A Nonlinear Realized Approach of SU(2) Chiral Symmetry Spontaneous Breaking and Properties of Nuclear Matter

Xiao-fu Lü\textsuperscript{1,2}, Bao-xi Sun\textsuperscript{1,4}, Yu-xin Liu\textsuperscript{3,1,5}\textsuperscript{*}, Hua Guo\textsuperscript{3}, and En-guang Zhao\textsuperscript{1,5}

\textsuperscript{1}Institute of Theoretical Physics, The Chinese Academy of Sciences, P. O. Box 2735, Beijing 100080, China
\textsuperscript{2}Department of Physics, Sichuan University, Chengdu 610064, China
\textsuperscript{3}School of Physics, Peking University, Beijing 100871, China
\textsuperscript{4}Institute of High Energy Physics, The Chinese Academy of Sciences, P.O. Box 918(4), Beijing 100039, China
\textsuperscript{5}Center of Theoretical Nuclear Physics, National Laboratory of Heavy ion Accelerator, Lanzhou 730000, China

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Abstract

A nonlinear realization of SU(2) chiral symmetry spontaneous breaking approach is developed in the composite operator formalism. A Lagrangian including quark, gluon and Goldstone boson degrees of freedom of the chiral quark model is obtained.

*Corresponding author
from the QCD Lagrangian. A way to link the chiral symmetry spontaneous breaking formalism at hadron level and that at quark level is predicted. too. The application to nuclear matter shows that the approach is quite successful in describing the properties of nuclear matter and the quark condensate in it.

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It has been known that chiral symmetry and its spontaneous breaking is the key points in understanding many features of the nature, such as the generation of some particles in strong interaction physics[1], the chiral property of many amino acids[2], and so on. Moreover, it has been widely used in guiding the design and synthesis of chemical compounds to control and improve their functions[3]. In low energy strong interaction physics, Goldstone bosons appear as the chiral symmetry is spontaneously broken. To carry out the constraint by the appearance of Goldstone bosons, several realization formalisms of the chiral symmetry spontaneous breaking ($\chi$SB) have been developed[1, 4, 5]. Furthermore, Considering the effective degrees of freedom to describing hadron structure, we know that several formalisms have been developed. For example, the one with constituent quarks and Goldstone bosons has been shown to be successful in describing some properties of hadrons[3, 6, 8] and hadron spectroscopy[4], even nucleon-nucleon interactions[10]. The one including constituent quarks and gluons are very powerful to describe hadron properties[11, 12, 13], too. Moreover, Manohar and Georgi proposed that, besides the constituent quarks, both the Goldstone bosons and gluons are necessary to describe hadron structure[14]. Such a scheme has also been widely used to investigate hadron spectroscopy and hadron interactions(see, for example, Refs.[15, 16]). However, which effective degrees of freedom have more sophisticated QCD foundation is still a significant topic in hadron physics and low energy QCD. In the sprit of $\chi$SB and the composite operator scheme[17] of the QCD, we propose, in this letter, a new approach to realize the $\chi$SB nonlinearly not only in formalism but also in practical calculation and obtain an effective Lagrangian including constituent quarks, Goldstone bosons and gluons. Then some light is shed on the effective degrees of freedom of hadron structure. We propose also a link between the formalism of $\chi$SB at hadron level and that at quark level. As an application, we take the approach to describe the properties of nuclear matter. By implementing the Hellmann-Feynman theorem[18], we evaluate the quark condensate in nuclear matter in the present formalism, too.
At first, we discuss the linear infinitesimal transformation in general

\[ \phi'_n(x) = \phi_n + i \epsilon \sum_m t^\alpha_{nm} \phi_m(x) \]  

(1)

where \( t^\alpha \) is a generator of the symmetry group of the Lagrangian and \( \phi_n \) is a spin-zero boson field or spin-zero composite operators. The quantum effective action is defined by

\[ \Gamma[\phi] \equiv - \int d^3x \sum_n \phi^n J_{\phi_n}(x) + W[J] \]  

(2)

This quantum effective action has the same symmetric property as the Lagrangian. It gives that

\[ \sum_{n,m} \int \frac{\delta \Gamma[\phi]}{\delta \phi_n} t^\alpha_{nm} \phi_m(x) d^4x = 0 \]  

(3)

