Neutron-skin thickness of finite nuclei in relativistic mean-field models with chiral limits

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We study several structure properties of finite nuclei using relativistic mean-field Lagrangians constructed according to the Brown-Rho scaling due to the chiral symmetry restoration at high densities. The models are consistent with current experimental constraints for the equations of state of symmetric matter at both normal and supra-normal densities and of asymmetric matter at sub-saturation densities. It is shown that these models can successfully describe the binding energies and charge radii of finite nuclei. Compared to calculations with usual relativistic mean-field models, these models give a reduced thickness of neutron skin in $^{208}\text{Pb}$ between 0.17 fm and 0.21 fm. The reduction of the predicted neutron skin thickness is found to be due to not only the softening of the symmetry energy but also the scaling property of $\rho$ meson required by the partial restoration of chiral symmetry.

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I. INTRODUCTION

The nuclear symmetry energy of isospin asymmetric nuclear matter not only plays a crucial role in a number of important issues in astrophysics, see, e.g., Refs. \cite{1, 2}, but is also important for understanding the structure of neutron- or proton-rich nuclei and the reaction dynamics of heavy-ion collisions, see, e.g., Ref. \cite{3–5}. However, the density dependence of the symmetry energy is still poorly known especially at supra-normal densities \cite{6, 7}. In a previous work \cite{8}, we constructed several relativistic mean-field (RMF) Lagrangians using in-medium hadron properties according to the Brown-Rho (BR) scaling due to the chiral symmetry restoration at high densities\cite{9, 10}. The scalings and associated parameters that describe the in-medium hadron properties are consistent with those from microscopic calculations or those extracted from recent experimental data. The symmetric part of the resulting equations of state (EOS) around normal density is consistent with the data of nuclear giant monopole resonances \cite{11} and at supra-normal densities it is constrained by the collective flow data from high energy heavy-ion reactions \cite{12}. Moreover, the density dependence of the symmetry energy at sub-saturation densities is in agreement with that extracted from the recent isospin diffusion data from intermediate energy heavy-ion reactions \cite{13–15}. An important feature of our models with the chiral limits is that the resulting symmetric nuclear matter EOS is soft at intermediate densities but stiff at high densities naturally, producing a maximum neutron star mass around $2.0M_\odot$ consistent with the recent astrophysical observations.

In the present work, we extend our well constrained RMF models to study ground-state properties of finite nuclei. In particular, we examine the neutron skin thickness of $^{208}\text{Pb}$. The size of the neutron skin of finite nuclei is determined by the competition between the neutron pressure and the surface tension. Indeed, it has been shown in many studies that the neutron skin thickness of $^{208}\text{Pb}$ is rather sensitive to the density dependence of the symmetry energy \cite{16–21}. The latter also influences the extraction of the incompressibility of nuclear matter from the experimental data of giant resonances in finite nuclei \cite{19, 22, 23}. The size of the neutron-skin in $^{208}\text{Pb}$ predicted
by these models thus also appears to be limited by both the symmetry energy and the nuclear matter incompressibility of the models used. For instance, non-relativistic models that give an incompressibility of about $\kappa = 220 \pm 15$ MeV and a symmetry energy at normal density less than 32 MeV produce a neutron skin thickness for $^{208}$Pb smaller than 0.2 fm, while the RMF models that have larger incompressibilities of 250-270 MeV and the symmetry energy at normal density larger than 34 MeV predict a neutron skin thickness for $^{208}$Pb as large as about 0.3 fm [21]. However, it is understood that it is the slope of the underlying symmetry energy that matters the most[21]. It was pointed out by Vretenar et.al. [19] that the incompressibility in RMF models was necessarily larger than 250 MeV for reproducing the nuclear structural properties. In a newly developed RMF model with the softened density dependence of the symmetry energy, Todd-Rutel and Piekarewicz obtained the incompressibility 230 MeV and predicted a neutron skin thickness of 0.21 fm for $^{208}$Pb [24].

