IC/79/133
INTERNAL REPORT
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19001

International Atomic Energy Agency
and
United Nations Educational Scientific and Cultural Organization

INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

ON THE THERMAL PROPERTIES OF POLARIZED NUCLEAR MATTER *

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ABSTRACT

The thermal properties of polarized nuclear matter are calculated using Skyrme III interaction modified by Dabrowski for polarized nuclear matter. The temperature dependence of the volume, isospin, spin and spin isospin pressure and energies are determined. The temperature, isospin, spin and spin isospin dependence of the equilibrium Fermi momentum is also discussed.

MIRAMARE - TRIESTE
August 1979

* To be submitted for publication.

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I. INTRODUCTION

Recently, much interest has been devoted to the study of the thermal properties of nuclear matter. Stocker and Burzlaff (1973) have presented an extended Fermi-gas model to describe volume, surface and asymmetry contributions to the nuclear excitation energy. They treated the potential part of the nuclear energy within Brueckner's energy density formalism and they included corrections due to surface and asymmetry effects. Kupper et al. (1974) have calculated the thermostatic properties of symmetric nuclear matter by extension of Thomas-Fermi approach to ground state nuclei by Myers and Swiatecki (1969). They have computed the free energy per nucleon $f(T, n)$ in Landau's quasiparticle approximation and derived the relevant thermostatic properties. Saur et al. (1976) have investigated the thermal properties of finite nuclei by applying thermal Hartree-Fock approximation and using Skyrme (1956, 1959) interaction. They have determined the surface effects by separating the bulk properties of nuclear matter, obtained from an equation of state, from that for finite nuclei. They gave explicit numerical expressions for temperature dependence of various quantities.

The ground-state energy of polarized nuclear matter was considered by Dabrowski and Haensel (1972a, 1972b) and Dabrowski (1976), using the K-matrix method and applying the Brueckner-Gammel-Thaler, Hamada-Johnson and Reid-soft-core nucleon-nucleon potentials. Dabrowski (1977) has analyzed the problem of spin instability of nuclear matter with the Skyrme interaction. The three-body part of the Skyrme interaction was replaced by a density dependent two-body force, and an explicit dependence of this force on the four densities of neutrons and protons with spin up and down was introduced .

In this work we generalize Dabrowski's (1977) calculations to finite temperatures.

The derivation of the equation of state for the case of polarized nuclear matter is given in section II. Section II contains also the derivation of the dependence of the equilibrium Fermi momentum on isospin, spin, spin-isospin and temperature. Section III contains the results and discussion.

II. THEORY

II.1 The Equation of State

We follow in this derivation a similar one introduced by Stocker (1973). The entropy per nucleon in Landau's quasi-particle approximation (Kupper et al.

1974) is given by

$$s = \frac{S}{A} = -n^{-1}\sum_{\tau} \frac{1}{(2\pi)^3}\int d\vec{k}\,\big(n_\tau(k)\ln n_\tau(k) + (1-n_\tau(k))\ln(1-n_\tau(k))\big). \tag{1}$$

where τ runs over spin and isotopic spin n_τ denotes the Fermi distribution of nuclear matter in equilibrium at temperature T,

$$n_\tau(k) = \left(1 + \exp\left(\frac{\hbar^2 k^2}{2M} + v_\tau(k) - \mu_\tau\right)/T\right)^{-1}, \tag{2}$$

n is the density of nuclear matter,

$$n = \frac{A}{V} = \sum_\tau \frac{1}{(2\pi)^3}\int d\vec{k}\; n_\tau(k) . \tag{3}$$

The Boltzmann constant is put equal to 1.

The single-particle potential, $v_\tau(k)$, in the effective mass approximation can be written as:

$$v_\tau(k) = v_\tau^{0} + k^2 v_\tau^{1} . \tag{4}$$

Introducing (3) and (2) into (1) we can perform the integration over $d\vec{k}$ as shown in Ref. [2] and the result is

$$s = \frac{\pi^2}{2} n^{-1} \sum_\tau n_\tau \frac{2M_\tau^*}{\hbar^2 k_F^{\tau 2}} T + O(T^3) , \tag{5}$$

$$n_\tau = \frac{k_F^{\tau 3}}{6\pi^2}\left(1 + \frac{\pi^2}{8}\left(\frac{2M_\tau^* T}{\hbar^2 k_F^{\tau 2}}\right)^2\right) , \tag{6}$$

