Higher-order effects on the incompressibility of isospin asymmetric nuclear matter

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Abstract

Analytical expressions for the saturation density of asymmetric nuclear matter as well as its binding energy and incompressibility at saturation density are given up to 4th-order in the isospin asymmetry $\delta = (\rho_n - \rho_p)/\rho$ using 11 characteristic parameters defined by the density derivatives of the binding energy per nucleon of symmetric nuclear matter, the symmetry energy $E_{\text{sym}}(\rho)$ and the 4th-order symmetry energy $E_{\text{sym,4}}(\rho)$ at the normal nuclear density $\rho_0$. Using an isospin- and momentum-dependent modified Gogny (MDI) interaction and the Skyrme-Hartree-Fock (SHF) approach with 63 popular Skyrme interactions, we have systematically studied the isospin dependence of the saturation properties of asymmetric nuclear matter, particularly the incompressibility $K_{\text{sat}}(\delta) = K_0 + K_{\text{sat,2}}\delta^2 + K_{\text{sat,4}}\delta^4 + O(\delta^6)$ at saturation density. Our results show that the magnitude of the higher-order $K_{\text{sat,4}}$ parameter is generally small compared to that of the $K_{\text{sat,2}}$ parameter. The latter essentially characterizes the isospin dependence of the incompressibility at saturation density and can be expressed as $K_{\text{sat,2}} = K_{\text{sym}} - 6L - \frac{J_0}{K_0}L$, where $L$ and $K_{\text{sym}}$ represent, respectively, the slope and curvature parameters of the symmetry energy at $\rho_0$ while $J_0$ is the third-order derivative parameter of symmetric nuclear matter at $\rho_0$. Furthermore, we have constructed a phenomenological modified Skyrme-like (MSL) model which can reasonably describe the general properties of symmetric nuclear matter and the symmetry energy predicted by both the MDI model and the SHF approach. The results indicate that the higher-order $J_0$ contribution to $K_{\text{sat,2}}$ generally cannot be neglected. In addition, it is found that there exists a nicely linear correlation between $K_{\text{sym}}$ and $L$ as well as between $J_0/K_0$ and $K_0$. These correlations together with the empirical constraints on $K_0$, $L$, $E_{\text{sym}}(\rho_0)$ and the nucleon effective mass lead to an estimate of $K_{\text{sat,2}} = -370 \pm 120$ MeV.

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I. INTRODUCTION

The equation of state (EOS) of nuclear matter is one of central issues in nuclear physics. For a cold nuclear matter, the EOS is usually defined as the binding energy per nucleon as a function of the density from which information on other thermodynamic properties of nuclear matter, such as its pressure and incompressibility can be obtained. With the establishment or construction of many radioactive beam facilities around the world, such as the Cooling Storage Ring (CSR) facility at HIRFL in China [1], the Radioactive Ion Beam (RIB) Factory at RIKEN in Japan [2], the FAIR/GSI in Germany [3], SPIRAL2/GANIL in France [4], and the Facility for Rare Isotope Beams (FRIB) in USA [5], it is possible in terrestrial laboratories to explore the EOS of an isospin asymmetric nuclear matter under the extreme condition of large isospin asymmetry, especially the density dependence of the nuclear symmetry energy. This knowledge, especially the latter, is important for understanding not only the structure of radioactive nuclei, the reaction dynamics induced by rare isotopes, and the liquid-gas phase transition in asymmetric nuclear matter, but also many critical issues in astrophysics [6, 7, 8, 9, 10, 11, 12, 13, 14, 15].

For symmetric nuclear matter with equal fractions of neutrons and protons, its EOS is relatively well-determined after about more than 30 years of studies by the nuclear physics community. In particular, the incompressibility of symmetric nuclear matter at its saturation density $\rho_0$ has been determined to be $240 \pm 20$ MeV from analyses of the nuclear giant monopole resonances (GMR) [16, 17, 18, 19, 20, 21, 22, 23, 24], and its EOS at densities of $2\rho_0 < \rho < 5\rho_0$ has also been constrained by measurements of collective flows in nucleus-nucleus collisions [8] and of subthreshold kaon production [25, 26] in relativistic nucleus–nucleus collisions. On the other hand, the EOS of asymmetric nuclear matter, especially the density dependence of the nuclear symmetry energy, is largely unknown. Although the nuclear symmetry energy at $\rho_0$ is known to be around 30 MeV from the empirical liquid-drop mass formula [27, 28], its values at other densities, especially at supra-saturation densities, are poorly known [6, 7]. Various microscopic and phenomenological models, such as the relativistic Dirac-Brueckner-Hartree-Fock (DBHF) [29, 30, 31, 32, 33, 34, 35] and the non-relativistic Brueckner-Hartree-Fock (BHF) [36, 37, 38, 39] approach, the relativistic mean-field (RMF) model based on nucleon-meson interactions [12, 40, 41, 42], and the non-relativistic mean-field model based on Skyrme-like interactions [43, 44, 45, 46, 47, 48, 49].
have been used to study the isospin-dependent properties of asymmetric nuclear matter, such as the nuclear symmetry energy, the nuclear symmetry potential, the isospin-splitting of the nucleon effective masses, etc., but the predicted results vary widely. In fact, even the sign of the symmetry energy above $3\rho_0$ is still uncertain \[52, 53\]. The theoretical uncertainties are mainly due to the lack of knowledge about the isospin dependence of in-medium nuclear effective interactions and the limitations in the techniques for solving the nuclear many-body problem.

Heavy-ion collisions, especially those induced by neutron-rich nuclei, provide a unique tool to investigate the EOS of asymmetric nuclear matter, especially the density dependence of the nuclear symmetry energy. During the last decade, significant progress has indeed been made both experimentally and theoretically on constraining the behavior of the symmetry energy at subsaturation density using heavy-ion reactions \[50, 54, 55, 56, 57, 58, 59, 60\] (See Ref. \[15\] for the most recent review). More recently, the IBUU04 transport model analysis of the FOPI data on the $\pi^-/\pi^+$ ratio in central heavy-ion collisions at SIS/GSI \[61\] energies suggests a very soft symmetry energy at the suprasaturation densities \[62\]. Information on the density dependence of the nuclear symmetry energy has also been obtained from the structure of finite nuclei and their excitations, such as the mass data \[63\], neutron skin in heavy nuclei \[64\], giant dipole resonances \[65\], pygmy dipole resonance \[66\], and so on. These studies have significantly improved our understanding of the EOS of asymmetric nuclear matter.

The incompressibility of asymmetric nuclear matter at its saturation density is a basic quantity to characterize its EOS. Since this quantity is largely undetermined, any constraint imposed on it would be extremely important. In the present work, we study the isospin dependence of the properties of asymmetric nuclear matter, including the saturation density as well as the binding energy and incompressibility at saturation density. In particular, we derive analytical expressions for these quantities up to the 4th-order in the isospin asymmetry $\delta = (\rho_n - \rho_p)/\rho$ and investigate the higher-order isospin asymmetry effects on the properties of asymmetric nuclear matter. For the incompressibility of an asymmetric nuclear matter at its saturation density, it can be written as $K_{\text{sat}}(\delta) = K_0 + K_{\text{sat},2}\delta^2 + K_{\text{sat},4}\delta^4 + O(\delta^6)$ with $K_{\text{sat},2} = K_{\text{sym}} - 6L - \frac{J_0}{K_0}L$, where $L$ and $K_{\text{sym}}$ represent, respectively, the slope and curvature parameters of the symmetry energy at $\rho_0$ while $J_0$ is the 3rd-order derivative parameter of symmetric nuclear matter at $\rho_0$. Therefore, the higher-order effects on $K_{\text{sat}}(\delta)$ also include
the 3rd-order density derivative parameter $J_0$. Our results indicate that higher-order isospin asymmetry effects on the incompressibility are usually not important but the higher-order $J_0$ contribution generally cannot be neglected. Furthermore, we construct a phenomenological modified Skyrme-like (MSL) model which can reasonably describe the general properties of symmetric nuclear matter and the symmetry energy predicted by both the MDI model and the Skyrme-Hartree-Fock (SHF) approach. We find that there exists a nicely linear correlation between $K_{\text{sym}}$ and $L$ as well as between $J_0/K_0$ and $K_0$. These correlations together with the empirical constraints on $K_0$, $L$ and $E_{\text{sym}}(\rho_0)$ lead to an estimate of $K_{\text{sat},2} = -370 \pm 120$ MeV.

The paper is organized as follows. In Section II, we discuss the general properties of asymmetric nuclear matter, and then give analytical expressions up to 4th-order terms in isospin asymmetry $\delta$ for the saturation density of asymmetric nuclear matter as well as its binding energy and incompressibility at saturation density. We then briefly introduce in Section III the three models used in the present paper, i.e., the isospin- and momentum-dependent MDI model, the SHF approach, and the phenomenological MSL model. The results and discussions are presented in Section IV. A summary is then given in Section V. For completeness, derivations of some important formula shown in Section II are briefly described in the Appendix.

II. SATURATION PROPERTIES OF ASYMMETRIC NUCLEAR MATTER

A. Equation of state of asymmetric nuclear matter

The EOS of isospin asymmetric nuclear matter, given by its binding energy per nucleon, can be generally expressed as a power series in the isospin asymmetry $\delta = (\rho_n - \rho_p)/\rho$, where $\rho = \rho_n + \rho_p$ is the baryon density with $\rho_n$ and $\rho_p$ denoting the neutron and proton densities, respectively. To the 4th-order in isospin asymmetry, it is written as

$$E(\rho, \delta) = E_0(\rho) + E_{\text{sym}}(\rho)\delta^2 + E_{\text{sym,4}}(\rho)\delta^4 + O(\delta^6),$$ (1)
where \( E_0(\rho) = E(\rho, \delta = 0) \) is the binding energy per nucleon of symmetric nuclear matter, and

\[
E_{\text{sym}}(\rho) = \frac{1}{2!} \frac{\partial^2 E(\rho, \delta)}{\partial \delta^2}, \quad (2) \\
E_{\text{sym,4}}(\rho) = \frac{1}{4!} \frac{\partial^4 E(\rho, \delta)}{\partial \delta^4} \bigg|_{\delta = 0}. \tag{3}
\]

In the above, \( E_{\text{sym}}(\rho) \) is the so-called nuclear symmetry energy and \( E_{\text{sym,4}}(\rho) \) is the 4th-order coefficient, which is called here the 4th-order nuclear symmetry energy. The absence of odd-order terms in \( \delta \) in Eq. (1) is due to the exchange symmetry between protons and neutrons in nuclear matter when one neglects the Coulomb interaction and assumes the charge symmetry of nuclear forces. The nuclear symmetry energy \( E_{\text{sym}}(\rho) \) thus corresponds to the lowest-order coefficient. The higher-order (including 4th-order) coefficients in \( \delta \) are usually very small and neglected, e.g., the magnitude of the \( \delta^4 \) term at the normal nuclear density \( \rho_0 \) (the saturation density of symmetric nuclear matter) is estimated to be less than 1 MeV in microscopic many-body approaches \[67, 68, 69\] and also in phenomenological models as shown later. Neglecting the contribution from higher-order terms in Eq. (1) leads to the well-known empirical parabolic law for the EOS of asymmetric nuclear matter, which has been verified by all many-body theories to date, at least for densities up to moderate values \[15\]. As a good approximation, the density-dependent symmetry energy \( E_{\text{sym}}(\rho) \) can thus be extracted from the parabolic approximation of \( E_{\text{sym}}(\rho) \approx E(\rho, \delta = 1) - E(\rho, \delta = 0) \).

It should be mentioned that the possible presence of the higher-order terms in \( \delta \) at supranormal densities can significantly modify the proton fraction in \( \beta \)-equilibrium neutron-star matter and the critical density for the direct Urca process which can lead to faster cooling of neutron stars \[70, 71\]. In addition, a recent study \[72, 73\] indicates that the higher-order terms in \( \delta \) are very important for determining the transition density and pressure at the inner edge separating the liquid core from the solid crust of neutron stars where the matter is extremely neutron-rich.

Around the nuclear matter saturation density \( \rho_0 \), the binding energy per nucleon in symmetric nuclear matter \( E_0(\rho) \) can be expanded, e.g., up to 4th-order in density, as

\[
E_0(\rho) = E_0(\rho_0) + L_0 \chi + \frac{K_0}{2!} \chi^2 \\
+ \frac{J_0}{3!} \chi^3 + \frac{I_0}{4!} \chi^4 + O(\chi^5), \tag{4}
\]
where $\chi$ is a dimensionless variable characterizing the deviations of the density from the saturation density $\rho_0$ of symmetric nuclear matter and it is conventionally defined as

$$\chi = \frac{\rho - \rho_0}{3\rho_0}. \quad (5)$$

The first term $E_0(\rho_0)$ on the right-hand-side (r.h.s) of Eq. (4) is the binding energy per nucleon in symmetric nuclear matter at the saturation density $\rho_0$ and the coefficients of other terms are

$$L_0 = 3\rho_0 \frac{dE_0(\rho)}{d\rho}|_{\rho=\rho_0}, \quad (6)$$

$$K_0 = 9\rho_0^2 \frac{d^2E_0(\rho)}{d\rho^2}|_{\rho=\rho_0}, \quad (7)$$

$$J_0 = 27\rho_0^3 \frac{d^3E_0(\rho)}{d\rho^3}|_{\rho=\rho_0}, \quad (8)$$

$$I_0 = 81\rho_0^4 \frac{d^4E_0(\rho)}{d\rho^4}|_{\rho=\rho_0}. \quad (9)$$

Obviously, we have $L_0 = 0$ according to the definition of the saturation density $\rho_0$ of symmetric nuclear matter and thus the second term on the r.h.s of Eq. (4) should vanish. The coefficient $K_0$ is the so-called incompressibility coefficient of symmetric nuclear matter and it characterizes the curvature of $E_0(\rho)$ at $\rho_0$. The coefficients $J_0$ and $I_0$ correspond to higher-order contributions and here we call them as 3rd-order and 4th-order incompressibility coefficients of symmetric nuclear matter, respectively. In the literature, one usually neglects the higher-order terms in Eq. (4) around the saturation density $\rho_0$ and obtain the following parabolic approximation to the EOS of symmetric nuclear matter:

$$E_0(\rho) = E_0(\rho_0) + \frac{K_0}{2} \chi^2 + O(\chi^3). \quad (10)$$

Around the normal nuclear density $\rho_0$, the nuclear symmetry energy $E_{\text{sym}}(\rho)$ can be similarly expanded, e.g., up to 4th-order in $\chi$, as

$$E_{\text{sym}}(\rho) = E_{\text{sym}}(\rho_0) + L\chi + \frac{K_{\text{sym}}}{2!} \chi^2 + \frac{J_{\text{sym}}}{3!} \chi^3 + \frac{I_{\text{sym}}}{4!} \chi^4 + O(\chi^5), \quad (11)$$

where $L$, $K_{\text{sym}}$, $J_{\text{sym}}$ and $I_{\text{sym}}$ are the slope parameter, curvature parameter, 3rd-order
coefficient, and 4th-order coefficient of the nuclear symmetry energy at \( \rho_0 \), i.e.,

\[
L = 3\rho_0 \frac{dE_{\text{sym}}(\rho)}{d\rho}|_{\rho=\rho_0},
\]

\[
K_{\text{sym}} = 9\rho_0^2 \frac{d^2E_{\text{sym}}(\rho)}{d\rho^2}|_{\rho=\rho_0},
\]

\[
J_{\text{sym}} = 27\rho_0^3 \frac{d^3E_{\text{sym}}(\rho)}{d\rho^3}|_{\rho=\rho_0},
\]

\[
I_{\text{sym}} = 81\rho_0^4 \frac{d^4E_{\text{sym}}(\rho)}{d\rho^4}|_{\rho=\rho_0}.
\]

The coefficients \( L, K_{\text{sym}}, J_{\text{sym}} \) and \( I_{\text{sym}} \) characterize the density dependence of the nuclear symmetry energy around the normal nuclear density \( \rho_0 \), and thus carry important information on the properties of nuclear symmetry energy at both high and low densities.

We can also expand the 4th-order nuclear symmetry energy \( E_{\text{sym},4}(\rho) \) around the normal nuclear density \( \rho_0 \) up to 4th-order in \( \chi \) as

\[
E_{\text{sym},4}(\rho) = E_{\text{sym},4}(\rho_0) + L_{\text{sym},4}\chi + \frac{K_{\text{sym},4}}{2}\chi^2
+ \frac{J_{\text{sym},4}}{3!}\chi^3 + \frac{I_{\text{sym},4}}{4!}\chi^4 + O(\chi^5),
\]

where \( L_{\text{sym},4}, K_{\text{sym},4}, J_{\text{sym},4} \) and \( I_{\text{sym},4} \) are the slope parameter, curvature parameter, 3rd-order coefficient, and 4th-order coefficient of the 4th-order nuclear symmetry energy \( E_{\text{sym},4}(\rho) \) at \( \rho_0 \), i.e.,

\[
L_{\text{sym},4} = 3\rho_0 \frac{dE_{\text{sym},4}(\rho)}{d\rho}|_{\rho=\rho_0},
\]

\[
K_{\text{sym},4} = 9\rho_0^2 \frac{d^2E_{\text{sym},4}(\rho)}{d\rho^2}|_{\rho=\rho_0},
\]

\[
J_{\text{sym},4} = 27\rho_0^3 \frac{d^3E_{\text{sym},4}(\rho)}{d\rho^3}|_{\rho=\rho_0},
\]

\[
I_{\text{sym},4} = 81\rho_0^4 \frac{d^4E_{\text{sym},4}(\rho)}{d\rho^4}|_{\rho=\rho_0}.
\]

In above Taylor expansions, we have kept all terms up to 4th-order in \( \delta \) or \( \chi \). The 14 coefficients, namely, \( E_0(\rho_0), K_0, J_0, I_0, E_{\text{sym}}(\rho_0), L, K_{\text{sym}}, J_{\text{sym}}, I_{\text{sym}}, E_{\text{sym},4}(\rho_0), L_{\text{sym},4}, K_{\text{sym},4}, J_{\text{sym},4} \) and \( I_{\text{sym},4} \), are well-defined, and they characterize the EOS of an asymmetric nuclear matter and its density dependence at the normal nuclear density \( \rho_0 \). Among these parameters, \( E_0(\rho_0), K_0, E_{\text{sym}}(\rho_0), L, \) and \( K_{\text{sym}} \) have been extensively studied in the literature.
and significant progress has been made over past few decades. Based on these well-defined quantities, we investigate in the present paper to what extend they can provide reliable information on the isospin dependence of the saturation properties of asymmetric nuclear matter as well as its properties at both high and low densities.

