LARGE-MASS NEUTRON STARS WITH HYPERONIZATION

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ABSTRACT

Within a density-dependent relativistic mean-field model and using in-medium meson–hadron coupling constants and meson masses, we explore the effects of in-medium hyperon interactions on the properties of neutron stars. We found that hyperonic constituents in large-mass neutron stars cannot be simply ruled out, while the recently measured mass of the millisecond pulsar J1614−2230 can significantly constrain in-medium hyperon interactions. In addition, we discuss the effects of nuclear symmetry energy on hyperonization in neutron stars.

Key words: dense matter – equation of state – stars: neutron

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1. INTRODUCTION

In-medium hyperon interactions play an important role in determining the properties of hypernuclei and hyperonization in neutron stars. Conversely, the observed properties of hypernuclei and neutron stars can be used to constrain in-medium hyperon interactions. For instance, it has been shown that the properties of hypernuclei are indeed very useful in extracting in-medium hyperon potentials at subsaturation densities (Friedman & Gal 2007). In bulk matter, such as neutron stars, hyperons can be produced by virtue of strong interactions. They can actually become important constituents of neutron stars and thus have important effects on astrophysical observations (for a review, see Glendenning 2001). In fact, it is well known that hyperonization can reduce the maximum mass of neutron stars by as much as 3/4 $M_\odot$ (Glendenning 1985; Glendenning & Moszkowski 1991; Jiang 2006). Currently, a number of phenomenological models, considering only the minimum compositions of nucleons and leptons using interaction parameters that are well calibrated using terrestrial nuclear laboratory data, can produce the maximum mass of neutron stars around or below 2 $M_\odot$ (Danielewicz et al. 2002; Piekarcewicz 2007). Interestingly, several neutron stars with large masses around 2 $M_\odot$ have recently been observed (Nice et al. 2005, 2008; Ozel 2006). In particular, the 2 $M_\odot$ pulsar J1614−2230 was measured rather accurately through the Shapiro delay (Demorest et al. 2010). Since the properties of neutron stars are determined by the nuclear equation of state (EOS) and hyperonization significantly reduces the maximum mass of neutron stars, it has been stated that these observations seem to rule out almost all currently proposed hyperon EOSs. Recent evidence for this can also be found in the work of the Brueckner approach (Schulze & Rijken 2011).

Though most of the hyperon EOSs produce lighter neutron stars, there were actually a few attempts in the past to stiffen the EOS either by invoking strong repulsions for hyperons or pushing the onset density of hyperons upward, leading to heavy neutron stars with hyperons (Hofmann et al. 2001; Takatsuka et al. 2002; Stone et al. 2007; Dexheimer & Schramm 2008). Recently, using nonlinear self-interacting terms involving a vector meson with hidden strangeness, Bednarek et al. obtained a stiff hyperon EOS, which is capable of producing the large mass of the PSR J1614−2230 (Bednarek et al. 2011). Usually, SU(6) relations are imposed to constrain the meson–hyperon coupling constants. The SU(6) relations were recently re-examined by Weissenborn et al. and an arbitrary breaking of these relations can also result in the stiffening of the hyperon EOS and an increase in the maximum mass of neutron stars with the inclusion of hyperons (Weissenborn et al. 2011a).

On the other hand, the quark deconfinement may occur in the medium as the spatial overlap of nucleons becomes sufficient to dissolve the boundary of color singlets with the increase of density. Of course, the quantitative understanding of such a color deconfinement in a cold medium is still model dependent. Typical effective QCD models include the NJL-like models, e.g., see Klähn et al. (2007), Ippolito et al. (2008), Pagliara & Schaffner-Bielich (2008), and Bonanno & Sedrakian (2012), and the widely used MIT bag models, e.g., see Prakash et al. (1997), Alford et al. (2005), and Weissborn et al. (2011b), as well as the Schwinger–Dyson approaches (Li et al. 2011). With the Maxwell or Gibbs construction for the hadron-quark phase coexistence, the resulting quark EOS may give rise to two possible types of stars: strange stars made entirely of absolutely stable strange matter (Glendenning 2000), and hybrid stars with a quark core and hadron outlayer. In order to be consistent with the recent observation of the 2 $M_\odot$ pulsar, the strong coupling and/or color superconductivity were shown to be necessary (Xu 2003; Alford et al. 2005; Klähn et al. 2007; Ippolito et al. 2008; Pagliara & Schaffner-Bielich 2008; Weissborn et al. 2011b; Bonanno & Sedrakian 2012). While the compositions of hybrid stars are rather model dependent, it is interesting to mention that Yasutake et al. suggested the hyperon suppression with the MIT model using a density-dependent bag constant (Yasutake et al. 2011).

