Antiferromagnetic spin phase transition in nuclear matter with effective Gogny interaction

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The possibility of ferromagnetic and antiferromagnetic phase transitions in symmetric nuclear matter is analyzed within the framework of a Fermi liquid theory with the effective Gogny interaction. It is shown that at some critical density nuclear matter with the D1S effective force undergoes a phase transition to the antiferromagnetic spin state (opposite directions of neutron and proton spins). The self-consistent equations of spin polarized nuclear matter with the D1S force have no solutions corresponding to ferromagnetic spin ordering (the same direction of neutron and proton spins) and, hence, the ferromagnetic transition does not appear. The dependence of the antiferromagnetic spin polarization parameter as a function of density is found at zero temperature.

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I. INTRODUCTION

The spontaneous appearance of spin polarized states in nuclear matter is a topic of a great current interest due to its relevance in astrophysics. In particular, the effects of spin correlations in the medium strongly influence the neutrino cross section and, hence, neutrino mean free path. Therefore, depending on whether nuclear matter is spin polarized or not, drastically different scenarios of supernova explosion and cooling of neutron stars can be realized. Another aspect relates to pulsars, which are considered to be rapidly rotating neutron stars, surrounded by strong magnetic field. There is still no general consensus regarding the mechanism to generate such a strong magnetic field of a neutron star. One of the hypotheses is that a magnetic field can be produced by a spontaneous ordering of spins in the dense stellar core.

The possibility of a phase transition of normal neutron and nuclear matter to the ferromagnetic spin state was studied by many authors [1–8], predicting the ferromagnetic transition at $\varrho \approx (2-4)\varrho_0$ for different parametrizations of Skyrme forces ($\varrho_0 = 0.16\ \mathrm{fm}^{-3}$ is the nuclear matter saturation density). In particular, the stability of strongly asymmetric nuclear matter with respect to spin fluctuations was investigated in Ref. [9], where it was shown that the system with localized protons can develop a spontaneous polarization, if the neutron-proton spin interaction exceeds some threshold value. This conclusion was confirmed also by calculations within the relativistic Dirac-Hartree-Fock approach to strongly asymmetric nuclear matter [10]. Competition between ferromagnetic (FM) and antiferromagnetic (AFM) spin ordering in symmetric nuclear matter with the Skyrme effective interaction was studied in Ref. [11], where it was clarified that the FM spin state is thermodynamically preferable to the AFM one for all relevant densities. However, strongly asymmetric nuclear matter with Skyrme forces undergoes a phase transition to a state with oppositely directed spins of neutrons and protons [12].

For the models with realistic nucleon-nucleon (NN) interaction, the ferromagnetic phase transition seems to be suppressed up to densities well above ϱ_0 [13–15]. In particular, no evidence of ferromagnetic instability has been found in recent studies of neutron matter [16] and asymmetric nuclear matter [17] within the Brueckner-Hartree-Fock approximation with realistic Nijmegen II, Reid93, and Nijmegen NSC97e NN interactions. The same conclusion was obtained in Ref. [18], where the magnetic susceptibility of neutron matter was calculated with the use of the Argonne v_{18} two-body potential and Urbana IX three-body potential.

Thus, the issue of appearance of spin polarized states in nuclear matter is a controversial one and models with effective Skyrme and realistic NN potentials predict different results. From this point of view, it is interesting to attract another type of NN interaction and to compare the results for this NN potential with the previous results. Here we continue the study of spin polarizability of nuclear matter with the use of an effective NN interaction, namely, we utilize the effective Gogny force [19,20]. In addition, the reason to choose the Gogny interaction is as follows. It is known that the Skyrme interaction is a density dependent zero-range NN potential. Its attractive advantage is its relative simplicity and successfulness in describing nuclei and their excited states. However, in many cases the finite range part of the nuclear interaction has the same importance as its density dependent zero-range part [21,22]. This disadvantage of the Skyrme interaction is overcome by the Gogny interaction due to its finite range character. We will find the phase diagram of spin polarized nuclear matter with the Gogny interaction and will compare it with the results for the Skyrme and realistic NN potentials.