**B. Saturation density**

One of the basic quantities for describing an asymmetric nuclear matter is its saturation density $\rho_{\text{sat}}(\delta)$ which is generally a function of the isospin asymmetry $\delta$. According to the definition of the saturation density $\rho_{\text{sat}}(\delta)$ of asymmetric nuclear matter, i.e.,

$$\left. \frac{\partial E(\rho, \delta)}{\partial \rho} \right|_{\rho=\rho_{\text{sat}}(\delta)} = 0,$$

it can be shown that up to 4th-order in $\delta$, the saturation density can be expressed as (see Appendix)

$$\rho_{\text{sat}}(\delta) = \left[ 1 - \frac{3L}{K_0} \delta^2 + \left( \frac{3K_{\text{sym}}L}{K_0^2} - \frac{3L_{\text{sym},4}}{K_0} \right) \delta^4 \right] \rho_0. \quad (22)$$

Therefore, $\rho_{\text{sat}}(\delta)$ can be written as

$$\rho_{\text{sat}}(\delta) = \rho_0 + \rho_{\text{sat},2} \delta^2 + \rho_{\text{sat},4} \delta^4 + O(\delta^6) \quad (23)$$

with

$$\rho_{\text{sat},2} = -\frac{3L}{K_0} \rho_0 \quad (24)$$

$$\rho_{\text{sat},4} = \left( \frac{3K_{\text{sym}}L}{K_0^2} - \frac{3L_{\text{sym},4}}{K_0} - \frac{3J_0 L^2}{2K_0^3} \right) \rho_0 \quad (25)$$

which reflects the shift of the saturation density of asymmetric nuclear matter relative to that of symmetric nuclear matter due to the finite isospin asymmetry.

**C. Binding energy at saturation density**

Another basic quantity of asymmetric nuclear matter is the binding energy per nucleon at saturation density, i.e., $E_{\text{sat}}(\delta)$, and it is generally a function of the isospin asymmetry $\delta$. 
According to Eq. (11) and Eq. (23), \( E_{\text{sat}}(\delta) \) can be expressed up to 4th-order in \( \delta \) as

\[
E_{\text{sat}}(\delta) = E_0(\rho_0) + \frac{K_0}{2 \times 9} \left( \frac{\rho_{\text{sat}} - \rho_0}{\rho_0} \right)^2 \\
+ \left[ E_{\text{sym}}(\rho_0) + \frac{L}{3} \left( \frac{\rho_{\text{sat}} - \rho_0}{\rho_0} \right) \right] \delta^2 + E_{\text{sym,4}}(\rho_0) \delta^4 \\
= E_0(\rho_0) + \frac{1}{2 \times 9} \frac{9L^2}{K_0} \delta^4 + E_{\text{sym}}(\rho_0) \delta^2 \\
- \frac{L}{3} \frac{3L}{K_0} \delta^4 + E_{\text{sym,4}}(\rho_0) \delta^4 + O(\delta^6) \\
= E_0(\rho_0) + E_{\text{sym}}(\rho_0) \delta^2 + \left( E_{\text{sym,4}}(\rho_0) - \frac{L^2}{2K_0} \right) \delta^4 \\
+ O(\delta^6),
\]

Therefore, \( E_{\text{sat}}(\delta) \) can be written as

\[
E_{\text{sat}}(\delta) = E_0(\rho_0) + E_{\text{sat,2}} \delta^2 + E_{\text{sat,4}} \delta^4 + O(\delta^6),
\]

with

\[
E_{\text{sat,2}} = E_{\text{sym}}(\rho_0) \quad (28) \\
E_{\text{sat,4}} = E_{\text{sym,4}}(\rho_0) - \frac{L^2}{2K_0}. \quad (29)
\]

The binding energy of asymmetric nuclear matter at saturation density is thus shifted from that of symmetry nuclear matter as a result of the shift of the saturation density when the isospin asymmetry is finite.

**D. Incompressibility at saturation density**

The incompressibility of asymmetric nuclear matter is an important quantity to describe its properties. It depends on the density and isospin asymmetry of the asymmetric nuclear matter and is characterized by the incompressibility coefficient defined as

\[
K(\rho, \delta) = 9 \frac{\partial P(\rho, \delta)}{\partial \rho} \\
= 18 \rho \frac{\partial E(\rho, \delta)}{\partial \rho} + 9\rho^2 \frac{\partial^2 E(\rho, \delta)}{\partial \rho^2} \\
= 18 \frac{P(\rho, \delta)}{\rho} + 9\rho^2 \frac{\partial^2 E(\rho, \delta)}{\partial \rho^2},
\]

10
where $P(\rho, \delta)$ is the pressure of asymmetric nuclear matter and it can be expressed as

$$P(\rho, \delta) = \rho^2 \frac{\partial E(\rho, \delta)}{\partial \rho}. \quad (31)$$

Conventionally, the incompressibility coefficient is defined at the saturation density where $P(\rho, \delta) = 0$. It is also called the isobaric incompressibility coefficient \[74\], and is given by

$$K_{\text{sat}}(\delta) = 9\rho_{\text{sat}}^2 \frac{\partial^2 E(\rho, \delta)}{\partial \rho^2} |_{\rho = \rho_{\text{sat}}}. \quad (32)$$

The isobaric incompressibility coefficient $K_{\text{sat}}(\delta)$ thus only depends on the isospin asymmetry $\delta$. One can show (see Appendix \[a\]) that up to 4th-order in $\delta$, the isobaric incompressibility coefficient $K_{\text{sat}}(\delta)$ can be expressed as

$$K_{\text{sat}}(\delta) = K_0 + (K_{\text{sym}} - 6L - \frac{J_0}{K_0} L)\delta^2 + (K_{\text{sym},4} - 6L_{\text{sym},4} - \frac{J_0 L_{\text{sym},4}}{K_0} + 9L^2 - \frac{J_{\text{sym}} L}{K_0} + \frac{I_0 L^2}{2K_0^2} + \frac{J_0 K_{\text{sym}} L}{K_0^2} + \frac{3J_0 L^2}{K_0^2} - \frac{J_{\text{sym}} L_{\text{sym}} L}{2K_0^2})\delta^4 + O(\delta^6), \quad (33)$$

which can be further written as

$$K_{\text{sat}}(\delta) = K_0 + K_{\text{sat},2}\delta^2 + K_{\text{sat},4}\delta^4 + O(\delta^6) \quad (34)$$

with

$$K_{\text{sat},2} = K_{\text{sym}} - 6L - \frac{J_0}{K_0} L, \quad (35)$$

$$K_{\text{sat},4} = K_{\text{sym},4} - 6L_{\text{sym},4} - \frac{J_0 L_{\text{sym},4}}{K_0} + 9L^2 - \frac{J_{\text{sym}} L}{K_0} + \frac{I_0 L^2}{2K_0^2} + \frac{J_0 K_{\text{sym}} L}{K_0^2} + \frac{3J_0 L^2}{K_0^2} - \frac{J_{\text{sym}} L_{\text{sym}} L}{2K_0^2}. \quad (36)$$

The coefficient $K_{\text{sat},2}$ and $K_{\text{sat},4}$ reflect the isospin dependence of the isobaric incompressibility of asymmetric nuclear matter.

If we use the parabolic approximation for the EOS of symmetric nuclear matter, i.e., Eq. \[10\], then the $K_{\text{sat},2}$ parameter is reduced to

$$K_{\text{asy}} = K_{\text{sym}} - 6L, \quad (37)$$
and this expression has been extensively used in the literature to characterize the isospin dependence of the incompressibility of asymmetric nuclear matter. Obviously, we have

\[ K_{\text{sat},2} = K_{\text{asy}} - \frac{J_0}{K_0} L, \]  

(38)

and the coefficient \( K_{\text{asy}} \) thus could be a good approximation to \( K_{\text{sat},2} \) if the 3rd-order derivative of the EOS of symmetric nuclear matter with respect to density, i.e., \( J_0 \), is negligible or the magnitude of the slope parameter \( L \) for the symmetry energy is very small. It should be noted that the higher-order derivatives of the EOS of symmetric nuclear matter with respect to density in Eq. (4), such as \( I_0 \), do not contribute to \( K_{\text{sat},2} \). In the following, we will check how the \( J_0 \) term affects the value of \( K_{\text{sat},2} \).

As shown in Appendix, the expressions Eq. (23) for the saturation density, Eq. (27) for the binding energy and Eq. (34) for the isobaric incompressibility coefficient are exact up to 4th-order in \( \delta \). It is thus interesting to see that with precision up to 4th-order in \( \delta \), we only need to know 11 coefficients \( E_0(\rho_0), K_0, J_0, I_0, E_{\text{sym}}(\rho_0), L, K_{\text{sym}}, E_{\text{sym},4}(\rho_0), L_{\text{sym},4}, K_{\text{sym},4} \) among the 14 coefficients \( E_0(\rho_0), K_0, J_0, I_0, E_{\text{sym}}(\rho_0), L, K_{\text{sym}}, J_{\text{sym}}, I_{\text{sym}}, E_{\text{sym},4}(\rho_0), L_{\text{sym},4}, K_{\text{sym},4}, J_{\text{sym},4}, I_{\text{sym},4} \) which are defined at the normal nuclear density \( \rho_0 \). The higher-order coefficient \( I_{\text{sym}} \) in Eq. (11) for the symmetry energy as well as \( J_{\text{sym},4} \) and \( L_{\text{sym},4} \) in Eq. (16) for the 4th-order symmetry energy do not affect the saturation density, the binding energy and the isobaric incompressibility coefficient with precision up to 4th-order in \( \delta \). Furthermore, the 6 coefficients \( E_0(\rho_0), K_0, J_0, E_{\text{sym}}(\rho_0), L, \) and \( K_{\text{sym}} \) determine completely the saturation density, the binding energy and the isobaric incompressibility coefficient with precision up to 2nd-order in \( \delta \).

III. THEORETICAL MODELS

In this section, we introduce the theoretical models used in the present work and also give some important expressions for completeness. These models include the modified finite-range Gogny effective interaction (MDI), the Hartree-Fock approach based on Skyrme interactions, and a phenomenological modified Skyrme-like (MSL) model. A very useful feature of these models is that analytical expressions for many interesting physical quantities in asymmetric nuclear matter at zero temperature can be obtained, and this makes it phys-
ically transparent and very convenient to check the higher-order effects on the properties of asymmetric nuclear matter.

A. Isospin- and momentum-dependent MDI interaction

The isospin- and momentum-dependent MDI interaction is based on the finite-range Gogny effective interaction [43]. In the MDI interaction, the potential energy density \( V(\rho, \delta) \) of an asymmetric nuclear matter at total density \( \rho \) and isospin asymmetry \( \delta \) is expressed as follows [43, 55],

\[
V(\rho, \delta) = A_u \rho_n\rho_p + \frac{A_l}{2\rho_0} (\rho_n^2 + \rho_p^2) + \frac{B}{\sigma + 1} \rho^\sigma + 1 \rho_0^\sigma \\
\times (1 - x\delta^2) + \frac{1}{\rho_0} \sum_{\tau, \tau'} C_{\tau, \tau'} \\
\times \int \int d^3 p d^3 p' f_\tau(\vec{r}, \vec{p}) f_\tau'(\vec{r}, \vec{p}') \frac{1}{1 + (\vec{p} - \vec{p}')^2 / \Lambda^2}.
\]

(39)

In the above, \( \tau = 1/2 \) \((-1/2)\) for neutrons (protons); \( \sigma = 4/3 \); \( f_\tau(\vec{r}, \vec{p}) \) is the phase-space distribution function of neutrons or protons at coordinate \( \vec{r} \) and momentum \( \vec{p} \). The parameters \( A_u(x), A_l(x), B, C_{\tau, \tau}, C_{\tau, -\tau} \) and \( \Lambda \) are obtained by fitting the momentum-dependence of the single-nucleon potential \( U(\rho, \delta, \vec{p}, \tau) \) to that predicted by the Gogny Hartree-Fock and/or the Brueckner-Hartree-Fock calculations, the saturation properties of symmetric nuclear matter, and the symmetry energy at the normal nuclear matter density \( \rho_0 \) [43]. The parameters \( A_u(x) \) and \( A_l(x) \) are given by

\[
A_u(x) = -95.98 - x \frac{2B}{\sigma + 1}, \quad A_l(x) = -120.57 + x \frac{2B}{\sigma + 1}
\]

(40)

in terms of the parameter \( x \), which is introduced to mimic various \( E_{\text{sym}}(\rho) \) predicted by different microscopic and/or phenomenological many-body theories [55]. By adjusting the \( x \) parameter, the \( E_{\text{sym}}(\rho) \) is varied without changing any property of symmetric nuclear matter and the symmetry energy at the saturation density \( E_{\text{sym}}(\rho_0) \), as the \( x \)-dependent \( A_u(x) \) and \( A_l(x) \) are adjusted accordingly. Using the definition in Eq. (2), we have \( E_{\text{sym}}(\rho_0) = 30.5 \) MeV at \( \rho_0 = 0.16 \) fm\(^{-3} \) while its value is 31.6 MeV within the parabolic approximation of \( E_{\text{sym}}(\rho) \approx E(\rho, \delta = 1) - E(\rho, \delta = 0) \). We note that the MDI interaction has been extensively used in the transport model for studying isospin effects in intermediate energy heavy-ion collisions induced by neutron-rich nuclei [14, 15, 44, 46, 53, 56, 79, 80, 81, 82, 83, 84], the
study of the thermal properties of asymmetric nuclear matter \[85, 86, 87\], and the transition density and pressure in neutron star crust \[72, 73\]. In particular, the isospin diffusion data from NSCL/MSU have constrained the value of \(x\) to between 0 and \(-1\) for nuclear densities from about \(0.3\rho_0\) to \(1.2\rho_0\) \[55, 56\].

With \(f_\tau(\vec{r}, \vec{p}) = \frac{2}{\hbar^3} \Theta(p_f(\tau) - p)\) for nuclear matter at zero temperature, the integrals in Eq. (39) can be evaluated analytically. In particular, we have \[42\]

\[
\int \int d^3p d^3p' \frac{f_\tau(\vec{r}, \vec{p}) f_\tau'(\vec{r}, \vec{p}')}{1 + (\vec{p} - \vec{p}')^2/\Lambda^2} = \frac{1}{6} \left(\frac{4\pi}{\hbar^3}\right)^2 \Lambda^2 \left\{ p_f(\tau) p_f(\tau') [3(p_f^2(\tau) + p_f^2(\tau')) - \Lambda^2] \right. \\
+ 4\Lambda (p_f^3(\tau) - p_f^3(\tau')) \arctan \frac{p_f(\tau) - p_f(\tau')}{\Lambda} \\
- (p_f^3(\tau) + p_f^3(\tau')) \arctan \frac{p_f(\tau) + p_f(\tau')}{\Lambda} \\
\left. + \frac{1}{4} \left[ \Lambda^4 + 6\Lambda^2(p_f^2(\tau) + p_f^2(\tau')) - 3(p_f^2(\tau) - p_f^2(\tau'))^2 \right] \right\}.
\]

The kinetic energy per nucleon of cold asymmetric nuclear matter is

\[
E_k(\rho, \delta) = \frac{1}{\rho} \int d^3p \left( \frac{p_n^2}{2m} f_n(\vec{r}, \vec{p}) + \frac{p_p^2}{2m} f_p(\vec{r}, \vec{p}) \right) \\
= \frac{4\pi}{5m\hbar^3} \rho (p_{fn}^5 + p_{fp}^5), 
\]

where \(p_{fn(p)} = \hbar(3\pi^2 \rho_{n(p)})^{1/3}\) is the Fermi momentum of neutrons (protons). The total energy per nucleon of cold asymmetric nuclear matter can be obtained as

\[
E(\rho, \delta) = \frac{V(\rho, \delta)}{\rho} + E_k(\rho, \delta).
\]

By setting \(\rho_n = \rho_p = \frac{\rho}{2}\) and \(p_{fn} = p_{fp} = p_f\), where \(p_f = \hbar(3\pi^2 \rho/2)^{1/3}\) is the fermi momentum of symmetric nuclear matter at density \(\rho\), we obtain following EOS for cold symmetric nuclear matter:
\[
E_0(\rho) = \frac{8\pi}{5m\hbar^3\rho} p_f^5 + \frac{\rho}{4\rho_0} (A_l + A_u) + \frac{B}{\sigma + 1} \left( \frac{\rho}{\rho_0} \right)^\sigma \\
+ \frac{1}{3\rho_0} (C_l + C_u) \left( \frac{4\pi}{\hbar^3} \right)^2 \Lambda^2 \\
\times \left\{ p_f^2 (6p_f^2 - \Lambda^2) - 8\Lambda p_f^3 \arctan \frac{2p_f}{\Lambda} \\
+ \frac{1}{4} \left[ \Lambda^4 + 12\Lambda^2 p_f^2 \right] \ln \frac{4p_f^2 + \Lambda^2}{\Lambda^2} \right\}. 
\]

Furthermore, from Eq. (43) the symmetry energy can be expressed as

\[
E_{\text{sym}}(\rho) = 1 \left( \frac{\partial^2 E}{\partial \delta^2} \right)_{\delta=0} \\
= \frac{8\pi}{9m\hbar^3\rho} p_f^5 + \frac{\rho}{4\rho_0} (A_l - A_u) - \frac{Bx}{\sigma + 1} \left( \frac{\rho}{\rho_0} \right)^\sigma \\
+ \frac{C_l}{9\rho_0\rho} \left( \frac{4\pi}{\hbar^3} \right)^2 \Lambda^2 \left[ 4p_f^4 - \Lambda^2 p_f^2 \ln \frac{4p_f^2 + \Lambda^2}{\Lambda^2} \right] \\
+ \frac{C_u}{9\rho_0\rho} \left( \frac{4\pi}{\hbar^3} \right)^2 \Lambda^2 \left[ 4p_f^4 - p_f^2 (4p_f^2 + \Lambda^2) \ln \frac{4p_f^2 + \Lambda^2}{\Lambda^2} \right], \quad (45)
\]

and the 4th-order nuclear symmetry energy can be written as

\[
E_{\text{sym,4}}(\rho) = \frac{1}{4!} \left( \frac{\partial^4 E}{\partial \delta^4} \right)_{\delta=0} \\
= \frac{8\pi}{3^5 m\hbar^3\rho} p_f^5 \\
- \frac{C_l}{3^5 \rho_0 \rho} \left( \frac{4\pi}{\hbar^3} \right)^2 \Lambda^2 \left[ 7\Lambda^2 p_f^2 \ln \frac{4p_f^2 + \Lambda^2}{\Lambda^2} - \frac{4(7\Lambda^4 p_f^4 + 42\Lambda^2 p_f^6 + 40p_f^8)}{(4p_f^2 + \Lambda^2)^2} \right] \\
- \frac{C_u}{3^5 \rho_0 \rho} \left( \frac{4\pi}{\hbar^3} \right)^2 \Lambda^2 \left[ (7\Lambda^2 p_f^2 + 16p_f^4) \ln \frac{4p_f^2 + \Lambda^2}{\Lambda^2} - 28p_f^4 - \frac{8p_f^6}{\Lambda^2} \right]. \quad (46)
\]

