Magnetic susceptibility of neutron matter

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(Received 5 November 1974)

The magnetic susceptibility of neutron matter is calculated in Brueckner theory with the Reid hard core potential and the Mongan nonlocal separable potential. The rearrangement contribution, calculated with suitable approximations, turns out to be important at the densities above normal nuclear density. The results obtained for the Reid hard core potential are similar to those obtained previously for the Reid soft-core potential. No evidence of a transition to ferromagnetic state is found.

I. INTRODUCTION

It is currently believed that pulsars are rapidly spinning neutron stars possessing intense magnetic fields,[1] and hence the magnetic properties of the neutron star matter are of a considerable interest. The matter of which real neutron stars are composed is a very complicated substance.[2] However, in the normal nuclear density region the neutron star matter is believed to consist mainly of neutrons, with a small admixture of protons, electrons, and muons. Some of the neutron star matter calculations[3] indicate that also at high densities ($k_F > 3$ fm^{-1}) the ground state of the cold dense matter could be nearly pure neutron matter. Thus, pure neutron matter appears to be a meaningful approximation of the real neutron star matter near and above nuclear density and does provide a starting point for the calculation which would take into consideration more complex effects.

Several attempts have been made in explaining the huge magnetic fields of the pulsars (10^{12}–10^{14} Gs) by means of the ferromagnetic transition of neutron matter at the densities much higher than the normal nuclear density. This possibility has been studied in Refs. 4–7 with rather oversimplified n-n potentials. In the calculations presented in Refs. 8–14 the magnetic susceptibility of neutron matter has been obtained using realistic n-n potentials. Clark[9] calculated the magnetic susceptibility of neutron matter within the frame of the Brueckner theory, assuming the Reid soft-core (RSC) n-n potential,[15] in the density region corresponding to $k_F < 2.02$ fm^{-1}. However, he neglected the rearrangement contribution resulting from the intrinsic density dependence of the effective n-n interaction in neutron matter. Pearson and Saunier[10] investigated the magnetic properties of neutron matter in the first order perturbation theory in the narrow density interval, corresponding to 1 fm$^{-1} < k_F < 1.8$ fm^{-1}, using the realistic effective n-n interaction. Pandharipande, Garde, and Srivastava[11] calculated the magnetic susceptibility of neutron matter in the wide density region, corresponding to $k_F < 5$ fm^{-1}, for the RSC potential, using lowest order variational method. In the paper of Pfarr[12] the magnetic properties of neutron matter were investigated for $k_F < 3$ fm^{-1} in the Hartree-Fock approximation by the method of the unitary transformation for the old hard-core Brueckner-Gammel-Thaler[16] potential and the RSC potential. Bäckman, Källman, and Sjöberg[13] applied the Landau theory of Fermi liquid. In their paper the Landau parameters were calculated from the RSC potential at $k_F < 2.5$ fm^{-1}. Moszkowski[14] calculated the magnetic susceptibility of neutron matter using the effective n-n interactions, constructed for the RSC potential by Sawada and Wong in their nuclear matter calculations. The density region considered in Moszkowski's calculation, corresponds to 1.25 fm$^{-1} < k_F < 3.78$ fm^{-1}.

In all Refs. 8–14, 17 no sign of the tendency to the ferromagnetic transition has been obtained. On the contrary, in all cases the calculated ratio χ_F/χ of the magnetic susceptibilities of neutron matter and Fermi gas of free neutrons is greater than 1 and increases monotonically with density.

No attempt of a complete calculation of χ_F/χ in Brueckner theory with realistic n-n interaction has been made. In Clark's calculation[9] the rearrangement contribution to χ_F/χ was neglected, while in Østgaard's paper[7] the calculation was not self-consistent and the potential used was unrealistic.

