Quark-Hybrid Matter in the Cores of Massive Neutron Stars

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Using a nonlocal extension of the SU(3) Nambu-Jona Lasinio model, which reproduces several of the key features of Quantum Chromodynamics, we show that mixed phases of deconfined quarks and confined hadrons (quark-hybrid matter) may exist in the cores of neutron stars as massive as around $2.1\,M_{\odot}$. The radii of these objects are found to be in the canonical range of $\sim 12-13$ km. According to our study, the transition to pure quark matter does not occur in stable neutron stars, but is shifted to neutron stars which are unstable against radial oscillations. The implications of our study for the recently discovered, massive neutron star PSR J1614–2230, whose gravitational mass is $1.97 \pm 0.04\,M_{\odot}$, are that this neutron star may contain an extended region of quark-hybrid matter at it center, but no pure quark matter.

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Introduction – Quantum Chromodynamics (QCD) has the properties of asymptotic freedom and confinement. The former implies that in the high-momentum transfer regime, the quarks behave essentially as free particles, i.e., the interaction between two quarks due to gluon exchange is very weak. This regime can therefore be treated using perturbation theory, where the quark-gluon coupling constant serves as an expansion parameter. For this momentum range, the dispersion processes can be calculated very accurately. By contrast, at low-momentum transfers ($\lesssim 1$ GeV) QCD becomes highly nonlinear, which prevents the use of perturbative methods. One of the renowned effective models that serves as a suitable approximation to QCD in the low-energy regime is the quark version of the Nambu Jona-Lasinio (NJL) model [1–3]. In this model chiral symmetry constraints are taken into account via effective interactions between quarks, through local four-point vertex interactions. The drawbacks of using local interactions are that the model must be regularized to avoid divergences in the loop integrals, and that the model is non-confining. The absence of confinement is essentially related to the fact that the dynamically generated constituent quark masses are momentum independent. Since the 1990's, there have been investigations proposing nonlocal interactions to solve these problems [4]. One interesting suggestion arises from the relationship between the NJL model and the model of one-gluon exchange where an effective gluon propagator is used to generate effective interactions between quarks. This provides a natural way to introduce a nonlocality in the quark-quark interaction, which can be characterized by a model-dependent form factor, $g(p)$ [5].

In this paper we analyze the global structure and composition of massive neutron stars in the framework of an extended version of the nonlocal SU(3) NJL model. Of particular interest is the question as to whether or not massive neutron stars, such as the recently discovered pulsar PSR J1614–2230 whose gravitational mass was found to be $1.97 \pm 0.04\, M_\odot$ [6], may contain stellar cores made of deconfined quark matter [7–9, 16]. Tendentially, one could argue that quark deconfinement may not occur in the cores of such high-mass neutron stars, since the underlying nuclear equation of state must be extremely stiff in order to support such high mass neutron stars. As will be shown in this paper, arguments along these lines appear premature.

Our study is based on the latest set of model parameters of the nonlocal SU(3) NJL model. The parameter fit is performed for a phenomenological value of the strange quark mass of $m_s = 140.7$ MeV. Details of the NJL formalism at zero chemical potential and finite temperature can be found in [10–12].

Over the last two decades, several authors have started to take into account the effect of the quark vector interaction in effective chiral models like NJL[8, 9, 13–17]. It is known that the repulsive character of the vector coupling in these models affects the quark-hadron phase transition and moves the chiral restoration to a larger value of the quark chemical potential [15]. Thus, if the quark deconfined transition in the cores of neutron stars is modeled by a NJL-like model, it is expected that the vector coupling contribution modifies the nuclear equation of state (EoS) and hence the mass-radius relationship of neutron stars. Most of the NJL studies of neutron stars are treating the interactions among quarks in terms of local fermion-fermion couplings and/or impose the condition of local electric charge neutrality on the stellar matter [8, 9, 16]. In this paper, we are using a generalized version of the NJL model where the interactions are nonlocal and momentum-dependent, and the condition of local charge neutrality is replaced with the more relaxed condition of global charge conservation, as neutron stars are characterized by two rather than one conserved charge [18]. As a consequence, the pressure in the mixed quark-hadron phase varies with density and is, therefore, not a priori excluded from neutron stars.

