Sequential deconfinement of quark flavors in neutron stars

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(Received 14 October 2008; revised manuscript received 24 November 2009; published 22 December 2009)

A scenario is suggested in which the three light quark flavors are sequentially deconfined under increasing pressure in cold asymmetric nuclear matter as found, for example, in neutron stars. The basis for this analysis is a chiral quark matter model of Nambu–Jona-Lasinio (NJL) type with diquark pairing in the spin-1 single-flavor, spin-0 two-flavor, and three-flavor channels. Nucleon dissociation sets in at about the saturation density, n_0, when the down-quark Fermi sea is populated (d-quark drip line) because of the flavor asymmetry induced by β equilibrium and charge neutrality. At about $3n_0$, u-quarks appear and a two-flavor color superconducting (2SC) phase is formed. The s-quark Fermi sea is populated only at still higher baryon density, when the quark chemical potential is of the order of the dynamically generated strange quark mass. Two different hybrid equations of state (EOSs) are constructed using the Dirac-Brueckner Hartree-Fock (DBHF) approach and the EOS of Shen *et al.* [H. Shen, H. Toki, K. Oyamatsu, and K. Sumiyoshi, Nucl. Phys. **A637**, 435 (1998)] in the nuclear matter sector. The corresponding hybrid star sequences have maximum masses of 2.1 and 2.0 $M_\odot$, respectively. Two- and three-flavor quark-matter phases exist only in gravitationally unstable hybrid star solutions in the DBHF case, whereas the Shen-based EOSs produce stable configurations with a 2SC phase component in the core of massive stars. Nucleon dissociation via d-quark drip could act as a deep crustal heating process, which apparently is required to explain superbursts and cooling of x-ray transients.

DOI: 10.1103/PhysRevC.80.065807 PACS number(s): 12.38.Mh, 11.10.St, 21.65.Cd, 26.60.−c

I. INTRODUCTION

The phenomenology of compact stars is intimately connected to the equations of state (EOSs) of matter at densities well beyond the nuclear saturation density, $n_0 = 0.16$ fm^{-3}. Compact stars are, therefore, natural laboratories for the exploration of baryonic matter under extreme conditions, complementary to those created in terrestrial experiments with atomic nuclei and heavy-ion collisions. Recent results derived from observations of compact stars provide serious constraints on the nuclear EOS (see Ref. [1] and references therein). A stiff EOS at high density is required to explain the high compact-star masses and radii, for which recent observations have provided growing evidence. A mass of $M \sim 2.0$ $M_\odot$ has been reported for some low-mass x-ray binaries (LMXBs; for example, 4U 1636-536 [2]) based on the assumption that the abrupt drop in the coherence of the lower kilohertz quasiperiodic oscillation (QPO) may be related to the innermost stable circular orbit (see also Ref. [3]). From observations of the bright isolated neutron star RX J1856.5-3754 (or RX J1856 in shorthand) in the optical and x-ray frequency ranges, a conservative lower limit of the apparent neutron star radius of $R_\infty = 16.5$ km is derived [4]. This value corresponds to a true (de-redshifted) radius of $R = 14$ km for a 1.4-$M_\odot$ neutron star or, equivalently, to a star mass of at least 2.1 $M_\odot$ when the radius does not exceed 12 km [1]. Another example is EXO 0748-676, a LMXB for which the compact-star mass *and* the radius have been constrained to $M \geqslant 2.10 \pm 0.28$ $M_\odot$ and $R \geqslant 13.8 \pm 0.18$ km [5] by a simultaneous measurement of the Eddington limit, the gravitational redshift, and the flux of thermal radiation. However, the status of the results for the latter object is unclear, because the gravitational redshift of $z = 0.35$ observed in the x-ray burst spectra [6] has not been confirmed, despite numerous attempts. Further constraints on the masses and radii of compact stars have been reported [7,8], but they deserve a careful discussion, which is beyond the scope of the present article. Although compact-star phenomenology apparently points toward a stiff EOS at high density, heavy-ion collision data for kaon production [9] and elliptic flow [10] set an upper limit on the stiffness of the EOS [1].

A key question regarding the structure of matter at high density is whether a phase transition to quark matter occurs inside compact stars and whether it is accompanied by unambiguous observable signatures. It has been argued that the observation of a compact star with high mass and large radius, as reported for EXO 0748-676, would be incompatible with a quark core [5], because quark deconfinement softens the EOS and lowers the maximum mass and corresponding radius. However, Alford *et al.* have demonstrated [11] with a few counterexamples that quark matter cannot be excluded by this argument. In particular, for a recently developed hybrid-star

