

Quark deconfinement in high-mass neutron stars

M. Orsaria*

*Department of Physics, San Diego State University,
5500 Campanile Drive, San Diego, CA 92182, USA and
CONICET, Rivadavia 1917, 1033 Buenos Aires, Argentina;
Gravitation, Astrophysics and Cosmology Group,
Facultad de Ciencias Astronómicas y Geofísicas, UNLP,
Paseo del Bosque S/N (1900), La Plata, Argentina*

H. Rodrigues†

*Centro Federal de Educação Tecnológica do Rio de Janeiro,
Av Maracanã 249, 20271-110, Rio de Janeiro, RJ, Brazil*

F. Weber‡

*Department of Physics, San Diego State University,
5500 Campanile Drive, San Diego, California 92182 and
Center for Astrophysics and Space Sciences, University of California,
San Diego, La Jolla, CA 92093, USA*

G. A. Contrera§

*CONICET, Rivadavia 1917, 1033 Buenos Aires, Argentina
IFLP, CONICET - Dpto. de Física, UNLP, La Plata, Argentina, and
Gravitation, Astrophysics and Cosmology Group,
Facultad de Ciencias Astronómicas y Geofísicas, UNLP,
Paseo del Bosque S/N (1900), La Plata, Argentina*

arXiv:1308.1657v1 [nucl-th] 7 Aug 2013

Abstract

In this paper, we explore whether or not quark deconfinement may occur in high-mass neutron stars such as J1614-2230 ($1.97 \pm 0.04 M_{\odot}$) and J0348+0432 ($2.01 \pm 0.04 M_{\odot}$). Our study is based on a non-local extension of the SU(3) Nambu Jona-Lasinio (n3NJL) model with repulsive vector interactions among the quarks. This model goes beyond the frequently used local version of the Nambu Jona-Lasinio model by accounting for several key features of QCD which are not part of the local model. We find that the n3NJL model predicts the existence of an extended region of a mixed phase of quarks and hadrons (quark-hybrid matter) in high-mass neutron stars of masses up to around $2.1 M_{\odot}$. Central stellar cores made of pure quark matter are not obtained. The radii of high-mass neutron stars with quark-hybrid matter in their centers are found to lie in the canonical range of 12 to 13 km.

* morsaria@rohan.sdsu.edu

† harg@cefet-rj.br

‡ fweber@mail.sdsu.edu

§ guscontrera@gmail.com

I. INTRODUCTION

White dwarf and neutron stars (NSs) are born in the aftermath of core-collapsing supernova explosions. Depending on NS mass and rotational frequency, gravity may compress the matter in the core regions of such stars up to more than ten times the density of ordinary atomic nuclei, thus providing a high-pressure environment in which numerous subatomic particle processes are likely to compete each other. Theoretical studies indicate that hyperons, boson condensates (pions, kaons, H-matter) and/or deconfined up, down and strange quarks may exist in the core regions of NSs (for an overview, see [1–12] and references therein). As it has been pointed out in the literature [13–15], there exists the very intriguing theoretical possibility that strange quark matter [16, 17] made of free up, down and strange quarks may be more stable than atomic nuclei. This is known in the literature as the strange quark matter (SQM) hypothesis. If true, it would have implications of fundamental importance for cosmology, the early universe, its evolution to the present day, and astrophysical compact objects such as neutron stars and white dwarfs (see [1, 2, 5, 6, 17–20] and references therein).

Based on qualitative considerations concerning the stiffness of the nuclear equation of state (EoS), one could argue that the observations of high-mass NSs, such as PSR J1614–2230 with a gravitational mass of $1.97 \pm 0.04 M_{\odot}$ [21] and PSR J0348+0432 with a mass of $2.01 \pm 0.04 M_{\odot}$ [22], do not contain deconfined quarks in their central core regions, since quark deconfinement would take away so much pressure that high-mass NSs could not exist. As shown in Ref. [23], conclusions of that kind are not necessarily correct.

Our paper builds on the investigations carried out in Ref. [23] for NSs containing deconfined quark matter, i.e. quark-hybrid stars (QHSs). The study is based on a generalized version of the Nambu Jona-Lasinio (NJL) model, which incorporates the basic symmetries of quantum chromodynamics (QCD). The NJL model [24] was proposed in 1961 to study the properties of hadrons, but during the 1980s it was modified to systems of interacting quarks. In this model, chiral symmetry constraints are taken into account via effective interactions between quarks, through local four point vertex interactions. Since the NJL model with local interactions between quarks is not renormalizable, it needs to be regularized to avoid divergences in the loop integrals. Furthermore, the model is non-confining, basically due to the fact that in the local NJL model the dynamically generated constituent quark masses are momentum independent. Since the 1990s, there have been studies [25] proposing

non-local interactions to solve these problems, at the expense of complications in the calculations. One interesting suggestion arises from the relationship between the NJL model and the one-gluon-exchange model, which incorporates an effective gluon propagator in order to carry describe effective interactions between quarks. This provides a natural way to introduce the non-locality in the quark-quark interaction [26] which can be characterized by a model-dependent form factor $g(p)$. Table I provides a comparison of the key features of the standard (local) NJL with those of the non-local NJL model.

In this work, we model the quark phase of a quark-hybrid neutron star using the non-local 3-flavor NJL (n3NJL) model of Refs. [30] and [31], which includes vector interactions among the quarks. In our previous work [23], we considered the vector interaction of the non-local NJL model in a phenomenological way, and found that the transition to pure quark matter occurs only in neutron stars which lie already on the gravitationally unstable branch of the stellar sequence. In this paper we improve on the n3NJL model by including the effect of the shift of the chemical potential on the form factor, when the vector interaction is considered. The results are compared with the those obtained by modeling the quark phase in the framework of the local SU(3) NJL model (l3NJL) described, for example, in Refs. [32–35]. Vector interactions are know to be important for the QCD phase diagram [36]. It is therefore interesting (if not mandatory) to explore the influence of the vector interactions on the EoS of NS/QHS matter and the structure of QHSs.

Global electric charge neutrality is imposed on the constituents of neutron star matter. Local NJL studies carried out for local electric charge neutrality have been reported recently in the Refs. [37] and [38]. In [37], it was found that massive hybrid stars with a quark core are compatible with the mass observation of PSR J1614–2230 if there is a smooth crossover between the hadronic and quark matter phases and strong repulsive correlations among the quarks in the quark matter phase exist. In reference [38], a non-linear Walecka model was used for the hadronic phase, using the parametrizations GM1 [39] and NL3 [40], and the l3NJL model for the quark phase. In that work it has been found that the observation of compact star masses a few percent larger than $1.97 \pm 0.04 M_{\odot}$ are hard to explain unless one uses a very stiff hadronic model for the EoS, such as NL3 with nucleons only, instead of the softer GM1 EoS.

This work is organized as follows. In Sect. II, we describe the local as well as the non-local extension of the SU(3) NJL model at zero temperature. In Sect. III, the non-linear relativistic

Walecka model, which is used for the confined hadronic matter, is briefly discussed. In Sect. IV, the construction of the mixed quark-hadron phase is analyzed by imposing global electric charge neutrality on the constituents of neutron star matter. Our results for the quark-hadron composition and bulk properties of neutron stars are presented in Sects. V and VI. Finally, a summary and discussion of our results is provided in Sect. VII.