To include the vector interaction in the NJL model we follow [19]. However, we shall consider three different vector fields, one for each quark flavor, instead of a single vector field for all quarks.

Description of quark matter phase in the framework of the nonlocal SU(3) NJL model – We start from the Euclidean effective action associated with the nonlocal SU(3) quark model,

$$
\begin{aligned}
S_E = \int d^4x \, \{ & \bar{\psi}(x) \left[-i\gamma_\mu \partial_\mu + \hat{m} \right] \psi(x) \\
& - \frac{G_s}{2} \left[j_a^S(x) \; j_a^S(x) + j_a^P(x) \; j_a^P(x) \right] \\
& - \frac{H}{4} \, T_{abc} \left[j_a^S(x) j_b^S(x) j_c^S(x) - 3 \; j_a^S(x) j_b^P(x) j_c^P(x) \right] \\
& - \frac{G_V}{2} j_{V,f}^\mu(x) j_{V,f}^\mu(x), \qquad (1)
\end{aligned}
$$

where ψ is a chiral $U(3)$ vector that includes the light quark fields, $\psi \equiv (u \; d \; s)^T$, and $\hat{m} = \mathrm{diag}(m_u, m_d, m_s)$ stands for the current quark mass matrix. For simplicity we consider the isospin symmetry limit, in which $m_u = m_d = \bar{m}$. The fermion kinetic term includes the convariant derivative $D_\mu \equiv \partial_\mu - iA_\mu$, where A_μ are color gauge fields, and the operator $\gamma_\mu \partial_\mu$ 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 $j_{V,f}^\mu(x)$ are given by

$$j_a^S(x) = \int d^4z \; \widetilde{g}(z) \; \bar{\psi}\left(x + \frac{z}{2}\right) \; \lambda_a \; \psi\left(x - \frac{z}{2}\right) \; ,$$
$$j_a^P(x) = \int d^4z \; \widetilde{g}(z) \; \bar{\psi}\left(x + \frac{z}{2}\right) \; i\,\gamma_5\lambda_a \; \psi\left(x - \frac{z}{2}\right) \; ,$$
$$j_{V,f}^{\mu}(x) = \int d^4z \; \widetilde{g}(z) \; \bar{\psi}_f\left(x + \frac{z}{2}\right) \; \gamma^{\mu} \; \psi_f\left(x - \frac{z}{2}\right), \tag{2}$$

where $\widetilde{g}(z)$ is a form factor responsible for the non-local character of the interaction, and the matrices λ_a, with $a = 0, .., 8$, are the usual eight Gell-Mann 3×3 matrices – generators of SU(3) – plus $\lambda_0 = \sqrt{2/3}\,\mathbb{I}_{3\times 3}$. Finally, the constants T_{abc} in the t'Hooft term accounting for flavor-mixing are defined by

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

After standard bosonization of Eq. (1), the integrals over the quark fields can be performed in the framework of the Euclidean four-momentum formalism. The grand canonical potential in the mean-field approximation at zero temperature, including the vector coupling, is then given by

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

where $N_c = 3$, $E_f = \sqrt{\vec{p}^{\,2} + m_f^2}$ and we have defined

$$\omega_f^2 = (\,p_0 \, + \, i\,\mu_f\,)^2 \, + \, \vec{p}^{\,2}. \tag{5}$$

The masses of free quarks are denoted by m_f, where $f = u, d, s$. The $\varpi_{V,f}$ mean field is related to the vector current density $j_{V,f}^{\mu}$ of Eq. (2).

The momentum-dependent constituent quark masses M_f depend explicitly on the quark mean fields $\bar{\sigma}_f$,

$$M_f(\omega_f^2) \;\; = \;\; m_f \, + \, \bar{\sigma}_f \, g(\omega_f^2), \tag{6}$$

where $g(\omega^2)$ denotes the Fourier transform of the form factor $\widetilde{g}(z)$.