In the present paper the magnetic susceptibility of neutron matter is calculated within the frame of Brueckner theory for the Reid hard-core (RHC) potential[15] and the nonlocal separable Mongan (M) potential.[18] The calculations were carried out for the densities corresponding to 1.0 fm$^{-1} < k_F < 2.6$ fm^{-1} in the case of the RHC potential and for 1.0 fm$^{-1} < k_F < 3.3$ fm^{-1} in the case of the M potential. In Sec. II general formulas for various terms of χ_F/χ are given. Numerical results and their discussion are presented in Sec. III.

II. GENERAL FORMULAS

Let us consider neutron matter composed of N neutrons contained in a periodicity box of volume Ω. In the presence of an external magnetic field, the two spin populations will no longer be equal in the ground state, and the total energy will be a function of the spin excess parameter

$$\alpha = (N_\uparrow - N_\downarrow)/N \ , \tag{1}$$

where $N_\uparrow$ and $N_\downarrow$ are the numbers of neutrons with spin up and spin down with respect to the direction of the applied field. The total energy of the system is the sum of the three terms,

$$E^H(N, \alpha) = E_{\rm kin}(N, \alpha) + E^n_{\rm pot}(N, \alpha) - \mu_n H N \alpha \ , \tag{2}$$

where $E_{\rm kin}$ is the kinetic energy of the neutron matter, $E^n_{\rm pot}$ is the nuclear potential energy of the system, and the last term is perturbation describing the interaction with the external magnetic field H, weak enough that its magnitude enters lineary into the energy of the system. The neutron magnetic moment is denoted by μ_n. The nuclear energy per particle,

$$\frac{1}{N} E^n(N, \alpha) = \frac{1}{N} (E_{\rm kin} + E^n_{\rm pot}) \ ,$$

expanded in powers of α takes the form

$$\frac{1}{N} E^n(N, \alpha) = \epsilon_0 + \tfrac{1}{2}\epsilon_\sigma \alpha^2 \ , \tag{3}$$

where ϵ_0 is the energy per neutron in the case of $\alpha = 0$, ϵ_σ is the spin symmetry energy of the neutron matter, and the spin excess parameter α is assumed to be small, so that the terms proportional to α^n, $n > 2$, can be neglected. The terms proportional to the odd powers of α vanish because of the condition

$$E^n(N, \alpha) = E^n(N, -\alpha) \ ,$$

resulting from the time reversal invariance of the neutron-neutron interaction. In the ground state the total energy of the system reaches the minimum for a fixed value of the spin excess parameter α_0,

$$\alpha_0 = \frac{\mu_n H}{\epsilon_\sigma} \ , \tag{4}$$

and thus the magnetic susceptibility of neutron matter is in our approximation given by

$$\chi = \frac{{\mu_n}^2 \rho}{\epsilon_\sigma} \ ,$$

where ρ is the density of neutron matter.

For the sake of convenience, we shall introduce a ratio of χ and the magnetic susceptibility of the Fermi gas of the free neutrons χ_F,

$$\chi_F = \frac{3}{2} {\mu_n}^2 \frac{\rho}{\epsilon_F} \ . \tag{5}$$

This ratio is given by the formula:

$$\frac{\chi_F}{\chi} = \frac{3}{2} \frac{\epsilon_\sigma}{\epsilon_F} \ , \tag{6}$$

where ϵ_F is the Fermi energy for unpolarized neutron matter with $N_\uparrow = N_\downarrow = \frac{1}{2}N$,

$$\epsilon_F = \frac{{k_F}^2}{2\mathfrak{M}} \ .$$

The Fermi momentum k_F, measured in units of $\hbar$, is given by

$$\rho = \frac{N}{\Omega} = \frac{{k_F}^3}{3\pi^2} \ ,$$

and $\mathfrak{M}$ is the neutron mass divided by $\hbar^2$. Thus, the calculation of the ratio χ_F/χ is equivalent to the calculation of the spin symmetry energy of the neutron matter ϵ_σ. Expansion (3) of $E_{\rm kin}$ gives us the value of the kinetic part of ϵ_σ ,

$$\epsilon_\sigma^{\rm kin} = \tfrac{2}{3}\epsilon_F \ ,$$

while the potential part of the spin symmetry energy is given by

$$\epsilon_\sigma^{\rm pot} = \frac{1}{N} \left(\frac{\partial^2 E^n_{\rm pot}}{\partial \alpha^2} \right)_{\rho, N} \Bigg|_{\alpha = 0} \ . \tag{7}$$