On one hand, our models have the characteristic of a small symmetry energy of 31.6 MeV and a small incompressibility of 230 MeV at saturation density $\rho_0 = 0.16$ fm$^{-3}$. One would thus expect a thin neutron skin in $^{208}$Pb. On the other hand, the cancellation of the medium effect from the vector meson mass and its coupling constant in our models results in an EOS in resemblance to those of the linear Walecka model and the best-fit model NL3 which have a much larger neutron skin thickness in $^{208}$Pb. Given the successful description of the high-density behavior of nuclear matter, it is necessary to extend our models to the low density region by studying the ground-state properties of finite nuclei, especially the neutron skin thickness of heavy nuclei. Moreover, it is important to have a better understanding about the various contributing factors to the neutron skin thickness within the RMF models.

The paper is organized as follows. In Section II, we briefly introduce the formalism of RMF with Brown-Rho scaling due to the chiral symmetry restoration at high densities. Results on the ground-state properties of finite nuclei, especially the neutron-skin thickness of finite nuclei are presented in Section III. A summary is finally given in Section IV.

II. A BRIEF SUMMARY OF THE FORMALISM

A lot of efforts had been devoted to describing well the binding energies and charge radii of finite nuclei simultaneously[25–28] within the RMF models. Inspired by the nonlinearity of $\sigma$ meson in chiral $\sigma$ model or induced by the model renormalization, nonlinear $\sigma$ meson terms were introduced to describe the medium effects that are important in the Brueckner theory. As a result, the binding energies and charge radii of finite nuclei were well reproduced, along with an improved description for the nuclear surface. Later on, the $\omega$ meson nonlinear term was introduced in compliance with the relativistic Brueckner results [29]. Some prominent representatives of the best-fit RMF models are NL3 [30], TM1 [29], FSUGold [24], and so on. Based on the microscopic relativistic Brueckner results, the density-dependent Hartree or Hartree-Fock approaches were developed to describe the finite nuclei (see [31, 32] and references therein). These approaches are usually parameter free but their predictions are not as accurate as the best-fit models. Our models used here have no nonlinear meson self-interactions, while the medium effects are given by the BR scaling.

In the present work, the model Lagrangian with the density-dependent couplings and meson masses
is written as

\[ \mathcal{L} = \overline{\psi} i \gamma_{\mu} \partial^\mu - M + g_{\sigma}^* \sigma - g_{0}^* \gamma_{\mu} \omega^\mu - g_{\lambda}^* \gamma_{\mu} \tau_3 b_0^\mu + \frac{1}{2} \left( 1 + \tau_3 \right) \gamma_{\mu} A^{\mu} \psi + \frac{1}{2} (\partial_{\nu} \sigma \partial^\nu \sigma - \sigma^2 ) \]

\[ - \frac{1}{4} F_{\mu \nu} F^{\mu \nu} + \frac{1}{2} m_0^2 \omega_{\mu} \omega^\mu - \frac{1}{4} B_{\mu \nu} B^{\mu \nu} + \frac{1}{2} m_0^2 b_0 \mu b_0^\mu - \frac{1}{4} A_{\mu \nu} A^{\mu \nu} \]  

(1)

where \( \psi, \sigma, \omega, \) and \( b_0 \) are the fields of the nucleon, scalar, vector, and isovector-vector mesons, with their masses \( M, m_0^2, m_0^2, \) and \( m_0^2, \) respectively. The meson coupling constants and masses with asterisks denote the density dependence, given by the BR scaling [8, 33, 34]. The Dirac equation of nucleons in the mean-field approximation is given by

\[ [i \gamma_{\mu} \partial^\mu - M + g_{\sigma}^* \sigma - g_{0}^* \gamma_{\mu} \omega^\mu - g_{\lambda}^* \gamma_{\mu} \tau_3 b_0^\mu - \frac{1}{2} \left( 1 + \tau_3 \right) \gamma_{\mu} A^{\mu} ] \psi = 0. \]  

(2)

The rearrangement term \( \Sigma^{R,0} \) is essential for the thermodynamic consistency to derive the pressure and plays role in modifying the single particle energy and total binding energy of finite nuclei. In the mean field approximation, the \( \Sigma^R \) is given by

\[ \Sigma^R_0 = - \frac{\partial \mathcal{L}}{\partial \rho} = m_0^2 \omega_0 \frac{\partial m_0^2}{\partial \rho} + m_0^2 b_0 \frac{\partial m_0^2}{\partial \rho} - m_0^2 \sigma \frac{\partial m_0^2}{\partial \rho} - m_0^2 \omega_0 \frac{\partial \rho_0}{\partial \rho} - m_0^2 \frac{\partial \rho_0}{\partial \rho} + \rho_0 \sigma \frac{\partial \rho_0}{\partial \rho}. \]  