### B. Skyrme-Hartree-Fock approach

In the standard Skyrme Hartree-Fock model [64, 88, 89, 90, 91, 92, 93, 94, 95, 96], the interaction is taken to have a zero-range, density- and momentum-dependent form with the interaction parameters determined from fitting the binding energies and charge radii of a large number of nuclei in the periodic table. In this approach, the EOS of asymmetric
nuclear matter can be expressed as \[\{91, 92, 95, 96\}\]

\[E(\rho, \delta) = \frac{3\hbar^2}{10m} \left( \frac{3\pi^2}{2} \right)^{2/3} \rho^{\frac{2}{3}} F_{5/3} + \frac{1}{8} t_0 [2x_0 + 2] - (2x_0 + 1) F_2 + \frac{1}{48} t_3 \rho^{\sigma+1} [2x_3 + 2] - (2x_3 + 1) F_2 + \frac{3}{40} \left( \frac{3\pi^2}{2} \right)^{2/3} \rho^{\frac{2}{3}} F_5 + 1 + 8 \left( 3 \times 10^{-6} \right) \frac{\Theta_s}{10} \left( \frac{3\pi^2}{2} \right)^{2/3} \rho^{\frac{2}{3}} F_{5/3} + 1 + 8 \left( 3 \times 10^{-6} \right) \frac{\Theta_s}{10} \left( \frac{3\pi^2}{2} \right)^{2/3} \rho^{\frac{2}{3}} F_5, \] (47)

where

\[F_m(\delta) = \frac{1}{2} \left[ (1 + \delta)^m + (1 - \delta)^m \right].\]

The EOS of symmetric nuclear matter can thus be written as

\[E_0(\rho) = \frac{3\hbar^2}{10m} \left( \frac{3\pi^2}{2} \right)^{2/3} \rho^{\frac{2}{3}} + \frac{3}{8} t_0 \rho + \frac{3}{80} \Theta_s \left( \frac{3\pi^2}{2} \right)^{2/3} \rho^{\frac{2}{3}} + \frac{1}{16} t_3 \rho^{\sigma+1}, \] (48)

with \(\Theta_s = 3t_1 + (5 + 4x_2)t_2\). Furthermore, the symmetry energy can be obtained as

\[E_{\text{sym}}(\rho) = \frac{1}{2} \left( \frac{\partial^2 E}{\partial \delta^2} \right)_{\delta=0} = \frac{\hbar^2}{6m} \left( \frac{3\pi^2}{2} \right)^{2/3} \rho^{\frac{2}{3}} - \frac{1}{8} t_0 (2x_0 + 1) \rho - \frac{1}{24} \left( \frac{3\pi^2}{2} \right)^{2/3} \Theta_{\text{sym}} \rho^{\frac{2}{3}} - \frac{1}{48} t_3 (2x_3 + 1) \rho^{\sigma+1}, \] (49)

with \(\Theta_{\text{sym}} = 3t_1 x_1 - t_2 (4 + 5x_2)\). Similarly, the 4th-order nuclear symmetry energy can be written as

\[E_{\text{sym,4}}(\rho) = \frac{1}{4!} \left( \frac{\partial^4 E}{\partial \delta^4} \right)_{\delta=0} = \frac{\hbar^2}{162m} \left( \frac{3\pi^2}{2} \right)^{2/3} \rho^{\frac{2}{3}} + \frac{1}{648} \left( \frac{3\pi^2}{2} \right)^{2/3} \Theta_{\text{sym,4}} \rho^{\frac{2}{3}}, \]
with \( \Theta_{\text{sym},4} = t_1(1 - x_1) + 3t_2(1 + x_2) \).

C. A phenomenological modified Skyrme-like (MSL) model

Following the energy density functional obtained from the Hartree-Fock approach with the zero-range, density- and momentum-dependent form with the Skyrme interaction, the binding energy per nucleon of a cold asymmetric nuclear matter at total density \( \rho \) and isospin asymmetry \( \delta \) in the modified Skyrme-like (MSL) model is parameterized as

\[
E_{\text{MSL}}(\rho, \delta) = \frac{\eta}{\rho} \left( \frac{\hbar^2}{2m_n^*} \rho^{5/3}_n + \frac{\hbar^2}{2m_p^*} \rho^{5/3}_p \right) \\
+ \frac{\alpha}{2} \frac{\rho}{\rho_0} + \frac{\beta}{\sigma + 1} \frac{\rho^\sigma}{\rho_0^\sigma} + E_{\text{loc}}^{\text{sym}}(\rho) \delta^2, \tag{50}
\]

where \( \eta = \frac{3}{5} (3\pi^2)^{2/3} \); \( \alpha, \beta \) and \( \sigma \) are parameters; and \( E_{\text{loc}}^{\text{sym}}(\rho) \) represents the local density-dependent part of the symmetry energy. The effective neutron and proton masses \( m_n^* \) and \( m_p^* \) are assumed to have forms similar to the standard SHF results \([91, 92, 95, 96]\), i.e.,

\[
\frac{\hbar^2}{2m_n^*} = \frac{\hbar^2}{2m} + \rho (C_{\text{eff}} + D_{\text{eff}} \delta) \tag{51}
\]

\[
\frac{\hbar^2}{2m_p^*} = \frac{\hbar^2}{2m} + \rho (C_{\text{eff}} - D_{\text{eff}} \delta), \tag{52}
\]

where \( C_{\text{eff}} \) and \( D_{\text{eff}} \) are constants. This implies that the single-nucleon potential depends quadratically on the nucleon momentum as in the standard SHF approach. In terms of the isoscalar effective mass \( m_s^* \) and the isovector effective mass \( m_v^* \) given by

\[
\frac{\hbar^2}{2m_s^*} = \frac{\hbar^2}{2m} + \rho C_{\text{eff}} \tag{53}
\]

\[
\frac{\hbar^2}{2m_v^*} = \frac{\hbar^2}{2m} + \rho (C_{\text{eff}} - D_{\text{eff}}), \tag{54}
\]

the nucleon effective mass can be written as \([97]\)

\[
\frac{\hbar^2}{2m_q^*} = \frac{2\rho_q}{\rho} \frac{\hbar^2}{2m_s^*} + \left( 1 - \frac{2\rho_q}{\rho} \right) \frac{\hbar^2}{2m_v^*}, \quad q = n, p. \tag{55}
\]

We note that the isovector effective mass \( m_v^* \) corresponds to the proton (neutron) effective mass in pure neutron (proton) matter. Also, we can easily obtain the following relation

\[
\frac{\hbar^2}{2m_n^*} - \frac{\hbar^2}{2m_p^*} = 2\delta \left( \frac{\hbar^2}{2m_s^*} - \frac{\hbar^2}{2m_v^*} \right). \tag{56}
\]
Experimentally, the isoscalar effective mass $m^*_s$ and the isovector effective mass $m^*_v$ at the normal nuclear density $\rho_0$ have been constrained to be $m^*_{s,0} \approx 0.8m$ and $m^*_{v,0} \approx 0.7m$, respectively [42, 91, 92, 98, 99, 100, 101, 102, 103]. With these constrained values for the isoscalar and isovector effective masses, Eq. (56) gives a larger neutron effective mass than the proton effective mass in a neutron-rich matter, which is consistent with experimental data on the isospin dependence of the nucleon optical potential and also recent microscopic and phenomenological many-body theory predictions [15, 33, 45].

The EOS of symmetric nuclear matter in the MSL model is then given by

$$E_0(\rho) = E^0_{\text{kin}} \left( \frac{\rho}{\rho_0} \right)^{2/3} + C \left( \frac{\rho}{\rho_0} \right)^{5/3} + \frac{\alpha}{2} \frac{\rho}{\rho_0} + \frac{\beta}{\sigma + 1} \left( \frac{\rho}{\rho_0} \right)^{\sigma},$$

where the first term represents the kinetic energy contribution with $E^0_{\text{kin}} = 3\hbar^2 \frac{m}{16m} \left( \frac{3\pi^2}{2} \right)^{2/3} \rho_0^{2/3}$ and the second term is due to the nucleon effective mass with the coefficient $C$ being a constant determined by the isoscalar effective mass $m^*_{s,0}$ as

$$C = \frac{m - m^*_{s,0}}{m^*_{s,0}} E^0_{\text{kin}}.$$  

The parameters $\alpha$, $\beta$ and $\sigma$ in the MSL model are determined by the binding energy per nucleon $E_0(\rho_0)$ and the incompressibility $K_0$ of cold symmetric nuclear matter at the saturation density $\rho_0$, and they can be expressed as

$$\alpha = -\frac{4}{3} E^0_{\text{kin}} - \frac{10}{3} C - \frac{2}{3} (E^0_{\text{kin}} - 3E_0(\rho_0) - 2C) \times \frac{K_0 + 2E^0_{\text{kin}} - 10C}{K_0 + 9E_0(\rho_0) - E^0_{\text{kin}} - 4C},$$

$$\beta = \left( E^0_{\text{kin}} - E_0(\rho_0) - \frac{2}{3} C \right) \times \frac{K_0 + 9E_0(\rho_0) + 5E^0_{\text{kin}} - 16C}{K_0 + 9E_0(\rho_0) - E^0_{\text{kin}} - 4C},$$

$$\sigma = \frac{K_0 + 2E^0_{\text{kin}} - 10C}{3E^0_{\text{kin}} - 9E_0(\rho_0) - 6C}.$$
In particular, for \( E_0(\rho_0) = -16 \text{ MeV}, \) \( m_{s,0}^* = 0.8m \) and \( \rho_0 = 0.16 \text{ fm}^{-3}, \) we have

\[
C = 5.53 \text{ (MeV)} \\
\alpha = -47.90 - 39.37 \frac{K_0 - 11.05}{K_0 - 188.21} \text{ (MeV)} \\
\beta = 19.68 \frac{K_0 + 166.11}{K_0 - 188.21} \text{ (MeV)} \\
\sigma = \frac{K_0 - 11.05}{177.16} ,
\]

where the units of \( K_0 \) is MeV.

The symmetry energy in the MSL model can be expressed as

\[
E_{\text{sym}}(\rho) = E_{\text{sym}}^{\text{kin}}(\rho_0) \left( \frac{\rho}{\rho_0} \right)^{2/3} + D \left( \frac{\rho}{\rho_0} \right)^{5/3} + E_{\text{sym}}^{\text{loc}}(\rho) ,
\]

where the first term is the kinetic energy contribution with \( E_{\text{sym}}^{\text{kin}}(\rho_0) = \frac{\hbar^2}{6m} \left( \frac{3\pi^2}{2} \rho_0 \right)^{2/3} \) and the second term is due to the contribution of the nucleon effective mass with the coefficient \( D \) determined by both \( m_{s,0}^* \) and \( m_{v,0}^* \) as

\[
D = \frac{5}{9} E_0^{\text{kin}} \left( 4 \frac{m}{m_{s,0}^*} - 3 \frac{m}{m_{v,0}^*} - 1 \right) .
\]

With \( m_{s,0}^* = 0.8m \) and \( m_{v,0}^* = 0.7m \) at \( \rho_0 = 0.16 \text{ fm}^{-3}, \) we have \( D = -3.51 \text{ MeV}. \) In Eq. (66), similarly to the momentum-independent MID model \[51\], the local density-dependent part \( E_{\text{sym}}^{\text{loc}}(\rho) \) is parameterized as

\[
E_{\text{sym}}^{\text{loc}}(\rho) = (1 - y) E_{\text{sym}}^{\text{loc}}(\rho_0) \frac{\rho}{\rho_0} + y E_{\text{sym}}^{\text{loc}}(\rho_0) \gamma_{\text{sym}} \left( \frac{\rho}{\rho_0} \right)^{\gamma_{\text{sym}}}
\]

with the constant \( E_{\text{sym}}^{\text{loc}}(\rho_0) \) determined by

\[
E_{\text{sym}}^{\text{loc}}(\rho_0) = E_{\text{sym}}(\rho_0) - E_{\text{sym}}^{\text{kin}}(\rho_0) - D .
\]

Obviously, we have \( E_{\text{sym}}^{\text{loc}}(\rho_0) = 21.2 \text{ MeV} \) following \( E_{\text{sym}}(\rho_0) = 30 \text{ MeV} \) and \( E_{\text{sym}}^{\text{kin}}(\rho_0) = 12.3 \text{ MeV} \) at \( \rho_0 = 0.16 \text{ fm}^{-3}. \) The default value for the \( \gamma_{\text{sym}} \) parameter is taken to be \( 4/3 \) in the MSL model following the \( E_{\text{sym}}(\rho) \) in the MDI interaction, namely, Eq. (45) (we will see how the \( \gamma_{\text{sym}} \) parameter affects the symmetry energy in the following). In particular, similarly to the \( x \) parameter in the MDI interaction, the dimensionless parameter \( y \) in the MSL model is introduced to mimic various \( E_{\text{sym}}(\rho) \) predicted by different microscopic and/or phenomenological many-body theories for a fixed \( \gamma_{\text{sym}} \) parameter. As we will show later, for
\( \gamma_{\text{sym}} = 4/3 \), adjusting the \( y \) value can nicely reproduce the \( E_{\text{sym}}(\rho) \) in the MDI interaction with \( x = -1, 0, \) and \( 1 \). Therefore, the symmetry energy density functional constructed in the MSL model is very flexible and can mimic very different density behaviors by varying only one parameter.

In the MSL model, similarly to the SHF approach, the 4th-order and higher-order nuclear symmetry energies only include contributions from the kinetic energy and the nucleon effective mass while the local density-dependent part of higher-order symmetry energies are neglected in Eq. (50). In particular, the 4th-order nuclear symmetry energy in the MSL model can be shown to be

\[
E_{\text{sym},4}(\rho) = \frac{1}{4!} \left( \frac{\partial^4 E}{\partial \delta^4} \right)_{\delta=0} = \frac{5}{243} E_{\text{kin}}^0 \left( \frac{\rho}{\rho_0} \right)^{2/3} + \frac{5}{243} C \left( \frac{\rho}{\rho_0} \right)^{5/3}, \tag{70}
\]

where the first term is the kinetic energy contribution while the second term is due to the contribution from the nucleon effective mass.

The MSL model is thus an extension of the momentum-independent MID model \([51]\) by including the effects of the nucleon effective mass. It provides a simple phenomenological parametrization of the EOS of asymmetric nuclear matter and is thus a convenient and transparent way to investigate the possible correlations among higher-order and lower-order characteristic parameters of asymmetric nuclear matter. In the MSL model, we have totally 8 free parameters, i.e., \( C, D, \alpha, \beta, \sigma, E_{\text{sym}}(\rho_0), y, \) and \( \gamma_{\text{sym}} \) which can be determined by empirical information on the EOS of symmetric nuclear matter, the nucleon effective mass and the density dependence of symmetry energy. In particular, the parameters \( C \) and \( D \) (or equivalently \( C_{\text{eff}} \) and \( D_{\text{eff}} \)) are determined by the isoscalar effective masses \( m_s^* \) and the isovector effective mass \( m_v^* \) at the normal nuclear density \( \rho_0 \), i.e., \( m_{s,0}^* \) and \( m_{v,0}^* \). The parameters \( \alpha, \beta \) and \( \sigma \) are determined by \( E_0(\rho_0) \), \( K_0 \) and \( \rho_0 \) while the parameters \( E_{\text{sym}}(\rho_0) \), \( y \), and \( \gamma_{\text{sym}} \) are introduced to mimic the density dependence of different symmetry energies predicted by microscopic and/or phenomenological many-body theories. As a default in the MSL model, we use \( m_{s,0}^* = 0.8m, m_{v,0}^* = 0.7m, \rho_0 = 0.16 \text{ fm}^{-3}, E_0(\rho_0) = -16 \text{ MeV}, \) \( K_0 = 240 \text{ MeV}, E_{\text{sym}}(\rho_0) = 30 \text{ MeV}, \) and \( \gamma_{\text{sym}} = 4/3 \) and vary the parameter \( y \) to describe different symmetry energies.
IV. RESULTS AND DISCUSSIONS

A. Characteristic parameters at normal nuclear density and EOS of asymmetric nuclear matter

As shown in Section II, the expressions for the saturation density Eq. (23), the binding energy Eq. (27) and the isobaric incompressibility coefficient Eq. (34) are exact up to 4th-order in $\delta$, and these expressions involve 11 characteristic parameters defined at the normal nuclear density $\rho_0$, i.e., $E_0(\rho_0)$, $K_0$, $J_0$, $I_0$, $E_{\text{sym}}(\rho_0)$, $L$, $K_{\text{sym}}$, $J_{\text{sym}}$, $E_{\text{sym},4}(\rho_0)$, $L_{\text{sym},4}$, and $K_{\text{sym},4}$. We summarize the values of $\rho_0$, $E_0(\rho_0)$, $K_0$, $J_0$, $I_0$, $K_{\text{asy}}$, $K_{\text{sat},2}$ and $K_{\text{sat},2}$ in Tables I, II, and III while the values of $E_{\text{sym}}(\rho_0)$, $L$, $K_{\text{sym}}$, $J_{\text{sym}}$, $E_{\text{sym},4}(\rho_0)$, $L_{\text{sym},4}$, $K_{\text{sym},4}$, $m_{s,0}^*/m$ and $m_{v,0}^*/m$ in Tables IV, V, and VI for the MDI interaction with $x = 1, 0, -1$ and the popular 63 standard Skyrme interactions with their saturation density and the symmetry energy satisfying $0.140 \, \text{fm}^{-3} < \rho_0 < 0.165 \, \text{fm}^{-3}$ and $25 \, \text{MeV} < E_{\text{sym}}(\rho_0) < 37 \, \text{MeV}$, respectively. For the 63 standard Skyrme interactions, the values in the Tables are sorted in the order of increasing values of $L$. It should be stressed here that the parameters of all Skyrme interactions are chosen to fit the binding energies and charge radii of a large number of nuclei in the periodic table. Detailed values of the parameters for these 63 Skyrme interactions can be found in Refs. [50, 89, 90, 91, 92, 93, 95, 96, 104, 105, 106, 107, 108, 109, 110, 111, 112, 113]. The selected ranges of $\rho_0$ and $E_{\text{sym}}(\rho_0)$ are essentially consistent with their empirical values inferred from experimental data. We note that here no constraints are imposed on $K_0$ and $L$ for selecting the Skyrme interactions as we will systematically explore the correlations of other physical quantities with $K_0$ or $L$.

