Physics Letters B 763 (2016) 145–150

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Symmetry energy and neutron star properties in the saturated Nambu–Jona-Lasinio model

Si-Na Wei, Wei-Zhou Jiang *, Rong-Yao Yang, Dong-Rui Zhang

Department of Physics, Southeast University, Nanjing 211189, China

ARTICLE INFO

Article history:
Received 25 May 2016
Received in revised form 9 October 2016
Accepted 10 October 2016
Available online 17 October 2016
Editor: W. Haxton

Keywords:
Symmetry energy
Neutron stars
NJL model

ABSTRACT

In this work, we adopt the Nambu–Jona-Lasinio (NJL) model that ensures the nuclear matter saturation properties to study the density dependence of the symmetry energy. With the interactions constrained by the chiral symmetry, the symmetry energy shows novel characters different from those in conventional mean-field models. First, the negative symmetry energy at high densities that is absent in relativistic mean-field (RMF) models can be obtained in the RMF approximation by introducing a chiral isovector–vector interaction, although it would be ruled out by the neutron star (NS) stability. Second, with the inclusion of the isovector–scalar interaction the symmetry energy exhibits a general softening at high densities even for the large slope parameter of the symmetry energy. The NS properties obtained in the present NJL model can be in accord with the observations. The NS maximum mass obtained with various isovector–scalar couplings and momentum cutoffs is well above the $2M_\odot$, and the NS radius obtained well meets the limits extracted from recent measurements. In particular, the significant reduction of the canonical NS radius occurs with the moderate decrease of the slope of the symmetry energy.

1. Introduction

The nuclear symmetry energy is important for understanding the reaction dynamics of heavy-ion collisions, the structures of neutron- and proton-rich nuclei, and properties of neutron stars (NS) [1–3]. Though the symmetry energy, which is the energy difference per nucleon between pure neutron matter and symmetric matter, is well constrained at saturation density to date [4–8], the density dependence of the symmetry energy is still poorly known especially at supra-normal densities [2,9]. The symmetry energy predicted by different models is rather diverse at high densities [10–17]. Unfortunately, the symmetry energy extracted from the data with various isospin diffusion models also suffers from the large uncertainty which diversifies in super-soft [18], soft [19], and stiff [20] forms at high densities. We note that new experiments to probe the high-density symmetry energy are also on the way [21]. While different high-density behaviors of the symmetry energy are usually classified by the magnitude of the slope of the symmetry energy at saturation density, we may raise the question: Are there new high-density behaviors of the symmetry energy that can't be simply elaborated by the slope parameter?

On the other hand, the super-soft symmetry energy which reaches the maximum and then turns to negative values at high densities can be obtained from some non-relativistic models [12, 13], while it can not be produced in the relativistic mean field (RMF) models [14–17]. For instance, the nonlinear RMF models [15], the density-dependent RMF models [16,22], and the point coupling RMF models [23–25] predict similar tendencies of symmetry energy, and no super-soft symmetry energy arises in these models [17]. Since the success of RMF models in interpreting the pseudospin symmetry [26–28] and analyzing polarization observables in proton–nuclei reactions [29,30] indicates that the relativistic dynamics that includes the large attractive scalar and repulsive vector [31–35] is of special importance, we may ask whether the super-soft symmetry energy is incompatible with the relativistic covariance, or it is hidden in some special interactions that are not included in usual RMF models.

To answer these questions, let's first recall the prime importance of the chiral symmetry in the strong interaction. In fact, the chiral symmetry has served as a cornerstone to construct the effective QCD models of the strong interaction [36,37]. In the development of RMF models, the chiral symmetry has also played an important role in guiding the nonlinear form of the meson self-interacting terms needed for the appropriate in-medium effects [38–42]. To explore the novel high-density behaviors of the

* Corresponding author.
E-mail address: wzjiang@seu.edu.cn (W.-Z. Jiang).

http://dx.doi.org/10.1016/j.physletb.2016.10.019

symmetry energy in the RMF approximation, it is appropriate to adopt chiral models and thus constrain the relevant interactions with the chiral symmetry. Among models respecting the chiral symmetry in bulk matter [38,41,43–45], the Nambu–Jona-Lasinio (NJL) model [43] and chiral-σ model [38,41] are two popular ones. The NJL model was originally proposed to realize the spontaneous symmetry breaking since the pion, as the Goldstone boson, can be derived dynamically. With the quark degrees of freedom, the NJL model is considered as an effective model for the QCD [46–48]. While it is not straightforward to construct the nucleons and describe nuclear matter due to the absence of the confinement in the NJL model [49], it is economic to realize in the NJL model the spontaneous breaking of the chiral symmetry with nucleonic degrees of freedom [50–53], like the chiral-σ model. In the hadron-level NJL model, the character of chiral symmetry is also measured by the chiral condensate in the non-perturbative vacuum. In this work, we thus study in the hadron-level NJL model the density dependence of the symmetry energy with the various interactions respecting the chiral symmetry.