(3)

The density dependence of parameters is described by the scaling functions that are the ratios of the in-medium parameters to those in the free space. We take the scaling functions for the coupling constants of scalar and vector mesons as [8]

\[ \Phi_{\sigma}(\rho) = \frac{1}{1 + \xi \rho / \rho_0}, \quad \Phi_{\omega}(\rho) = \frac{1 - y_{\rho} \rho / \rho_0}{1 + y_{\omega} \rho / \rho_0}. \]  

(4)

For hadron masses, the scaling function reads

\[ \Phi(\rho) = 1 - y_{\rho} \rho / \rho_0. \]  

(5)

The energy per nucleon is given by [35]

\[ E = \sum_\alpha (2j_\alpha + 1) \epsilon_\alpha - \frac{1}{2} \int d^3 \vec{r} \left[ -g_0^* \sigma \rho_0 + g^*_0 \omega_0 \rho + g^*_0 b_0 \rho_3 + e A_0 \rho_0 - 2\Sigma^R_0(\rho) \right] + E_{c.m.}. \]  

(6)

where \( \epsilon \) is the eigen energy of nucleon, \( \alpha \) denotes the occupied state, and \( j \) is the quantum number of total angular momentum. The energy from the center of mass correction is \( E_{c.m.} = -3/4 \times 414^{-1/3} \) [28]. The equation of motion of mesons can be written out according to the Euler-Lagrange equation. The equations of nucleons and mesons are coupled nonlinearly and can be solved by iterations. The detailed procedure can be found in numerous literatures. Here, we just emphasize that the meson mass in the Green function is that in the free space. As an example, the \( \sigma \) meson field is integrated out as follows

\[ \sigma(r) = \int \phi_1 r_1^2 G(r, r_1, m_\sigma) [ -g_0^*(r_1) \rho_0(r_1) - (m_0^2 - m_\sigma^2(r_1)) \sigma(r_1) ] \]  

(7)

where \( m_\sigma \) is the free \( \sigma \) meson mass, \( g_0^*(r_1) = g_0^*(\rho(r_1)) \), and \( m_\sigma^*(r_1) = m_\sigma^*(\rho(r_1)) \). The expression of \( G \) can be found in Ref. [36].

### III. RESULTS AND DISCUSSIONS

The details in determining model parameters can be found in our previous work [8]. Here we will firstly do calculations based on the parameter set SL1. In Table I, we tabulate the parameter set
TABLE I: Parameter sets used in the calculations. The vacuum hadron masses are $M = 938\text{MeV}$, $m_\omega = 783\text{MeV}$ and $m_\rho = 770\text{MeV}$. The coupling constants given here are those at zero density. The incompressibility and the symmetry energy are respectively $230\text{MeV}$ and $31.6\text{MeV}$ at saturation density $\rho_0 = 0.16\text{fm}^{-3}$ for all models. The binding energy per nucleon $(B/A)$, $m_\sigma$, and the incompressibility $\kappa$ are in unit of MeV.

| Model | $m_\sigma$ | $g_\sigma$ | $g_\omega$ | $g_\rho$ | $y$ | $x$ | $B/A$ | $M^+/M$ | $y_\rho$ | $y_\rho$ |
|-------|------------|------------|------------|----------|-----|-----|-------|--------|----------|----------|
| SL1   | 600        | 10.3665    | 10.4634    | 3.7875   | 0.126| 0.234| -16.0 | 0.679  | -        | -        |
| SLC   | 590        | 10.1408    | 10.3261    | 3.8021   | 0.126| 0.239| -16.3 | 0.685  | -        | -        |
| SLCd  | 590        | 10.1408    | 10.3261    | 5.7758   | 0.126| 0.239| -16.3 | 0.685  | 0.5191   | 0.5191   |