In the lowest order Brueckner theory the nuclear potential energy of the neutron matter is given by the expression

$$E^n_{\rm pot} = \frac{1}{2} \sum_{s_3 s_3'} \sum_{\vec{m}_1}^{(s_3)} \sum_{\vec{m}_2}^{(s_3')} \{(\vec{m}_1 s_3 \vec{m}_2 s_3' \,|\, K(\kappa\lambda) |\, \vec{m}_1 s_3 \vec{m}_2 s_3') - \text{exchange}\}, \tag{8}$$

where the sum $\sum_{\vec{m}}^{(s_3)}$ denotes summation over all momenta states occupied by neutrons with the third component of the spin equal to s_3, and the intrinsic dependence of the matrix on the two Fermi momenta for neutrons with spin up and with spin down,

$$\kappa^3 = {k_F}^3(1 + \alpha) \ , \tag{9a}$$

$$\lambda^3 = {k_F}^3(1 - \alpha) \ , \tag{9b}$$

has been explicitly shown. The order in which the Fermi momenta appear as arguments of K will be kept the same throughout this paper; namely, the first argument always stands for the Fermi momentum of spin-up neutrons. Thus, nuclear potential energy depends on α in two ways: first, through the upper limits of the sums over $\vec{m}_1$ and $\vec{m}_2$ in Eq. (8) and, secondly, through the intrinsic dependence of K on the two Fermi momenta, related to α by Eqs. (9). Hence, when we calculate the second derivative of $E^n_{\rm pot}$ indicated in Eq. (7) we get two parts of $\epsilon^{\rm pot}_\sigma$,

$$\epsilon^{\rm pot}_\sigma = \epsilon^{(0)\,\rm pot}_\sigma + \Delta\epsilon_\sigma \,, \tag{10}$$

the first part, $\epsilon^{(0)\,\rm pot}_\sigma$, resulting from the first type of the dependence, and the second part, $\Delta\epsilon_\sigma$, resulting from the second type of the dependence of $E^n_{\rm pot}$ on α. The term $\Delta\epsilon_\sigma$ will be referred to as the rearrangement part of ϵ_σ, and the nonrearrangement part of ϵ_σ will be denoted by

$$\epsilon^{(0)}_\sigma = \tfrac{2}{3}\epsilon_F + \epsilon^{(0)\,\rm pot}_\sigma \,.$$

The calculation of $\epsilon^{\rm pot}_\sigma$, i.e., the calculation of the second derivative in Eq. (7), is lengthy but straightforward. The final result for $\epsilon^{(0)\rm pot}$ is

$$\epsilon^{(0)\,\rm pot}_\sigma = N\int \frac{d\hat{k}'_F}{4\pi}\Big(\vec{k}_F\vec{k}'_F\Big|K(11;k_F) - \frac{1}{2}\sum_s K(s0;k_F)\Big|\vec{k}_F\vec{k}'_F\Big) + \tfrac{1}{3}k_F\left(\frac{\partial V(m)}{\partial m}\right)_{m=k_F}, \tag{11}$$

where $K(s\,m_s;k_F)$ and $V(m)$ are, respectively, the effective interaction in the representation of the total spin s of the neutron pair and the model single-particle potential, calculated in the case of neutron matter with $N_\uparrow = N_\downarrow = \frac{1}{2}N$. For the rearrangement part of the spin symmetry energy we obtain

