Sequential deconfinement of quark flavors in neutron stars

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We suggest a scenario where the three light quark flavors are sequentially deconfined under increasing pressure in cold asymmetric nuclear matter as found, e.g., in neutron stars. The basis for our analysis is a chiral quark matter model of Nambu–Jona-Lasinio (NJL) type with diquark pairing in the spin-1 single flavor (CSL), spin-0 two flavor (2SC) and three flavor (CFL) channels. We find that nucleon dissociation sets in at about the saturation density, $n_0$, when the down-quark Fermi sea is populated (d-quark dripline) due to the flavor asymmetry induced by $\beta$-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. We construct two different hybrid equations of state (EoS) using the Dirac-Brueckner Hartree-Fock (DBHF) approach and the EoS by Shen et al. in the nuclear matter sector. The corresponding hybrid star sequences have maximum masses of, respectively, 2.1 and 2.0 $M_\odot$. Two- and three-flavor quark-matter phases exist only in gravitationally unstable hybrid star solutions in the DBHF case, while the Shen-based EoS 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 superbusts and cooling of X-ray transients.

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I. INTRODUCTION

The phenomenology of compact stars is intimately connected to the EoS 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 1 and references therein. A stiff EoS at high density is needed to explain the high compact-star masses and radii, for which there is growing evidence from recent observations. A mass of $M \sim 2.0 M_\odot$ has been reported for some low-mass X-ray binaries (LMXBs), e.g., 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 3. From observations of the bright isolated neutron star RX J1856.5-3754 (shorthand: RX J1856) in the optical and X-ray frequency range, a conservative lower limit of the apparent neutron star radius of $R_\infty = 16.5$ km is derived 4. This 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 5. Another example is EXO 0748-676, an LMXB for which the compact-star mass and radius have been constrained to $M \geq 2.10 \pm 0.28 M_\odot$ and $R \geq 13.8 \pm 0.18$ km 6 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 $z = 0.35$ observed in the X-ray burst spectra 7 has not been confirmed, despite numerous attempts. Further constraints on the masses and radii of compact stars have been reported (see, e.g., 8, 9), but they deserve a careful discussion which is beyond the scope of the present paper. While compact-star phenomenology apparently points towards a stiff EoS at high density, heavy-ion collision data for kaon production 10 and elliptic flow 11 set an upper limit on the stiffness of the EoS 12.

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 13, because quark deconfinement softens the EoS and lowers the maximum mass and corresponding radius. However, Alford et al. have demonstrated 14 with a few counter examples that quark matter cannot be excluded by this argument. In particular, for a recently developed
hybrid star EoS [12], based on the DBHF approach in the nuclear sector and a three-flavor chiral quark model [13], stable hybrid stars with masses ranging from 1.2 M⊙ to 2.1 M⊙ is obtained, in accordance with modern mass-radius constraints, see also [14]. In this model, a sufficiently low critical density for quark deconfinement has been achieved by a strong diquark coupling, while a repulsive vector meanfield 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. In the present work we discuss a new scenario that comprises a sequential transition from nuclear matter to deconfined quark matter, which could play an important role in asymmetric matter, in particular for the phenomenology of compact stars.

Chiral quark models of NJL type with dynamical 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, \( \mu_f = \mu_{e,f} \approx m_f \), 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 equations of state, because they are essentially flavor blind.

Under the \( \beta \)-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 increasing the baryochemical potential the d-quark chemical potential is therefore the first to reach the critical value, \( \mu_{e,d} \approx m_d \), where the chiral symmetry gets (approximately) restored in a first-order transition and deconfined d-quarks appear. Due to the finite value of \( \mu_e \) the u-quark chemical potential is still below \( \mu_{e,u} \approx m_u \), while the s-quark density is zero due to the high s-quark mass. A single-flavor d-quark phase therefore forms in co-existence with the positively charged nuclear-matter medium.

Why has this interesting scenario been left 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 dynamical approaches like NJL models which in practice are used mainly 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. Since 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 on 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 dripline) and the effect may therefore be called the d-quark dripline. We approximate the quark and nucleon components as subphases, which are described by separate models.