SL1 which is a little different from that in Ref. [8] because of a larger $m_\sigma$ here. The nuclear matter property does not change at all by the variation of $m_\sigma$ because the ratio $g_\sigma/m_\sigma$ keeps unchanged, while the description for finite nuclei is modified. The latter can be seen in Eq. (7) and is correct for all cases irrespective of whether the coupling constant and meson mass are density dependent or not. For the linear Walecka model, the binding energies and charge radii can not be described well simultaneously by adjusting the $m_\sigma$ [27]. The second term in the integrand of Eq. (7) is of prime importance in improving the description of the binding energies and charge radii of finite nuclei. Besides, another improvement comes from the inclusion of the rearrangement term appearing in Eqs. (2) and (6). Using $m_\sigma = 500\text{MeV}$ in the original version of SL1, the total binding energy for $^{40}\text{Ca}$ is smaller than the experimental value by $70\text{MeV}$, and for $^{208}\text{Pb}$ the deviation is about $230\text{MeV}$. These results are improved appreciably by shifting the $m_\sigma$ up to $600\text{MeV}$, as shown in Table II. The relative error of charge radii is less than 1.5% and that of binding energy is less than 3%.

The accuracy of predictions can be much improved by setting $B/A = -16.3\text{MeV}$ and $m_\sigma = 590\text{MeV}$. Such a slight shift changes very little properties of infinite nuclear matter since the incompressibility is kept the same as that of the SL1 parameter set. The new parameter set called SLC is also listed in Table I. Except for $^{16}\text{O}$, excellent agreement with the experimental data is achieved using the model SLC, as shown in Table II, though the binding energy of medium-mass nuclei is slightly underestimated by about 1.5%. To see the effects of the density dependence of the symmetry energy on the properties of finite nuclei, we derive the parameter set SLCd. The parameter set SLCd is different from the SLC only for the $\rho$ meson coupling constant and its density dependence which is modified by an additional coefficient $y_\rho$, as given in Table I. The results obtained by using the SLCd parameter set are given in Table II, where for comparison the FSUGold results are also listed. Fig. 1 shows the density dependence of the symmetry energy with the SLC, SLCd and SL1 parameter sets. For comparison, we also include the MDI(x=0) and MDI(x=-1) results which represent the experimental constraints on the density dependence of the symmetry energy at sub-saturation densities, derived by studying the isospin diffusion data within a transport model [14, 15]. As shown in Fig. 1, no obvious difference between the SLC and SL1 is visible for the density dependence of the symmetry energy. Interestingly, all the symmetry energies from the parameter sets SLC, SLCd and SL1 are consistent with that constrained by the isospin diffusion data.

It is now well known that the neutron skin thickness $(r_n - r_p)$ of heavy nuclei depends sensitively on the slope of the symmetry energy at normal density. A stiffer density dependence of the symmetry energy results in a thicker neutron skin. The difference between the symmetry energies with the SLC and SLCd, shown in Fig. 1, is responsible for the variation of the neutron skin in neutron-rich nuclei.
FIG. 1: (Color online) The symmetry energy as a function of density for different models. The SL1 and SLC results overlap completely.

TABLE II: The binding energy per nucleon (B/A), charge radii \( r_c \) in fm, and neutron skins \( r_n - r_p \) obtained from different models. The available experimental data are from [37, 38].

