Aspects of short-range correlations in a relativistic model

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In the present work short-range correlations are introduced for the first time in a relativistic approach to the equation of state of the infinite nuclear matter in the framework of the Hartree-Fock approximation using an effective Hamiltonian derived from the σ-ω Walecka model. The unitary correlation method is used to introduce short-range correlations. The effect of the correlations in the ground state properties of the nuclear matter is discussed.

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Nuclear physics is an effective theory of the nucleus regarded as a system of nucleons. In this theory the essential degrees of freedom are the center of mass coordinates of the nucleons as well as their spins and isospins, the interactions being expressed as nucleon-nucleon forces. It is usually accepted, however, that the nucleon-nucleon interaction becomes strongly repulsive at short distances in the relative coordinate of two nuclear particles. It is seen that the phase shifts for $^1S_0$ and $^3S_1$ are positive at low energies and become negative at higher energies [1]. This indicates a repulsive core at short distances and attraction at long distances. There have been attempts to derive the nucleon-nucleon force using chiral perturbation theory [2]. In this approach, one- and two-pion exchange contributions are taken into account up to the third chiral order. However, to reproduce the nucleon scattering data and the D-wave phase shifts by this method, an *ad hoc* contact interaction (which represents the short-range force) must also be included. The realistic nucleon-nucleon forces are basically phenomenological. The Bonn interactions [3] are based on meson exchange treated in a relativistic nonlocal manner. The Argonne interactions [4], on the other hand, describe the pion exchange in a local approximation, the short- and medium-range nuclear interactions being controlled by phenomenological parameters.

Nonrelativistic calculations based on realistic *NN* potentials predict equilibrium points that are not able to describe simultaneously the correct binding energy and saturation density; either the saturation density is correct but the binding energy is too small, or the correct binding energy is obtained at a too high density. This behavior is normally referred to as "Coester line" [5]. This problem is generally circumvented through the introduction of a three-body repulsive force [6] or density-dependent repulsive mechanisms. A mechanism of this kind is already present in relativistic models. Because of Lorentz covariance and self-consistency, as the nuclear density increases, the nucleon effective mass decreases. As a result there is a reduction of the attractive force and a net increase of the repulsive force. The relativistic mean field (RMF) theory formulated by Teller and collaborators [7] and by Walecka [8,9] is successful in describing both infinite nuclear matter and finite nuclei. In fact, it has been shown in [9] that important relativistic effects present in RMF theory are equivalent to the inclusion of three-body or higher-order repulsive potentials in nonrelativistic calculations. Moreover, the relativistic Hartree-Fock (HF) approximation [10,11] has also been used for the description of finite nuclei and infinite nuclear matter.

In this paper, we focus on the treatment of short-range correlations in dense nuclear matter using for the first time relativistic equations of motion. Although there are several procedures which may be used to introduce short-range correlations into the model wave function, we preferred to work with the unitary operator method as proposed by Villars [12]. There are several advantages in using a unitary model operator. In particular, one automatically guarantees that the correlated state is normalized. Further, we shall see that the calculation of correlation corrections to the matrix elements of one- and two-body operators is a simple task. The general idea has existed for a long time [13–15] but has not been pursued for the relativistic case, because in nonrelativistic models the interaction arises from the interplay between a long-range attraction and a very strong short-range repulsion so that, indeed, it is indispensable to take short-range correlations into account. In RMF models, the parameters are phenomenological, fitted to the saturation properties of nuclear matter. Short-range correlation effects may be included to some extent in the model parameters. We expect that the values of these parameters are more fundamental, or less artificial, when correlations are not neglected.

We start with the effective Hamiltonian as

$$H = \int \psi_\alpha^\dagger(\vec{x})(-i\vec{\alpha}\cdot\vec{\nabla} + \beta M)_{\alpha\beta}\psi_\beta(\vec{x})\,d\vec{x} + \frac{1}{2}\int \psi_\alpha^\dagger(\vec{x})\psi_\gamma^\dagger(\vec{y})V_{\alpha\beta,\gamma\delta}(|\vec{x}-\vec{y}|)\psi_\delta(\vec{y})\psi_\beta(\vec{x})\,d\vec{x}\,d\vec{y}, \tag{1}$$

with

$$V_{\alpha\beta,\gamma\delta}(r) = (\beta)_{\alpha\beta}(\beta)_{\gamma\delta}V_\sigma(r) + (\delta_{\alpha\beta}\delta_{\gamma\delta} - \vec{\alpha}_{\alpha\beta}\cdot\vec{\alpha}_{\gamma\delta})V_\omega(r), \tag{2}$$