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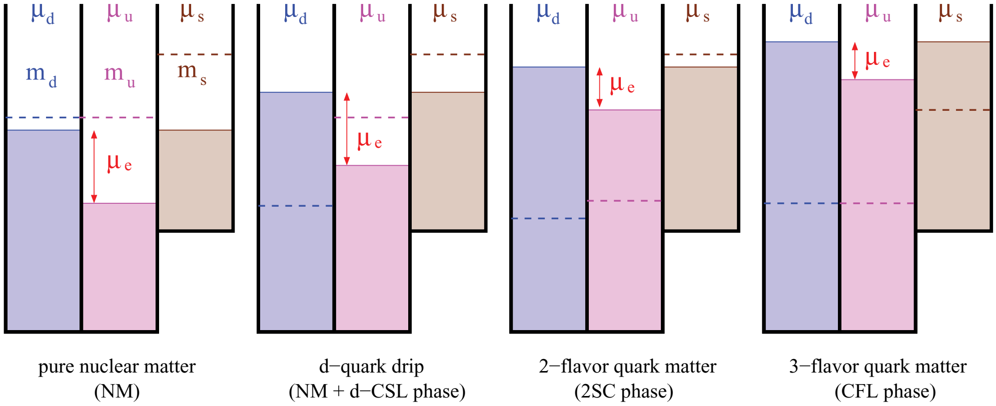

FIG. 1. (Color online) Schematic of chemical potentials (columns) and sequential deconfinement of quarks with increasing baryon density (from left to right). The flavor-dependent thresholds for chiral symmetry restoration (deconfinement) are approximately given by the dynamically generated quark masses m_f, $f = u, d, s$ (dashed lines). With increasing quark chemical potential, $\mu = (\mu_u + \mu_d)/2$, the d-quark chemical potential is the first to reach the threshold in isospin asymmetric matter. Nucleon dissociation therefore sets in as d-quarks are deconfined. Still higher μ is needed to form two-flavor and three-flavor quark matter phases.

EOS [12], based on the Dirac-Brueckner Hartree-Fock (DBHF) approach in the nuclear sector and a three-flavor chiral quark model [13], stable hybrid stars with masses ranging from 1.2 to 2.1 $M_\odot$ is obtained, in accordance with modern mass-radius constraints (see also Ref. [14]). In this model, a sufficiently low critical density for quark deconfinement has been achieved via a strong diquark coupling, while a repulsive vector mean field in the quark-matter sector resulted in sufficient stiffness to achieve a high maximum mass of the compact-star sequence. The corresponding hybrid EOS for symmetric matter was shown to fulfill the constraints derived from elliptic flow in heavy-ion collisions. The present work discusses a new scenario that comprises a sequential transition from nuclear matter to deconfined quark matter, which could play an important role in asymmetric matter, particularly for the phenomenology of compact stars.

Chiral quark models of Nambu–Jona-Lasinio (NJL) type with dynamic chiral-symmetry breaking have the property that the symmetry is restored (and quarks are deconfined) separately for each flavor. When solving the gap and charge-neutrality equations self-consistently, the chiral-symmetry restoration for a given flavor occurs when the chemical potential of that flavor reaches a critical value that is approximately equal to the dynamically generated quark mass, $\mu_f = \mu_{c,f} \approx m_f$, where $f = u, d, s$. In asymmetric matter, the quark chemical potentials are different. Consequently, the NJL model behavior suggests that the critical density of deconfinement is flavor dependent (see Fig. 1). In this scenario, the down-quark flavor is the first to drip out of nucleons when the density increases, followed by the up quark flavor and eventually also by strange quarks. This behavior is absent in simple and commonly applied bag-model EOSs, because they are essentially flavor blind.

Under the β-equilibrium condition in compact stars, the chemical potentials of quarks and electrons are related by $\mu_d = \mu_s$ and $\mu_d = \mu_u + \mu_e$. The mass difference between the strange and the light quark flavors, $m_s \gg m_u, m_d$, has two consequences: (1) the down and strange quark densities are different, so charge neutrality requires a finite electron density and, consequently, (2) $\mu_d > \mu_u$. When the baryon chemical potential is increased, the d-quark chemical potential is therefore the first to reach the critical value, $\mu_{c,d} \approx m_d$, where the chiral symmetry becomes (approximately) restored in a first-order transition and deconfined d-quarks appear. Due to the finite value of μ_e, the u-quark chemical potential is still below $\mu_{c,u} \approx m_u$, whereas the s-quark density is zero due to the high s-quark mass. A *single-flavor* d-quark phase, therefore, forms in coexistence with the positively charged nuclear-matter medium.

Why has this interesting scenario gone unnoticed? One reason is that bag models, which are commonly used to describe quark matter in compact-star interiors, cannot address sequential deconfinement. Another reason is that the single-flavor d-quark phase is negatively charged and cannot be neutralized with leptons. It was therefore disregarded in dynamic approaches such as NJL models, which in practice are used to describe the deconfined and "pure" quark matter phase only. In the following we discuss the single-flavor phase for the first time, under the natural assumption that the neutralizing background is nuclear matter. Because nucleons are bound states of quarks, a mixed phase of nucleons and free d-quarks could naturally arise when nucleonic bound states dissociate (Mott effect).