Introducing an effective potential \( V(\phi) \) as \( \Gamma[\phi] = -vV[\phi] \), where \( v \) is the space-time volume, one can rewrite Eq.(3) as

\[ \sum_{n,m} \frac{\partial V(\phi)}{\partial \phi_n} t^\alpha_{nm} \phi_m = 0, \]  

(4)

where \( \phi \) is independent of \( x \). Taking the differentiation with respect to \( \phi_l \), we obtain

\[ \sum_{n} \frac{\partial V(\phi)}{\partial \phi_n} t^\alpha_{nl} + \sum_{n,m} \frac{\partial^2 V(\phi)}{\partial \phi_n \partial \phi_l} t^\alpha_{mn} \phi_m = 0 \]  

(5)

The symmetry spontaneous breaking appears when \( \frac{\partial V(\phi)}{\partial \phi_m} = 0 \) and \( \phi_m \neq 0 \) where \( \phi_m = \langle 0 | \phi_m | 0 \rangle \). In accordance with the condition of the symmetry spontaneous breaking, Eq.(5) can be written as

\[ \sum_{n,m} \left( \frac{\partial^2 V(\phi)}{\partial \phi_n \partial \phi_l} \right)_{\phi=m} t^\alpha_{nm} \phi_m = 0. \]  

(6)

It is apparent that, if the symmetry is broken, \( \sum_m t^\alpha_{nm} \phi_m \) should not identically equate to zero. The massless eigenvectors of the mass matrix \( \frac{\partial^2 V(\phi)}{\partial \phi_n \partial \phi_l} \) span then a linear space in which the Goldstone bosons lie and the remainder has to be perpendicular to it. This constraint can lead various forms of the realization of the \( \chi_{SB} SU_L(2) \times SU_R(2) \supset SU(2) \). In this paper we shall discuss the nonlinear realization of \( SU_L(2) \times SU_R(2) \supset SU(2) \) on the quark level at first. Then we propose a link between the \( \chi_{SB} \) at quark level and that at hadron level.
Constructing four composite operators in terms of the current quark fields $u$ and $d$ as

$$\psi_4 = \bar{q} q, \quad \psi_i = i \bar{q} \gamma^5 \tau^i q$$

where $\tau^i$ is the Pauli matrix, we can verify that the transformation of $\psi_n (n=1,2,3,4)$ under the chiral symmetry transformation of the quarks $e^{i \gamma_5 \epsilon^i \tau^i}$ can be written as

$$\delta \psi_4 = 2 \epsilon^i \psi_i, \quad \delta \psi_i = -2 \epsilon^i \psi_4$$

where $\epsilon^i$ are the infinitesimal parameters.

Analyses in QCD sum rules\cite{19} and calculations in lattice QCD\cite{20} have shown that the vacuum expectation value of quarks (quark condensate) is not zero, i.e., $\langle 0 | \bar{q}(x)q(x) | 0 \rangle \neq 0$. It indicates that the chiral symmetry $SU_L(2) \times SU_R(2)$ is spontaneously broken and chiral condensate appears. As a consequence, the $\psi_n$ has to be separated into two parts: one contains Goldstone bosons and the other does not. This can be done by writing

$$\psi_n = \sum_m (e^{2i \xi^i T^i})_{nm} \bar{\psi}_m$$

where $T^{\alpha\beta}$ are the generators of the four dimensional rotation group. To make the $\bar{\psi}_n$ do not contain Goldstone bosons, the $\bar{\psi}_n$ can be taken as $\bar{\psi}_n = (\bar{\psi}_4, 0, 0, 0)$. In the infinitesimal form, we have then

$$\delta \psi_i = 2 \xi^i \bar{\psi}_4$$

Comparing Eq.(10) with Eq.(8), we obtain the relation between the constituent quarks and the current quarks as

$$\bar{q} = e^{-i \gamma_5 \xi^i \tau^i} q$$

In order to simplify the factor $e^{-i \gamma_5 \xi^i \tau^i}$ we introduce the Goldstone bosons as

$$\xi^i = (\tan^{-1} |\vec{\zeta}|) \frac{\vec{\xi}^i}{|\vec{\zeta}|}$$

Using Eq.(12), we can obtain the nonlinear realization of the $\chi\text{SB}$ as

$$e^{-i \gamma_5 (\tan^{-1} |\vec{\zeta}|) \vec{\xi} \cdot \vec{\zeta}/|\vec{\zeta}|} = \frac{1 - i \gamma_5 \vec{\tau} \cdot \vec{\zeta}}{\sqrt{1 + \vec{\zeta}^2}}.$$
Since the theory leaves the SU(2) invariant, the \( \tilde{q} \) constructs a representation of SU(2).

