Collective modes of asymmetric nuclear matter in quantum hadrodynamics

V. Greco,[1] M. Colonna,[1] M. Di Toro,[1] and F. Matera[2]
[1]*Laboratorio Nazionale del Sud, Via S. Sofia 44, I-95123 Catania, Italy and Dipartimento di Fisica, Università degli Studi di Catania, Catania, Italy*
[2]*Dipartimento di Fisica, Università degli Studi di Firenze and Istituto Nazionale di Fisica Nucleare, Sezione di Firenze, Via G. Sansone 1, I-50019, Sesto F.no (Firenze), Italy*
(Received 15 May 2002; published 15 January 2003)

We discuss a fully relativistic Landau Fermi liquid theory based on the quantum hadrodynamics effective field picture of nuclear matter. From the linearized kinetic equations we get the dispersion relations of the propagating collective modes. We focus our attention on the dynamical effects of the interplay between scalar and vector channel contributions. An interesting "mirror" structure in the form of the dynamical response in the isoscalar-isovector degree of freedom is revealed, with a complete parallelism in the role respectively played by the compressibility and the symmetry energy. All that seems to support the introduction of an explicit coupling to the scalar-isovector channel of the nucleon-nucleon interaction. In particular we study the influence of this coupling (to a δ-meson-like effective field) on the collective response of asymmetric nuclear matter (ANM). Interesting contributions are found on the propagation of isovectorlike modes at normal density and on an expected smooth transition to isoscalarlike oscillations at high baryon density. Important "chemical" effects on the neutron-proton structure of the mode are shown. For dilute ANM we have the isospin distillation mechanism of the unstable isoscalarlike oscillations, while at high baryon density we predict an almost pure neutron wave structure of the propagating sounds.

DOI: 10.1103/PhysRevC.67.015203 PACS number(s): 24.10.Cn, 24.10.Jv, 21.30.Fe, 24.30.Cz

I. INTRODUCTION

The quantum hadrodynamics (QHD) effective field model represents a very successful attempt to describe, in a fully consistent relativistic picture, equilibrium and dynamical properties of nuclear systems at the hadronic level [1–3]. Very nice results have been obtained for the nuclear structure of finite nuclei [4–6], for the nuclear matter (NM) equation of state and liquid-gas phase transitions [7] and for the dynamics of nuclear collisions [8,9]. Relativistic random-phase-approximation (RRPA) theories have been developed to study the nuclear collective response [10–15].

In this paper we present a relativistic linear response theory with the aim of a transparent connection between the collective dynamics and the coupling to various channels of the nucleon-nucleon interaction. In particular we will focus our attention on the dynamical response of asymmetric nuclear matter since one of the main points of our discussion is the relevance of the coupling to a scalar isovector channel, the virtual $\delta[a_0(980)]$ meson, not considered in the usual dynamical studies. A related point of interest is the dynamical treatment of the Fock terms.

The isospin physics is assuming more and more relevance in connection to the new radioactive beam facilities and to nuclear astrophysics. The introduction of the isovector-scalar channel in covariant approaches can play a key role in the effective interaction in asymmetric matter; see Ref. [16]. This point has not received great attention before for two main reasons.

(i) The δ channel has not been considered *a priori*, just on the basis of the weak contribution to the free nucleon-nucleon interaction [17,18]. But in the spirit of the *effective field theory* as a relativistic *density functional theory* (the EFT-DFT framework; see [19]), the relevance of this channel could be completely different in nuclear matter, due to medium and many-body effects. In particular we can expect a large contribution from exchange terms of the strongly coupled isoscalar channels; see the discussion in Ref. [20] as well as in [16].

(ii) This extension is not well supported by the existing set of data, as remarked in Refs. [6,21–23]. Clearly these negative outcomes are mainly derived from the lack of information on observables more sensitive to the density dependence of the symmetry term.

We like to note that very recently (see the conclusions of Ref. [24]), the δ channel has been reconsidered as an interesting improvement of covariant approaches, in the framework of the EFT-DFT philosophy. One of the main tasks of our work is just to try to select the dynamical observables more sensitive to it; see also the conclusions of Ref. [16].

In this respect the results reported here on the collective response can be useful in order to solve the open problem of the determination of the scalar-isovector coupling. As already remarked, contributions to this channel mainly come from correlation effects. Therefore the correct microscopic approach should be to derive the relative coupling constant, in a QHD mean field framework, from Dirac-Brueckner-Hartree-Fock calculations. Several attempts have been recently performed (see Refs. [23,25,26]), but the results have been up to now not fully model independent.

An important outcome of our work is to show that the two effective couplings, vector and scalar, in the isovector channel influence in a different way the static (symmetry energy) and dynamic (collective response) properties of asymmetric nuclear matter. This will open new possibilities for a phenomenological determination of these fundamental quantities.

In the paper we will derive transparent analytical results.

TABLE I. Parameter sets.

Parameter	Set I	Set II	NLHF	NL3
f_σ (fm^2)	11.27	same	9.15	15.73
f_ω (fm^2)	6.48	same	3.22	10.53
f_ρ (fm^2)	1.0	2.8	1.9	1.34
f_δ (fm^2)	0.00	2.0	1.4	0.00
A (fm^{-1})	0.022	same	0.098	0.01
B	-0.0039	same	-0.021	-0.003

In order to show also some quantitative effects of the dynamical contribution of the δ channel we have to fix in some way the corresponding coupling. We have used a constant value (see Table I) extracted from the analysis of Ref. [23], where it actually appears not strongly density dependent in a wide range of baryon densities. Some results are also presented with the dynamics of Fock correlations explicitly accounted for [nonlinear Hartree-Fock (NLHF) case; see Sec. II]. Now all the coupling constants will acquire some density dependence [20].

A relativistic extension of the Landau linear response theory of Fermi liquids has been considered before just starting from the relativistic form of the Landau parameters [27–29]. We will show that the full dispersion relations obtained from the relativistic kinetic equations present some interesting corrections that cannot be neglected.

The main physics results are the following.

(i) The important effect of a δ-meson coupling on the isovector collective mode at saturation baryon density. This is of interest for the relativistic study of the giant dipole resonance in heavy finite nuclei. It is important to note that the inclusion of Fock terms acts in the same direction.

(ii) The presence of noticeable "chemical effects" in the propagating collective oscillations; i.e., the charge symmetry of the "waves" is quite different from the asymmetry of the initial equilibrium matter. The effect is opposite for the unstable modes present at low densities, more proton rich, and leads to the isospin distillation effect, and for the stable propagating sounds at high baryon density which appear mostly like "pure neutron waves."

(iii) A stimulating "mirror" structure of the isoscalar and isovector linear response, with the restoring forces given by the potential part respectively of the compressibility and the symmetry energy. The interplay between the scalar and vector meson effective fields in the dynamics is very similar for the two degrees of freedom, as already observed for static properties [16]. The conclusion is that a formally consistent relativistic effective field model should include on the same footing isoscalar and isovector meson fields, *both* scalar and vector.

Our results around normal density can be used as general guidelines in predicting the behavior of volume collective modes in finite β-unstable nuclei. Similar study for asymmetric NM have been performed in Ref. [30] using Skyrme-like interactions. Apart the difference in the interactions used, in particular for the symmetry terms, we will see similar results and interesting new relativistic effects.

In Sec. II we derive the kinetic equations in the general case of nonlinear self-interacting terms, including the Fock corrections. We discuss the inclusion of the δ-meson channel, also on the model parameters. In Sec. III we present the relativistic linear response equations. In Sec. IV we have a general discussion on the formal structure of the dispersion relations, the role played by the scalar-vector mesons, and the comparison to nonrelativistic cases. Results for isovector(*like*) collective modes are presented in Sec. V, in particular for the asymmetry and baryon density effects. The isoscalar(*like*) response is analyzed in Sec. VI. Conclusions and outlooks can be found in Sec. VII.

II. KINETIC EQUATIONS FROM A QHD EFFECTIVE THEORY

We start from the QHD effective field picture of the hadronic phase of nuclear matter [1–3]. In order to include the main dynamical degrees of freedom of the system we will consider the nucleons coupled to the isoscalar scalar σ and vector ω mesons and to the isovector scalar δ and vector ρ mesons.

The Lagrangian density for this model, including nonlinear isoscalar-scalar σ terms [31], is given by

$$\begin{aligned}\mathcal{L}=&\bar{\psi}[\gamma_\mu(i\partial^\mu-g_\omega\mathcal{V}^\mu-g_\rho\boldsymbol{\mathcal{B}}^\mu\cdot\boldsymbol{\tau})-(M-g_\sigma\phi-g_\delta\boldsymbol{\tau}\cdot\boldsymbol{\delta})]\psi\\&+\frac{1}{2}(\partial_\mu\phi\partial^\mu\phi-m_s^2\phi^2)-\frac{a}{3}\phi^3-\frac{b}{4}\phi^4-\frac{1}{4}W_{\mu\nu}W^{\mu\nu}\\&+\frac{1}{2}m_v^2\mathcal{V}_\nu\mathcal{V}^\nu+\frac{1}{2}(\partial_\mu\boldsymbol{\delta}\cdot\partial^\mu\boldsymbol{\delta}-m_\delta^2\boldsymbol{\delta}^2)-\frac{1}{4}\mathbf{G}_{\mu\nu}\cdot\mathbf{G}^{\mu\nu}\\&+\frac{1}{2}m_\rho^2\boldsymbol{\mathcal{B}}_\nu\cdot\boldsymbol{\mathcal{B}}^\nu,\end{aligned}\tag{1}$$

where $W^{\mu\nu}(x)=\partial^\mu\mathcal{V}^\nu(x)-\partial^\nu\mathcal{V}^\mu(x)$ and $\mathbf{G}^{\mu\nu}(x)=\partial^\mu\boldsymbol{\mathcal{B}}^\nu(x)-\partial^\nu\boldsymbol{\mathcal{B}}^\mu(x)$.

Here $\psi(x)$ is the nucleon fermionic field, and $\phi(x)$ and $\mathcal{V}^\nu(x)$ represent neutral scalar and vector boson fields, respectively. $\boldsymbol{\delta}(x)$ and $\boldsymbol{B}^\nu(x)$ are the charged scalar and vector fields, and $\boldsymbol{\tau}$ denotes the isospin matrices.