Following Ref. [19], the quark vector interaction shifts the quark chemical potential according to

$$\widetilde{\mu}_f = \mu_f - \varpi_{V,f}. \tag{7}$$

Note that the shifting of the quark chemical potential does not affect the nonlocal form factor $g(\omega_f^2)$ as discussed in [19]. The mean-field values of the auxiliary fields $\bar{S}_f$ are given by [10]

$$\bar{S}_f = -\,16\,N_c \int_0^{\infty} dp_0 \int_0^{\infty} \frac{dp}{(2\pi)^3} \, g(\omega_f^2) \, \frac{M_f(\omega_f^2)}{\omega_f^2 + M_f^2(\omega_f^2)} \; . \tag{8}$$

In this paper we adopt a Gaussian form for the nonlocal form factor g,

$$g(\omega^2) = \exp\left(-\omega^2/\Lambda^2\right), \tag{9}$$

where Λ plays a role for the stiffness of the chiral transition. This parameter, together with the current quark mass $\bar{m}$ of up and down quarks and the coupling constants G_s and H in Eq. (4), have been fitted to the pion decay constant, f_π, and meson masses m_π, m_η, and $m_{\eta'}$, as described in [11, 12]. The result of this fit is $\bar{m} = 6.2$ MeV, $\Lambda = 706.0$ MeV, $G_s\Lambda^2 = 15.04$, $H\Lambda^5 = -337.71$. The strange quark current mass is treated as a free parameter and was set to $m_s = 140.7$ MeV. The strength of the vector interaction G_V is usually expressed in terms of the strong coupling constant G_s. To account for the uncertainty in the theoretical predictions for the ratio G_V/G_s, we treat the vector coupling constant as a free parameter [20–22], which varies from 0 to $0.1\,G_s$.

Using these parametrizations, the fields $\bar{\sigma}_f$ and $\varpi_{V,f}$ can be determined by minimizing Eq. (4),

$$\frac{\partial\Omega^{NL}}{\partial\bar{\sigma}_f} = \frac{\partial\Omega^{NL}}{\partial\varpi_{V,f}} = 0. \tag{10}$$

Description of confined hadronic matter – The hadronic phase is described in the framework of the non-linear relativistic mean field theory [23–26], where baryons (neutrons, protons, hyperons) interact via the exchange of scalar, vector and isovector mesons (σ, ω, $\vec{\rho}$, respectively). The Lagrangian of the theory is given by

$$\begin{aligned}\mathcal{L} = &\sum_{B=n,p,\Lambda,\Sigma,\Xi} \bar{\psi}_B \big[\gamma_\mu(i\partial^\mu - g_\omega\omega^\mu - g_\rho\vec{\rho}^\mu) \\ &- (m_N - g_\sigma\sigma)\big]\psi_B + \frac{1}{2}(\partial_\mu\sigma\partial^\mu\sigma - m_\sigma^2\sigma^2) \\ &- \frac{1}{3}b_\sigma m_N(g_\sigma\sigma)^3 - \frac{1}{4}c_\sigma(g_\sigma\sigma)^4 - \frac{1}{4}\omega_{\mu\nu}\omega^{\mu\nu} \\ &+ \frac{1}{2}m_\omega^2\,\omega_\mu\omega^\mu + \frac{1}{2}m_\rho^2\,\vec{\rho}_\mu\cdot\vec{\rho}^\mu \\ &- \frac{1}{4}\vec{\rho}_{\mu\nu}\vec{\rho}^{\mu\nu} + \sum_{\lambda=e^-,\mu^-}\bar{\psi}_\lambda(i\gamma_\mu\partial^\mu - m_\lambda)\psi_\lambda\,,\end{aligned} \tag{11}$$

where B sums all baryon states which become populated in neutron star matter [25, 26]. The quantities g_ρ, g_σ, and g_ω are the meson-baryon coupling constants. Non-linear σ-meson self-interactions are taken into account in Eq. (11) via the terms proportional to b_σ and c_σ [25, 26]. We have solved the equations of motion for the baryon and meson field equations, which follow from Eq. (11), for the relativistic mean-field approximation [25, 26]. For this approximation the meson fields σ, ω, ρ are approximated by their respective mean-field values $\bar{\sigma} \equiv \langle\sigma\rangle$, $\bar{\omega} \equiv \langle\omega\rangle$, and $\bar{\rho} \equiv \langle\rho_{03}\rangle$ [25, 26]. The parameters of the model, labeled GM1, are adjusted to the properties of nuclear matter at saturation density. They are taken from [27].

