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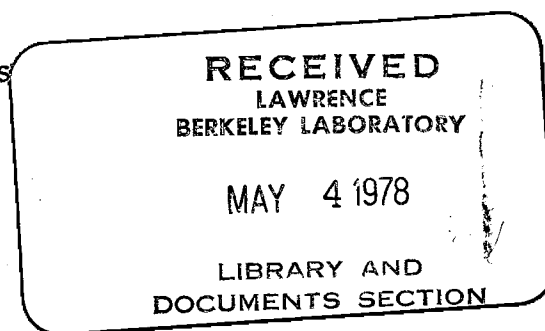
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D. Vautherin and M. Gyulassy

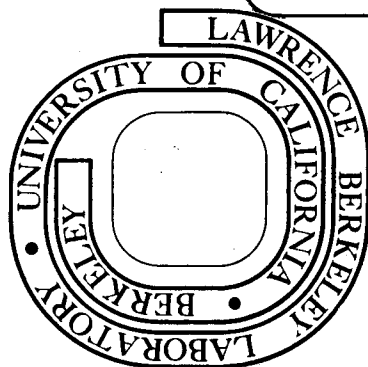
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NUCLEAR MATTER, PCAC, and PION CONDENSATION⁺D. Vautherin^{*} and M. Gyulassy

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ABSTRACT

An effective Lagrangian is constructed with strong chiral symmetry breaking terms that leads to both normal nuclear saturation as well as PCAC and current algebra. We then show that pion condensation is consistent with both these physical constraints.

⁺This work was done with support from Division of Nuclear Physics of U.S. Department of Energy.

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To investigate possible exotic properties of nuclear systems at high densities, many groups¹⁻⁴ have considered Lorentz invariant Lagrangian theories in the mean field approximation. The existing theories can be grouped into one of two broad classes:

1. Theories^{1,3} with parameters constrained to reproduce normal nuclear saturation,

$$E/A = W(\rho_0) = -16 \text{ MeV} \quad (1)$$

at $\rho_0 = 0.17 \text{ fm}^{-3}$.

2. Theories^{2,4} constrained to satisfy current algebra and Partially Conserved Axial Currents (PCAC),

$$\partial^\mu A_\mu^i(x) = c_\pi m_\pi^2 \pi_i(x) \quad (2)$$

with $c_\pi = 94 \text{ MeV}$ and π_i being the pion field. Having imposed the physical constraint of nuclear physics, Eq. (1), or particle physics, Eq. (2), the objective of both classes of theories is to compute the equation of state $W(\rho)$ of nuclear matter at high densities $\rho > \rho_0$. The resulting $W(\rho)$ is hoped to shed light on the possibility of phase transitions in dense nuclear systems. Particular interest has focused on phase transitions characterized by σ and/or π condensation.

A disturbing feature of both classes of theories is that the two constraints, Eqs. (1,2), are not satisfied simultaneously by any of the existing models. The class of theories^{1,3} that have good nuclear saturation properties do not satisfy PCAC and, hence, ignore many cherished properties of strong interactions such as the Goldberger-Trieman and the Adler-Weisberger relations. On the other hand, the

class of theories^{2,4} motivated by PCAC and current algebra do not lead to normal nuclear saturation and, hence, ignore nuclear physics. The point of this paper is to show that a simple Lagrangian can be constructed that incorporates both physical constraints, Eqs. (1,2). With that Lagrangian we then investigate the possibility of pion condensation. We find that pion condensation survives both constraints.

Our approach is based on the following point of view. We want to construct an effective relativistic Lagrangian $\mathcal{L}^{\text{eff}}$, that yields Eq. (1) in the mean field approximation (MFA) as in Refs. (1,3), while at the same time yields Eq. (2) as in Ref. (2,4). We then use $\mathcal{L}^{\text{eff}}$ as an extrapolation tool to explore properties of nuclear systems at densities away from $\rho = \rho_0$.

The $\mathcal{L}^{\text{eff}}$ is used here in the mean field approximation in the same spirit as pseudo-potentials are used in the Born approximation elsewhere in nuclear physics. We are then not explaining either eq. (1) or eq. (2) but merely demand that our extrapolation Lagrangian incorporates them. Imposing these constraints on $\mathcal{L}^{\text{eff}}$ is designed to increase our confidence in extrapolations away from $\rho = \rho_0$.