II. QUARK MATTER PHASE

A. The local 3-flavor NJL model with vector coupling (l3NJL)

The NJL model accounts for the interactions between constituent quarks and provides a simple scheme for studying spontaneous chiral symmetry breaking, a key feature of quantum chromodynamics (QCD) in the low temperature and density domain, and its manifestations in hadron physics, such as dynamical quark mass generation, the appearance of quark pair condensates, and the role of pions as Goldstone bosons. In its simplest form, the NJL model uses only a local scalar-isoscalar and pseudoscalar-isovectorial interaction between fermions. However, in the framework of the l3NJL model that we use, the effective action is given by

$$\begin{aligned}
S_E = \int d^4x \left\{ \bar{\psi}(x)(i\rlap{\not{D}} - \hat{m})\psi(x) + \frac{1}{2} G_S [(\bar{\psi}(x)\lambda_a\psi(x))^2 + (\bar{\psi}(x)i\gamma_5\lambda_a\psi(x))^2] \right. \\
+ H [\det[\bar{\psi}(x)(1 + \gamma_5)\psi(x)] + \det[\bar{\psi}(x)(1 - \gamma_5)\psi(x)]] \\
\left. - G_V [(\bar{\psi}(x)\gamma^\mu\lambda_a\psi(x))^2 + (\bar{\psi}(x)i\gamma^\mu\gamma_5\lambda_a\psi(x))^2] \right\}, \quad (1)
\end{aligned}$$

where ψ is a chiral U(3) vector that includes the light quark fields, $\psi \equiv (u, d, s)^T$, $\hat{m} = \text{diag}(m_u, m_d, m_s)$ is the current quark mass matrix, λ_a with $a = 1, \dots, 8$ stands for the eight generators of SU(3), and $\lambda_0 = \sqrt{2/3} \mathbb{1}_{3 \times 3}$. The values for the coupling constants G_S and H as well as the strange quark mass m_s and the three-momentum ultraviolet cutoff parameter Λ , are usually taken as model parameters. In this work we consider the values determined in Ref. [33], i.e., $m_u = m_d = 5.5$ MeV, $m_s = 140.7$ MeV, $\Lambda = 602.3$ MeV, $G_S\Lambda^2 = 3.67$ and $H\Lambda^5 = -12.36$. The vector coupling constant G_V is treated as a free parameter.

At the mean-field level, the thermodynamic potential at zero temperature reads

$$\begin{aligned} \Omega^L(M_f, \mu) = & G_S \sum_{f=u,d,s} \langle \bar{\psi}_f \psi_f \rangle^2 + 4H \langle \bar{\psi}_u \psi_u \rangle \langle \bar{\psi}_d \psi_d \rangle \langle \bar{\psi}_s \psi_s \rangle \\ & - 2N_c \sum_{f=u,d,s} \int_{\Lambda} \frac{d^3 p}{(2\pi)^3} E_f - \frac{N_c}{3\pi^2} \sum_{f=u,d,s} \int_0^{p_{Ff}} dp \frac{p^4}{E_f} - G_V \sum_f \rho_f^2 \end{aligned} \quad (2)$$

where $E_f = \sqrt{\mathbf{p}^2 + M_f^2}$ and $p_{Ff} = \sqrt{\mu_f^2 - M_f^2}$. The constituent quark masses M_f are given by

$$M_f = m_f - 2G_S \langle \bar{\psi}_f \psi_f \rangle - 2H \langle \bar{\psi}_j \psi_j \rangle \langle \bar{\psi}_k \psi_k \rangle, \quad (3)$$

with $f, j, k = u, d, s$ indicating cyclic permutations. The vector interaction shifts the quark chemical potential according to

$$\mu_f \rightarrow \mu_f - 2G_V \rho_f, \quad (4)$$

where ρ_f is the quark number density corresponding to the flavor f in the mean field approximation, that is,

$$\rho_f = \frac{N_c}{3\pi} [(\mu_f - 2G_V \rho_f)^2 - M_f^2]^{3/2}. \quad (5)$$

The mean field values are determined from the solution of the ‘‘gap’’ equations, obtained by minimizing the thermodynamic potential with respect to the quark condensates $\langle \bar{\psi}_f \psi_f \rangle$, that is,

$$\frac{\partial \Omega^L}{\partial \langle \bar{\psi}_f \psi_f \rangle} = 0, \quad f = u, d, s. \quad (6)$$

B. The non-local 3-flavor model with vector coupling (n3NJL)

In this section we briefly describe the non-local extension of the SU(3) Nambu Jona-Lasinio (n3NJL) model. The Euclidean effective action for the quark sector, including vector coupling, is given by

$$\begin{aligned} S_E = \int d^4 x \left\{ \bar{\psi}(x) [-i\hat{\not{D}} + \hat{m}] \psi(x) - \frac{G_S}{2} [j_a^S(x) j_a^S(x) + j_a^P(x) j_a^P(x)] \right. \\ \left. - \frac{H}{4} T_{abc} [j_a^S(x) j_b^S(x) j_c^S(x) - 3 j_a^S(x) j_b^P(x) j_c^P(x)] - \frac{G_V}{2} j_V^\mu(x) j_V^\mu(x) \right\}, \end{aligned} \quad (7)$$

where ψ and \hat{m} stands for the light quark fields and the current quark mass matrix, respectively. For simplicity, we consider the isospin symmetric limit in which $m_u = m_d = \bar{m}$. The

operator $\not{\partial} = \gamma_\mu \partial_\mu$ in Euclidean space is defined as $\vec{\gamma} \cdot \vec{\nabla} + \gamma_4 \frac{\partial}{\partial \tau}$, with $\gamma_4 = i\gamma_0$. The currents $j_a^{S,P}(x)$ and $j_V^\mu(x)$ are given by

$$\begin{aligned} j_a^S(x) &= \int d^4z \tilde{g}(z) \bar{\psi} \left(x + \frac{z}{2} \right) \lambda_a \psi \left(x - \frac{z}{2} \right) , \\ j_a^P(x) &= \int d^4z \tilde{g}(z) \bar{\psi} \left(x + \frac{z}{2} \right) i \gamma_5 \lambda_a \psi \left(x - \frac{z}{2} \right) , \\ j_V^\mu(x) &= \int d^4z \tilde{g}(z) \bar{\psi} \left(x + \frac{z}{2} \right) \gamma^\mu \psi \left(x - \frac{z}{2} \right) , \end{aligned} \quad (8)$$

where $\tilde{g}(z)$ is a form factor responsible for the non-local character of the interaction and λ_a represent the same matrices as in the local model.

Finally, the constants T_{abc} in the t'Hooft term accounting for flavor-mixing are defined by

$$T_{abc} = \frac{1}{3!} \epsilon_{ijk} \epsilon_{mnl} (\lambda_a)_{im} (\lambda_b)_{jn} (\lambda_c)_{kl} . \quad (9)$$

After the standard bosonization of Eq. (7), the integrals over the quark fields can be performed in the framework of the Euclidean four-momentum formalism. The grand canonical potential, for the mean-field approximation at zero temperature, can then be written as

$$\begin{aligned} \Omega^{NL}(M_f, 0, \mu_f) &= -\frac{N_c}{\pi^3} \sum_{f=u,d,s} \int_0^\infty dp_0 \int_0^\infty dp \ln \left\{ [\hat{\omega}_f^2 + M_f^2(\omega_f^2)] \frac{1}{\omega_f^2 + m_f^2} \right\} \\ &\quad -\frac{N_c}{\pi^2} \sum_{f=u,d,s} \int_0^{\sqrt{\mu_f^2 - m_f^2}} dp p^2 [(\mu_f - E_f)\theta(\mu_f - m_f)] \\ &\quad -\frac{1}{2} \left[\sum_{f=u,d,s} (\bar{\sigma}_f \bar{S}_f + \frac{G_S}{2} \bar{S}_f^2) + \frac{H}{2} \bar{S}_u \bar{S}_d \bar{S}_s \right] - \sum_{f=u,d,s} \frac{\varpi_f^2}{4G_V} , \end{aligned} \quad (10)$$

where $N_c = 3$, $E_f = \sqrt{p^2 + m_f^2}$, and $\omega_f^2 = (p_0 + i\mu_f)^2 + p^2$. The constituent quark masses M_f are treated as momentum-dependent quantities, given by

$$M_f(\omega_f^2) = m_f + \bar{\sigma}_f g(\omega_f^2) , \quad (11)$$

where $g(\omega_f^2)$ is the Fourier transform of the form factor $\tilde{g}(z)$. The vector mean fields ϖ_f are associated to the vector current densities $j_V^\mu(x)$, where a different vector field for each quark flavor f has been considered. The inclusion of vector interactions shifts the quark chemical potential as

$$\mu_f \rightarrow \hat{\mu}_f = \mu_f - g(\omega_f^2) \varpi_f . \quad (12)$$

This modifies the quantity ω_f according to

$$\omega_f^2 \rightarrow \widehat{\omega}_f^2 = (p_0 + i\widehat{\mu}_f)^2 + p^2. \quad (13)$$

We followed the prescriptions given in [41] to include the vector interaction. Note that the quark chemical potential shift does not affect the non-local form factor $g(\omega_f^2)$, as discussed in [41, 42], avoiding a recursive problem.

