HYBRID STARS IN THE LIGHT OF THE MASSIVE PULSAR PSR J1614–2230

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ABSTRACT

We perform a systematic study of hybrid star configurations using several parameterizations of a relativistic mean-field hadronic equation of state (EoS) and the Nambu–Jona–Lasinio (NJL) model for three-flavor quark matter. For the hadronic phase we use the stiff GM1 and TM1 parameterizations, as well as the very stiff NL3 model. In the NJL Lagrangian we include scalar, vector, and ’t Hooft interactions. The vector coupling constant $g_v$ is treated as a free parameter. We also consider that there is a split between the deconfinement and the chiral phase transitions which is controlled by changing the conventional value of the vacuum pressure $-\Omega_0$ in the NJL thermodynamic potential by $-(\Omega_0 + \delta \Omega_0)$, with $\delta \Omega_0$ a free parameter. We find that, as we increase the value of $\delta \Omega_0$, hybrid stars have a larger maximum mass but are less stable, i.e., hybrid configurations are stable within a smaller range of central densities. For large enough $\delta \Omega_0$, stable hybrid configurations are not possible at all. The effect of increasing the coupling constant $g_v$ is very similar. We show that stable hybrid configurations with a maximum mass larger than the observed mass of the pulsar PSR J1614–2230 are possible for a large region of the parameter space of $g_v$ and $\delta \Omega_0$ provided the hadronic EoS contains nucleons only. When the baryon octet is included in the hadronic phase, only a very small region of the parameter space allows an explanation of the mass of PSR J1614–2230. We compare our results with previous calculations of hybrid stars within the NJL model. We show that it is possible to obtain stable hybrid configurations also in the case $\delta \Omega_0 = 0$ that corresponds to the conventional NJL model for which the pressure and density vanish at zero temperature and chemical potential.

Key words: equation of state – pulsars: individual (PSR J1614–2230) – stars: neutron

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1. INTRODUCTION

The recent determination of the mass of the pulsar PSR J1614–2230 as $1.97 \pm 0.04 \, M_\odot$ by Demorest et al. (2010) renewed discussions about the possibility of exotic matter being present at the core of neutron stars. Since the description of matter at densities beyond nuclear saturation is model dependent, several works have explored different aspects of the fact that the maximum neutron star mass implied by any equation of state (EoS) must exceed the mass of PSR J1614–2230 (Weissenborn et al. 2011a, 2011b, 2012; Bednarek et al. 2011; Lastowiecki et al. 2012; Bonanno & Sedrakian 2012). Some authors have revisited the role of hyperons in the EoS, showing that it is possible to construct stiff EoS with hyperons that are compatible with up-to-date hypernuclear data (Bednarek et al. 2011). Others have investigated the role of the vector meson–hyperon coupling, going from the SU(6) quark model to a broader SU(3) symmetry (Weissenborn et al. 2011a), and of hyperon potentials (Weissenborn et al. 2012) in order to determine their impact on the maximum mass of neutron stars.