Certainly, in the future, additional constraints stemming from details of πN and $\pi\pi$ interactions must also be incorporated in $\mathcal{L}^{\text{eff}}$. However, for now we impose only eq. (1,2) with the idea that it is only worthwhile adding further constraints if predictions of exotic phenomena survive at least these two important constraints.

We consider the following $\mathcal{L}^{\text{eff}}$ involving the nucleon ψ , sigma σ , pion π_i and omega ω_μ fields:

$$\begin{aligned}
 \mathcal{L}(x) = & \bar{\psi} \left\{ (i\partial_\mu - g_v \omega_\mu) \gamma^\mu - g_s (\sigma + i\gamma_5 \tilde{\tau} \cdot \tilde{\pi}) \right\} \psi \\
 & + \frac{1}{2} \left\{ (\partial_\mu \sigma)^2 + (\partial_\mu \pi)^2 \right\} \\
 & - \frac{1}{2} \left\{ \frac{1}{2} F_{\mu\nu} F^{\mu\nu} + m_v^2 \omega_\mu^2 \right\} \\
 & - U(\sigma, \pi_i)
 \end{aligned} \tag{3}$$

where the potential density U is parameterized as

$$\begin{aligned}
 U(\sigma, \pi_i) = & \frac{a}{2} \left[\left(\frac{\sigma}{\sigma_0} \right)^2 + \left(\frac{\pi}{\sigma_0} \right)^2 \right] + \frac{b}{4} \left[\left(\frac{\sigma}{\sigma_0} \right)^2 + \left(\frac{\pi}{\sigma_0} \right)^2 \right]^2 \\
 & - (a+b) \frac{\sigma}{\sigma_0} + \frac{a}{2}(c-1) \left(\frac{\pi}{\sigma_0} \right)^2 + \frac{a}{2} + \frac{3}{4}b
 \end{aligned} \tag{4}$$

The vacuum is then specified by $\langle \sigma \rangle = \sigma_0$ and $\langle \pi_i \rangle = 0$ at which point

$$U = \partial U / \partial \sigma = \partial U / \partial \pi = 0$$

The divergence of the axial current $A_\mu^i(x)$ is computed from the variation of $\mathcal{L}$ under a chiral transformation:

$$(\sigma \rightarrow \sigma + \tilde{\epsilon} \cdot \tilde{\pi}, \quad \tilde{\pi} \rightarrow \tilde{\pi} - \tilde{\epsilon} \sigma, \quad \psi \rightarrow (1 + i\gamma_5 \tilde{\epsilon} \cdot \tilde{\tau}) \psi),$$

$$\partial^\mu A_\mu^i = - \frac{\partial U}{\partial \sigma} \pi_i + \frac{\partial U}{\partial \pi_i} \sigma \tag{5}$$

Equation (5) shows that $\partial^\mu A_\mu^i = 0$ if $\mathcal{L}$ were chiral symmetric, i.e.

$U = U(\sigma^2 + \pi^2)$. Note that the form of the current $A_\mu^i = \delta \mathcal{L} / \delta (\partial_\mu \epsilon_i)$

does not depend on U since U involves no derivatives. Hence, current

algebra is not effected by the specific form of U . However, PCAC requires

$\partial^\mu A_\mu^i$ to be given by Eq. (2). The PCAC constraint on U is then

$$\left[-\frac{\partial U}{\partial \sigma} \pi_i + \frac{\partial U}{\partial \pi_i} \sigma \right]_{\sigma=\sigma_0} = c_\pi m_\pi^2 \pi_i, \quad (6)$$

Observe that Eq. (6) is only a local constraint on U in the (σ, π) plane near the vacuum point $\sigma = \sigma_0$, $\pi_i = 0$. This is a reminder of the well known (but sometimes forgotten) fact that while approximate chiral symmetry is sufficient for PCAC, it is by no means necessary. Note that PCAC has only been tested in conditions where $\sigma \approx \sigma_0$.