In this work we use a Gaussian form for the non-local form factor,

$$g(\omega_f^2) = \exp(-\omega_f^2/\Lambda^2), \quad (14)$$

where Λ is a parameter of the model, playing the role of an ultraviolet cut-off momentum scale. This parameter, as well as the quark current masses and coupling constants in Eq. (7), can be chosen so as to reproduce the phenomenological values of pion decay constant f_π , and the meson masses m_π , m_η , $m_{\eta'}$, as described in Refs. [30, 31]. In this work, for the n3NJL model we use the parameters listed in Table II, taken from Ref. [23].

Within the stationary phase approximation, the mean-field values of the auxiliary fields \bar{S}_f turn out to be related to the mean-field values of the scalar fields $\bar{\sigma}_f$ [43] and are given by

$$\bar{S}_f = -16 N_c \int_0^\infty dp_0 \int_0^\infty \frac{dp}{(2\pi)^3} g(\omega_f^2) \frac{M_f(\omega_f^2)}{\widehat{\omega}^2 + M_f^2(\omega_f^2)}. \quad (15)$$

The mean field values of $\bar{\sigma}_u$, $\bar{\sigma}_s$ and ϖ_f are obtained from the corresponding ‘‘gap’’ equations

$$\frac{\partial\Omega^{NL}}{\partial\bar{\sigma}_f} = 0, \quad \frac{\partial\Omega^{NL}}{\partial\varpi_f} = 0. \quad (16)$$

III. HADRONIC MATTER PHASE

The hadronic phase is described in the framework of non-linear relativistic field theory [44, 45], where baryons (neutrons, protons, hyperons, delta states) interact via the exchange of scalar, vector and isovector mesons (σ , ω , ρ , respectively). The parametrizations used in our study are GM1 and GM3 given in Ref. [39], and NL3 of Ref. [40]. They are summarized in Table III.

The lagrangian is given by

$$\mathcal{L} = \mathcal{L}_H + \mathcal{L}_\ell, \quad (17)$$

with

$$\mathcal{L}_\ell = \sum_{\lambda=e^-, \mu^-} \bar{\psi}_\lambda (i\gamma_\mu \partial^\mu - m_\lambda) \psi_\lambda \quad (18)$$

the leptonic lagrangian, and

$$\begin{aligned} \mathcal{L}_H = & \sum_{B=n,p,\Lambda,\Sigma,\Xi} \bar{\psi}_B [\gamma_\mu (i\partial^\mu - g_\omega \omega^\mu - g_\rho \vec{\rho}_\mu) - (m_N - g_\sigma \sigma)] \psi_B \\ & + \frac{1}{2} (\partial_\mu \sigma \partial^\mu \sigma - m_\sigma^2 \sigma^2) - \frac{1}{3} b_\sigma m_N (g_\sigma \sigma)^3 - \frac{1}{4} c_\sigma (g_\sigma \sigma)^4 \\ & - \frac{1}{4} \omega_{\mu\nu} \omega^{\mu\nu} + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu + \frac{1}{2} m_\rho^2 \vec{\rho}_\mu \cdot \vec{\rho}^\mu - \frac{1}{4} \vec{\rho}_{\mu\nu} \vec{\rho}^{\mu\nu}, \end{aligned} \quad (19)$$

being the hadronic lagrangian. The quantity B sums all hyperon states which become populated in neutron star matter at a given density [1, 2]. Intriguingly (see Sect. VI), we find that, aside from hyperons, the Δ^- state becomes populated in neutron star matter at densities which are encountered in the cores of stable neutron stars. Simpler treatments of the quark-hadron phase transition based on the MIT bag model [46, 47] do not predict the occurrence of the Δ^- state. The quantities g_ρ , g_σ , and g_ω in Eq. (19) are meson-baryon coupling constants whose values are summarized in in Table III.

Table IV lists the properties of symmetric nuclear matter computed from Eq. (17) for the relativistic mean-field approximation. The most important differences between the three parameter sets concern the values of the nuclear incompressibility and the asymmetry energy.

IV. QUARK-HADRON MIXED PHASE

For chemically equilibrated quark matter at finite density, the basic particle processes involved are given by the strong process $u + d \leftrightarrow u + s$ as well as the weak processes $d(s) \rightarrow u + e^-$ and $u + e^- \rightarrow d(s)$. We are assuming that neutrinos, once created by the weak reactions, leave the system, which is equivalent to supposing that the neutrino chemical potential is equal to zero. Then, the chemical potential for each quark flavor f is given by

$$\mu_f = \mu_b - Q \mu_e, \quad (20)$$

where $Q = \text{diag}(2/3, -1/3, -1/3)$ in flavor space and $\mu_b = 1/3 \sum_f \mu_f$ is the baryonic chemical potential.

The contribution from free degenerate leptons to the quark phase is given by

$$\Omega_{\lambda=e^-, \mu^-}(\mu_e) = -\frac{1}{\pi^2} \int_0^{P_{F_\lambda}} p^2 \left(\sqrt{p^2 + m_\lambda^2} - \mu_e \right) dp. \quad (21)$$

Muons occur in the system if the electron chemical potential $\mu_e = \mu_\mu$ is greater than the muon rest mass, $m_\mu = 105.7$ MeV. For electrons we have $m_e = 0.5$ MeV. The total thermodynamical potential for the quarks phase is given by Eq. (2) for the l3NJL model or by Eq. (10) for the n3NJL model, supplemented in either case with the leptonic contribution of Eq. (21).

If the dense interior of a neutron star is indeed converted to quark matter, it must be three-flavor quark matter since it has lower energy than two-flavor quark matter. And just as for the hyperon content of neutron stars, strangeness is not conserved on macroscopic time scales, which allows neutron stars to convert confined hadronic matter to three-flavor quark matter until equilibrium brings this process to a halt. As first realized by Glendenning [1, 46, 47], the presence of quark matter enables the hadronic regions of the mixed phase to arrange to be more isospin symmetric than in the pure phase by transferring charge to the quark phase in equilibrium with it. The symmetry energy will be lowered thereby at only a small cost in rearranging the quark Fermi surfaces. The electrons play only a minor role when neutrality can be achieved among the baryon-charge carrying particles. The stellar implication of this charge rearrangement is that the mixed phase region of the star will have positively charged regions of nuclear matter and negatively charged regions of quark matter [1, 46, 47]. First studies of the transport properties of such matter have been reported in [48, 49]

To determine the mixed phase region of quarks and hadrons, we start from the Gibbs condition for pressure equilibrium between confined hadronic (P^H) matter and deconfined quark (P^q) matter. The Gibbs condition is given by

$$P^H(\mu_b^H, \mu_e^H, \{\phi\}) = P^q(\mu_b^q, \mu_e^q, \{\psi\}), \quad (22)$$

with $\mu_b^H = \mu_b^q$ for the baryon chemical potentials and $\mu_e^H = \mu_e^q$ for the electron chemical potentials in the hadronic (H) and quark (q) phase. By definition, the quark chemical potential is given by $\mu_b^q = \frac{\mu_n}{3}$ where μ_n is the chemical potential of the neutron. The quantities $\{\phi\}$ and $\{\psi\}$ in Eq. (22) stand collectively for the field variables and Fermi momenta that characterize the solutions to the equations of confined hadronic matter and

deconfined quark matter, respectively. In the mixed phase, the baryon number density, n_b , and the energy density, ε , are given by

$$n_b = (1 - \chi)n_b^H + \chi n_b^q, \quad (23)$$

and

$$\varepsilon = (1 - \chi)\varepsilon^H + \chi\varepsilon^q, \quad (24)$$

where n_b^H (ε^H) and n_b^q (ε^q) denote the baryon number (energy) densities of the hadron and quark phase, respectively. The quantity $\chi \equiv V_q/V$ denotes the volume proportion of quark matter, V_q , in the unknown volume V . By definition, χ therefore varies between 0 and 1, depending on how much confined hadronic matter has been converted to quark matter. The condition of global electric charge neutrality is given by the equation

$$(1 - \chi) \sum_{i=B,l} q_i^H n_i^H + \chi \sum_{i=q,l} q_i^q n_i^q = 0. \quad (25)$$

where q_i is the electric charge of the i -th specie in units of the electron charge. In this work we have chosen global rather than local electric charge neutrality, since the latter is not fully consistent with the Einstein-Maxwell equations and the micro physical condition of chemical equilibrium and relativistic quantum statistics, as shown in [50]. In contrast to local electric charge neutrality, the global neutrality condition put a new positive electric charge on hadronic matter, rendering it more isospin symmetric, and a net negative electric charge on quark matter, allowing neutron star matter to settle down in a lower energy state that otherwise possible [46, 47].