Concerning quark matter, it is known that models of strange stars made of absolutely stable quark matter satisfy comfortably the new constraint if color superconductivity is taken into account (Lugones & Horvath 2003; Horvath & Lugones 2004). However, it is not straightforward to construct models of hybrid stars with more than two solar masses (Benhar & Cipollone 2011; Weissenborn et al. 2011b). A recurrent difficulty is that most hybrid EoS do not have at the same time a stable quark matter core and a sufficiently large maximum mass. For instance, most versions of the widely used Nambu–Jona–Lasinio (NJL) model are too soft to meet any of the above requirements. In a recent work, Benhar & Cipollone (2011) performed a systematic study of the role of the vector and instanton-induced terms in the NJL Lagrangian and their connection with the properties of hybrid stars. They explored a broad region of the parameter space showing that the instanton-induced interaction does not affect the stiffness of the quark matter EoS, whereas the effect of the repulsive vector interaction is significant. However, according to these authors, no values of the corresponding coupling constants allow for the formation of a stable core of quark matter (Benhar & Cipollone 2011). These conclusions are in qualitative agreement with previous results using a similar EoS for the hadronic phase but including color superconductivity and neglecting the vector interaction term in the NJL EoS (Baldo et al. 2003), and with NJL models that implement the use of a density-dependent cutoff (Baldo et al. 2007). Also, Coelho et al. (2010) use an SU(2) NJL model with a vector term but they still find unstable hybrid stars. Within a different picture, Lenzi et al. (2010) show that it is possible to obtain a stable sequence of compact stars with a quark core using an NJL model with an ad hoc momentum cutoff that depends on the baryon chemical potential. However, within this approach, the maximum mass is still below the mass of PSR J1614–2230. According to Benhar & Cipollone (2011), their results are not essentially affected by the assumption that the hadronic phase consists of nucleons only, or by the formation of mixed phases. However, more recent work by Bonanno & Sedrakian (2012) has succeeded in obtaining very massive stable hybrid configurations using a three-flavor NJL model that contains two free parameters: the transition density from hadronic matter to quark matter and the vector coupling of quarks. They show that high-mass stable configurations with color-superconducting quark cores can be constructed if vector interactions are included in the quark phase and if a very stiff hadronic EoS is employed (e.g., the NL3 model with hyperons or the GM3 model with nucleons).
In this work, we perform an extensive study of hybrid star masses using several parameterizations of a relativistic mean-field hadronic EoS together with a typical three-flavor NJL model with scalar, vector, and ‘t Hooft interactions. Within this approach, the hadronic and quark–gluon degrees of freedom are derived from different Lagrangians and the deconfinement transition is associated with the point where both models have the same free energy. Thus, by construction, chiral symmetry is restored at the conventional value of the vacuum pressure $-\Omega_0$, we shall use a value $\Omega_0$ which is in general different to the $\Omega_0$ that is in general different to the same free energy. Thus, by construction, chiral symmetry transition is associated with the point where both models have the same free energy.

Table 1

| Set  | GM1  | TM1  | NL3  |
|------|------|------|------|
| $m_\rho$ (MeV) | 1.78 | 1.78 | 1.76 |
| $m_\omega$ (MeV) | 783 | 783 | 782.5 |
| $m_\rho$ (MeV) | 770 | 770 | 763 |

\[ \begin{align*}
\Omega_H &= \sum_{B} \bar{\psi}_B \gamma_5 (i \partial^\mu - g_{\omega B} \omega^\mu - \frac{1}{2} g_{\rho B} \tau^\mu \rho^\mu) \\
&\quad - (m_B - g_{\omega B} \sigma) \psi_B + \frac{1}{2} \left( \partial^\mu \sigma \partial^\sigma - m_\sigma^2 \sigma^2 \right) \\
&\quad - \frac{1}{4} \omega_\mu \omega^\mu + \frac{1}{2} m_\omega^2 \omega^\mu \omega^\nu + \frac{1}{4} \rho_\mu \rho^\nu \\
&\quad + \frac{1}{2} m_\rho^2 \rho^\mu \rho^\nu - \frac{1}{3} b m_\rho (g_\sigma \sigma^3) - \frac{1}{4} c (g_\sigma \sigma)^4 \\
&\quad + \sum_{L} \bar{\psi}_L [i \gamma_\mu \partial^\mu - m_L] \psi_L,
\end{align*} \]

2. HADRPHIC PHASE

The relativistic mean-field model is widely used to describe hadronic matter in compact stars. In this paper we adopt the following standard Lagrangian (Boguta & Bodmer 1977; Glendenning & Moszkowski 1991):

3. QUARK PHASE

3.1. The Model

To describe the quark matter phase we use the SU(3) NJL model with scalar–pseudoscalar, isoscalar–vector, and ‘t Hooft six-fermion interactions. The Lagrangian density of the model is

\[ \begin{align*}
\mathcal{L}_Q &= \bar{\psi}(i \gamma_\mu \partial^\mu - m_\psi \psi \\
&\quad + g_\psi \sum_{a=0}^{8} [(\bar{\psi} \lambda^a \psi)^2 + (\bar{\psi} i \gamma_5 \lambda^a \psi)^2] \\
&\quad - g_v \sum_{a=0}^{8} [(\bar{\psi} \gamma_\mu \lambda^a \psi)^2 + (\bar{\psi} \gamma_5 \gamma_\mu \lambda^a \psi)^2] \\
&\quad + g_\psi \{ \text{det} [\bar{\psi} (1 + \gamma_5)] + \text{det} [\bar{\psi} (1 - \gamma_5)] \},
\end{align*} \]

where $\psi = (u, d, s)$ denotes the quark fields, $\lambda^a (0 \leq a \leq 8)$ are the U(3) flavor matrices, $m = \text{diag}(m_u, m_d, m_s)$ is the quark current mass, and $g_\psi, g_\sigma$, and $g_\rho$ are coupling constants.