$$\Delta\epsilon_\sigma = -\tfrac{2}{3}V_R(k_F) + 2N\int\frac{d\hat{k}_F}{4\pi}\int_m\left(\vec{k}_F\,\vec{m}\Big|k_F\left(\frac{\partial}{\partial\kappa}-\frac{\partial}{\partial\lambda}\right)K(11;\kappa\lambda)\Big|\vec{k}_F\,\vec{m}\right) + \tfrac{1}{2}N\int_m\int_{m'}\left(\vec{m}\vec{m}'\Big|k_F^2\left(\frac{\partial}{\partial\kappa}-\frac{\partial}{\partial\lambda}\right)^2\left[K(11;\kappa\lambda)+\frac{1}{2}\sum_s K(s0;\kappa\lambda)\right]\Big|\vec{m}\vec{m}'\right), \tag{12}$$

where all the derivatives are to be calculated at the point $\kappa=\lambda=k_F$, and where

$$\int_m = \frac{1}{4\pi k_F{}^3}\int_{|\vec{m}|<k_F} d^3m \,.$$

In Eq. (12) V_R is the rearrangement potential, calculated at the Fermi surface in the case of $N_\uparrow = N_\downarrow = \frac{1}{2}N$,

$$V_R(k_F) = \frac{3N}{4}\int_m\int_{m'}\left(\vec{m}\vec{m}'\Big|k_F\frac{d}{dk_F}\sum_{s m_s}K(s\,m_s;k_F)\Big|\vec{m}\vec{m}'\right).$$

In all the above formulas the notation of Ref. 17 has been used. Let us notice that in Eq. (12) we omitted the cross terms, resulting from the matrix elements of the effective interaction between the spin singlet and spin triplet states. These elements of the K matrix vanish in the approximation, in which the exclusion principle operator in the Bethe-Goldstone equation for the K matrix is replaced by the value averaged over the directions of the total momentum of the neutron pair.

Although it is possible to calculate the K matrix which depends on two different Fermi momenta, this calculation is very tedious. For instance, such a calculation for nuclear matter with a neutron excess has been performed in Refs. 19–21. On the other hand, an approximation introduced by Brueckner and Gammel[22] enables one to calculate $\Delta\epsilon_\sigma$ using the K matrix for unpolarized neutron matter. This approximation has the form

$$K(11;\kappa\lambda)\approx K(11;\kappa)\,, \tag{13a}$$

$$\sum_s K(s0;\kappa\lambda)\approx\sum_s K(s0;\omega)\,, \tag{13b}$$

where

$$\omega = 2^{-1/2}(\kappa^2+\lambda^2)^{1/2}\,, \tag{14}$$

and where the K matrices on the right-hand side of Eqs. (13) are calculated in the case of $\alpha=0$ with indicated value of the Fermi momentum. When expressions (13) are introduced into Eq. (12) we get the following approximate formulas for the rearrangement part of the spin symmetry energy

of neutron matter:

$$\Delta\epsilon_\sigma = -\tfrac{2}{3}V_R(k_F) + 2N\int \frac{d\hat{k}_F}{4\pi}\int_m \left(\vec{k}_F\,\vec{m}\,\middle|\, k_F\frac{d}{dk_F}K(11;k_F)\middle|\vec{k}_F\,\vec{m}\right)$$
$$+ \tfrac{1}{2}N\int_m\int_{m'}\left(\vec{m}\vec{m}'\middle|\left[k_F{}^2\frac{d^2}{dk_F{}^2}K(11;k_F)+\tfrac{1}{2}k_F\frac{d}{dk_F}\sum_s K(s0;k_F)\right]\middle|\vec{m}\vec{m}'\right). \quad (15)$$