Considering recent studies on the consequences of various quark EOSs, in this work we examine the consistency of the hyperon EOS with the recent observation of PSR J1614−2230. Despite impressive progress made in recent decades in...
constraining the nuclear EOS using both astrophysical observations and nuclear reaction data, see, e.g., Youngblood et al. (1999), Danielewicz et al. (2002), and Li et al. (2008), many uncertainties still remain. In-medium hyperon interactions are among the most uncertain ingredients in neutron star models. We shall thus seek in-medium hyperon interactions that can produce the observed maximum mass of neutron stars. In view of the fact that some microscopic theories, such as the Brueckner approach, still have difficulty obtaining the 2 $M_\odot$ of hyperonized neutron stars, we resort to the phenomenological models developed in Jiang et al. (2007a, 2007b) to analyze the effects of various in-medium hyperon interactions on properties of neutron stars.

2. Density-Dependent Relativistic Mean-Field Models and Parameterizations

In order to conveniently study the in-medium interactions for hyperons, we look for density-dependent relativistic models without nonlinear interactions. In our previous works (Jiang et al. 2007a, 2007b), we constructed density-dependent relativistic mean-field (RMF) models using in-medium hadron properties according to the Brown–Rho (BR) scaling because of the chiral symmetry restoration at high densities (Brown & Rho 1991, 2005a, 2005b; Song 2001; Brown et al. 2007). In these models, the symmetric part of the resulting EOSs around normal density is consistent with the data of nuclear giant monopole resonances (Youngblood et al. 1999). At supra-normal densities, it is constrained by the collective flow data from high-energy heavy-ion reactions (Danielewicz et al. 2002), while the resulting density dependence of the symmetry energy at sub-saturation densities agrees with what was extracted from the isospin diffusion data from intermediate-energy heavy-ion reactions (Tsang et al. 2004; Chen et al. 2005; Li & Chen 2005).

Our models with the chiral limits are soft at intermediate densities but naturally stiff at high densities, producing a heavy maximum neutron star mass around 2 $M_\odot$. It is interesting to see that the most recent EOS extracted from the celestial observations features similar characters (Steiner et al. 2010). Apart from the usual studies on minimum constituents in neutron stars with electrons, protons, and neutrons, in this work we include the hyperonic degrees of freedom to study the in-medium interactions for hyperons with the constraints of recent celestial observations. The model Lagrangian with the density-dependent parameters is written as

$$\mathcal{L} = \bar{\psi}_B \left( [i \gamma \mu \partial^\mu - M_B^2 + g_{\sigma B} \sigma - g_{\omega B} \omega \mu - g_{\rho B} \gamma_\mu \partial^\mu \right] \psi_B + \frac{1}{2} \left( \partial_\mu \sigma \partial^\mu \sigma - m_\sigma^2 \sigma^2 \right) - \frac{1}{4} F_{\mu \nu} F^{\mu \nu} + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu - \frac{1}{4} B_{\mu \nu} B^{\mu \nu} + \frac{1}{2} m_\rho^2 \rho_\mu \rho^\mu + \mathcal{L}_Y + \mathcal{L}_I, \tag{1}$$

where $\psi_B$, $\sigma$, $\omega$, and $b_0$ are the fields of the baryons, scalar, vector, and isovector-vector mesons, with their masses $M_B^2$, $m_\sigma^2$, $m_\omega^2$, and $m_\rho^2$, respectively. $F_{\mu \nu}$ and $B_{\mu \nu}$ are the strength tensors of the $\omega$ and $\rho$ mesons, respectively. The meson coupling constants and masses with asterisks denote the density dependence, given by the BR scaling (Jiang et al. 2007a, 2007b). $\mathcal{L}_Y$ and $\mathcal{L}_I$ are the Lagrangians for leptons and hyperons, respectively. The parameters for strange mesons $\sigma^+$ (i.e., $f_0$, 975 MeV) and $\phi$ (1020 MeV) in $\mathcal{L}_Y$ are assumed to be density independent.