Note that we have not included a diquark interaction term in the Lagrangian. As shown by Rüster et al. (2005), color-superconducting phases (at $T = 0$) are favored in the regime of strong diquark coupling, $g_d/g_\sigma \approx 1$. However, in the regime of intermediate diquark coupling strength, $g_d/g_\sigma = 3/4$, color...
superconductivity appears only above a chemical potential \( \mu \sim 3 \times 440 \text{ MeV} = 1320 \text{ MeV} \) (see Figure 1 of Rüster et al. 2005). For weaker diquark coupling, color superconductivity is shifted to very large densities that are not present at neutron star cores. Since the case of strong diquark coupling has already been considered by Bonanno & Sedrakian (2012) and Pagliara & Schaffner-Bielich (2008), we shall focus here on a case where color superconductivity is negligible.

The mean-field thermodynamic potential density \( \Omega \) for a given baryon chemical potential \( \mu \) at \( T = 0 \) is given by

\[
\Omega = -\eta N_f \sum_i \int_{p_i}^{\Lambda} \frac{p^2 dp}{2\pi^2} \sqrt{p^2 + M_i^2} + 2 g \sum_i \langle \bar{\psi} \psi \rangle^2
- 2 g \sum_i \langle \bar{\psi} i \gamma^5 \psi \rangle^2 + 4 g \langle \bar{\tau} u \rangle \langle \bar{\tau} d \rangle \langle \bar{\tau} s \rangle
- \eta N_f \sum_i \mu_i \int_{p_i}^{\Lambda} \frac{p^2 dp}{2\pi^2} - \Omega_0,
\]

where the sum is over the quark flavor \((i = u, d, s)\), the constants \( \eta = 2 \) and \( N_f = 3 \) are the spin and color degeneracies, and \( \Lambda \) is a regularization ultraviolet cutoff to avoid divergences in the medium integrals. The Fermi momentum of the particle \( i \) is given by \( k_F = \sqrt{\mu_i^2 + M_i^2} - M_i \), where \( \mu_i \) is the quark chemical potential modified by the vectorial interaction, i.e., \( \mu_{u,d,s} = \mu_{u,d,s} - 4g_i \langle \bar{\tau} i \gamma^5 \tau \rangle \langle \bar{\tau} i \gamma^5 \tau \rangle \).

In this work we consider the following set of parameters (Kunihiro 1989; Ruivo et al. 1999): \( \Lambda = 631.4 \text{ MeV}, g_s A_s^2 = 1.829, g_s A_s^2 = -9.4, m_u = m_d = 5.6 \text{ MeV}, m_s = 135.6 \text{ MeV} \) in order to fit the vacuum values for the pion mass, the pion decay constant, the kaon mass, the kaon decay constant, and the quark condensates: \( m_{\pi} = 139.0 \text{ MeV}, f_{\pi} = 93.0 \text{ MeV}, m_{K} = 495.7 \text{ MeV}, f_{K} = 98.9 \text{ MeV}, \langle \bar{u}u \rangle^{1/3} = \langle \bar{d}d \rangle^{1/3} = -246.7 \text{ MeV}, \langle \bar{s}s \rangle^{1/3} = -266.9 \text{ MeV} \). The value of the vector coupling constant \( g \) is treated as a free parameter because the masses of the vector mesons are not dictated by chiral symmetry. In order to obtain the EoS, we assume that matter is charge neutral and in equilibrium under weak interactions.

3.2. \( \Omega_0 \) as a Free Parameter

The conventional procedure for fixing the \( \Omega_0 \) term in Equation (3) is to assume that the grand thermodynamic potential \( \Omega \) must vanish at zero \( \mu \) and \( T \). For the parameterization quoted above, this assumption leads to the value \( \Omega_0 = 5076.2 \text{ MeV fm}^{-3} \). Nevertheless, this prescription is no more than an arbitrary way to uniquely determine the EoS of the NJL model without any further assumptions (Schertler et al. 1999). Furthermore, in the MIT bag model, for instance, the pressure in the vacuum is non-vanishing. In view of this, Pagliara & Schaffner-Bielich (2008) adopt a different strategy. They fix a bag constant for the hadron–quark deconfinement to occur at the same chemical potential as the chiral phase transition. This method leads to a significant change in the EoS with respect to the conventional procedure.

