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γ-ray bursts and the QCD phase diagram

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The quantum chromodynamics phase diagram shows the phase transitions that can take place in matter at different temperatures and densities. In this work we discuss the possibility that γ-ray bursts might result from a phase change in the interior of a neutron star and calculate the energy released in the conversion of a metastable star into a stable star. We consider several different initial and several different final configurations. Initial metastable stars are taken as hadronic, hybrid, and quark stars with unpaired quarks; possible stable stars are hybrid and quark stars, taken both with unpaired and paired phases to study the deconfinement phase transition and normal quark matter–superconducting quark matter phase transition within a large number of relativistic models used to describe compact stars. The models used for the hadron matter are the nonlinear Walecka model and the quark-meson coupling model with and without hyperons. For the quark matter we have used the MIT bag model, the bag model with paired quarks to which we refer as the color-flavor-locked phase model, and the Nambu–Jona-Lasinio model. Within this mechanism we obtain energies of the order of $10^{50}-10^{53}$ erg, accounting for both long and short γ-ray bursts.

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I. INTRODUCTION

The possibility that a quark-gluon plasma (QGP) could be formed in heavy-ion collisions [1–3] arose when quantum chromodynamics (QCD) at finite temperature and high densities became a topic of increasing interest owing to the discovery of asymptotic freedom [4] about 30 years ago. In cosmology, the relevant conditions for QGP formation occur $10\,\mu$s after the Big Bang; nuclear matter first appears after about 1 ms. In laboratory searches for QGP, in large colliders around the world (RHIC/BNL, ALICE/CERN, GSI, etc.), experimentalists are trying to do the opposite: to convert hadronic matter at sufficiently high temperatures into QGP. The QGP diagram is shown in Fig. 1. In this case the conserved quantities during the phase transition are the strangeness, the isospin, the electric charge, and the baryonic number. The Gibbs conditions for phase coexistence require that the temperature, the pressure, and the chemical potentials in both phases be the same. There are three chemical potentials that determine the equation of state (EOS) of matter: the ones associated with the three lightest quarks or other three baryon chemical potentials, namely, for the proton, the neutron, and the lambda hyperon. All baryon chemical potentials can be written in terms of these three chemical potentials. The phase transition from hadronic matter to a deconfined quark matter is expected to take place at around a few times the nuclear saturation density (at zero temperature), and this could occur in the interiors of neutron stars [5,6]. Suggestions in the literature on the interior composition of neutron stars can be classified into pure hadronic matter (hadronic stars) with or without hyperons [5,7,8]; a mixed phase of hadrons and quarks (hybrid stars) [5,9–12]; a mixed phase of hadrons and pion or kaon condensates [13–15] (also hybrid stars); and deconfined quarks [14,16,17] (strange or quark stars). According to Ref. [18], one possibility does not exclude the others since the analysis of different astrophysical phenomena associated with compact X-ray sources show that there are stars with radii in the range of 10–12 km and also stars with smaller radii, around 6–9 km. At temperatures of the order of $\simeq 0-40$ MeV, which are the relevant temperatures in compact stars [10], there are two possibilities for phase transitions, as can be seen from the QGP diagram in Fig. 1. As the density increases, hadron matter first converts into QGP or into either a crystalline quark matter or a two-flavor superconducting phase, and subsequently into a color-flavor-locked superconducting phase. A possible transition from a hadron phase directly to a color-flavor-locked phase, which describes superconducting matter, has already been investigated in the context of hybrid stars [19,20].

In this paper we discuss the possibility that γ-ray bursts (GRBs) might be a manifestation of a phase change in the interior of a neutron star. There are two classes of GRBs [21]: The short bursts, which are generally hard (SGRBs) and the long bursts, generally soft (LGRBs). They are distinguished mainly by their duration and the energies released [22]. The long-soft γ-ray bursts occur in star-forming galaxies at cosmological distances with a red shift of the order of $z = 1$ and they produce X-ray and optical afterglows. They are believed to be associated with the explosion of massive stars [23]. For a long time the energies associated with SGRBs, as well as their sources, were unknown. Recently, accurate results have been obtained for two SGRBs: GRB 050509B, located at $z = 0.225$ [24], and GRB 050709, located at $z = 0.16$ [25]. These results seem to indicate that the SGRBs are caused by stellar merging into final compact binary systems. The total isotropic energy released in the first few hundred seconds is of the order of 10^{50} erg, which is two or three orders of magnitude

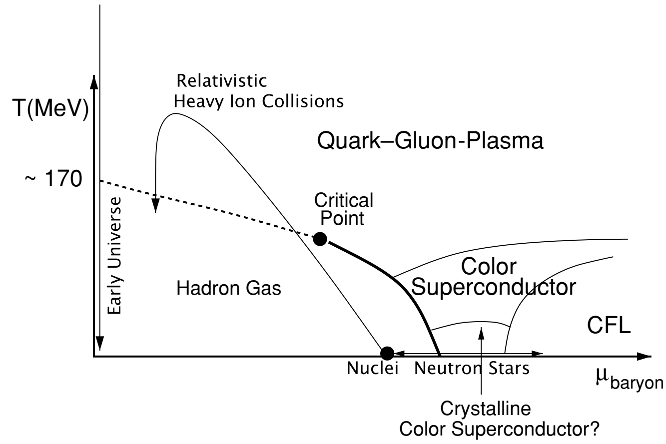

FIG. 1. QGP phase diagram.

smaller than seen in LGRBs. According to Ref. [26], LGRBs are associated with supernovae of Type Ic, which present two distinct kinds of bursts: Some generate soft and weak bursts with no evidence for ultarelativistic jets and some generate strong GRBs with ultrarelativistic jets. These bursts can be explained by the conversion of a hadronic star into a quark star, as proposed in Ref. [27], where an energy release of 10^{53} erg was calculated. This value is one order of magnitude larger than the result of previous calculations [28–30], which could not account for the energy needed to explain the most energetic GRBs.