The condition of weak equilibrium requires the presence of electrons and muons, which are treated as free relativistic quantum gases, as described by the last term on the right-hand-side of Eq. (11). Neutron star matter is characterized by the conservation of electric and baryon number. This feature leads to the chemical equilibrium condition

$$\mu_i = B_i\,\mu_n - Q_i\,\mu_e\,, \tag{12}$$

where μ_n and μ_e denote the chemical potentials of neutrons and electrons, respectively. The quantities B_i and Q_i stand for the baryon numbers and the electric charges of the mesons and baryons of Eq. (11). Equation (12) greatly simplifies the mathematical analysis, since only the knowledge of two independent chemical potentials, μ_n and μ_e, is necessary. The latter are obtained from [25, 26]

$$\begin{aligned}\mu_B &= g_\omega\bar{\omega} + g_\rho\bar{\rho_{03}}I_B^3 + \sqrt{k_B^2 + m_B^{*2}}\,, \\ \mu_\lambda &= \sqrt{k_\lambda^2 + m_\lambda^2}\,,\end{aligned} \tag{13}$$

where $m_B^* = m_B - g_\sigma\bar{\sigma}$ denote the effective medium-modified baryon masses, k_B and k_λ are the Fermi momenta of baryons and leptons, respectively, and I_B^3 is the third component of the isospin vector of a baryon of type B. Finally, aside from chemical equilibrium, the condition of electric charge neutrality for confined hadronic matter is also of critical importance for the composition of neutron star matter. This condition is given by [25, 26]

$$\sum_B Q_i\,(2J_B+1)\,\frac{k_B^3}{6\pi^2} - \sum_\lambda \frac{k_\lambda^3}{3\pi^2} = 0\,, \tag{14}$$

where J_B denotes the spin of baryon B.