Since the 11 characteristic parameters are defined at the normal nuclear density $\rho_0$, it is of interest to study the extent these characteristic parameters can provide information on the properties of asymmetric nuclear at sub-saturation and supra-saturation density regions. As an example, we show in Fig. I the energy per nucleon of symmetry nuclear matter from the MDI interaction as a function of its density. Also shown in Fig. I are the results obtained by using Eq. (11) including terms up to $\chi^2$, $\chi^3$, and $\chi^4$, respectively. It is seen that Eq. (11) with terms up to $\chi^2$, i.e., the parabolic approximation Eq. (10) which involves only two characteristic parameters, i.e., $E_0(\rho_0)$ and $K_0$, can approximate very well the EOS of
TABLE I: The saturation density $\rho_0$ and the characteristic parameters $E_0(\rho_0)$ (MeV), $K_0$ (MeV), $J_0$ (MeV), $I_0$ (MeV), $K_{asy}$ (MeV), $K_{sat,2}$ (MeV) and $K_{sat,2}$ (MeV) at saturation density for the MDI interaction with $x = -1, 0, 1$ and the SHF predictions with 63 standard Skyrme interactions. The small differences from Table I of Ref. [73] for some Skyrme interactions are due to the use of 0.17 (0.33) as an approximation of 1/6 (1/3) for the $\sigma$ parameter in Ref. [73]. The results shown here are thus more accurate than those in Ref. [73].

| Force  | $\rho_0$ | $E_0(\rho_0)$ | $K_0$  | $J_0$  | $I_0$  | $K_{asy}$ | $K_{sat,2}$ | $K_{sat,4}$ |
|--------|----------|---------------|--------|--------|--------|-----------|-------------|-------------|
| MDI(1) | 0.160    | -16.2         | 212.5  | -447.6 | 2160.8 | -352.0    | -321.1      | -8.4        |
| MDI(0) | 0.160    | -16.2         | 212.5  | -447.6 | 2160.8 | -443.1    | -316.3      | 52.9        |
| MDI(-1)| 0.160    | -16.2         | 212.5  | -447.6 | 2160.8 | -534.3    | -311.4      | 214.4       |
| Z      | 0.159    | -16.0         | 330.3  | -65.0  | -348.2 | -359.6    | -369.4      | 100.5       |
| E$_\sigma$ | 0.163  | -16.0         | 248.6  | -352.4 | 1337.1 | -236.6    | -288.9      | 57.9        |
| E      | 0.159    | -16.1         | 333.5  | -63.0  | -356.3 | -383.1    | -389.1      | 41.1        |
| Z$_\sigma$ | 0.163  | -15.9         | 233.3  | -369.0 | 1546.0 | -225.1    | -271.6      | 43.3        |
| SVI    | 0.143    | -15.8         | 363.6  | 153.5  | -1107.4| -427.3    | -424.2      | -4.4        |
| Z$_\sigma^*$ | 0.162 | -16.0         | 234.9  | -369.2 | 1544.4 | -305.5    | -312.6      | -7.1        |
| SkSC4  | 0.161    | -15.9         | 234.7  | -380.8 | 1549.2 | -316.5    | -320.1      | -4.1        |
| SI     | 0.155    | -16.0         | 370.4  | 152.3  | -1129.5| -469.2    | -469.7      | -14.2       |
| BSk3   | 0.157    | -15.8         | 234.8  | -380.9 | 1529.8 | -347.6    | -336.6      | -9.2        |
| BSk1   | 0.157    | -15.8         | 231.3  | -385.6 | 1588.7 | -325.0    | -313        | -8.4        |
| SIII   | 0.145    | -15.9         | 355.4  | 101.4  | -903.0 | -453.2    | -456.0      | -20.2       |
| BSk2   | 0.157    | -15.8         | 233.7  | -380.1 | 1542.4 | -344.8    | -331.9      | -9.7        |
| MSk7   | 0.157    | -15.8         | 231.2  | -385.4 | 1587.3 | -331.0    | -315.4      | -8.6        |
| BSk4   | 0.157    | -15.8         | 236.8  | -367.2 | 1466.9 | -341.2    | -321.7      | -11.4       |
| BSk8   | 0.159    | -15.8         | 230.3  | -372.4 | 1578.2 | -310.0    | -286.0      | -16.8       |
| BSk6   | 0.157    | -15.8         | 229.1  | -370.6 | 1571.3 | -316.3    | -289.0      | -16.8       |
| BSk7   | 0.157    | -15.8         | 229.3  | -370.9 | 1572.8 | -317.3    | -288.2      | -17.0       |
| SKP    | 0.163    | -16.0         | 201.0  | -435.6 | 2127.8 | -384.3    | -341.9      | 7.4         |
| BSk5   | 0.157    | -15.8         | 237.2  | -367.9 | 1470.3 | -368.8    | -335.6      | -7.5        |
| SKXm   | 0.159    | -16.0         | 238.1  | -380.4 | 1542.2 | -435.3    | -384.0      | 10.3        |
| RATP   | 0.160    | -16.0         | 239.4  | -349.7 | 1451.5 | -385.5    | -338.2      | -6.6        |
| Force  | $\rho_0$ | $E_0(\rho_0)$ | $K_0$ | $J_0$ | $I_0$ | $K_{asy}$ | $K_{sat,2}$ | $K_{sat,A}$ |
|--------|---------|----------------|-------|-------|-------|------------|-------------|-------------|
| SKX    | 0.155   | -16.1          | 271.1 | -297.4| 904.0 | -451.2     | -414.8      | 2.3         |
| SKXce  | 0.156   | -15.9          | 268.2 | -294.6| 892.9 | -439.3     | -402.5      | 3.8         |
| BSk15  | 0.159   | -16.0          | 241.6 | -363.1| 1457.0| -395.9     | -345.4      | -2.7        |
| BSk16  | 0.159   | -16.1          | 241.7 | -363.6| 1459.9| -396.6     | -344.2      | -1.7        |
| BSk10  | 0.159   | -15.9          | 238.8 | -370.3| 1479.6| -418.3     | -360.6      | 12.7        |
| SGI   | 0.158   | -15.6          | 214.7 | -380.9| 1741.8| -371.7     | -304.9      | 17.8        |
| BSk12  | 0.159   | -15.9          | 238.1 | -369.1| 1474.9| -419.4     | -360.5      | 14.0        |
| BSk11  | 0.159   | -15.9          | 238.1 | -369.2| 1475.3| -420.0     | -360.5      | 14.6        |
| SLy10  | 0.156   | -15.9          | 229.7 | -358.3| 1559.6| -374.7     | -314.2      | -24.6       |
| BSk13  | 0.159   | -15.9          | 238.1 | -369.2| 1475.2| -420.8     | -360.6      | 15.5        |
| BSk9   | 0.159   | -15.9          | 231.4 | -374.9| 1591.5| -384.7     | -320.0      | -3.1        |
| BSk14  | 0.159   | -15.9          | 239.3 | -358.7| 1434.8| -415.5     | -349.7      | 14.2        |
| SLy230a| 0.160   | -16.0          | 229.9 | -364.2| 1593.6| -364.1     | -293.9      | -32.1       |
| SLy6   | 0.159   | -15.9          | 229.9 | -360.2| 1568.8| -383.7     | -312.9      | -13.2       |
| SLy8   | 0.160   | -16.0          | 229.9 | -363.2| 1587.1| -388.4     | -316.8      | -12.0       |
| SLy4   | 0.160   | -16.0          | 229.9 | -363.1| 1586.9| -392.1     | -320.5      | -12.7       |
| SLy0   | 0.161   | -16.0          | 230.2 | -365.2| 1598.7| -389.2     | -317.2      | -11.6       |
| SLy3   | 0.160   | -16.0          | 229.9 | -363.4| 1588.0| -395.4     | -323.4      | -12.9       |
| SKM*   | 0.160   | -15.8          | 216.6 | -386.1| 1768.9| -430.6     | -349.0      | 37.3        |
| SLy230b| 0.160   | -16.0          | 229.9 | -363.1| 1586.8| -395.5     | -322.9      | -12.2       |
| SLy7   | 0.158   | -15.9          | 229.7 | -359.2| 1562.9| -402.7     | -327.6      | -10.4       |
| SLy2   | 0.160   | -15.9          | 229.2 | -362.7| 1585.5| -406.1     | -328.9      | -7.0        |
| SLy1   | 0.160   | -16.0          | 229.8 | -364.3| 1594.4| -408.7     | -331.3      | -7.3        |
| SKM    | 0.160   | -15.8          | 216.6 | -386.1| 1768.9| -444.9     | -356.9      | 45.1        |
| SII    | 0.148   | -16.0          | 341.4 | 15.8 |-567.5 | -565.9     | -568.2      | 23.2        |
| SLy5   | 0.161   | -16.0          | 229.9 | -364.1| 1592.7| -413.7     | -334.0      | -4.6        |
TABLE III: Continued with Table II

| Force  | \(\rho_0\) | \(E_0(\rho_0)\) | \(K_0\) | \(J_0\) | \(I_0\) | \(K_{asy}\) | \(K_{sat,2}\) | \(K_{sat,4}\) |
|--------|--------|----------------|--------|--------|--------|------------|------------|------------|
| SLy9   | 0.151  | -15.8          | 229.8  | -350.4 | 1511.4 | -413.7     | -329.2     | 1.7        |
| SkI6   | 0.159  | -15.9          | 248.2  | -326.6 | 1251.8 | -402.2     | -324.3     | 15.5       |
| SkI4   | 0.160  | -15.9          | 248.0  | -331.2 | 1280.0 | -402.9     | -322.2     | 21.9       |
| SIV    | 0.151  | -16.0          | 324.6  | -68.8  | -234.9 | -517.7     | -504.2     | 64.8       |
| SGI    | 0.154  | -15.9          | 261.8  | -297.9 | 1005.0 | -435.2     | -362.5     | 61.8       |
| SkO*   | 0.160  | -15.7          | 222.1  | -390.2 | 1706.7 | -495.6     | -373.2     | 98.3       |
| SkMP   | 0.157  | -15.6          | 231.0  | -338.4 | 1425.2 | -468.9     | -366.8     | 101.5      |
| SKa    | 0.155  | -16.0          | 263.2  | -300.1 | 1014.4 | -526.2     | -441.1     | 105.0      |
| SKO    | 0.160  | -15.8          | 222.8  | -391.3 | 1712.3 | -519.1     | -379.5     | 143.9      |
| \(R_\sigma\) | 0.158  | -15.6          | 237.4  | -348.4 | 1377.2 | -523.3     | -397.5     | 184.5      |
| \(G_\sigma\) | 0.158  | -15.6          | 237.2  | -348.7 | 1379.5 | -550.1     | -411.9     | 230.6      |
| SKT4   | 0.159  | -16.0          | 235.5  | -383.0 | 1562.5 | -589.3     | -436.2     | 201.2      |
| SV     | 0.155  | -16.0          | 305.7  | -175.8 | 183.5  | -552.4     | -497.1     | 176.2      |
| SkI3   | 0.158  | -16.0          | 258.2  | -303.9 | 1088.3 | -530.1     | -411.8     | 152.1      |
| SkI2   | 0.158  | -15.8          | 240.9  | -339.7 | 1351.4 | -555.3     | -408.2     | 226.9      |
| SkI5   | 0.156  | -15.8          | 255.8  | -302.0 | 1083.7 | -616.4     | -463.7     | 347.2      |

symmetric nuclear matter from about \(0.5\rho_0\) to \(1.5\rho_0\). Including higher-order terms of \(\chi^3\) and \(\chi^4\) with the characteristic parameters \(J_0\) and \(I_0\) in Eq. \((\text{II})\) improves significantly the approximation to the EOS at low densities and that up to about \(2\rho_0\). To describe reasonably the EOS of symmetric nuclear matter above \(2\rho_0\), one needs to include higher-order terms in \(\chi\).

Fig. 2 displays the density dependence of symmetry energy using the MDI interaction with \(x = 1, 0,\) and \(-1\) together with corresponding results obtained by using Eq. \((\text{III})\) up to \(\chi, \chi^2, \chi^3,\) and \(\chi^4,\) respectively. It is seen that the importance of the contributions from higher-order terms in \(\chi\) to the density dependence of symmetry energy depends on the stiffness of symmetry energy. For a supra-soft symmetry energy \((x = 1)\), terms up to \(\chi^3\) are needed to describe reasonably the symmetry energy from sub-saturation densities to about \(2\rho_0\). The situation is similar for the case of modestly soft symmetry energy \((x = 0)\). For the
TABLE IV: The characteristic parameters $E_{\text{sym}}(\rho_0)$ (MeV), $L$ (MeV), $K_{\text{sym}}$ (MeV), $J_{\text{sym}}$ (MeV),
$E_{\text{sym},4}(\rho_0)$ (MeV), $L_{\text{sym,4}}$ (MeV), $K_{\text{sym,4}}$ (MeV), $m^*_{s,0}/m$ and $m^*_{v,0}/m$ at saturation density for the
MDI interaction with $x = -1, 0, 1$ and the SHF predictions with 63 standard Skyrme interactions.

| Force   | $E_{\text{sym}}(\rho_0)$ | $L$  | $K_{\text{sym}}$ | $J_{\text{sym}}$ | $E_{\text{sym},4}(\rho_0)$ | $L_{\text{sym,4}}$ | $K_{\text{sym,4}}$ | $m^*_{s,0}/m$ | $m^*_{v,0}/m$ |
|---------|--------------------------|------|------------------|------------------|-----------------------------|-------------------|------------------|----------------|----------------|
| MDI(1)  | 30.5                     | 14.7 | -264.0           | 660.0            | 0.62                        | 0.53              | -4.82            | 0.67           | 0.54           |
| MDI(0)  | 30.5                     | 60.2 | -81.7            | 295.3            | 0.62                        | 0.53              | -4.82            | 0.67           | 0.54           |
| MDI(-1) | 30.5                     | 105.8| 100.7            | -69.3            | 0.62                        | 0.53              | -4.82            | 0.67           | 0.54           |
| E       | 26.4                     | -36.9| -457.8           | 880.0            | 0.88                        | 3.04              | 2.38             | 0.84           | 0.70           |
| Z       | 26.8                     | -49.7| -657.9           | 495.2            | 0.78                        | 2.56              | 1.47             | 0.84           | 0.73           |
| $E_{\sigma}$ | 26.7         | -29.4| -401.4           | 883.1            | 0.91                        | 3.14              | 2.58             | 0.78           | 0.66           |
| SVI     | 26.9                     | -7.3 | -471.3           | 146.0            | 0.67                        | 2.08              | 0.76             | 0.95           | 0.81           |
| $Z_{\sigma}$ | 28.8            | -4.5 | -332.6           | 725.1            | 0.92                        | 3.24              | 2.79             | 0.77           | 0.65           |
| SkSC4   | 28.8                     | -2.2 | -329.5           | 708.3            | 0.46                        | 0.91              | -1.85            | 1.00           | 1.00           |
| SI      | 29.2                     | 1.2  | -461.8           | 141.4            | 0.70                        | 2.16              | 0.73             | 0.91           | 0.80           |
| BSk3    | 27.9                     | 6.8  | -306.9           | 550.3            | 0.71                        | 2.19              | 0.76             | 1.12           | 0.89           |
| BSk1    | 27.8                     | 7.2  | -281.8           | 606.4            | 0.43                        | 0.79              | -2.04            | 1.05           | 1.05           |
| BSk2    | 28.0                     | 8.0  | -297.0           | 557.9            | 0.71                        | 2.18              | 0.74             | 1.04           | 0.86           |
| MSk7    | 27.9                     | 9.4  | -274.6           | 592.1            | 0.43                        | 0.79              | -2.04            | 1.05           | 1.05           |
| SHII    | 28.1                     | 9.9  | -393.7           | 130.4            | 0.83                        | 2.89              | 2.34             | 0.76           | 0.66           |
| BSk4    | 28.0                     | 12.5 | -265.9           | 558.4            | 0.61                        | 1.70              | -0.22            | 0.92           | 0.85           |
| BSk8    | 28.0                     | 14.9 | -220.9           | 624.9            | 0.43                        | 0.78              | -2.09            | 0.80           | 0.87           |
| BSk6    | 28.0                     | 16.8 | -215.2           | 603.5            | 0.45                        | 0.89              | -1.85            | 0.80           | 0.86           |
| BSk7    | 28.0                     | 18.0 | -209.4           | 598.2            | 0.42                        | 0.77              | -2.08            | 0.80           | 0.87           |
| SKP     | 30.0                     | 19.6 | -266.8           | 508.6            | 0.94                        | 3.33              | 2.96             | 1.00           | 0.74           |
| BSk5    | 28.7                     | 21.4 | -240.3           | 499.9            | 0.64                        | 1.83              | 0.04             | 0.92           | 0.84           |
| SKXm    | 31.2                     | 32.1 | -242.8           | 428.7            | 0.88                        | 3.02              | 2.40             | 0.97           | 0.75           |
| RATP    | 29.2                     | 32.4 | -191.2           | 440.6            | 1.06                        | 3.94              | 4.21             | 0.67           | 0.56           |
| Force   | $E_{\text{sym}}(\rho_0)$ | $L$  | $K_{\text{sym}}$ | $J_{\text{sym}}$ | $E_{\text{sym,4}}(\rho_0)$ | $L_{\text{sym,4}}$ | $K_{\text{sym,4}}$ | $m^*_s/m^*_s,0/m$ |
|---------|-----------------|-----|----------------|---------------|-----------------|----------------|---------------|-----------------|
| SKX     | 31.1            | 33.2| -252.1        | 379.7         | 0.89            | 3.10           | 2.61          | 0.99            |
| SKXce   | 30.1            | 33.5| -238.4        | 356.9         | 0.89            | 3.12           | 2.65          | 1.01            |
| BSk15   | 30.0            | 33.6| -194.4        | 466.5         | 0.64            | 1.84           | 0.03          | 0.80            |
| BSk16   | 30.0            | 34.9| -187.4        | 461.9         | 0.60            | 1.66           | -0.31         | 0.80            |
| BSk10   | 30.0            | 37.2| -194.9        | 397.0         | 0.69            | 2.08           | 0.50          | 0.92            |
| SGII    | 26.8            | 37.6| -145.9        | 330.4         | 0.87            | 3.01           | 2.39          | 0.79            |
| BSk12   | 30.0            | 38.0| -191.4        | 392.5         | 0.68            | 2.02           | 0.40          | 0.92            |
| BSk11   | 30.0            | 38.4| -189.8        | 390.1         | 0.67            | 2.01           | 0.37          | 0.92            |
| SLy10   | 32.0            | 38.7| -142.2        | 591.2         | 0.37            | 0.50           | -2.60         | 0.68            |
| BSk13   | 30.0            | 38.8| -187.9        | 386.6         | 0.67            | 2.00           | 0.37          | 0.92            |
| BSk9    | 30.0            | 39.9| -145.3        | 475.8         | 0.37            | 0.47           | -2.71         | 0.80            |
| BSk14   | 30.0            | 43.9| -152.0        | 388.3         | 0.60            | 1.65           | -0.33         | 0.80            |
| SLy230a | 32.0            | 44.3| -98.2         | 602.9         | 0.06            | -1.07          | -5.80         | 0.70            |
| SLy6    | 31.2            | 45.2| -112.5        | 511.3         | 0.38            | 0.55           | -2.54         | 0.69            |
| SLy8    | 31.4            | 45.3| -116.5        | 511.4         | 0.40            | 0.64           | -2.39         | 0.70            |
| SLy4    | 31.8            | 45.4| -119.9        | 521.0         | 0.40            | 0.63           | -2.40         | 0.69            |
| SLy0    | 31.5            | 45.4| -116.8        | 510.6         | 0.40            | 0.64           | -2.39         | 0.70            |
| SLy3    | 32.1            | 45.5| -122.1        | 526.2         | 0.40            | 0.63           | -2.41         | 0.70            |
| SKM*    | 30.0            | 45.8| -155.9        | 330.5         | 0.94            | 3.32           | 2.97          | 0.79            |
| SLy230b | 32.0            | 46.0| -119.7        | 521.5         | 0.40            | 0.61           | -2.43         | 0.69            |
| SLy7    | 32.4            | 48.1| -114.3        | 516.6         | 0.38            | 0.55           | -2.53         | 0.69            |
| SLy2    | 32.3            | 48.8| -113.5        | 502.9         | 0.40            | 0.63           | -2.39         | 0.70            |
| SLy1    | 32.5            | 48.8| -115.7        | 508.5         | 0.40            | 0.63           | -2.40         | 0.70            |
| SKM     | 30.7            | 49.3| -148.8        | 323.3         | 0.91            | 3.19           | 2.71          | 0.79            |
| SII     | 34.2            | 50.0| -265.7        | 104.7         | 1.10            | 4.21           | 4.94          | 0.58            |
| SLy5    | 32.7            | 50.3| -111.9        | 499.2         | 0.40            | 0.63           | -2.40         | 0.70            |
stiffer symmetry energy \(x = -1\), including terms up to \(\chi^2\) already give a good description of the symmetry energy from sub-saturation densities to about 3\(\rho_0\).