II. PHASE TRANSITION TO QUARK MATTER: NUCLEON DISSOCIATION

The task to develop a unified description of the phase transition from nuclear matter to quark matter on the quark level, as a dissociation of three-quark bound states into their constituents in the spirit of a Mott transition, has not yet been solved. Only some aspects have been studied within a nonrelativistic potential model [15,16] and within the NJL model [17]. Here we consider a chemical equilibrium reaction of the form $n + n \leftrightarrow p + 3d$, which results in a mixed phase of nucleons and down quarks once the d-quark chemical potential exceeds the critical value. This scenario is analogous to the dissociation of nuclear clusters in the crust of neutron stars (neutron drip line) and the effect may therefore be called the *d-quark drip line*. The quark and nucleon components are approximated as subphases, which are described by separate models.

For the nuclear matter subphase we use two alternatives: (1) the DBHF approach [18–22] with the relativistic Bonn A potential, where the nucleon self-energies are based on a T-matrix obtained from the Bethe-Salpeter equation in the ladder approximation, and (2) the EOS by Shen *et al.* [23], which is based on relativistic mean-field theory and includes the contribution of heavy nuclei, described within the Thomas-Fermi approximation. Despite its drawbacks, this EOS is instructive because it is available for a large enough range of densities, temperatures, and isospin asymmetries that it qualifies for applications in studies of supernova collapse and protoneutron star evolution. Only very recently could significant progress be made, for example, in generalizing the nuclear-statistical equilibrium approach [24] and in implementing a quantum-statistical description for cluster formation and dissociation (the Mott effect) [25]. The quark-matter phase is described within a three-flavor NJL-type model, which includes diquark pairing channels [13,26–28]. This approach is justified because the $\mu > 0$ domain of the quantum chromodynamics (QCD) phase diagram is rather poorly understood. A more fundamental approach, like solving the in-medium QCD Schwinger-Dyson equations in a concrete QCD model [29–31], is demanding and, therefore, beyond scope of this work. The path-integral representation of the NJL partition function is given by

$$Z(T,\hat{\mu}) = \int \mathcal{D}\bar{q}\mathcal{D}q \exp\left\{\int_0^\beta d\tau \times \int d^3x[\bar{q}(i\partial\!\!\!/ - \hat{m} + \hat{\mu}\gamma^0)q + \mathcal{L}_{\rm int}]\right\}, \quad (1)$$

$$\mathcal{L}_{\rm int} = G_S\left\{\sum_{a=0}^{8}[(\bar{q}\tau_a q)^2 + (\bar{q}i\gamma_5\tau_a q)^2] + \eta_{D0}\sum_{A=2,5,7} j^\dagger_{D0,A}j_{D0,A} + \eta_{D1}j^\dagger_{D1}j_{D1}\right\}, \quad (2)$$

where $\hat{\mu} = \frac{1}{3}\mu_B + \mathrm{diag}_f(\frac{2}{3}, -\frac{1}{3}, -\frac{1}{3})\mu_Q + \lambda_3\mu_3 + \lambda_8\mu_8$ is the diagonal quark chemical potential matrix and $\hat{m} = \mathrm{diag}_f(m_u, m_d, m_s)$ is the current-quark mass matrix. For $a = 0$, $\tau_0 = (2/3)^{1/2}\mathbf{1}_f$; otherwise τ_a and λ_a are Gell-Mann matrices acting in, respectively, flavor and color spaces. $C = i\gamma^2\gamma^0$ is the charge conjugation operator and $\bar{q} = q^\dagger\gamma^0$. The scalar quark-antiquark current-current interaction is given explicitly and has coupling strength G_S. The three-momentum cutoff, Λ, is fixed by low-energy QCD phenomenology (see Table I of Ref. [32]). The spin-0 and spin-1 diquark currents are $j_{D0,A} = q^T iC\gamma_5\tau_A\lambda_A q$ and $j_{D1} = q^T iC(\gamma_1\lambda_7 + \gamma_2\lambda_5 + \gamma_3\lambda_2)q$. Although the relative coupling strengths η_{D0} and η_{D1} are essentially free parameters, we restrict the discussion to the Fierz values, $\eta_{D0} = 3/4$ and $\eta_{D1} = 3/8$ (see Ref. [33]). Color superconducting phases in QCD with one flavor were first discussed in Refs. [34–36], where it was also pointed out