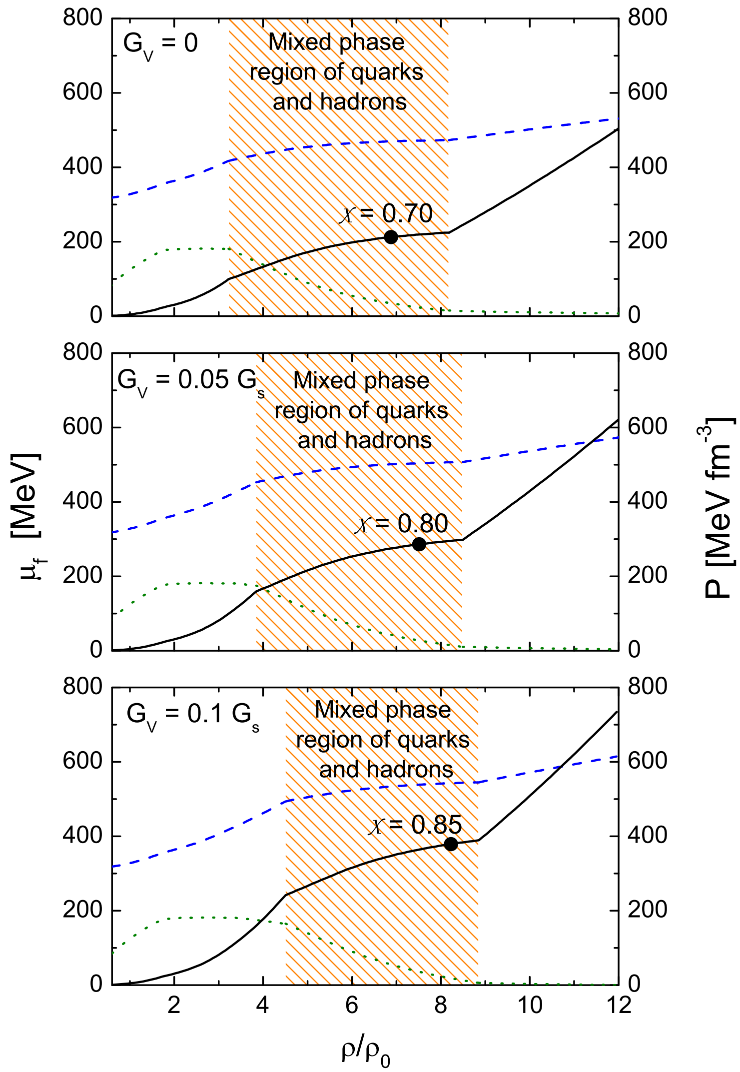


FIG. 1. (Color online) Pressure, P (solid lines), baryon chemical potential, μ_b (dashed lines), and electron chemical potential, μ_e (dotted lines) as a function of baryon number density, ρ, in units of the normal nuclear matter density, $\rho_0 = 0.16\ \mathrm{fm}^{-3}$. The hatched areas denote the mixed phase regions where confined hadronic matter and deconfined quark matter coexist. The solid dots indicate the central densities of the associated maximum-mass stars, shown in Fig. 3, and χ is the respective fraction of quark matter inside of them. The results are computed for three different values of the vector coupling constant, ranging from 0, to 0.05 G_s, to 0.1 G_s.

For the quark phase, the chemical potentials associated with quarks and electrons follow from Eq. (12) as $\mu_u = \mu_b - 2\mu_e/3$ and $\mu_d = \mu_s = \mu_b + \mu_e/3$, where $\mu_b = \mu_n/3$ stands for the baryon chemical potential, expressed in terms of the chemical potential of neutrons.

Description of the mixed phase of quarks and hadrons – To determine the mixed phase region of quarks and hadrons, we start from the Gibbs condition for phase equilibrium between hadronic (H) and quark (Q) matter,

$$P_H(\mu_n, \mu_e, \{\phi\}) = P_Q(\mu_n, \mu_e)\,, \tag{15}$$

where P_H and P_Q denote the pressures of hadronic matter and quark matter, respectively [18]. The quantity $\{\phi\}$ in Eq. (15) stands collectively for the field variables ($\bar{\sigma}$, $\bar{\omega}$, $\bar{\rho}$) and Fermi momenta (k_B, k_λ) that characterize a solution to the equations of confined hadronic matter. We use the symbol $\chi \equiv V_Q/V$ to denote the volume proportion of quark matter, V_Q, in the unknown volume V. By definition, χ then varies between 0 and 1, depending on how much confined hadronic matter has been converted to quark matter. Equation (15) is to be supplemented with the conditions of global baryon charge conservation and global electric charge conservation. The global conservation of baryon charge

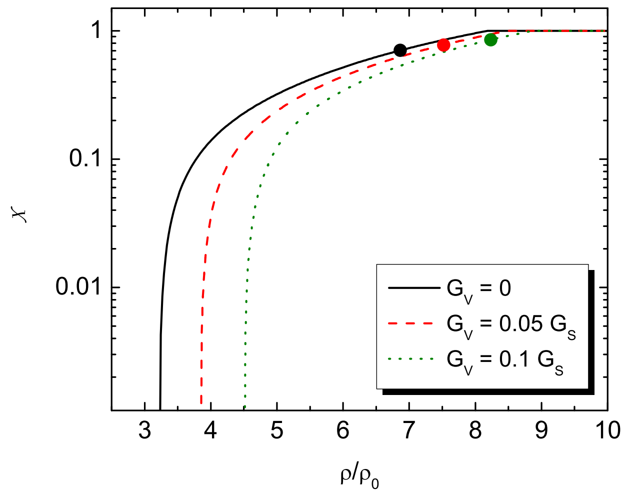


FIG. 2. (Color online) Volume fraction, χ, of quark phase as a function of baryon number density, ρ, in units of normal nuclear matter density, $\rho_0 = 0.16$ fm^{-3}. The solid dots indicate the central densities of the respective maximum-mass stars shown in Fig. 3.

is expressed as [18]

$$\rho_b = \chi\, \rho_Q(\mu_n, \mu_e) + (1-\chi)\, \rho_H(\mu_n, \mu_e, \{\phi\})\,, \tag{16}$$