In Fig. 3, we show the density dependence of the 4th-order symmetry energy \(E_{\text{sym,4}}(\rho)\) using the MDI interaction together with corresponding results obtained by using Eq. \(\chi_1, \chi_2, \chi_3,\) and \(\chi_4\), respectively. We note here that the \(E_{\text{sym,4}}(\rho)\) does not depend on the \(x\) parameter in the MDI interaction as shown in Eq. \(\chi_1\). Firstly, one can see from Fig. 3 that the magnitude of \(E_{\text{sym,4}}(\rho)\) is very small compared to that of \(E_{\text{sym}}(\rho)\). As shown in Table IV, the value of \(E_{\text{sym,4}}(\rho)\) at the normal nuclear density \(\rho_0\) is about 0.62 MeV which is consistent with the predictions from the SHF approach using different Skyrme interactions as shown in Tables IV, V, and VI where one can see that only 5 Skyrme interactions, i.e., RATP, SII, SIV, SKa, and SV, among the 63 Skyrme interactions have \(E_{\text{sym,4}}(\rho_0)\) larger than 1 MeV (but still less than 2 MeV). The value of \(E_{\text{sym,4}}(\rho_0) = 0.62\) MeV is further consistent with

| Force   | \(E_{\text{sym}}(\rho_0)\) | \(L\)  | \(K_{\text{sym}}\) | \(J_{\text{sym}}\) | \(E_{\text{sym,4}}(\rho_0)\) | \(L_{\text{sym,4}}\) | \(K_{\text{sym,4}}\) | \(m_s^*/m\) | \(m_{v,0}/m\) |
|---------|-----------------------------|------|-------------------|------------------|-----------------------------|-------------------|-------------------|------------|------------|
| SLy9    | 32.1                        | 55.4 | -81.3             | 461.8            | 0.33                        | 0.32              | -2.88             | 0.67       | 0.80       |
| SkI6    | 29.9                        | 59.2 | -46.8             | 378.1            | 0.28                        | 0.04              | -3.55             | 0.64       | 0.80       |
| SkI4    | 29.5                        | 60.4 | -40.6             | 351.2            | 0.30                        | 0.15              | -3.36             | 0.65       | 0.80       |
| SIV     | 31.2                        | 63.5 | -136.7            | 79.5             | 1.37                        | 5.51              | 7.50              | 0.47       | 0.41       |
| SGI     | 28.3                        | 63.9 | -52.0             | 194.5            | 0.86                        | 3.00              | 2.36              | 0.61       | 0.57       |
| SKO*    | 32.1                        | 69.7 | -77.5             | 221.4            | 0.55                        | 1.38              | -0.90             | 0.90       | 0.87       |
| SkMP    | 29.7                        | 69.8 | -50.3             | 159.7            | 0.93                        | 3.31              | 3.00              | 0.65       | 0.58       |
| SKa     | 32.9                        | 74.6 | -78.5             | 174.5            | 1.13                        | 4.33              | 5.06              | 0.61       | 0.51       |
| SKO     | 32.0                        | 79.6 | -42.3             | 130.0            | 0.59                        | 1.57              | -0.53             | 0.89       | 0.85       |
| R\(\sigma\) | 30.6                    | 85.7 | -9.1              | 22.2             | 0.85                        | 2.88              | 2.13              | 0.78       | 0.68       |
| G\(\sigma\) | 31.4                    | 94.0 | 14.0              | -26.7            | 0.85                        | 2.87              | 2.12              | 0.78       | 0.68       |
| SKT4    | 35.5                        | 94.1 | -24.5             | 97.8             | 0.45                        | 0.91              | -1.83             | 1.00       | 1.00       |
| SV      | 32.8                        | 96.1 | 24.2              | 48.0             | 1.07                        | 0.78              | 4.07              | 0.38       | 0.33       |
| SkI3    | 34.8                        | 100.5| 73.0              | 211.5            | 0.12                        | -0.74             | -5.10             | 0.58       | 0.82       |
| SkI2    | 33.4                        | 104.3| 70.7              | 51.6             | 0.37                        | 0.48              | -2.66             | 0.68       | 0.80       |
| SkI5    | 36.6                        | 129.3| 159.6             | 11.7             | 0.12                        | -0.72             | -5.04             | 0.58       | 0.80       |
FIG. 1: (Color online) Energy per nucleon as a function of density for symmetric nuclear matter in the MDI interaction. Also included are results obtained by using Eq. (4) up to $\chi^2$, $\chi^3$, and $\chi^4$, respectively.

The above results thus indicate that generally one needs higher-order terms in $\chi$ (higher than 4th-order) to describe the EOS of asymmetric nuclear matter at high density region (above $2\rho_0$). We note that above conclusions obtained with the MDI interaction are also valid for the SHF approach and the MSL model. These features imply that it is very difficult to obtain correct information on the high density behaviors of the EOS for asymmetric nuclear matter based on the characteristic parameters obtained at the normal nuclear density $\rho_0$. At this point, it should be mentioned that the transport model analysis of heavy-ion collisions at intermediate and high energies as well as the astrophysical observations, especially on
the properties of compact stars, provide unique tools to extract information on the EOS of asymmetric nuclear matter at high densities [8, 12, 15, 26, 114].

**B. Isospin dependence of the saturation properties of asymmetric nuclear matter**

In the following, we show the results on the saturation properties of asymmetric nuclear matter, i.e., the saturation density as well as the binding energy and incompressibility at saturation density. Especially, we investigate their isospin dependence and study if the higher-order terms in the isospin asymmetry $\delta$ ($\delta^4$ term) are important for the description of the saturation properties of asymmetric nuclear matter.

1. The saturation density

The saturation density is a basic quantity of asymmetric nuclear matter. To see the symmetry energy dependence of the saturation density of asymmetric nuclear matter, we
FIG. 3: (Color online) Density dependence of the 4th-order symmetry energy $E_{\text{sym},4}(\rho)$ using the MDI interaction together with results obtained by using Eq. (16) up to $\chi$, $\chi^2$, $\chi^3$, and $\chi^4$, respectively.

use here the MDI interaction with $x = 1, 0, \text{ and } -1$. The density dependence of the symmetry energy from this interaction is shown in Fig. 4. Also included in Fig. 4 are the results from the widely used APR (Akmal-Pandharipande-Ravenhall) prediction [115] and the phenomenological MSL model prediction with $y = -15, -7.5, \text{ and } 0.75$ (We will discuss the MSL results later). It is seen that the APR prediction for the symmetry energy resembles very well that from the MDI interaction with $x = 0$ up to about $3.5 \rho_0$.

Using the MDI interaction with $x = 1, 0, \text{ and } -1$, we have calculated the density and isospin asymmetry dependence of the binding energy per nucleon of asymmetric nuclear matter, and the results are shown in Fig. 5. Further indicated in Fig. 5 are corresponding saturation points in the $E-\rho$ plane. One can see that different symmetry energies lead to different EOS of pure neutron matter ($\delta = 1$) as expected. In particular, the EOS of pure neutron matter for $x = -1$ is bounded at low densities. In addition, different symmetry energies lead to rather different behaviors for the saturation points in the $E-\rho$ plane.

In order to see more clearly the isospin dependence of the saturation density, we show in Fig. 6 the saturation density $\rho_{\text{sat}}(\delta)$ as a function of $\delta^2$ in the MDI interaction with
FIG. 4: (Color online) Density dependence of the symmetry energy from the MDI interaction with \( x = 1, 0, \) and \(-1\). The results from the widely used APR (Akmal-Pandharipande-Ravenhall) prediction [115] and the phenomenological MSL model prediction with \( y = -15, -7.5, \) and \(0.75\) are also included for comparison.

\( x = 1, 0, \) and \(-1, \) respectively. The exact saturation density \( \rho_{\text{sat}}(\delta) \) is obtained from Eqs. (21) and (43). Corresponding results from Eq. (23) including terms up to \( \delta^2 \) and up to \( \delta^4 \), respectively, are also included for comparison. The results indicate that the saturation density generally decreases with isospin asymmetry and more neutron-rich nuclear matter has lower saturation density. In addition, for the stiffer symmetry energy \( (x = -1) \), the nuclear matter can be bounded even for pure neutron matter (The corresponding saturation density is about \( 0.3\rho_0 \)). The inset in Fig. 6 displays corresponding results at smaller isospin asymmetries with \( \delta^2 \leq 0.1 \) which is relevant to the properties of finite nuclei. In the small isospin asymmetry region \( (\delta^2 \leq 0.1) \), the saturation density \( \rho_{\text{sat}}(\delta) \) displays a linear dependence on \( \delta^2 \) and Eq. (23) including terms up to \( \delta^2 \) thus approximates very well the exact \( \rho_{\text{sat}}(\delta) \). Furthermore, how \( \rho_{\text{sat}}(\delta) \) decreases with increasing \( \delta \) depends on the stiffness of symmetry energy with the softer symmetry energy having a weaker dependence, and this feature can be easily understood from Eq. (24) which indicates that the slope of \( \rho_{\text{sat}}(\delta) \)
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FIG. 5: (Color online) The density and isospin asymmetry dependence of the binding energy per nucleon for asymmetric nuclear matter in the MDI interaction with \( x = 1, 0, \) and \(-1\). The saturation points at different isospin asymmetries are also indicated.

with respect to \( \delta^2 \) is proportional to \(-\frac{3L_{K_0}}{K_0}\rho_0\). In addition, at larger isospin asymmetries with \( \delta^2 \geq 0.3 \), including higher-order \( \delta \) terms up to \( \delta^4 \) in Eq. (23) still deviates significantly from the exact \( \rho_{\text{sat}}(\delta) \) and higher-order terms of \( \delta \) are thus necessary (except for the case of \( x = 0 \) where Eq. (23) with terms up to \( \delta^4 \) gives a good approximation to the exact \( \rho_{\text{sat}}(\delta) \) in the whole \( \delta \) region where the asymmetric matter can still be bounded). These features imply that higher-order terms in \( \delta \) may be important for the determination of the saturation density of nuclear matter at very neutron-rich nuclear environment, such as inside a neutron star.

2. Binding energy at saturation density

The isospin dependence of the binding energy per nucleon of asymmetric nuclear matter at saturation density \( E_{\text{sat}}(\delta) \) is shown in Fig. 7 as a function of \( \delta^2 \) for the MDI interaction with \( x = 1, 0, \) and \(-1\). The exact \( E_{\text{sat}}(\delta) \) is obtained from Eqs. (21) and (43). Corresponding results from Eq. (27) including terms up to \( \delta^2 \) and up to \( \delta^4 \), respectively, are also included.
FIG. 6: (Color online) The saturation density $\rho_{\text{sat}}$ as a function of $\delta^2$ in the MDI interaction with $x = 1, 0, -1$. Corresponding results from Eq. (23) including terms up to $\delta^2$ and up to $\delta^4$, respectively, are also included for comparison. The inset in Fig. 6 displays corresponding results at smaller isospin asymmetries with $\delta^2 \leq 0.1$.

for comparison. It is seen that $E_{\text{sat}}(\delta)$ generally increases with increasing isospin asymmetry. The results at smaller isospin asymmetries with $\delta^2 \leq 0.1$ are shown in the inset of Fig. 7 and it is seen that $E_{\text{sat}}(\delta)$ displays a linear dependence on $\delta^2$ and therefore can be very well approximated by Eq. (27) including terms up to $\delta^2$. We note that the rate at which $E_{\text{sat}}(\delta)$ increases with $\delta^2$ at small $\delta^2$ is determined uniquely by $E_{\text{sym}}(\rho_0)$ as shown in Eq. (27). Also, it is seen that including higher-order $\delta$ terms up to $\delta^4$ in Eq. (27) gives a good approximation to the exact $E_{\text{sat}}(\delta)$ in the whole $\delta$ region (except for the case of $x = -1$ where Eq. (27) including terms up to $\delta^4$ still deviates significantly from the exact $E_{\text{sat}}(\delta)$ at larger isospin asymmetries with $\delta^2 \geq 0.3$ and higher-order terms of $\delta$ are thus needed).

3. Incompressibility at saturation density

Shown in Fig. 8 is the incompressibility at saturation density $K_{\text{sat}}(\delta)$ as a function of $\delta^2$ for the MDI interaction with $x = 1, 0, -1$. The exact $K_{\text{sat}}(\delta)$ is obtained from
Eqs. (21), (32), and (43), and corresponding results from Eq. (34) including terms up to $\delta^2$ and up to $\delta^4$ are also included for comparison. It is seen that $K_{\text{sat}}(\delta)$ generally decreases with increasing isospin asymmetry and more neutron-rich nuclear matter has smaller incompressibility. This feature is consistent with earlier calculations based on microscopic many-body approaches [116]. The softening of the incompressibility of asymmetric nuclear matter with increasing isospin asymmetry may have important implications in understanding the mechanism for supernovae explosions [116, 117]. Corresponding results at smaller isospin asymmetries with $\delta^2 \leq 0.1$ are given in the inset of Fig. 8, and it shows that Eq. (34) including terms up to $\delta^2$ approximates very well the exact $K_{\text{sat}}(\delta)$ as $K_{\text{sat}}(\delta)$ displays a good linear correlation with $\delta^2$. As to the decreasing rate of $K_{\text{sat}}(\delta)$ with increasing $\delta^2$ at small $\delta^2$, it is determined by the parameter $K_{\text{sat},2}$ as shown in Eq. (34) which depends on the characteristic parameters $J_0$, $K_0$, $L$, and $K_{\text{sym}}$. In addition, including higher-order $\delta$ terms up to $\delta^4$ in Eq. (34) is seen to give a good approximation to the exact $K_{\text{sat}}(\delta)$ in the whole $\delta$ region (except for the case of $x = -1$ where Eq. (34) including terms up to $\delta^4$ deviates significantly from the exact $K_{\text{sat}}(\delta)$ at larger isospin asymmetries with $\delta^2 \geq 0.3$ and higher-order terms of $\delta$ are thus important).

The above results indicate that the saturation properties of asymmetric nuclear matter,
FIG. 8: (Color online) Same as Fig. 6 but for the incompressibility at saturation density $K_{\text{sat}}$.

i.e., the saturation density as well as the binding energy and incompressibility at saturation density, exhibit a good linear correlation with $\delta^2$ at smaller isospin asymmetries with $\delta^2 \leq 0.1$ which is relevant to the properties of finite heavy nuclei. On the other hand, depending on the stiffness of nuclear symmetry energy, higher-order terms in $\delta$ ($\delta^4$ and higher-order terms) may become important for describing reasonably the saturation properties of asymmetric nuclear matter at larger isospin asymmetries with $\delta^2 \geq 0.3$. The importance of higher-order isospin asymmetry terms for the stiffer symmetry energy has also been observed in previous studies on the transition density in neutron stars $^{72,73}$. In addition, the saturation density and the incompressibility at saturation density generally decrease with the magnitude of isospin asymmetry while the binding energy at saturation density shows an opposite behavior.

Again, we note that above conclusions obtained from the MDI interaction are also valid for the SHF approach and the MSL model. Our results are further consistent with the very recent study based on the RMF model $^{118}$.  