III. NUMERICAL RESULTS AND DISCUSSION

The K matrix has been calculated self-consistently within the frame of the Brueckner theory over a range of neutron matter densities for the RHC and M neutron-neutron potentials. The calculation was performed using the angle-averaged Pauli exclusion principle operator and the effective mass approximation for the hole states. Pure kinetic energies in the intermediate states were used. This standard choice of the single-particle potential in the intermediate states could not be used at the high density region and hence the calculation was restricted to $k_F < 3.3$ fm^{-1} in the case of the M potential and $k_F < 2.6$ fm^{-1} in the case of the RHC potential. The details of the calculation are presented in Ref. 23. The differentiation of the K matrix was carried out numerically, by making finite shifts in the Fermi momentum, $k_F \to k_F + h$, and repeating the whole self-consistent calculation. Then the derivatives were determined from the resulting shifts in the matrix, using the five point Lagrange interpolation formula. The contribution to ϵ_σ from the $J > 2$ partial waves was estimated in the first Born approximation for one-pion exchange potential.

The quotient χ_F/χ can be split into two parts,

$$\chi_F/\chi = (\chi_F/\chi)_0 + (\chi_F/\chi)_R ,$$

where $(\chi_F/\chi)_0$ and $(\chi_F/\chi)_R$ are, respectively, the nonrearrangement and rearrangement parts of χ_F/χ. The values of $(\chi_F/\chi)_0$ and $(\chi_F/\chi)_R$ for the RHC and M potentials are given in Table I. Near and below the normal nuclear density ($k_F < 2$ fm^{-1}) the values of χ_F/χ for the RHC and M potentials are approximatively the same. At higher densities ($k_F > 2$ fm^{-1}) the values of the nonrearrangement parts of χ_F/χ stay nearly the same for both potentials, while the values of the rearrangement parts differ significantly. Thus, the difference between the RHC and the M values of χ_F/χ results mainly from the differences between the rearrangement parts. In Fig. 1 our values of χ_F/χ are plotted versus k_F. The results of Pfarr,[12] Moszkowski,[14] and Pandharipande *et al.*[11] for the RSC potential are shown for comparison. Our results for the RHC potential agree quite well with those obtained by Pandharipande *et al.* The discrepancy between the Moszkowski's, Pandharipande's, and Pfarr's results for the same neutron-neutron interaction

TABLE I. Values of the nonrearrangement and rearrangement parts of χ_F/χ for the RHC and M potentials.

k_F (fm^{-1})	$(\chi_F/\chi)_0$ M	$(\chi_F/\chi)_0$ RHC	$(\chi_F/\chi)_R$ M	$(\chi_F/\chi)_R$ RHC	χ_F/χ M	χ_F/χ RHC
1.0	1.62	1.76	0.03	−0.02	1.65	1.74
1.2	1.66	1.82	0.19	0.03	1.85	1.85
1.4	1.73	1.89	0.27	0.05	2.00	1.94
1.6	1.84	1.92	0.32	0.14	2.16	2.06
1.8	1.95	1.97	0.34	0.28	2.29	2.25
2.0	2.05	2.04	0.35	0.38	2.40	2.42
2.2	2.15	2.12	0.33	0.47	2.48	2.59
2.5	2.27	2.24	0.21	0.57	2.47	2.81
2.7	2.29		0.08		2.37	
2.9	2.30		−0.03		2.27	
3.1	2.32		−0.09		2.23	
3.3	2.33		−0.16		2.17	