| Nucleus | Expt. | SL1 | SLC | SLCd | FSUGold |
|---------|-------|-----|-----|------|---------|
| \(^{16}\text{O}\) | B/A  | 7.98 | 8.03 | 8.07 | 8.07 | 7.96 |
| r_c     | 2.70 | 2.72 | 2.74 | 2.74 | 2.69 |
| r_n - r_p | -0.03 | -0.03 | -0.03 | -0.03 |
| \(^{40}\text{Ca}\) | B/A  | 8.55 | 8.43 | 8.54 | 8.54 | 8.54 |
| r_c     | 3.48 | 3.43 | 3.45 | 3.44 | 3.42 |
| r_n - r_p | -0.05 | -0.05 | -0.05 | -0.05 |
| \(^{48}\text{Ca}\) | B/A  | 8.67 | 8.41 | 8.52 | 8.46 | 8.58 |
| r_c     | 3.47 | 3.45 | 3.46 | 3.47 | 3.44 |
| r_n - r_p | -0.20 | -0.20 | -0.18 | -0.20 |
| \(^{90}\text{Zr}\) | B/A  | 8.71 | 8.44 | 8.59 | 8.56 | 8.68 |
| r_c     | 4.26 | 4.22 | 4.23 | 4.23 | 4.25 |
| r_n - r_p | 0.10 | 0.10 | 0.08 | 0.09 |
| \(^{114}\text{Sn}\) | B/A  | 8.52 | 8.24 | 8.40 | 8.38 | 8.49 |
| r_c     | 4.61 | 4.57 | 4.58 | 4.59 | 4.60 |
| r_n - r_p | 0.09 | 0.09 | 0.07 | 0.09 |
| \(^{132}\text{Sn}\) | B/A  | 8.36 | 8.18 | 8.35 | 8.23 | 8.34 |
| r_c     | 4.68 | 4.68 | 4.69 | 4.70 | 4.71 |
| r_n - r_p | 0.27 | 0.27 | 0.23 | 0.27 |
| \(^{208}\text{Pb}\) | B/A  | 7.87 | 7.68 | 7.87 | 7.79 | 7.89 |
| r_c     | 5.50 | 5.46 | 5.47 | 5.48 | 5.52 |
| r_n - r_p | 0.21 | 0.21 | 0.17 | 0.21 |

as seen in Table II. Noting that the model FSUGold has almost the same symmetry energy in the whole density region as that of the MDI(x=0) and has larger symmetry energies at low densities than the SL1 and SLC, the neutron-skin thickness in \(^{208}\text{Pb}\) with the SL1 and SLC would thus be expected to be much larger than that given by the FSUGold. However, from Table II, we can see surprisingly
that the SL1, SLC and FSUGold predict almost the same neutron skin for magic nuclei listed in Table II. Therefore, given the same density dependence of the symmetry energy as with the best-fit model FSUGold, the present models with the BR scaling predict smaller values for the neutron-skin thickness for finite nuclei. Compared to other best-fit RMF models, such as the NL1 and TM1 that predict the thickness of neutron skin in $^{208}$Pb to be 0.28fm and 0.26fm, respectively, this reduction of the predicted neutron-skin thickness is more appreciable. To understand the underlying physics for the above observation, we may write down the integration form for the isovector potential that is a product of the $\rho$ meson field and the coupling constant $g^*_\rho$

$$V_\rho(r) = g^*_\rho(r) \int \rho(r_1) \rho_3(r_1) \left[ -g^*_\rho(r_1) \rho_3(r_1) - (m^2_\rho - m^2_\rho r^2(r_1)) b_0(r_1) \right]. \quad (8)$$

For the SL1 and SLC, the ratio $C_\rho = g^*_\rho/m^*_\rho$ is a constant. Properties of asymmetric nuclear matter in the RMF approximation depend on the constant ratio $C_\rho$ rather than the respective values of the coupling constant and the mass. This is the same as in the linear Walecka model and the best-fit model NL3. However, the linear Walecka model and the NL3 model predict much thicker neutron skin of 0.27fm and 0.28fm in 208Pb, respectively [27, 30]. As the symmetry energy changes from 35MeV to 31.6MeV in the linear Walecka model, the neutron thickness is still as large as 0.25fm. It is similar for the NL3 model, but a reduction in the symmetry energy to 31.6MeV results in large deviations of calculated nucleus masses from the experimental values [19], which is strongly unfavored by the best-fit model. The linear Walecka model, the NL3 and the FSUGold all have a much smaller effective nucleon mass $M^*$ at the saturation density than ours. It should be noted that it is not the larger $M^*$ that results in the small neutron skin since a small neutron skin of 0.22fm in 208Pb is still obtained with our model that is adjusted to have a small $M^* = 0.6M$ with the same symmetry energies as given by the SL1 and SLC parameter sets. It is interesting that the scaling property of the $\rho$ meson that is hidden behind calculations for nuclear matter is now displayed clearly in finite nuclei. The second term in the integrand of Eq. (8) is induced by the $\rho$ mass scaling. Together with the density dependent coupling constant, it is responsible for the reduction of the neutron skin thickness in neutron-rich nuclei. The drop of the neutron skin thickness with the SLCd, compared to the SL1 and SLC, can just be attributed to the softening of the symmetry energy. We note that the neutron skin is totally insensitive to the value of $m_\sigma$, though a better fit of both binding energies and charge radii relies on the choice of $m_\sigma$. The thickness of neutron skin in $^{208}$Pb varies from 0.21fm to 0.17fm by changing from the SLC to the SLCd. This range is consistent with current measurements using the X-ray cascade from antiprotonic atoms: 0.16 $\pm$ 0.06fm [39] and also agrees well with that obtained from the analysis of the isospin diffusion data [20]: 0.22 $\pm$ 0.04. Interestingly, the predicted range of neutron skin in $^{208}$Pb between 0.21fm and 0.17fm is very close to that predicted by using some Skyrme interactions [21].