The connection between the chiral and the deconfinement phase transitions along the QCD phase diagram has received considerable attention in recent years (see Fukushima & Hatsuda 2011, and references therein). For zero chemical potential, lattice results show that both transitions occur at the same temperature (Karsch 2002; Laermann & Philipson 2003). At finite baryon chemical potential, this coincidence is an open question (Fukushima & Hatsuda 2011). However, since a chiral critical endpoint may exist in the plane of \( T \) and \( \mu \) (Stephanov 2005), it has been conjectured that the deconfinement and chiral transitions split from one another at that point (McLerran & Pisarski 2007). As a consequence, a confined but chiral symmetric phase, called the quarkyonic phase, can exist in the region of high baryon density. Since this conjecture is based on arguments that are valid in the large-\( N_c \) limit, it is not clear whether this quarkyonic phase can exist in the real QCD phase diagram. In the present paper, we are not modeling the quarkyonic matter because our confined phase is described by a hadronic model and the chiral transition is restricted to the quark phase. However, we may explore the above possibility of having chiral restoration and deconfinement occurring at different densities. To this end, we shall substitute \( \Omega_0 \) in Equation (3) by the new value \( \Omega_0 + \delta \Omega_0 \), where \( \delta \Omega_0 \) is a free parameter:

\[
\Omega_0 \rightarrow \Omega_0 + \delta \Omega_0 \quad \text{in Equation (3).}
\]

With this change, the thermodynamic potential \( \Omega \) can be non-vanishing at zero \( \mu \) and \( T \), and the \( \mu \) of the deconfinement transition can be tuned.\(^1\) Clearly, \( \delta \Omega_0 \) has a minimum value because the phase transition cannot be shifted to a pressure regime where the NJL model describes the vacuum. That is, we fix a minimum limit to \( \delta \Omega_0 \) for which the phase transition occurs at the chiral symmetry restoration point as performed by Pagliara & Schaffner-Bielich (2008). On the other hand, in principle there is no maximum value for \( \delta \Omega_0 \) since the phase transition can be shifted to arbitrarily large pressures.

In order to illustrate the dependence of the EoS on the new parameter \( \delta \Omega_0 \), we depict in Figure 1 the pressure as a function of the chemical potential for different values of \( \delta \Omega_0 \) and the pressure of the deconfinement transition \( P_{\text{ph}} \) as a function of \( \delta \Omega_0 \). Note that a small change in the value of \( \delta \Omega_0 \) may result in a significant modification of the phase transition density, and consequently in a very different hybrid EoS.

4. RESULTS

We have solved the Tolman–Oppenheimer–Volkoff equations for spherically symmetric and static stars in order to investigate the influence of \( g \) and \( \delta \Omega_0 \) on the maximum mass of hybrid stars.

In Figures 2 and 3 we show the EoS for some specific parameterizations and the corresponding stellar configurations in a diagram of mass \( M \) versus central energy density \( \epsilon_c \). The plateaus represent a first-order hadron–quark phase transition where both phases have the same pressure and Gibbs free energy per baryon. We consider that charge neutrality is fulfilled locally, i.e., each phase is separately charge neutral. This leads to a sharp discontinuity in the density profile of the star. In the left panel of Figure 2 we note that as we increase the value of the vector coupling constant \( g \), the density of the phase transition also increases, and therefore the hybrid star has a smaller quark core and a larger hadronic contribution. This leads to larger maximum masses because the hadronic EoS is stiffer than the quark EoS. At the same time, there is a larger density jump between the two

\(^1\) The chemical potential at which the chiral transition occurs is determined from the solution of the gap equations for the constituent masses and therefore it does not depend on the value of \( \Omega_0 \). However, the chemical potential for the deconfinement phase transition depends on \( \Omega_0 \) because it is determined by matching the pressures of the hadronic and quark phases. As a consequence, tuning \( \Omega_0 \) is an easy way to control the splitting between both chemical potentials.
phases, which tends to destabilize the configuration. Due to these two effects, together with the fact that the vector term stiffens the NJL EoS, models with a larger $g_v$ give larger maximum masses but have stable quark cores within a smaller range of central densities (see the right panel of Figure 2). In the left panel of Figure 3, we show the effect of changing the magnitude of the shift between the deconfinement and the chiral phase transitions. As we increase $\delta \Omega_0$ from negative to positive values, we increase the density of the phase transition as well as the density jump between the two phases. However, the NJL EoS becomes slightly softer because there is a larger contribution to the EoS of the regime with a partially restored chiral symmetry. Since the latter effect is not so strong, the impact of increasing $\delta \Omega_0$ is analogous to increasing $g_v$, i.e., models with a larger $\delta \Omega_0$ result in larger maximum masses but the quark cores are stable within a smaller range of central densities (see the right panel of Figure 3).