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As a framework of consideration, we choose a Fermi liquid (FL) description of nuclear matter [23–25]. We explore the possibility of FM and AFM phase transitions in nuclear matter, when the spins of protons and neutrons are aligned in the same direction or in the opposite direction, respectively. In contrast to the approach, based on the calculation of magnetic susceptibility, we obtain the self-consistent equations for the FM and AFM spin order parameters and solve them at zero temperature. This allows us not only to determine the critical density of instability with respect to spin fluctuations, but also to establish the density dependence of the order parameters and to clarify the question of thermodynamic stability of various phases.

Note that we consider the thermodynamic properties of spin polarized states in nuclear matter up to the high density region relevant for astrophysics. Nevertheless, we take into account the nucleon degrees of freedom only, although other degrees of freedom, such as pions, hyperons, kaons, or quarks could be important at such high densities.

II. BASIC EQUATIONS

The normal states of nuclear matter are described by the normal distribution function of nucleons $f_{\kappa_1\kappa_2}=\mathrm{Tr}\varrho a^+_{\kappa_2}a_{\kappa_1}$, where $\kappa\equiv(\mathbf{p},\sigma,\tau)$, $\mathbf{p}$ is the momentum, $\sigma(\tau)$ is the projection of spin (isospin) on the third axis, and ϱ is the density matrix of the system. The self-consistent matrix equation for determining the distribution function f follows from the minimum condition of the thermodynamic potential [23] and is

$$f=\{\exp(Y_0\varepsilon+Y_4)+1\}^{-1}\equiv\{\exp(Y_0\xi)+1\}^{-1}. \tag{1}$$

Here the nucleon single particle energy ε is defined through the energy functional of the system $E(f)$ as

$$\varepsilon_{\kappa_1\kappa_2}(f)=\frac{\partial E(f)}{\partial f_{\kappa_2\kappa_1}}. \tag{2}$$

In Eq. (1), the quantities ε and Y_4 are matrices in the space of κ variables, with $Y_{4\kappa_1\kappa_2}=Y_{4\tau_1}\delta_{\kappa_1\kappa_2}$ $(\tau_1=n,p)$, $Y_0=1/T$, $Y_{4n}=-\mu_n/T$, and $Y_{4p}=-\mu_p/T$ being the Lagrange multipliers, μ_n and μ_p being the chemical potentials of neutrons and protons, and T being the temperature.

Taking into account the possibility of FM and AFM phase transitions in nuclear matter, the normal distribution function f of nucleons and single particle energy ε can be expanded in the Pauli matrices σ_i and τ_k in spin and isospin spaces

$$f(\mathbf{p})=f_{00}(\mathbf{p})\sigma_0\tau_0+f_{30}(\mathbf{p})\sigma_3\tau_0+f_{03}(\mathbf{p})\sigma_0\tau_3+f_{33}(\mathbf{p})\sigma_3\tau_3, \tag{3}$$

$$\varepsilon(\mathbf{p})=\varepsilon_{00}(\mathbf{p})\sigma_0\tau_0+\varepsilon_{30}(\mathbf{p})\sigma_3\tau_0+\varepsilon_{03}(\mathbf{p})\sigma_0\tau_3+\varepsilon_{33}(\mathbf{p})\sigma_3\tau_3. \tag{4}$$

Expressions for the distribution functions f_{00}, f_{30}, f_{03}, f_{33} in terms of the quantities ε read

$$f_{00}=\frac{1}{4}\{n(\omega_{+,+})+n(\omega_{+,-})+n(\omega_{-,+})+n(\omega_{-,-})\},$$

$$f_{30}=\frac{1}{4}\{n(\omega_{+,+})+n(\omega_{+,-})-n(\omega_{-,+})-n(\omega_{-,-})\}, \tag{5}$$

$$f_{03}=\frac{1}{4}\{n(\omega_{+,+})-n(\omega_{+,-})+n(\omega_{-,+})-n(\omega_{-,-})\},$$

$$f_{33}=\frac{1}{4}\{n(\omega_{+,+})-n(\omega_{+,-})-n(\omega_{-,+})+n(\omega_{-,-})\}.$$