Then it is necessary to have
\[
e^{-i\gamma_5 \vec{r} \cdot \vec{\tau}} e^{-i\gamma_5 (\tan^{-1} |\vec{r}|) \vec{r} \cdot \vec{r} / |\vec{r}|} e^{-i\gamma_5 (\tan^{-1} |\vec{r}|) \vec{r} \cdot \vec{r} / |\vec{r}|} e^{-i\vec{r} \cdot \vec{\theta}}
\]

After some algebraic calculation we obtain the relation between the Goldstone bosons and the constituent quarks as
\[
\delta \tilde{q} = -i \vec{\tau} \cdot (\vec{\epsilon} \times \vec{\zeta}),
\]
\[
\delta \vec{\zeta} = \vec{\epsilon}(1 - \vec{\zeta}^2) + 2\vec{\zeta}(\vec{\epsilon} \cdot \vec{\zeta}),
\]
\[
\vec{\theta} = \vec{\epsilon} \times \vec{\zeta},
\]
\[
(15a)
\]
\[
(15b)
\]
\[
(15c)
\]

Under these transformations, the QCD Lagrangian can be effectively expressed in terms of the constituent quark, gluon and Goldstone bosons as
\[
\begin{align*}
\mathcal{L}_{\text{eff}} &= \sum \bar{\tilde{q}} \gamma^\mu \left( \partial_\mu + i \frac{\vec{\zeta} \cdot \vec{\tau}}{1 + \vec{\zeta}^2} - \frac{i\gamma_5 \vec{\tau} \cdot \partial_\mu \vec{\zeta}}{1 + \vec{\zeta}^2} \right) \tilde{q} + i\bar{\tilde{q}} \gamma^\mu \left( -ig^a \frac{\lambda^a}{2} A^a_{\mu \rho} \right) \tilde{q} \\
&- m\bar{\tilde{q}} - \frac{1}{4} G^a_{\mu \rho} G^{a \mu \rho} + \mathcal{L}_M.
\end{align*}
\]

where \( m \) is the constituent quark mass resulting from quark condensates. \( \mathcal{L}_M \) is the Lagrangian of the Goldstone boson fields \( \vec{\zeta} \) and the \( \tilde{\psi}_4 \). In this sense, the effective degrees of freedom to describe hadron structure should involve constituent quarks, gluons and Goldstone bosons. One can then recognize that the constituent quark model of hadron structure involving constituent quarks, Goldstone bosons and gluons\(^{[14]}\) has a quite solid QCD foundation. Furthermore, the double counting and the spurious state involved in Manohar-Georgi’s formalism\(^{[14]}\) disappears in the present approach.

\(^{[14]}\)

From the effective action \( \Gamma \) in Eq.(2) we have
\[
\mathcal{L}_M = -\frac{1}{2A} \partial_\mu \tilde{\psi}_4 \partial^\mu \tilde{\psi}_4 - \frac{2}{A} \tilde{\psi}_4 \tilde{\psi}_4 \tilde{D}_\mu \tilde{D}^\mu + \mathcal{L}(\tilde{\psi}_4)
\]

where \( D_\mu^i = \frac{\partial_\mu \sigma^i}{1 + \tau_2} \). \( A \) is a dimensional constant. Considering the pseudoscalar field property of \( \zeta_i \) and the assumption of the saturation of vacuum, we obtain that the \( \mathcal{L}_M \) can be written as the superposition of the following two parts
\[
\mathcal{L}_\sigma = -\frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - \frac{1}{2} m^2 \tau_2 \sigma^2 - \frac{g_2}{3} \sigma^3 - \frac{g_3}{4} \sigma^4,
\]
\[
(17)
\]
\[ L_\pi = \frac{1}{2} \frac{\partial^\mu \pi}{1 + \frac{\pi^2}{f^2}} \cdot \frac{\partial^\nu \pi}{1 + \frac{\pi^2}{f^2}} - 2 f_\pi \sigma \frac{\partial^\mu \pi \cdot \partial^\nu \pi}{(1 + \frac{\pi^2}{f^2})^2} - 2 \sigma^2 \frac{\partial^\mu \pi \cdot \partial^\nu \pi}{(1 + \frac{\pi^2}{f^2})^2}. \] 

(19)

And the parameters \( g_2 \) and \( g_3 \) hold relation \( g_2 = \frac{3m_\pi \sqrt{A}}{m_\tau^2} g_3 \) with \( m_\tau \) being the current quark mass.