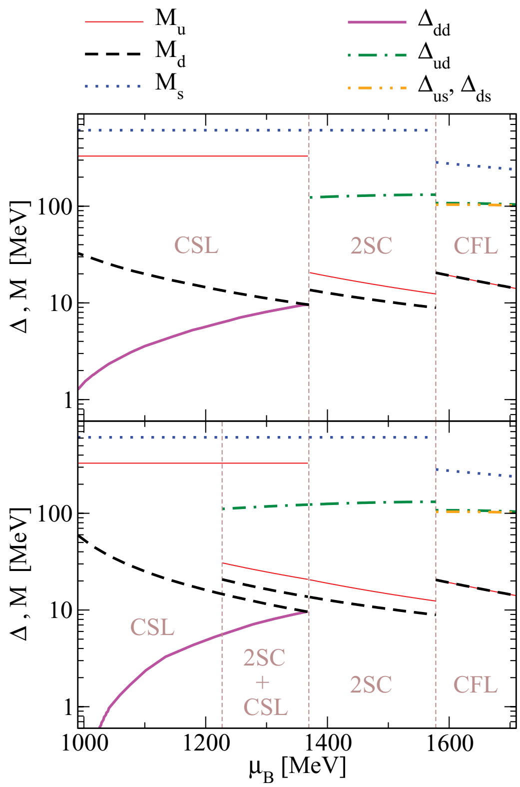

FIG. 2. (Color online) Solution of the NJL gap equations for isospin-asymmetric charge-neutral matter. The upper (lower) panel corresponds to the hybrid EOS based on the DBHF (Shen) nuclear EOS. The asymmetry at a given value of the baryon chemical potential, μ_B, is different in the two cases because the charge density of nuclear matter depends on the model used.

that the gap is of order 1 MeV in the spin-1 color-spin-locking (CSL) phase. This feature of the CSL phase is robust. See Ref. [37] for an analysis of our isotropic ansatz for the spin-1 diquark current, Ref. [38] for its generalization to the nonlocal case, and Ref. [39] for a self-consistent Dyson-Schwinger approach.

The gaps and the renormalized masses are determined by minimization of the mean-field thermodynamic potential under the constraints of charge neutrality and β equilibrium (see Fig. 2; for further details, see Refs. [13,26–28]). Different phases are characterized by different values of the order parameters (masses, gaps, etc.) and correspond to different local minima of the thermodynamic potential. For a particular choice of the baryon chemical potential there may be several local minima of the thermodynamic potential. The physical solution is that with the lowest free energy or, equivalently, the highest pressure. Figure 3 shows the pressure of various phase constructions. Because we use separate models for the confined and deconfined states of quarks, the dissociation of nucleons does not appear automatically within the model. Instead, for a given value of the baryon number chemical potential, three different phase constructions are considered: (1) the homogenous charge-neutral and β-equilibrated nuclear matter phase; (2) the homogenous charge-neutral and β-equilibrated quark matter phase; and

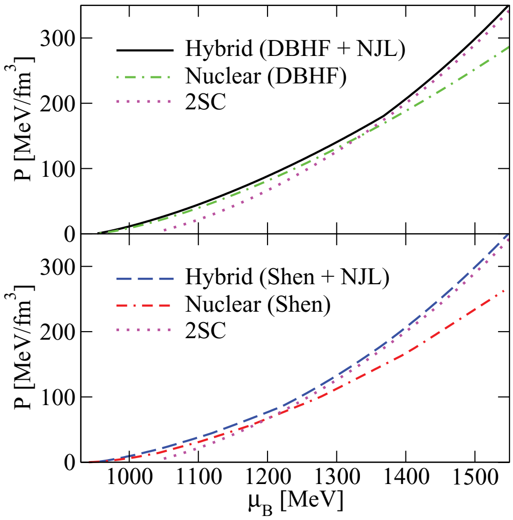

FIG. 3. (Color online) Pressure of matter in β equilibrium for different nuclear-matter models (DBHF, upper panel; Shen, lower panel) and phase transition constructions with color superconducting quark matter (NJL model). Here DBHF (or Shen) + NJL refers to the mixed phase of nuclear matter and quark matter. At low chemical potential the quark-matter phase component is the negatively charged d-CSL phase, which lowers the asymmetry of the system and thereby gives a higher pressure for the mixed phase. At higher chemical potentials (1370 MeV for DBHF + NJL and 1230 for Shen + NJL), there is a transition in the quark sector to a two-flavor color superconducting (2SC) phase component.