Here $n(\omega)=\{\exp(\omega/T)+1\}^{-1}$ and

$$\begin{aligned}\omega_{+,+}&=\xi_{00}+\xi_{30}+\xi_{03}+\xi_{33},\\ \omega_{+,-}&=\xi_{00}+\xi_{30}-\xi_{03}-\xi_{33},\\ \omega_{-,+}&=\xi_{00}-\xi_{30}+\xi_{03}-\xi_{33},\\ \omega_{-,-}&=\xi_{00}-\xi_{30}-\xi_{03}+\xi_{33},\end{aligned} \tag{6}$$

where

$$\xi_{00}=\varepsilon_{00}-\mu_{00},\quad \xi_{30}=\varepsilon_{30},$$

$$\xi_{03}=\varepsilon_{03}-\mu_{03},\quad \xi_{33}=\varepsilon_{33},$$

$$\mu_{00}=\frac{\mu_n+\mu_p}{2},\quad \mu_{03}=\frac{\mu_n-\mu_p}{2}.$$

The quantity $\omega_{\pm,\pm}$, being the exponent in the Fermi distribution function n, plays the role of the quasiparticle spectrum. In the general case, the spectrum is fourfold split due to the spin and isospin dependence of the single particle energy $\varepsilon(\mathbf{p})$ in Eq. (4). The branches $\omega_{\pm,+}$ correspond to neutrons with spin up and spin down, and the branches $\omega_{\pm,-}$ correspond to protons with spin up and spin down.

The distribution functions f should satisfy the normalization conditions

$$\frac{4}{\mathcal{V}}\sum_{\mathbf{p}}f_{00}(\mathbf{p})=\varrho, \tag{7}$$

$$\frac{4}{\mathcal{V}}\sum_{\mathbf{p}}f_{03}(\mathbf{p})=\varrho_n-\varrho_p\equiv\alpha\varrho, \tag{8}$$

$$\frac{4}{\mathcal{V}}\sum_{\mathbf{p}}f_{30}(\mathbf{p})=\varrho_\uparrow-\varrho_\downarrow\equiv\Delta\varrho_{\uparrow\uparrow}, \tag{9}$$

$$\frac{4}{\mathcal{V}}\sum_{\mathbf{p}}f_{33}(\mathbf{p})=(\varrho_{n\uparrow}+\varrho_{p\downarrow})-(\varrho_{n\downarrow}+\varrho_{p\uparrow})\equiv\Delta\varrho_{\uparrow\downarrow}. \tag{10}$$

Here α is the isospin asymmetry parameter, $\varrho_{n\uparrow}$, $\varrho_{n\downarrow}$ and $\varrho_{p\uparrow}$, $\varrho_{p\downarrow}$ are the neutron and proton number densities with spin up and spin down, respectively; $\varrho_\uparrow=\varrho_{n\uparrow}+\varrho_{p\uparrow}$ and $\varrho_\downarrow=\varrho_{n\downarrow}+\varrho_{p\downarrow}$ are the nucleon densities with spin up and spin down. The quantities $\Delta\varrho_{\uparrow\uparrow}$ and $\Delta\varrho_{\uparrow\downarrow}$ play the roles of FM and AFM spin order parameters [12].