$$n = \sum n_\tau \tag{7}$$

then s, to first order in T, is

$$s = \frac{\pi^2}{4}\frac{M^* T}{\hbar^2 k_F^3}(\varkappa_n + \lambda_n + \varkappa_p + \lambda_p) , \tag{8}$$

where m^* denotes the effective mass (M_τ^*) at the Fermi surface and is assumed to be the same for all spin and isotopic spin states, $\varkappa_n$, $\varkappa_p$, λ_n and λ_p are the Fermi momenta (k_F^τ) of neutron (proton) with spin-up and neutron (proton) with spin-down,

$$\left.\begin{matrix}\varkappa_n^3\\ \lambda_n^3\end{matrix}\right\} = k_F^3\,(1 + \alpha_\tau \pm \alpha_\sigma \pm \alpha_{\sigma\tau}) . \tag{9}$$

$$\left.\begin{matrix}\varkappa_p^3\\ \lambda_p^3\end{matrix}\right\} = k_F^3\,(1 - \alpha_\tau \pm \alpha_\sigma \mp \alpha_{\sigma\tau}) . \tag{10}$$

$$\alpha_\tau = (N-Z)/A , \quad \alpha_n = (N\uparrow - N\downarrow)/A ,$$

$$\alpha_p = (Z\uparrow - Z\downarrow)/A , \quad \alpha_\sigma = \alpha_n + \alpha_p ,$$

$$\alpha_{\sigma\tau} = \alpha_n - \alpha_p \tag{11}$$

$N\uparrow(N\downarrow)$ denotes neutrons with spin-up (down) and $Z\uparrow$ ($Z\downarrow$) denotes protons with spin-up (down). Knowing the entorpy s we can determine the specific heat c_v from the relation [2],

$$c_v = n T \,\partial_T^{n_\tau} s . \tag{12}$$

Equations (7), (11) give

$$c_v = \frac{M^* T}{6\hbar^2}(\varkappa_n + \lambda_n + \varkappa_p + \lambda_p) . \tag{13}$$

We can use expression (13) for c_v to determine the integral energy per particle through the relation

$$e(T,n) = e(T=0,n) + \frac{1}{n}\int_0^T dT' C_v(T') , \quad (14)$$

where $e(T = 0 , n)$ is calculated using Skyrme III interaction (Beiner et al. 1975) modified by Dabrowski (1977). The modification consists in replacing the three body term by a density dependent two body interaction of the form:

$$V_{12}^{(3)} = \frac{1}{6} t_3 (1+x_3 P^{\sigma}) \Big[\rho + (\rho_{n\uparrow}+\rho_{n\downarrow}-\rho_{p\uparrow}-\rho_{p\downarrow}) T_3$$

$$+ (\rho_{n\uparrow}-\rho_{n\downarrow}+\rho_{p\uparrow}-\rho_{p\downarrow}) S_3 + (\rho_{n\uparrow}-\rho_{n\downarrow}-\rho_{p\downarrow}+\rho_{p\uparrow}) Y \Big], \quad (15)$$

where the densities for $\rho_{n\uparrow}$, $\rho_{n\downarrow}$, $\rho_{p\uparrow}$, $\rho_{p\downarrow}$ correspond to $N\uparrow$ neutrons with spin up, $N\downarrow$ neutrons with spin down, $Z\uparrow$ protons with spin up, $Z\downarrow$ protons with spin down,

$$T_3 = (\vec{\tau}_1 + \vec{\tau}_2)_3/2 , \quad S_3 = (\vec{\sigma}_1 + \vec{\sigma}_2)_3/2 ,$$

and

$$Y = (\tau_{13}\sigma_{13} + \tau_{23}\sigma_{23})/2 . \quad (16)$$

In the case of polarized nuclear matter the entropy, specific heat, integral energy can be expanded up to second order in α_τ, α_σ and $\alpha_{\sigma\tau}$ to have the form

$$s = s_v + \frac{1}{2}(\alpha_\tau^2 s_\tau + \alpha_\sigma^2 s_\sigma + \alpha_{\sigma\tau}^2 s_{\sigma\tau}) \quad (17)$$

and similar relations for c_v and e where

$$s_v = \pi^2 \frac{M^*}{\hbar^2 k_F^2} T , \quad (18a)$$

$$\bar{s}_i = -\frac{2}{9}\pi^2 \frac{M^*}{\hbar^2 k_F^2} T , \quad i = \tau, \sigma, \sigma\tau . \quad (18b)$$