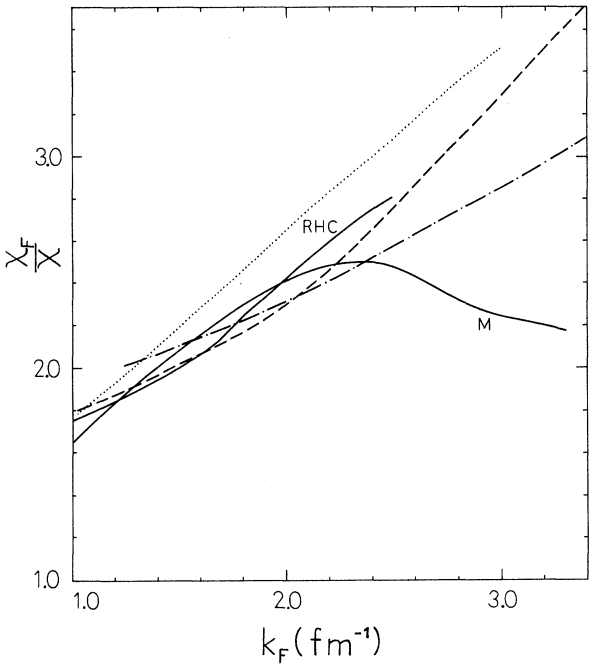

FIG. 1. Values of χ_F/χ versus k_F for the RHC and M potentials (solid lines). The results of Pfarr (Ref. 12) (dotted line), Pandharipande *et al.* (Ref. 11) (dashed line), and Moszkowski (Ref. 14) (dash-dotted line, corresponding to the SW effective interaction) for RSC potential are shown for comparison.

should be attributed to different approximations used by these authors in their calculations of the effective neutron-neutron interaction in neutron matter.

In Pfarr's paper[11] the energy of neutron matter was calculated in the Hartree-Fock approximation with an effective neutron-neutron interaction obtained by the method of the unitary transformation from the RSC potential. This approach is in the first order equivalent to the Jastrow method but avoids complications of the normalization. The parameters of the transformation were not calculated self-consistently but were determined phenomenologically from the condition of the best agreement with the saturation properties of symmetric nuclear matter. The resulting effective nucleon-nucleon interaction was extrapolated to the case of the neutron matter. In view of all these approximations, the approximate agreement of Pfarr's values of χ_F/χ with those obtained in Refs. 12 and 14 should be considered as very good.

In Ref. 11 the energy of the polarized neutron matter has been computed in Jastrow variational theory, using the lowest-order cluster expansion. Subsidiary healing conditions were imposed on the Jastrow correlation functions to simulate the exclusion principle effect. Many approximations were introduced to simplify the calculational procedure. E.g., for the sake of simplicity the spin-orbit and tensor forces were neglected, and the central part of the $^3P_2$-$^3F_2$ RSC potential was assumed in all odd l states.

The effective neutron-neutron interaction, used by Moszkowski,[14] has been developed by Sawada and Wong in their nuclear matter calculation, and utilized for several values of the density and neutron excess. The main simplification in Sawada and Wong calculations concerned the choice of the single-particle energies, which is not justified from the point of view of the complete many-body calculation. Generally, the single-particle energies used by Sawada and Wong, and in consequence their effective interaction (SW), should be treated rather as an "educated guess" for the form of the effective interaction in neutron matter, especially at high densities. On the other hand, the SW effective interaction was determined for unpolarized nuclear matter and the use of its matrix elements for the completely polarized ($\alpha=1$) neutron matter may be criticized.

Our results obtained for the M potential are very characteristic. They agree with our RHC and the quoted RSC results below and near normal nuclear density ($k_F<2$ fm^{-1}). At higher densities the values of χ_F/χ for the M potential are substantially lower than those obtained for the RHC and RSC potentials. Contrary to the monotonic increase in χ_F/χ for the RHC and RSC potentials, for the M potential χ_F/χ decreases for $k_F>2.5$ fm^{-1}. Generally, the differences between the results for the RHC and M potentials are much more pronounced in the case of χ_F/χ than the differences between the values of the energy per neutron, calculated in Ref. 23. Thus, in our case the calculation of the magnetic susceptibility of neutron matter represents a very sensitive test for the assumed form of the nuclear interaction.