In Figure 4, background colors represent the maximum mass of hybrid stars for different values of $g_v$ and $\delta \Omega_0$. We used the hadronic EoS of Table 1 with nucleons and electrons. Within each panel we show contour lines indicating specific values of the maximum mass. The black line represents the limit between parameters that allow for stable hybrid stars and those that always give unstable hybrid stars. The value 1.97 $M_{\odot}$, corresponding to the observed mass of PSR J1614$-$2230 (Demorest et al. 2010), is shown with a red dashed line. An interesting feature of Figure 4 is that large masses are situated on the right-upper corner but stable configurations are located on the left-lower corner of the figure (or left side of the figure in the case of NL3). This clearly illustrates the difficulty of obtaining stable hybrid stars with arbitrarily large masses. Concerning the effect of the hadronic model, we see that stable hybrid stars have higher values of the maximum mass for the stiffer hadronic EoS.

The observed mass of PSR J1614$-$2230 can be explained by parameters within the large region located between the red dashed line and the solid black line in each panel of Figure 4. However, a hypothetical future observation of a neutron star with a mass $\sim 10\%$ larger than the mass of PSR J1614$-$2230 will be hard to explain within hybrid star models using the GM1 and TM1 EoS (see panels (a) and (b) of Figure 4) and will require a very stiff hadronic model such as NL3.

The effect of hyperons is shown in Figure 5, where we consider the NL3 parameterization with the inclusion of the baryon octet. Compared with the case without hyperons, the maximum mass values are altered by a few percent. This follows from the fact that the deconfinement phase transition occurs at relatively low densities, i.e., in regions where the baryon octet has a minor contribution. Nevertheless, hyperons have a large effect on the possibility of finding stable hybrid stars of large enough mass. When we increase the value of $g_v$ or increase $\delta \Omega_0$, the deconfinement transition is shifted to larger densities and the hadronic EoS with hyperons tends to be favored.

Figure 1. (a) Pressure as a function of the chemical potential for different values of the parameter $\delta \Omega_0$. (b) Pressure of the deconfinement phase transition as a function of $\delta \Omega_0$ for different values of the coupling constant $g_v$. Note that a small change in $\delta \Omega_0$ can produce a significant change in the pressure of the phase transition. (A color version of this figure is available in the online journal.)

Figure 2. (a) Pressure as a function of the baryon number density in units of the nuclear saturation density $\rho_0$ (we assumed $\rho_0 = 0.17$ fm$^{-3}$). (b) Mass of hybrid stars as a function of the central mass-energy density $\epsilon_c$. We use $\delta \Omega_0 = 0$ and different values of $g_v$. (A color version of this figure is available in the online journal.)

Figure 3. Same as Figure 2 but adopting $g_v/g_s = 0.2$ and different values of $\delta \Omega_0$ (labels for $\delta \Omega_0$ are in MeV fm$^{-3}$). (A color version of this figure is available in the online journal.)

Figure 4. Background colors represent the maximum mass of hybrid stars for different values of $g_v$ and $\delta \Omega_0$. We used the hadronic EoS of Table 1 with nucleons and electrons. Within each panel we show contour lines indicating specific values of the maximum mass. The black line represents the limit between parameters that allow for stable hybrid stars and those that always give unstable hybrid stars. The value 1.97 $M_{\odot}$, corresponding to the observed mass of PSR J1614$-$2230 (Demorest et al. 2010), is shown with a red dashed line. An interesting feature of Figure 4 is that large masses are situated on the right-upper corner but stable configurations are located on the left-lower corner of the figure (or left side of the figure in the case of NL3). This clearly illustrates the difficulty of obtaining stable hybrid stars with arbitrarily large masses. Concerning the effect of the hadronic model, we see that stable hybrid stars have higher values of the maximum mass for the stiffer hadronic EoS.