where

$$V_\sigma(r) = -\frac{g_\sigma^2}{4\pi}\frac{e^{-m_\sigma r}}{r}, \quad V_\omega(r) = \frac{g_\omega^2}{4\pi}\frac{e^{-m_\omega r}}{r}, \tag{3}$$

and $\vec{\alpha}$ are the Dirac matrices. In Eq. (1), ψ is the nucleon field interacting through the scalar and vector potentials. The equal time quantization condition for the nucleons is given by

$$[\psi_\alpha(\vec{x},t),\psi_\beta(\vec{y},t)^\dagger]_+ = \delta_{\alpha\beta}\delta(\vec{x}-\vec{y}), \tag{4}$$

where α and β refer to the spin indices. We now also have the field expansion for the nucleons ψ at time $t=0$ given as [16]

$$\psi(\vec{x}) = \frac{1}{\sqrt{V}}\sum_{r,k}\left[U_r(\vec{k})c_{r,\vec{k}} + V_r(-\vec{k})\tilde{c}^\dagger_{r,-\vec{k}}\right]e^{i\vec{k}\cdot\vec{x}}, \tag{5}$$

where U_r and V_r are given by

$$U_r(\vec{k}) = \begin{pmatrix} \cos\frac{\chi(\vec{k})}{2} \\ \vec{\sigma}\cdot\hat{k}\sin\frac{\chi(\vec{k})}{2}\end{pmatrix}u_r;\; V_r(-\vec{k}) = \begin{pmatrix} -\vec{\sigma}\cdot\hat{k}\sin\frac{\chi(\vec{k})}{2} \\ \cos\frac{\chi(\vec{k})}{2}\end{pmatrix}v_r. \tag{6}$$

For free spinor fields, we have $\cos\chi(\vec{k}) = M/\epsilon(\vec{k})$, $\sin\chi(\vec{k}) = |\vec{k}|/\epsilon(\vec{k})$ with $\epsilon(\vec{k}) = \sqrt{\vec{k}^2+M^2}$. However, we will deal with interacting fields so that we take the ansatz $\cos\chi(\vec{k}) = M^*(\vec{k})/\epsilon^*(\vec{k})$, $\sin\chi(\vec{k}) = |\vec{k}^*|/\epsilon^*(\vec{k})$, with $\epsilon^*(\vec{k}) = \sqrt{\vec{k^*}^2 + M^{*2}(\vec{k})}$, where $\vec{k}^*$ and $M^*(\vec{k})$ are the effective momentum and effective mass respectively and $\chi(\vec{k})$ will be determined variationally. The equal time anticommutation conditions are $[c_{r,\vec{k}}, c^\dagger_{s,\vec{k}'}]_+ = \delta_{rs}\delta_{\vec{k},\vec{k}'} = [\tilde{c}_{r,\vec{k}}, \tilde{c}^\dagger_{s,\vec{k}'}]_+$. The vacuum $|0\rangle$ is defined through $c_{r,\vec{k}}|0\rangle = \tilde{c}^\dagger_{r,\vec{k}}|0\rangle = 0$; one-particle states are written $|\vec{k},r\rangle = c^\dagger_{r,\vec{k}}|0\rangle$; two- and three-particle uncorrelated states are written, respectively, as $|\vec{k},r;\vec{k}',r'\rangle = c^\dagger_{r,\vec{k}}c^\dagger_{r',\vec{k}'}|0\rangle$, and $|\vec{k},r;\vec{k}',r';\vec{k}'',r''\rangle = c^\dagger_{r,\vec{k}}c^\dagger_{r',\vec{k}'}c^\dagger_{r'',\vec{k}''}|0\rangle$, and so on.