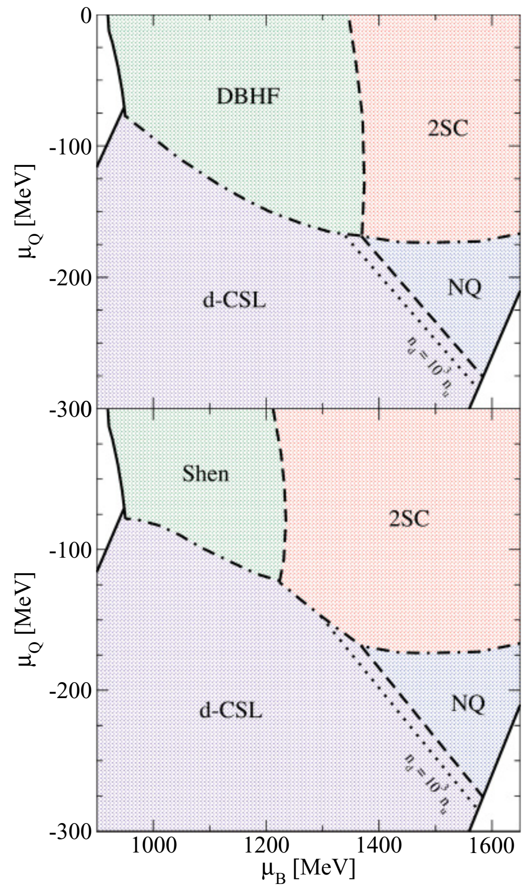

FIG. 4. (Color online) Phase diagrams in the plane of baryon and charge chemical potential. The dash-dotted line denotes the border between oppositely charged phases. The nuclear-matter EOSs are DBHF (upper panel) and Shen *et al.* (lower panel). Only one- and two-flavor solutions are displayed, because three-flavor matter is charge neutral for $\mu_Q \sim 0$. The transition to the nearly symmetric three-flavor CFL phase occurs at $\mu_B \sim 1600$ MeV (see Fig. 2). At low densities the mixture of nuclear matter with one-flavor quark matter (d-CSL) is favored. At higher densities, beyond the up-quark threshold, a mixture with two-flavor quark matter is favored. The two-flavor phases considered here are the normal quark-matter (NQ) and the superconducting (2SC) phases.

(3) a charge-neutral equilibrium mixture of nuclear matter and quark matter, or two different quark-matter phases. For the models considered here we find that the asymmetry in three-flavor (color-flavor-locked [CFL]) quark matter is so small that it makes little sense to consider inhomogenous phase constructions. For mixtures of nuclear matter with one-flavor (d-CSL) or two-flavor quark matter in the 2SC or normal quark-matter (NQ) phase, however, the asymmetry is significantly lower and the pressure is higher when compared with the homogeneous phases.

Figure 4 plot the thermodynamically favored phase in the plane of baryon and charge chemical potentials. The hybrid EOS corresponds to the dash-dotted lines in Fig. 4 (i.e., the borders between positively and negatively charged phases), and they are constructed such that the corresponding mixture of nuclear matter, quark matter, and leptons is charge neutral. At low densities a mixture with one-flavor quark matter is favored. At higher density, beyond the up-quark threshold, a mixture of one-flavor and two-flavor quark matter is favored. The strange flavor occurs at still higher densities. Note that the 2SC phase cannot persist at high $|\mu_Q|$ because the large difference in the Fermi levels of u- and d-quarks prevents their pairing and the two-flavor quark matter is therefore in the normal phase (NQ).

Using the hybrid EOS, we calculate compact-star sequences by solving the Oppenheimer-Volkoff equations for hydrostatic equilibrium. The hybrid-star sequences fulfill all modern constraints on the mass-radius relationship (see Fig. 5). For the DBHF hybrid EOS, all stars with DBHF + CSL matter in the core are stable equilibrium solutions, whereas the appearance of u-quarks and the associated formation of a 2SC subphase renders the sequence unstable. The situation is somewhat different for the Shen hybrid EOS, because in addition to Shen + CSL stars there are stable solutions with 2SC + CSL matter in the core. In both cases, configurations with strange quarks in the core are unstable. The hybrid star sequences "masquerade" as neutron stars [40] because the mechanical properties are similar to those of nuclear-matter stars and the transition from nuclear matter to the mixed phase is associated with a relatively small discontinuity in the density. Unmasking neutron star interiors may therefore require observables based on transport properties, which could be strongly modified in the presence of color superconductivity. It has been suggested that such tests of the structure of matter at high density be based on an analysis of the cooling behavior [41–44] or the