From the Lagrangian, Eq. (1), with the Euler procedure a set of coupled equations of motion for the meson and nucleon fields can be derived. The basic approximation in nuclear matter applications consists of neglecting all the terms containing derivatives of the meson fields, with respect to the mass contributions. Then the meson fields are simply connected to the operators of the nucleon scalar and current densities by the following equations:

$$\hat{\Phi}/f_\sigma+A\hat{\Phi}^2+B\hat{\Phi}^3=\bar{\psi}(x)\psi(x)\ \equiv\widehat{\rho_S},\tag{2}$$

$$\hat{\mathcal{V}}^\mu(x)=f_\omega\bar{\psi}(x)\gamma^\mu\psi(x)\equiv f_\omega\hat{j}_\mu,$$

$$\hat{\mathbf{B}}^\mu(x)=f_\rho\bar{\psi}(x)\gamma^\mu\boldsymbol{\tau}\psi(x),$$

$$\hat{\boldsymbol{\delta}}(x)=f_\delta\bar{\psi}(x)\boldsymbol{\tau}\psi(x),\tag{3}$$

where $\hat{\Phi}=g_\sigma\phi$, $f_\sigma=(g_\sigma/m_\sigma)^2$, $A=a/g_\sigma^3$, $B=b/g_\sigma^4$, $f_\omega=(g_\omega/m_\omega)^2$, $f_\rho=(g_\rho/2m_\rho)^2$, and $f_\delta=(g_\delta/m_\delta)^2$.

For the nucleon fields we get a Dirac-like equation. Indeed after substituting Eqs. (2) and (3) for the meson field operators, we obtain an equation which contains only nucleon field operators. All the equations can be consistently solved in a relativistic mean field (RMF) approximation, where most applications have been performed, in particular in the Hartree scheme [3,5].

The inclusion of Fock terms is conceptually important [32,33] since it automatically leads to contributions to various channels, also in the absence of explicit coupling terms. We will discuss this point later. A thorough study of the Fock contributions in a QHD approach with nonlinear self-interacting terms has been recently performed [34], in particular for asymmetric matter [20].

The present approximation implies that retardation and finite range effects in the exchange of mesons between nucleons are neglected. Nevertheless, thanks to the small Compton wavelengths of the mesons σ, ω, ρ, and δ, the assumptions expressed by Eqs. (2) and (3) are quite reasonable. For light mesons such as pions this approximation is not justified. However, in this case a perturbative expansion in the pion-nucleon coupling constant seems to be reasonable [32]. Moreover, it has been shown that the inclusion of pions does not change qualitatively the description of nuclear matter around normal conditions [32].

We remark that the kinetic approach discussed here is fully consistent with the previous approximation. We are concerned with a semiclassical description of nuclear dynamics, so that the nuclear medium is supposed to be in states for which the nucleon scalar and current densities are smooth functions of the space-time coordinates.

Within a mean field picture of the QHD model we focus our analysis on a description of the many-body nuclear system in terms of one-body dynamics. This is enough for the scope of the paper. Correlation effects can be effectively included at the level of coupling constants, as noted in the discussion of the results.

We will perform the many-body calculations in quantum phase space, introducing the Wigner transform of the one-body density matrix for the fermion field [35,36].

The one-particle Wigner function is defined as

$$[\hat{F}(x,p)]_{\alpha\beta}=\frac{1}{(2\pi)^4}\int d^4R e^{-ip\cdot R}\times\left\langle :\bar{\psi}_\beta\left(x+\frac{R}{2}\right)\psi_\alpha\left(x-\frac{R}{2}\right):\right\rangle,$$

where α and β are double indices for spin and isospin. The angular brackets denote statistical averaging and the colons denote normal ordering. The Wigner function is a matrix in spin and isospin spaces; in the case of asymmetric NM it is useful to decompose it into neutron and proton components. Following the treatment of the Fock terms in the nonlinear QHD model introduced in Refs. [34,20], we obtain for the Wigner function the following kinetic equation:

$$\frac{i}{2}\partial_\mu\gamma^\mu\hat{F}^{(i)}(x,p)+\gamma^\mu p^*_{\mu i}\hat{F}^{(i)}(x,p)-M_i^*\hat{F}^{(i)}(x,p)+\frac{i}{2}\Delta[\tilde{f}_\omega j_\mu(x)\gamma^\mu\pm\tilde{f}_\rho j_{3\mu}(x)\gamma^\mu-\tilde{f}_\sigma\rho_S(x)\mp\tilde{f}_\delta\rho_{S3}(x)]\hat{F}^{(i)}(x,p)=0,\qquad i=n,p,\tag{4}$$

where $\Delta=\partial_x\cdot\partial_p$, with ∂_x acting only on the first term of the products. Here $\rho_{S3}=\rho_{Sp}-\rho_{Sn}$ and $j_{3\mu}(x)=j_\mu^p(x)-j_\mu^n(x)$ are the isovector scalar density and the isovector baryon current, respectively. We have defined the kinetic momentum and effective masses as

$$p^*_{\mu i}=p_\mu-\tilde{f}_\omega j_\mu(x)\pm\tilde{f}_\rho j_{3\mu}(x),$$
$$M_i^*=M-\tilde{f}_\sigma\rho_S(x)\pm\tilde{f}_\delta\rho_{S3}(x),\tag{5}$$

with the effective coupling functions given by

$$\tilde{f}_\sigma=\frac{\Phi}{\rho_S}-\frac{1}{8}\frac{d\Phi(x)}{d\rho_S(x)}-\frac{1}{2\rho_S}\mathrm{Tr}\hat{F}^2(x)\frac{d^2\Phi(x)}{d\rho_S^2(x)}+\frac{1}{2}f_\omega+\frac{3}{2}f_\rho-\frac{3}{8}f_\delta,$$

$$\tilde{f}_\omega=\frac{1}{8}\frac{d\Phi(x)}{d\rho_S(x)}+\frac{5}{4}f_\omega+\frac{3}{4}f_\rho+\frac{3}{8}f_\delta,$$

$$\tilde{f}_\delta=-\frac{1}{8}\frac{d\Phi(x)}{d\rho_S(x)}+\frac{1}{2}f_\omega-\frac{1}{2}f_\rho+\frac{9}{8}f_\delta,$$

$$\tilde{f}_\rho=\frac{1}{8}\frac{d\Phi(x)}{d\rho_S(x)}+\frac{1}{4}f_\omega+\frac{3}{4}f_\rho-\frac{1}{8}f_\delta,\tag{6}$$

where $8\,\mathrm{Tr}\hat{F}^2(x)=\rho_S^2+j_\mu j^\mu+\rho_{S3}^2+j_{3\mu}j^{3\mu}$. We remind the reader that we are dealing with a transport equation so the currents and densities, in general, are varying functions of the space-time, at variance with the case of nuclear matter at equilibrium.

The expression of Eq. (5) for the effective mass embodies an isospin contribution from Fock terms also without a direct inclusion of the δ meson in the Lagrangian. The usual RMF approximation (Hartree level) is covered by the Hartree-Fock results; one has has only to change the coupling functions $\tilde{f}_i(i=\sigma,\omega,\rho,\delta)$, Eqs. (6), with the coupling constants f_i.

Equilibrium properties: The nuclear equation of state

We will focus our study on the collective modes. In order to analyze the results it is essential to relate them to the equation of state (EOS), which we will briefly discuss in the following. In particular, for the collective response in asymmetric nuclear matter the behavior of the symmetry energy E_{sym} is important.

The energy density and pressure for symmetric and asymmetric nuclear matter and the n,p effective masses can be self-consistently calculated just in terms of the four boson coupling constants $f_i \equiv (g_i^2/m_i^2)$, $i=\sigma,\omega,\rho,\delta$, and the two parameters of the σ self-interacting terms, $A \equiv a/g_\sigma^3$ and $B \equiv b/g_\sigma^4$; see Refs. [34,20].

The isoscalar meson parameters are fixed from symmetric nuclear matter properties at $T=0$: saturation density $\rho_0 = 0.16\ \mathrm{fm}^{-3}$, binding energy $E/A=-16$ MeV, nucleon effective mass $M^*=0.7M_N$ ($M_N=939$ MeV) [37], and incompressibility $K_V=240$ MeV at ρ_0. The fitted f_σ, f_ω, A, B parameters are reported in Table I. They have quite standard values for these minimal nonlinear RMF models. Set I and set II correspond to the best parameters within a nonlinear Hartree calculation, respectively, with the $\rho(\text{set I,NLH}-\rho)$ and with the $\rho+\delta$ $(\text{set II,NLH}-(\rho+\delta))$ couplings in the isovector channel (see the discussion in Ref. [16]). NLHF stands for the nonlinear Hartree-Fock scheme described before.

In the table we report also the NL3 parametrization, widely used in nuclear structure calculations [39]. We remind the reader that the NL3-saturation properties for symmetric matter are chosen as $\rho_0=0.148\ \mathrm{fm}^{-3}$, $M^*=0.6M_N$, and $K_V=271.8$ MeV. The symmetry parameter is $a_4 = 37.4$ MeV.

The symmetry energy in ANM is defined from the expansion of the energy per nucleon $E(\rho_B,\alpha)$ in terms of the asymmetry parameter α defined as

$$\alpha \equiv -\frac{\rho_{B3}}{\rho_B} = \frac{\rho_{Bn}-\rho_{Bp}}{\rho_B} = \frac{N-Z}{A}.$$

We have

$$E(\rho_B,\alpha) \equiv \frac{\epsilon(\rho_B,\alpha)}{\rho_B} = E(\rho_B) + E_{sym}(\rho_B)\alpha^2 + O(\alpha^4) + \cdots \tag{7}$$

and so, in general,

$$E_{sym} \equiv \frac{1}{2}\left.\frac{\partial^2 E(\rho_B,\alpha)}{\partial\alpha^2}\right|_{\alpha=0} = \frac{1}{2}\rho_B \left.\frac{\partial^2\epsilon}{\partial\rho_{B3}^2}\right|_{\rho_{B3}=0}. \tag{8}$$

In the Hartree case an explicit expression for the symmetry energy can be easily derived [40,16]:

$$\begin{aligned} E_{sym}(\rho_B) &= \frac{1}{6}\frac{k_F^2}{E_F^*} + \frac{1}{2} f_\rho \rho_B - \frac{1}{2} f_\delta \frac{M^{*2}\rho_B}{E_F^{*2}[1+f_\delta A(k_F,M^*)]} \\ &\equiv E_{sym}^{kin} + E_{sym}^{pot}, \end{aligned} \tag{9}$$

where k_F is the nucleon Fermi momentum corresponding to ρ_B, $E_F^* \equiv \sqrt{(k_F^2+M^{*2})}$ and M^* is the effective nucleon mass in symmetric NM, $M^*=M_N-g_\sigma\phi$.