V. MODELS FOR THE ULTRA-DENSE PART OF THE EOS OF NEUTRON STAR MATTER

Figures 1 and 2 show our models for the EoS of neutron star matter. The hadronic contributions are computed for the lagrangian given in Eq. (17). For the quark matter phase the local NJL model (Sec. II A) as well as the non-local NJL model (Sec. II B.) have been used.

For the hadronic phase we consider the parameter sets GM1 and NL3. For the strength of the vector repulsion we set the values $G_V^L/G_S = 0$ and $G_V^L/G_S = 0.3$ for the l3NJL model, and $G_V^{NL}/G_S = 0$ and $G_V^{NL}/G_S = 0.09$ for the n3NJL model. Figures 3 and 4 show the

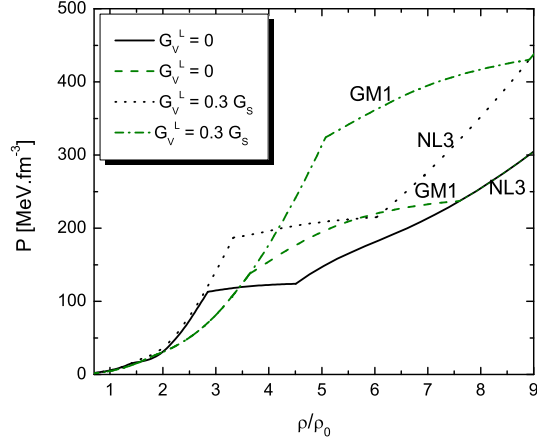


Figure 1. (Color online) Pressure as a function of baryon number density for the local NJL model, l3NJL. The hadronic parametrizations are GM1 and NL3, and the vector repulsion strengths are $G_V^L/G_S = 0$ and $G_V^L/G_S = 0.3$.

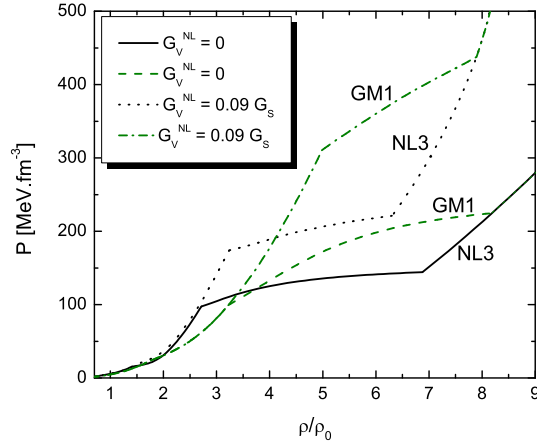


Figure 2. (Color online) Same as Fig. 1, but for the non-local NJL model, n3NJL, and vector repulsion strengths $G_V^{NL}/G_S = 0$ and $G_V^{NL}/G_S = 0.09$.

equations of state of Figs. 1 and 2 in the 3-space spanned by the neutron chemical potential, electron chemical potential, and pressure.

For pure hadronic matter, depending on the EoS, maximum masses for neutron stars of up to $\sim 2.77 M_\odot$ are obtained. This is shown in Fig. 5, which displays the mass-radius relationship of neutron stars for the parametrizations (Table III) studied in this work.

Fig. 6 shows the mass-radius relationship of neutron stars for three parametrizations

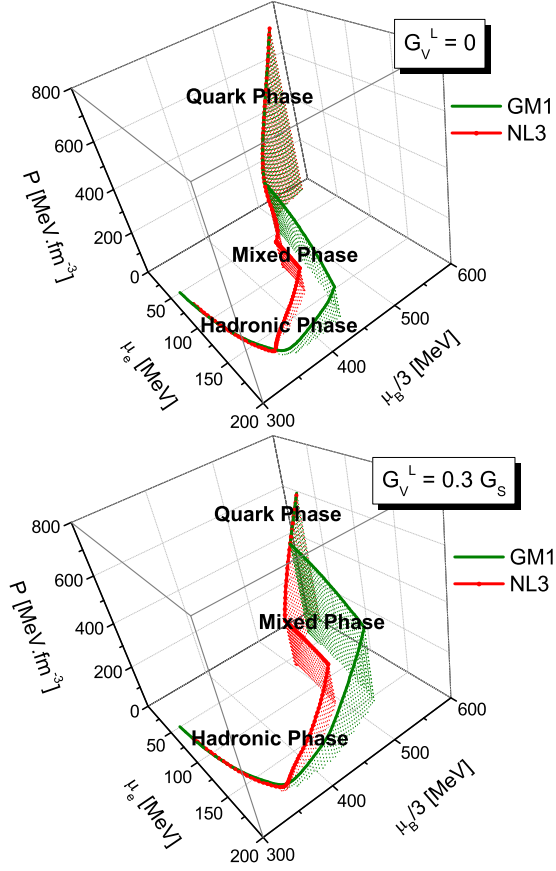


Figure 3. (Color online) Pressure as a function of neutron ($\mu^n = \mu_B/3$) and electron (μ^e) chemical potential for the local NJL model, l3NJL, hadronic parametrizations GM1 and NL3, and vector repulsions $G_V^L/G_S = 0$ and $G_V^L/G_S = 0.3$.

of the non-linear relativistic mean-field approximation including quark matter but without considering the effect of the vector interactions among quarks. The quark matter softens the EoS, which reduces the maximum mass of the neutron star. The masses of the recently discovered high-mass neutron stars J1614-2230 ($1.97 \pm 0.04 M_\odot$) [21] and J0348+0432 ($2.01 \pm 0.04 M_\odot$) [22] are shown for comparison. Depending on the (unknown) strength of the vector repulsion, both stars could contain quark-hybrid matter in their cores.

VI. QUARK-HYBRID COMPOSITION

In Figs. 7, 8, 9, 10 we show the relative particle fractions as a function of baryon number density for both the local and non-local NJL model.

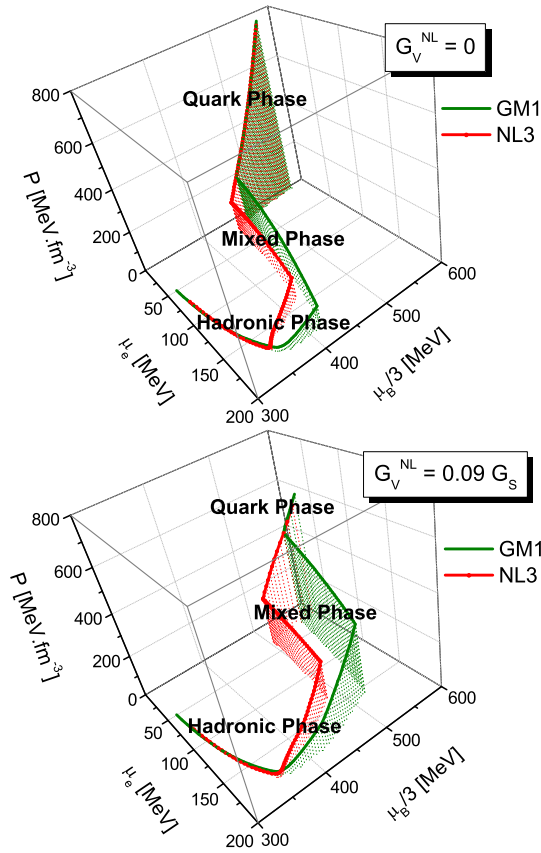


Figure 4. (Color online) Same as Fig. 3, but for the non-local NJL model, n3NJL, hadronic parametrizations GM1 and NL3, and vector repulsions $G_V^{NL}/G_S = 0$ and $G_V^{NL}/G_S = 0.09$.

Table V shows the densities at which quark-hybrid matter gets depletonized, which depends on the ratio G_V/G_S .

Since we are imposing the condition of global charge neutrality on the quark-hadron phase transition, the pressure varies monotonically with the proportion of the phases in equilibrium, as it is showed in Fig. 11. The hatched areas denote the mixed phase regions where confined hadronic matter and deconfined quark matter coexist. Our calculations show that, in the non-local case, similarly to the observation in Ref. [23], the inclusion of the quark vector coupling contribution shifts the onset of the phase transition to higher densities and narrows the width of the mixed quark-hadron phase, when compared to the case $G_V = 0$. On the contrary, when the quark matter phase is represented by the local l3NJL model, the width of the mixed phase tends to be larger for finite G_V values. This effect can be seen both in Fig. 11 as well as in Tables VI and VII.

