Effect of trapped neutrinos in the hadron matter to quark matter transition

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We study the transition from hadron matter to quark matter in the presence of a gas of trapped electron and muon neutrinos. We show that trapped neutrinos make the densities of hadron matter deconfinement noticeably higher than in the case in which neutrinos are not present. We discuss the possible consequences of this effect in supernova explosions and protoneutron star evolution. [S0556-2821(98)09218-2]

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I. INTRODUCTION

Since long ago, the phase transition from hadron matter to quark matter has been expected to occur in a variety of physically different situations such as those found in heavy ion collisions and astrophysics. We shall concentrate on this problem in the astrophysical context.

There are two situations in which the phase transition from hadron to quark matter has been studied in stellar objects: neutron stars [1,2] and proto-neutron stars (PNSs) [3]. Neutron stars are cold objects that have lived long enough to get rid of their initial neutrino content ($\mu_{\nu_e}=0$ where μ_{ν_e} is the electron neutrino chemical potential). On the other hand, PNSs are objects formed in the gravitational collapse of a massive star [4,5]. Such objects have very high temperatures in their interior (typically few tens of MeV) and, more importantly, a large amount of trapped neutrinos. After the release of its neutrino content, the PNS evolves to become a neutron star. To be more specific, neutrinos are trapped in the sense that they have a mean free path much shorter than the size of the PNS (see e.g. Ref. [6]). This forces the material to be in weak equilibrium with such neutrino plasma.

In this paper we are interested in the deconfinement transition from hadron to quark matter in the presence of a gas of trapped neutrinos. As we shall see below, the role of trapped neutrinos is fundamental in fixing the deconfinement conditions because they push the deconfinement transition densities of hadron matter to higher values than those that appear when neutrinos are not present. In a PNS, the neutrino content is radiated away in a timescale of typically ~ 10 sec. This radiation may be the very reason for massive stars to explode as supernovae. Also, it has dramatic consequences on the conditions at which the deconfinement transition to quark matter is expected to occur.

First of all, we think that the way in which hadron matter to quark matter transition should be treated needs some clarification. Glendenning [7] has developed a technique for the treatment of first order phase transitions in complex systems (systems with more than one conserved charge). The basic assumption of this development is that [7] ‘‘for complex systems the conserved charges can be shared by the two phases in equilibrium, in different concentrations in each phase. As a consequence, at a given temperature there is a mixed phase in which the pressure and the density of each phase vary continuously with the proportion of phases in equilibrium, and are not a constant as it is for simple one-component systems.’’

This technique has been applied [3,7,8] to study a phase transition from hadron matter to strange quark matter (i.e. quark matter in equilibrium under weak interactions). This transition has two conserved charges (baryon and electric) and so, the above mentioned results are obtained: it appears a phase where quark and nuclear matter in β-equilibrium coexist together with a uniform gas of electrons for a finite range of pressures. Also, the effect of trapped neutrinos has been taken into account in the study of the thermodynamics of this mixed phase [3].

This treatment is correct in the study of the mixed phase of an hybrid star with a strange quark matter core and a hadron matter envelope but not in the study of the transition process itself. This is because, as it is well known, the transition to strange quark matter is not direct, but it occurs in two steps. First, hadrons composing nuclear matter deconfine in a strong interaction time scale of $\sim 10^{-23}$ seconds to quark matter, leaving a quark gas which is *not* in equilibrium under weak interactions. The usually called two flavor quark matter is produced. Later, weak interactions will chemically equilibrate the system in a timescale of $\sim 10^{-8}$ sec. Thus, the study of a mixed phase is meaningful only after β-equilibration of the quark plasma and more importantly, *if* matter reaches the conditions for deconfinement to occur. As a result of these weak decays, the so-called strange quark matter is produced, temperature is significantly increased, and a great amount of additional neutrinos are produced [9]. This should be important not only for the supernova event itself, but also for the PNS evolution and for the neutrino signal that, in principle, may be observed. These stages of the PNS evolution have not been studied in detail yet.

In this work we shall be concerned with the process of deconfinement. We shall study it as a first order phase transition in which the abundances per baryon of each quark flavor and of each kind of lepton remain unchanged (in both the confined phase and the deconfined one) since weak interactions cannot operate in a strong interactions time scale. So, electric charge, baryon number and strangeness are automatically conserved in the transition, forcing the deconfinement process to behave as the phase transition of a simple system

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with one conserved charge and not as a complex one.

Perhaps we should clarify the last statement. The essentially different character of a first-order transition in a simple body as compared with one in a complex body originates in the possibility that in the latter the conserved charges can be shared by the two phases in equilibrium in different concentrations in each phase than those with which the body was prepared, consistent with the conservation laws [7].