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C. Constraining the $K_{\text{sat},2}$ parameter from the phenomenological MSL model

1. General information on the $K_{\text{sat},2}$ parameter

As shown in the above, the $K_{\text{sat},2}$ parameter essentially characterizes the isospin dependence of the incompressibility of asymmetric nuclear matter at saturation density and the higher-order parameters (e.g., $K_{\text{sat},4}$) are only important for extremely neutron-rich (or proton-rich) nuclear matter with stiffer nuclear symmetry energies. Actually, it can be seen from Tables I, II, and III that the magnitude (absolute values) of $K_{\text{sat},2}$ is generally much larger than that of $K_{\text{sat},4}$ for the MDI interaction with $x = 1, 0, -1$ and the 63 standard Skyrme interactions considered in the present work. Shown in Fig. 9 are the absolute values of $K_{\text{sat},2}$ and $K_{\text{asy}}$ as functions of $L$ for the MDI interaction with $x = 1, 0, -1$ and the 63 standard Skyrme interactions considered in the present work. It is seen that these values of $K_{\text{sat},2}$ can be nicely expressed as $-400 \pm 120$ MeV. For the magnitude of the $K_{\text{asy}}$ parameter, it is generally larger than that of the $K_{\text{sat},2}$ parameter, especially
for the stiffer symmetry energies (larger $L$ values), which indicates that the higher-order $J_0$ parameter is important as discussed later. Furthermore, the absolute values of $K_{\text{sat},2}$ are clearly much larger than that of $K_{\text{sat},4}$ except that at very large $L$ values the absolute values of $K_{\text{sat},4}$ may become larger and comparable with that of $K_{\text{sat},2}$. This feature is consistent with the results shown in Fig. 8 where the higher-order terms are seen to be only important for the stiffer symmetry energies.

It is generally believed that information on $K_{\text{sat},2}$ can be extracted experimentally by measuring the GMR in neutron-rich nuclei [119]. Usually, one can define a finite nucleus incompressibility $K_A(N, Z)$ for a nucleus with $N$ neutrons and $Z$ protons ($A = N + Z$) by the energy of GMR $E_{\text{GMR}}$, i.e.,

$$E_{\text{GMR}} = \sqrt{\frac{\hbar^2 K_A(N, Z)}{m \langle r^2 \rangle}}, \quad (71)$$

where $m$ is the nucleon mass and $\langle r^2 \rangle$ is the mean square mass radius of the nucleus in the ground state. Similarly to the semi-empirical mass formula, the finite nucleus incompressibility $K_A(N, Z)$ can be expanded as [119]

$$K_A(N, Z) = K_0 + K_{\text{surf}} A^{-1/3} + K_{\text{curv}} A^{-2/3} + (K_\tau + K_{\text{ss}} A^{-1/3}) \left( \frac{N - Z}{A} \right)^2 + K_{\text{Coul}} \frac{Z^2}{A^{4/3}} + \cdots, \quad (72)$$

Neglecting the $K_{\text{curv}}$ term, the $K_{\text{ss}}$ term and other higher-order terms in Eq. (72), one can express the finite nucleus incompressibility $K_A(N, Z)$ as

$$K_A(N, Z) = K_0 + K_{\text{surf}} A^{-1/3} + \tau A^{-2/3} \left( \frac{N - Z}{A} \right)^2 + K_{\text{Coul}} \frac{Z^2}{A^{4/3}}, \quad (73)$$

where $K_0$, $K_{\text{surf}}$, $K_\tau$, and $K_{\text{coul}}$ represent the volume, surface, symmetry, and Coulomb terms, respectively. The $K_\tau$ parameter is usually thought to be equivalent to the $K_{\text{sat},2}$ parameter. However, we would like to stress here that the $K_{\text{sat},2}$ parameter is a theoretically well-defined physical property of asymmetric nuclear matter as shown previously while the value of the $K_\tau$ parameter may depend on the details of the truncation scheme in Eq. (72). As shown in Ref. [120], $K_\tau$ may be related to the isospin-dependent part of the surface
properties of finite nuclei, especially the surface symmetry energy. Therefore, cautions are
needed to interpret the $K_\tau$ parameter as the $K_{\text{sat},2}$ parameter (we will go back to this point
later).

Earlier attempts based on the above method have given widely different values for the
$K_\tau$ parameter. For example, a value of $K_\tau = -320 \pm 180$ MeV with a large uncertainty was
obtained in Ref. [121] from a systematic study of the GMR in the isotopic chains of Sn and
Sm. In this analysis, the value of $K_0$ was found to be $300 \pm 25$ MeV, which is somewhat
larger than the commonly accepted value of $240 \pm 20$ MeV. In a later study, an even less
stringent constraint of $-566 \pm 1350 < K_\tau < 139 \pm 1617$ MeV was extracted from the GMR
of finite nuclei, depending on the mass region of nuclei and the number of parameters used
in parameterizing the incompressibility of finite nuclei [122]. More recently, a much more
stringent constraint of $K_\tau = -550 \pm 100$ MeV has been obtained in Ref. [22, 23] from
measurements of the isotopic dependence of the GMR in even-A Sn isotopes.

2. Correlation between $J_0$ and $K_0$

As shown in Eq. (35), the $K_{\text{sat},2}$ parameter is completely determined by the 4 char-
acteristic parameters $K_0$, $J_0$, $L$, and $K_{\text{sym}}$ at the normal nuclear density. It thus would
be interesting to estimate the possible value of $K_{\text{sat},2}$ from knowledge on $K_0$, $J_0$, $L$, and
$K_{\text{sym}}$. For the incompressibility of symmetric nuclear matter at its saturation density $\rho_0$,
the transport model analyses on experimental data from subthreshold kaon production in
heavy-ion collisions favor a soft equation of state [26, 123, 124]. More recently, the value of
$K_0$ has been more stringently determined to be $240 \pm 20$ MeV from the nuclear GMR data
[16, 21, 22, 23, 24].

While the $K_0$ parameter has been relatively well determined, the $J_0$ parameter is poorly
known, and there is actually no experimental information on the $J_0$ parameter. In the MSL
model, one can easily calculate from Eq. (57) the $J_0$ parameter as

$$ J_0 = 27 \rho^3 \frac{\partial^3 E_0(\rho)}{\partial \rho^3} |_{\rho=\rho_0} $$

$$ = \frac{1}{E_{\text{kin}}^0 - 3E_0 - 2C} \left[ K_0^2 + (18E_0 - 2E_{\text{kin}}^0 - 8C)K_0 
+ 12E_{\text{kin}}^0E_0 + 6C(3E_{\text{kin}}^0 - 25E_0) \right]. \quad (74) $$
FIG. 10: (Color online) The third derivative parameter $J_0$ and its ratio $J_0/K_0$ to the incompressibility $K_0$ of symmetry nuclear matter at saturation density as functions of $K_0$ from the MSL model, the MDI interaction and the SHF prediction with the 63 Skyrme interactions shown in Tables I, II, and III.

In particular, for $m_{s,0}^* = 0.8m$, we have

$$J_0 = \frac{1}{59.1} \left( K_0^2 - 376.4K_0 + 11214.645 \right) \text{ (MeV)}, \quad (75)$$

while we have

$$J_0 = \frac{1}{70.1} \left( K_0^2 - 332.2K_0 - 4243.2 \right) \text{ (MeV)}, \quad (76)$$

for $m_{s,0}^* = m$ (and then $C = 0$) and

$$J_0 = \frac{1}{51.2} \left( K_0^2 - 408K_0 + 22252 \right) \text{ (MeV)} \quad (77)$$

for $m_{s,0}^* = 0.7m$ and $C = 9.47$ MeV. In above equations, the unit of $K_0$ is MeV. Therefore, the $J_0$ parameter in the phenomenological MSL model is a quadratic function of the $K_0$ parameter.
Shown in Fig. [10] are $J_0$ and $J_0/K_0$ as functions of $K_0$ in the MSL model for different values of $m_{s,0}^* = m$, 0.9$m$, 0.8$m$ and 0.7$m$. Also included in Fig. [10] are corresponding results from the MDI interaction and the SHF prediction with the 63 Skyrme interactions shown in Tables [I] [II] and [III]. It is seen that the correlation between $J_0$ and $K_0$ is similar among these three different models. Also, all three models show an approximately linear correlation between $J_0/K_0$ and $K_0$. This linear correlation can be easily understood from Eqs. (75), (76) and (77). For example, on the r.h.s of Eq. (75), the last term is small compared with the first two terms, and thus one has $J_0 \approx \frac{1}{59.1} (K_0^2 - 376.4K_0)$, and then $J_0/K_0 \approx \frac{1}{59.1} (K_0 - 376.4)$ with the units of $K_0$ in MeV. The linear dependence of $J_0/K_0$ on $K_0$ becomes better for larger values of $m_{s,0}^*$ which usually leads to smaller values for the last two terms on the r.h.s of Eq. (74) as shown in Eq. (76). We note that the correlation between $J_0$ and $K_0$ obtained in the present work is also consistent with the early finding by Pearson [125]. In addition, it is seen from Fig. [10] that the $m_{s,0}^*$ only has visible effects on the correlation between $J_0$ and $J_0/K_0$ for smaller $K_0$ values while it has almost no effects for larger $K_0$ values. While there do not exist any empirical constraints on the $J_0$ parameter, we assume in the present study the correlation between $J_0$ and $K_0$ from the MSL model is valid and then determine $J_0/K_0$ from the experimental constraint on $K_0$.

3. Correlation between $L$ and $K_{sym}$

The parameters $L$ and $K_{sym}$ are determined by the density dependence of the symmetry energy around saturation density. In recent years, significant progress has been made both experimentally and theoretically in extracting the information on the behaviors of nuclear symmetry energy at sub-saturation density. Using the isospin- and momentum-dependent IBUU04 transport model with in-medium NN cross sections, the isospin diffusion data were found to be consistent with the symmetry energy from the MDI interaction with $x$ between 0 and −1, which can be parameterized by a density-dependent symmetry energy of $E_{sym}(\rho) \approx 31.6(\rho/\rho_0)^\gamma$ with $\gamma = 0.69 - 1.05$ at subnormal density ($\rho \leq \rho_0$) [50, 55, 56], leading to the extraction of $61 \text{ MeV} \leq L \leq 111 \text{ MeV}$ and $-82 \text{ MeV} \leq K_{sym} \leq 101 \text{ MeV}$ [50, 55, 56]. Using the Skyrme interactions consistent with the EOS obtained from the MDI interaction with $x$ between 0 and −1, the neutron-skin thickness of heavy nuclei calculated within the Hartree-Fock approach is consistent with available experimental data [50, 126] and also
FIG. 11: (Color online) Correlation between $K_{\text{sym}}$ and $L$ from the MSL model with $\gamma_{\text{sym}} = 4/3$ and 5/3, the MDI interaction and the SHF prediction with the 63 Skyrme interactions shown in Tables IV, V, and VI. The results from two simple one-parameter symmetry energies in Eqs. (81) and (82) are also shown for comparison.

that from a relativistic mean-field model based on an accurately calibrated parameter set that reproduces the GMR in $^{90}$Zr and $^{208}$Pb as well as the isovector giant dipole resonance of $^{208}$Pb [127]. The extracted symmetry energy further agrees with the symmetry energy $E_{\text{sym}}(\rho) = 31.6(\rho/\rho_0)^{0.69}$ recently obtained from the isoscaling analyses of isotope ratios in intermediate energy heavy ion collisions [57], which gives $L \approx 65$ MeV and $K_{\text{sym}} \approx -61$ MeV. Furthermore, the above limited range of $E_{\text{sym}}(\rho)$ at subsaturation density is essentially consistent with the symmetry energy $E_{\text{sym}}(\rho) = 12.5(\rho/\rho_0)^{2/3} + 17.6(\rho/\rho_0)^{\gamma}$ with $\gamma = 0.4 - 1.05$, extracted very recently from the ImQMD (Improved QMD) model analyses of both the isospin diffusion data and the double neutron/proton ratio [59]. The symmetry energy $E_{\text{sym}}(\rho) = 12.5(\rho/\rho_0)^{2/3} + 17.6(\rho/\rho_0)^{\gamma}$ with $\gamma = 0.4 - 1.05$ leads to the constraints of $46 \text{ MeV} \leq L \leq 80 \text{ MeV}$ and $-63 \text{ MeV} \leq K_{\text{sym}} \leq -17 \text{ MeV}$.

It should be noted that all above constraints on $L$ and $K_{\text{sym}}$ are based on some unique energy density functionals and thus special correlation between $L$ and $K_{\text{sym}}$ has been implicitly assumed. It is thus of interest to study if there exists a universal correlation between $L$ and
$K_{\text{sym}}$ in different models. For the MDI interaction, it is implied from Eq. (45) that the $L$ and $K_{\text{sym}}$ both change linearly with the parameter $x$. Therefore they are linearly correlated by varying the parameter $x$ that changes the density dependence of the symmetry energy. For the MSL model, we can obtain from Eq. (66) following expressions

$$L = 2E_{\text{sym}}^{\text{kin}}(\rho_0) + 5D + 3 \left[ E_{\text{sym}}(\rho_0) - E_{\text{sym}}^{\text{kin}}(\rho_0) - D \right]$$

$$+ 3y(\gamma_{\text{sym}} - 1) \left[ E_{\text{sym}}(\rho_0) - E_{\text{sym}}^{\text{kin}}(\rho_0) - D \right]$$

$$K_{\text{sym}} = 9y\gamma_{\text{sym}}(\gamma_{\text{sym}} - 1) \left[ E_{\text{sym}}(\rho_0) - E_{\text{sym}}^{\text{kin}}(\rho_0) - D \right]$$

$$+ 10D - 2E_{\text{sym}}^{\text{kin}}(\rho_0).$$

(78)

(79)

Therefore, for the MSL model, Eq. (78) and Eq. (79) show that the $L$ and $K_{\text{sym}}$ both change linearly with the parameter $y$, and thus they are also linearly correlated by varying the parameter $y$ to change the density dependence of the symmetry energy. In particular, we have

$$K_{\text{sym}} = 3\gamma_{\text{sym}}L + E_{\text{sym}}^{\text{kin}}(\rho_0)(3\gamma_{\text{sym}} - 2)$$

$$+ 2D(5 - 3\gamma_{\text{sym}}) - 9\gamma_{\text{sym}}E_{\text{sym}}(\rho_0).$$

(80)

Also, $L$ and $K_{\text{sym}}$ are expected to be correlated within the SHF energy density functional.

Shown in Fig. 11 are the correlation between $K_{\text{sym}}$ and $L$ from the MSL model with $\gamma_{\text{sym}} = 4/3$ and 5/3. (We used here the default values of $E_{\text{sym}}^{\text{kin}}(\rho_0) = 12.3$ MeV, $E_{\text{sym}}(\rho_0) = 30$ MeV, and $D = -3.51$ MeV from $m_{s,0}^* = 0.8m$ and $m_{v,0}^* = 0.7m$), the MDI interaction and the SHF prediction with the 63 Skyrme interactions shown in Tables IV, V, and VI. It is seen that the $K_{\text{sym}}$ parameter indeed displays approximately a linear correlation with the $L$ parameter for the SHF predictions with the 63 Skyrme interactions and this linear correlation is nicely reproduced by the MDI interaction and the MSL model with $\gamma_{\text{sym}} = 4/3$. For the MSL model, one can see from Eq. (80) that the $\gamma_{\text{sym}}$ parameter controls the shape (slope) of the linear correlation between $L$ and $K_{\text{sym}}$. Furthermore, Fig. 11 shows that results from some Skyrme interactions deviate from the linear correlation obtained by the MDI interaction and the MSL model with $\gamma_{\text{sym}} = 4/3$. In order to consider the uncertainty in the shape (slope) for the correlation between $L$ and $K_{\text{sym}}$, we also include the result from the MSL model with $\gamma_{\text{sym}} = 5/3$. The correlation between $K_{\text{sym}}$ and $L$ from the SHF predictions with the 63 Skyrme interactions is consistent nicely with that from the MSL model with $\gamma_{\text{sym}} = 4/3$ and 5/3. Furthermore, we find that the isospin-dependent nucleon
effective mass has very small effects on the correlation between $K_{\text{sym}}$ and $L$. This can be easily understood from Eq. \(80\) since the value of $2D(5 - 3\gamma_{\text{sym}})$ is only about 7 MeV for $\gamma_{\text{sym}} = 4/3$ (it is zero for $\gamma_{\text{sym}} = 5/3$). The linear correlation between $K_{\text{sym}}$ and $L$ implies that one can extract $K_{\text{sym}}$ from $L$.

As mentioned above, an one-parameter parametrization for the symmetry energy is sometimes used for simplicity, i.e.,

$$E_{\text{sym}}(\rho) = E_{\text{sym}}(\rho_0) \left( \frac{\rho}{\rho_0} \right)^\gamma$$ \hspace{1cm} (81)

or

$$E_{\text{sym}}(\rho) = E_{\text{sym}}^{\text{kin}}(\rho_0) \left( \frac{\rho}{\rho_0} \right)^{2/3} \left[ E_{\text{sym}}(\rho_0) - E_{\text{sym}}^{\text{kin}}(\rho_0) \right] \left( \frac{\rho}{\rho_0} \right)^\gamma. \hspace{1cm} (82)$$

For the parametrization of Eq. \(81\), the $L$ and $K_{\text{sym}}$ can be obtained as

$$L = 3\gamma E_{\text{sym}}(\rho_0)$$ \hspace{1cm} (83)

$$K_{\text{sym}} = 9\gamma(\gamma - 1)E_{\text{sym}}(\rho_0)$$

and furthermore we have

$$K_{\text{sym}} = \frac{L^2}{E_{\text{sym}}(\rho_0)} - 3L,$$

which indicates that the $K_{\text{sym}}$ parameter is quadratically correlated with $L$. The result obtained by varying the parameter $\gamma$ is shown in Fig. 11 by the dotted line. Similarly, for the parametrization of Eq. \(82\), the $L$ and $K_{\text{sym}}$ can be expressed as

$$L = 2E_{\text{sym}}^{\text{kin}}(\rho_0) + 3\gamma \left[ E_{\text{sym}}(\rho_0) - E_{\text{sym}}^{\text{kin}}(\rho_0) \right]$$ \hspace{1cm} (84)

$$K_{\text{sym}} = 9\gamma(\gamma - 1) \left[ E_{\text{sym}}(\rho_0) - E_{\text{sym}}^{\text{kin}}(\rho_0) \right] - 2E_{\text{sym}}^{\text{kin}}(\rho_0),$$ \hspace{1cm} (85)

and the $K_{\text{sym}}$ parameter is again quadratically correlated with $L$, and this is illustrated by the dash-dot-dotted line in Fig. 11 obtained by varying the value of $\gamma$.