The existence of models with strongly broken chiral symmetry satisfying Eq. (2) is the key to reconciling constraints Eqs. (1) and (2). The point is that Eq. (1) also requires^{1,3} only a limited knowledge of $U(\sigma, \pi)$. The addition of any function of π^2 to U has no effect on the calculation leading to Eq. (1). Since Eq. (1) only constrains $U(\sigma, 0)$ while Eq. (2) only constrains the derivatives of U near the vacuum, it is not surprising that a form of $U(\sigma, \pi)$ can be found that leads to both Eqs. (1) and (2).

To illustrate the differences between models satisfying Eq. (1) and/or Eq. (2), we plot in Fig. 1 equipotential surfaces of $U(\sigma, \pi)$ in three models. Figure 1(a) corresponds to the σ model^{2,4} with $a = -507 \text{ MeV/fm}^3$, $b = 530 \text{ MeV/fm}^3$, $c = 1$ in Eq. (4). The approximate chiral invariance is reflected in the approximate rotation symmetry of U about $\sigma = \pi = 0$. As discussed in Refs. (2,3), normal nuclei (Eq. (1)) are however unstable in this model. In Fig. 1(b) a Walecka type model¹ is illustrated for which $a = 379 \text{ MeV/fm}^3$, $b = 0$, $c = 1$. In this case Eq. (1) is satisfied, but $\partial^\mu A_\mu^i = (a/\sigma_0)\pi_i$. To recover Eq. (2) requires

$\sigma_0 = 1580$ MeV, which then implies $m_\sigma = m_\pi = 34$ MeV! If, on the other hand, m_π is taken as the physical pion mass then $\partial^\mu A_\mu^i = 4 c_\pi m_\pi^2 \pi_i$ which is inconsistent with Goldberger-Trieman relation, pion decay width, etc. In this model chiral symmetry is clearly broken.

Finally, Fig. 1(c) illustrates a $U(\sigma, \pi)$ satisfying both Eqs. (1) and (2). This model⁵ differs from Walecka's model only in that $1/c = 16$ which leads to $m_\sigma/m_\pi = 4$ and $\sigma_0 = c_\pi$.

Comparing Figs. 1(a), 1(c) near the vacuum point, one can see how the local property, Eq. (6), can be simultaneously satisfied by both models. Comparing Figs. 1(b), 1(c) shows that $U(\sigma, 0)$ is identical and, hence, both models lead to the same condition Eq. (1). The interesting point to observe is that Eq. (1) was made consistent with Eq. (2) simply by retaining both of the commonly employed versions of symmetry breaking⁴: $\mathcal{L}_{SB}^{(1)} = c_1 \sigma$ and $\mathcal{L}_{SB}^{(2)} = c_2 \pi^2$. In this model then, chiral symmetry is strongly broken on a global scale in the (σ, π) plane.

We have then constructed one $\mathcal{L}^{eff}$ leading to both Eqs. (1,2) by dispensing with approximate chiral invariance. While $\mathcal{L}^{eff}$ does not have then the aesthetic appeal of chiral models, we emphasize again that $\mathcal{L}^{eff}$ is only an extrapolation tool and not a theory in itself. Furthermore we note that $\mathcal{L}^{eff}$ is of course not uniquely determined by Eqs. (1,2) alone. There must be a large class of $\mathcal{L}^{eff}$ that satisfy Eqs. (1,2). In particular, it is possible that an $\mathcal{L}^{eff}$ can be constructed retaining approximate chiral invariance that is also consistent with both Eqs. (1,2). Such an approach (most likely including $\Delta_{33}(1236)$ and ρ fields) will probably be necessary in order to incorporate the correct S-wave πN and $\pi\pi$ scattering in the future.

However, for now we will investigate neutral pion condensation using $\mathcal{L}^{\text{eff}}$ in Eq. (3) with $a = 379 \text{ MeV/fm}^3$, $b = 0$, $c = 1/16$, and $g_s^2 = 100$ and $g_v^2/m_v^2 = 195.7/m_N^2$. While the free space S-wave πN interaction is not given correctly with Eq. (3), in isospin symmetric nuclear matter it vanishes to lowest order and therefore does not concern us here. Also because we will consider only small amplitude pion fields, the $\pi\pi$ interaction can safely be neglected. Of course, the strong P-wave πN interaction is correctly given in Eq. (3).

