HADRON–QUARK CROSSOVER AND MASSIVE HYBRID STARS WITH STRANGENESS

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Received 2012 June 2; accepted 2012 December 5; published 2013 January 21

ABSTRACT

Using the idea of smooth crossover from hadronic matter with hyperons to quark matter with strangeness, we show that the maximum mass ($M_{\text{max}}$) of neutron stars with quark matter cores can be larger than those without quark matter cores. This is in contrast to the conventional softening of the equation of state due to exotic components at high density. The essential conditions for reaching our conclusion are that (1) the crossover takes place at relatively low densities, around three times the normal nuclear density and (2) the quark matter is strongly interacting in the crossover region. From these, the pressure of the system can be greater than that of purely hadronic matter at a given baryon density in the crossover density region and leads to $M_{\text{max}}$ greater than 2 solar mass. This conclusion is insensitive to the different choice of the hadronic equation of state with hyperons. We remark upon several implications of this result to the nuclear incompressibility, the hyperon mixing, and the neutrino cooling.

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

Online-only material: color figures

1. INTRODUCTION

Since neutron stars (NSs) are gravitationally bound, observation of the NS mass ($M$) provides a precious probe for the equation of state (EOS) and the relevant composition and structure of dense matter. The larger $M$ requires a stiffer EOS and thereby provides a stringent condition on the state of matter, since the various new phases proposed thus far, such as pion condensation, kaon condensation, hyperon ($Y$) mixing, deconfined quark matter, and so on, are believed to soften the EOS.

For example, the EOS is softened dramatically when $Y$ takes part in NS cores and the theoretical maximum mass fails to exceed the observed mass $M_{\text{obs}} = 1.44 M_{\odot}$ for PSR1913+16 with $M_{\odot}$ being the solar mass. This contradiction between theory and observation may indicate a missing repulsion in hyperonic matter such as the three-body interaction (TNI) acting universally in the nucleon and hyperon sectors (Nishizaki et al. 2001; Takatsuka et al. 2006, 2008).

Very recently, the observation of heavy NS with $M = (1.97 \pm 0.04) M_{\odot}$ for PSR J1614–2230 (Demorest et al. 2010) has been reported, which causes an even stronger constraint on the EOS. In particular, it stimulates intensive discussions on whether or not such a heavy NS can have exotic components in the central core (Özcel et al. 2010; Kurkela et al. 2010; Kim et al. 2011; Klahn et al. 2012; Lattimer & Prakash 2010; Weissenborn et al. 2011; Bonanno & Sedrakian 2012; Chen et al. 2012; Schramm et al. 2012; Whittenbury et al. 2012).

In particular, in many of the previous works treating the transition from the hadron phase to the quark phase at high density, point-like hadrons and quarks are treated as independent degrees of freedom, and Gibbs phase equilibrium conditions are imposed. However, quantum chromodynamics (QCD) tells us that such a description is not fully justified, since all hadrons are extended objects composed of quarks and gluons. Indeed, one may expect a gradual onset of quark degrees of freedom in dense matter associated with the percolation of finite-size hadrons, i.e., a smooth crossover from hadronic matter to quark matter; see seminal works on hadron percolation (Baym 1979; Celik et al. 1980), recent works on hadron–quark crossover (Baym et al. 2008; Maeda et al. 2009), and a recent review on phase transitions in dense QCD (Fukushima & Hatsuda 2011).

The aim of this paper is to perform a phenomenological study on the mechanism realizing hybrid stars (NSs with a quark core) compatible with 2 $M_{\odot}$, from the point of view of the smooth crossover between hadronic matter with hyperons and quark matter with strangeness. Unlike the traditional approach to the hadron–quark transition, the crossover picture gives us a novel tool for studying whether or not a hybrid star with a quark core is compatible with massive NSs (Takatsuka et al. 2012).