We make six points relevant to phase transitions in neutrons stars. (1) For a phase transition in a stellar interior, strangeness conservation is not required, charge conservation is restricted to the neutral case, and β equilibrium has to be enforced. The Gibbs conditions remain the same, but only two chemical potentials determine the EOS, the neutron and the electron ones. (2) Based on the Bodmer-Witten hypothesis [28,31–33] that "strange matter" is the true ground state of all matter, it was shown in Ref. [29] that once there is a seed of strange quark matter inside a neutron star, the whole star converts to a strange quark star. The quark matter front would absorb neutrons, protons, and hyperons, liberating their quark constituents. This conversion should take from 1 ms to 1 s [34], which agrees with the typical duration of a SGRB. (3) In Ref. [27] the energy released in the conversion of neutron (hadronic) stars to strange stars was calculated and shown to be of the order of $(1-4) \times 10^{53}$ erg. For the hadronic stars two possible EOSs were considered: one with protons and neutrons only and another one with the eight lightest baryons. For the strange quark stars, the usual MIT bag model was used to describe quark matter and different choices for the bag parameter were used. (4) In Refs. [18,35] it was found that above a threshold value of the gravitational mass, a pure hadronic star is metastable to a conversion into a hybrid or quark star. (5) In Ref. [36] the energy released in the conversion of a hadronic star into a hybrid or quark star with superconduction phase was estimated to be of order 10^{53} erg. (6) In Refs. [35] and [18] the finite-size effects between the deconfined and confined phases are taken into account and the mean life of the metastable star is calculated. In the first paper β equilibrium is considered during the phase transition; in the second case it is supposed that only strong forces act during the phase transition.

Here we generalize the approach of Ref. [27], in particular, considering the stellar conversions in which baryon number is conserved. We are not concerned with the lifetime of the star but are interested in the energy liberated during the stellar conversion. We calculate the energy released in various possible conversions from a metastable to a stable star and check whether they can account for LGRBs or SGRBs. The authors of Ref. [27] considered conversions from a metastable star into a stable star, with the stable configuration assumed to be a hybrid or quark star. Here we consider several different initial and several different final configurations. Initial metastable stars are taken as hadronic, hybrid, and quark stars with unpaired quarks and possible stable stars are hybrid and quark stars, taken both with unpaired and paired phase to study the deconfinement phase transition and normal quark matter–superconducting quark matter phase transition within a large number of relativistic models used to describe compact stars.

In Sec. II of the present paper we outline the formalism used in our calculations, in Sec. III the results are shown and discussed, and in Sec. IV the conclusions are drawn.

II. FORMALISM

In this section we give a brief summary of the models used in our calculations. More information about the general formalism can be obtained from the references mentioned in the introduction. The models used for the hadron matter are the nonlinear Walecka model (NLWM) [37], in which we consider the GM1 parametrization of [38], except when the δ mesons are considered in which case we use the parametrizations discussed in Ref. [11], and the quark-meson coupling (QMC) model [39]. For the quark matter we have used the Nambu–Jona-Lasinio (NJL) model [40], the MIT bag model [41], and the bag model with paired quarks to which we refer as the color-flavor-locked (CFL) phase model [19].

A. Hadronic phase

1. Nonlinear Walecka model

The Lagrangian density of the NLWM with δ mesons and hyperons reads [37,38,42]

$$\begin{aligned}\mathcal{L} = \sum_B \bar{\psi}_B[\gamma_\mu(i\partial^\mu - g_{vB}V^\mu - g_{\rho B}\mathbf{t}\cdot\vec{b}^\mu)\\ -(M - g_{sB}\phi - g_{\delta B}\mathbf{t}\cdot\vec{\delta})]\psi_B\\ +\frac{1}{2}\left(\partial_\mu\phi\partial^\mu\phi - m_s^2\phi^2 - \frac{1}{3}\kappa\phi^3 - \frac{1}{12}\lambda\phi^4\right)\\ -\frac{1}{4}\Omega_{\mu\nu}\Omega^{\mu\nu} + \frac{1}{2}m_v^2V_\mu V^\mu - \frac{1}{4}\vec{B}_{\mu\nu}\cdot\vec{B}^{\mu\nu}\\ +\frac{1}{2}m_\rho^2\vec{b}_\mu\cdot\vec{b}^\mu + \frac{1}{2}(\partial_\mu\vec{\delta}\partial^\mu\vec{\delta} - m_\delta^2\vec{\delta}^2),\end{aligned} \quad (1)$$

with $\sum_B$ extending over the eight baryons and where g_{iB} and m_i are, respectively, the coupling constants of the mesons