We now introduce short-range correlations through a unitary operator method. The correlated wave function [17] is $|\Psi\rangle = e^{iS}|\Phi\rangle$, where $|\Phi\rangle$ is a Slater determinant and S is, in general, an n-body Hermitian operator, splitting into a two-body part, a three-body part, etc. Consider now the expectation value of H,

$$E = \frac{\langle\Psi|H|\Psi\rangle}{\langle\Psi|\Psi\rangle} = \frac{\langle\Phi|e^{-iS}He^{iS}|\Phi\rangle}{\langle\Phi|\Phi\rangle}. \tag{7}$$

We restrict ourselves to the two-body correlation diagrams shown in Fig. 1. Let us denote the two-body correlated wave function by

$$|\overline{\vec{k},r;\vec{k}',r'}\rangle = e^{iS}|\vec{k},r;\vec{k}',r'\rangle \approx f_{12}|\vec{k},r;\vec{k}',r'\rangle, \tag{8}$$

where f_{12} is the short-range correlation factor, the so-called Jastrow factor [18]. For simplicity, we have written $f_{12} = f(\vec{r}_{12})$, $\vec{r}_{12} = \vec{r}_1 - \vec{r}_2$. We make the assumption that $f(r) = 1-(\alpha+\beta r)e^{-\gamma r}$ where α, β, and γ are parameters.

At this point some remarks on the choice of the Jastrow factor are appropriate. The important effect of the short-range correlations is the replacement, in the expression for the ground state energy, of the interaction matrix element $\langle\vec{k},r;\vec{k}',r'|V_{12}|\vec{k},r;\vec{k}',r'\rangle$ by $\langle\overline{\vec{k},r;\vec{k}',r'}|V_{12}+t_1+t_2|\overline{\vec{k},r;\vec{k}',r'}\rangle - \langle\vec{k},r;\vec{k}',r'|t_1+t_2|\vec{k},r;\vec{k}',r'\rangle$, where t_i is the kinetic energy operator of particle i. As argued by Moszkowski [19] and Bethe [20], it is expected that the true ground state wave function of the nucleus containing correlations coincide with the independent particle, or Hartree-Fock wave function, for interparticle distances $r \geqslant r_{\mathrm{heal}}$, where $r_{\mathrm{heal}} \approx 1$ fm is the so-called healing distance. This behavior is a consequence of the constraints imposed by the Pauli principle. Moreover, although in general the correlation factor $f(r)$ may depend on the isospin and spin quantum numbers of the two-body channel, we assume, for simplicity, that it is a plain, state-independent, Jastrow factor. In the present calculation, we are only studying the effect of correlations in the infinite symmetric nuclear matter and we are not even taking into account the contributions of the π and ρ mesons. This justifies the simple form for $f(r)$. A natural consequence of having the correlations introduced by an unitary operator is the occurrence of an normalization constraint on $f(r)$,

$$\int[f^2(r)-1]\,d^3r = 0. \tag{9}$$

The correlated ground-state energy becomes

$$\begin{aligned}\mathcal{E} = {} & \frac{\nu}{\pi^2}\int_0^{k_F}k^2\,dk\,[|k|\sin\chi(k)+M\cos\chi(k)] + \frac{\tilde{F}_\sigma(0)}{2}\rho_s^2 + \frac{\tilde{F}_\omega(0)}{2}\rho_B^2 + C_0\rho_B\frac{\nu}{\pi^2}\int_0^{k_F}k^2\,dk\,[|k|\sin\chi(k)+M\cos\chi(k)] \\ & -\frac{4}{(2\pi)^4}\int_0^{k_f}k^2\,dk\,k'^2\,dk'\{[|k|\sin\chi(k)+2M\cos\chi(k)]I(k,k') + |k|\sin\chi(k')J(k,k')\} \\ & +\frac{1}{(2\pi)^4}\int_0^{k_f}k\,dk\,k'dk'\left[\sum_{i=\sigma,\omega}A_i(k,k') + \cos\chi(k)\cos\chi(k')\sum_{i=\sigma,\omega}B_i(k,k') + \sin\chi(k)\sin\chi(k')\sum_{i=\sigma,\omega}C_i(k,k')\right],\end{aligned} \tag{10}$$

where $C_0 = \int[f^2(r)-1]\,d^3r$ so that, according to (9), $C_0 = 0$, and A_i, B_i, C_i, I, and J are exchange integrals and ν is the spin degeneracy factor. In the equation (10) for the energy density, the first term results from the kinetic contributions, the second and third terms come respectively from the σ and ω direct contributions from the potential energy with correlations, the fourth and fifth ones from the direct and exchange correlation contributions from the kinetic energy, and the last one from the $\sigma+\omega$ exchange contributions from the potential energy with correlations. The angular integrals are given by

$$A_i(k,k') = B_i(k,k') = 2\pi\frac{g_i^2}{4\pi}\int_0^\pi d(\cos\theta)\tilde{F}_i(k,k',\cos\theta),$$

$$C_i(k,k') = 2\pi\frac{g_i^2}{4\pi}\int_0^\pi \cos\theta\,d(\cos\theta)\tilde{F}_i(k,k',\cos\theta),$$