The energy density and pressure in the RMF approximation read, respectively,

$$\mathcal{E} = \frac{1}{2} m_\omega^2 \omega_0^2 + \frac{1}{2} m_\rho^2 \rho_0^2 + \frac{1}{2} m_\sigma^2 \sigma_0^2 + \frac{1}{2} \sum_i \frac{2}{(2 \pi)^3} \int d^3 k E_i^2,$$

$$p = \frac{1}{2} m_\omega^2 \omega_0^2 + \frac{1}{2} m_\rho^2 \rho_0^2 + \frac{1}{2} m_\sigma^2 \sigma_0^2 - \frac{1}{2} m_\sigma^2 \sigma^2 - \frac{1}{2} m_\omega^2 \omega^2 - \frac{1}{2} m_\rho^2 \rho^2 - \sum_i \frac{2}{(2 \pi)^3} \int d^3 k E_i^2,$$
from both terrestrial experiments and astrophysical observations (Lattimer & Lim 2012). By comparing the calculations of the SLC and SLCd parameter sets, we examine the effects of the symmetry energy on the mass–radius correlation of neutron stars, though the SLCd model can describe the available observations more consistently with conclusions based on other analyses in the literature. The parameterization for hyperons is elaborated in the following. The coupling of mesons with hyperons can generally be given in terms of the parameters $X_{\sigma Y}$, $X_{\omega Y}$, and $X_{\rho Y}$, which are ratios of the meson coupling with hyperons to that with nucleons. Because of the lack of strong constraints on these parameters (Weissenborn et al. 2011a), a variety of choices of these parameters varying roughly from 0.2 to unity were used in practical studies (Glendenning & Moszkowski 1991; Avancini & Menezes 2006). Considering that the nucleonic sector of our models respects chiral limits at high densities, we assume two cases for hyperons: the usual case (UC) that the hyperons have a similar medium effect to nucleons, and the separable case (SC) that the meson–hyperon coupling constant is separated into density-dependent and density-independent parts regardless of the chiral limit constraint on the strange sector in hyperons. Nevertheless, the hyperon potentials (Millener et al. 1988; Hausmann & Weise 1989; Fukuda et al. 1998),

\[ U_A^{(N)} = -30 \text{ MeV} = -U_{\Sigma}^{(N)}, \quad U_{\Xi}^{(N)} = -18 \text{ MeV}, \quad (4) \]
in nuclear matter at saturation density are used to preserve the relation between the vector and scalar meson coupling constants. Note that the repulsive $\Sigma$ hyperon potential is invoked here (Mares et al. 1995; Nouni et al. 2002). For the strange mesons, we adopt the density-independent coupling constants in both cases for simplicity. The potentials for the $\Lambda$ and $\Xi$ hyperons in $\Xi$ matter $U_{\Xi}^{(\Xi)} = U_{\Xi}^{(\Xi)} = U_{\Xi}^{(\Xi)} = -40 \text{ MeV}$ are used to obtain the coupling constants of the strange mesons (for details, see Jiang 2006; Schaffner-Bielich & Gal 2000).

The ratio parameters for coupling constants in the UC are assumed to be density independent. In the SC, we consider the following scaling functions for the meson–hyperon coupling constants that consist of two terms:

\[ \Phi_{\omega A}(\rho) = \Phi_{\omega A}(\rho_0) + \frac{1}{3} \Phi_{\omega N}(\rho_0) + \frac{2}{3} \Phi_{\omega N}(\rho), \]

\[ \Phi_{\omega \Xi}(\rho) = \frac{2}{3} \Phi_{\omega N}(\rho_0) + \frac{1}{3} \Phi_{\omega N}(\rho), \]

\[ \Phi_{\sigma Y}(\rho) = (1 - f_{\sigma Y}) \Phi_{\sigma N}(\rho_0) + f_{\sigma Y} \Phi_{\sigma N}(\rho), \quad (5) \]