In order to characterize spin ordering in the neutron and proton subsystems, it is convenient to introduce neutron and proton spin polarization parameters

$$\Pi_n = \frac{\varrho_{n\uparrow} - \varrho_{n\downarrow}}{\varrho_n}, \quad \Pi_p = \frac{\varrho_{p\uparrow} - \varrho_{p\downarrow}}{\varrho_p}. \tag{11}$$

The self-consistent equations for the components of the single particle energy have the form [11,12]

$$\xi_{00}(\mathbf{p}) = \varepsilon_0(\mathbf{p}) + \tilde{\varepsilon}_{00}(\mathbf{p}) - \mu_{00}, \quad \xi_{30}(\mathbf{p}) = \tilde{\varepsilon}_{30}(\mathbf{p}),$$

$$\xi_{03}(\mathbf{p}) = \tilde{\varepsilon}_{03}(\mathbf{p}) - \mu_{03}, \quad \xi_{33}(\mathbf{p}) = \tilde{\varepsilon}_{33}(\mathbf{p}). \tag{12}$$

Here $\varepsilon_0(\mathbf{p}) = \hbar^2\mathbf{p}^2/2m_0$ is the free single particle spectrum, m_0 is the bare mass of a nucleon, and $\tilde{\varepsilon}_{00}$, $\tilde{\varepsilon}_{30}$, $\tilde{\varepsilon}_{03}$, $\tilde{\varepsilon}_{33}$ are the FL corrections to the free single particle spectrum, related to the normal FL amplitudes $U_0(\mathbf{k}), \ldots, U_3(\mathbf{k})$ by formulas

$$\tilde{\varepsilon}_{00}(\mathbf{p}) = \frac{1}{2\mathcal{V}}\sum_{\mathbf{q}} U_0(\mathbf{k}) f_{00}(\mathbf{q}), \quad \mathbf{k} = \frac{\mathbf{p}-\mathbf{q}}{2},$$

$$\tilde{\varepsilon}_{30}(\mathbf{p}) = \frac{1}{2\mathcal{V}}\sum_{\mathbf{q}} U_1(\mathbf{k}) f_{30}(\mathbf{q}),$$

$$\tilde{\varepsilon}_{03}(\mathbf{p}) = \frac{1}{2\mathcal{V}}\sum_{\mathbf{q}} U_2(\mathbf{k}) f_{03}(\mathbf{q}),$$

$$\tilde{\varepsilon}_{33}(\mathbf{p}) = \frac{1}{2\mathcal{V}}\sum_{\mathbf{q}} U_3(\mathbf{k}) f_{33}(\mathbf{q}). \tag{13}$$

Further we do not take into account the effective tensor forces, which lead to coupling of the momentum and spin degrees of freedom [26–28], and, correspondingly, to anisotropy in the momentum dependence of FL amplitudes with respect to the spin polarization axis.

To obtain numerical results, we use the effective Gogny interaction. The amplitude of the NN interaction in this case reads

$$\hat{v}(\mathbf{p},\mathbf{q}) = t_0(1 + x_0\hat{P}_\sigma)\varrho^\gamma + \pi^{3/2}\sum_{i=1}^{2}\mu_i^3(W_i + B_i\hat{P}_\sigma - H_i\hat{P}_\tau - M_i\hat{P}_\sigma\hat{P}_\tau)e^{-(\mathbf{p}-\mathbf{q})^2\mu_i^2/4}, \tag{14}$$

where $\hat{P}_\sigma$ and $\hat{P}_\tau$ are the spin and isospin exchange operators, and t_0, x_0, μ_i, W_i, B_i, H_i, and M_i are some phenomenological constants, characterizing a given parametrization of the Gogny forces. In numerical calculations we shall utilize the D1S potential [20]. Using the same procedure as in Ref. [24], it is possible to find expressions for the normal FL amplitudes in terms of Gogny force parameters

$$U_0(\mathbf{k}) = 6t_0\varrho^\gamma + 2\pi^{3/2}\sum_{i=1}^{2}\mu_i^3(2B_i - 2H_i - M_i + 4W_i) - 2\pi^{3/2}\sum_{i=1}^{2}e^{-\mathbf{k}^2\mu_i^2}\mu_i^3(2B_i - 2H_i - 4M_i + W_i),$$

$$U_1(\mathbf{k}) = -2t_0\varrho^\gamma(1-2x_0) + 2\pi^{3/2}\sum_{i=1}^{2}\mu_i^3(2B_i - M_i) + 2\pi^{3/2}\sum_{i=1}^{2}e^{-\mathbf{k}^2\mu_i^2}\mu_i^3(2H_i - W_i),$$