2. HADRONIC MATTER WITH STRANGENESS

Let us start with an example of hadronic EOS with hyperons (H-EOS; Nishizaki et al. 2001, 2002; Takatsuka et al. 2006) obtained using the following procedure. (1) Effective two-body and three-body potentials $V_{BB}$ ($B = n, p, \Lambda, \Sigma^-$) are constructed on the basis of the $G$-matrix formalism to take into account their density dependence. (2) A phenomenological three-body nucleon force expressed in the form of a two-body potential (Friedman & Pandharipande 1981) is introduced to reproduce the saturation of symmetric nuclear matter (the saturation density $\rho_0 = 0.17$ fm$^{-3}$ and the binding energy $E_0 = -16$ MeV) and the incompressibility $\kappa$. (3) We assume universal three-body repulsion even for the hyperons through replacement, $\hat{U}_{NN} \rightarrow \hat{U}_{BB}$. (4) By using $\hat{V}_{BB} + \hat{U}_{BB}$, we obtain the hadronic EOS under charge neutrality and $\beta$-equilibrium and calculate particle composition $\gamma_i$ ($i = n, p, \Lambda, \Sigma^-, e^-$, and $\mu^-$) as a function of total baryon density $\rho$. This gives us our H-EOS of NS matter with hyperon mixing.

The stiffness of H-EOS is specified by $\kappa$; the case with $\kappa = 300$ (250) MeV is denoted by TNI3u (TNI2u). Here, the subscript “u” means that the three-body force is introduced universally including the hyperon sector. The maximum mass ($M_{\text{max}}$) of the NSs with hyperon-mixed core with universal...
three-body force becomes $M_{\text{max}} \simeq 1.82(1.52)M_{\odot}$ for TNI3u (TNI2u). On the other hand, with only the nucleon three-body force (TNI3), it is reduced to $M_{\text{max}} \simeq 1.1M_{\odot}$, which is not even compatible with $M_{\text{obs}} = 1.44M_{\odot}$. One of the effects of the universal three-body force is to delay the onset of hyperon mixing, e.g., $\rho_{u,d,s} \simeq \rho_{u,d,s} \simeq 4\rho_{0}$, for the TNI3u and TNI2u cases, while $\rho_{u,d,s} \simeq 2.2\rho_{0}$, $\rho_{u,d,s} \simeq 2.5\rho_{0}$ for TNI3.

To check the model dependence of the H-EOS, we consider two alternative EOSs with hyperons: AV18+TBF+ΛΣ (Baldo et al. 2000), which is based on the $G$-matrix approach with the AV18 nucleon–nucleon potential Urbana-type three-body nucleon potential and the Nijmegen soft-core nucleon–hyperon potential, and SCL3ΛΣ (Tsubakihara et al. 2010), which is a relativistic mean-field model with chiral SU(3) symmetry. The maximum masses (central baryon densities, radius) of NSs composed only of hadronic EOSs, AV18+TBF+ΛΣ and SCL3ΛΣ, are $M_{\text{max}} \sim 1.22 M_{\odot}$ (7.35$\rho_{0}$, 10.46 km) and $M_{\text{max}} \sim 1.36 M_{\odot}$ (5.89$\rho_{0}$, 11.42 km), respectively.

In Figure 1, we plot the energy per particle ($E/A$) of the hadronic EOS with hyperons (TNI3u, TNI2u, TNI3, and SCL3ΛΣ) together with the nuclear EOS without hyperons (APR; Akmal et al. 1998). We do not show AV18+TBF+ΛΣ since it is almost the same as TNI3. The filled circles on each line denote the density where the hyperons start to mix. From this figure, one can see that (1) the mixture of hyperons softens the EOS relative to APR and (2) onset of the hyperon mixture is shifted to higher density if we consider the universal three-body interaction. In the following, we adopt TNI3u as a typical example of the H-EOS and comment on the results with other H-EOSs at the end.

3. QUARK MATTER WITH STRANGENESS

Next, we consider the EOS of quark matter with strangeness (Q-EOS). Around the baryon density of a few times $\rho_{0}$, where the hadron–quark crossover is expected to take place at zero temperature, the deconfined (or partially deconfined) quarks would be still strongly interacting. An analogous situation at high temperature has been expected theoretically and evidence of a strongly interacting quark–gluon plasma (sQGP) was found experimentally in relativistic heavy-ion collisions (see, e.g., a review Fukushima & Hatsuda 2011).