Recently, remarkable progresses in NS observations have been achieved. Accurate mass measurements determined two large-mass NS's: the radio pulsar J1614-2230 with mass of $M = 1.97 \pm 0.04M_\odot$ [54] and the J0348+0432 with mass of $M = 2.01 \pm 0.04M_\odot$ [55]. However, there is no consensus on the extracted NS radius [56] reported in the literature [57–63], due to the systematic uncertainties involved in the distance measurements and theoretical analyses of the light spectrum [64–67]. In this work, we will then investigate whether the parametrizations of the present saturated NJL model can satisfy the NS mass constraint and provide some useful comparisons with various NS radius constraints. In the following, we will in turn present the formalism, analyze the results, and give the summary.

2. Formalism

The original NJL model that only contains scalar, pseudoscalar, vector and axial vector interactions can not reproduce saturation properties of nuclear matter. In order to obtain the saturation property, the scalar–vector (SV) interaction, which also respects the chiral symmetry, was introduced [50,51]. This is similar to the chiral-σ model, where the saturation is fulfilled by introducing the scalar–vector coupling [41,68]. Similar efforts were also made to study the nuclear matter saturation and the phase diagram in the NJL model [52,53]. The Lagrangian of the saturated NJL model can then be written as [51]:

$$\begin{aligned}\mathcal{L}_0 = {}& \bar\psi(i\gamma_\mu\partial^\mu - m_0)\psi + \frac{G_S}{2}[(\bar\psi\psi)^2 - (\bar\psi\gamma_5\tau\psi)^2] \\ & - \frac{G_V}{2}[(\bar\psi\gamma_\mu\psi)^2 + (\bar\psi\gamma_\mu\gamma_5\psi)^2] \\ & + \frac{G_{SV}}{2}[(\bar\psi\psi)^2 - (\bar\psi\gamma_5\tau\psi)^2]\cdot[(\bar\psi\gamma_\mu\psi)^2 \\ & + (\bar\psi\gamma_\mu\gamma_5\psi)^2],\end{aligned} \tag{1}$$

where m_0 is the bare nucleon mass. G_S, G_V and G_{SV} are the scalar, vector and scalar–vector coupling constants, respectively. It is easy to see that the Lagrangian is chiral symmetric when $m_0 = 0$. In order to investigate the density dependence of the symmetry energy, we introduce the isovector, isovector–vector and isovector–scalar interactions in the Lagrangian which are written as:

$$\begin{aligned}\mathcal{L}_{IV} = {}& \frac{G_\rho}{2}[(\bar\psi\gamma_\mu\tau\psi)^2 + (\bar\psi\gamma_\mu\gamma_5\tau\psi)^2] \\ & + \frac{G_{\rho V}}{2}[(\bar\psi\gamma_\mu\tau\psi)^2 + (\bar\psi\gamma_\mu\gamma_5\tau\psi)^2] \\ & \cdot[(\bar\psi\gamma_\mu\psi)^2 + (\bar\psi\gamma_\mu\gamma_5\psi)^2] \\ & + \frac{G_{\rho S}}{2}[(\bar\psi\gamma_\mu\tau\psi)^2 + (\bar\psi\gamma_\mu\gamma_5\tau\psi)^2] \\ & \cdot[(\bar\psi\psi)^2 - (\bar\psi\gamma_5\tau\psi)^2],\end{aligned} \tag{2}$$

where G_ρ, $G_{\rho V}$ and $G_{\rho S}$ are the isovector, isovector–vector and isovector–scalar coupling constants, respectively. $\mathcal{L}_{IV}$ is also chirally symmetric. Using the mean-field approximation,

$$\begin{aligned}(\bar\psi A\psi)(\bar\psi B\psi) = {}& (\bar\psi A\psi) < \bar\psi B\psi > + < \bar\psi A\psi > (\bar\psi B\psi) \\ & - < \bar\psi A\psi >< \bar\psi B\psi >,\end{aligned} \tag{3}$$