$$U_2(\mathbf{k}) = -2t_0\varrho^\gamma(1+2x_0) - 2\pi^{3/2}\sum_{i=1}^{2}\mu_i^3(2H_i + M_i) - 2\pi^{3/2}\sum_{i=1}^{2}e^{-\mathbf{k}^2\mu_i^2}\mu_i^3(2B_i + W_i),$$

$$U_3(\mathbf{k}) = -2t_0\varrho^\gamma - 2\pi^{3/2}\sum_{i=1}^{2}\mu_i^3M_i - 2\pi^{3/2}\sum_{i=1}^{2}e^{-\mathbf{k}^2\mu_i^2}\mu_i^3W_i. \tag{15}$$

Thus, with account of expressions (5) for the distribution functions f, we obtain the self-consistent equations (12) and (13) for the components of the single particle energy $\xi_{00}(\mathbf{p})$, $\xi_{30}(\mathbf{p})$, $\xi_{03}(\mathbf{p})$, $\xi_{33}(\mathbf{p})$, which should be solved jointly with the normalization conditions (7)–(10), determining the chemical potentials μ_{00}, μ_{03} and FM and AFM spin order parameters $\Delta\varrho_{\uparrow\uparrow}, \Delta\varrho_{\uparrow\downarrow}$. Since the FL amplitudes in Eqs. (15) contain two Gaussian terms, the self-consistent equations represent, in fact, the set of coupled integral equations, which can be solved iteratively using the Gaussian mesh points in the momentum space.

III. PHASE TRANSITIONS IN SYMMETRIC NUCLEAR MATTER

Early research on spin polarizability of nuclear matter was based on the calculation of magnetic susceptibility and finding its pole structure [4,5], determining the onset of instability with respect to spin fluctuations. Here we shall solve directly the self-consistent equations for the FM and AFM spin order parameters at zero temperature. In this study we consider the case of symmetric nuclear matter ($\varrho_n = \varrho_p$).

The FM spin ordering corresponds to the case $\Delta\varrho_{\uparrow\uparrow} \neq 0$, $\xi_{30}(\mathbf{p}) \neq 0$, $\Delta\varrho_{\uparrow\downarrow} = 0$, $\xi_{33}(\mathbf{p}) = 0$, and there are two unknown parameters μ_{00}, $\Delta\varrho_{\uparrow\uparrow}$ and two unknown functions $\xi_{00}(\mathbf{p})$, $\xi_{30}(\mathbf{p})$ [$\mu_{03} = 0$, $\varepsilon_{03}(\mathbf{p}) = 0$ as a consequence of isospin symmetry]. The AFM spin ordering corresponds to the case $\Delta\varrho_{\uparrow\downarrow} \neq 0$, $\xi_{33}(\mathbf{p}) \neq 0$, $\Delta\varrho_{\uparrow\uparrow} = 0$, $\xi_{30}(\mathbf{p}) = 0$ and we should find two unknown parameters μ_{00}, $\Delta\varrho_{\uparrow\downarrow}$ and two unknown functions $\xi_{00}(\mathbf{p})$, $\xi_{33}(\mathbf{p})$.

In the FM spin state of symmetric nuclear matter we have $\varrho_{n\uparrow} = \varrho_{p\uparrow}$, $\varrho_{n\downarrow} = \varrho_{p\downarrow}$, nucleons with spin up fill the Fermi surface of radius k_2, and nucleons with spin down occupy the Fermi surface of radius k_1, which satisfy the relationships

$$\frac{1}{3\pi^2}(k_2^3 - k_1^3) = \Delta\varrho_{\uparrow\uparrow}, \quad \frac{1}{3\pi^2}(k_1^3 + k_2^3) = \varrho. \tag{16}$$

Since at zero temperature there are no spin up nucleons with $k > k_2$ and there are no spin down nucleons with $k > k_1$, then, as follows from Eq. (6), $\omega_{+,+}(k_2) = \omega_{+,-}(k_2) = 0$, $\omega_{-,+}(k_1) = \omega_{-,-}(k_1) = 0$.