The integral

$$A(k_F,M^*) \equiv \frac{4}{(2\pi)^3}\int d^3k \frac{k^2}{(k^2+M^{*2})^{3/2}} = 3\left(\frac{\rho_S}{M^*} - \frac{\rho_B}{E_F^*}\right). \tag{10}$$

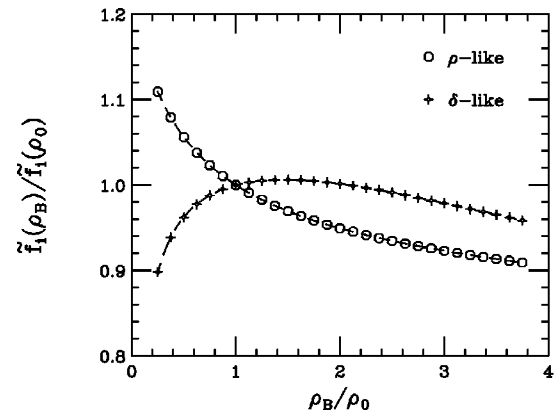

FIG. 1. Baryon density variation of the isovector effective coupling when the Fock terms are included.

We remark that $A(k_F,M^*)$ is certainly very small at low densities, and actually it can be still neglected up to a baryon density $\rho_B \simeq 3\rho_0$ (see Ref. [16]).

Then in the density range of interest here we can use, at the leading order, a much simpler form of the symmetry energy, with transparent δ-meson effects:

$$E_{sym}(\rho_B) = \frac{1}{6}\frac{k_F^2}{E_F^*} + \frac{1}{2}\left[f_\rho - f_\delta\left(\frac{M^*}{E_F^*}\right)^2\right]\rho_B. \tag{11}$$

We see that, when the δ is included, the observed a_4 value actually assigns the combination $[f_\rho - f_\delta(M/E_F)^2]$ of the (ρ,δ) coupling constants. If $f_\delta \neq 0$, we have to increase the ρ coupling (see Fig. 1 of Ref. [40]). In our calculations we use the value $a_4=32$ MeV.

In Table I, set I corresponds to $f_\delta=0$. In set II, f_δ is chosen as 2.0 fm^2, roughly derived from the analysis of Ref. [23]. As already noted in the Introduction this choice is not essential for our discussion: the aim of our work is just to show the new dynamical effects of the δ-meson coupling and to select the corresponding most sensitive observables.

In order to have the same a_4 we must increase the ρ-coupling constant of a factor of 3, up to $f_\rho=2.8$ fm^2. Now the symmetry energy at saturation density is actually built from the balance of scalar (attractive) and vector (repulsive) contributions, with the scalar channel becoming weaker with increasing baryon density [16]. This is indeed the isovector counterpart of the saturation mechanism occurring in the isoscalar channel for the symmetric nuclear matter. From such a scheme we get further support for the introduction of the δ coupling in the symmetry energy evaluation.

When the δ "channel" is included the behavior of the symmetry energy is stiffer at high baryon density for the relativistic mechanism discussed before. When the Fock terms are evaluated the new "effective" couplings, Eqs. (6), naturally acquire a density dependence. This is shown in Fig. 1 for the isovector terms. The decrease of the "effective" ρ coupling at high density accounts for a slight softening of the symmetry energy. Details of the calculation can be found in Refs. [20,16].

III. LINEAR RESPONSE EQUATIONS

In this section we study collective oscillations that propagate in cold nuclear matter due to the mean field dynamics. In some sense we follow a relativistic extension of the method introduced by Landau to study liquid ^{3}He [41–43] and recently applied to investigate stable and unstable modes in nuclear matter [30,44,45]. The starting point is the kinetic transport equation (4). We look for solutions corresponding to small oscillations of $\hat{F}(x,p)$ around the equilibrium value. Therefore we put

$$\hat{F}(x,p)=\hat{H}(p)+\hat{G}(x,p), \tag{12}$$

where $\hat{H}(p)$ is the Wigner function at equilibrium (see the Appendix) and $\hat{G}(x,p)$ represent its fluctuations.

For the equilibrium state the Wigner function contains only the isoscalar term and the third component of the isovector term. We limit ourselves to studying excited states in these channels. Therefore we consider isovector density fluctuations with $m_T=0$ only; i.e., we do not study processes where a neutron converts into a proton or vice versa. In a linear response scheme oscillations in the aforementioned channels are decoupled from the remaining ones ($m_T=\pm 1$).

In the linear approximation, i.e., neglecting terms of second order in $\hat{G}(x,p)$, the equations for the Wigner functions become

$$\frac{i}{2}\partial_\mu \gamma^\mu \hat{G}_{(1)}(x,p)+(\Pi_\mu-\tilde{f}_\rho b_\mu)\gamma^\mu \hat{G}_{(1)}(x,p) - M_1^* \hat{G}_{(1)}(x,p)=\left(1-\frac{i}{2}\Delta\right)[\hat{\mathcal{F}}(x)+\hat{\mathcal{F}}_3(x)]\hat{H}_{(1)}(p), \tag{13}$$

for protons, and

$$\frac{i}{2}\partial_\mu \gamma^\mu \hat{G}_{(2)}(x,p)+(\Pi_\mu+\tilde{f}_\rho b_\mu)\gamma^\mu \hat{G}_{(2)}(x,p) - M_2^* \hat{G}_{(2)}(x,p)=\left(1-\frac{i}{2}\Delta\right)[\hat{\mathcal{F}}(x)-\hat{\mathcal{F}}_3(x)]\hat{H}_{(2)}(p), \tag{14}$$

for neutrons, where $M_1^*=M-\tilde{f}_\sigma\rho_S-\tilde{f}_\delta\rho_{S3}$ and $M_2^*=M-\tilde{f}_\sigma\rho_S+\tilde{f}_\delta\rho_{S3}$. The quantities $\hat{\mathcal{F}}(x)$ and $\hat{\mathcal{F}}_3(x)$ are the isoscalar and isovector components of the self-consistent field:

$$\hat{\mathcal{F}}(x)=-8\tilde{f}_\sigma G(x)+8\tilde{f}_\omega\gamma_\mu G^\mu(x)-8\frac{\partial\tilde{f}_\sigma}{\partial\rho_S}\rho_S G(x) -8\frac{\partial\tilde{f}_\sigma}{\partial j_\mu}\rho_S G^\mu(x)-8\frac{\partial\tilde{f}_\sigma}{\partial\rho_{S3}}\rho_S G_3(x)-8\frac{\partial\tilde{f}_\sigma}{\partial j_{3\mu}}\rho_S G_3^\mu(x) +8\frac{\partial\tilde{f}_\omega}{\partial\rho_S}\gamma_\mu j^\mu G(x), \tag{15}$$

$$\hat{\mathcal{F}}_3(x)=-8\tilde{f}_\delta G_3(x)+8\tilde{f}_\rho\gamma_\mu G_3^\mu(x)-8\frac{\partial\tilde{f}_\delta}{\partial\rho_S}\ \rho_{S3}\ \ G(x) +8\frac{\partial\tilde{f}_\rho}{\partial\rho_S}\gamma_\mu j_3^\mu G(x). \tag{16}$$

The Hartree approximation is recovered by vanishing all the derivatives inside the quantities $\hat{\mathcal{F}}(x)$ and $\hat{\mathcal{F}}_3(x)$, except $\partial\tilde{f}_S/\partial\rho_S$, since still $\tilde{f}_\sigma=\Phi(\rho_S)/\rho_S$.

In order to obtain the equations for the collective oscillations we multiply Eqs. (13) and (14) by γ_λ. After performing the traces, we equate to zero both the real and imaginary parts of the result [10,11]. Furthermore, by Fourier transforming and integrating over the four-momentum, we get the set of equations for the scalar and vector fluctuations of each species ($i=1,2$ for proton, neutron, respectively):

$$\sum_{j=1}^{2}\left\{\left[\delta_{i,j}+\left(\frac{\rho_{Si}}{M_i^*}-4\,C^{(i)}(k)\right)D_{ij}^S+4\,C^{\mu(i)}(k)B_{\mu ij}^V\right]G_{(j)}(k) +\left[4\,C_\mu^{(i)}(k)D_{ij}^V+\left(\frac{\rho_{Si}}{M_i^*}-4\,C^{(i)}(k)\right)B_{\mu ij}^S\right]G_{(j)}^\mu(k)\right\}=0, \tag{17}$$

$$\sum_{j=1}^{2}\left\{\left[\delta_{i,j}g_\mu^\lambda+\left(\frac{\rho_{Si}}{M_i^*}g_\mu^\lambda+4\frac{C_\mu^{(i)\lambda}(k)}{M_i^*}\right)D_{ij}^V -4\,C^{(i)\lambda}(k)B_{\mu ij}^S\right]G_{(j)}^\mu(k)-\left[4\,C^{(i)\lambda}(k)D_{ij}^S -\left(\frac{\rho_{Si}}{M_i^*}g_\mu^\lambda+4\frac{C_\mu^{(i)\lambda}(k)}{M_i^*}\right)B_{\mu ij}^V\right]G_{(j)}(k)\right\}=0, \tag{18}$$

with

$$D_{ij}^S=\tilde{f}_\sigma+(-1)^{i+j}\tilde{f}_\delta+\frac{\partial\tilde{f}_\sigma}{\partial\rho_S}\rho_S-(-1)^j\frac{\partial\tilde{f}_\sigma}{\partial\rho_3}\rho_S -(-1)^i\frac{\partial\tilde{f}_\delta}{\partial\rho_S}\rho_3,$$

$$D_{ij}^V=\tilde{f}_\omega+(-1)^{i+j}\tilde{f}_\rho+\frac{\partial\tilde{f}_\omega}{\partial\rho_S}\rho_S-(-1)^i\frac{\partial\tilde{f}_\rho}{\partial\rho_S}\rho_3,$$

$$B_{\mu ij}^V=\frac{\partial\tilde{f}_\omega}{\partial\rho_S}j_\mu-(-1)^i\frac{\partial\tilde{f}_\rho}{\partial\rho_S}j_{3\mu},$$

$$B_{\mu ij}^S=\frac{\partial\tilde{f}_\sigma}{\partial j^\mu}\rho_S-(-1)^j\frac{\partial\tilde{f}_\delta}{\partial j_3^\mu}\rho_S.$$

The explicit expressions of the coefficients $C^{(i)}(k)$, $C_\lambda^{(i)}(k)$, and $C_{\lambda\mu}^{(i)}(k)$, together with some of their properties, can be found in the Appendix.