Since numerical simulations of QCD on the spacetime lattice at high baryon density are not available due to the notorious sign problem, we treat the strongly interacting quark matter (sQM) at zero temperature using the $(2+1)$-flavor Nambu–Jona-Lasinio (NJL) model (see the reviews by Vogl & Weise 1991; Klevansky 1992; Hatsuda & Kunihiro 1994; Buballa 2005). It is an effective theory of QCD and is particularly useful for taking into account the non-perturbative phenomena such as the partial restoration of chiral symmetry at high density. The Lagrangian model reads

$$\mathcal{L}_{\text{NJL}} = \overline{q}(i\gamma\partial - m)q + \frac{G_{5}^{2}}{2} \sum_{a=0}^{8} (\overline{q}^{a}\lambda^{a}q)^{2} + (\overline{q}^{a}\gamma_{5}\lambda^{a}q)^{2},$$

where the quark field $q_{i}$ ($i = u, d, s$) has three colors and three flavors with the current quark mass $m_{q}$. The term proportional to $G_{5}$ is a $U(3)_{u} \times U(3)_{s}$ symmetric four-fermi interaction, where $\lambda^{a}$ are the Gell-Mann matrices with $\lambda^{0} = \sqrt{2/3}I$. The term proportional to $G_{D}$ is the Kobayashi–Maskawa–t’Hooft six-fermi interaction which breaks $U(1)_{B}$ symmetry. The third term proportional to $g_{v}$ is a phenomenological vector-type interaction. It has some varieties depending on its flavor-structure: here we use the form given in Equation (1) which leads to an universal flavor-independent repulsion among quarks.

In the mean-field approximation, the constituent quark masses $M_{i}$ ($i = u, d, s$) are generated dynamically through the NJL interactions ($G_{5}, G_{D}$), $M_{i} = m_{i} - 2G_{5}\sigma_{i} - 2G_{D}\sigma_{j}\sigma_{k}$, where $\sigma_{i} = (\langle \bar{q}_{i}q_{i} \rangle)$ is the quark condensate in each flavor, and $(i, j, k)$ corresponds to the cyclic permutation of $u, d, s$. On the other hand, the vector interaction ($g_{v}$) leads to an effective chemical potential (Asakawa & Yazaki 1989), $\mu^{(\text{eff})}_{i} = \mu_{i} - g_{v} \sum_{j} \langle q_{j}q_{j} \rangle$. The basic parameters of the NJL model are determined from hadron phenomenology. In this paper, we adopt the so-called HK parameter set (Hatsuda & Kunihiro 1994), $\Lambda = 631.4$ MeV, $G_{5}A_{5}^{2} = 1.835$, $G_{D}A_{5}^{2} = 9.29$, $m_{u,d} = 5.5$ MeV, $m_{s} = 135.7$ MeV, where $\Lambda$ is the three-momentum cutoff. The magnitude of $g_{v}$ has not been well determined, but recent studies of the model applied to the QCD phase diagram suggest that it can be comparable to or even larger than $G_{5}$ (Bratovic et al. 2012; Lourenco et al. 2012); therefore, we change its value in the range,

$$0 \leq \frac{g_{v}}{G_{5}} \leq 1.5.$$ (2)

The EOS of sQM with strangeness is obtained from the above model under charge neutrality and $\beta$-equilibrium with $u, d, s, e^{-}$, and $\mu^{-}$. It turns out that the strange quarks start to appear at $\rho \simeq 4\rho_{0}$ regardless of the magnitude of $g_{v}$. Also, $\mu^{-}$ never appears under the presence of the strangeness.

4. HADRON–QUARK CROSSOVER

To realize the smooth transition between the hadronic matter and the quark matter, we first define a “crossover density
region” characterized by its central value $\bar{\rho}$ and its width $\Gamma$. The description of the matter in terms of the pure hadronic EOS is accurate for $\rho \ll \bar{\rho} - \Gamma$, while the description in terms of the pure quark EOS is accurate for $\rho \gg \bar{\rho} + \Gamma$. In the crossover region, $\bar{\rho} - \Gamma \leq \rho \leq \bar{\rho} + \Gamma$, both hadrons and quarks are strongly interacting, such that neither the pure hadronic EOS nor the pure quark EOS is reliable. In this situation, we perform a smooth interpolation of the two descriptions phenomenologically with a procedure similar to the one at high temperature (Asakawa & Hatsuda 1997):

$$P = P_H \times f_- + P_Q \times f_+,$$

$$f_\pm = \frac{1}{2} \left( 1 \pm \tanh \left( \frac{\rho - \bar{\rho}}{\Gamma} \right) \right),$$

where $P_H$ and $P_Q$ are the pressure in the pure hadron matter and in the pure quark matter, respectively. One should not confuse our interpolated pressure in Equation (3) with that of the “phase equilibrium condition” such as $P_Q = P_H$ at fixed chemical potential. (For phenomenological attempts to interpolate H-EOS and Q-EOS within the conventional picture of first-order phase transition, see, e.g., Burgio et al. 2002, and references therein.)