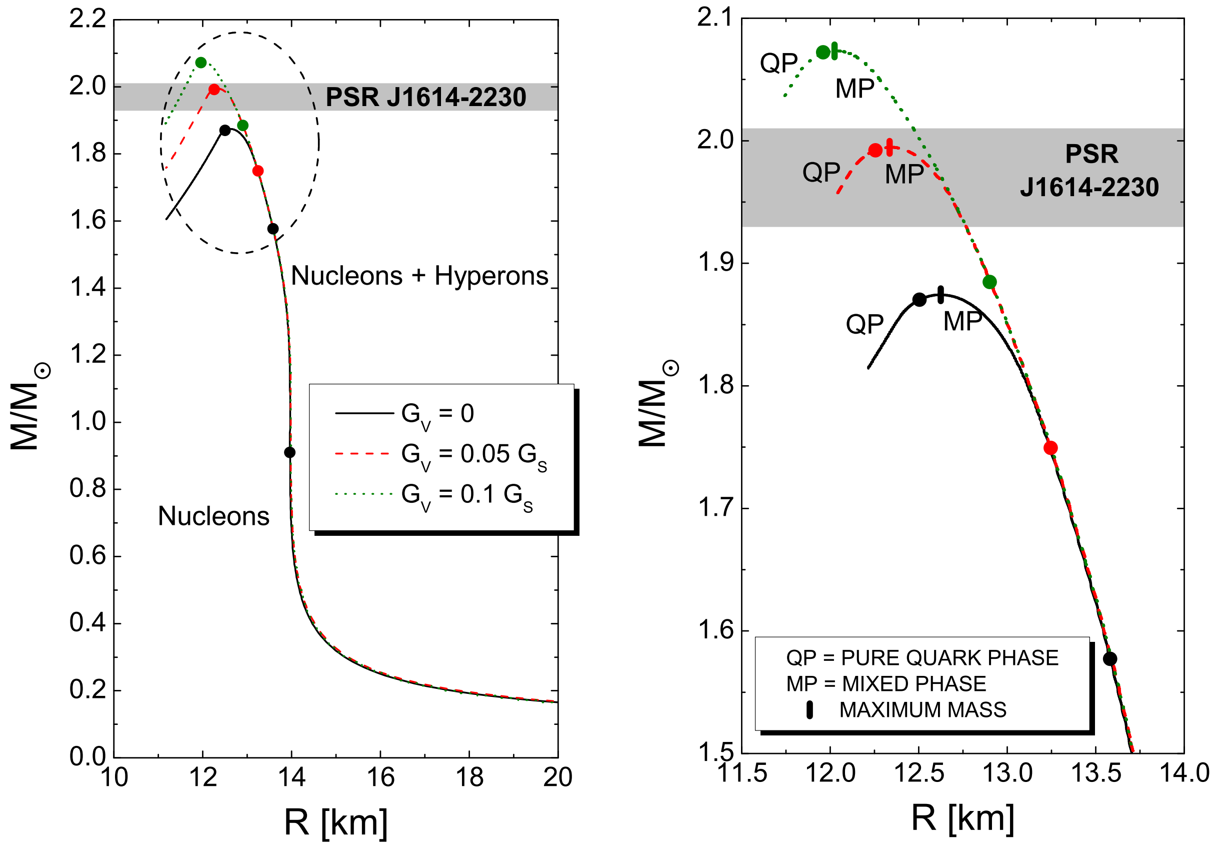
where ρ_Q and ρ_H denote the baryon number densities of the quark phase and hadronic phase, respectively. The global neutrality of electric charge is given by [18]

$$0 = \chi\ q_Q(\mu_n, \mu_e) + (1-\chi)\ q_H(\mu_n, \mu_e, \{\phi\})\,, \tag{17}$$

with q_Q and q_H denoting the electric charge densities of the quark phase and hadronic phase, respectively. 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 β–equilibrium and relativistic quantum statistics, as shown in [28]. Local NJL studies carried out for local electric charge neutrality have been reported recently in Refs. [8, 9]. In Ref. [9] the nonlinear relativistic mean-field model GM1 and the local SU(3) NJL model with vector interaction were used to describe the hadronic and quark matter phases, respectively. The authors found that the observation of neutron stars with masses a few percent higher than the $1.97 \pm 0.04\, M_\odot$ would be hard to explain unless, instead of GM1, one uses a very stiff model for the hadronic EOS with nucleons only. They obtain neutron stars with stable pure quark matter cores in their centers. In contrast to this, we find that such neutron stars are not be stable if the nonlocal NJL model is used instead of the local model and the less stringent condition of global electric charge neutrality is imposed on the composition of the stellar matter.