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Figure 4. Background colors represent the maximum mass of hybrid stars for different parameterizations of the NJL model (i.e., different values of $g_v$ and $\delta \Omega_0$). In each panel, we use a different hadronic EoS (without hyperons): (a) GM1, (b) TM1, and (c) NL3 (see Table 1). Note that the color scale is different for each panel. The solid contour lines indicate specific values of the maximum mass. The black solid line represents the boundary between parameterizations that allow for stable hybrid stars and parameterizations that do not. The red dashed line indicates the value $1.97 M_\odot$ corresponding to the observed mass of PSR J1614$-$2230 (Demorest et al. 2010). The mass of PSR J1614$-$2230 can be explained by parameters within the region between the red dashed line and the solid black line.

(A color version of this figure is available in the online journal.)

Figure 5. Same as panel (c) of Figure 4 but for the NL3 model with hyperons. Hybrid stars are not possible for the set of parameters within the white region. Only parameters in a very small region near the upper-left corner of the colored region explain the mass of PSR J1614$-$2230.

(A color version of this figure is available in the online journal.)

5. DISCUSSION AND CONCLUSIONS

In this work, we performed a systematic study of hybrid star configurations using a relativistic mean-field hadronic EoS and the NJL model for three-flavor quark matter. For the hadronic phase we used the stiff GM1 and TM1 parameterizations, as well as the very stiff NL3 model. In the NJL Lagrangian we included scalar, vector, and 't Hooft interactions. The vector coupling constant $g_v$ was treated as a free parameter. We also considered that there is an arbitrary split between the deconfinement and the chiral phase transitions. This split can be adjusted by a redefinition of the constant parameter $\Omega_0$ in the NJL thermodynamic potential, i.e., by making the replacement $\Omega_0 \rightarrow \Omega_0 + \delta \Omega_0$ in Equation (3), where $\delta \Omega_0$ is a free parameter. We find that, as we increase the value of $\delta \Omega_0$, hybrid stars have a larger maximum mass but are less stable (i.e., hybrid configurations are stable within a smaller range of central densities). For large enough $\delta \Omega_0$, stable hybrid configurations are not possible at all (see Figure 3). The effect of increasing the coupling constant $g_v$ is very similar (see Figure 2). These effects are clear in Figure 4, where we show the maximum mass of static spherically symmetric stars in the parameter space of $g_v$ and $\delta \Omega_0$.

Almost everywhere in the star, above a certain limit there is no deconfinement transition at all, i.e., the star is always hadronic (see the white region in Figure 5). As a consequence, the mass of PSR J1614$-$2230 can be attained for models within a very small region of the parameter space located near the upper-left corner of the colored region in Figure 5.
should apply to a broader range of stiff hadronic models, e.g., models including microscopic three-body forces among nucleons (see Schulze 2011, and references therein). The main effect of hyperons is that they preclude the deconfinement transition in the region of the parameter space that allows large maximum masses (see Figure 5). Only a very small area near the upper-left corner of the colored region of Figure 5 explains the mass of PSR J1614–2230. According to recent studies, one does not expect substantial changes when introducing refinements in the description of hyperons, such as hyperon–hyperon potentials, hyperonic three-body forces, etc. (see, e.g., Vidaña et al. 2011; Schulze 2011, and references therein).

It is worth summarizing the main assumptions and results of recent work using the NJL model to describe hybrid stars. Schertler et al. (1999) use several parameterizations of an extended relativistic mean-field model to describe the hadronic phase with hyperons. For the deconfined quark phase they use the NJL model with three flavors, including scalar and ‘t Hooft terms, and using the RKH parameter set (Rehberg et al. 1996). Their conclusion is that typical neutron stars with masses around 1.4 solar masses do not possess any deconfined quark matter in their center. More recently, Baldo et al. (2007) use a SU(2) NJL model with the RKH parameterization and a cutoff that depends on the chemical potential. For the hadronic phase they adopt a nucleonic EoS obtained within the Brueckner–Bethe–Goldstone approach using the Argonne v18 two-body potential, supplemented by the Urbana phenomenological three-body force. They are not able to obtain stable hybrid stars. Coelho et al. (2010) use a SU(2) NJL model with the RKH parameterization but adding a vector term together with the GM1 parameterization of a relativistic mean-field hadronic model (Glendenning & Moszkowski 1991). The density of the phase transition increases with $g_8$, but they still find unstable hybrid stars. Benhar & Cipollone (2011) work with SU(3) and the RKH parameterization. They use several values of $g_8$ and they also vary the coupling constant of the ‘t Hooft term. For the hadronic phase they use a phenomenological Hamiltonian including the Argonne v18 nucleon–nucleon potential. They conclude that no values of the corresponding coupling constants allow for the formation of a stable core of quark matter.