The energy density $\varepsilon$ as a function of $\rho$ is obtained by integrating the thermodynamical relation $P = \rho^2 (\varepsilon/\rho)/\partial \rho$, which guarantees the thermodynamical consistency for any values of $\bar{\rho}$ and $\Gamma$. To explore the relation between the interpolated EOS and the maximum mass of NSs, we change the parameters ($\bar{\rho}$, $\Gamma$) under two conditions: (1) the system is always thermodynamically stable $dP/d\rho > 0$ and (2) the normal nuclear matter is described dominantly by $P_H$, i.e., the condition $\bar{\rho} - 2\Gamma > \rho_0$.

Shown in Figure 2 is the pressure $P$ as a function of baryon density $\rho$ obtained by interpolating TNI3u H-EOS (the dashed line) and NJL Q-EOS (the dash-dotted line) with $g_v/G_S = 0, 1.0, 1.5$. The dashed line denotes hadronic matter with TNI3u interaction. (A color version of this figure is available in the online journal.)

Figure 2. Solid line: pressure ($P$) as a function of baryon density $\rho$ obtained by interpolating H-EOS and Q-EOS with $g_v = G_S$ in the crossover region (the shaded region in density). Dash-dotted line: Q-EOS in the NJL model with $g_v = G_S$. Dashed line: H-EOS with the TNI3u interaction. The inset shows EOSs interpolated with $g_v/G_S = 0, 1.0, 1.5$. (A color version of this figure is available in the online journal.)

In Figure 3, we plot the pressure $P$ as a function of the energy density $\varepsilon$ for the three EOSs interpolated in the crossover region, ($\bar{\rho}$, $\Gamma$) = (3$\rho_0$, $\rho_0$). These curves are the basic input when we solve the Tolman–Oppenheimer–Volkov (TOV) equation. The case for H-EOS with TNI3u is also plotted with the dashed line for comparison. The existence of a region where $P$ (the solid line) becomes larger than $P_H$ (the dashed line) is essential for having a large $M_{\text{max}}$, as shown below.

Figure 3. Solid lines: $P$ vs. energy density ($\varepsilon$) for the interpolated EOSs with $g_v/G_S = 0, 1.0, 1.5$. The dashed line denotes hadronic matter with TNI3u interaction. (A color version of this figure is available in the online journal.)

5. HYBRID STAR STRUCTURE

In Figure 4, the solid lines indicate NS masses as a function of the central density $\rho_c$ which is obtained by solving the TOV equation with the EOSs given in the inset of Figure 2, i.e., interpolated EOSs between TNI3u H-EOS and NJL Q-EOS with $g_v/G_S = 0, 1.0, 1.5$. Results with the hadronic EOS only (TNI3u, AV18+TBF+ΛΣ, and SCL3ΛΣ) are also plotted with the dashed lines for comparison. The filled circle indicates a point beyond which the strangeness appears in the central core.
of the star. The cross symbols indicate the points where \( M_{\text{max}} \) is achieved.

For TNI3u H-EOS only, we have \( M_{\text{max}} = 1.82 M_\odot \), while for the interpolated EOS with \( g_y/G_0 \geq 1 \), we have \( M_{\text{max}} > 2 M_\odot \). This indicates that the existence of a quark core inside NSs is compatible with the observation of a high-mass NS as long as (1) there is a smooth crossover between the hadronic phase and the quark phase and (2) there exists a strong quark correlation with a repulsive nature inside the quark matter.

In Figure 5, the solid lines indicate the mass \( (M)\)–radius \( (R) \) relation of the NSs with the interpolated EOSs given in the inset of Figure 2. The results with the H-EOS only (TNI3u, AV18+TBF+ΛΣ, and SCL3ΛΣ) are also plotted for comparison with the dashed lines. The gray band shows \( M = (1.97 \pm 0.04) M_\odot \) corresponding to PSR J1614–2230 (Demorest et al. 2010). Because of the stiff EOS at high density for \( g_y/G_0 \geq 1 \), the radius \( R \) for \( 0.5 \lesssim M/M_\odot \lesssim 2 \) takes almost a constant value, 11–11.5 km (Özel et al. 2012). Also, strangeness only emerges for massive stars with \( M \gtrsim 2 M_\odot \).