the Lagrangian can be simplified to be

$$\begin{aligned}\mathcal{L} = \mathcal{L}_0 + \mathcal{L}_{IV} = {}& \bar\psi[i\gamma_\mu\partial^\mu - m(\rho,\rho_S) - \gamma^0\Sigma(\rho,\rho_S,\rho_3)]\psi \\ & - U(\rho,\rho_S,\rho_3),\end{aligned} \tag{4}$$

where m, Σ and U are defined as

$$m(\rho,\rho_S) = m_0 - (G_S + G_{SV}\rho^2 + G_{\rho S}\rho_3^2)\rho_S, \tag{5}$$

$$\begin{aligned}\Sigma(\rho,\rho_S,\rho_3) = {}& G_V\rho + G_\rho\rho_3\tau_3 - G_{SV}\rho_S^2\rho - G_{\rho V}\rho_3^2\rho \\ & - G_{\rho V}\rho_3\rho^2\tau_3 - G_{\rho S}\rho_3\rho_S^2\tau_3,\end{aligned} \tag{6}$$

$$\begin{aligned}U(\rho,\rho_S,\rho_3) = {}& \frac{1}{2}(G_S\rho_S^2 - G_V\rho^2 - G_\rho\rho_3^2 + 3G_{SV}\rho_S^2\rho^2 \\ & + 3G_{\rho V}\rho_3^2\rho^2 + 3G_{\rho S}\rho_3^2\rho_S^2).\end{aligned} \tag{7}$$

Eq. (5) is the gap equation for the nucleon effective mass in the NJL model. Here $\rho =< \bar\psi\gamma^0\psi >$, $\rho_3 =< \bar\psi\gamma^0\tau_3\psi >$ and $\rho_S = < \bar\psi\psi >$ are vector, isovector and scalar densities, respectively. From the energy–momentum tensor, we may obtain the following energy density and pressure

$$\begin{aligned}\epsilon = {}& -\sum_{i=p,n}\nu_i\int_{p_{F_i}}^{\Lambda}\frac{d^3p}{(2\pi)^3}(p^2+m^2)^{1/2} + \frac{G_V\rho^2}{2} + \frac{G_\rho\rho_3^2}{2} + \frac{G_S\rho_S^2}{2} \\ & + \frac{G_{SV}\rho^2\rho_S^2}{2} - \frac{G_{\rho V}{\rho_3}^2\rho^2}{2} + \frac{G_{\rho S}{\rho_3}^2\rho_S^2}{2} + \epsilon_0,\end{aligned} \tag{8}$$

$$\begin{aligned}P = {}& -\sum_{i=p,n}\frac{\nu_i}{3}\int_{p_{F_i}}^{\Lambda}\frac{d^3k}{(2\pi)^3}\frac{k^2}{\sqrt{k^2+m^2}} + \frac{G_V\rho^2}{2} + \frac{G_\rho\rho_3^2}{2} - \frac{G_S\rho_S^2}{2} \\ & - \frac{3G_{SV}\rho_S^2\rho^2}{2}\end{aligned} \tag{9}$$

$$-\frac{3G_{\rho V}\rho_3^2\rho^2}{2} - \frac{3G_{\rho S}{\rho_3}^2\rho_S^2}{2} - \frac{2\Lambda^3\sqrt{\Lambda^2+m^2}}{3\pi^2} - \epsilon_0, \tag{10}$$

where Λ is the momentum cutoff, and the ϵ_0 is introduced to give the vanishing energy density of the vacuum state. From the energy density, we can derive the symmetry energy as

$$\begin{aligned}E_{sym}(\rho) = {}& \frac{1}{2}\frac{\partial^2(\epsilon/\rho)}{\partial\delta^2}\bigg|_{\delta=0} \\ = {}& \frac{p_F^2}{6E_F} + \frac{1}{2}G_\rho\rho - \frac{1}{2}G_{\rho V}\rho^3 - \frac{1}{2}G_{\rho S}\rho_S^2\rho,\end{aligned} \tag{11}$$

where $\delta = (\rho_n - \rho_p)/\rho$ is the isospin asymmetry parameter and $E_F = \sqrt{p_F^2 + m^2}$. The symmetry energy has a term linear in ρ^3 due to the isovector–vector interaction. The slope of the symmetry energy at saturation density is defined as

$$L = 3\rho_0\frac{\partial E_{sym}}{\partial\rho}\bigg|_{\rho_0}. \tag{12}$$