For the nuclear matter subphase we use two alternatives: 1) The DBHF approach [18, 19, 20, 21, 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. 2) The EoS by Shen et al. [23], which is based on relativistic mean-field theory and includes the contribution of heavy

FIG. 1: (Color online) Schematic picture 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 \( \mu \) is needed to form 2-flavor and 3-flavor quark matter phases.
nuclei, described within the Thomas-Fermi approximation. Despite its drawbacks this EoS is instructive since 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, significant progress could be made, e.g., in generalizing the nuclear statistical equilibrium approach [24] and in implementing a quantum statistical description for cluster formation and dissociation (Mott effect) [25]. The quark matter phase is described within a three-flavor NJL-type model, which includes diquark pairing channels [13, 26, 27, 28]. This approach is justified since the $\mu > 0$ domain of the 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, 30, 31] is demanding and is therefore beyond scope of this work. The path-integral representation of the NJL partition function is given by

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

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

where $\hat{m} = \frac{1}{6} \mu_B + \text{diag}(\frac{3}{2}, -\frac{3}{2}, -\frac{3}{2}) \mu_Q + \lambda_3 \mu_3 + \lambda_8 \mu_8$ is the diagonal quark chemical potential matrix and $\bar{\eta} = \text{diag}(\eta_u, \eta_d, \eta_s)$ is the current-quark mass matrix. For $a = 0$, $\tau_0 = (2/3)^{1/2} 1_f$, otherwise $\tau_a$ and $\lambda_a$ 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 3-momentum cutoff, $\Lambda$, is fixed by low-energy QCD phenomenology (see table I of [32]). The spin-0 and spin-1 diquark currents are $j_{D0, A} = q^T i C C \gamma_5 \tau_A A q$ and $j_{D1} = q^T i C (\gamma_1 \lambda_7 + \gamma_2 \lambda_5 + \gamma_3 \lambda_2) q$. While the relative coupling strengths $\eta_{D0}$ and $\eta_{D1}$ are essentially free parameters, we restrict the discussion to the Fierz values, $\eta_{D0} = 3/4$ and $\eta_{D1} = 3/8$, see [33]. Color superconducting phases in QCD with one flavor were first discussed in Refs. [34, 35, 36], where it was also pointed out 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 [37] for an analysis of our isotropic ansatz for the spin-1 diquark current, [38] for its generalization to the non-local case, and [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 $\beta$-equilibrium, see Fig. 2. For further details, see [13, 26, 27, 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 lowest free energy, or, equivalently, highest pressure. In Fig. 3 we show the pressure of various phase constructions. Since 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 $\beta$-equilibrated nuclear matter phase. 2) The homogenous charge-neutral and $\beta$-equilibrated quark matter phase. 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 quark matter (CFL) is so small that it makes little sense to consider inhomogeneous 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 higher when compared with the homogeneous phases. In Fig. 4 we 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 $|\mu_Q|$ since 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).
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, $\mu_B$, is different in the two cases, because the charge density of nuclear matter depends on the model used.

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, while 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 might therefore require observables based on transport properties, which could be strongly modified in presence of color superconductivity. It has been suggested to base such tests of the structure of matter at high density on an analysis of the cooling behavior [41, 42, 43, 44] or the 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 we have illustrated above, 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 realise the d-CSL phase in that case. A potential consequence of this 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. As the origin of the $U(1)_A$ anomaly is unknown, see, e.g., the discussion in [13], and the critical asymmetry depends on the parametrization of the NJL model and on the nuclear matter model used, a definite answer whether the d-CSL phase is realised is a matter of further investigation. Other effects of the inhomogenous phase mixture, e.g., Coulomb interactions and surface tension

FIG. 3: (Color online) Pressure of matter in beta-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 of 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 2SC phase component.
for the phenomenology of compact stars.

d-CSL phase, which could have important consequences of traditional phases such as the nuclear matter phase. In increasing asymmetry, in direct contrast to the behaviour cause three-flavor matter is charge-neutral for the free energy of the d-CSL phase decreases with in-
tion. Irrespective of these unsettled issues it is clear that only one-flavor and two-flavor solutions are displayed, be-
clear matter EoS is DBHF (upper) and Shen

\[ \mu \] note the border between oppositely charged phases. The nu-
should also be considered in a future detailed investiga-
tion. Irrespective of these unsettled issues it is clear that the free energy of the d-CSL phase decreases with in-
creasing asymmetry, in direct contrast to the behaviour 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.

FIG. 4: (Color online) Phase diagrams in the plane of baryon and charge chemical potential. The dash-dotted line denote the border between oppositely charged phases. The nu-
clear matter EoS is DBHF (upper) and Shen et al. (lower). Only one-flavor and two-flavor solutions are displayed, be-
cause three-flavor matter is charge-neutral for \( \mu \) \( \sim \) 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) phase.