$i = s, v, \rho, \delta$ with the hyperons and their masses. Self-interacting terms for the σ meson are also included; κ and λ denote the corresponding coupling constants and $\mathbf{t}$ is the isospin operator. The set of constants is defined by $g_{sB} = x_{sB} g_s$, $g_{vB} = x_{vB} g_v$, $g_{\rho B} = x_{\rho B} g_\rho$, and $g_{\delta B} = x_{\delta B} g_\delta$, and $x_{sB}, x_{vB}, x_{\rho B}$, and $x_{\delta B}$ are equal to 1 for the nucleons. We have chosen $x_s = 0.7$ and $x_v = x_\rho = 0.783$ and assumed that the couplings to the Σ and Ξ are equal to those of the Λ hyperon [5,38]. For consistency we have taken $x_\delta = x_s$ [11].

2. QMC model

In the QMC model [39], the nucleon in nuclear medium is assumed to be a static spherical MIT bag in which quarks interact with the scalar and vector fields, σ, ω, and ρ, and these fields are treated as classical fields in the mean field approximation. The quark field, $\psi_q(x)$, inside the bag then satisfies the equation of motion

$$\left[i\,\partial\!\!\!/ - \left(m_q^0 - g_\sigma^q \sigma_0\right) - \gamma^0\left(g_\omega^q \omega_0 + \tfrac{1}{2} g_\rho^q \tau_{3q} b_{03}\right)\right]\psi_q(x) = 0, \quad q = u, d, s, \tag{2}$$

where σ_0, ω_0, and b_{03} are the classical meson fields for σ, ω, and ρ mesons. m_q^0 is the current quark mass, τ_{3q} is the third component of the Pauli matrices, and g_σ^q, g_ω^q, and g_ρ^q are the quark couplings with σ, ω, and ρ mesons, respectively. The energy of a static bag describing baryon B, consisting of three ground-state quarks, can be expressed as

$$E_B^{\text{bag}} = \sum_q n_q \frac{\Omega_q}{R_B} - \frac{Z_B}{R_B} + \frac{4}{3}\pi R_B^3 B_B, \tag{3}$$

where $\Omega_q \equiv \sqrt{x_q^2 + (R_B m_q^*)^2}$, $m_q^* = m_q^0 - g_\sigma^q \sigma$, R_B is the bag radius of the baryon B, Z_B is a parameter that accounts for zero-point motion, and B_B is the bag constant. The set of parameters used in the present work is given in Ref. [12] for the bag value $B_B^{1/4} = 210.854$ MeV, $m_u^0 = m_d^0 = 5.5$ MeV, and $m_s^0 = 150$ MeV. The effective mass of a nucleon bag at rest is taken to be $M_B^* = E_B^{\text{bag}}$. One can then compute the total energy density and the pressure including leptons and hyperons.

B. Quark phase

1. The Nambu–Jona-Lasinio model

The SU(3) version of the NJL model [40] includes most of symmetries of QCD, including chiral symmetry, and its breaking, which is essential in treating the lightest hadrons. The NJL model also includes a scalar-pseudoscalar interaction and the t' Hooft six-fermion interaction that models the axial $U(1)_A$ symmetry breaking. It is defined by the Lagrangian density

$$\begin{aligned} L = {} & \bar{q}\,(i\gamma^\mu \partial_\mu - m) q + g_S \sum_{a=0}^{8} [(\bar{q}\,\lambda^a q)^2 \\ & + (\bar{q}\, i\,\gamma_5 \lambda^a q)^2] + g_D \{\det[\bar{q}_i (1+\gamma_5) q_j] \\ & + \det[\bar{q}_i (1-\gamma_5) q_j]\}, \end{aligned} \tag{4}$$

where $q = (u, d, s)$ are the quark fields and λ_a $(0 \leqslant a \leqslant 8)$ are the U(3) flavor matrices. The model parameters are $m = \text{diag}(m_u, m_d, m_s)$, the current quark mass matrix ($m_d = m_u$), the coupling constants g_S and g_D, and the cutoff in three-momentum space, Λ.The NJL model is valid only for quark momenta smaller than the cutoff Λ.

The set of parameters is chosen to fit the values in vacuum for the pion mass, the pion decay constant, the kaon mass, and the quark condensates. We consider the following set of parameters [43,44]: $\Lambda = 631.4$ MeV, $g_S \Lambda^2 = 1.824$, $g_D \Lambda^5 = -9.4$, $m_u = m_d = 5.6$ MeV, and $m_s = 135.6$ MeV, which are fitted to the following properties: $m_\pi = 139$ MeV, $f_\pi = 93.0$ MeV, $m_K = 495.7$ MeV, $f_K = 98.9$ MeV, $\langle \bar{u}u \rangle = \langle \bar{d}d \rangle = -(246.7\ \text{MeV})^3$, and $\langle \bar{s}s \rangle = -(266.9\ \text{MeV})^3$.