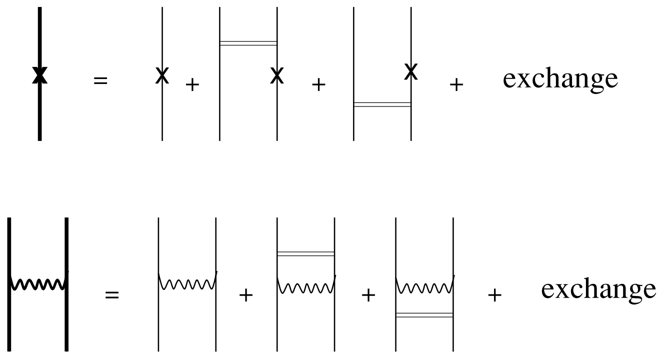

FIG. 1. The kinetic and potential contributions to the two-body cluster diagrams. Double lines represent two-body correlations.

$$I(k,k') = 2\pi \int_0^\pi d(\cos\theta)\tilde{C}_1(k,k',\cos\theta),$$

and

$$J(k,k') = 2\pi \int_0^\pi \cos\theta\, d(\cos\theta)\tilde{C}_1(k,k',\cos\theta),$$

where

$$\tilde{F}_i(\vec{k},\vec{k}') = \int [f(r)V_i(r)f(r)] e^{i(\vec{k}-\vec{k}')\cdot\vec{r}}\, d\vec{r} \quad \text{and} \quad \tilde{C}_1(\vec{k},\vec{k}') = \int [f^2(r)-1]\, e^{i(\vec{k}-\vec{k}')\cdot\vec{r}}\, d\vec{r}. \tag{11}$$

More explicitly, the first terms of the above angle integrals read for $f(r)=1$

$$A_\sigma = g_\sigma^2\theta_\sigma, \qquad B_\sigma = g_\sigma^2\theta_\sigma, \qquad C_\sigma = -2g_\sigma^2\phi_\sigma,$$

$$A_\omega = 2g_\omega^2\theta_\omega, \qquad B_\omega = -4g_\omega^2\theta_\omega, \qquad C_\omega = -4g_\omega^2\phi_\omega,$$

$$\theta_i(p,p') = \log\left(\frac{(k+k')^2+m_i^2}{(k-k')^2+m_i^2}\right),$$

$$\phi_i(p,p') = \frac{k^2+k'^2+m_i^2}{4kk'}\theta_i(p,p') - 1.$$

The baryon density and the scalar density are

$$\rho_B = \frac{2\nu}{(2\pi)^3}\int_0^{k_f} d\vec{k} = \frac{2\nu k_f^3}{6\pi^2}, \quad \rho_s = \frac{2\nu}{(2\pi)^3}\int_0^{k_f}\cos\chi(\vec{k})\, d\vec{k}. \tag{12}$$

The effective mass is given by

$$M^* = M + \tilde{F}_\sigma(0)\rho_s - \frac{2\times 2M}{(4\pi)^2}\int k'^2\, dk'\, I(k,k') + \frac{1}{(4\pi)^2 k}\int k'\, dk' \cos\chi(k')\sum_i B_i. \tag{13}$$

The parameters of the model have to be fixed. They are the couplings g_σ, g_ω, the meson masses m_σ and m_ω, and three more parameters from the short-range correlation function, α, β, and γ. The couplings are chosen to satisfy the ground-state properties of the nuclear matter and are given in Table I. We choose $m_\sigma = 550$ MeV and take $m_\omega = 783$ MeV. The normalization condition (9) determines β. We fix α either by imposing the condition $f(0)=0$, which appears to be a natural choice from our experience with the nonrelativistic case, or by minimizing the energy. We choose γ so that the correct healing distance [19] is reproduced. If α is chosen to be 1, we assume a density-independent parameter γ equal to 750 MeV (HF+corr-IV in Table I). On the other hand, if we choose to determine α variationally, we assume that the parameter γ increases linearly with the Fermi momentum. The last choice is consistent with the idea that the healing distance decreases as k_F increases. Of course there are other possible choices for these parameters. The parameters we have used are tabulated in Table I together with the compressibility K, the relative effective mass M^*/M, the kinetic energy $\mathcal{T}/\rho_B - M$, the direct and exchange parts of the potential energy ($\mathcal{V}_d/\rho_B$ and $\mathcal{V}_e/\rho_B$ repectively) with correlation, and the correlation contribution to the kinetic energy $\mathcal{T}^C/\rho_B$, all computed at the saturation point. We next compute the range of the correlations, obviously related to the healing distance, also included in Table I and defined as [20]