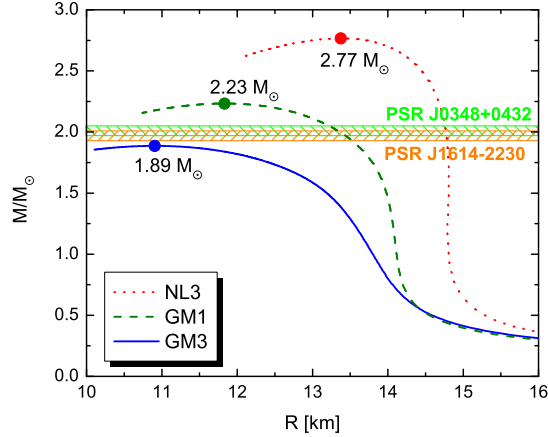


Figure 5. (Color online) Mass-radius relationships of neutron stars for the pure hadronic (no quark matter) EoS studied in this work.

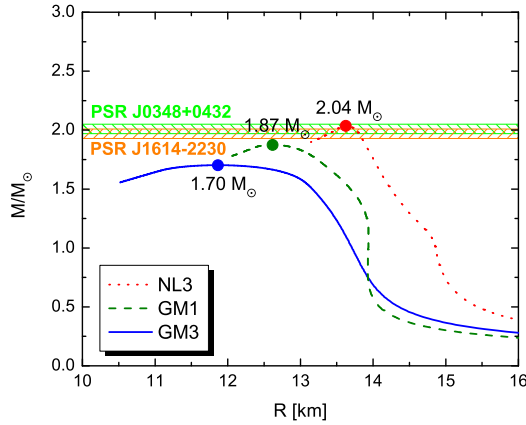


Figure 6. (Color online) Same as Fig. 5, but with quark matter included.

To account for the uncertainty in the theoretical predictions for the ratio G_V/G_S [51–54], we treat the vector coupling constant as a free parameter. The n3NJL model is more sensitive to the increase of G_V/G_S than the local model. For $G_V/G_S > 0.09$ we have a shift of the onset of the quark-hadron phase transition to higher and higher densities, preventing quark deconfinement in the cores of NSs. However, we can reach values of G_V/G_S up to 0.3 for the local NJL case. Next we determine the bulk properties of spherically symmetric neutron stars for the collection of equations of state for both the local and non-local NJL models. The properties are performed by solving the Tolmann-Oppenheimer-Volkoff (TOV)

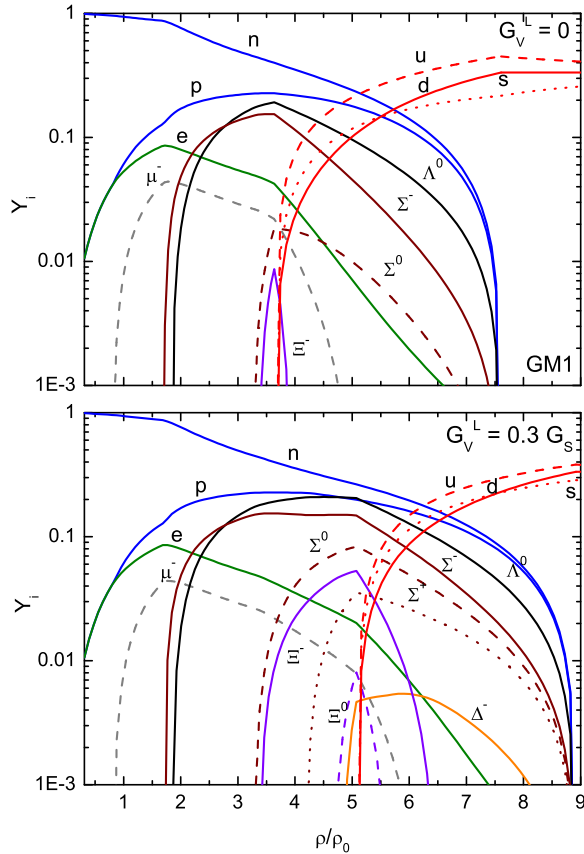


Figure 7. (Color online) Particle population of neutron star matter computed for the local NJL model. The hadronic model parametrization is GM1.

equations of general relativity theory [2, 55]. The outcomes are shown in Figs. 12 and 13, respectively. The green and orange bands in Fig. 12 show the masses of the recently discovered, massive neutron stars PSR J1614-2230 [21] and J0348+0432 [22].

As a star becomes increasingly massive, the stellar core compositions consist of either only nucleons, nucleons and hyperons, a mixed phase of quarks and hadrons (MP), or a pure quark matter phase (QP). Nevertheless, when the EoS is constructed with the non-local model considering the NL3 parametrization for the hadronic phase, we find that the latter only exist in NSs which lie to the left of the stars' mass peak. This can be seen in more detail in Figs. 12(b) and 13(b). Such stars are unstable against radial oscillations and thus cannot exist stably in the universe. In contrast, all NSs on the MP branches up to the mass peaks are stable against radial perturbations. The heavier stars of the stellar sequences contain mixed phases of quarks and hadrons in their centers. But the densities inside these

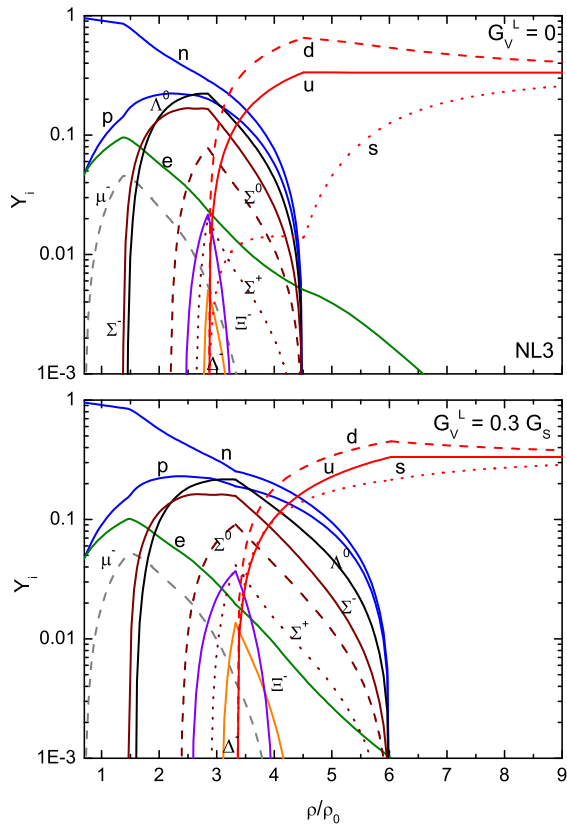


Figure 8. (Color online) Same as Fig. 7, but for the hadronic model parametrization NL3.

stars are not high enough to generate pure quark matter in their cores. Such matter can only be formed in neutron stars which are located on the gravitationally unstable branch of the mass-radius relationship. However, when the EoS is constructed for the local NJL model, we obtain stars with pure quark matter in their cores. The results for the masses and radii of the maximum-mass models are shown in Table VIII.

VII. SUMMARY AND CONCLUSIONS

In this work we use extensions of the local (l3NJL) and non-local (n3NJL) Nambu-Jona Lasinio model to analyze quark deconfinement in the cores of neutron stars. We constructed the phase transition from hadronic matter to quark matter via the Gibbs conditions, imposing global electric charge neutrality. We have calculated the mass-radius relationship of such stars. Depending on the strength of quark vector repulsion, we find that mixed phases of confined hadronic matter and deconfined quarks can exist in neutron stars as massive as

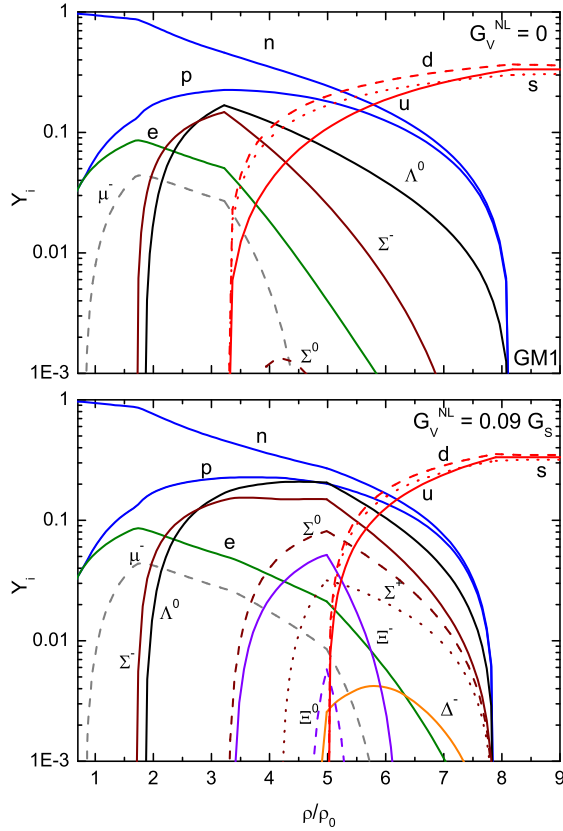


Figure 9. (Color online) Particle population of neutron star matter computed for n3NJL. The hadronic model parametrization is GM1.

around $2.1 - 2.4 M_{\odot}$. The radii of these objects are between 12 and 13 km, as expected for neutron stars.