2. The MIT bag model

The MIT bag model [41] has been extensively used to describe quark matter. In its simplest form, the quarks are considered to be free inside a bag and the thermodynamic properties are derived from the Fermi gas model. The energy density, the pressure, and the quark q density, respectively, are given by

$$\mathcal{E} = 3 \times 2 \sum_{q=u,d,s} \int \frac{d^3p}{(2\pi)^3} \sqrt{\mathbf{p}^2 + m_q^2} + Bag, \tag{5}$$

$$P = \frac{1}{\pi^2} \sum_q \int dp \frac{\mathbf{p}^4}{\sqrt{\mathbf{p}^2 + m_q^2}} - Bag, \tag{6}$$

$$\rho_q = 3 \times 2 \int \frac{d^3p}{(2\pi)^3}, \tag{7}$$

where 3 stands for the number of colors, 2 stands for the spin degeneracy, m_q are the quark masses, and Bag represents the bag pressure. We use $m_u = m_d = 5.5$ MeV, $m_s = 150.0$ MeV, and different choices for the Bag parameter.

3. Color-flavor-locked quark phase

Recently, many authors [19] have discussed the possibility that the quark matter is in a color-superconducting phase, in which quarks near the Fermi surface are paired, forming Cooper pairs, which condense and break the color gauge symmetry. At sufficiently high density the favored phase is called CFL, in which quarks of all three colors and all three flavors are allowed to pair. In this model, the quark matter is treated as a Fermi sea of free quarks with an additional contribution to the pressure arising from the formation of the CFL condensates.

The CFL phase can be described with the thermodynamic potential

$$\Omega_{\text{CFL}}(\mu_q, \mu_e) = \Omega_{\text{quarks}}(\mu_q) + \Omega_{\text{GB}}(\mu_q, \mu_e) + \Omega_l(\mu_e), \tag{8}$$

with $\mu_q = \mu_n/3$, where μ_n is the neutron chemical potential, and

$$\Omega_{\rm quarks}(\mu_q) = \frac{6}{\pi^2}\int_0^{\nu} p^2 dp (p-\mu_q) + \frac{3}{\pi^2} \times \int_0^{\nu} p^2 dp\left(\sqrt{p^2+m_s^2}-\mu_q\right) - \frac{3\Delta^2\mu_q^2}{\pi^2} + B, \tag{9}$$

with $m_u = m_d = 5\,\text{MeV}$ and $\nu = 2\mu_q - \sqrt{\mu_q^2 + m_s^2/3}$, $\Omega_{\rm GB}(\mu_q, \mu_e)$ is a contribution from the Goldstone bosons arising from chiral symmetry breaking in the CFL phase [19],

$$\Omega_{\rm GB}(\mu_q, \mu_e) = -\frac{1}{2} f_\pi^2 \mu_e^2 \left(1 - \frac{m_\pi^2}{\mu_e^2}\right)^2, \tag{10}$$

where $f_\pi^2 = (21 - 8\ln 2)\mu_q^2/36\pi^2$, and $m_\pi^2 = m_s(m_u + m_d) \times 3\Delta^2/\pi^2 f_\pi^2$, $\Omega_l(\mu_e)$ is the negative of the leptonic pressure, and the quark number densities are equal (i.e., $\rho_u = \rho_d = \rho_s = \nu^3 + 2\Delta^2\mu_q/\pi^2$.) In these expressions Δ, the gap parameter, is taken to be 100 MeV.

The electric charge density carried by the pion condensate is given by

$$Q_{\rm CFL} = f_\pi^2 \mu_e \left(1 - \frac{m_\pi^4}{\mu_e^4}\right). \tag{11}$$

For the case of strange stars, charge neutrality has to be enforced. Hence, we take μ_e in these expressions so that $Q_{\rm CFL}$ vanishes. We note that the CFL phase model is a variant of the MIT bag model in which the quarks are allowed to pair. Hereafter we refer to the model where pairing is present as the CFL phase model.

C. β equilibrium

After deleptonization takes place, the charge neutrality condition has to be enforced. Moreover, since the time scale of a star is effectively infinite compared to the weak interaction time scale, strangeness conservation is violated. The strangeness quantum number is therefore not conserved in a star. In a star with hadron matter, the net strangeness is determined by the condition of β equilibrium, which for baryon B is then given by $\mu_B = \mu_n - q_B\mu_e$, where μ_B is the chemical potential of baryon B and q_B its electric charge. Thus the chemical potential of any baryon can be obtained from the two independent chemical potentials μ_n and μ_e of neutron and electron, respectively.

In a star with quark matter $\mu_s = \mu_d = \mu_u + \mu_e$ and $\mu_e = \mu_\mu$.