In the AFM spin state $\varrho_{n\uparrow}=\varrho_{p\downarrow}$, $\varrho_{n\downarrow}=\varrho_{p\uparrow}$, neutrons with spin up and protons with spin down fill the Fermi surface of radius k_2, and neutrons with spin down and protons with spin up occupy the Fermi surface of radius k_1, satisfying the equations

$$\frac{1}{3\pi^2}(k_2^3-k_1^3)=\Delta\varrho_{\uparrow\downarrow}, \quad \frac{1}{3\pi^2}(k_1^3+k_2^3)=\varrho. \qquad (17)$$

At zero temperature there are no spin up neutrons and spin down protons with $k>k_2$ and there are no spin down neutrons and spin up protons with $k>k_1$. Hence, as follows from Eq. (6), $\omega_{+,+}(k_2)=\omega_{-,-}(k_2)=0$, $\omega_{+,-}(k_1)=\omega_{-,+}(k_1)=0$.

In the totally FM polarized state we have $\Delta\varrho_{\uparrow\uparrow}=\varrho$, $k_2=k_F\equiv(3\pi^2\varrho)^{1/3}$ and the degrees of freedom, corresponding to nucleons with spin down are frozen ($k_1=0$). For totally AFM polarized nuclear matter we have $\Delta\varrho_{\uparrow\downarrow}=\varrho$, $k_2=k_F$, where k_2 is given by the same expression as in the FM case, since now the degrees of freedom, corresponding to neutrons with spin down and protons with spin up, are frozen ($k_1=0$).

Now we present the results of the numerical solution of the self-consistent equations with the D1S Gogny effective force. The main qualitative feature is that for the D1S force there are solutions corresponding to AFM spin ordering and there are no solutions corresponding to FM spin ordering. The reason is that the sign of the multiplier $t_0(2x_0-1)$ in the density dependent term of the FL amplitude U_1, determining spin-spin correlations, is positive, and, hence, the corresponding term increases with increase of nuclear matter density, preventing instability with respect to spin fluctuations. Contrarily, the density dependent term $-2t_0\varrho^{\gamma}$ in the FL amplitude U_3, describing spin-isospin correlations, is negative, leading to spin instability with oppositely directed spins of neutrons and protons at higher densities. Here the situation is similar to that with the Skyrme effective forces SLy4 and SLy5 in strongly asymmetric nuclear matter [12], where analogous behavior of the FL amplitudes U_1 and U_3 in the high density domain prohibits the formation of the state with the same direction of neutron and proton spins and leads to the appearance of the state with the oppositely directed spins of neutrons and protons at high densities. However, the results with the Gogny effective interaction are in contrast with the results of microscopic calculations with a realistic NN interaction [17], predicting that the FM spin state is always preferable over the AFM one for all relevant densities, but is less favorable compared to the normal state.

In Fig. 1 it is shown the dependence of the AFM spin polarization parameter $\Delta\varrho_{\uparrow\downarrow}/\varrho$ as a function of density at zero temperature. The AFM spin order parameter arises at density $\varrho\approx3.8\varrho_0$ for the D1S potential. A totally antiferromagnetically polarized state ($\Delta\varrho_{\uparrow\downarrow}/\varrho=1$) is formed at $\varrho\approx4.3\varrho_0$. The neutron and proton spin polarization parameters for the AFM spin ordered state are opposite in sign and equal to