We notice, however, that the formalism developed in the Hartree-Fock approximation allows us to achieve the set of equations valid also for an approach to QHD based on a dependence on the scalar density of all couplings [46]. In particular in our case the coupling functions to the scalar-isoscalar channel depend on all the isoscalar and isovector densities and currents (for details, see [34]); then, it is even more general.

The set of equations developed in the Hartree-Fock approximation recover the ones corresponding to the usual Hartree approximation (RMF). As already mentioned, it is easily obtained by considering the coupling $\tilde{f}_i$ to each channel equal to the coupling constant of the corresponding meson. The result is an appreciably plainer structure of the equations due to the constant value of all couplings, except $\tilde{f}_\sigma$ [47]. Equations (17) and (18) can be reduced to

$$\sum_{j=1}^{2}\left\{\left[\delta_{i,j}+\left(\frac{\rho_{Si}}{M_i^*}-4\,C^{(i)}(k)\right)D_{ij}^S\right]G_{(j)}(k)+[4\,C_\mu^{(i)}(k)D_{ij}^V]G_{(j)}^\mu(k)\right\}=0, \tag{19}$$

$$\sum_{j=1}^{2}\left\{\left[\delta_{i,j}g_\mu^\lambda+\left(\frac{\rho_{Si}}{M_i^*}g_\mu^\lambda+4\,\frac{C_\mu^{(i)\lambda}(k)}{M_i^*}\right)D_{ij}^V\right]G_{(j)}^\mu(k)-[4\,C^{(i)\lambda}(k)D_{ij}^S]G_{(j)}(k)\right\}=0, \tag{20}$$

with $D_{ij}^S=d\Phi/d\rho_S+(-1)^{i+j}f_\delta$, $D_{ij}^V=f_\omega+(-1)^{i+j}f_\rho$.

The physics effects appear more transparent and we will follow the Hartree scheme in the next sections, keeping well in mind that the Fock contributions can be easily included. We expect to have some extra contributions in the various interaction channels without qualitative modifications of the physical response.

The normal collective modes are plane waves, characterized by the wave vector $[k^\mu=(k^0,0,0,|\mathbf{k}|)]$; they are determined by solving the set of homogeneous equations (19) and (20). The solutions correspond only to longitudinal waves and do not depend on k^0 and $|\mathbf{k}|$ separately, but only on the ratio

$$v_s=\frac{k^0}{|\mathbf{k}|}.$$

The sound velocities are given by values of v_s for which the relevant determinant of the set [Eqs. (17) and (18)] vanishes, i.e., the dispersion relations. In correspondence the neutron-proton structure of the eingenvectors (normal modes) can be derived. It should be remarked that in asymmetric nuclear matter isoscalar and isovector components are mixed in the normal modes. Here this can be argued by the fact that in each of the equations (19) and (20) both proton-neutron densities and currents appear.

However, we remind the reader that one can still identify isovectorlike excitations as the modes where neutrons and protons move out of phase, while isoscalarlike modes are characterized by neutrons and protons moving in phase [30,48].

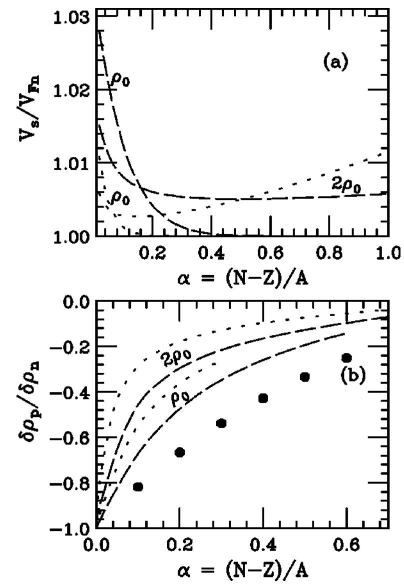

FIG. 2. Isovectorlike modes: (a) Ratio of zero-sound velocities to the neutron Fermi velocity V_{Fn} as a function of the asymmetry parameter α for two values of baryon density. Long-dashed line: NLH$-\rho$. Dotted line: NLH$-(\delta+\rho)$. (b) Corresponding ratios of proton and neutron amplitudes. All lines are labeled with the baryon density $\rho_0=0.16$ fm^{-3}. The solid circles in panel (b) represent the trivial behavior of $-(\rho_p/\rho_n)$ vs α.

IV. ROLE OF SCALAR-VECTOR FIELDS IN THE DYNAMICAL RESPONSE

Before showing numerical results for the dynamical response of asymmetric nuclear matter, in various baryon density regions and using the different effective interactions, we would like to analyze in detail the structure of the relativistic linear response theory in order to clearly pin down the role of each meson coupling.

A. Isovector response

One may expect that once a_4 is fixed, the velocity of sound is also fixed [27]. On the other hand, our results clearly indicate a different dynamical response with or without the δ-meson channel, for interactions which give *exactly the same* a_4 *parameter* (see Figs. 2 and 3) in the following. In order to get a clear understanding of this effect we will consider the case of symmetric nuclear matter in the Hartree scheme, where the dispersion relations are assuming a transparent analytical form.

For symmetric NM the densities, the effective masses, and the coefficients $C^{(i)}(k)$, $C_\lambda^{(i)}(k)$, and $C_{\lambda\mu}^{(i)}(k)$ are equal for protons and neutrons. Now it is also possible to decouple the collective modes into *pure* isoscalar and isovector oscillations [30]. After a straightforward rearrangement of Eqs. (19) and (20), we have, for the isovector modes,

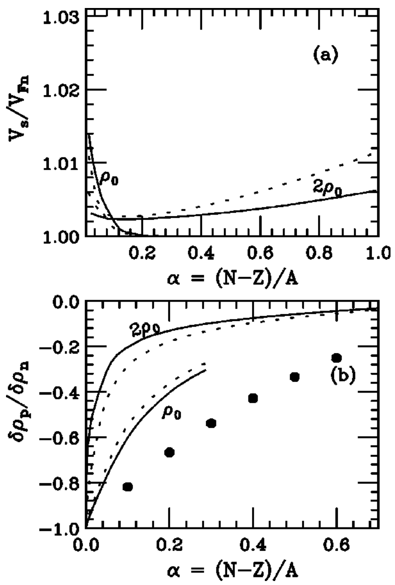

FIG. 3. The same as Fig. 2 but for NLHF (solid lines) and NLH$-(\delta+\rho)$ (dotted lines).

$$\delta\rho_3+\left[\frac{\rho_S}{M^*}-8\frac{C_{00}(k)}{M^*}(1-v_s^2)\right]f_\rho\,\delta\rho_3-8\,f_\delta C_0(k)\,\delta\rho_{S3}=0,$$

$$\delta\rho_{S3}+\left[\frac{\rho_S}{M^*}-8\,C(k)\right]f_\delta\,\delta\rho_{S3}+8\,f_\rho C_0(k)(1-v_s^2)\,\delta\rho_3=0. \tag{21}$$

We stress that the structure is the same for the isoscalar excitations; of course, one has to change the isovector fluctuations with the isoscalar ones ($\delta\rho_B$, $\delta\rho_S$) and the coupling constants of isovector mesons with those of the isoscalar mesons [47].

Note that in this case to find the zero-sound velocity one has to evaluate determinant of a 2×2 matrix (and not a 4×4); hence the condition for having a solution can be written as

$$1+N_F\left[f_\rho(1-v_s^2)-f_\delta\frac{M^{*2}}{E_F^{*2}}\left(1-f_\delta A(k_F,M^*)-f_\rho\frac{\rho_S}{M^*}v_s^2\right)\right]\varphi(s)=0. \tag{22}$$

Here $N_F=2K_FE_F^*/\pi^2$ is the density of states at the Fermi surface and $s\equiv v_s/v_F$. To get Eq. (22) we have used the expression for $C(k)$, $C_0(k)$, and $C_{00}(k)$ in terms of the Lindhard function $\varphi(s)$ (see the Appendix). The quantity $A(K_F,M^*)$ is the same integral discussed in Eqs. (9)–(11).

At this point we can make the approximation

$$v_s^2\simeq v_F^2=\frac{k_F^2}{E_F^{*2}}$$

to evaluate the expression inside the square brackets. Looking at Figs. 2, 3, and 6 this is a good approximation within 3%. Equation (22) assumes a quite clear form

$$1+\frac{6\,E_F^*}{k_F^2}\left[E_{sym}^{pot}-\frac{f_\rho}{2}\frac{k_F^2}{E_F^{*2}}\left(1-f_\delta\frac{M^*}{E_F^{*2}}\rho_S\right)\rho_B\right]\varphi(s)=0\,, \tag{23}$$

where the potential part of the symmetry energy explicitly appears in the dispersion relations, but *joined to an important correction term* which shows a different f_ρ, f_δ structure with respect to that of E_{sym}^{pot}, Eqs. (9) and (11). *We can easily have interactions with the same a_4 value at normal density but with very different isovector response*. E.g., when we include the δ channel we know that we have to increase the f_ρ coupling in order to have the same a_4 [see the discussion of Eqs. (9) and (11)], but now the "restoring force" [coefficient of the Lindhard function in Eq. (23)] will be strongly reduced.