where $\rho_0$ is the saturation density and $f_{\sigma Y}$ is an adjustable constant. The scaling functions $\Phi_{\omega A}$ and $\Phi_{\omega \Xi}$ for the $\rho$ meson are calculated the same way as those of the $\omega$ meson. Note that $\Phi_{\omega N}(\rho_0)$ and $\Phi_{\sigma N}(\rho_0)$ are just constants. The factors before $\Phi_{\sigma N}(\rho_0)$ and $\Phi_{\omega N}(\rho_0)$, implied from constituent quark compositions, play a role in averaging the coupling constant on the hadron level, between the density-dependent part originating from the chiral limit, and the density-independent part that is presumably attributed to the strange sector in hyperons. The form in Equation (5) produces the relation $\Phi_{\omega Y} \equiv \Phi_{\omega \Xi}$ at saturation density. Thus, we do not need to readjust the parameter $g_{\sigma Y}(\rho_0)$ as $f_{\sigma Y}$ changes. For a few choices of $f_{\sigma Y}$, we plot the meson–$\Lambda$ hyperon coupling constants in Figure 1. It is seen that the larger the $f_{\sigma Y}$, the smaller the $g_{\sigma Y}$ at high densities.

3. RESULTS AND DISCUSSIONS

Since the parameters for the nucleonic sector of the SLC and SLCd are clearly given in Jiang et al. (2007b), we only list the parameters for the hyperonic sector in Table 1. In the same case, all the parameters but $g^*_{\sigma \Xi}$ are the same for models SLC and SLCd in Table 1, because the only difference between models SLC and SLCd is that the latter has a softer symmetry energy than the former. Since the neutron star properties are rather insensitive to the coupling parameters of $\Sigma$ and $\Xi$ hyperons owing to their small fractions in the core of neutron stars, for simplicity, we take $f_{\sigma \Xi} = f_{\sigma \Xi} = f_{\sigma \Lambda}$ for the SC in the calculation. For a similar reason, we prefer the choice $X_{\omega \Lambda} = X_{\omega \Xi} = X_{\omega \Lambda}$ in the UC calculation, unless otherwise noted. For the $\rho$ meson in the UC, the usual relation $X_{\rho \Xi} = 2X_{\rho \Xi} = 2$ is used. In both the SC and UC, the $g^*_{\rho \Lambda}$ is zero. Note that all parameters used in this work for hyperons meet relation (4).

As a well-known consequence, the emergence of the hyperon degree of freedom results in the softening of the EOS. While models SLC and SLCd were constructed based on the BR scaling, the softening turns out to be too appreciable at high densities to stabilize the neutron star in the UC with relatively small $X_{\omega \Lambda}$. As shown in the upper panel of Figure 2, this is related to the rapid decrease of the nucleon effective mass, corresponding to an increasingly large scalar field that provides the attraction. As discovered earlier (Jiang et al. 2007a), the vector coupling constant is decisive in generating a stiff EOS at high densities. Thus, the stiffening of the EOS can follow from increasing the parameter $X_{\omega \Lambda}$. With a larger $X_{\omega \Lambda}$, for instance $X_{\omega \Lambda} = 0.9$, the EOS stiffens to recover the stability of neutron stars. For the SC, the EOS can be stiffened by increasing the parameter $f_{\sigma Y}$, since the latter results in a decrease of the scalar coupling constant, as shown in Figure 1. Generally, the EOS obtained with the SC is much stiffer than that with the UC. Meanwhile, the accelerating decrease of the baryon effective mass due to the inclusion of hyperons can be greatly suppressed in the SC, as shown in the lower panel of Figure 2. In Figure 2, the $\Lambda$ and $\Xi$ hyperon effective masses are also displayed. The upward shift of hyperon masses at high densities in the lower
panel of Figure 2 is due to the decrease of the source term (namely, the hyperon density) of the strange mesons, also seen below. The hyperon effective mass is much larger than the nucleon one. This would justify the use of different in-medium interactions for hyperons and nucleons in the SC.

The drop in the baryon effective mass in chemically equilibrated matter is closely associated to the onset density and fractions of hyperons. In Figure 3, we display the particle fractions as a function of density. We see that the fractions of hyperons. In Figure 3, we display the particle fractions in the UC and SC with the model SLC as a function of density. We see that the fractions of hyperons.