D. Hybrid stars

To describe hybrid stars, a mixed phase with hadrons and quarks is built, in which charge neutrality is not imposed locally but only globally [5,45]. This means that the quark and hadronic phases are not neutral separately, but rather, the system will prefer to rearrange itself so that

$$\chi\rho_c^{\rm QP} + (1-\chi)\rho_c^{\rm HP} + \rho_c^l = 0,$$

where ρ_c^{iP} is the electric charge density of the phase i, χ is the volume fraction occupied by the quark phase, $(1-\chi)$ is the volume fraction occupied by the hadron phase, and ρ_c^l is the electric charge density of leptons. We consider a uniform background of leptons in the mixed phase since Coulomb interaction has not been taken into account. According to the Gibbs conditions for phase coexistence, the baryon chemical potentials, temperatures, and pressures have to be identical in both phases; $\mu_{\rm HP} = \mu_{\rm QP}$, $T_{\rm HP} = T_{\rm QP}$, and $P_{\rm HP}(\mu_{\rm HP}, T) = P_{\rm QP}(\mu_{\rm QP}, T)$, reflecting the needs of chemical, thermal, and mechanical equilibrium, respectively. As a consequence, the energy density and total baryon density in the mixed phase read

$$\langle\mathcal{E}\rangle = \chi\mathcal{E}^{\rm QP} + (1-\chi)\mathcal{E}^{\rm HP} + \mathcal{E}^l \tag{12}$$

and

$$\langle\rho\rangle = \chi\rho^{\rm QP} + (1-\chi)\rho^{\rm HP}. \tag{13}$$

III. RESULTS

Given the EOS, the next step is to solve the Tolman-Oppenheimer-Volkoff equations [46]. The main properties of the stars, including their gravitational and baryonic masses, their radii, and central energy densities are then computed.

Following Ref. [27] we assume that the conservation of the baryonic number can be approximated by the conservation of the baryonic mass of the star, where the baryonic mass is the mass of the equivalent number of baryons if the star was dissociated into neutrons at infinity. It is obtained by integrating the baryon number density over the volume and multiplying by the neutron mass, as in equation (8.28) of Ref. [5]. The quark phase will be a more compressible phase and therefore the radius of a deconfined quark star is smaller than the radius of a neutron star with equivalent baryon number, as shown in Ref. [14]. The energy released is determined by the change in the gravitational energy during the conversion. Thus we identify the energy released as $\Delta E = [M_G(\text{MS}) - M_G(\text{SS})]$, where $M_G(\text{MS})$ is the gravitational mass of the metastable star and $M_G(\text{SS})$ is the gravitational mass of the stable star. In c.g.s. units, this energy is given by

$$\Delta E = [(M_G(\text{MS}) - M_G(\text{SS}))/M_\odot]\times 17.88\times 10^{53}\,\text{erg}, \tag{14}$$

where $M_\odot$ is the solar mass. Notice that for ΔE to be positive, the gravitational mass of the metastable star, at fixed baryonic mass, has to be larger than the gravitational mass of the stable star.

In Table I we present a series of model hadronic, hybrid, and quark stars and calculate all possible released energies in the conversion mechanism. There are many works on the different types of stars, some of them mentioned in the introduction and many of them described in detail in Refs. [5,6]. In the table we just cite the references from which the calculation of the stellar gravitational and baryonic masses were taken. The typical gravitational mass versus radius plots for some of the hadronic, hybrid, and quark stars are given in Fig. 2. Only bare quark stars are considered in this paper. The value of the bag parameter is given beside the quark model; for instance, MIT 160 stands for the MIT bag model with a bag

TABLE I. Metastable, stable stars, and released energy.

MS	Model (MS)	SS	Model (SS)	M_b ($M_\odot$)	ΔE (10^{53} erg)	Figures
hadronic [11]	NLWMδ (p,n)	quark [16]	MIT 160	1.56	1.91	3
hadronic [11]	NLWMδ (p,n)	quark [16]	MIT 180	1.25	0.081	3
hadronic [11]	NLWMδ (8b)	quark [16]	MIT 160	1.56	0.94	3
hadronic [11]	NLWMδ (8b)	quark [16]	MIT 180	—	<0	3
hadronic [11]	NLWMδ (p,n)	quark [16]	CFL 160	1.56	3.73	3
hadronic [11]	NLWMδ (8b)	quark [16]	CFL 160	1.56	2.84	3
hadronic [8]	NLWM (8b)	quark [16]	MIT 160	1.56	0.95	3
hadronic [8]	NLWM (8b)	quark [16]	MIT 160	1.8	1.34	3
hadronic [8]	NLWM (8b)	quark [16]	CFL 160	1.56	2.84	3
hadronic [8]	NLWM (8b)	quark [16]	CFL 160	2.11	3.95	3
hadronic [11]	NLWMδ (p,n)	quark [17]	NJL	—	<0	3
hadronic [11]	NLWMδ (8b)	quark [17]	NJL	—	<0	3
hadronic [12]	QMC (p,n)	quark [16]	MIT 160	1.56	1.15	4
hadronic [12]	QMC (8b)	quark [16]	MIT 160	1.56	1.20	4
hadronic [12]	QMC (p,n)	quark [16]	CFL 160	1.56	2.97	4
hadronic [12]	QMC (8b)	quark [16]	CFL 160	1.56	3.01	4
hadronic [12]	QMC (p,n)	hybrid [15]	QMC+kaons	1.56	0.088	5
hadronic [12]	QMC (8b)	hybrid [15]	QMC+kaons	—	0.0	5
hadronic [11]	NLWMδ (8b)	hybrid [11]	NLWMδ(8b)+MIT 180	1.56	0.071	6
hadronic [8]	NLWM (8b)	hybrid [47]	NLWM(8b)+MIT 170	1.56	0.42	6
hadronic [8]	NLWM (8b)	hybrid [47]	NLWM(8b)+MIT 160	1.56	0.58	6
hadronic [8]	NLWM (8b)	hybrid [47]	NLWM(8b)+CFL 200	1.56	0.008	6
hadronic [8]	NLWM (8b)	hybrid [9,10]	NLWM(8b)+NJL	1.52	0.027	6
hybrid [11]	NLWMδ(8b)+MIT 180	quark [16]	MIT 160	1.56	0.92	7
hybrid [11]	NLWMδ(8b)+MIT 180	quark [16]	CFL 160	1.56	2.75	7
hybrid [9]	NLWM(8b)+MIT 170	quark [16]	MIT 170	—	<0	7
hybrid [9]	NLWM(8b)+MIT 160	quark [16]	MIT 160	1.56	0.46	7
hybrid [12]	QMC(8b)+CFL 200	quark [16]	CFL 160	1.56	2.89	7
hybrid [12]	QMC(8b)+CFL 200	quark [16]	CFL 160	1.8	3.31	7
hybrid [11]	NLWM(8b)+NJL	quark [17]	NJL	—	<0	7
hybrid [47]	NLWM(8b)+CFL 200	quark [16]	CFL 160	1.56	2.90	7
quark [16]	MIT 160	quark [16]	CFL 160	1.56	1.82	4