Thus far, we have taken \((\bar{\rho}, \Gamma) = (3\rho_0, \rho_0)\) as characteristic parameters for the crossover region. In Table 1, we show how \( M_{\text{max}} \) and \( \rho_c \) depend on the choice of these parameters for fixed values of \( g_y \). As the crossover region becomes lower and/or wider in baryon density, the EOS characteristic to the NS core \((\rho = (2-4)\rho_0)\) becomes stiffer. This is why \( M_{\text{max}} > 2 M_\odot \) becomes possible for \( \bar{\rho}/\rho_0 = 3 \) and 4.

Now, we examine how \( M_{\text{max}} \) depends on the different choices of H-EOSs. Shown in Table 2 are the maximum and central density for different H-EOSs with \( \bar{\rho} = 3\rho_0 \) and \( \Gamma/\rho_0 = 1 \). As far as \( M_{\text{max}} \) is concerned, the result is insensitive to the choice of the hadronic EOS with hyperons.

6. SUMMARY AND CONCLUDING REMARKS

In this paper, we have investigated the mass \( (M)\)–radius \( (R) \) relation of the NSs with the interpolated EOSs given in the inset of Figure 2. The results with the H-EOS only (TNI3u, AV18+TBF+ΛΣ, and SCL3ΛΣ) are also plotted for comparison with the dashed lines. The gray band shows \( M = (1.97 \pm 0.04) M_\odot \) corresponding to PSR J1614–2230 (Demorest et al. 2010). Because of the stiff EOS at high density for \( g_y/G_0 \geq 1 \), the radius \( R \) for \( 0.5 \lesssim M/M_\odot \lesssim 2 \) takes almost a constant value, 11–11.5 km (Özel et al. 2012). Also, strangeness only emerges for massive stars with \( M \gtrsim 2 M_\odot \).

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We note that TNI3, A V18+TBF+M lighter than 1.92 due to the universal three-body force, such that the hybrid stars and thereby leads to smaller $M_{\text{max}}$.

Although we have mainly focused on the results using the TNI3u H-EOS corresponding to $\kappa = 300$ MeV, very similar results are obtained for TNI2u, as shown in Table 2. This introduces an interesting possibility that we can reconcile the existence of massive NSs ($M > 2M_\odot$) with the experimental nuclear incompressibility $\kappa = (240 \pm 20)$ MeV (Shlomo et al. 2006). We note that TNI3, AV18+TBF+$\Lambda\Sigma$, and SCL3$\Lambda\Sigma$ give almost the same $M_{\text{max}}$ as that of TNI3u and TNI2u. This does not, however, imply the irrelevance of the universal three-body force; for example, the interpolated EOS with H-EOS (TNI3u) and Q-EOS ($g_s = G_s$) leads to strangeness mixing only above $4\rho_0$ due to the universal three-body force, such that the hybrid stars lighter than $1.92 M_\odot$ do not have strangeness (see Figure 4). Then, the hyperon direct Urca process (e.g., $\Lambda \rightarrow p + e^- + \bar{\nu}_e$, $p + e^- \rightarrow \Lambda + \nu_e$), which leads to extremely rapid neutrino-cooling and contradicts observations (Takatsuka et al. 2006), does not take place except for massive NSs.

Our massive hybrid stars with $M > 2M_\odot$ originate from the transition from the soft EOS to stiff EOS in the density region $\rho \sim (2-4)\rho_0$ due to the hadron–quark crossover. Therefore, it is important to explore such a crossover using independent laboratory experiments with medium-energy heavy-ion collisions.

Finally, we note that the crossover region may have rich non-perturbative phases such as color superconducting phases, inhomogeneous phases, the quarkyonic phase, and so on (Fukushima & Hatsuda 2011). How these structures affect the basic conclusion of this paper would be an interesting future problem to examine.

K.M. and T.H. thank Wolfram Weise for discussions. T.T. thanks Ryozo Tamagaki, Toshitaka Tatsumi, and Shigeru Nishizaki for discussions and interest in this work. We also thank K. Tsukahara and A. Ohnishi for providing us with the numerical data for the SCL3 EOS. This research was supported in part by MEXT Grant-in-Aid for Scientific Research on Innovative Areas (No. 2004:20105003) by JSPS Grant-in-Aid for Scientific Research (B) No. 22340052, and by RIKEN 2012 Strategic Programs for R&D.

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