Remarkably, with the increase of density, the hyperon fractions in the SC trend downward until they disappear after reaching the maximum, as shown in the right panel of Figure 3. The occurrence of this phenomenon is due to the fact that the vector meson–hyperon coupling constant $g_{hV}$ has a weaker density dependence (see Equation (5)) than the meson–nucleon coupling constant $g_{hN}$. At high densities, $g_{hV}$ exceeds $g_{hN}$, and so do the vector potentials. The chemical equilibrium thus makes the hyperon Fermi momenta and fractions lower. As an application in studying properties of neutron stars, this actually results in the exclusion of hyperons in the core of neutron stars and accordingly the re-stiffening of the EOS at high densities. In addition to our scheme, we note that there have been other attempts to decrease the number of hyperons in neutron stars. For instance, different couplings of a new boson to hyperons and nucleons were proposed to decrease the number of hyperons in neutron stars. In addition to our scheme, we note that there have been other attempts to decrease the number of hyperons in neutron stars. In addition to our scheme, we note that there have been other attempts to decrease the number of hyperons in neutron stars.
up to $4.5 \rho_0$ of this density can be obtained. However, binding relation (4) should then be reduced to $U_{\Lambda}^N(\rho_0) = 0$ MeV. In this case, the hyperon fraction becomes so small that the hyperonic constituents become unimportant. On the quark level, it is interesting to see that the mechanism of small numbers of strange quarks in hybrid stars was explored within a specific quark model (Buballa et al. 2004).

We mention that the baryon fractions in neutron stars are also sensitive to symmetry energy. Jiang (2006) illustrated that the onset density for hyperons increases moderately with the softening of the symmetry energy. For the same reason, the onset density for hyperons in the SLCd is about $0.2 \rho_0$ larger than those in the SLC. To save space herein, we do not display particle fractions for the SLCd in a figure similar to Figure 3.

Although the emergence of hyperons is a cause for softening the nuclear EOS, the specific behavior relies on in-medium interactions, as clearly shown in Figure 4. Softening persists at high densities for the UC. However, the softening is succeeded by a stiffening in the SC where hyperons feel a different in-medium interaction from nucleons. Looking at the right panel of Figure 3, we see that the stiffening of the EOS occurs with the suppression of hyperon fractions. As the hyperon vanishes, the EOS returns to the normal EOS without hyperons. As shown in Figure 4, the EOS evolves to become stiffer than normal with increasing density. The neutron star matter thus transits to the normal isospin-asymmetric matter prior to the vanishing of hyperons. This eventually leaves a limited density window allowing the existence of hyperons in neutron stars. Interestingly, we find that the influence of the hyperonization in the SLCd model decreases clearly, as compared to the SLC model. This is mainly attributed to a larger onset density of hyperons with the SLCd due to a softening of the symmetry energy, as mentioned above.

In the UC, since the nucleon effective mass vanishes at a certain critical density, we need to consider the EOS beyond the critical density. Though the relationship between the chiral restoration and deconfinement occurrence is still discussed, it is nevertheless a convenient and usual way to neglect the distinction between them. Beyond the critical density, we thus adopt a quark matter EOS described by the MIT model with the appropriate bag parameter [(179.5 MeV)²] to smoothly connect the hadron matter and quark matter EOSs. The connection to the quark matter further softens the EOS, as shown in Figure 4.

We now turn to the consequences of hyperonizations on properties of neutron stars. In particular, it is interesting to see whether our hyperonization model can give neutron star properties that are compatible with the recent observation of the millisecond pulsar J1614−2230. The mass of this pulsar was accurately determined to be $1.976 \pm 0.04 M_\odot$ (Demorest et al. 2010). It was concluded that such a large mass can rule out almost all currently proposed hyperon equations of state. The mass–radius relation of neutron stars is obtained by solving the standard Tolman-Oppenheimer-Volkoff (TOV) equation with the specified nuclear EOS above. For the low-density crust, we adopt a quark matter EOS described by the MIT model with the appropriate bag parameter [(179.5 MeV)²] to smoothly connect the hadron matter and quark matter EOSs. The connection to the quark matter further softens the EOS, as shown in Figure 4.

Figure 4. EOS of isospin-asymmetric matter with models SLC and SLCd. Curves are obtained with the parameters listed in Table 1. The arrow indicates the critical point to quark matter. The result, including the muon but without hyperons, is also depicted for comparisons.

Figure 5. Mass–radius relation of neutron stars in the UC and SC with the SLC and SLCd. In the SC, two curves are obtained with $f_\sigma Y = 0.8$ and 0.9, respectively. The hatched areas give the probability distributions with 1σ (red) and 2σ (green) confidence limits for $r_{\mu h} \gg R$ summarized in Steiner et al. (2010). (A color version of this figure is available in the online journal.)
relation to the SLCd is almost independent of the parameter $f_{\sigma Y}$ in the SC, as shown in Figure 5. For the UC, it looks undoubtedly that the EOS is ruled out by the recent observation (Demorest et al. 2010), since the maximum mass of neutron stars in this case just sprints to $1.4 M_\odot$ with a much more compact size than the canonical neutron star without hyperons.

