser, and W. Weise, Nucl. Phys. **A624**, 527 (1997).
[12] C. Ernst, S. A. Bass, M. Belkacem, H. Stöcker, and W. Greiner, Phys. Rev. C **58**, 447 (1998).
[13] B. D. Serot and J. D. Walecka, Adv. Nucl. Phys. **16**, 1 (1986); S. A. Chin, Ann. Phys. (N.Y.) **108**, 301 (1977).
[14] H. Shiomi and T. Hatsuda, Phys. Lett. B **334**, 281 (1994).
[15] T. Hatsuda, H. Shiomi, and H. Kuwabara, Prog. Theor. Phys. **95**, 1009 (1996).
[16] H.-C. Jean, J. Piekarewicz, and A. G. Williams, Phys. Rev. C **49**, 1981 (1994); K. Saito, K. Tsushima, A. W. Thomas, and A. G. Williams, Phys. Lett. B **433**, 243 (1998).
[17] Jan-e Alam, S. Sarkar, P. Roy, B. Dutta-Roy, and B. Sinha, Phys. Rev. C **59**, 905 (1999).
[18] A. Mishra, J. C. Parikh, and W. Greiner, J. Phys. G **28**, 151 (2002).
[19] A. Mishra, J. Reinhardt, H. Stoecker, and W. Greiner, Phys. Rev. C **66**, 064902 (2002).
[20] A. Mishra, P. K. Panda, S. Schramm, J. Reinhardt, and W. Greiner, Phys. Rev. C **56**, 1380 (1997); A. Mishra, P. K. Panda, and W. Greiner, J. Phys. G **27**, 1561 (2001); **28**, 67 (2002).
[21] P. Papazoglou, D. Zschiesche, S. Schramm, J. Schaffner-Bielich, H. Stöcker, and W. Greiner, Phys. Rev. C **59**, 411 (1999).
[22] D. Zschiesche, P. Papazoglou, S. Schramm, C. Beckmann, J. Schaffner-Bielich, H. Stocker, and W. Greiner, *Symmetries in Intermediate and High Energy Physics*, edited by A. Faessler, T. S. Kosmos, and G. K. Leontaris, Vol. 163 of Springer Tracts in Modern Physics (Springer, Berlin, 2000), p. 129.
[23] R. J. Furnstahl, B. D. Serot, and H-B. Tang, Nucl. Phys. **A615**, 441 (1997); **A618**, 446 (1997).
[24] J. J. Sakurai, *Currents and Mesons* (University of Chicago Press, Chicago, 1969).
[25] S. Gasiorowicz and D. Geffen, Rev. Mod. Phys. **41**, 531 (1969).
[26] P. K. Mitter and L. J. Swank, Nucl. Phys. **B8**, 205 (1968).
[27] Y. Sugahara and H. Toki, Nucl. Phys. **A579**, 557 (1994).
[28] J. Schechter and Y. Ueda, Phys. Rev. D **3**, 168 (1971).
[29] J. Schechter, Phys. Rev. D **21**, 3393 (1980).
[30] B. D. Serot and J. D. Walecka, Int. J. Mod. Phys. E **6**, 515 (1997).
[31] M. Asakawa, C. M. Ko, P. Levai, and X. J. Qiu, Phys. Rev. C **46**, R1159 (1992); M. Herrmann, B. L. Friman, and W. Nörenberg, Nucl. Phys. **A560**, 411 (1993); G. Chanfray and P. Shuck, *ibid.* **A545**, 271c (1992).
[32] W. Grein, Nucl. Phys. **B131**, 255 (1977); W. Grein and P. Kroll, Nucl. Phys. **A338**, 332 (1980).
[33] J. J. Sakurai, *Currents and Mesons* (University of Chicago Press, Chicago, 1969).
[34] M. Gell-Mann, D. Sharp, and W. D. Wagner, Phys. Rev. Lett. **8**, 261 (1962).
[35] B. A. Li, Phys. Rev. D **52**, 5165 (1995).
[36] F. Klingl, N. Kaiser, and W. Weise, Z. Phys. A **356**, 193 (1996).
[37] Ö. Kaymakcalan, S. Rajeev and I. Schechter, Phys. Rev. D **30**, 594 (1984).
[38] N. K. Glendenning, F. Weber, and S. A. Moszkowski, Phys. Rev. C **45**, 844 (1992).
[39] G. Mao, Phys. Rev. C **67**, 044318 (2003).
[40] H. Kim, S. Schramm, and S. H. Lee, Phys. Rev. C **56**, 1582 (1997).
[41] C. Gale and J. I. Kapusta, Nucl. Phys. **B357**, 65 (1991).
[42] T. Renk and A. Mishra, Phys. Rev. C **69**, 054905 (2004).
[43] A. Mishra, K. Balazs, D. Zschiesche, S. Schramm, H. Stöcker, and W. Greiner, Phys. Rev. C **69**, 024903, 2004.
[44] A. Mishra, K. Balazs, D. Zschiesche, S. Schramm, H. Stöcker, and W. Greiner (unpublished).
[45] D. Zschiesche, S. Schramm, J. Schaffner-Bielich, H. Stöcker, and W. Greiner, Phys. Lett. B **547**, 7 (2002).