parameter $Bag^{1/4} = 160\,\text{MeV}$. ($p$,$n$) means that only protons and neutrons are considered in the EOS and (8b) means that the eight lightest baryons are taken into account. The numbers associated with the figures displaying the baryonic versus gravitational masses of the stars considered in each case are also shown in Table I.

To obtain ΔE we consider stars with baryonic masses of the order of $1.56M_\odot$, which correspond to metastable stars with

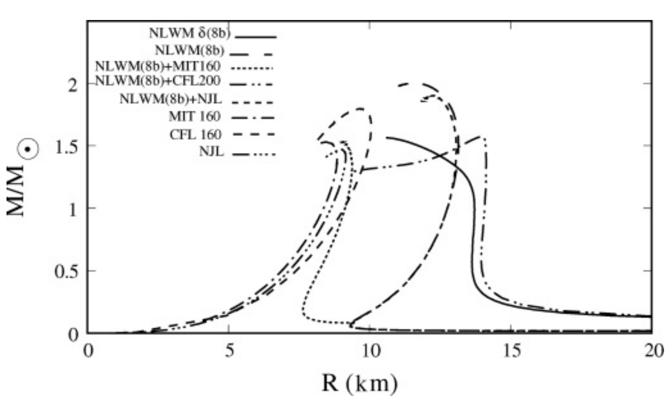

FIG. 2. Mass to radius relation for the families of stars described by some of the hadronic, hybrid, and quark stars described in the text.

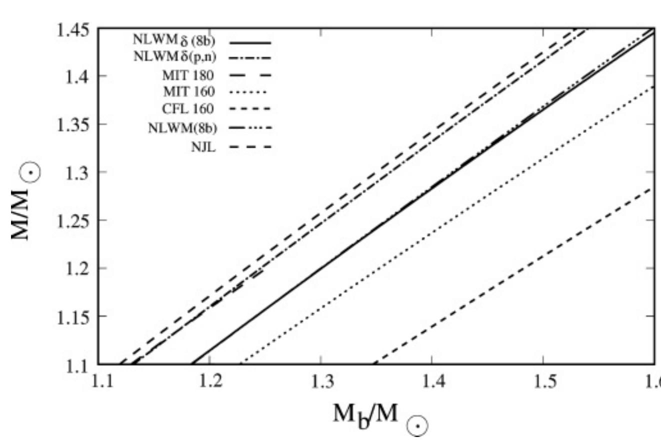

FIG. 3. Gravitational vs baryonic masses for the families of stars described by hadronic and quark matter.

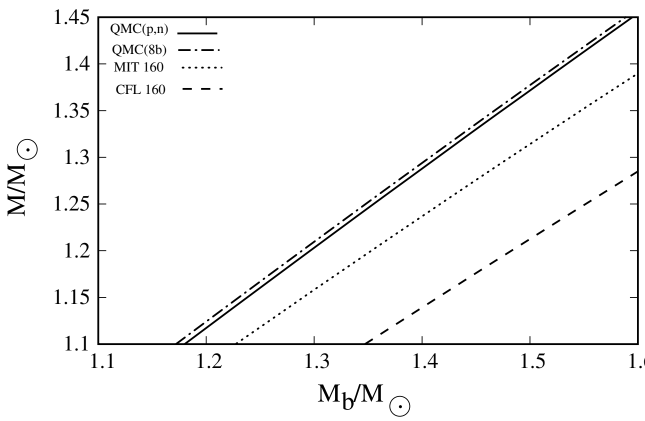

FIG. 4. Gravitational vs baryonic masses for the families of stars described by hadronic and quark matter.