With our $\mathcal{L}^{\text{eff}}$, the ground state energy of symmetric nuclear matter in MFA is

$$E = \int d^3r \left\{ \sum_{i=1}^A \bar{\psi}_i(r) [-i\vec{\gamma} \cdot \vec{\nabla} + g_s(\sigma + i\gamma_5 \vec{\tau} \cdot \vec{\pi}) + g_v \gamma_0 V(r)] \psi_i(r) + \frac{1}{2} [(\nabla\sigma)^2 + (\nabla\pi)^2 - (\nabla V)^2] - \frac{1}{2} m_v^2 V^2 + U(\sigma, \vec{\pi}) \right\} \quad (7)$$

where A denotes the nucleon number, ψ_i the normalized single-particle nucleon states, and V the time-component of the vector meson field.

Since we are interested in a static solution we have assumed¹ that

$$\omega_\mu = V\delta_{\mu 0} \quad [1].$$

The ground-state values of σ , $\vec{\pi}$, V and the Dirac spinors ψ_i are determined by requiring that the ground-state energy E is stationary with respect to small variations of these variables i.e.

$$\delta E/\delta\sigma = \delta E/\delta\vec{\pi} = \delta E/\delta V = \delta E/\delta\psi_i = 0 \quad (8)$$

In infinite nuclear matter there exists a solution of Eq. (8) such that

$$\vec{\nabla}_\sigma = \vec{\nabla}_V = \vec{\pi} = 0 \quad (9)$$

In this case the single-particle wave functions ψ_i are plane waves

$$\psi(k,s;r) = \omega(k,s) \exp(i\vec{k} \cdot \vec{r}) / \sqrt{\Omega} \quad (10)$$

where Ω is the total volume. The index $s = 1, 2$ corresponds to the spin degree of freedom and the four-component spinor ω satisfies

$$(\vec{\alpha} \cdot \vec{k} + \beta m^*) \omega(\vec{k}, s) = \varepsilon(\vec{k}) \omega(\vec{k}, s) \quad , \quad (11)$$

where the effective nucleon mass is

$$m^* = g_s \sigma \quad (12)$$

and $\varepsilon(k) = \sqrt{k^2 + m^{*2}} - g_v V$. Thus

$$\omega(\vec{k}, s) = \sqrt{\frac{\varepsilon(k) + m^*}{2\varepsilon(k)}} \begin{pmatrix} \chi(s) \\ \frac{\vec{\sigma} \cdot \vec{k}}{\varepsilon(k) + m^*} \chi(s) \end{pmatrix} \quad , \quad (13)$$

where $\chi(s)$ are the usual Pauli spinors.

With Eq. (10), minimizing E with respect to V gives⁽¹⁾

$$E(k_F)/\Omega = 4 \int \frac{d^3 p}{(2\pi)^3} \sqrt{p^2 + m^{*2}} \varepsilon(k_F - |\vec{p}|) + U(\sigma = \frac{m^*}{g_s} \quad , \quad \pi = 0) + \frac{1}{2} g_v^2 \rho^2 / m_v^2 \quad , \quad (14)$$

where $\rho = 2k_F^3 / (3\pi^2)$ and m^* is obtained¹ by minimizing $E(k_F)$ in Eq. (14) with respect to σ .

To investigate whether the homogeneous solution (9) in infinite nuclear matter is stable against neutral pion-condensation, we consider a small external pion field

$$\pi_0(\vec{r}) = a \cos \vec{q} \cdot \vec{r} \quad \pi_+(\vec{r}) = \pi_-(\vec{r}) = 0 \quad (15)$$

In order to check that Eq. (9) corresponds to a minimum, we calculate the change in the total energy E to second order in a . We will keep the values of σ and V fixed (which is valid to second order in a) and look only for the response of the nucleons to the external field Eq. (15). The relation between the present method and the method using the pion propagator is discussed in the article by Campbell et al [4].