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 \([49, 50]\). 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 [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 2SC+CSL phase has been calculated self-consistently in [51] and is based on the flavor-changing weak processes of electron capture and \( \beta \) decay

\[
\mu + e^- \rightarrow d + \nu_e , \quad d \rightarrow u + e^- + \bar{\nu}_e .
\]

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

\[
\varepsilon_0 \approx \frac{914}{1680} \frac{\pi}{G_F^2} \mu_e \mu_u \mu_d T^6 \theta_{ue}^2 .
\]

Here \( G_F \) is the weak coupling constant and \( \theta_{ue} \) is the angle between the up-quark and electron momenta, which is obtained from momentum conservation in the matrix element, see Fig. 4. The triangle of momentum conservation holds for 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 \( \theta_{de} \), the result is

\[
p_{F,u} - p_{F,u} - p_{F,e} \approx - \frac{1}{2} p_{F,e} \theta_{de}^2 .
\]

For small angles \( \theta_{de} \approx \theta_{ue} \), 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, \( \alpha_s \), \([4]\) or the effect of finite masses \( e \):

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

\[
\mu_i \approx p_{F,i} \left[ 1 + \frac{1}{2} \left( \frac{n_i}{p_{F,i}} \right)^2 \right] , \quad i = u, d, e
\]

From \([4]\)-\( 7 \) and the \( \beta \)-equilibrium condition, \( \mu_d = \mu_u + \mu_e \), the angle \( \theta_{de} \) that determines the emissivity \([4]\) is
FIG. 5: (Color online) Compact star sequences (a) and hybrid equations of state (b). 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 \[3\].

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

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 due to the large difference in the masses of baryons and deconfined quarks.

IV. MECHANISM FOR DEEP CRUSTAL HEATING

Superbursts are rare, puzzling phenomena observed as a extremely long (4-14 hours) and energetic \((\sim 10^{42}\text{erg})\) 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 \[54\] superbursts could originate from accreting strange stars with a thin crust and a core of three-flavor quark matter in the color-flavor-locked (CFL) phase. The suppression of the neutrino emissivity and heat conductivity in the CFL phase \[55, 56, 57\], 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 \((0.5 - 3) \times 10^{12}\ \text{g cm}^{-2}\). 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 \(6 \times 10^8\ \text{K}\) and column depths of about \(10^{12}\ \text{g cm}^{-2}\). Such high temperatures in the crust at a certain depth
are caused by deep crustal heating [50, 51, 52]. The important ingredients for the strange-star model of superbursts [53] are a thin baryonic crust of thickness 100 to 400 m, an energy release of 1 to 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, $\kappa$, of quark matter in the range $10^{19} - 10^{22}$ erg cm$^{-1}$ s$^{-1}$ K$^{-1}$.

In Fig. 7 we show 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, like the CFL phase, can suppress fast neutrino emission processes of all quark flavors and are able to fulfill the fusion ignition condition. This is the case also for the single-flavor CSL phase. As we have shown above the fast direct URCA process is not possible at all in this phase, while slow neutrino cooling processes like bremsstrahlung of electrons and d-quarks exist.

We want to estimate the order of magnitude of the energy release $\Delta E$ 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 thermodynamical 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 transition should be energetically favored. At the present stage of our work we cannot quantify this statement due to the absence of confining forces between color charges in our quark matter model. An estimate which we would suggest here is to determine the fraction $\chi$ of dCSL matter at the d-quark dripline in the vicinity of the crust-core boundary. Then one multiplies the change in energy per baryon due to the process $n \to dddu$ with $\chi$ as an estimate for the probability of this process to occur per accreted nucleon. A rough estimate (see Fig. 7) gives $0.001 \leq \chi \leq 0.01$ which for a jump of the d-quark mass gap by 300 MeV (see Fig. 2) at the chiral transition (d-quark dripline) results in $0.6 \leq \Delta E [\text{MeV}] \leq 6$ 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 $\kappa$.

Therefore, a strange matter core is not necessarily needed 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 provide a consistent explanation also 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 paper we suggest a new quark-nuclear hybrid model for compact star applications that fulfills modern constraints from observations of compact stars. Due to isospin asymmetry, down-quarks may “drip out” from nucleons and form a single-flavor color superconducting (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 thermodynamically favored also 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.
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