All the above versions of the NJL model use the conventional procedure of imposing that the pressure and density must vanish at zero temperature and chemical potential. However, a different prescription is used by Pagliara & Schaffner-Bielich (2008). First, they introduce a hadronic EoS and compute the transition to quark matter by a Maxwell construction. To fix the bag constant they assume that deconfinement occurs at the same chemical potential as the chiral phase transition, i.e., they require that the pressure of quark matter is equal to the pressure of the hadronic matter at the critical chemical potential for which chiral symmetry is restored. The bag value obtained with this assumption is the lowest possible value for the bag constant in the NJL model because it allows use of the NJL EoS just starting from the chemical potential of the chiral phase transition. Using this procedure, Pagliara & Schaffner-Bielich (2008) computed the EoS of quark matter within the NJL model by including effects from the chiral condensates, the diquark coupling pattern, and a repulsion vector term. They find that hybrid stars containing a CFL core are stable but the maximum mass is $\sim1.8 M_\odot$, i.e., incompatible with PSR J1614–2230. More recently, Bonanno & Sedrakian (2012) use an NJL model supplemented by the ‘t Hooft and vector interactions and consider the 2SC and CFL color-superconducting phases. For the hadronic phase they use a relativistic mean-field model and adopt the NL3 parameterization with hyperons and the GM3 parameterization with nucleons only. In both cases they are able to obtain maximum masses above the mass of PSR J1614–2230 because they use the non-conventional procedure of treating the density of the deconfinement transition as a free parameter.

Our analysis is related to that of Bonanno & Sedrakian (2012) because we can control the density of the phase transition via the parameter $\delta \Omega_0$ in the thermodynamic potential. However, since our parameter space is constructed in terms of $\delta \Omega_0$, and not in terms of the deconfinement density as in Bonanno & Sedrakian (2012), the connection with conventional NJL models is more transparent in our case. Additionally, our procedure includes the non-conventional prescription of Pagliara & Schaffner-Bielich (2008) as a special case. It is also worth noticing that we use the parameterization of the NJL EoS given by Kunihiro (1989) and Ruivo et al. (1999) while the above authors use the RKK one which is somewhat softer (see Buballa 2004 for more details on the parameterizations). Another difference is that they consider quark matter in the 2SC and CFL phases, whereas we do not consider color superconductivity in the quark phase. In this sense, these works are complementary because Pagliara & Schaffner-Bielich (2008) and Bonanno & Sedrakian (2012) work in the regime of strong diquark coupling ($g_8/g_8 \approx 1$) where color superconductivity is strongly favored, whereas we work in the weak diquark coupling regime ($g_8/g_8 \to 0$) for which color superconductivity is shifted to very large densities that are not present at neutron star cores (see Rüster et al. 2005 for more details).

In summary, our results show that hybrid configurations with maximum masses equal to or larger than the observed mass of PSR J1614–2230 are possible for a significant region of the parameter space of $g_8$ and $\delta \Omega_0$ provided a stiff enough hadronic EoS without hyperons is used. It is also worth highlighting the fact that we can obtain compact stars with stable quark cores without having to perform any modification to the NJL model (i.e., setting $\delta \Omega_0 = 0$). This is in contrast to the results obtained by Baldo et al. (2007), Benhar & Cipollone (2011), and Coelho et al. (2010), who are unable to obtain stable hybrid star configurations with the NJL model. The difference arises because these authors use the softer parameterization of Rehberg et al. (1996) while we use the stiffer parameterization of Kunihiro (1989) and Ruivo et al. (1999), i.e., the use of the latter parameterization allows us to reproduce the mass of the pulsar PSR J1614–2230 without using exceptionally stiff parameterizations of the hadronic EoS and keeping the conventional procedure for fixing $\Omega_0$. It is also interesting to note that we have typically $\delta \Omega_0/\Omega_0 \sim 0.1\%$, i.e., small departures of the conventional $\Omega_0$ are sufficient to obtain massive enough stable hybrid stars. Finally, we emphasize that the observation of compact star masses a few percent larger than the mass of PSR J1614–2230 will be hard to explain within hybrid star models using the GM1 and TM1 EoS, and will require a very stiff hadronic model such as NL3 with nucleons only.

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