It is worth noting that a much smaller neutron-skin thickness in $^{208}$Pb can be obtained within the RMF models provided that smaller values of symmetry energy are given. For instance, the ES25 model [2] constructed with the incompressibility of 211.7 MeV and the symmetry energy of 25 MeV at saturation density predicts a neutron skin as thin as 0.138 fm for $^{208}$Pb. With the same incompressibility and symmetry energies as the ES25 model, our model predicts an even smaller neutron skin of 0.126 fm for $^{208}$Pb. Of course, the magnitude of the reduction due to the in-medium properties of the $\rho$ meson is accordingly smaller.

The deficiency of our models mainly lies in the unsatisfactory description for spin-orbit splittings. For instance, the 1p spin-orbit splitting in $^{16}$O is 3.7 and 3.5MeV with the SL1 and SLC, respectively,
which is smaller than the experimental value by about 2.5 MeV. This deficiency stems mainly from a much larger Dirac mass $M^*$ in our present models, i.e., about 0.68M. Usually, a Dirac mass of about 0.5-0.6M can reasonably describe the spin-orbit splitting in RMF models (see, e.g., Ref. [40]). As a result, the spin-orbit potential of the central part is much suppressed in our models, compared to that given by the best-fit models. The deviation of spin-orbit splittings from the available data is comparable to that given by the density dependent approaches based on the relativistic Brueckner results [32]. The right spin-orbit splitting in our calculations can be obtained by reducing the Dirac mass moderately. On the other hand, the larger Dirac mass leads to larger Landau mass (about 0.74 in our models), which gives a more reasonable description for the level density [40].

**IV. SUMMARY**

In summary, we have studied ground-state properties of finite nuclei within the relativistic mean-field models that are constructed according to the Brown-Rho scaling due to the chiral symmetry restoration at high densities. Not only can the constructed models SLC and SLCd be consistent with current experimental results for the equation of state of symmetric matter at normal and supra-normal densities and of asymmetric matter at sub-saturation densities, but also present a fairly satisfactory description for the ground state properties of finite nuclei. The binding energies and charge radii of a variety of magic nuclei are nicely reproduced. Our results indicate that the neutron-skin thickness of finite nuclei depends not only on the density dependence of the symmetry energy but also on the scaling property of the $\rho$ meson. Compared to calculations of usual RMF models a reduction of neutron skin thickness of neutron-rich nuclei is observed in our models. We find that the scaling property of the $\rho$ meson that is hidden in nuclear matter calculations is shown up clearly by the reduction of the neutron skin in $^{208}$Pb. The neutron skin thickness of $^{208}$Pb is predicted to be between 0.17 fm to 0.21 fm consistent with the current measurements.