It is worth adding some discussion about the influence of the symmetry energy on the radii of neutron stars. It is now well established that the maximum mass of neutron stars is dominated by the high-density behavior of the EOS, while the radius is primarily determined by the slope of the symmetry energy at intermediate densities ($1-3\rho_0$; Lattimer & Prakash 2001, 2004; Steiner et al. 2005; Xu et al. 2009). The large difference between the radii of low-mass neutron stars obtained with SLC and SLCd can be attributed to the difference in the slope parameter $L$. As discussed in detail earlier in Fattoyev & Piekarewicz (2010), the central density of an 18–20 km star is near $\rho_0$, and the crust ends approximately at $(1/3-1/2)\rho_0$. In this density range, the pressure is dominated by the symmetry energy and not the incompressibility of symmetric nuclear matter. Our results shown in Figure 5 are consistent with those in Fattoyev & Piekarewicz (2010). For the $1.4 M_\odot$ neutron stars in the SC, we see, however, that the inclusion of hyperons moderately reduces the range of the neutron star radius. For instance, the radius ranges from 11.4 km with the SLCd to 12.3 km with the SLC for $f_{\sigma Y} = 0.9$, which is well situated in the domain extracted by Steiner et al. (2010). With the emergence of hyperons, the sensitivity of the star radius to the symmetry energy is clearly reduced. This is because the $\Lambda$ hyperon, being the dominant component of hyperons, is an isospin scalar. The suppression of the isovector potential $g_{\sigma Y}^p\rho_0$ due to the appearance of $\Lambda$ hyperons (Jiang 2006) results in the reduction of the asymmetric matter pressure and thus the star radius, while the magnitude of the suppression depends on the specific values of the symmetry energy and hyperon fraction. As a result, we observed in our calculations that the star radius is less sensitive to differences in the symmetry energy in both the SC and UC. In Figure 5, we also include the constraints of mass–radius trajectories for $r_{ph} \gg R$ obtained by Steiner et al. (2010). Our results in the SC are either within (for SLCd) or not very far off from (for SLC) the constrained region, though the inclusion of hyperons seems to tilt the vertical trajectories. For low-mass neutron stars, the radii with the SLC are predicted to go beyond the optimal region extracted very recently by Steiner et al. (2012), unless some other scenarios that allow a loose extension of the radii are invoked to extract the radius constraints (Suleimanov et al. 2011; Zhang et al. 2007). However, the maximum mass of neutron stars is almost independent of the slope $L$ and the radii of low-mass neutron stars. In our cases, the maximum mass is only reduced by about 1% when the symmetry energy is softened from model SLC to SLCd. On the other hand, it is known that the measurements of the neutron star radii are far less precise than the mass measurements, e.g., Suleimanov et al. (2011) and Zhang et al. (2007), and references therein. We note that an option of $r_{ph} = R$ can cause a visible slanting of the vertical trajectories (Steiner et al. 2010). The agreement of our results are thus better with the constraints obtained for $r_{ph} = R$.

We note that there have been a few endeavors in the past to involve hyperons in heavy neutron stars either by invoking strong repulsions for hyperons or pushing upward the onset density of hyperons (Hofmann et al. 2001; Takatsuka et al. 2002; Stone et al. 2007; Dexheimer & Schramm 2008; Bednarek et al. 2011; Weissenborn et al. 2011a). The main purpose of our work is to constraint the in-medium hyperon potentials using the two solar mass constraints of neutron stars. This indeed requires strong repulsions for hyperons. Due to different density dependencies of nucleonic and hyperonic potentials in the SC, the hyperonic vector potential exceeds the nucleonic one, leading to a very significant suppression of hyperons at high densities. As a result, the hyperons would exist in a shell of a neutron star core even with a small $L$ value.