Equation (23) suggests to define an effective symmetry energy like

$$E_{sym}^*=E_{sym}^{pot}-\frac{f_\rho}{2}\frac{k_F^2}{E_F^{*2}}\left(1-f_\delta\rho_S\frac{M^*}{E_F^{*2}}\right)\rho_B\,, \tag{24}$$

which acts as a restoring force for the isovector mode. We can see that once the symmetry energy is fixed its effect on the dynamical response depends on the strength of each isovector field. In particular we can easily verify that the f_δ factor inside the brackets in Eq. (24) is a second-order correction and the "leading contribution" to the reduction $\Delta E_{sym}^*=E_{sym}^{pot}-E_{sym}^*$ is essentially given by the coupling of the ρ-meson field. On the other hand, we know [20] that once the symmetry energy at saturation density a_4 is fixed, the change of f_ρ is only due to the strength of f_δ. We have seen from Table I that f_ρ can go from 1 fm^2, if we switch off the δ channel, to $f_\rho=2.8$ fm^2. In terms of the effective symmetry energy this means (if we consider the dynamical response at ρ_0), $\Delta E_{sym}^*\sim4$ MeV if $f_\delta\sim2.0$ fm^2. This "softening" of the restoring force easily accounts for the decrease of the sound velocity (v_s/v_{Fn} seen in Fig. 2) for symmetric nuclear matter, $\alpha=0$, when we pass from NLH$-\rho$ and NLH$-(\rho+\delta)$.

We like to note that a similar effect has been pointed out from a detailed nonrelativistic Skyrme-RPA study of the giant dipole resonance in heavy nuclei (^{208}Pb) using effective interactions with various isovector terms [49]. A different sensitivity of the average resonance frequencies on the symmetry energy a_4 and on its slope has been found. In a covariant scheme we can see from Eq. (24) that such behavior can be achieved only using two isovector fields, at the lowest order. Another interest of this result is that a dynamical observable can be more sensitive to the microscopic structure of the isovector interaction. For instance, in a careful study of the neutron distributions (see Ref. [24] already quoted in the Introduction), it is clearly shown that these observables are almost equally correlated to the value, slope, and curvature of the symmetry term.

B. Isoscalar response

As already remarked, we like to note that for symmetric NM there is a tight analogy between the isoscalar and isovector responses in the RMF approach. In the isoscalar degree of freedom the compressibility will play the same role as the symmetry energy in the dispersion relation equations. Also in this case we will have an important correction term coming from the interplay of the scalar and vector channels.

Equation (22) now becomes [47]

$$1+N_F\left[f_\omega(1-v_s^2)-f_\sigma\frac{M^{*2}}{E_F^{*2}}\left(1-f_\sigma A(k_F,M^*)-f_\omega\frac{\rho_S}{M^*}v_s^2\right)\right]\varphi(s)=0, \tag{25}$$

which can be reduced to the isoscalar equivalent of Eq. (23):

$$1+\frac{E_F^*}{3\,k_F^2}\left[K_{NM}^{pot}-9\,f_\omega\frac{k_F^2}{E_F^{*2}}\left(1-f_\sigma\frac{M^*}{E_F^{*2}}\rho_S\right)\rho_B\right]\varphi(s)=0, \tag{26}$$

where the K_{NM}^{pot} is the potential part of the nuclear matter compressibility, which in the Hartree scheme has the simple structure [47] [see also Eq. (16) of Ref. [27]]

$$K_{NM}(\rho_B)=\frac{3\,k_F^2}{E_F^*}+9\left[f_\omega-f_\sigma\left(\frac{M^*}{E_F^*}\right)^2\right]\rho_B\equiv K_{NM}^{kin}+K_{NM}^{pot}. \tag{27}$$

By means of such an analogy, the previous discussion can be extended to isoscalar oscillations with the role of E_{sym} now "played" by the compressibility. In this case, however, one always takes into account both the scalar and vector channels in any RMF models. However, the coupling constant f_ω can assume very different values depending on the required value for effective masses M_0^*. This is easy to understand since in the RMF limit the saturation binding energy has the simple form

$$E/A(0)=E_F^*+f_\omega\rho_B(0)-M_N,$$

where M_N is the bare nucleon mass. So we see that in order to have the same saturation values of $\rho_B(0)$, $E/A(0)$ when we decrease M_0^*, we have to increase f_ω. We derive then the natural conclusion that if two EOS have different effective masses, even if the compressibility is equal, the dynamical behavior is expected to be different. This is a very general feature present also in nonrelativistic approaches.

From studies of monopole resonances in finite nuclei with the RMF model it seems that a higher value of compressibility is required respect to nonrelativistic calculations. Many authors state that this certainly demands for a clarification [13,14]. Even if the monopole resonance is not directly connected to the isoscalar collective mode in nuclear matter, our discussion nicely suggests to look at the interplay between the effective mass and compressibility. For example, we can estimate by means of Eq. (26) that we can have a shift between the compressibility and the "effective compressibility" of the order of $\sim$100 MeV among different parametrizations with the same K_{NM}. Therefore model with $K\sim 300A$ MeV can reproduce the same frequencies of other models with $K\sim 200A$ MeV (and a slightly larger M_0^*).

C. Landau parameters

We would like to briefly discuss the relativistic equations for collective modes in terms of the Landau parameters. Interesting features will appear from the comparison to the nonrelativistic analogous case. We will focus first on the isovector response, but as already clearly shown before, the structure of the results will be absolutely similar in the isoscalar channel.

The general nonrelativistic expression for the isovector modes can be found in Ref. [42]:

$$1+\left[F_0^a+\frac{F_1^a}{1+1/3\,F_1^a}s^2\right]\varphi(s)=0, \tag{28}$$

where F_0^a is the "isovector" combination of the Landau F_0 parameters for neutrons and protons $F_0^a=F_a^{nn}-F_0^{np}$, which can be expressed in terms of density variations of the chemical potentials:

$$F_0^{qq'}\equiv\frac{\partial\mu_q}{\partial\rho_{q'}}N_q-\delta_{qq'},\quad N_q\equiv\frac{k_{Fq}E_{Fq}^*}{\pi^2},\quad q=n,p. \tag{29}$$

F_1^a are the equivalent for the momentum-dependent part of the mean field. In the relativistic approach, for symmetric nuclear matter, we get

$$F_0^a=F_\rho-F_\delta\frac{M^{*2}}{E_F^{*2}}\frac{1}{1+f_\delta A(k_F,M^*)},$$

$$F_1^a=-F_\rho\frac{v_F^2}{1+\dfrac{1}{3}F_\rho v_F^2}, \tag{30}$$

where $F_i=N_Ff_i(i=\rho,\delta)$ with $N_F=2N_{n,p}$. Note that the F_1^a contribution comes only from the vector coupling. By using the expression E_{sym}^{pot}, Eq. (9), we can write Eq. (28) in the same form as Eqs. (22) and (23). The result is a similar expression but with the lack of the term in f_δ inside the brackets in Eq. (23). As said, this is not the leading term; however, around saturation density it amounts to about 10% of the total correction.

Moreover, turning to the analogy with isoscalar channel the coupling of the σ field is now much larger and this purely relativistic contribution could be up to 20%. We underline this point because generally the linear response in the RMF model is discussed calculating the Landau parameters and then using these estimations directly in the nonrelativistic expression for collective modes [27,29].

In conclusion from the analysis in terms of the Landau parameters, we can describe the effect of the scalar-vector coupling competition previously discussed in the following way. The symmetry energy fixes the F_0^a, in fact

$$E_{sym}=\frac{k_F^2}{6\,E_F^*}(1+F_0^a), \tag{31}$$

but in the dynamical response the F_1^a also enters, linked to the momentum dependence of the mean field, mostly given by the vector meson coupling. The results are completely analogous in the isoscalar channel, with the compressibility given by

$$K_{NM}=\frac{3\,k_F^2}{E_F^*}(1+F_0^s), \tag{32}$$

with the ''isoscalar'' combination $F_0^s=F_0^{nn}+F_0^{np}$. The relativistic forms of the isoscalar Landau parameters are exactly the same as in Eq. (30), just substituting the δ,ρ coupling constants with the σ,ω ones [47].

V. ISOVECTOR COLLECTIVE MODES IN ASYMMETRIC NUCLEAR MATTER

In this section we discuss results for the isovector collective oscillations which are driven by the symmetry energy terms of the nuclear EOS. The aim is mainly to investigate the effect of the scalar-isovector channel. This is normally not included in studying the isovector modes and in general the properties of symmetric matter in a relativistic approach, while it should be naturally present on the basis of the analysis shown in the previous section (and in Ref. [16] for equilibrium properties). Moreover, we stress again that Hartree-Fock scheme embodies in any case the presence of a scalar-isovector channel, *even without the inclusion of the δ-meson field* [20].

We will first show results obtained in the Hartree scheme (NLH) including either both the isovector ρ and δ mesons or only the ρ meson. Even if the Hartree approximation has a simpler structure, it contains all the physical effects we want to point out. Finally from the complete Hartree-Fock (NLHF) calculations we will confirm the dynamical contribution of the scalar isovector channel.

For NLH calculations we use the parametrizations of Table I, set I (ρ) and set II ($\rho+\delta$). In the Hartree-Fock case the coupling constant f_δ is adjusted to the value $\tilde{f}_\delta(\rho_0)=2.0\ \mathrm{fm}^2$ of the NLHF model, Eqs. (6).

A. Hartree results

Let us start by considering isovectorlike excitations. In Fig. 2(a) we show the sound velocities in the Hartree approximation, as a function of the asymmetry parameter α for different baryon densities. We actually plot the sound velocities in units of the neutron Fermi velocities. This is physically convenient: when the ratio approaches 1.00 we can expect that this ''zero'' sound will not propagate due to the strong coupling to the ''chaotic'' single-particle motions (''Landau damping''). This quantity then will also directly give a measure of the ''robustness'' of the collective mode we are considering.