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[1] J. M. Lattimer and M. Prakash, Phys. Rep. 333, 121 (2000); Astrophys. J. 550, 426 (2001); Science 304, 536 (2004).
[2] A. W. Steiner, M. Prakash, J. M. Lattimer and P. J. Ellis, Phys. Rep. 411, 325 (2005).
[3] Isospin Physics in Heavy-Ion Collisions at Intermediate Energies, Eds. Bao-An Li and W. Udo Schröder (Nova Science Publishers, Inc, New York, 2001).
[4] C. J. Horowitz, J. Piekarewicz, Phys. Rev. Lett. 86, 5647 (2001).
[5] W. Z. Jiang, Y. L. Zhao, Phys. Lett. B617,33 (2005).
[6] B. A. Li, Phys. Rev. Lett. 85, 4221 (2000); ibid. 88, 192701 (2002).
[7] L.W. Chen, C.M. Ko, B.A. Li and G.C. Yong, arXiv:0704.2340 [nucl-th], Frontiers of Physics in China, Vol. 2, 327(2007).
[8] W. Z. Jiang, B. A. Li and L. W. Chen, Phys. Lett. B653, 184 (2007).
[9] G. E. Brown, M. Rho, Phys. Rev. Lett. 66, 2720 (1991).
[10] G. E. Brown, J. W. Holt, C.-H. Lee, M. Rho, Phys. Rep. 439, 161 (2007).
[11] D.H. Youngblood, H.L. Clark, and Y.W. Lui, Phys. Rev. Lett. 82, 691 (1999).
[12] P. Danielewicz, R. Lacey, W. G. Lynch, Science 298, 1592, (2002).
[13] M. B. Tsang et.al. Phys. Rev. Lett. 92, 062701 (2004).
[14] L. W. Chen, C. M. Ko, B. A. Li, Phys. Rev. Lett. 94, 032701 (2005).
[15] B. A. Li, L. W. Chen, Phys. Rev. C72, 064611 (2005).
[16] B.A. Brown, Phys. Rev. Lett. 85, 5296 (2000).
[17] S. Typel and B.A. Brown, Phys. Rev. C 64, 027302 (2001).
[18] J. Piekarewicz, Phys. Rev. C66, 034305 (2002).
[19] D. Vretenar, T. Niksic, and P. Ring, Phys. Rev. C68, 024310 (2003).
[20] L. W. Chen, C. M. Ko, B. A. Li, Phys. Rev. C72, 064309 (2005).
[21] R. J. Furnstahl, Nucl. Phys. A706, 85 (2002).
[22] J. Piekarewicz, Phys. Rev. C69, 041301 (2004).
[23] G. Colo', N. Van Giai, J. Meyer, K. Benaceur and P. Bonche, Phys. Rev. C70, 024307 (2004).
[24] B. G. Todd-Rutel, J. Piekarewicz, Phys. Rev. Lett. 95, 122501 (2005).
[25] S. A. Chin and J. D. Walecka, Phys. Lett. B52, 24 (1974); S. A. Chin, Ann. of Phys. 108, 301 (1977).
[26] J. Boguta, A. R. Bodmer, Nucl. Phys. A292 (1977) 413;
[27] C. J. Horowitz, B. D. Serot, Nucl. Phys. A368, 503 (1981).
[28] P. G. Reinhard, M. Rüfa, J. Maruhn, W. Greiner, J. Friedrich Z. Phys. A323, 13 (1986).
[29] Y. Sugahara and H. Toki, Nucl. Phys. A579, 557 (1994).
[30] G. A. Lalazissis, J. König, and P. Ring, Phys. Rev. C55, 540 (1997).
[31] Z. Y. Ma, L. Liu, Phys. Rev. C66, 024321 (2002).
[32] F. Hofmann, C. M. Keil, H. Lenske, Phys. Rev. C64, 034314 (2001).
[33] G. E. Brown and M. Rho, nucl-th/0509001; ibid. nucl-th/0509002.
[34] C. Song, Phys. Rep. 347, 289 (2001).
[35] C. Fuchs, H. Lenske, H. H. Wolter, Phys. Rev. C52, 3043 (1995).
[36] W. Z. Jiang, Y. L. Zhao, Z. Y. Zhu, S. F. Shen, Phys. Rev. C72, 024313 (2005).
[37] G. Audia, A.H. Wapstra, C. Thibault, Nucl. Phys. A729, 337(2003).
[38] I. Angeli, Atom. Data Nucl. Data Tab. 87, 185 (2004).
[39] B. Klos, A. Trzcinska, J. Jastrzebski, et. al., nucl-ex/0702016.
[40] S. Typel, Phys. Rev. C71, 064301 (2005).