The free energy per particle is given by

$$f(T,n) = e(T,n) - T s(T,n) . \quad (19)$$

The pressure p is given by

$$p = n^2 \left(\frac{\partial f}{\partial n}\right)_T . \quad (20)$$

We can make a similar expansion of s to p to get

$$p = p_v + \frac{1}{2}(\alpha_\tau^2 P_\tau + \alpha_\sigma^2 P_\sigma + \alpha_{\sigma\tau}^2 P_{\sigma\tau}) . \quad (21)$$

where for Skyrme III interaction we get

$$P_v = \frac{1}{5}\frac{\hbar^2 k_F^2}{M} n + \frac{3}{8} t_0 n^2 + \frac{1}{8} t_3 n^3 + \frac{1}{16}(3t_1 + 5t_2) n^2 k_F^2$$

$$+ \frac{2}{9}\frac{M}{\hbar^2} k_F T^2 \quad (22)$$

and

$$P_i = \frac{2}{9}\frac{\hbar^2 k_F^2}{M} n + \frac{5}{9} t_2 k_F^2 n^2 - \frac{4}{81}\frac{M}{\hbar^2} k_F T^2$$

$$+ \begin{cases} -\frac{1}{4} t_0 (1+2x_0) n^2 - \frac{1}{4} t_3 n^3 & \\ -\frac{1}{4} t_0 (1-2x_0) n^2 + \frac{5}{12} t_3 n^3 & \text{for} \\ -\frac{1}{4} t_0 n^2 + \frac{1}{12} t_3 n^3 & \end{cases} \begin{cases} i = \tau \\ i = \sigma \\ i = \sigma\tau \end{cases} \quad (23)$$

II.2 The Temperature and the Spin-Isospin Dependence of the Equilibrium Density of Nuclear Matter

The equilibrium Fermi momentum was discussed by Weiss and Cameron (1969), Bethe (1971) and by Dworzecka (1966) for zero temperature nuclear matter with neutron excess. We will now generalize this discussion to the case of non-zero temperature and polarized nuclear matter. The equilibrium Fermi momentum k_F for fixed values of α_τ , α_σ , $\alpha_{\sigma\tau}$ and T is treated as a solution of the equation

$$\frac{\partial f(\alpha_\tau,\alpha_\sigma,\alpha_{\sigma\tau},T,k_F)}{\partial k_F} = 0 \ . \qquad (24)$$

We assume the knowledge of the solution to Eq. (24) in the case $\alpha_\tau = \alpha_\sigma = \alpha_{\sigma\tau} = 0$ and $T = 0$ and we denote this solution as k_{F0} , where k_{F0} satisfies the condition:

$$\left.\frac{\partial f(\alpha_\tau=\alpha_\sigma=\alpha_{\sigma\tau}=0, T=0, k_F)}{\partial k_F}\right|_{k_{F0}} = 0 \ . \qquad (25)$$

Then we expand the derivative in Eq. (24) in power series of k_F about k_{F0} neglecting terms of higher order than first. We obtain:

$$\frac{\partial f(\alpha_\tau,\alpha_\sigma,\alpha_{\sigma\tau},T,k_F)}{\partial k_F} = \left.\frac{\partial f(\alpha_\tau,\alpha_\sigma,\alpha_{\sigma\tau},T,k_F)}{\partial k_F}\right|_{k_{F0}}$$

$$+ \Delta k_F \left.\frac{\partial^2 f(\alpha_\tau,\alpha_\sigma,\alpha_{\sigma\tau},T,k_F)}{\partial k_F^2}\right|_{k_{F0}} = 0 \ . \qquad (26)$$

To solve this equation we have to expand the derivatives in power series of α_τ , α_σ , $\alpha_{\sigma\tau}$ and T . Neglecting terms of higher order than α_τ^2 , α_σ^2 , $\alpha_{\sigma\tau}^2$ and T^2 and making use of equations 13, 14, 17 we get:

$$\Delta k_F = -\frac{k_F^2}{K}\left[\frac{1}{2}\left(\alpha_\tau^2 \frac{\partial e_\tau(0,n_0)}{\partial k_F} + \alpha_\sigma^2 \frac{\partial e_\sigma(0,n_0)}{\partial k_F} + \alpha_{\sigma\tau}^2 \frac{\partial e_{\sigma\tau}(0,n_0)}{\partial k_F}\right) + \frac{\pi^2 M^*}{\hbar^2 k_F^3} T^2\right] , \qquad (27)$$

where K is the compressibility and n_0 is the equilibrium density.