It has been well known that there exists \( \chi \)SB formalism at hadron level. However the relation between the formalism and that at quark level is not clear enough. To link the \( \chi \)SB formalism at hadron level with the \( \chi \)SB formalism at quark level described above, we construct a set of composite operators which have the quantum numbers of proton and neutron as

\[
p(x) = \varepsilon_{abc} \left( u^{aT}(x), d^{aT}(x) \right) \gamma^\mu u^b(x) \gamma_5 \gamma_\mu i \tau_2 \begin{pmatrix} u^c(x) \\ d^c(x) \end{pmatrix},
\]

\[
n(x) = \varepsilon_{abc} \left( u^{aT}(x), d^{aT}(x) \right) \gamma^\mu d^b(x) \gamma_5 \gamma_\mu i \tau_2 \begin{pmatrix} u^c(x) \\ d^c(x) \end{pmatrix},
\]

(20a)

(20b)

where \( a, b, c \) are the color indices, \( \tau_2 \) acts only on the isospin indices and the row and column matrix express the doublet of the isospin. Since the nucleon in this doublet are composed of current quarks, we can refer them as current nucleon. Implementing the chiral transformation of quarks \( \tilde{q} = \exp(-i \gamma_5 \xi \tau_i)q \), we have the transition among nucleons as

\[
\tilde{N} = e^{-i \gamma_5 \xi \tau_i} N,
\]

(21)

where \( N = \begin{pmatrix} p(x) \\ n(x) \end{pmatrix} \), \( \tilde{N} = \begin{pmatrix} \tilde{p}(x) \\ \tilde{n}(x) \end{pmatrix} \) in which

\[
\tilde{p}(x) = \varepsilon_{abc} \left( \tilde{u}^{aT}(x), \tilde{d}^{aT}(x) \right) \gamma^\mu \tilde{u}^b(x) \gamma_5 \gamma_\mu i \tau_2 \begin{pmatrix} \tilde{u}^c(x) \\ \tilde{d}^c(x) \end{pmatrix},
\]

\[
\tilde{n}(x) = \varepsilon_{abc} \left( \tilde{u}^{aT}(x), \tilde{d}^{aT}(x) \right) \gamma^\mu \tilde{d}^b(x) \gamma_5 \gamma_\mu i \tau_2 \begin{pmatrix} \tilde{u}^c(x) \\ \tilde{d}^c(x) \end{pmatrix},
\]

(22a)

(22b)

The Eq. (21) shows that we can obtain the nucleon from the current nucleon. The nucleon observed in experiments in the low energy region can be formed only through Eq. (22).
Taking the same measure as described for the quarks and considering the Goldberg-Treiman relation and the partial conservation of the axial current (PCAC), we obtain the Lagrangian density with the nonlinear realization of the SU(2) $\chi$SB at hadron level as

$$
\mathcal{L}_{ch} = \bar{\psi} \left( i\gamma_\mu \partial^\mu - M_N - \frac{\bar{\tau} \cdot (\bar{\pi} \times \bar{\phi})}{f_\pi^2 (1 + \frac{\bar{\pi}^2}{f_\pi^2})} + \frac{2iM_N g_A \gamma_5 \bar{\tau} \cdot \bar{\pi}}{f_\pi (1 + \frac{\bar{\pi}^2}{f_\pi^2})} - \frac{2g_A \gamma_5 \bar{\tau} \cdot \bar{\pi} \bar{\phi}}{f_\pi^2 (1 + \frac{\bar{\pi}^2}{f_\pi^2})^2} \right) \psi
$$

$$
-\sigma^2 \bar{\psi} \left( \frac{8i g' M_N \gamma_5 \bar{\tau} \cdot \bar{\pi}}{f_\pi (1 + \frac{\bar{\pi}^2}{f_\pi^2})} + \frac{8g' \gamma_5 \bar{\tau} \cdot \bar{\pi} \bar{\phi}}{f_\pi^3 (1 + \frac{\bar{\pi}^2}{f_\pi^2})^2} \right) \psi - \frac{\delta M_p + \delta M_n}{2} \left( \frac{1 - \frac{\bar{\pi}^2}{f_\pi^2}}{1 + \frac{\bar{\pi}^2}{f_\pi^2}} \right) \bar{\psi} \psi - g_\sigma \bar{\psi} \sigma \psi
$$