It is very interesting to see from Fig. 11 that for larger $L$ values ($L \geq 45$ MeV which is consistent with the constraint from heavy-ion collision data shown later), all above symmetry energy functionals from different models and parameterizations give consistent predictions for the $K_{\text{sym}}$-$L$ correlation. This nice feature implies that using these different models and
parameterizations for the symmetry energy will not influence significantly the determination of the $K_{sat,2}$ parameter. On the other hand, the $K_{sym}$-$L$ correlation from the two one-parameter parameterizations on the symmetry energies in Eqs. (81) and (82) deviate significantly from the MDI, MSL and SHF predictions for small $L$ values. Actually, the two forms of one-parameter parametrization for the symmetry energy in Eqs. (81) and (82) may be too simple to describe a softer symmetry energy. As shown in Ref. [55], Eq. (81) cannot describe correctly the density dependence of the symmetry energy from the MDI interaction with $x = 1$ (the Gogny interaction). On the other hand, as shown in Fig. 4, the MSL model can give a nice description of the density dependence of the symmetry energy from the very soft to the very stiff. Although there is no direct empirical information on the $K_{sym}$ parameter and some uncertainty on the $K_{sym}$-$L$ correlation still exist, we assume here that the correlation between $K_{sym}$ and $L$ from the MSL model with $\gamma_{sym} = 4/3$ and $5/3$ is valid and then use the experimental constraint on $L$ to extract the value of $K_{sym}$.

4. Phenomenological MSL model constraint on the $K_{sat,2}$ parameter

From above analyses, we can now extract information on the $K_{sat,2}$ parameter from the experimental constraints on the $K_0$ parameter and the $L$ parameter within the phenomenological MSL model. As pointed out previously, the value of $K_0$ has been relatively well determined to be $240 \pm 20$ MeV from the nuclear GMR [16, 21, 22, 23, 24]. The slope parameter $L$ has been found to correlate linearly with the neutron-skin thickness of heavy nuclei and thus can in principle be determined from measured thickness of the neutron skin of such nuclei [50, 64, 126, 128, 129, 130, 131, 132]. Unfortunately, because of the large uncertainties in the experimental measurements, this has not yet been possible so far. The proposed experiment of parity-violating electron scattering from $^{208}$Pb, i.e., Parity Radius Experiment (PREx) at the Jefferson Laboratory is expected to give an independent and accurate measurement of its neutron skin thickness (within 0.05 fm) [133, 134] and thus to impose a stringent constraint on the slope parameter $L$ in future. On the other hand, as mentioned previously, heavy-ion collisions, especially those induced by neutron-rich nuclei, provide a unique tool to explore the density dependence of the symmetry energy and thus the $L$ parameter. Actually, as discussed previously, the $L$ parameter has been already limited significantly by heavy-ion collisions data.

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From Eqs. (35) and (80), we have in the MSL model

\[
K_{\text{sat},2} = K_{\text{sym}} - 6L - \frac{J_0}{K_0}L \\
= -\left(\frac{J_0}{K_0} + 6 - 3\gamma_{\text{sym}}\right)L \\
+ (3\gamma_{\text{sym}} - 2)E_{\text{sym}}(\rho_0) + 2D(5 - 3\gamma_{\text{sym}}) \\
- 9\gamma_{\text{sym}}E_{\text{sym}}(\rho_0).
\]

(86)

Results on \(K_{\text{sat},2}\) as a function of \(L\) are shown in Fig. 12 for \(\gamma_{\text{sym}} = 4/3\) (panel (a)) and 5/3 (panel (b)) and with \(m_{s,0}^* = 0.8m\) and \(m_{v,0}^* = 0.7m\) for different values of \(K_0\) and \(E_{\text{sym}}(\rho_0)\). The shaded region indicates constraints within the MSL model with \(K_0 = 240 \pm 20\) MeV, \(E_{\text{sym}}(\rho_0) = 30 \pm 5\) MeV, and \(46\) MeV \(\leq L \leq 111\) MeV limited by the heavy-ion collision data. The results from the SHF approach with 63 Skyrme interactions are also included for comparison. In addition, the constraint of \(K_{\tau} = -550 \pm 100\) MeV obtained in Ref. [22, 23] from measurements of the isotopic dependence of the GMR in even-A Sn isotopes is also indicated.

From Eqs. (35) and (80), we have in the MSL model

\[
K_{\text{sat},2} = K_{\text{sym}} - 6L - \frac{J_0}{K_0}L \\
= -\left(\frac{J_0}{K_0} + 6 - 3\gamma_{\text{sym}}\right)L \\
+ (3\gamma_{\text{sym}} - 2)E_{\text{sym}}(\rho_0) + 2D(5 - 3\gamma_{\text{sym}}) \\
- 9\gamma_{\text{sym}}E_{\text{sym}}(\rho_0).
\]

(86)

Results on \(K_{\text{sat},2}\) as a function of \(L\) are shown in Fig. 12 for \(\gamma_{\text{sym}} = 4/3\) (panel (a)) and 5/3 (panel (b)) and with \(K_0 = 220, 240,\) and 260 MeV. In Eq. (86), the value of \(J_0\) is obtained from Eq. (74) using the value of \(K_0\). For other quantities in the MSL model, the default values \(E_{\text{sym}}(\rho_0) = 12.3\) MeV, \(E_{\text{sym}}(\rho_0) = 30\) MeV, and \(D = -3.51\) MeV from \(m_{s,0}^* = 0.8m\) and \(m_{v,0}^* = 0.7m\) have been used in the calculations. Since the value of \(2D(5 - 3\gamma_{\text{sym}})\) is only about 7 MeV for \(\gamma_{\text{sym}} = 4/3\) and is zero for \(\gamma_{\text{sym}} = 5/3\), as mentioned before, the
contribution of the $D$ parameter to $K_{\text{sat,2}}$ is thus small. Also, the correlation between $K_{\text{sym}}$ and $L$ depends on $E_{\text{sym}}(\rho_0)$ and $D$ as shown in Eq. [80]. To take into consideration of the uncertainty in the value of $E_{\text{sym}}(\rho_0)$, we also include in Fig. 12 the results with $K_0 = 220$ MeV and $E_{\text{sym}}(\rho_0) = 25$ MeV as well as $K_0 = 260$ MeV and $E_{\text{sym}}(\rho_0) = 35$ MeV, which represent, respectively, the upper and lower bounds for a fixed value of $L$. The shaded region in Fig. 12 further shows the constrained $L$ values from heavy-ion collision data, namely, $46$ MeV $\leq L \leq 111$ MeV. The lower limit of $L = 46$ MeV is obtained from the lower bound in the ImQMD analyses of the isospin diffusion data and the double neutron/proton ratio [59] while the upper limit of $L = 111$ MeV corresponds to the upper bound of $L$ from the IBUU04 transport model analysis of the isospin diffusion data [50, 54, 55, 56]. The constraint $46$ MeV $\leq L \leq 111$ MeV is consistent with the analyses of the pygmy dipole resonances [66], the giant dipole resonance (GDR) of $^{208}$Pb analyzed with Skyrme interactions [65], the very precise Thomas-Fermi model fit to the binding energies of 1654 nuclei [63], and the recent neutron-skin analysis [78]. These empirically extracted values of $L$ represent the best and most stringent phenomenological constraints available so far on the nuclear symmetry energy at sub-saturation densities.

It is seen from Fig. 12 that the $K_{\text{sat,2}}$ parameter decreases with increasing $L$ for $\gamma_{\text{sym}} = 4/3$ while the opposite behavior is observed for $\gamma_{\text{sym}} = 5/3$. This feature can be easily understood from Eq. [86]. For $\gamma_{\text{sym}} = 4/3$, Eq. [86] is reduced to

$$K_{\text{sat,2}} = -(J_0/K_0 + 2)L - 12E_{\text{sym}}(\rho_0) + 17.6 \text{ (MeV)},$$

(87)

while for $\gamma_{\text{sym}} = 5/3$, it is reduced to

$$K_{\text{sat,2}} = -(J_0/K_0 + 1)L - 15E_{\text{sym}}(\rho_0) + 36.9 \text{ (MeV)}. \quad (88)$$

For $K_0 = 240 \pm 20$ MeV, Fig. 11 (or Eq. [74]) shows that the value of $J_0/K_0$ varies from about $-1.9$ to $-1.2$ when $m_{*0}$ changes from $m$ to $0.7m$. Eq. [87] (Eq. [88]) thus shows that $K_{\text{sat,2}}$ decreases (increases) with increasing $L$ for $\gamma_{\text{sym}} = 4/3 (5/3)$.

An interesting feature observed from Fig. 12 is that the $K_{\text{sat,2}}$ parameter depends significantly on the symmetry energy at the normal nuclear density $E_{\text{sym}}(\rho_0)$. This can be seen more clearly from Eqs. [87] and [88] which show that changing $E_{\text{sym}}(\rho_0)$ by 5 MeV leads to a variation of $60 - 75$ MeV for $K_{\text{sat,2}}$. This feature indicates that an accurate determination of $E_{\text{sym}}(\rho_0)$ is important for determining the value of $K_{\text{sat,2}}$. In addition, we have used
$E_{\text{sym}}(\rho_0) = 30$ MeV in Fig. 11 and one can easily see that a variation of 5 MeV for $E_{\text{sym}}(\rho_0)$ would lead to a shift of $60 - 75$ MeV for the $K_{\text{sym}}$ parameter for a fixed $L$. Therefore, the MSL model with $\gamma_{\text{sym}} = 4/3$ and $5/3$ together with $E_{\text{sym}}(\rho_0) = 30 \pm 5$ MeV may provide a nice estimate for both the $K_{\text{sym}}-L$ correlation and its uncertainty.

From the shaded region indicated in Fig. 12, we find that for $\gamma_{\text{sym}} = 4/3$, we have $-437$ MeV $\leq K_{\text{sat,2}} \leq -292$ MeV for $L = 46$ MeV while $-487$ MeV $\leq K_{\text{sat,2}} \leq -306$ MeV for $L = 111$ MeV. For $\gamma_{\text{sym}} = 5/3$, we have $-477$ MeV $\leq K_{\text{sat,2}} \leq -302$ MeV for $L = 46$ MeV while $-461$ MeV $\leq K_{\text{sat,2}} \leq -251$ MeV for $L = 111$ MeV. These results indicate that based on the MSL model with $4/3 \leq \gamma_{\text{sym}} \leq 5/3$, $K_0 = 240 \pm 20$ MeV, $25$ MeV $\leq E_{\text{sym}}(\rho_0) \leq 35$ MeV, and $46$ MeV $\leq L \leq 111$ MeV, the $K_{\text{sat,2}}$ parameter can vary from $-251$ MeV to $-487$ MeV. The results shown in Fig. 12 is obtained from a $J_0/K_0$ value that is evaluated with the default value $m_{s,0}^* = 0.8m$. Similar analyses indicate that the $K_{\text{sat,2}}$ parameter varies from $-261$ MeV to $-489$ MeV if we use $m_{s,0}^* = 0.7m$ while it varies from $-245$ MeV to $-485$ MeV if $m_{s,0}^* = 0.9m$ is used. These results indicate that the extracted value for $K_{\text{sat,2}}$ is not sensitive to the variation of the nucleon effective mass. The MSL model analyses with $4/3 \leq \gamma_{\text{sym}} \leq 5/3$, $K_0 = 240 \pm 20$ MeV, $E_{\text{sym}}(\rho_0) = 30 \pm 5$ MeV, $46$ MeV $\leq L \leq 111$ MeV, and $m_{s,0}^* = 0.8 \pm 0.1m$ thus lead to an estimate of $K_{\text{sat,2}} = -370 \pm 120$ MeV.

Also included in Fig. 12 are the results from the SHF approach with the 63 Skyrme interactions. It is seen that among the 63 Skyrme interactions shown in Tables II, III and IV there are only 19 Skyrme interactions with predicted $K_{\text{sat,2}}$ consistent with the constraints obtained here (the shaded regions in Fig. 12 (a) or (b)). The 19 Skyrme interactions are SKa, SKT4, G$_\sigma$, SkI3, SkI2, R$_\sigma$, SKO, SKO*, SkMP, SGI, SKM, SLy5, SLy1, SLy9, SLy2, SLy7, SkI6, SLy230b, and SkI4. From Tables II - VI one can see that all these 19 Skyrme interactions predict $0.151$ fm$^{-3} \leq \rho_0 \leq 0.161$ fm$^{-3}$, $-16.0$ MeV $\leq E_0(\rho_0) \leq -15.6$ MeV, $217$ MeV $\leq K_0 \leq 263$ MeV, $28$ MeV $\leq E_{\text{sym}}(\rho_0) \leq 35$ MeV, and $46$ MeV $\leq L \leq 104$ MeV, which are consistent with the empirical information. However, the value of the nucleon effective mass predicted by the 19 Skyrme interactions varies widely, i.e., from $0.58m$ to $m$ for $m_{s,0}^*$ while from $0.51m$ to $m$ for $m_{e,0}^*$. If we further impose the constraints of $m_{s,0}^* = 0.8 \pm 0.1m$ and $m_{s,0}^* > m_{e,0}^*$, then only five Skyrme interactions, namely, G$_\sigma$, R$_\sigma$, SKM, SKO and SKO*, remain to be consistent with all known empirical constraints except that the SKM interaction gives a little smaller incompressibility $K_0 = 216.6$ MeV than the empirical constrained values of $K_0 = 240 \pm 20$ MeV. These features imply that the constraint on the $K_{\text{sat,2}}$ parameter
does not significantly limit the nucleon effective mass as the former has previously been shown to be insensitive to the latter.

5. Discussions on parameters $K_{\text{sat},2}$, $K_{\text{asy}}$, and $K_{\tau}$

As shown in Eq. (38), the $K_{\text{asy}}$ parameter corresponds to the $K_{\text{sat},2}$ parameter when $J_0$ is zero, i.e., the parabolic approximation to the EOS of symmetric nuclear matter (Eq. (10)) is valid. In the MSL model, a vanishing $J_0$ requires a $K_0$ value of about 340 MeV, which is significantly larger than the empirical value of $240 \pm 20$ MeV. Since $J_0/K_0 \approx -1.5$ for $K_0 = 240$ MeV in the MSL model, we have $K_{\text{sat},2} \approx K_{\text{asy}} + 1.5L$ or $K_{\text{asy}} \approx K_{\text{sat},2} - 1.5L$. Therefore, the difference between the $K_{\text{asy}}$ parameter and the $K_{\text{sat},2}$ parameter depends on the slope parameter $L$ of the symmetry energy and a larger $L$ value (stiffer symmetry energy) would lead to a larger difference. It should be stressed, however, that the $K_{\text{asy}}$ parameter is completely determined by the density dependence of symmetry energy regardless of the EOS of symmetric nuclear matter.

Based on the IBUU04 transport model analysis of the isospin diffusion data, a value of $K_{\text{asy}} = -500 \pm 50$ MeV has been extracted from the symmetry energy obtained by the MDI interaction with the $x$ parameter between 0 and $-1$. The extracted $K_{\text{asy}} = -500 \pm 50$ MeV is essentially consistent with the symmetry energy extracted from the ImQMD model analyses of both the isospin diffusion data and the double neutron/proton ratio [59], which predict a $K_{\text{asy}}$ value from about $-500$ MeV to $-340$ MeV, and the symmetry energy $E_{\text{sym}}(\rho) = 31.6(\rho/\rho_0)^{0.69}$ obtained from the isoscaling analyses [57], which gives $K_{\text{asy}} \approx -453$ MeV. Furthermore, the constraint $K_{\text{asy}} = -500 \pm 50$ MeV is consistent with the very recent constraint of $K_{\text{asy}} \approx -500^{+125}_{-100}$ MeV obtained from the study of neutron skin of finite nuclei [78]. In the MDI interaction, one finds $J_0/K_0 = -2.1$ from Table I and we thus have $-311$ MeV $\leq K_{\text{sat},2} \leq -316$ MeV, corresponding to the prediction from the MDI interaction with the $x$ parameter between 0 and $-1$. Therefore, the $K_{\text{sat},2}$ value for the MDI interaction is significantly larger than the $K_{\text{asy}}$ value and is insensitive to the density dependence of symmetry energy. On the other hand, the $J_0$ parameter in the MDI interaction is important for the determination of the $K_{\text{sat},2}$ parameter. The importance of the higher-order $J_0$ parameter to $K_{\text{sat},2}$ can be further seen from Fig. [9] which shows that for the 63 Skyrme interactions considered here, the magnitude of the $K_{\text{asy}}$ parameter is generally
larger than that of the $K_{\text{sat,2}}$ parameter, especially, for the stiffer symmetry energies (larger $L$ values). These results indicate that the higher-order $J_0$ contribution to $K_{\text{sat,2}}$ generally cannot be neglected.

In Fig. 12 the constraint $K_\tau = -550 \pm 100$ MeV obtained in Ref. [22, 23] from recent measurements of the isotopic dependence of the GMR in even-A Sn isotopes is also shown. While the estimate of $K_{\text{sat,2}} = -370 \pm 120$ MeV obtained in the present work has small overlap with the constraint of $K_\tau = -550 \pm 100$ MeV, the latter still has significantly larger magnitude than the former. According to Fig. 12 there are only 6 Skyrme interactions, namely, SII, SIV, SV, SI, SkI5, SIII that predict $K_{\text{sat,2}} = -550 \pm 100$ MeV. As shown in Tables I, II, and III the values for $K_0$ from SI, SII, SIII SIV, and SV are all larger than 305 MeV which are obviously inconsistent with the empirical constraint of $K_0 = 240 \pm 20$ MeV. For SkI5, we have $L = 129.3$ MeV which is significantly larger than the empirical constraint of $46$ MeV $\leq L \leq 111$ MeV. Therefore, none of the 63 Skyrme interactions considered in the present work is consistent simultaneously with $K_{\text{sat,2}} = -550 \pm 100$ MeV and the empirical constraints of $K_0 = 240 \pm 20$ MeV and $46$ MeV $\leq L \leq 111$ MeV. These features imply that the $K_\tau$ parameter extracted from Eq. (73) based on the GMR may not fully reflect the $K_{\text{sat,2}}$ parameter. As mentioned before, the $K_\tau$ parameter may depend on the detailed truncation scheme in Eq. (72). The constraint $K_\tau = -550 \pm 100$ MeV obtained in Ref. [22, 23] is based on Eq. (73) and thus neglects contributions from the $K_{\text{curv}}$ term, $K_{\text{as}}$ and other higher-order terms in Eq. (72). It is expected that including contributions from higher-order terms in Eq. (73) may change the extracted value for the $K_\tau$ parameter as found in Ref. [122].