The numerical results and the discussion will be found in the next section.

III. RESULTS AND DISCUSSION

The parameters of the Skyrme force employed in the calculations are (Beiner et al. 1975)

$t_0 = -1128.75$ MeV Fm^3 , $t_1 = 395.0$ MeV Fm^3 ,
$t_2 = -95.0$ MeV Fm^5 , $t_3 = 14000.0$ MeV Fm^6 ,
$x_0 = 0.45$

This force leads to saturation of nuclear matter at $k_F = 1.29$ Fm^{-1} . The effective mass is considered to be equal to the free particle mass. This choice of the effective mass is justified by the fact that the effective mass at the Fermi surface in nuclear matter may well be approximated by the free nucleon mass. This is in agreement with the results of investigations of heavy nuclei (Brown 1971).

Fig. (1) shows the (p_v , n) diagrams at different temperatures and they reproduce almost the same features as the corresponding one obtained by Sauer et al. The symmetry pressures P_τ , P_σ and $P_{\sigma\tau}$ depend very slightly on temperature and therefore they are reproduced in Fig. (2) only for $T = 0$. The value of P_τ is always negative, however, the values of P_σ and $P_{\sigma\tau}$ are always positive which indicates that they have opposite effects on the total pressure. To the best of our knowledge these curves have not been reproduced before.

The equilibrium volume, isospin, spin and spin-isospin energies can be written in the form

$$\epsilon_i(T) = \epsilon_i(0) + \gamma_i T^2 \quad , i = v, \tau, \sigma, \sigma\tau \ , \qquad (28)$$

where

$\epsilon_V = -15.9$ MeV , $\epsilon_\tau = 56$ MeV , $\epsilon_\sigma = 81$ MeV , $\epsilon_{\sigma\tau} = 69$ MEV

$$\gamma_V = \frac{\pi^2}{2} \frac{M^*}{\hbar^2 k_F^2} = 0.07 \ \mathrm{MeV}^{-1} \left(\text{for } k_F = 1.29\,\mathrm{F}^{-1}\right). \quad (29)$$

$$\gamma_i = -\frac{2}{9}\gamma_V = -.016 \ \mathrm{MeV}^{-1} \ , \text{for } i = \tau, \sigma, \sigma\tau . \quad (30)$$

The values of γ_V and γ_τ obtained by Sauer et al. are $\gamma_V = .055\ \mathrm{MeV}^{-1}$, $\gamma_\tau = .012\ \mathrm{MeV}^{-1}$.

The comparison between our results and those of Sauer et al. shows that a better agreement between them can be achieved if we insert for the effective mass $m^* = 0.77$ m . This will account for our approximation of taking the entropy as that of a free particle entropy. However, the value of γ_τ obtained by Sauer et al. has the opposite sign to the value we have obtained.

Equation (25) for Δk_F gives

$$k_F = k_{F0}\left(1 - .03\,\alpha_\tau^2 - 0.54\,\alpha_\sigma^2 - 0.28\,\alpha_{\sigma\tau}^2 - 4.01\times10^{-4}\,T^2\right). \quad (31)$$

This relation gives us the dependence of the equilibrium Fermi momentum on α_τ , α_σ , $\alpha_{\sigma\tau}$ and T for the parameters of Skyrme III interaction. The density can easily be calculated, and we get

$$\rho = \rho_0\left(1 - .09\,\alpha_\tau^2 - 1.62\,\alpha_\sigma^2 - .84\,\alpha_{\sigma\tau}^2 - .12\times10^{-2}\,T^2\right). \quad (32)$$

The coefficient of T^2 agrees very well with the value $.126 \times 10^{-2}$ obtained before by Sauer et al. (1976). The coefficient of α_τ^2 is much smaller than the value obtained before by Dworzecka (1966) namely -0.49 and that of Wiess and Cameron (1969) namely -0.45. The value of this parameter depends on the value of $\frac{d\epsilon_\tau(0,n_0)}{dk_F}$ which was calculated by Farine et al. (1978) ($\equiv$ 2L in the notation of Farine et al.) for a number of different interactions. It was found that it changes much from one set of Skyrme interaction to another.