Dotted lines refer to calculations including $(\rho+\delta)$ mesons; long-dashed lines correspond to the case with only the ρ meson. Calculations are performed at $\rho_B=\rho_0$ and $\rho_B=2\rho_0$. We stress that the results of the two calculations differ already at zero asymmetry, $\alpha=0$. At normal density (ρ_0 curves), in spite of the fact that the symmetry energy coefficient $a_4=E_{sym}(\rho_0)$ is exactly the same in the two cases, significant differences are observed in the response of the system. From Fig. 2(a) we can expect a reduction of the frequency for the bulk isovector dipole mode in stable nuclei when the scalar isovector channel (δ *like*) is present. Moreover, we note that, in the NLH$-\rho$ case, the excitation of isovector modes persists up to higher asymmetries at saturation density.

These are nontrivial features, related to the different way scalar and vector fields enter in the dynamical response of the nuclear system. Such behaviors are therefore present in both collective responses: isoscalar and isovector. We have devoted the whole previous section (Sec. IV) to a complete discussion of this effect.

Differences are observed even at $\rho_B=2\rho_0$, where, however, also the symmetry energy is different. A larger E_{sym} is obtained in the case including the δ meson, and this leads to a compensation of the effect observed at normal nuclear density. In particular, at higher asymmetries α the collective excitation becomes more robust for NLH$-(\rho+\delta)$. Differences are observed also in the ''chemical'' structure of the mode, represented by the ratio $\delta\rho_p/\delta\rho_n$, plotted in Fig. 2(b). The ratio of the out-of-phase n,p oscillations does not follow the ratio of the n,p densities for a fixed asymmetry, given by the solid circles in the figure. We systematically see a larger amplitude of the neutron oscillations. The effect is more pronounced when the δ (scalar-isovector) channel is present (dotted lines).

B. Hartree-Fock results

We have also performed the calculation in the more general case of the Hartree-Fock approximation (NLHF), whose formalism has been presented before, Eqs. (17) and (18). We have fitted the same properties of symmetric NM at the saturation density as for the Hartree case (NLH). In particular at ρ_0 the value of the isovector coupling is fixed in order get the same symmetry energy (the a_4 parameter) of the NLH$-(\rho+\delta)$ case.

In Fig. 3 we can see that quite similar results are obtained in Hartree-Fock calculations, with respect to the Hartree results including ρ and δ mesons, especially at the normal density. This can be understood by considering that in Hartree-Fock calculation the effective density-dependent couplings associated with the isovector channels are tuned in such a way to roughly reproduce, at normal density, the values of the coupling constants f_ρ and f_δ of the Hartree scheme: then not only is a_4 the same, but also its internal structure. Since such a tuning can be done only at a given density value, some differences are observed at $\rho_B=2\rho_0$,

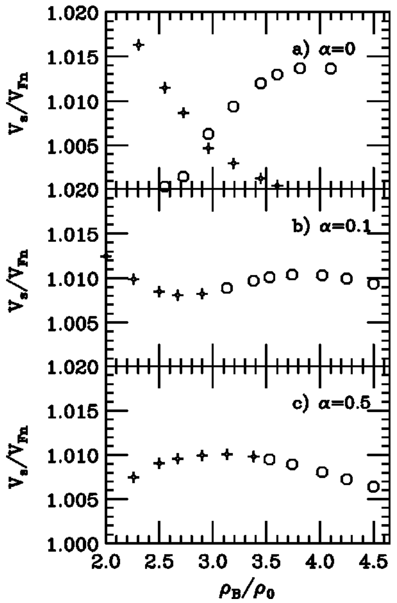

FIG. 4. Sound phase velocities of the propagating collective mode vs the baryon density (NLH+ρ case). Crosses: isovectorlike. Open circles: isoscalarlike. (a) Symmetric matter. (b) Asymmetric matter, $\alpha=0.1$. (c) Asymmetric matter, $\alpha=0.5$.

due to the density dependence of the effective coupling constants of the NLHF scheme; see Fig. 1. In particular a slightly smaller value of the sound velocity is expected at higher baryon densities.

C. Disappearance of the isovector modes

For asymmetric matter we have found that, in all the calculation schemes, with increasing baryon density the isovector modes disappear: we call such densities ρ_B^{cross}. E.g., from Figs. 2(b) and 3(b) we see that the ratio $\delta\rho_p/\delta\rho_n$ tends very quickly to zero with increasing baryon density, almost for all asymmetries. Around this transition density we expect to have an almost *pure neutron wave* propagation of the sound. Here we show the results of the NLH+ρ case (see Figs. 4 and 5), but the effect is clearly present in all the models.

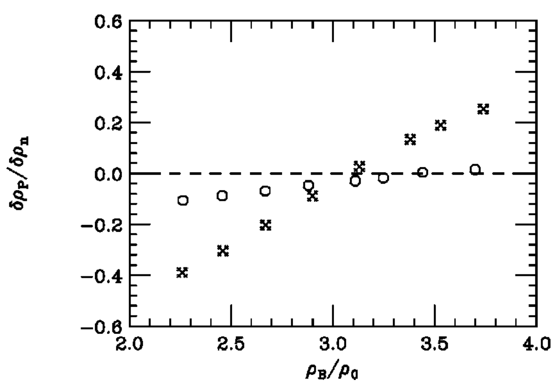

FIG. 5. Ratio of protons and neutron amplitudes in the propagating mode, for different asymmetries, as a function of the baryon density around the ρ_B^{cross}. Crosses: $\alpha=0.1$, Fig. 4(b). Open circles: $\alpha=0.5$, Fig. 4(c).

For symmetric matter we have a real crossing of the two phase velocities, isoscalar and isovector, as shown in Fig. 4(a). Above ρ_B^{cross} the isoscalar mode is the most robust.

For asymmetric matter we observe a transition in the structure of the propagating normal mode, from isovectorlike to isoscalarlike, Figs. 4(b) and 4(c). Similar effects have been seen in a nonrelativistic picture [30].

For a given asymmetry α the value of ρ_B^{cross} is different for the three models considered, as can be argued by the behavior of $\delta\rho_p/\delta\rho_n$ at $2\rho_0$ in Figs. 2(b) and 3(b). E.g., for $\alpha=0.1$, NLHF has the lower value ($\rho_B^{cross}\simeq 2.4\rho_0$), while NLH$-\rho$ has the higher one ($\rho_B^{cross}\simeq 3.0\rho_0$). This is again related to the reduction of the isovector restoring force when the scalar-isovector channel (δ like) is present; see Sec. IV.

From Fig. 5 we see that the proton component of the propagating sound is quite small in a relatively wide region around the ‘‘transition’’ baryon density, a feature becoming more relevant with increasing asymmetry; see the open circle line. This is quite interesting since it could open the possibility of an experimental observation of the *neutron wave* effect.

VI. ISOSCALAR COLLECTIVE MODES IN ASYMMETRIC NUCLEAR MATTER

So far we have focused our discussion on the isovectorlike response of the asymmetric nuclear matter. However, it is well known that isoscalarlike modes can exist also in asymmetric nuclear matter see [30,48], and references therein.

A. Exotic high baryon density modes

From the previous analysis we have seen the isoscalarlike excitations to become dominant at high baryon density, above the ρ_B^{cross} introduced before.

Some results are shown in Fig. 6. It should be noticed that the frequency of the isoscalarlike modes is essentially related to the compressibility of the system at the considered density. In Fig. 6(a) we display the sound velocity obtained in Hartree and Hartree-Fock calculations at $\rho_B=3.5\rho_0$, as a function of the asymmetry α. The differences observed among calculations performed within the Hartree or Hartree-Fock scheme are due to a different behavior of the associated equation of state at high density.

At $\alpha=0$ the two Hartree models have exactly the same isoscalar mean fields, but for asymmetric nuclear matter the different behavior of the symmetry energy leads to a different compressibility. The case NLH$-(\rho+\delta)$ which has the stiffer E_{sym} (resulting in a greater incompressibility for $\alpha>0$) with respect to NLH$-(\rho)$ shows also a greater increase of v_s/v_{Fn} with density. Instead, NLHF, even if it has the same compressibility K_{NM} at saturation density, shows a different v_s. This should be due to the density dependence of the coupling function arising from exchange terms, which leads to different values of K_{NM} out of ρ_0 (even for $\alpha=0$).

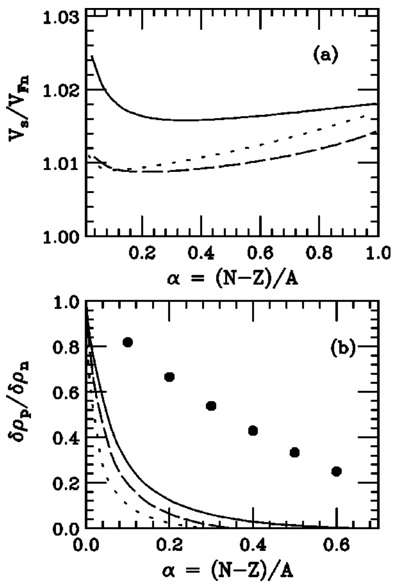

FIG. 6. The same of Fig. 2 for isoscalarlike modes, at $\rho_B = 3.5\rho_0$. Solid line: NLHF. Long-dashed line: NLH$-\rho$. Dotted line: NLH$-(\rho+\delta)$. The solid circles in panel (b) represent the behavior of ρ_p/ρ_n vs α.

Some differences are observed also in the chemical composition of the mode [Fig. 6(b)]. The black spots show the behavior of ρ_p/ρ_n vs α. Note the *pure neutron wave* structure of the propagating sound, since the oscillations of protons appear strongly damped ($\delta\rho_p/\delta\rho_n \ll \rho_p/\rho_n$); unfortunately this is an effect not experimentally accessible (at present) (see also the discussion at the end of the previous section).

Before closing this discussion we have to remark that the isoscalarlike modes at high baryon density are vanishing if the nuclear EOS becomes softer. This is indeed the results of two recent models, Refs. [29,50], where the nuclear compressibility is decreasing at high baryon density for a reduction of the isoscalar vector channel contribution. In [29] this is due to self-interacting high order terms for the ω meson, while in [50] it is due to a reduced f_ω coupling with increasing baryon density.