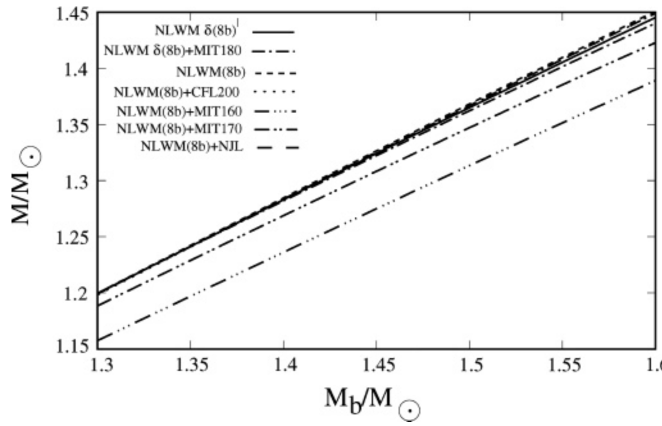

FIG. 6. Gravitational vs baryonic masses for the families of stars described by hadronic and hybrid matter.

gravitational masses of the order of $1.4M_\odot$, except in a few cases where the maximum baryonic masses were smaller or too close to this value. After conversion the gravitational mass of the stable star is, of course, smaller. In most cases we could have chosen larger values for the fixed baryonic mass and they would yield ΔE values larger than but of the same order of magnitude as the ones shown in Table I. Some examples are also shown in the table. For the EOSs studied, we could have chosen hadronic stars with larger baryonic masses but they would decay into black holes, releasing larger amounts of energy than given in the table. Notice that this choice is limited by the value of the maximum baryonic masses obtained with each EOS and these values generally lie between $2M_\odot$ and $2.5M_\odot$.

It is worth pointing out that we choose the bag parameter for the MIT and CFL phase models as $Bag^{1/4} = 160\,\text{MeV}$ because, according to Ref. [20], quark matter in β equilibrium has an energy per baryon smaller than the neutron mass only for $144\text{ MeV} \leqslant Bag^{1/4} \leqslant 162\,\text{MeV}$, in which case a star consisting only of quarks is possible. For larger *Bag* values, quarks can only exist in the core of the neutron stars. Smaller *Bag* values would predict u, d-quark matter more stable than iron. For the CFL phase model, it is possible to obtain strange stars within a larger range of *Bag* values for matter in β equilibrium. Nonetheless, in the same reference, the authors suggest that a pre-existing phase with quark matter not in β equilibrium has to exist for a short period during the deconfinement phase transition. In this case $Bag^{1/4} < 176\,\text{MeV}$ for $1.6M_\odot$ stars. We have considered $Bag^{1/4} > 160\,\text{MeV}$ in some cases to verify whether a conversion would take place when the restrictions mentioned here were not obeyed. In all calculations with the CFL phase model, the Δ parameter was taken equal to 100 MeV. The dependence of the EOS of strange quark matter on Δ has been studied in Refs. [19,20]. For a given *Bag* value if Δ is not large enough, unpaired quark matter is more stable than quark matter in the CFL phase [20].

We first analyze the results obtained when the metastable star is hadronic and the stable star is a strange quark star. The results depend sensitively on the models chosen for each kind of star, but they are of the order of 10^{53} erg, the only exception coming from a case where the maximum baryonic mass of the considered star (MIT 180) was too low and hence would not correspond to a stable star in an unpaired phase with a gravitational mass of $\simeq 1.4M_\odot$.

Next we consider the deconfined quark matter described within the NJL model. This model allows us to investigate in which way the correct description of the chiral symmetry may

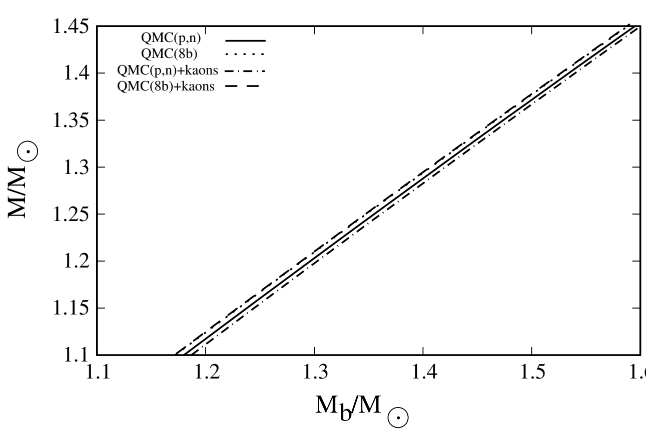

FIG. 5. Gravitational vs baryonic masses for the families of stars described by hadronic and hybrid matter.