The results for the mixed phase region are shown in Figs. 1 and 2. Our calculations show that the inclusion of the quark vector coupling contribution shifts the onset of the phase transition to higher densities, and also narrows the width of the mixed quark-hadron phase, when compared to the case $G_V = 0$. The mixed phases range from $3.2 - 8.2\rho_0$, $3.8 - 8.5\rho_0$, and $4.5 - 8.9\rho_0$ for vector coupling constants $G_V/G_s = 0,\ 0.05,\ 0.1$, respectively. We note that there is considerable theoretical uncertainty in the ratio of G_V/G_s [29] since a rigorous derivation of the effective couplings from QCD is not possible. Combining the ratios of G_V/G_s from the molecular instanton liquid model and from the Fierz transformation, the value of G_V/G_s is expected to be in the range $0 \leqslant G_V/G_s \leqslant 0.5$ [30]. For our model, values of $G_V/G_s > 0.1$ shift the onset of the quark-hadron phase transition to such high densities that quark deconfinement can not longer occur in the cores of neutron stars.

Next we determine the bulk properties of spherically symmetric neutron stars for the collection of equations of state shown in Fig. 1. The properties are determined by solving the Tolmann-Oppenheimer-Volkoff (TOV) equation of general relativity theory [31]. The outcome is shown in Fig. 3. One sees that depending on the value of the vector coupling constant, G_V, the maximum neutron star masses increase from $1.87\, M_\odot$ for $G_V = 0$, to $2.00\, M_\odot$ for $G_V = 0.05\, G_s$, to $2.07\, M_\odot$ for $G_V = 0.1\, G_s$. The heavier stars of all three stellar sequences contain mixed phases of quarks and hadrons in their centers. The densities in these stars are however not high enough to generate pure quark matter in the cores. Such matter forms only in neutron stars which are already located on the gravitationally unstable branch of the mass-radius relationships. Another intriguing finding is that neutron stars with canonical masses of around $1.4\ M_\odot$ do not posses a mixed phase of quarks and hadrons but are made entirely of confined hadronic matter.



(a)Mass-radius relationships of neutron stars with quark-hybrid matter in their centers.

(b)Enlargement of the circled region of Fig. 3(a). The labels 'MP' and 'QP' stand for mixed phase and pure quark phase, respectively. The solid vertical bars denote maximum-mass stars.

FIG. 3. (Color online) Depending on the strength of the vector repulsion (G_V) of the nonlocal NJL model, maximum masses up to $2.1\,M_\odot$ are obtained. With increasing stellar mass, 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). The latter, however, exists only in neutron stars which lie to the left of their respective mass peaks. 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. The gray band denotes the 1-σ error bar of the $M = (1.97 \pm 0.04)M_\odot$ neutron star PSR J1614-2230 [6].

Summary and Conclusions – In this paper, we show that high-mass neutron star, such as PSR J1614–2230 with a gravitational mass of $1.97 \pm 0.04\,M_\odot$ [6], may contain mixtures of quarks and hadrons in their central regions. Our analysis is based on a nonlocal extension of the SU(3) Nambu-Jona Lasinio model, which reproduces some of the key features of Quantum Chromodynamics at densities relevant to neutron stars. Critical is the inclusion of the quark vector coupling contribution in the nonlocal SU(3) NJL model. Our results also show that the transition to pure quark matter occurs only in neutron stars which lie already on the gravitationally unstable branch of the mass-radius relationship. The existence of pure quark matter in massive (as well as in all other, less massive) neutron stars would thus be ruled out by our study.

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