Finally we note that all causality violation problems (superluminal sound velocities) observed in the nonrelativistic results at high baryon density (see [27] and Fig. 3c in Ref. [30]) are completely absent in the relativistic approach (see the high density trends in Fig. 4).

B. Isospin distillation in dilute matter

We have also investigated the response of the system in the region of spinodal instability associated with the liquid-gas phase transition, which occurs at low densities. It is known that in this region an isoscalar unstable mode can be found, with imaginary sound velocity, which gives rise to an exponential growth of the fluctuations. The latter can represent a dynamical mechanism for the multifragmentation process observed in heavy-ion collisions. We have found this kind of solution in the present approach. In Fig. 7 we show the ratio $\delta\rho_p/\delta\rho_n$ as function of the initial asymmetry for

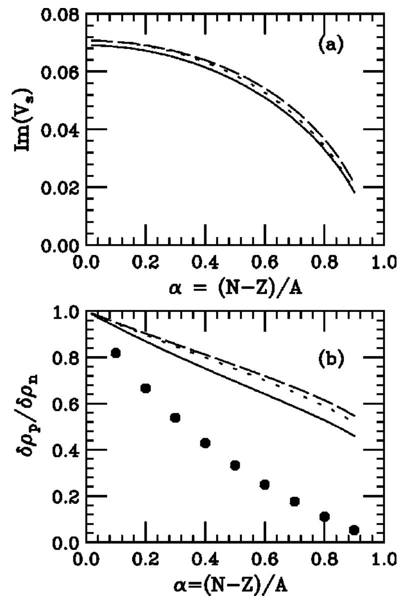

FIG. 7. Isoscalarlike unstable modes at $\rho_B = 0.4\rho_0$: Imaginary sound velocity (a), in c units, and ratio of proton and neutron amplitudes (b) as a function of the asymmetry α. Solid line: NLHF. Dotted line: NLH$-(\rho+\delta)$. Long-dashed line: NLH$-\rho$. The solid circles in panel (b) represent the behavior of $\delta\rho_p/\delta\rho_n$ vs α.

such a collective mode. For all the interactions this ratio is different from the corresponding ρ_p/ρ_n of the initial asymmetry α. This is exactly the chemical effect associated with the new instabilities in dilute asymmetric matter [7,48].

In particular it is found that, when isoscalarlike modes become unstable, the ratio $\delta\rho_p/\delta\rho_n$ becomes *larger* that the ratio ρ_p/ρ_n (at variance with the stable modes at high densities; see Fig. 6). Hence proton oscillations are relatively larger than neutron oscillations, leading to a more symmetric liquid phase and to a more neutron-rich gas phase, during the disassembly of the system. This is the so-called isospin distillation effect in fragmentation, and signatures of this effect could be searched by looking at the ratio N/Z of fragments produced in dissipative heavy-ion collisions [51,52].

We note here that in dilute asymmetric NM we can distinguish two regions of instability: mechanical (cluster formation) and chemical (component separation). There is, however, no discontinuity in the structure of the unstable modes which are developing. For all realistic effective nuclear interactions (relativistic and nonrelativistic) the nature of the unstable normal modes at low densities is always *isoscalarlike*, i.e., with neutrons and protons oscillating in phase, although with a distillation effect discussed before (see Ref. [48] for a fully detailed study of this important property of asymmetric nuclear matter).

In Fig. 7 we observe that Hartree results (with and without the δ meson) are very similar and, indeed, at low density the symmetry energy behavior is nearly the same in the two cases. On the other hand, differences are observed with respect to the Hartree-Fock case. In fact, the NLHF symmetry energy presents a softer behavior (around $\rho_B = 0.4\rho_0$), which leads to a smaller distillation effect. We remark that the equality between the NLH symmetrics with and without δ is in agreement with the analysis in terms of the generalized Landau parameters associated with normal modes developed in Ref. [16]. We can conclude that there are essentially no effects of the scalar isovector channel on isospin distillation in the spinodal decomposition.

VII. CONCLUSIONS AND OUTLOOK

We have developed a linear response theory starting from relativistic kinetic equations deduced within a quantum hadrodynamics effective field picture of the hadronic phase of nuclear matter. In the asymmetric case we consider as the main dynamical degrees of freedom the nucleon fields coupled to the *isoscalar*, scalar σ and ω, and to the *isovector*, scalar δ and vector ρ, mesons.

Using the Landau procedure we derive the dispersion relations which give the sound phase velocity and the internal structure of the normal collective modes, stable and unstable. We have focused our attention on the effect of the isovector mesons on the collective response of asymmetric (neutron-rich) matter. In order to better understand the dynamical role of the different mesons, the results are obtained in the Hartree approximation, which has a simpler and more transparent form. The contribution of Fock terms is also discussed.

We have singled out some qualitative new effects of the δ-meson-like channel on the dynamical response of ANM. Essentially, our investigation indicates that even if the symmetry energy is fixed, the dynamical response is affected by its internal structure, i.e., the presence or not of an isovector-scalar field. This is implemented by the explicit introduction of an effective δ meson and/or by the Fock term contributions. Both mechanisms are absent in the present relativistic RPA calculations for finite nuclei. In the spirit of the EFT-DFT approach [19] it would be interesting to see the effect of an isovector scalar field extension, at the lowest order, on the existing covariant RPA results. The richer sensitivity of the collective response on the density dependence of the symmetry term can be of large importance for our knowledge of the isovector part of the in medium interaction.

We like to remark that the same interplay between scalar (σ-meson) and vector (ω-meson) contributions can be seen in the dynamical isoscalar modes. In general we clearly show a close analogy in the structure of the linear response equations: (i) same form of the dispersion relations [cf. Eqs. (22) and (25)], (ii) parallel role of E_{sym}^{pot} and K_{NM}^{pot} in the determination of the restoring force [Eqs. (23) and (27)], and (iii) parallel structure of the corrections due to the scalar-vector meson competition [Eqs. (23) and (27)].

This appears to be a beautiful ‘‘mirror’’ structure of the relativistic approach that seems to nicely support the introduction of a δ-meson-like coupling in the isovector channel, at least from a formal point of view. We like to remind the reader that the same ‘‘mirror’’ structure of the relativistic picture has been recently stressed in Ref. [16] for equilibrium properties—saturation binding and symmetry energy—the a_1 and a_4 parameters of the Weiszaecker mass formula.

The relativistic dispersion relations have been compared with the nonrelativistic ones of the Landau Fermi liquid theory expressed in terms of Landau parameters evaluated in a relativistic scheme [29]. The corrections appear to be not negligible, particularly for the isoscalar response.

From the numerical results on the collective response of ANM some general features are qualitatively present in all the effective interactions in the isovector channel.

(i) In asymmetric matter we have a mixing of pure isoscalar and pure isovector oscillations which leads to a *chemical* effect on the structure of the propagating collective mode: the ratio of the neutron/proton density oscillations $\delta\rho_n/\delta\rho_p$ is different from the initial ρ_n/ρ_p of the matter at equilibrium. However, we can still classify the nature of the excited collective motions as *isoscalarlike* (when neutron and protons are oscillating in phase) and *isovectorlike* (out of phase). Note that similar effects can be obtained also using nonrelativistic effective forces [30,48].

(ii) For a given asymmetry the isovectorlike mode is the most robust at low baryon density, always showing a larger neutron component in the oscillations. With increasing baryon density we observe a smooth transition, at a $\rho_B^{cross} \simeq (2-3)\rho_0$, to an isoscalarlike branch, still with a dominant $\delta\rho_n$. In the region of the transition we predict a propagation of almost *pure neutron waves*. For relatively large asymmetries $[\alpha \equiv (N-Z)/(N+Z) = 0.5, N=3Z]$ this behavior is present in a wide interval of densities around ρ_B^{cross}. All that seems to suggest the possibility of an experimental observation of related effects in intermediate-energy heavy-ion collisions with exotic beams. If the compressibility of nuclear matter is decreasing at high baryon density, also these exotic isoscalarlike modes will disappear. This could be a nice signature of the softening of the nuclear EOS at high densities.

(iii) The isoscalarlike motions become unstable at subsaturation densities still with a strong chemical effect, now in the opposite direction with respect to the one discussed before, present in the stable high density modes. Now the unstable oscillation is more proton rich, eventually leading to the formation of more symmetric clusters versus a very neutron-rich gas phase. This is the *neutron distillation* effect [7,30,45,48,51,52], a new important feature of the liquid-gas phase transition in asymmetric nuclear systems.

ACKNOWLEDGMENTS

We warmly thank Hermann H. Wolter and Stephan Typel for several pleasant and stimulating discussions.

APPENDIX

The Wigner matrix $\hat{H}(p)$ for matter at equilibrium saturated in spin has the following form:

$$\hat{H}(p) = H(p) + \gamma_\mu H^\mu(p).$$

From the kinetic equation one obtains the relation for nuclear matter at equilibrium between the scalar and vector parts,

$$H^\mu(p) = \frac{p_\mu^*}{M^*} H^\mu(p),$$

where the zero component of the vector part is proportional to the Fermi Dirac distribution function:

$$H_0^{(i)}(p)=\frac{1}{4}\frac{1}{(2\pi)^3}\Theta(E_{Fi}^*-E_{ki}^*)\,\delta(p_0^{*(i)}-E_{ki}^*),$$

where $i=n,p$ and $E_{ki}^*=(k^2+M_i^{*2})^{1/2}$.

The coefficients $C^{(i)}(k)$, $C_\lambda^{(i)}(k)$, and $C_{\lambda\mu}^{(i)}(k)$ introduced in Sec. II are given by the integrals

$$C^{(i)}(k)=M_i^*\int d^4p\frac{H'^{\lambda}_{(i)}(p)k_\lambda}{p_\rho^{*i}k^\rho}, \tag{A1a}$$

$$C_\lambda^{(i)}(k)=\int d^4p\frac{H'^{\mu}_{(i)}(p)k_\mu p_\lambda^{*i}}{p_\rho^{*i}k^\rho}, \tag{A1b}$$

$$C_{\lambda\mu}^{(i)}(k)=\int d^4p\frac{H'^{\nu}_{(i)}(p)k_\nu p_\lambda^{*i}p_\mu^{*i}}{p_\rho^{*i}k^\rho}, \tag{A1c}$$

where $H'^{\lambda}_{(i)}(p)=\partial^\lambda_{(p)}H_{(i)}(p)$. The index i specifies the kind of nucleon: $i=1$ for protons and $i=2$ for neutrons. The frequency k^0 includes an imaginary part $i\epsilon$ with ϵ positive infinitesimal.