According to this study, for the n3NJL model, a transition to pure quark matter occurs only in neutron stars which lie on the gravitationally unstable branch of the stellar sequences if the parametrization NL3 for the hadronic matter is used. This is different for the l3NJL model, for which we obtain stable neutron stars with pure quark matter cores in their centers. With increasing stellar mass, the stellar core compositions consist either of nucleons only, nucleons and hyperons, a mixed phase of quarks and hadrons (MP), or a pure quark matter phase (QP). The latter, however, exists only in neutron stars which lie beyond their mass peaks when the EoS is constructed with the n3NJL model. Such stars are unstable against radial oscillations and thus cannot exist stably in the universe. In contrast to this, all neutron stars on the MP branches up to the mass peaks are stable.

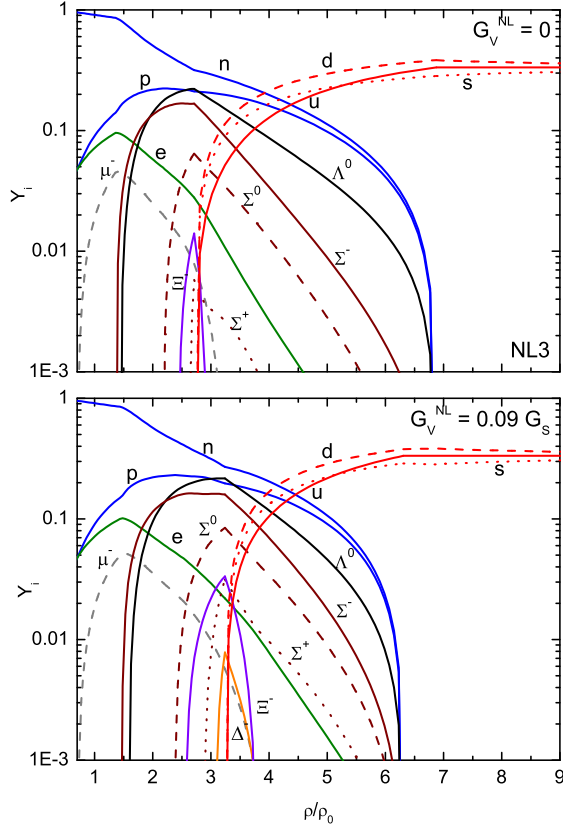


Figure 10. (Color online) Same as Fig. 9, but for the hadronic model parametrization NL3.

ACKNOWLEDGMENTS

M. Orsaria thanks N. N. Scoccola for helpful discussions and T. Hell for answer some questions about his previous works on nonlocal models. M. Orsaria and G. Contrera thank CONICET for financial support. H. Rodrigues thanks CAPES for financial support under contract number BEX 6379/10-9. F. Weber is supported by the National Science Foundation (USA) under Grant PHY-0854699.

-
- [1] N. K. Glendenning, *Compact Stars, Nuclear Physics, Particle Physics, and General Relativity*, 2nd ed. (Springer-Verlag, New York, 2000).
 - [2] F. Weber, *Pulsars as Astrophysical Laboratories for Nuclear and Particle Physics*, High Energy Physics, Cosmology and Gravitation Series (IOP Publishing, Bristol, Great Britain, 1999).

- [3] *Physics of Neutron Star Interiors*, ed. by D. Blaschke, N. K. Glendenning, and A. Sedrakian, Lecture Notes in Physics **578** (Spring-Verlag, Berlin, 2001).
- [4] J. M. Lattimer and M. Prakash, *Astrophys. J.* **550**, 426 (2001).
- [5] F. Weber, *Prog. Part. Nucl. Phys.* **54**, 193 (2005).
- [6] D. Page and S. Reddy, *Ann. Rev. Nucl. Part. Sci.* **56**, 327 (2006).
- [7] P. Haensel, A. Y. Potekhin, and D. G. Yakovlev, *Neutron Stars 1*, Astrophysics and Space Science Library, (Springer-Verlag, New York, 2006).
- [8] T. Klähn *et al.*, *Phys. Rev. C* **74**, 035802 (2006).
- [9] A. Sedrakian, *Prog. Part. Nucl. Phys.* **58**, 168 (2007).
- [10] T. Klähn *et al.*, *Phys. Lett. B* **654**, 170 (2007).
- [11] M. G. Alford, A. Schmitt, K. Rajagopal, and T. Schäfer, *Rev. Mod. Phys.* **80**, 1455 (2008).
- [12] *Strongly Interacting Matter - The CBM Physics Book*, Lecture Notes in Physics **814** (Springer, 2011).
- [13] A. Bodmer, *Phys. Rev. D* **4**, 1601 (1971).
- [14] E. Witten, *Phys. Rev. D* **30**, 272 (1984).
- [15] H. Terazawa, INS-Report-338 (INS, Univ. of Tokyo, 1979); *J. Phys. Soc. Japan* **58**, 3555 (1989); *ibid.*, **58**, 4388 (1989); *ibid.*, **59**, 1199 (1990).
- [16] E. Farhi and R. L. Jaffe, *Phys. Rev. D* **30**, 2379 (1984).
- [17] C. Alcock, E. Farhi, and A. V. Olinto, *Astrophys. J.* **310**, 261 (1986); C. Alcock and A. V. Olinto *Ann. Rev. Nucl. Part. Sci.* **38**, 161 (1988).
- [18] Proc. Int. Workshop on *Strange Quark Matter in Physics and Astrophysics*, J. Madsen and P. Haensel (eds.), *Nucl. Phys. B (Proc. Suppl.)* **24B** (1991).
- [19] J. Madsen, *Lecture Notes in Physics* **516**, 162 (1999).
- [20] I. Sagert, M. Wietoska, and J. Schaffner-Bielich *J. Phys. G.* **32**, S241 (2006).
- [21] P. B. Demorest, T. Pennucci, S. M. Ranson, M. S. E. Roberts, and J. W. T. Hessels, *Nature* **467**, 1081 (2010).
- [22] J. Antoniadis *et al.*, *Science* **340**, no. 6131 (2013).
- [23] M. Orsaria, H. Rodrigues, F. Weber and G. A. Contrera, *Phys. Rev. D* **87**, 023001 (2013).
- [24] Y. Nambu and G. Jona-Lasinio, *Phys. Rev.* **122**, 345 (1961); *Phys. Rev.* **124**, 246 (1961).
- [25] G. Ripka, *Quarks bound by chiral fields* (Oxford University Press, Oxford, 1997).
- [26] S. M. Schmidt, D. Blaschke, and Yu. L. Kalinovsky, *Phys. Rev. C* **50**, 435 (1994).