$$\Pi_n=-\Pi_p=\frac{\Delta\varrho_{\uparrow\downarrow}}{\varrho}.$$

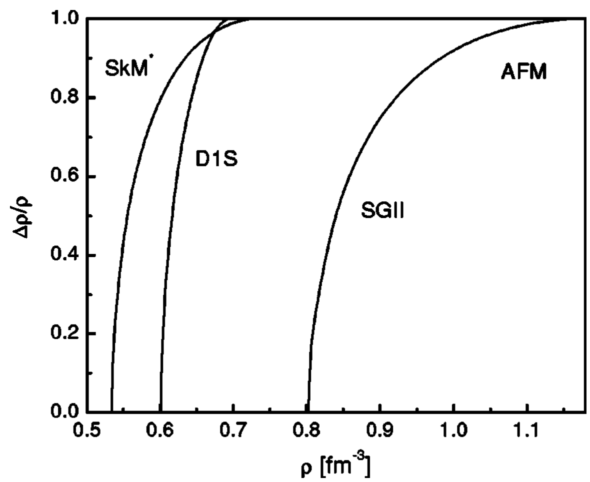

FIG. 1. AFM spin polarization parameter as a function of density at zero temperature for the D1S Gogny force and the SkM*, SGII Skyrme forces.

For comparison, we plot in Fig. 1 the density dependence of the AFM spin polarization parameter for the Skyrme effective forces SkM* and SGII [11]. It is seen that the results with the D1S potential are close to those with the SkM* potential (for the D1S force the AFM spin polarization parameter is saturated within a narrower density domain than for the SkM* force).

To check the thermodynamic stability of the spin ordered state with oppositely directed spins of neutrons and protons, it is necessary to compare the free energies of this state and the normal state. In Fig. 2, the difference between the total energies per nucleon of the spin ordered and normal states is shown as a function of density at zero temperature. One can see that nuclear matter in the model with the D1S effective force undergoes at some critical density a phase transition to the AFM spin state.

In Fig. 3, the difference between the total energies per nucleon of the spin polarized and normal states is decom-

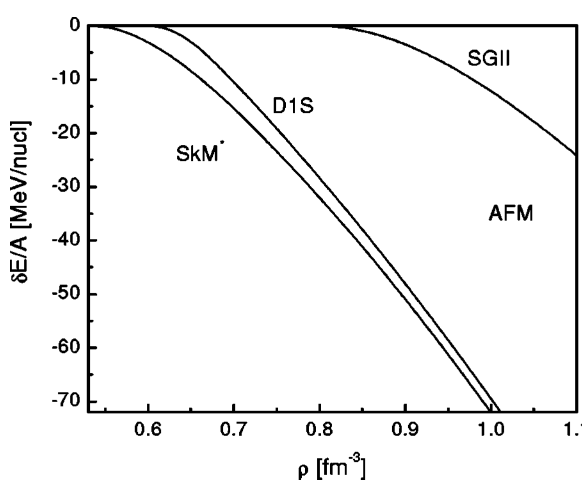

FIG. 2. Total energy per nucleon, measured from its value in the normal state, for the AFM spin state as a function of density at zero temperature for the D1S Gogny force and the SkM*, SGII Skyrme forces.

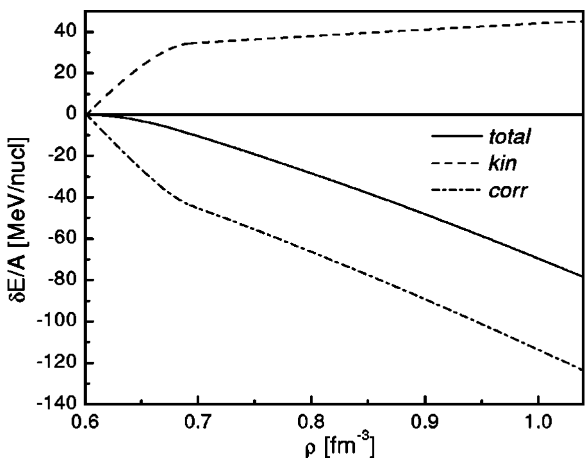

FIG. 3. Total energy per nucleon, measured from its value in the normal state, for the AFM spin state as a function of density at zero temperature for the D1S Gogny force, decomposed into kinetic and correlation parts.

posed into two parts, the kinetic and correlation ones. In spite of increasing the kinetic energy per nucleon in the AFM spin state, the AFM spin state becomes thermodynamically preferable over the normal state due to the energy gain caused by medium correlations, mainly by spin-isospin correlations, leading to AFM instability of the ground state.