$$
- \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - U(\sigma) - \left( \frac{1}{2} + 2f_\pi \sigma + 2\sigma^2 \right) \frac{\partial_\mu \bar{\tau} \cdot \partial^\mu \bar{\pi}}{\bar{\pi} \bar{\pi}} - \frac{1}{2} m_\pi^2 \bar{\pi}^2,
$$

where $M_N$ is the mass of the nucleon caused by the $\chi$SB, $\delta M_p$ and $\delta M_n$ are the masses of the current proton and neutron, respectively. The $U(\sigma)$ stands for

$$
U(\sigma) = \frac{1}{2} m_\sigma^2 \sigma^2 + \frac{1}{3} g_2 \sigma^3 + \frac{1}{4} g_3 \sigma^4,
$$

where the $g_2$ and $g_3$ have relation $g_2 = -\frac{3 \delta M_N \sqrt{A}}{m_\pi^2} g_3$ with $\delta M_N = \delta M_n = \delta M_p$.

To represent the repulsive interaction among nucleons and the isospin symmetry breaking in nuclear matter, we can include $\omega$ and $\rho$ mesons in the way having been widely used [21]. Then in the mean field approximation, we have the Lagrangian

$$
\mathcal{L}_{RMF} = \bar{\psi} \left( i\gamma_\mu \partial^\mu - M_N \right) \psi - \delta M_N \left( \frac{1 - \frac{\bar{\pi}^2}{f_\pi^2}}{1 + \frac{\bar{\pi}^2}{f_\pi^2}} \right) \bar{\psi} \psi
$$

$$
- g_\sigma \bar{\psi} \psi \sigma_0 - g_\omega \bar{\psi} \gamma^0 \psi \omega_0 - g_\rho \bar{\psi} \gamma^0 t_3 \rho_{03} \psi
$$

$$
- \frac{1}{2} m_\sigma^2 \sigma_0^2 - \frac{g_2}{3} m_\pi^2 \sigma_0^3 - \frac{g_3}{4} \sigma_0^4 - \frac{1}{2} m_\pi^2 \omega_0^2 + \frac{1}{2} m_\rho \rho_{03}^2,
$$

where the $\sigma_0$ is the expectation value of the isoscalar scalar field, $\pi_0^2$ is that of $\bar{\pi}^2$, $\omega_0$ is that of the temporal component for the isoscalar vector field since there is no spatial direction for a uniform nuclear matter at rest, $\rho_{03}$ is that of $\rho$ meson in the nuclear matter.

By fitting the saturation properties of nuclear matter, we obtain the parameters with the restriction $M_N + \delta M_N = 938$ MeV. Two sets of the parameters are listed in Table
1. For the parameter set $C_1$, we get the saturation density of $0.152 \text{fm}^{-3}$, binding energy of $15.297 \text{MeV}$, a compression modulus of $349.10 \text{MeV}$, symmetry energy coefficient $33.645 \text{MeV}$ and the effective mass of a nucleon of $0.687M_N$ for symmetric nuclear matter. The parameter set $C_2$ gives a saturation density $0.151 \text{fm}^{-3}$, a binding energy of nucleon $15.040 \text{MeV}$, a compression modulus $K = 298.88 \text{MeV}$, a symmetry energy coefficient $32.593 \text{MeV}$ and an effective mass of nucleon $0.736M_N$ for symmetric nuclear matter. Meanwhile the curves for the equation of states (EOS) are obtained too. The numerical results show that the EOSs for the two sets of parameters are quite close to each other. we display then only the equation of states for the parameter set $C_2$ in Fig. 1. The figure shows evidently that the symmetric nuclear matter can exist stably. However stable pure neutron matter does not exist.