V. SUMMARY AND CONCLUSIONS

We have studied in the present paper higher-order effects on the properties of isospin asymmetric nuclear matter when the EOS and saturation properties of asymmetric nuclear matter are expanded in powers of the isospin asymmetry parameter $\delta = (\rho_n - \rho_p)/\rho$ and the dimensionless variable $\chi = (\rho - \rho_0)/3\rho_0$ that characterizes the deviations of the density from the normal nuclear density $\rho_0$. Analytical expressions for the saturation density of asymmetric nuclear matter as well as its binding energy and incompressibility at saturation density have been derived and given exactly up to 4th-order in the isospin asymmetry $\delta = (\rho_n - \rho_p)/\rho$ using 11 characteristic parameters defined at the normal nuclear density $\rho_0$.  


by the density derivatives of the binding energy per nucleon of symmetric nuclear matter, the symmetry energy $E_{\text{sym}}(\rho)$ and the 4th-order symmetry energy $E_{\text{sym},4}(\rho)$, namely, $E_0(\rho_0)$, $K_0$, $J_0$, $I_0$, $E_{\text{sym}}(\rho_0)$, $L$, $K_{\text{sym}}$, $J_{\text{sym}}$, $E_{\text{sym},4}(\rho_0)$, $L_{\text{sym},4}$, and $K_{\text{sym},4}$. Our method/recipe to derive the analytical expressions for the saturation properties of asymmetric nuclear matter is quite general, and in principle, the higher-order coefficients (higher than $\delta^4$) in the isospin asymmetry $\delta$ can be easily obtained. Using the isospin- and momentum-dependent MDI interaction, the SHF approach with 63 popular Skyrme interactions, and the phenomenological MSL model, we have systematically studied the higher-order effects on the properties of asymmetric nuclear matter.

Firstly, our results indicate that including terms up to $\chi^4$ can give a good description of the EOS of symmetric nuclear matter $E_0(\rho)$, the density-dependent symmetry energy $E_{\text{sym}}(\rho)$, and the density-dependent 4th-order symmetry energy $E_{\text{sym},4}(\rho)$ for densities less than about $2\rho_0$ while higher-order terms in $\chi$ (higher than 4th-order and thus more characteristic parameters defined at $\rho_0$ are necessary) are generally needed to describe reasonably the $E_0(\rho)$, $E_{\text{sym}}(\rho)$, and $E_{\text{sym},4}(\rho)$ at higher density region (above $2\rho_0$). These features imply that it is very difficult to predict the high density behaviors of the EOS of asymmetric nuclear matter based on the characteristic parameters obtained at the normal nuclear density $\rho_0$. Therefore, it is of great interest to use the transport model analysis of heavy-ion collisions at intermediate and high energies as well as the astrophysical observations of compact stars as tools to extract information on the EOS of asymmetric nuclear matter at high densities.

Secondly, we have systematically studied the isospin dependence of the saturation properties of asymmetric nuclear matter, i.e., the saturation density $\rho_{\text{sat}}(\delta)$ as well as the binding energy $E_{\text{sat}}(\delta)$ and incompressibility $K_{\text{sat}}(\delta)$ at saturation density. In particular, we have compared their exact results with the analytical expressions expanded up to the $\delta^2$ term and the $\delta^4$ term, respectively. Our results indicate that the saturation properties of asymmetric nuclear matter, i.e., $\rho_{\text{sat}}(\delta)$, $E_{\text{sat}}(\delta)$ and $K_{\text{sat}}(\delta)$, exhibit a good linear dependence on $\delta^2$ at smaller isospin asymmetries with $\delta^2 \leq 0.1$ which is relevant to the structure of finite heavy nuclei. This feature implies that the higher-order terms (the $\delta^4$ term and higher-order terms in $\delta$) are not important for the saturation properties of asymmetric nuclear matter, at least for an asymmetric nuclear matter that is not extremely neutron-rich (proton-rich). On the other hand, for asymmetric nuclear matter with extremely large isospin asymmetries with $\delta^2 \geq 0.3$ and depending on the stiffness of the nuclear symmetry energy, the higher-order
terms in $\delta$ ($\delta^4$ and higher-order terms) may become important for describing reasonably the saturation properties. It is further found that the saturation density and the incompressibility at saturation density generally decrease with the magnitude of the isospin asymmetry while the binding energy at saturation density shows the opposite behavior.

Finally, we have studied in detail the second-order isospin coefficient $K_{\text{sat},2}$ of the incompressibility of an asymmetric nuclear matter at its saturation density. It is found that the magnitude of the higher-order $K_{\text{sat},4}$ parameter is generally small compared to that of the $K_{\text{sat},2}$ parameter, so the latter essentially characterizes the isospin dependence of the incompressibility at saturation density. Furthermore, we have found that the $K_{\text{sat},2}$ parameter is uniquely determined by $L$, $K_{\text{sym}}$ and $J_0/K_0$. Since there is no experimental information on the $J_0$ parameter and the $K_{\text{sym}}$ parameter, we have thus used the MSL model, which can reasonably describe the general properties of symmetric nuclear matter and the symmetry energy predicted by both the MDI model and the SHF approach, to estimate the value of $K_{\text{sat},2}$. Our results indicate that generally the higher-order $J_0$ contribution to $K_{\text{sat},2}$ cannot be neglected, especially for larger $L$ values. Interestingly, it is found that there exists a nicely linear correlation between $K_{\text{sym}}$ and $L$ as well as between $J_0/K_0$ and $K_0$ for the three different models used here, i.e., the MDI interaction, the MSL interaction, and the SHF approach with 63 Skyrme interactions. For the MSL model, the correlation between $K_{\text{sym}}$ and $L$ is further found to depend significantly on the value of $E_{\text{sym}}(\rho_0)$ but not on variations of the nucleon effective mass. These correlations and features have allowed us to extract the values of the $J_0$ parameter and the $K_{\text{sym}}$ parameter from the empirical information on $K_0$, $L$ and $E_{\text{sym}}(\rho_0)$. In particular, using the empirical constraints of $K_0 = 240 \pm 20$ MeV, $E_{\text{sym}}(\rho_0) = 30 \pm 5$ MeV, $46$ MeV $\leq L \leq 111$ MeV and $m_s,0 = 0.8 \pm 0.1 m$ in the MSL model leads to an estimate of $K_{\text{sat},2} = -370 \pm 120$ MeV.

While the estimated value of $K_{\text{sat},2} = -370 \pm 120$ MeV in the present work has small overlap with the constraint of $K_\tau = -550 \pm 100$ MeV obtained in Refs. \[22, 23\] from recent measurements of the isotopic dependence of the GMR in even-A Sn isotopes, its magnitude is significantly smaller than that of the constrained $K_\tau$. Recently, there are several studies \[118, 135, 136\] on extracting the value of the $K_{\text{sat},2}$ parameter based on the idea initiated by Blaizot and collaborators that the values of both $K_0$ and $K_{\text{sat},2}$ should be extracted from the same consistent theoretical model that successfully reproduces the experimental GMR energies of a variety of nuclei. These studies show that no single model (interaction) can
simultaneously describe correctly the recent measurements of the isotopic dependence of the
GMR in even-A Sn isotopes and the GMR data of $^{90}$Zr and $^{208}$Pb nuclei, and this makes
it difficult to accurately determine the value of $K_{\text{sat,2}}$ from these experimental data. Also,
a very recent study [137] indicates that the effect due to the nuclear superfluidity may also
affect the extraction of the $K_{\text{sat,2}}$ parameter from the nuclear GMR. As pointed out in [118],
these features suggest that the $K_r = -550 \pm 100$ MeV obtained in Ref. [22, 23] may suffer
from the same ambiguities already encountered in earlier attempts to extract the $K_0$ and
$K_{\text{sat,2}}$ of infinite nuclear matter from finite-nuclei extrapolations. This problem remains an
open challenge, and both experimental and theoretical insights are needed in future studies.

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**APPENDIX: DERIVATION OF ANALYTICAL EXPRESSIONS FOR THE SAT-
URATION DENSITY AND INCOMPRESSIBILITY TO HIGHER-ORDER IN
ISOSPIN ASYMMETRY**

For completeness, we present in this appendix a simple derivation of the analytical ex-
pressions for the saturation density of asymmetric nuclear matter and its incompressibility
at saturation density up to higher-order terms in the isospin asymmetry $\delta = (\rho_n - \rho_p)/\rho$.
In particular, analytical expressions for the saturation density and incompressibility at sat-
uration density are given exactly up to 4th-order in $\delta$.

The binding energy per nucleon $e(\rho, \delta)$ of asymmetric nuclear matter can be expanded
up to 2nth-order in $\delta$ as

$$e(\rho, \delta) = e_0(\rho) + e_2(\rho)\delta^2 + e_4(\rho)\delta^4 + \cdots + e_{2n}(\rho)\delta^{2n}. \quad (A.1)$$

The first term $e_0(\rho)$ represents the EOS of symmetry nuclear matter with $\delta = 0$. We further expand each $e_i(\rho)$ in Eq. (A.1) up to nth-order in a dimensionless variable $z = (\rho - \rho_0) / \rho_0$, i.e.

$$e_0(\rho) = e_0(\rho_0) + a_{01}z + a_{02}z^2 + \cdots + a_{0n}z^n$$
$$e_2(\rho) = e_2(\rho_0) + a_{11}z + a_{12}z^2 + \cdots + a_{1n}z^n$$
$$e_4(\rho) = e_4(\rho_0) + a_{21}z + a_{22}z^2 + \cdots + a_{2n}z^n$$
$$\vdots$$
$$e_{2n}(\rho) = e_{2n}(\rho_0) + a_{n1}z + a_{n2}z^2 + \cdots + a_{nn}z^n. \quad (A.2)$$

The coefficients $a_{ij}$ represent the characteristic parameters of asymmetric nuclear matter defined at the normal nuclear density $\rho_0$. The saturation density $\rho_{\text{sat}}$ is then determined by the following equation

$$\left. \frac{\partial e}{\partial \rho} \right|_{\rho=\rho_{\text{sat}}} = 0 \quad (A.3)$$

or equivalently

$$\left. \frac{\partial e}{\partial z} \right|_{z=z_{\text{sat}}} = 0 \quad (A.4)$$

with $z_{\text{sat}} = (\rho_{\text{sat}} - \rho_0) / \rho_0$

Substituting Eq. (A.1) and Eq. (A.2) into Eq. (A.4) leads to the following equation

$$a_{01} + 2a_{02}z + 3a_{03}z^2 + \cdots + na_{0n}z^{n-1} +$$
$$+ (a_{11} + 2a_{12}z + 3a_{13}z^2 + \cdots + na_{1n}z^{n-1})\delta^2 +$$
$$+ (a_{21} + 2a_{22}z + 3a_{23}z^2 + \cdots + na_{2n}z^{n-1})\delta^4 +$$
$$+ \cdots +$$
$$+ (a_{n1} + 2a_{n2}z + 3a_{n3}z^2 + \cdots + na_{nn}z^{n-1})\delta^{2n} = 0. \quad (A.5)$$

Obviously, we have $a_{01} = 0$ as the saturation density of symmetry nuclear matter is determined by

$$\left. \frac{\partial e}{\partial \rho} \right|_{\delta=0, \rho=\rho_0} = 0.$$

In order to obtain the solution to Eq. (A.5), i.e., $z_{\text{sat}}$ with precision of $O(\delta^{2k+2})$, we assume

$$z_{\text{sat}} = A_2\delta^2 + A_4\delta^4 + \cdots + A_{2k}\delta^{2k}, \quad (A.6)$$

where $A_{2m}$ ($m = 1, 2, \ldots, k$) are coefficients to be determined from comparing the coefficients of $\delta^{2j}$ terms ($j = 1, 2, \ldots, m$) on the two sides of Eq. (A.5) after substituting Eq. (A.6) into
Eq. (A.5). For example, for $k = 1$, i.e., $z_{\text{sat}} = A_2 \delta^2$, the coefficient $A_2$ can be determined by the following equality

$$2a_{02}A_2 + a_{11} = 0.$$  

(A.7)

Similarly, for $k = 2$, i.e., $z_{\text{sat}} = A_2 \delta^2 + A_4 \delta^4$, the coefficient $A_4$ can be determined by

$$2a_{02}A_4 + 3a_{03}A_2^2 + 2a_{12}A_2 + a_{21} = 0.$$  

(A.8)

Furthermore, for $k = 3$ and 4, the coefficients $A_6$ and $A_8$ can be determined, respectively, by

$$2a_{02}A_6 + 6a_{03}A_2A_4 + 4a_{04}A_2^3 + 2a_{12}A_4 + 3a_{13}A_2^2 + 2a_{22}A_2 + a_{31} = 0$$  

(A.9)

and

$$2a_{02}A_8 + 3a_{03}(2A_2A_6 + A_4^2) + 12a_{04}A_2^2A_4 + 5a_{05}A_2^4 + 2a_{12}A_6 + 6a_{13}A_2A_4 + 4a_{14}A_2^3 + 2a_{22}A_4 + 3a_{23}A_2^2 + 2a_{32}A_2 + a_{41} = 0.$$  

(A.10)

From above analyses, the coefficient of the $\delta^{2m}$ term, $A_{2m}$, can be determined by all lower-order coefficients $A_2, A_4, A_{2m-2}$ ($m = 1, 2, 3, \ldots, n$). The $A_{2m}$ obtained above are complete and precise up to the order of $\delta^{2m}$, and the higher-order characteristic parameters $a_{ij}$ do not contribute to $A_{2m}$. In order to write down the general expressions for $A_{2m}$, we define the following symbol

$$B_{n,m} = \sum A_2^{i_1}A_4^{i_2} \cdots A_{2n}^{i_n},$$  

(A.11)
where we have \(i_1, i_2, \ldots, i_n \in \{0, N_+\}\) and they satisfy following conditions

\[
\sum_{j=1}^{n} j i_j = m, \quad \sum_{j=1}^{n} i_j = n.
\] (A.12)

Then the coefficient of the \(\delta^{2m}\) term satisfies the equation

\[
\sum_{j=0}^{m-1} \sum_{i=2}^{m-j+1} ia_{ji} B_{i-1,m-j} + a_{m1} = 0
\] (A.13)

from which we can obtain the coefficient \(A_{2m}\) as

\[
A_{2m} = -\frac{1}{a_{02}} \left( \sum_{i=3}^{m+1} ia_{0i} B_{i-1,m} + \sum_{j=1}^{m-1} \sum_{i=2}^{m-j+1} ia_{ji} B_{i-1,m-j} + a_{m1} \right).
\] (A.14)

Therefore, one can obtain the saturation density of asymmetry nuclear matter to any order of \(\delta\).

Taking the precision to the order of \(\delta^4\), we then have

\[
A_2 = -\frac{a_{11}}{2a_{02}}, \quad A_4 = \frac{a_{12} a_{11}}{2a_{02}} - \frac{3a_{03} a_{11}^2}{8a_{02}^2} - \frac{a_{21}}{2a_{02}}.
\] (A.15)

Converting the coefficients \(a_{ij}\) into the conventional forms, i.e.,

\[
a_{02} = \frac{K_0}{18}, a_{03} = \frac{J_0}{162}, a_{04} = \frac{J_0}{24 \times 81},
\]

\[
a_{11} = \frac{L}{3}, a_{12} = \frac{K_{\text{sym}}}{18}, a_{13} = \frac{J_{\text{sym}}}{162},
\]

\[
a_{21} = \frac{L_{\text{sym},4}}{3}, a_{22} = \frac{K_{\text{sym},4}}{18},
\] (A.16)

we then obtain

\[
z_{\text{sat}} = -\frac{3L}{K_0} \delta^2 + \left( \frac{3K_{\text{sym}} L}{K_0^2} - \frac{3L_{\text{sym},4}}{K_0} - \frac{3J_0 L^2}{2K_0^3} \right) \delta^4.
\] (A.17)

So the saturation density can be obtained as

\[
\rho_{\text{sat}} = \rho_0 \left[ 1 - \frac{3L}{K_0} \delta^2 + \left( \frac{3K_{\text{sym}} L}{K_0^2} - \frac{3L_{\text{sym},4}}{K_0} - \frac{3J_0 L^2}{2K_0^3} \right) \delta^4 + O(\delta^6) \right],
\] (A.18)

which is exactly Eq. (22).
The incompressibility coefficient \( K_{\text{sat}} \) of asymmetry nuclear matter at the saturation density is defined as

\[
K_{\text{sat}} = 9\rho_0^2 \frac{\partial^2 e}{\partial \rho^2} \bigg|_{\rho = \rho_{\text{sat}}}. \tag{A.19}
\]

Substituting Eq. \( \text{(A.1)} \) and Eq. \( \text{(A.2)} \) into Eq. \( \text{(A.19)} \) and using Eq. \( \text{(A.18)} \), we can easily obtain

\[
K_{\text{sat}} = 9\rho_0^2 \left[ 1 - \frac{3L}{K_0} \delta^2 + \left( \frac{3K_{\text{sym}} L}{K_0^2} - \frac{3L_{\text{sym}}}{K_0} - \frac{3J_0 K_0^2}{2K_0^3} \right) \delta^4 \right]^{\frac{1}{2}} \nonumber
\]

\[
\times \left[ \frac{K_0}{9\rho_0^2} + \frac{J_0}{27\rho_0^2} \rho - \rho_0 + \frac{I_0}{162} \left( \rho - \rho_0 \right)^2 \right] \delta^2 \nonumber
\]

\[
+ \left( \frac{K_{\text{sym}}}{9\rho_0^2} + \frac{J_{\text{sym}}}{27\rho_0^2} \rho - \rho_0 + \frac{I_{\text{sym}}}{162} \left( \rho - \rho_0 \right)^2 \right) \delta^4 \nonumber
\]

\[
+ \left( \frac{K_{\text{sym}}}{9\rho_0^2} + \frac{J_{\text{sym}}}{27\rho_0^2} \rho - \rho_0 + \frac{I_{\text{sym}}}{162} \left( \rho - \rho_0 \right)^2 \right) \delta^4 + O(\delta^6) \bigg|_{\rho = \rho_{\text{sat}}} \nonumber
\]

\[
= K_0 + \left( K_{\text{sym}} - 6L - \frac{J_0}{K_0} L \right) \delta^2 
\]

\[
+ \left( K_{\text{sym}} - 6L_{\text{sym}} - \frac{J_0}{K_0} L_{\text{sym}} \right) \delta^2 + \frac{9L^2}{K_0} - \frac{J_{\text{sym}} L}{K_0} 
\]

\[
+ \frac{I_0 L^2}{2K_0^2} + \frac{J_0 K_{\text{sym}} L}{K_0^2} + \frac{3J_0 L^2}{K_0^2} - \frac{J_0^2 L^2}{2K_0^3} \delta^4 + O(\delta^6), \tag{A.20}
\]

which is exactly Eq. \( \text{(33)} \). We note that the expression \( K_{\text{sat}} = K_0 + \left( K_{\text{sym}} - 6L - \frac{J_0}{K_0} L \right) \delta^2 \) was originally given in Ref. [119]. For terms higher than \( \delta^4 \), they can be straightforwardly obtained following above derivation but many more higher-order characteristic parameters \( a_{ij} \) would be needed.

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