The change ΔE_M in the meson energy is

$$\Delta E_M = \int d^3r \left\{ \frac{1}{2} (\vec{\nabla} \pi_0)^2 + \frac{1}{2} m_\pi^2 \pi_0^2 \right\} = \frac{\Omega}{4} a^2 (q^2 + m_\pi^2) \quad (16)$$

Assuming $a \ll m^*/g_s$, the change ΔE_F in the energy of the fermions can be evaluated by second-order perturbation theory. From Eq. (7), the unperturbed Hamiltonian is $H_0 = \vec{\alpha} \cdot \vec{p} + \beta m^* + g_v V$, and the perturbation is

$$W = ig_s (a \cos \vec{q} \cdot \vec{r}) \gamma_0 \gamma_5 \tau_z \quad (17)$$

Since the first order term vanishes, we have

$$\Delta E_F = \sum_{n \neq 0} \frac{|\langle 0 | W | n \rangle|^2}{E_n - E_0} \quad (18)$$

where $|0\rangle$ is the Slater determinant describing the ground state and $|n\rangle$ an excited state. Because W is a one-body operator $|n\rangle$ must be of the form of a one particle-one hole state

$$|n\rangle = a^\dagger(k', s', \tau') a(k, s, \tau) |0\rangle \quad (19)$$

where s, τ are spin and isospin indices, and where a and $a^\dagger$ are destruction and creation operators, and where (k, s, τ) corresponds to an occupied level while (k', s', τ') corresponds to an empty state. Explicitly

$$\Delta E_F = \sum_{s, s'} \sum_{\vec{k}} g_s^2 a^2 \frac{|\omega^\dagger(\vec{k}, s) \gamma_0 \gamma_5 \omega(\vec{k}+\vec{q}, s')|^2}{\epsilon(\vec{k}) - \epsilon(\vec{k}+\vec{q})} \quad (20)$$

where the $\vec{k}$ summation is restricted to

$$|\vec{k}| \leq k_F \quad \text{and} \quad |\vec{k}+\vec{q}| \geq k_F \quad (21)$$

Notice that we have excluded the sum over negative energy states in Eq. (20). This (divergent) sum is a contribution to the vacuum fluctuation energy.

The matrix elements involved in Eq. (20) are

$$\omega^\dagger(\vec{k}, s) \gamma_0 \gamma_5 \omega(\vec{k}', s') = N(k) N(k') (s | \vec{\sigma} | s') \cdot \left(\frac{\vec{k}'}{\epsilon(k') + m^*} - \frac{\vec{k}}{\epsilon(k) + m^*} \right) \quad (22)$$

where $N(k) = ((\epsilon(k) + m^*)/2\epsilon(k))^{1/2}$.

For simplicity we use the non-relativistic limit of Eq. (22). As will be seen below this approximation is justified for the range of densities which will be of interest to us. Assuming $\vec{q}$ to be parallel to the z -axis that matrix element reduces to

$$|\omega^\dagger(\vec{k}, s) \gamma_0 \gamma_5 \omega(\vec{k}+\vec{q}, s')| \approx \delta_{ss'} \frac{|\vec{q}|}{2m^*} \quad (23)$$

The expression for ΔE_F therefore reduces to

$$\frac{\Delta E_F}{A} = - \frac{3g_s^2 a^2}{32m^*} \phi(x) \quad , \quad (24)$$

where A is the nucleon number, $x = q/k_F$ and

$$\phi(x) = x^2 + x \left(1 - \frac{x^2}{4} \right) \log \left| \frac{2+x}{2-x} \right| \quad (25)$$

is the familiar Lindhard function⁵. By comparing Eq. (24) with the expression for the change in the meson-energy

$$\frac{\Delta E_M}{A} = \frac{3\pi^2 a^2}{8k_F} \left(x^2 + \frac{m_\pi^2}{k_F^2} \right) \quad , \quad (26)$$

it is clear that at low density, nuclear matter will be stable against neutral pion condensation. On the other hand for some critical value k_c of the Fermi momentum an instability will appear. The critical value k_c is given by

$$k_c = \inf_x \left\{ k_0(x) \right\} \quad , \quad (27)$$

where k_0 denotes, for a given x , the solution of the equation

$$\Delta E_M + \Delta E_F = 0 \quad , \quad \text{i.e.}$$

$$\frac{1}{k_0} \left(x^2 + \frac{m_\pi^2}{k_0^2} \right) = g_s^2 \frac{\phi(x)}{4\pi^2 m^*(k_0)} \quad (28)$$

Notice that the effective nucleon mass m^* is a function of the Fermi momentum through Eqs. (12,14). This function can be shown to decrease continuously from the nucleon mass m to zero. As a consequence, Eq. (28) has one real solution only.