Making use of the Hellmann-Feynman theorem\cite{18}, one can obtain the relation between the quark condensate in nuclear matter $Q(\rho) = \langle 0 | \bar{q}(x)q(x) | 0 \rangle_{\rho \neq 0}$ and that in vacuum $Q(0) = \langle 0 | \bar{q}(x)q(x) | 0 \rangle_{\rho = 0}$ as

$$
\frac{Q(\rho)}{Q(0)} = 1 + \frac{1}{2Q(0)} \left( \sum_{N=n,p} \frac{\partial \varepsilon}{\partial M_N} \frac{dM_N}{dm_q} + \frac{\partial \varepsilon}{\partial M_\sigma} \frac{dM_\sigma}{dm_q} + \frac{\partial \varepsilon}{\partial M_\pi} \frac{dM_\pi}{dm_q} \right)
= 1 - \frac{2\sigma_N}{(m_\pi f_\pi)^2} (\rho_s(p) + \rho_s(n)) - \frac{\pi_0^2}{f_\pi^2} - \frac{3g_3 m_\pi^2 \sigma_0^2}{2m_\sigma^4}
$$

where $\sigma_N$ is the nucleon $\sigma$ term, $\rho_s = \langle \bar{\psi}_s \psi \rangle$ is the scalar density of the nucleons.

With the parameter sets $C_2$ determined above and $\sigma_N = 45 \text{MeV}$, we evaluate the ratio of the in-medium quark condensate to that in vacuum. The obtained results are represented in Fig. 2. From the figure one can easily know that the quark condensate in nuclear matter decreases, in general, as the nuclear matter density increases. It manifests that the “upturn” problem in the other approaches at hadron level\cite{22, 23, 24} is removed in the present $\chi$SB approach. Investigating the figure more carefully, one may know that there exists a density $\rho = 0.235 \text{fm}^{-3}$ at which the decreasing rate is enhanced. Such a behavior is consistent with the result given in Ref.\cite{25}. Relevant investigation shows that pion condensate in strong interaction appears as the nuclear matter density reaches $0.235 \text{fm}^{-3}$ and beyond. It indicates that the sudden increase of the decreasing rate or
the discontinuous decrease of the quark condensate in nuclear matter may come from the appearance of pion condensate.

In summary, a nonlinear realization of SU(2) $\chi$SB approach is developed in the composite operator formalism of QCD. A Lagrangian including quark, gluon and Goldstone bosons of chiral quark model is obtained from the QCD Lagrangian. It indicates that, besides the constituent quarks, both the gluons and Goldstone bosons are the effective degrees of freedom to describe hadrons. We constructed also a way link the $\chi$SB at hadron level and that at quark level. As an application, the properties of symmetric nuclear matter and pure neutron matter are investigated in the mean-field approximation. The calculation shows that the present approach is quite powerful in describing the properties of nuclear matter. Meanwhile the quark condensate is evaluated. It shows that the quark condensate decreases with the increasing of nuclear matter density. The chiral symmetry is then restored gradually in nuclear matter as the density increases.

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Table 1. The parameters in the calculation of the SU(2) chiral symmetry spontaneous breaking model with $m_\pi = 139.57 \text{MeV}$ and $f_\pi = 130 \text{MeV}$ ($m_i (i = \sigma, \omega, \rho)$, $M_N$ and $\delta M_N$ in MeV, $g_2$ in fm$^{-1}$).

|   | $g_\sigma(N)$ | $m_\sigma$ | $g_\omega(N)$ | $m_\omega$ | $g_\rho(N)$ | $m_\rho$ | $g_2$ | $g_3$ | $M_N$ | $\delta M_N$ |
|---|--------------|------------|--------------|------------|-------------|---------|------|------|------|------------|
| C1 | 9.111        | 540        | 10.587       | 783        | 8.480       | 770     | -4.0 | 20.0 | 888  | 50         |
| C2 | 9.111        | 570        | 9.573        | 783        | 8.480       | 770     | -9.0 | 37.5 | 888  | 50         |
Figure Captions

**Fig. 1.** Obtained average energy per nucleon $\varepsilon/\rho_N - M_N$ as a function of nucleon density $\rho_N$ for the parameter sets $C_2$. The solid line denotes symmetric nuclear matter, and the dashed line is for pure neutron matter.

**Fig. 2.** Obtained ratio of the quark condensate in nuclear matter to that in vacuum as a function of nucleon density for the parameter sets $C_2$. The solid line denotes that in symmetric nuclear matter, and the dashed line is that for pure neutron matter.
$\epsilon / \rho_N - (M_N + \delta M) \text{ (MeV)}$ vs $\rho_N \text{ (fm}^{-3}\text{)}$
\( \frac{Q(\rho)}{Q(0)} \)

\( \rho_N \) (fm\(^{-3}\))