In principle, the magnetic susceptibility could be determined from the difference between the values of the energy per neutron in neutron matter with $\alpha=0$ and $\alpha=1$, assuming that $E^n_{\rm pot}$ is a linear function of α^2. This is a very good approximation.[11,12] However, the generalization of the standard calculation in lowest order Brueckner theory to the case of completely polarized neutron matter is not trivial. In the case of completely polarized neutron matter the model single-particle potential has the form[24]

$$V(\vec{m}_0)=V_0(m_0)+V_2(m_0)P_2(\hat{m}_0)\ ,$$

where $\vec{m}_0$ is the neutron momentum and $P_2(\hat{m}_0)$ is the Legendre polynomial of the cosine of the angle between $\vec{m}_0$ and the direction of the spin polarization. The tensor part of V does not enter directly into the potential energy of the system. However, it should be taken into account when considering the dispersive effects of the medium. This introduces some new difficulties to the self-consistent calculation.

Summarizing, one may conclude, that the results for the RHC potential, obtained within the frame of the Brueckner theory, are quite similar to those obtained for the RSC potential using different formalisms.[11,12,14] The results obtained for the nonlocal M potential, especially the density dependence of the magnetic susceptibility of neutron matter, are substantially different from those obtained for the modern local neutron-neutron potentials. No evidence of a transition to ferromagnetic state is found in the density region considered.

Note added in proof: The values of χ_F/χ for the RSC potential have been obtained just recently by the author, using the methods and approximations of the present paper. These results, which are in agreement with those obtained within the frame of the Landau theory of Fermi liquids in Ref. 13, will be published elsewhere.

The author is indebted to Professor J. Dabrowski for directing his attention to this problem and for many stimulating and helpful discussions.

[1]A. Hewish, Annu. Rev. Astron. Astrophys. 8, 265 (1970).
[2]H. A. Bethe, Annu. Rev. Nucl. Sci. 21, 93 (1971).
[3]V. R. Pandharipande and V. K. Garde, Phys. Lett. 33B, 608 (1972).
[4]D. H. Brownell, Jr. and J. Callaway, Nuovo Cimento 60B, 169 (1969).
[5]M. J. Rice, Phys. Lett. 29A, 637 (1969).
[6]S. D. Silverstein, Phys. Rev. Lett. 23, 139 (1969).
[7]E. Østgaard, Nucl. Phys. A154, 202 (1970).
[8]J. W. Clark and N.-C. Chao, Nuovo Cimento Lett. 2, 185 (1969).
[9]J. W. Clark, Phys. Rev. Lett. 23, 1463 (1969).
[10]J. M. Pearson and G. Saunier, Phys. Rev. Lett. 24, 325 (1970).
[11]V. R. Pandharipande, V. K. Garde, and J. K. Srivastava, Phys. Lett. 38B, 485 (1972).
[12]J. Pfarr, Z. Phys. 251, 152 (1972).
[13]S.-O. Bäckman, C.-G. Källman, and O. Sjöberg, Phys. Lett. 43B, 263 (1973).
[14]S. A. Moszkowski, Phys. Rev. C 9, 1613 (1974).
[15]R. V. Reid, Ann. Phys. (N.Y.) 50, 411 (1968).
[16]J. L. Gammel, R. S. Christian, and R. M. Thaler, Phys. Rev. 105, 311 (1957).
[17]J. Dabrowski and P. Haensel, Phys. Rev. C 7, 916 (1973).
[18]R. Mongan, Phys. Rev. 178, 1537 (1969).
[19]K. A. Brueckner, S. A. Coon, and J. Dabrowski, Phys. Rev. 168, 1184 (1968).
[20]P. J. Siemens, Nucl. Phys. A141, 225 (1970).
[21]O. Sjöberg, Nucl. Phys. A222, 161 (1974).
[22]K. A. Brueckner and J. L. Gammel, Phys. Rev. 109, 1040 (1958).
[23]P. Haensel, Nucl. Phys. (to be published).
[24]J. Dabrowski and P. Haensel, Can. J. Phys. 52, 1768 (1974).