To solve the previous set of equations we have first constructed a table of values of the function $m^*(k_F)$ by minimizing numerically the total energy, Eq. (14). We have used the same value of g_s, g_v as in Walecka's model [1] in order to have the same saturation properties.

We have then solved Eq. (28) using the iteration procedure

$$k_0 = 4m^*(k_0) \frac{\pi^2}{g_s} \frac{1}{\Phi(x)} \left(x^2 + \frac{m_\pi^2}{k_0^2} \right) \quad (29)$$

for various values of x . Using $g_s = 10$, $m_\pi = 0.7 \text{ fm}^{-1}$, the critical value of k_F is found to be

$$k_c = 1.08 \text{ fm}^{-1} \quad (30)$$

for which $m^*(k_c) = 0.79 m$. This value is obtained for $x = 1.30$. The momentum of the condensate is therefore

$$q = 1.40 \text{ fm}^{-1} \quad (31)$$

Since in Walecka's model saturation is obtained at $k_F = 1.42 \text{ fm}^{-1}$ i.e. $\rho_0 = 0.193 \text{ nucleons/fm}^3$ the critical density is

$$\rho_c = 0.44 \rho_0 \quad (32)$$

At this density nucleons in nuclear matter are definitely non-relativistic which justifies a posteriori, the non-relativistic approximation which was used above.

The values of ρ_c and q we have derived, should still be corrected for the inclusion of the Δ_{33} resonance as well as short range correlations. Our results are similar to those of non-relativistic models [6]. They are also similar to the results obtained by Dautry [7] in the sigma model using

the Thomas-Fermi approximation, and by Celenza and Pirner [7] who use a field theoretic model reproducing low energy πN data.

The inclusion of $\pi N \Delta$ as well as NN and $N \Delta$ correlations⁷, is expected to increase ρ_c beyond ρ_0 . Therefore it will most likely be necessary to look at high energy nuclear collision for evidence of pion condensation.⁸

In conclusion we have been able to construct a Lagrangian which is consistent with both saturation properties of nuclear matter and particle properties summarized by current algebra and PCAC. It does not predict any transition towards an abnormal nuclear matter phase at high density [4] since it has the same saturation properties as Walecka's model [1]. On the other hand it still predicts a possible pionic instability in nuclear matter.

Stimulating discussions with F. Dautry, R. E. Peierls, H. J. Pirner, M. Rho, and D. Campbell are gratefully acknowledged.

References

- 1) J. D. Walecka, Ann. Phys. 83, 491 (1974); S. A. Chin and J. D. Walecka, Phys. Letters 52B, 24 (1974).
- 2) A. K. Kerman and L. D. Miller, Contribution to Heavy Ion Summer Study, LBL, July 1974, MIT-CTP Pub. 449.
- 3) S. A. Moszkowski and C. G. Källman, Nucl. Phys. A287, 495 (1977).
J. Boguta and A. R. Bodmer, Nucl. Phys. A292, 413 (1977).
See E. B. Nyman, M. Rho, Nucl. Phys. A268, 408 (1976) for discussion of quantum corrections to the sigma model.
- 4) T. D. Lee, Rev. Mod. Phys. 47, 267 (1975).
D. Campbell, R. Dashen and J. Mannassah, Phys. Rev. D12, 979, 1010 (1975).
M. Chanowitz and P. J. Siemens, Phys. Lett. 70B, 175 (1977).
- 5) A. L. Fetter and J. D. Walecka, Quantum Theory of Many-particle Systems, McGraw-Hill, New York, 1971.
- 6) A. B. Migdal, The Pion Field in Nuclear Matter, Rev. Mod. Phys., to be published.
W. Weise and G. E. Brown, Phys. Lett. 27C, 1 (1976).
- 7) F. Dautry, Chiral Symmetry and Neutral Pion Condensation, Les Houches Summer School 1977.
L. Celenza and H. J. Pirner, Nucl. Phys. A294, 375 (1978).
- 8) M. Gyulassy, W. Greiner, Ann. Phys. (NY) 109, 485 (1977).