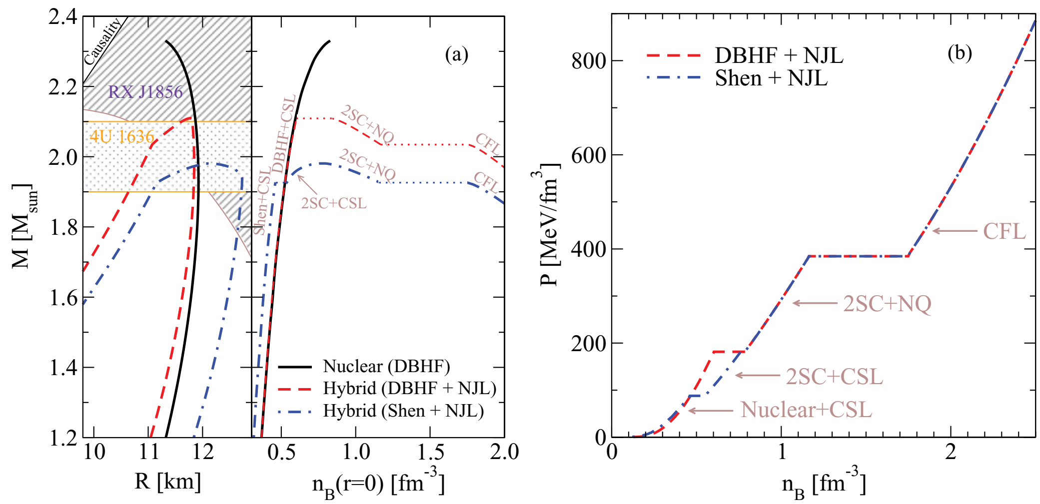

FIG. 5. (Color online) (a) Compact star sequences and (b) hybrid EOSs. The phase structure at the center, $r = 0$, changes with increasing density, as indicated in the figures. Constraints on the compact-star mass come from 4U 1636 [2] and on the mass-radius relation from RX J1856 [4].

stability of rapidly spinning stars against r-modes [45,46]. It has turned out that these phenomena are sensitive to the details of color superconductivity in quark matter.

The down-quark chemical potential exceeds that of up quarks in asymmetric nuclear matter and, as illustrated earlier, this could lead to sequential deconfinement of the quark flavors. We have checked that the breaking of the $U(1)_A$ symmetry with a six-point 't Hooft interaction does not rule out the single-flavor d-CSL solution, but the phase border is shifted to higher $|\mu_Q|$ (i.e., more asymmetry is needed to realize the d-CSL phase in that case). A potential consequence is that the asymmetry of charge-neutral nuclear matter is less than that of d-CSL matter with broken $U(1)_A$ symmetry, and that the nuclear phase, therefore, is thermodynamically favored. Because the origin of the $U(1)_A$ anomaly is unknown (see, e.g., the discussion in Ref. [13]), and the critical asymmetry depends on the parametrization of the NJL model and on the nuclear-matter model used, a definite answer for whether the d-CSL phase is realized is a matter of further investigation. Other effects of the inhomogenous phase mixture (e.g., Coulomb interactions and surface tension) should also be considered in a future detailed investigation. Irrespective of these unsettled issues, it is clear that the free energy of the d-CSL phase decreases with increasing asymmetry, in direct contrast to the behavior of traditional phases such as the nuclear-matter phase. In the following we discuss another interesting feature of the d-CSL phase, which could have important consequences for the phenomenology of compact stars.

III. BULK VISCOSITY AND URCA EMISSIVITY OF THE SINGLE-FLAVOR CSL PHASE

Rotating compact stars would be unstable against r-modes in the absence of viscosity [47,48]. Constraints on the composition of compact-star interiors can therefore be obtained from observations of millisecond pulsars [45,46]. In such investigations the bulk viscosity is a key quantity and constraints on matter phases in neutron-star interiors can be based on its value. Here we consider some relevant aspects of the bulk viscosity for color superconducting phases, starting with the 2SC phase and following the approach described in Ref. [49]. Note that the 2SC phase considered in Ref. [50] is a three-flavor phase, for which the nonleptonic process $u + d \leftrightarrow u + s$ is the dominant contribution. This process is not relevant for the 2SC phase discussed here, where the strange-quark Fermi sea is not occupied.

The temperature-dependent bulk viscosity for the $2\mathrm{SC} + \mathrm{CSL}$ phase was calculated self-consistently in Ref. [51] and is based on the flavor-changing weak processes of electron capture and β-decay:

$$u + e^- \to d + \nu_e, \quad d \to u + e^- + \bar{\nu}_e. \tag{3}$$

It has been shown that the bulk viscosity is related to the direct Urca emissivity, which for normal quark matter was first calculated by Iwamoto [52] and can be expressed as

$$\varepsilon_0 \simeq \frac{914\pi}{1680} G_F^2 \mu_e \mu_u \mu_d T^6 \theta_{ue}^2, \tag{4}$$

where G_F is the weak coupling constant and θ_{ue} is the angle between the up-quark and electron momenta, which is obtained from momentum conservation in the matrix element (see Fig. 6). The triangle of momentum conservation holds for

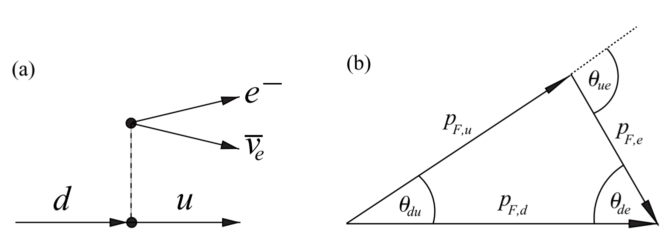

FIG. 6. Direct Urca process in quark matter (a) and triangle of momentum conservation for it (b).