We may conclude that our calculations indicate the dependence of the thermal properties of nuclear matter with an excess of neutrons, of spin-up neutrons and of spin-up protons, and they are reliable at low temperatures. However, at higher temperatures we have to use the exact form of Fermi integrals and take into account the dependence of Skyrme parameters on density which depends on temperature (Kupper et al. 1974).

ACKNOWLEDGMENTS

One of the authors (M.Y.M.H.) would like to thank Professor Abdus Salam, the International Atomic Energy Agency and UNESCO for hospitality at the International Centre for Theoretical Physics, Trieste.

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Figure 1

Nuclear matter isotherm in the (n, Pv) diagram.

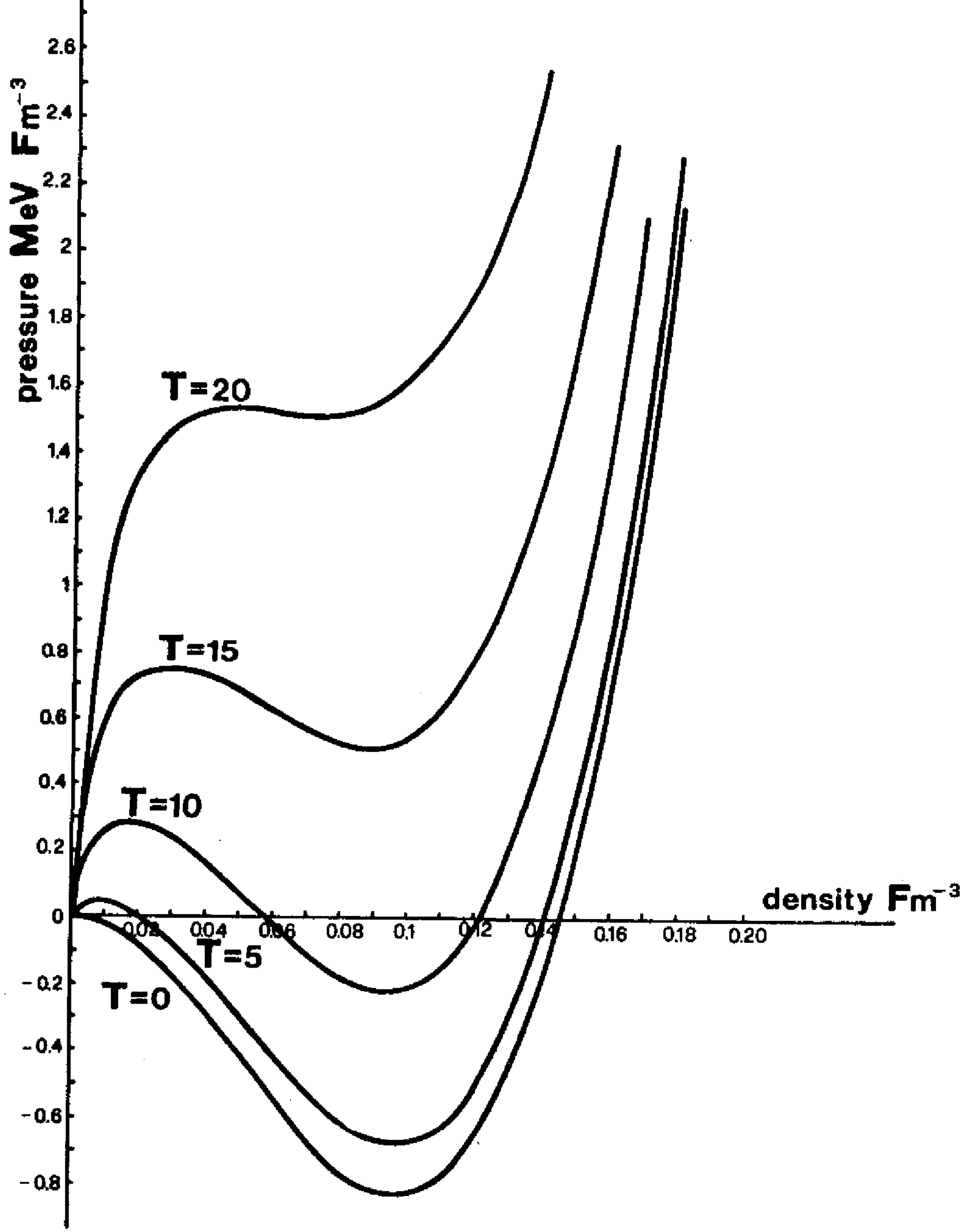
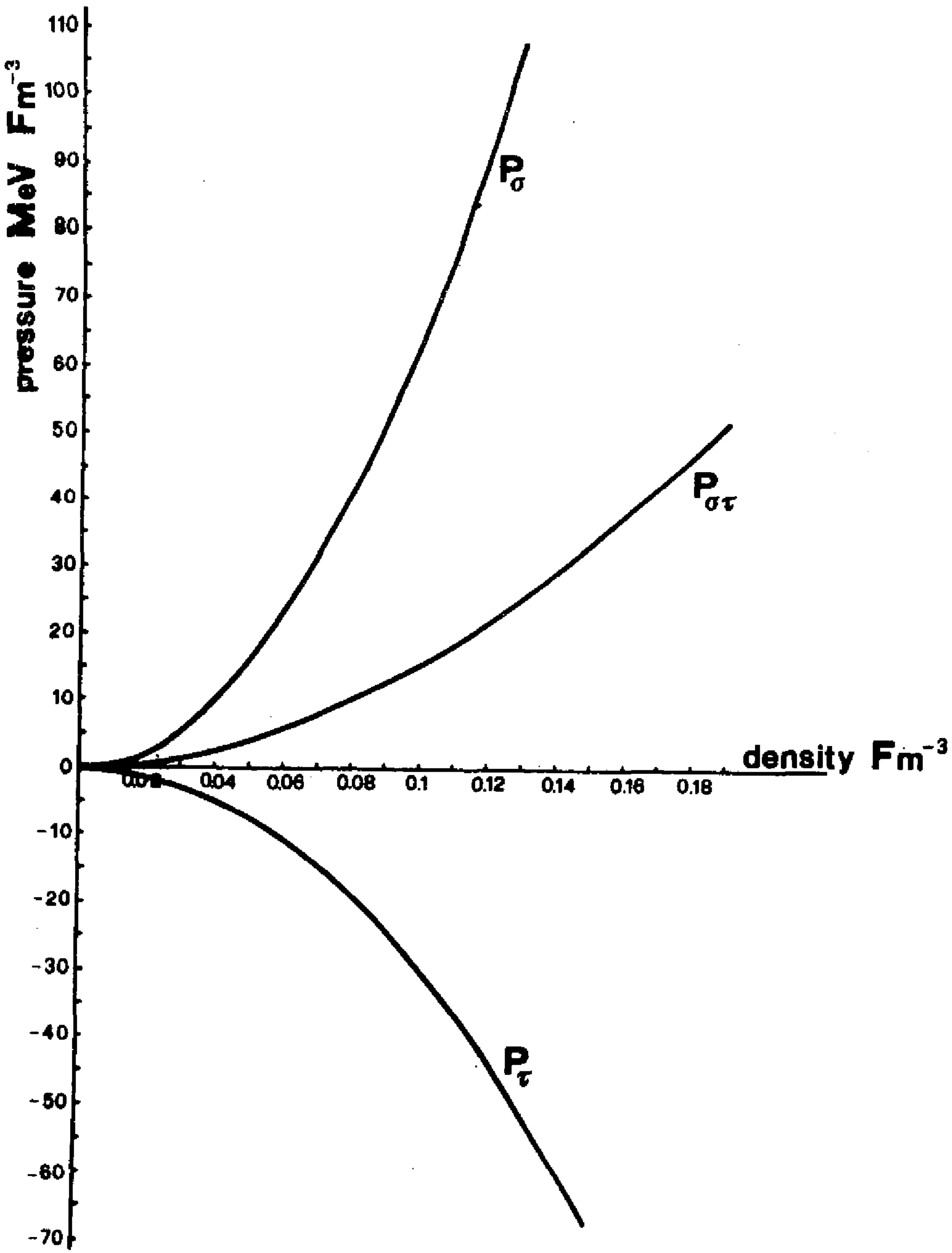


Figure 2

Nuclear matter isotherms in the (n, P_i), $i = \tau, \sigma, \sigma\tau$.
diagram



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