By using the definitions (A1) it can be easily checked that

$$k^\lambda C_\lambda^{(i)}(k)=0,\qquad k^\lambda C_{\lambda\mu}^{(i)}(k)=-k_\mu\frac{\rho_{Si}}{4},$$

$$\sum_\mu C_\mu^{(i)\mu}(k)=M_i^*C^{(i)}(k)-\frac{1}{2}\rho_{Si}\,. \tag{A2}$$

In order to be more specific we choose the z axis in the direction of the wave vector $\mathbf{k}$. As a consequence, the following coefficients identically vanish:

$$C_1^{(i)}(k),\ \ C_2^{(i)}(k),\ \ C_{10}^{(i)}(k),\ \ C_{20}^{(i)}(k),\ \ C_{lm}^{(i)}(k)$$

for $l\neq m$ (l and m are space indices). In addition, for symmetry reasons,

$$C_{11}^{(i)}(k)=C_{22}^{(i)}(k).$$

The integrals in Eqs. (A1) can be evaluated analytically. They give

$$C^{(i)}(k)=-\frac{1}{2}\frac{1}{M_i^*}\rho_{Si}+\frac{3}{4}\frac{\rho_{Bi}}{E_{Fi}^*}-\frac{1}{4}N_i\frac{{M_i^*}^2}{E_{F_i}^*}\varphi(s_i), \tag{A3a}$$

$$C_0^{(i)}=+\frac{1}{4}N_i\varphi(s_i), \tag{A3b}$$

$$C_{00}^{(i)}=-\frac{1}{4}\rho_{Si}+\frac{1}{4}N_iM_i^*\varphi(s_i), \tag{A3c}$$

$$C_{11}^{(i)}(k)=\frac{1}{4}\rho_{Si}-\frac{3}{8}\frac{M_i^*}{E_{Fi}^*}\rho_{Bi}+\frac{3}{8}\frac{M_i^*}{E_{Fi}^*}(s_i^2-1)\varphi(s_i), \tag{A3d}$$

where v_{F_i} is the Fermi velocity, $s_i=k^0/(v_{F_i}|\mathbf{k}|)$, N_i are the density of states at Fermi surface, and

$$\varphi(s_i)=1-\frac{s_i}{2}\ln\left|\frac{s_i+1}{s_i-1}\right|+\frac{i}{2}\pi s_i\,\theta(1-s_i)$$

is the Lindhard function. The remaining coefficients $C_3^{(i)}(k)$, $C_{03}^{(i)}(k)$, and $C_{33}^{(i)}(k)$ can be evaluated by means of the relations (A2).

[1] J.D. Walecka, Ann. Phys. (N.Y.) **83**, 491 (1974).
[2] B.D. Serot and J.D. Walecka, in *Advances in Nuclear Physics*, edited by J. W. Negele and E. Vogt (Plenum, New York, 1986), Vol. 16.
[3] B.D. Serot and J.D. Walecka, Int. J. Mod. Phys. E **6**, 515 (1997).
[4] M.M. Sharma, M.A. Nagarajan, and P. Ring, Ann. Phys. (N.Y.) **231**, 110 (1994); G.A. Lalazissis, M.M. Sharma, P. Ring, and Y.K. Gambhir, Nucl. Phys. **A608**, 202 (1996).
[5] P. Ring, Prog. Part. Nucl. Phys. **37**, 193 (1996).
[6] S. Typel and H.H. Wolter, Nucl. Phys. **A656**, 331 (1999).
[7] H. Müller and B.D. Serot, Phys. Rev. C **52**, 2072 (1995).
[8] B. Blattel, V. Koch, and U. Mosel, Rep. Prog. Phys. **56**, 1 (1993).
[9] C.M. Ko and G.Q. Li, J. Phys. G **22**, 1673 (1996).
[10] A. Dellafiore and F. Matera, Phys. Rev. C **44**, 2456 (1991).
[11] F. Matera and V.Yu. Denisov, Phys. Rev. C **49**, 2816 (1994).
[12] D. Vretenar, G.A. Lalazissis, R. Behnsch, W. Poeschl, and P. Ring, Nucl. Phys. **A621**, 853 (1997).
[13] Z.Y. Ma, N. Van Giai, H. Toki, and M. L'Huillier, Phys. Rev. C **55**, 2385 (1997).
[14] Z.Y. Ma, N. Van Giai, A. Wandelt, D. Vretenar, and P. Ring, Nucl. Phys. **A686**, 173 (2001), and references therein.
[15] C.J. Horowitz and K. Wehrberger, Nucl. Phys. **A531**, 665 (1991).
[16] B. Liu, V. Greco, V. Baran, M. Colonna, and M. Di Toro, Phys. Rev. C **65**, 045201 (2002).
[17] B.A. Nikolaus, T. Hoch, and D.G. Madland, Phys. Rev. C **46**, 1757 (1992).
[18] R.J. Furnstahl and J.J. Rusnak, Nucl. Phys. **A627**, 495 (1997).
[19] R.J. Furnstahl and B.D. Serot, Comments Nucl. Part. Phys. **2**, A23 (2000), and references therein.
[20] V. Greco, M. Colonna, M. Di Toro, G. Fabbri, and F. Matera, Phys. Rev. C **64**, 045203 (2001).
[21] R.J. Furnstahl and B.D. Serot, Nucl. Phys. **A671**, 447 (2000).
[22] J. Buervenich, D.G. Madland, J.A. Marhun, and P.G. Reinhard, Phys. Rev. C **65**, 044308 (2002).
[23] F. Hofmann, C.M. Keil, and H. Lenske, Phys. Rev. C **64**, 034314 (2001).

[24] R.J. Furnstahl, Nucl. Phys. **A706**, 85 (2002).
[25] E. Schiller and H. Muether, Eur. Phys. J. A **11**, 15 (2001).
[26] Z. Ma and L. Liu, Phys. Rev. C **66**, 024321 (2002).
[27] T. Matsui, Nucl. Phys. **365**, 365 (1981).
[28] C. Song, Phys. Rep. **347**, 289 (2001).
[29] J.C. Caillon, P. Gabinski, and J. Labarsoque, Nucl. Phys. **A696**, 623 (2001).
[30] M. Colonna, M. Di Toro, and A.B. Larionov, Phys. Lett. B **428**, 1 (1998).
[31] J. Boguta and C.E. Price, Nucl. Phys. **A505**, 123 (1989).
[32] C.J. Horowitz and B.D. Serot, Nucl. Phys. **A399**, 529 (1983).
[33] A. Bouyssy, J.F. Mathiot, Nguyen Van Giai, and S. Marcos, Phys. Rev. C **36**, 380 (1987).
[34] V. Greco, F. Matera, M. Colonna, M. Di Toro, and G. Fabbri, Phys. Rev. C **63**, 035202 (2001).
[35] S.R. de Groot, W.A. van Lee, and Ch.G. van Weert, *Relativistic Kinetic Theory* (North-Holland, Amsterdam, 1980).
[36] R. Hakim, Nuovo Cimento **6**, 1 (1978).
[37] This ''Dirac mass'' choice at normal density is actually slightly above the ''window'' (0.55–0.66)M_N compatible with effective spin-orbit potentials in finite nuclei; see [38], and references therein. This choice is dictated from the fact that with a smaller value we start to have some convergence numerical problems in our procedure to evaluate Fock terms. In any case our dynamical response will be not qualitatively modified by a reduction to 0.65M_N of the nucleon Dirac mass in symmetric matter. Moreover, as discussed in detail in Ref. [38], an upwards shift of the above ''window'' can be expected if explicit tensor coupling terms are included in the effective Lagrangian.
[38] R.J. Furnstahl, J.J. Rusnak, and B.D. Serot, Nucl. Phys. **A632**, 607 (1998).
[39] G.A. Lalazissis, J. Konig, and P. Ring, Phys. Rev. C **55**, 540 (1997).
[40] S. Kubis and M. Kutschera, Phys. Lett. B **399**, 191 (1997).
[41] L.D. Landau and E.M. Lifshitz, *Statistical Physics* (Pergamon Press, Oxford, 1980), p. 288.
[42] A.A. Abrikosov and I.M. Khalatnikov, Rep. Prog. Phys. **22**, 329 (1959).
[43] G. Baym and C.J. Pethick, in *The Physics of Liquid and Solid Helium*, edited by K.H. Bennemann and J.B. Ketterson (Wiley, New York, 1978), Vol. 2, p. 1.
[44] C.J. Pethick and D.G. Ravenhall, Ann. Phys. (N.Y.) **183**, 131 (1988).
[45] V. Baran, M. Colonna, M. Di Toro, and A.B. Larionov, Nucl. Phys. **A632**, 287 (1998).
[46] Ch. Fuchs, H. Lenske, and H.H. Wolter, Phys. Rev. C **52**, 3043 (1995).
[47] We remind the reader that if ϕ field has self-interaction terms instead of f_σ, one has to use $d\phi/d\rho_S$.
[48] V. Baran, M. Colonna, M. Di Toro, and V. Greco, Phys. Rev. Lett. **86**, 4492 (2001).
[49] P.G. Reinhard, Nucl. Phys. **A649**, 305c (1999).
[50] S. Typel and H.H. Wolter (private communication); ''Collective Modes of Nuclear Matter in Relativistic Mean Field Theories,'' Report LMU-Physik Section, Munich, 2002.
[51] H.S. Xu *et al.*, Phys. Rev. Lett. **85**, 716 (2000).
[52] M. Di Toro *et al.*, Nucl. Phys. **A681**, 426c (2001); V. Baran, M. Colonna, V. Greco, M. Di Toro, M. Zielinska Pfabé, and H.H. Wolter, *ibid.* **A706**, 603 (2002).