the late cooling stage, when the temperature is below 1 MeV and neutrinos are untrapped. Trigonometric relations are used to find an analytical expression for momentum conservation. To lowest order in θ_{de} the result is

$$p_{F,d} - p_{F,u} - p_{F,e} \simeq -\tfrac{1}{2} p_{F,e}\theta_{de}^2. \tag{5}$$

For small angles, $\theta_{de} \simeq \theta_{ue}$, so it is possible to obtain an expression for the matrix element of the direct Urca process. Following Iwamoto [52], one has to account either for quark-quark interactions to lowest order in the strong coupling constant, α_s [Eq. (6)], or the effect of finite masses [Eq. (7)]:

$$\mu_i = p_{F,i}\left(1 + \frac{2}{3\pi}\alpha_s\right), \quad i = u, d, \tag{6}$$

$$\mu_i \simeq p_{F,i}\left[1 + \frac{1}{2}\left(\frac{m_i}{p_{F,i}}\right)^2\right], \quad i = u, d, e. \tag{7}$$

From Eqs. (5)–(7) and the β-equilibrium condition, $\mu_d = \mu_u + \mu_e$, the angle θ_{de} that determines the emissivity [Eq. (4)] is obtained [Eq. (5)]:

$$\theta_{de}^2 \simeq \begin{cases} \dfrac{4}{3\pi}\alpha_s \\ \dfrac{m_d^2}{p_{F,e}p_{F,d}}\left[1 - \left(\dfrac{m_u}{m_d}\right)^2\left(\dfrac{p_{F,d}}{p_{F,u}}\right) - \left(\dfrac{m_e}{m_d}\right)^2\left(\dfrac{p_{F,d}}{p_{F,e}}\right)\right] \end{cases}. \tag{8}$$

If interactions and masses are neglected, or the Fermi sea of one species is closed as in the single-flavor CSL phase, it follows that the triangle of momentum conservation in Fig. 6 degenerates to a line or cannot be closed. In that case, the matrix element vanishes with the consequence that the direct Urca process does not occur, and also the bulk viscosity is zero. However, in the mixed nuclear + CSL phase, there could be important friction and pair-breaking/formation processes, which we have not yet studied in detail. This could be an interesting issue for further investigation because of the large difference in the masses of baryons and deconfined quarks.

IV. MECHANISM FOR DEEP CRUSTAL HEATING

Superbursts are rare, puzzling phenomena observed as extremely long (4–14 h) and energetic ($\sim 10^{42}$ erg) type-I x-ray bursts from LMXBs. They take place if the accreted hydrogen and helium at the surface burns in an unstable manner, which is the normal case [53]. As suggested in Ref. [54], superbursts could originate from accreting strange stars with a thin crust and a core of three-flavor quark matter in the CFL phase. The suppression of the neutrino emissivity and heat conductivity in the CFL phase [55–57], caused by pairing gaps that affect all flavors, is of particular importance in this superburst scenario. Following Cumming *et al.* [58], the underlying mechanism is unstable thermonuclear burning of carbon in the crust, at column depths of about $(0.5–3) \times 10^{12}$ g cm^{-2}. Carbon is a remnant of accreted hydrogen and helium at the surface. Observed superburst light curves suggest that the burning takes place at a depth where the crust reaches temperatures of about 6×10^8 K and column depths of about 10^{12} g cm^{-2}. Such high temperatures in the crust at a certain depth are caused by deep crustal heating [59–61]. The important ingredients for the strange-star model of superbursts [54] are a thin baryonic crust of thickness 100–400 m, an energy release of 1–100 MeV per accreted nucleon by conversion into strange matter, a suppression of the fast direct Urca neutrino emissivity to the order of 10^{21} erg cm^{-3} s^{-1}, and a thermal conductivity, κ, of quark matter in the range 10^{19}–10^{22} erg cm^{-1} s^{-1} K^{-1}.

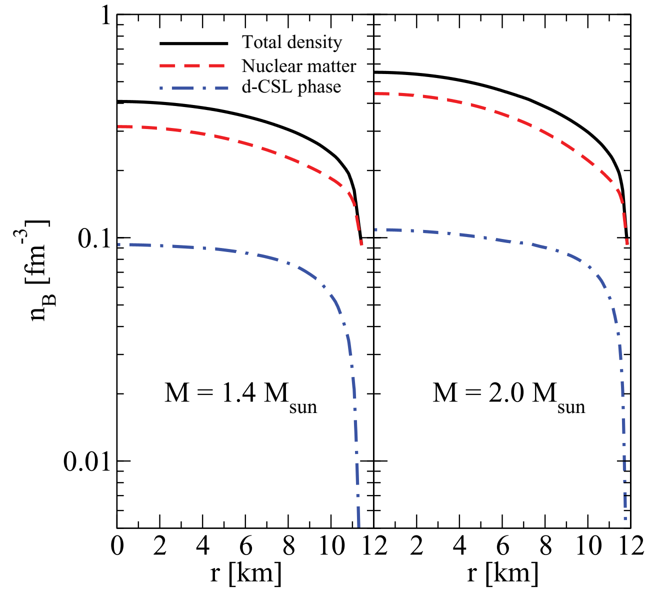

FIG. 7. (Color online) Density profiles of two stars with masses 1.4 and 2.0 M$_\odot$. In the model adopted here, the mixed phase of d-CSL quark matter with nuclear matter extends up to the crust-core boundary.