- [27] R. D. Bowler and M. C. Birse, Nucl. Phys. A **582**, 655 (1995); R. S. Plant and M. C. Birse, Nucl. Phys. A **628**, 607 (1998).
- [28] M. B. Parappilly, P. O. Bowman, U. M. Heller, D. B. Leinweber, A. G. Williams, and J. B. Zhang, Phys. Rev. D **73**, 054504 (2006).
- [29] D. Blaschke, Yu L. Kalinovsky, G. Röpke, S. Schmidt, and M. K. Volkov, Phys. Rev. C **53**, 2394 (1996).
- [30] G. A. Contrera, D. Gomez Dumm, and N. N. Scoccola, Phys. Lett. B **661** 113 (2008).
- [31] G. A. Contrera, D. Gomez Dumm, and N. N. Scoccola, Phys. Rev. D **81**, 054005 (2010).
- [32] T. Hatsuda and T. Kunihiro, Phys. Rep. **247**, 22 (1994).
- [33] P. Rehberg, S. P. Klevansky, and J. Hufner, Phys. Rev. C **53**, 410 (1996).
- [34] C. Ratti, Europhys. Lett. **61**, 314 (2003).
- [35] G. Y. Shao, Phys. Lett. B **704**, 343 (2011).
- [36] M. Kitazawa, T. Koide, T. Kunihiro, and Y. Nemoto, Prog. Theor. Phys. **108**, 929 (2002).
- [37] K. Masuda, T. Hatsuda, and T. Takatsuka, Astrophys. J. **764**, 12 (2013).
- [38] C. H. Lenzi and G. Lugones, Astrophys. J. **759**, 57 (2012).
- [39] N. K. Glendenning and S. A. Moszkowski, Phys. Rev. Lett. **67**, 2414 (1991).
- [40] G. A. Lalazissis, J. Konig, and P. Ring, Phys. Rev. C **55**, 540 (1997).
- [41] D. B. Blaschke, D. Gomez Dumm, A. G. Grunfeld, T. Klähn, and N. N. Scoccola, Phys. Rev. C **75**, 065804 (2007).
- [42] K. Kashiwa, T. Hell, and W. Weise, Phys. Rev. D **84**, 056010 (2011).
- [43] A. Scarpettini, D. Gomez Dumm, and N. N. Scoccola, Phys. Rev. D **69**, 114018 (2004).
- [44] J. D. Walecka, Ann. Phys. **83**, 497 (1974).
- [45] B. D. Serot and J. D. Walecka, Adv. Nucl. Phys. **16**, 1 (1986).
- [46] N. K. Glendenning, Phys. Rev. D **46**, 1274 (1992).
- [47] N. K. Glendenning, Phys. Rep. **342**, 393 (2001).
- [48] S. Reddy, G. Bertsch, and M. Prakash, Phys. Lett. B **475**, 1 (2000).
- [49] X. Na, R. Xu, F. Weber and R. Negreiros, Phys. Rev. D **86**, 123016 (2012).
- [50] M. Rotondo, Jorge A. Rueda, R. Ruffini, and S.-S. Xue, Phys. Lett. B **701** 667 (2011).
- [51] C. Sasaki, B. Friman, and K. Redlich, Phys. Rev. D **75**, 054026 (2007).
- [52] K. Fukushima, Phys. Rev. D **77**, 114028 (2008) [Erratum-ibid. Phys. Rev. D **78**, 039902 (2008)].

- [53] N. M. Bratovic, T. Hatsuda, and W. Weise, Phys. Lett. B **719**, 131 (2013).
- [54] G. A. Contrera, A. G. Grunfeld, and D. B. Blaschke, arXiv:1207.4890 [hep-ph].
- [55] R. C. Tolman, Phys. Rev. **55**, 364 (1939); J. R. Oppenheimer and G. M. Volkoff, Phys. Rev. **55**, 374 (1939).

Table I. Comparison of the key features of the local NJL and non-local NJL model.

Local NJL	Non-local NJL model
<ul style="list-style-type: none"> • Lack of confinement 	<ul style="list-style-type: none"> • Confinement with a proper choice of the non-local regulator and model parameters [27]
<ul style="list-style-type: none"> • Quark-quark scalar-isoscalar and pseudoscalar-isovectorial local interaction 	<ul style="list-style-type: none"> • Quark-quark interaction through phenomenologically effective quark propagator
<ul style="list-style-type: none"> • Non-renormalizable. Ultra-violet (UV) cutoff Λ is needed 	<ul style="list-style-type: none"> • UV divergences are fixed [25]. Model dependent form factor $g(p)$
<ul style="list-style-type: none"> • Dynamical quark masses are momentum independent 	<ul style="list-style-type: none"> • Dynamical quark masses are momentum dependent as also found in lattice calculations of QCD [28]
<ul style="list-style-type: none"> • Divergences in the meson loop integrals. Extra cutoffs are needed. 	<ul style="list-style-type: none"> • The momentum dependent regulator makes the theory finite to all orders in the $1/N_c$ expansion [29]
<ul style="list-style-type: none"> • The cutoff Λ is turned off at high momenta, limiting the applicability of the model at high densities 	<ul style="list-style-type: none"> • The form factor provides a natural cutoff that falls off at high momenta

Table II. Parameter set used for the non-local NJL model (n3NJL) calculations presented in this paper.

Parameters	n3NJL
$m_u = m_d$	6.2 MeV
m_s	140.7 MeV
Λ	706.0 MeV
$G_S\Lambda^2$	15.04
$H\Lambda^5$	-337.71

Table III. Parameter values for the hadronic lagrangian of Eq. (19).

Coupling	Parametrizations		
	GM1	GM3	NL3
constants			
g_σ	8.910	8.175	10.217
g_ω	10.610	8.172	12.868
g_ρ	8.196	8.259	8.948
b_σ	0.002947	0.008659	0.002055
c_σ	-0.001070	-0.002421	-0.002651

Table IV. Properties of symmetric nuclear matter at saturation density for the parameters listed in Table III. Shown are the saturation density ρ_0 , energy per baryon E/N , nuclear incompressibility K , effective nucleon mass m_N^* , and asymmetry energy a_{sy} .

Properties	Parametrizations		
	GM1	GM3	NL3
ρ_0 (fm $^{-3}$)	0.153	0.153	0.148
E/N (MeV)	-16.3	-16.3	-16.3
K (MeV)	300	240	272
m^*/m_N	0.78	0.70	0.60
a_{sy} (MeV)	32.5	32.5	37.4

Table V. Deleptonization densities for the local and non-local NJL model, for the GM1 parametrization.

G_V/G_S	Deleptonization density (ρ_0)	
	Local NJL	Non-local NJL
0 (GM1)	6.65	5.85
0 (NL3)	6.61	4.55
0.09 (GM1)	—	7.03
0.09 (NL3)	—	5.33
0.30 (GM1)	7.45	—
0.30 (NL3)	6.26	—

Table VI. Width of the mixed phase, central densities of the associated maximum-mass star of the QHS in the Local NJL model case.

G_V^L	Mixed phase	Central density of M_{\max}
	(ρ_0)	(ρ_0)
0 (GM1)	3.64 – 7.60	7.21
0.30 (GM1)	5.07 – 8.92	8.81
0 (NL3)	2.85 – 4.51	5.96
0.30 (NL3)	3.33 – 6.03	6.52

Table VII. Widths of the mixed phases and central densities of the associated maximum-mass stars for the non-local NJL model.

G_V^{NL}	Mixed phase	Central density of M_{\max}
	(ρ_0)	(ρ_0)
0 (GM1)	3.22 – 8.18	6.87
0.09 (GM1)	4.83 – 7.89	8.69
0 (NL3)	2.71 – 6.87	5.68
0.09 (NL3)	3.24 – 6.31	6.28

Table VIII. Maximum masses and radii of neutron stars made of quark-hybrid matter for different vector repulsion (G_V/G_S).

NJL Model	G_V/G_S	R_{\max} (km)	M_{\max}/M_{\odot}
Local	0	12.49	1.96
GM1	0.30	11.80	2.11
Local	0	13.68	2.07
NL3	0.30	13.53	2.37
Non-local	0	12.62	1.87
GM1	0.09	11.81	2.11
Non-local	0	13.62	2.04
NL3	0.09	13.56	2.35