$$R^2 = \frac{\int r^2[f(r)-1]^2\, d^3r}{\int [f(r)-1]^2\, d^3r}. \tag{14}$$

In Fig. 2 we plot the binding energies as functions of the density for the Hartree, HF, quark-meson-coupling (QMC) model [21], a nonlinear Walecka model NL3 [22], and the four choices of parameters in our calculation, as given in Table I. From this figure one can see that in all cases the inclusion of correlations makes the equation of state (EOS)

TABLE I. Model parameters and ground state properties of nuclear matter at saturation density. These results were obtained with fixed $M = 939$ MeV, $m_\sigma = 550$ MeV, and $m_\omega = 783$ MeV at $k_{F0} = 1.3$ fm^{-1} with binding energy $E_B = \varepsilon/\rho - M = -15.75$ MeV. We used four choices of the parameters for the correlation: HF+corr-I corresponds to $\gamma = 700 + 200k_F/k_{F0}$ MeV; HF+corr-II to $\gamma = 600 + 300k_F/k_{F0}$ MeV; HF+corr-III to $\gamma = 600 + 400k_F/k_{F0}$ MeV; and HF+corr-IV to $\alpha = 1$ and $\gamma = 750$ MeV (fixed).

	g_σ	g_ω	α	β (MeV)	γ (MeV)	R fm	K (MeV)	M^*/M	$\mathcal{T}/\rho_B - M$ (MeV)	$\mathcal{V}_d/\rho_B$ (MeV)	$\mathcal{V}_e/\rho_B$ (MeV)	$\mathcal{T}^C/\rho_B$ (MeV)
Hartree	11.079	13.806					540	0.540	8.11	−23.86		
HF	10.432	12.223					585	0.515	5.87	−37.45	15.83	
HF+corr-I	4.3852	2.5157	13.656	−2068.381	900	0.303	475.	0.602	15.58	−82.75	23.53	27.89
HF+corr-II	4.3852	2.5157	13.656	−2068.381	900	0.303	463.	0.602	15.58	−82.75	23.53	27.89
HF+corr-III	4.4559	2.6098	13.855	−2252.448	1000	0.275	429.	0.625	15.95	−73.12	20.46	19.96
HF+corr-IV	11.836	13.608	1.0	−244.678	750	0.318	573.	0.515	7.16	−45.17	17.94	4.31

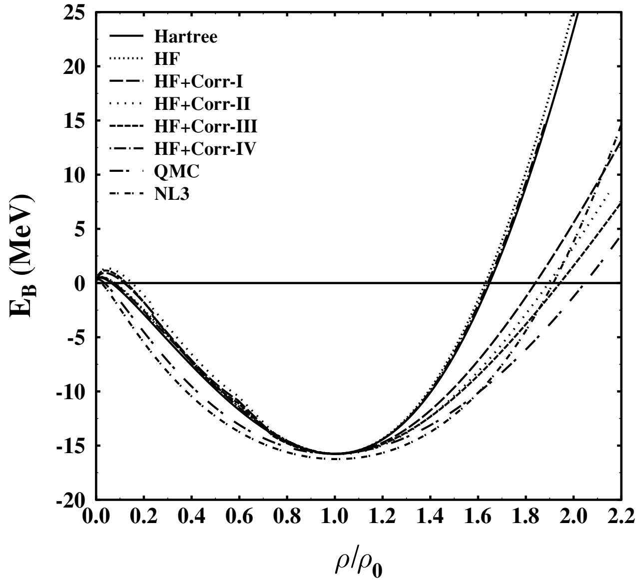

FIG. 2. EOS for different parametrizations as defined in Table I.