Figure 7 shows that d-CSL quark matter (in the mixed phase with nuclear matter) extends up to the crust-core boundary, as strange quark matter does in the case of strange stars. One of the main arguments for strange matter in the context of a superburst mechanism is the fact that superconducting phases, such as the CFL phase, can suppress fast neutrino emission processes of all quark flavors and are able to fulfill the fusion ignition condition. This case is also true for the single-flavor CSL phase. As shown earlier, the fast direct Urca process is not possible at all in this phase, whereas slow neutrino cooling processes like bremsstrahlung of electrons and d-quarks exist.

We want to estimate the order of magnitude of the energy release ΔE due to partial conversion of ordinary nuclear matter to DBHF + CSL hybrid matter at the crust-core boundary. As we apply the Gibbs construction of a phase transition, a density-dependent volume fraction of the d-quark admixture in the nuclear + CSL phase results, varying from zero to unity. The caveat of this construction is that all thermodynamic quantities at the onset of the phase transition vary continuously. However, in reality an infinitesimally small fraction of the d-quark subphase would imply that large residual color forces between d-quarks should occur. Therefore, a solution of the phase admixture problem with a finite jump of the d-quark admixture at the onset of the tranistion should be energetically favored. At the present stage of our work we cannot quantify this statement because of the absence of confining forces between color charges in our quark-matter model. An estimate which we suggest here is to determine the fraction χ of d-CSL matter at the d-quark drip line in the vicinity of the crust-core boundary. Then one would multiply the change in energy per baryon due to the process $n \rightarrow ddu$ with χ as an estimate for the probability of this process to occur per accreted nucleon. A rough estimate (see Fig. 7) gives $0.001 \leqslant \chi \leqslant 0.01$, which for a jump of the d-quark mass gap by 300 MeV (see Fig. 2) at the chiral transition (d-quark drip line) results in $0.6 \leqslant \Delta E \leqslant 6$ MeV. This meets well the estimated range $\Delta E \sim 1 - 100$ MeV [54,58] and could thus, in principle, explain burst ignition at appropriate depths for a suitable value of κ.

Therefore, a strange-matter core is not necessarily required to resolve the superburst puzzle, because a hybrid-star model with quark matter in the d-CSL phase could have similar properties. Stejner *et al.* [53] show that deep crustal heating mechanisms at the crust-core boundary (e.g., conversion of baryonic matter to strange quark matter), which can fulfill the constraints of the superburst scenario, also provide a consistent explanation for the cooling of soft x-ray transients. Along the lines of this argument, we suggest that the d-quark drip effect, which leads to a mixture of nuclear matter with single-flavor quark matter in the CSL phase, can serve as a deep crustal heating mechanism. Superbursts and the cooling of x-ray transients are not only consistent with quark matter in compact stars but may qualify as a signature for its occurrence.

V. CONCLUSIONS

In this article, a new quark-nuclear hybrid model for compact star applications was suggested that fulfills modern constraints from observations of compact stars. Because of isospin asymmetry, down quarks may "drip out" from nucleons and form a single-flavor CSL phase that is mixed with nuclear matter already at the crust-core boundary in compact stars. The CSL phase has interesting cooling and transport properties that are in accordance with constraints from the thermal and rotational evolution of compact stars [51]. It remains to be investigated whether this new compact-star composition could lead to unambiguous observational consequences, and whether it is also thermodynamically favored when effects like Coulomb screening and surface tension are accounted for. We conjecture that the d-quark drip may serve as an effective deep crustal heating mechanism for the explanation of the puzzling superburst phenomenon and the cooling of x-ray transients.

ACKNOWLEDGMENTS

D.B. is supported in part by the Polish Ministry of Science and Higher Education under Grants No. N N 202 0953 33 and N N 202 2318 37 and by the Russian Fund for Basic Research under Grant No. 08-02-01003-a. T.K. is grateful for partial support from the Department of Energy, Office of Nuclear Physics, Contract No. DE-AC02-06CH11357. The work of F.S. was supported by the Belgian Fund for Scientific Research (FNRS). D.B. and F.S. acknowledge support from CompStar, a Research Networking Programme of the European Science Foundation.

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