Besides the effect of hyperonization on static properties, another consequence of hyperonization is on the thermal evolution of neutron stars. For the SC in SLC, we see from Figure 5 that $\Lambda$ hyperons start to appear above $2.5\rho_0$. At this central density the neutron star mass is about $1.2 M_\odot$. For most observed neutron stars that have larger masses, it appears that the direct Urca (DU) process with nucleons (Lattimer et al. 1991) and/or hyperons (Prakash et al. 1992) would occur when the proton fraction in the neutron star matter with hyperonization can be in excess of the DU threshold (14%) in the SC. According to the thermal evolution of observed neutron stars analyzed in Page et al. (2004) and Yakovlev & Pethick (2004), the fast cooling with the DU processes seemed to be excluded in most neutron stars except the massive ones. Slower cooling is possible when the neutrino emissivity can be suppressed by the superfluidity of constituent particles, such as nucleons and hyperons when the temperature falls below a critical temperature. Page et al. set a stringent requirement for the critical temperature of neutron superfluidity, without which enhanced cooling from DU processes may be needed in at least half of the observed young cooling neutron stars (Page et al. 2009). While the DU cooling involving only nucleons was regarded to be too fast, it was pointed out by Tsuruta et al. that the DU cooling with hyperons in neutron stars can be compatible with the observations, provided that hyperon superfluidity is appropriately accounted for (Tsuruta et al. 2009). In our models with the SC, hyperonization can thus be favorable in making up the potential incompatibility in the thermal evolution by considering hyperon superfluidity. On the other hand, the threshold mass of neutron stars which allow the DU process increases with the softening of the symmetry energy due to which the proton fraction exceeds the threshold value at larger densities. For instance, the mass is around $1.3 M_\odot$ with the SC of the model SLCd, and the threshold density for the DU process is about $4\rho_0$, which is coincidently the onset density of hyperons in Tsuruta et al. (2009).

Finally, we discuss some details concerning the in-medium interactions for hyperons. Constrained by the large mass of observed pulsars, it is favorable for us to select the in-medium interactions for hyperons in the SC. It is, however, interesting to find that the single-particle potentials for hyperons in the UC and SC almost overlap in a large density domain ranging from zero density to intermediate density. A significant departure appears only at high densities (roughly $\geq 4\rho_0$). At lower densities, the SC and UC descriptions give rather limited differences in single-particle properties. This thus indicates that without compromising the success in describing the properties of hypernuclei, one can significantly constrain the high-density hyperon interactions with the large mass of neutron stars. In addition, we have noticed that many studies using various quark models with the postulate of strong interactions and/or color superconductivity can also lead to large-mass hybrid stars (Alford et al. 2005; Klähn et al. 2007; Ippolito et al. 2008; Weissenborn et al. 2011b; Bonanno & Sedrakian 2012) with stiff EOSs at high densities. Our models, respecting chiral limit, possess a similarly stiff EOS at high densities. On the other
hand, neutron stars may be composed of more complicated constituents including quarks and meson condensates. It is useful to consider these non-baryonic degrees of freedom and their interplay with hyperons to more quantitatively constrain the in-medium interactions for hyperons. However, this is beyond the scope of the present work.

4. SUMMARY

The density-dependent RMF model, which was constructed to respect the chiral symmetry restoration at high densities in terms of in-medium hadron properties according to the BR scaling, is extended to include the hyperonization in isospin-asymmetric matter. We examined the effects of in-medium hyperon interactions on the properties of neutron stars. We found that the maximum mass of neutron stars can significantly constrain in-medium hyperon interactions. In particular, assuming two categories of in-medium interactions for hyperons, we investigated their distinct roles in hyperonizations and properties of neutron stars. With different in-medium interactions for hyperons and nucleons (i.e., the SC case), a maximum neutron star mass of $2M_\odot$ can be obtained with the model where the nucleonic EOS is consistent with the terrestrial nuclear laboratory data. The result in the SC is compatible with recent observations of the millisecond pulsar J1614−2230 mass. In this scheme, the number of hyperons is limited in neutron stars. Interestingly, we also found that the softening of the symmetry energy can play an important role in further reducing hyperons in neutron stars. On the contrary, the scheme that adopts similar in-medium interactions for hyperons, we investigated their distinct roles in hyperonizations and properties of neutron stars. We thank William Newton for his very helpful and constructive comments. J.W.Z. thanks Ang Li and Zhao-Qing Feng for useful discussions and expresses his sincere gratitude to Dr. Song Dan at the affiliated Hospital of Nanjing Medical University for his great help during the time of this work.
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