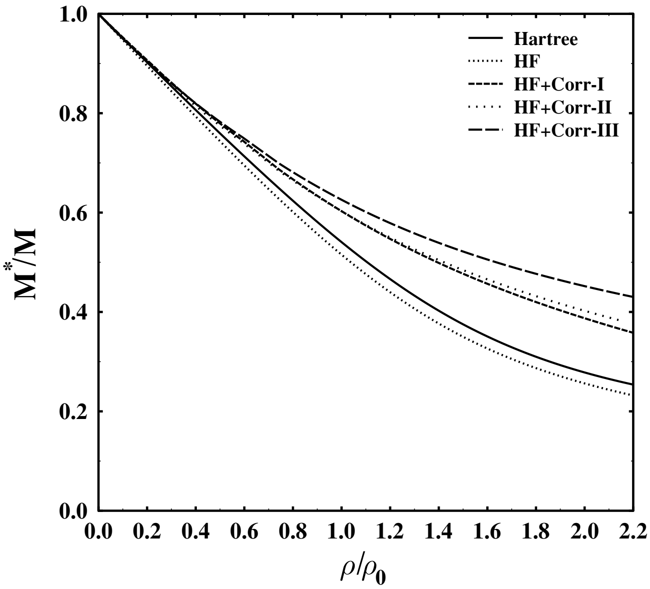

FIG. 4. Effective mass for different parametrizations as defined in Table I.

softer than the Hartree or HF calculations. NL3 and QMC also provide softer EOS around nuclear matter saturation density, but around two times saturation density, some of the EOS with correlations are softer than NL3. Correlations always tend to soften the EOS, except when $\alpha = 1$ is fixed. In this case, the EOS almost coincides with the curve obtained with the Hartree calculation.

In Fig. 3 we plot the individual contributions to the energy density as functions of density for the Hartree, HF, and HF+correlations (set II). Notice that the correlation contribution is of the same order as the exchange contribution. Hence, it cannot be disregarded. In fact if we compare the coupling constants g_σ and g_ω obtained for the different calculations, we conclude that the introduction of the correlations reduces the coupling constants. Correlation effects in the Hartree and HF calculations may be taken partially into account by a correct choice of the coupling constants. However, the explicit introduction of correlations has other effects such as softening the EOS.

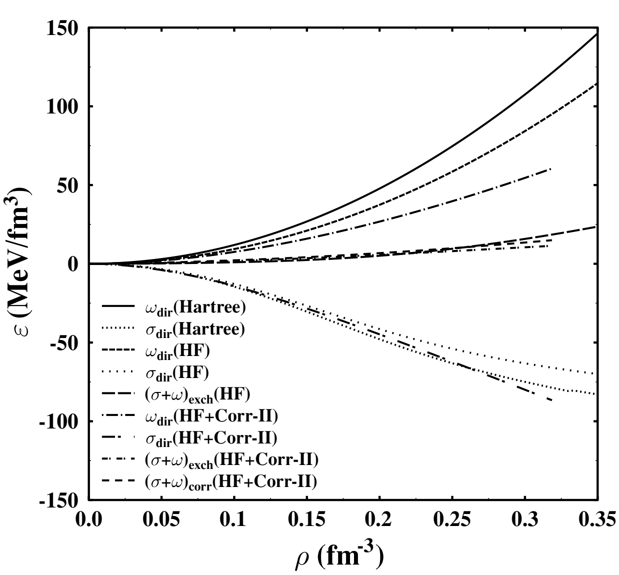

FIG. 3. Individual contributions for the energy density.

The effective mass as defined by Eq. (13) for $k = k_F$ is plotted as a function of density in Fig. 4. While HF reduces the Hartree effective mass for a given density, HF with correlations gives higher values of the effective mass. The parametrization chosen for γ determines the behavior: a higher value of γ at densities $\rho \geqslant \rho_0$ gives a softer EOS (Fig. 2) and a larger effective mass. Mean field theory approaches with low effective masses at saturation density have proved to be inadequate if one wants to extend the model for high densities with the inclusion of hyperons [23]. These extensions are very important in nuclear astrophysics. Hence, approximations that provide higher values of the effective mass at saturation density can be very convenient in describing compact stellar objects. Notice also that higher values of γ mean smaller healing distances. In fact, at higher densities, the correlated two-particle wave-function will not mix with Pauli blocked states (below the Fermi level) only if γ is large enough.

In conclusion, we have included two-body correlations in the σ-ω relativistic model. In the present approach, we have negelected retardation effects, since we are just dealing with a description of nuclear matter in equilibrium. We have verified numerically that the contribution of retardation terms in HF is negligible. The two-body correlation contribution corresponds to an important amount of the energy density and provides a softer EOS as compared with the HF approximation. Furthermore, the effect of the correlations is similar to the role played by the nonlinear terms in relativistic models or the three-body forces in nonrelativistic potential models. The inclusion of three-body correlations is currently under investigation.

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