Figure Captions

Fig. 1 Contour plot of $U(\sigma/\sigma_0, \pi/\sigma_0)$, Eq. (4), in steps of 20 MeV/fm^3 between $U = 10$ and 230 MeV/fm^3 .

- a) Model satisfying Eq. (2) but not Eq. (1);
- b) Model satisfying Eq. (1) but not Eq. (2);
- c) Model satisfying both Eqs. (1) and (2). See discussion in text.

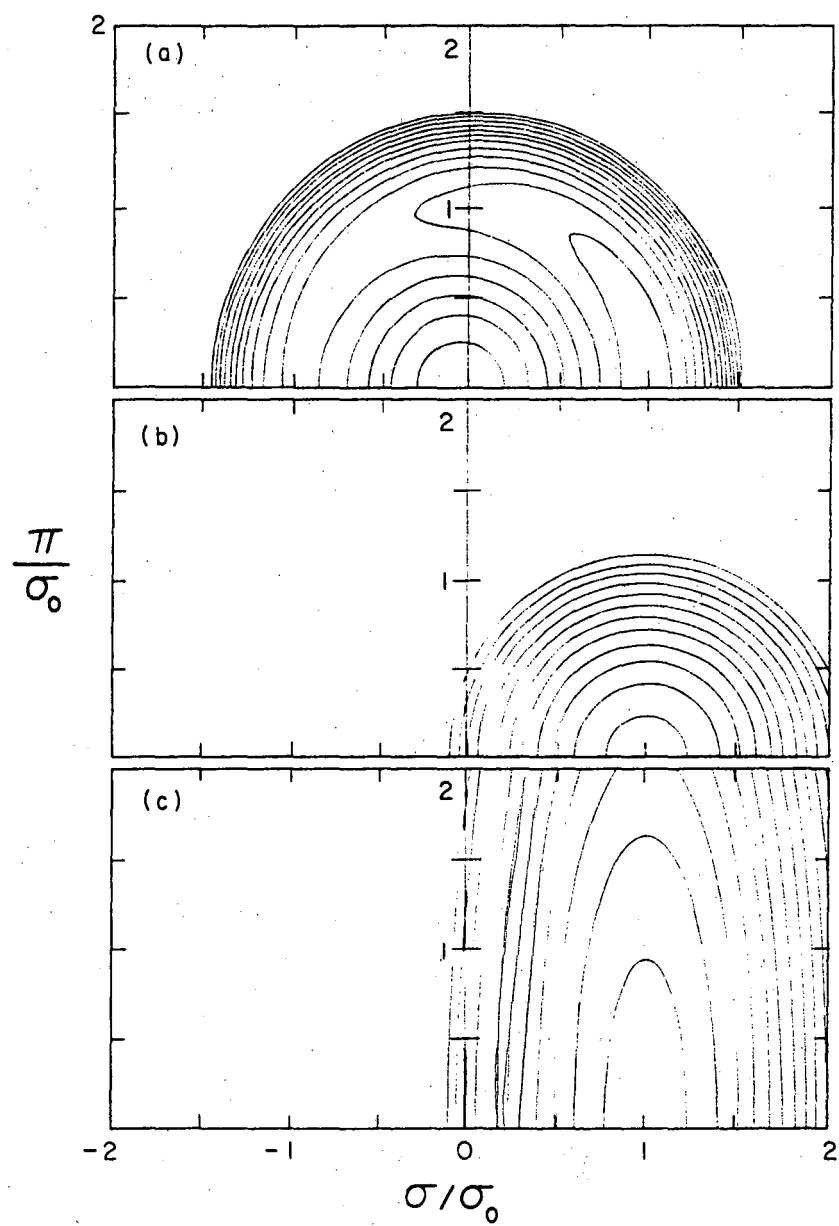


Fig. 1

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