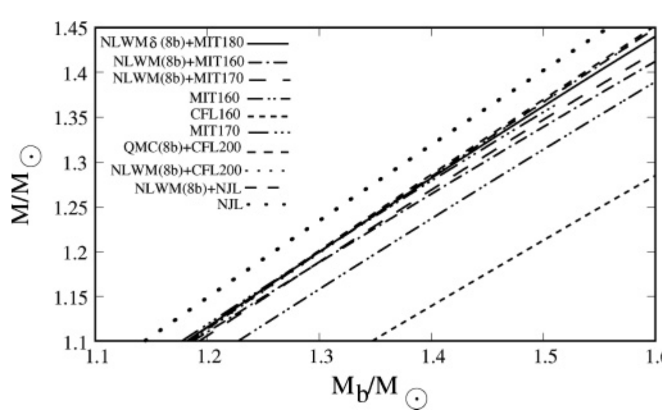

FIG. 7. Gravitational vs baryonic masses for the families of stars described by hybrid and quark matter.

affect the star properties. When the NJL model is chosen to describe the quark matter, the energy is negative, and hence conversion is not possible.This is due to the large constituent quark masses, which only approach the current quark masses at high densities and generate a lower binding energy than in the hadronic star. In fact, it was shown in Ref. [48] that for reasonable properties of the pion and kaon the NJL does not predict matter more stable than iron at zero temperature. Within the NJL model the quark EOS is more stable than the hadronic one at quite high densities. In Ref. [9] we have also shown that the effective bag constant in the NJL model is larger than $(160\ \mathrm{MeV})^4$ for $\rho > 2\rho_0$, where ρ_0 is the saturation density.

A negative value of ΔE is also obtained if a large value for the bag parameter is taken either in the MIT or in the CFL phase model. Similar values for ΔE were also mentioned in Ref. [27]. If a star with only protons and neutrons in the NLWM is considered, when it converts to a strange quark star, the energy released is always larger than in the conversion of a hadronic star with all eight lightest baryons, because if density is high enough hadronic matter with hyperons is more stable than proton-neutron hadronic matter. If the star is described by the QMC model, stars with proton and neutrons only and stars with eight baryons yield very similar results. It has been already discussed in Ref. [12] that QMC predicts a softer EOS than NLWM, and including hyperons within the QMC approach does not affect much the behavior of the EOS.

The importance of including the scalar isovector virtual $\delta(a_0(980))$ field in hadronic effective field theories when asymmetric nuclear matter is studied [42] has been stressed. Its presence introduces in the isovector channel the structure of relativistic interactions, where a balance between a scalar (attractive) and a vector (repulsive) potential exists. The δ and ρ mesons give rise to the corresponding attractive and repulsive potentials in the isovector channel. The introduction of the δ meson mainly affects the behavior of the system at high densities, when, owing to Lorentz contraction, its contribution is reduced, leading to a harder EOS at densities larger than $\sim 1.5\rho_0$ [42]. However, the influence of the δ meson in the present calculation is negligible since results with the NLWMδ and NLWM are practically identical within the precision of our calculations, although the EOSs are somewhat different owing to the different parametrizations used.

Generally speaking, the value of the released energy depends on the choice of the bag parameter used in the EOS of the quark matter. Larger ΔE values are obtained for smaller bag parameters and, within the same description, larger amounts of energy are released from the decay of more massive baryonic stars.

We now consider conversions from hadronic to hybrid stars, for which we have found much lower released energies, of the order of $10^{50}-10^{52}$ erg. The inclusion of hyperons softens the EOS at high densities but not so much as the softening from deconfinement. A conversion from a hadronic to a hybrid star with kaons is possible, but the released energy is only measurable for the case without hyperons. The neutron star charge neutrality favors kaon condensation because the neutrons at the top of the Fermi sea decay into protons plus electrons, which, in turn, have an increase in the energy as the density increases. When the electron chemical potential equals the effective kaon mass, the kaons are favored in helping with the conservation of charge neutrality because they are bosons and can condense in the lowest energy state. This process is not so effective if hyperons are included.

The choice of the parameters affects the size of the core of quarks in a hybrid star: the smaller the bag parameter, the larger the core. As a result, we find that the smaller the bag parameter, the larger the energy released in the conversion of hadronic to hybrid stars.

We also consider conversions from a hybrid star to a quark star. When the quark star is described by matter in the CFL phase, the released energy is two to three times larger than if a conversion takes place to a quark star with unpaired quarks. The energy released is of the same order of magnitude as obtained for a conversion of a hadronic to a quark star if the *Bag* value for the quark star is taken smaller than the one used for the hybrid star.

Finally, the conversion from a quark star with unpaired quarks to another one with a paired phase seems also to be possible, as a phase transition from a QGP to a color superconducting phase is possible in a QCD phase diagram. The energy released is of the order of 10^{53} erg and depends on the choice of the *Bag* and Δ values.

IV. CONCLUSIONS

We have calculated the energy released in the possible conversion of a metastable compact star into a more stable compact star, in the framework of different relativistic mean-field models for the hadronic and the quark matter. The results depend heavily on the parametrizations considered for the quark matter. Hadronic models are fitted to properties of nuclear matter at saturation. We have considered two different types of hadronic models: the NLWM and the QMC. The quark EOS was described within the NJL model and the MIT bag model. In the last case both an unpaired and a CFL phase have been considered. The parameters of the NJL model have been fitted to meson properties. For the MIT bag model we have tested different parametrizations.

Our results indicate that this kind of conversion mechanism from a metastable to a stable star releases energies of the order of $10^{50}-10^{53}$ erg, accounting for GRBs in general.

According to Refs. [24,25], the origin of SGRBs is certainly different from the origin of LGRBs. Notice, however, that the conclusions drawn in Ref. [25] are based on the present sample of bursts and their associated galaxy redshift measurements. Bursts coming from a source with a larger redshift would have a correspondingly larger intrinsic energy.

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