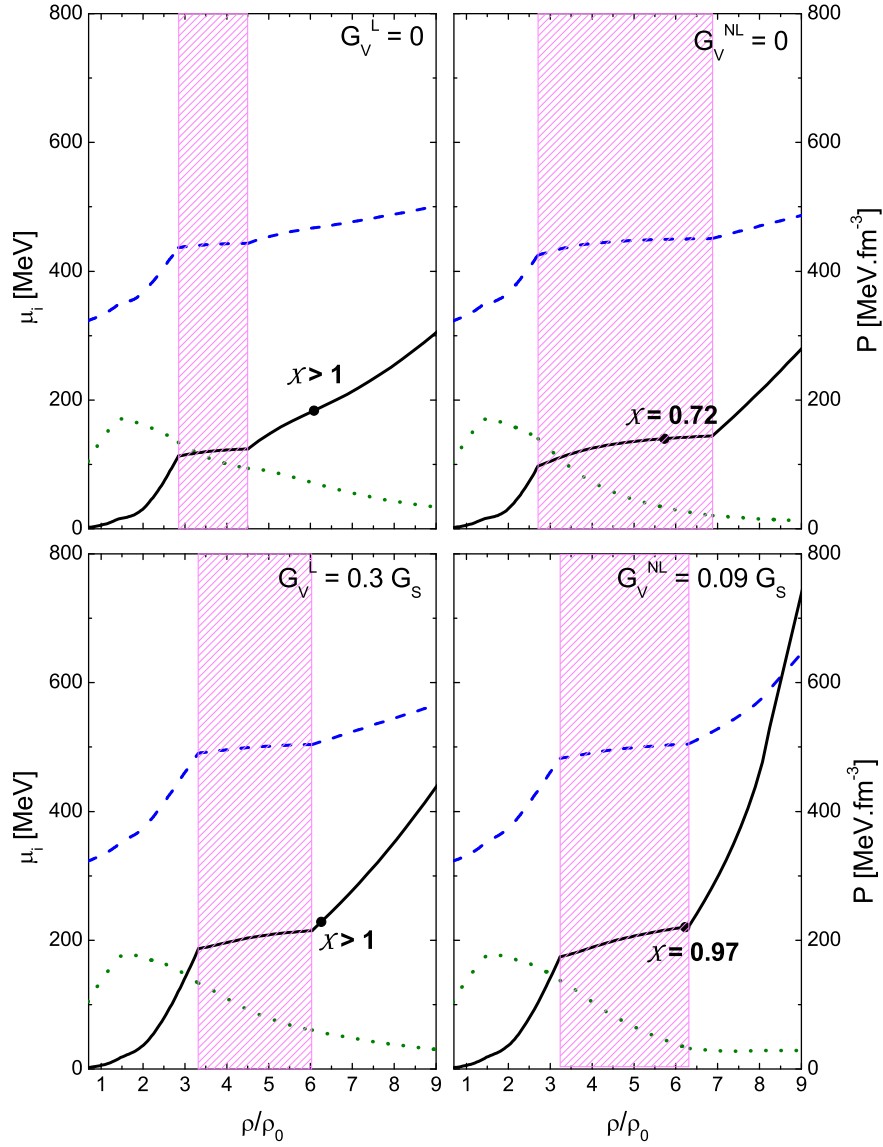
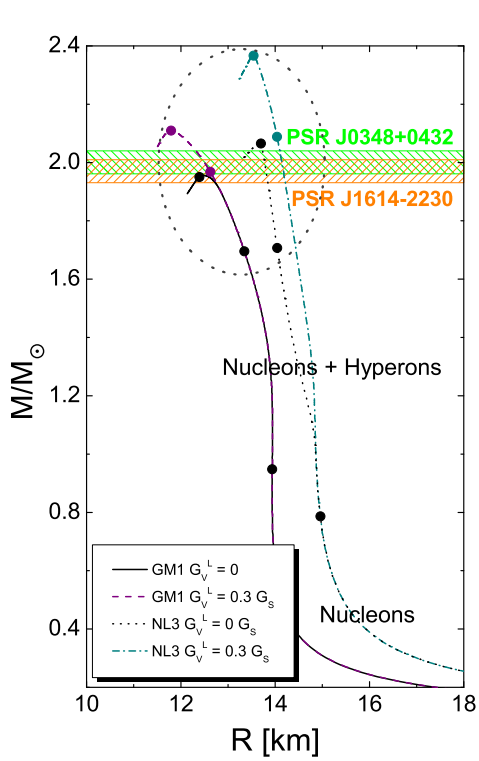
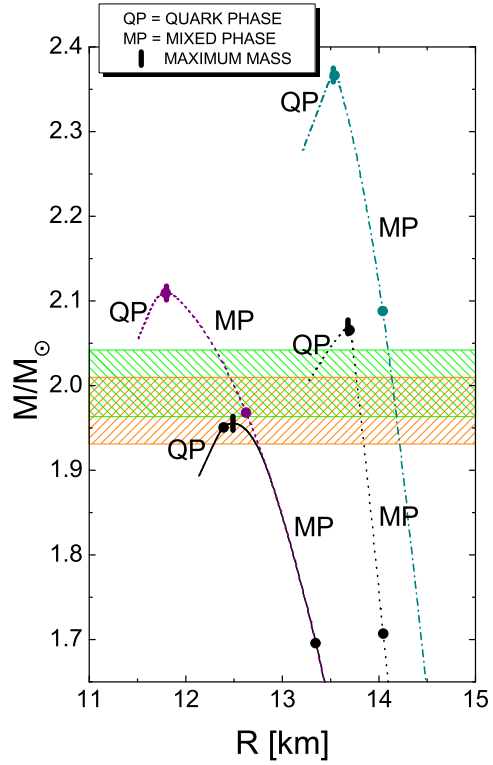


Figure 11. (Color online) Pressure P (solid lines), baryonic chemical potential μ_B (dashed lines), and electron chemical potential μ_e (dotted lines) as a function of baryon number density in units of normal nuclear matter density, $\rho_0 = 0.16 \text{ fm}^{-3}$ for parametrization NL3. The hatched areas denote the mixed phase regions where confined hadronic matter and deconfined quark matter coexist in chemical equilibrium with each other. The results are computed for the vector coupling constant $G_V = 0$ and the highest values of the ratio G_V/G_S considered for both the local (left panels) and non-local (right panels) NJL model.

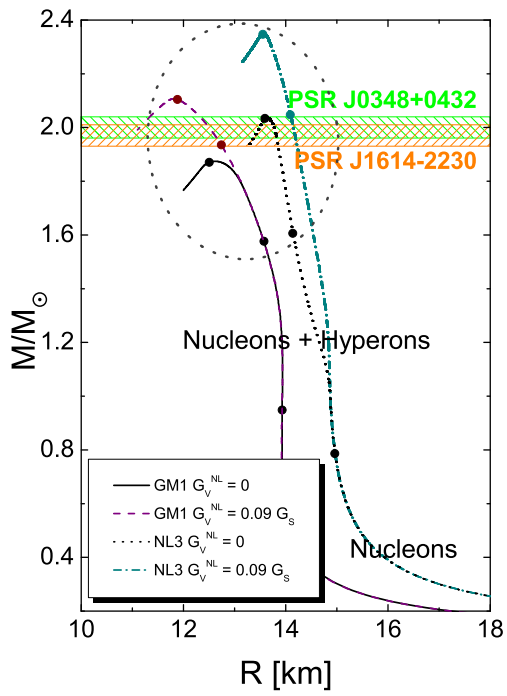


(a) Mass-radius relationships of neutron stars made of quark-hybrid matter.

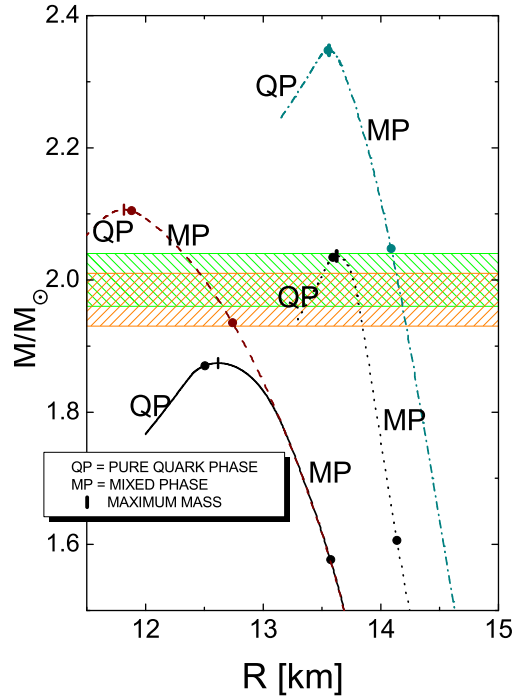


(b) Enlargement of the circled region of Fig. 12(a).

Figure 12. (Color online) Quark-hybrid matter inside of neutron stars, computed for the standard NJL model and hadronic model parametrization GM1 and NL3. The symbols 'MP' and 'QP' stand for mixed phase and pure quark phase, respectively. The vertical bars denote the maximum-mass star of each stellar sequence.



(a) Mass-radius relationships of neutron stars made of quark-hybrid matter.



(b) Enlargement of the circled region of Fig. 13(a).

Figure 13. (Color online) Same as Fig. 12, but computed for the non-local NJL model and hadronic model parametrizations GM1 and NL3.