IV. DISCUSSION AND CONCLUSIONS

Spin instability is a common feature associated with a large class of Skyrme models, but is not realized in more microscopic calculations. In this respect, it is interesting to study the possibility of appearance of spin polarized states in nuclear matter utilizing another type of NN potential. Here we used the finite range Gogny effective interaction, which is successful in describing nuclei and their excited states. The force parameters are determined empirically by calculating the ground state in the Hartree-Fock approximation and by fitting the observed ground state properties of nuclei and nuclear matter. The analysis based on the Gogny interaction shows that the self-consistent equations of symmetric nuclear matter have solutions corresponding to AFM spin ordering and have no solutions at all corresponding to FM spin ordering. This result is in contrast to calculations with the Skyrme interaction, predicting a FM phase transition in symmetric nuclear matter [11]. In the last case, the SkM*, SGII parametrizations were used, developed to fit the properties of nucleon systems with small isospin asymmetry. However, if the spin ordered state with oppositely directed spins of neutrons and protons for the Gogny interaction survives in the domain of high isospin asymmetry, then this result can coincide with the calculations for the Skyrme interaction, predicting the appearance of this state in strongly asymmetric nuclear matter [12]. In the last case, the calculations were done with the SLy4,SLy5 parametrizations, developed to reproduce the properties of nuclear matter with high isospin asymmetry.

In a microscopic approach, one starts with the bare interaction and obtains an effective particle-hole interaction by solving iteratively the Bethe-Goldstone equation. In contrast to the Skyrme and Gogny models, calculations with realistic NN potentials predict more repulsive total energy per particle for a polarized state comparing to the nonpolarized one for all relevant densities, and, hence, give no indication of a phase transition to a spin ordered state.

It must be emphasized that the interaction in the spin and isospin dependent channels is a crucial ingredient in calculating spin properties of nuclear matter and different behavior at high densities of the interaction amplitudes describing spin-spin and spin-isospin correlations lies behind this divergence in calculations with effective and realistic potentials, from one side, and calculations with different types of effective forces, from the other side. Since our calculations with the Gogny interaction predict the AFM spin state as the ground state of symmetric nuclear matter, this emphasizes the role of spin-isospin correlations in the high density region. Due to the antiferromagnetic spin polarization, some neutrons and protons with opposite spins, e.g., spin up neutrons and spin down protons, fill the Fermi surface with the larger radius and others, spin down neutrons and spin up protons, occupy the Fermi surface with the smaller radius. When density increases, some neutrons and protons undergo spin-flip transitions from the inner Fermi surface to the outer one due to increase of spin-isospin correlations. The usual way to constrain the interaction parameters of spin dependent amplitudes is based on the data on isoscalar [29] and isovector (giant Gamow-Teller) [30] spin-flip resonances. However, it is necessary to note that in order to get robust results for the spin polarization phenomena, these constraints should be obtained for the high density region of nuclear matter. Probably, such constraints can be obtained from the data on the time decay of the magnetic field of isolated neutron stars [31].

In summary, we have considered the possibility of phase transitions into spin ordered states of symmetric nuclear matter within the Fermi liquid formalism, where the NN interaction is described by the D1S Gogny effective force. In contrast to the previous considerations, where the possibility of formation of FM spin polarized states was studied on the basis of calculation of the magnetic susceptibility, we obtain self-consistent equations for the FM and AFM spin order parameters and solve them at zero temperature. It has been shown in the model with the D1S effective force that symmetric nuclear matter undergoes a phase transition to the spin polarized state with oppositely directed spins of neutrons and protons, while the state with the same direction of the neutron and proton spins does not appear. The AFM spin order parameter arises at density $\varrho \approx 3.8\varrho_0$ and is saturated at $\varrho \approx 4.3\varrho_0$. These results may be of importance for the adequate description of spin related phenomena in the interior of neutron stars.

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