The individual abundances are conserved in a deconfinement transition (flavor conservation). As we shall show in Sec. III, in this case flavor conservation implies electric charge conservation. That is, if the hadron phase is assumed to be electrically neutral, this forces the quark phase to be electrically neutral too.

Moreover, due to flavor conservation, there is not any charge (such as strangeness, electric charge, etc.) conserved by the two phases in equilibrium in different concentrations in each phase than those with which the body was prepared. So, the deconfinement transition behaves as the transition of a simple system. Since the total pressure, and the density of each phase in equilibrium is independent of the proportion of phases in a simple system, the phases will be separated by any external field that distinguishes between them, such as gravity, which distinguishes their different densities [7]. As a consequence, there should not be possible any "mixed phase" in strong equilibrium.

However, note that the deconfined phase is unstable under weak decays and so, in practice, gravity may not have the time to separate both phases. But, what is important to note, is that there is an essential difference between a deconfinement transition and the phase transition of a system with more than one conserved charge. This difference is a consequence of the validity of flavor conservation condition in deconfinement transition. It is not possible to choose a state of the whole system in which one phase has some value of a given charge different of its original value and different of the value of the charge in the other phase.

Note that although the deconfined phase has a very short life, the study we present here is essential in determining whether the onset of the transition to a quark matter state will occur or not in astrophysical conditions. The subsequent time evolution of just deconfined quark matter is out of the scope of the present paper and will be studied in future work. A study of the decay of pure two flavor quark matter into strange quark matter has been carried out in [9].

We make two approaches in the study of the deconfinement transition. First, for simplicity, we neglect the work associated with bubble formation, that is, we consider both phases in bulk. Second, in most of this work we consider that only electron neutrinos are present. This is representative of the conditions reached in the gravitational collapse of a massive star. For completeness, we shall also address the properties of the transition when muon neutrinos are degenerate, however we should note that such case should not occur in the astrophysical objects we are interested in here, because most of muon neutrinos come from pair formation, and so, in equilibrium, they should have a zero chemical potential.

The here considered transition may be relevant in the explosion of supernovae. The actual mechanism responsible for type II supernova triggering has remained elusive up to now (for a comprehensive review on this topic see Ref. [10]). However, if the transition to quark matter occurs in the core, it could provide neutrino luminosities large enough for the success of the delayed explosion mechanism. Also, it may produce a detonation wave that could carry enough energy to blow away the envelope of a massive star [11].

The remainder of this paper is organized as follows. In Sec. II we describe the equations of state we employ here for hadron matter and quark matter. In Sec. III we study the transition in a very simple model at zero temperature, and in a more complete model at finite temperature. Finally we discuss in Sec. IV the implications of the results for collapsed stellar configurations.

II. EQUATIONS OF STATE

A. Hadron matter

In order to make a detailed thermodynamical description of the nuclear phase, we shall assume the Glendenning's mean field model equation of state which incorporates n, p, Λ, Σ^+, Σ^0, Σ^-, Ξ^-, Ξ^0, μ, and e [12]. Baryons interact by means of the exchange of an attractive scalar field σ, a repulsive vector field ω, and an isovector ρ-meson field which are allowed to acquire density-dependent average values in the relativistic mean field approximation. In order to adapt such treatment to the actual conditions prevailing in PNSs, we also include a Fermi gas of electron neutrinos ν_e and muon neutrinos ν_μ in β-equilibrium with the above particles and consider the mixture to be at finite temperature.

The total pressure P and mass-energy density ρ for a system composed by baryons $B=n,\ p,\ \Lambda,\ \Sigma^+,\ \Sigma^0,\ \Sigma^-,\ \Xi^-$, Ξ^0 and leptons $L=\mu,\ e,\ \nu_e,\ \nu_\mu$ are given by

$$P=\sum_{i=B,L} P_i+\frac{1}{2}\left(\frac{g_\omega}{m_\omega}\right)^2\rho_B'^2-\frac{1}{2}\left(\frac{g_\sigma}{m_\sigma}\right)^{-2}(g_\sigma\sigma)^2 -\frac{1}{3}bm_n(g_\sigma\sigma)^3-\frac{1}{4}c(g_\sigma\sigma)^4+\frac{1}{2}\left(\frac{g_\rho}{m_\rho}\right)^2\rho_{I_3}'^2, \tag{1}$$

$$\rho=\sum_{i=B,L} \rho_i+\frac{1}{2}\left(\frac{g_\omega}{m_\omega}\right)^2\rho_B'^2+\frac{1}{2}\left(\frac{g_\sigma}{m_\sigma}\right)^{-2}(g_\sigma\sigma)^2 +\frac{1}{3}bm_n(g_\sigma\sigma)^3+\frac{1}{4}c(g_\sigma\sigma)^4+\frac{1}{2}\left(\frac{g_\rho}{m_\rho}\right)^2\rho_{I_3}'^2, \tag{2}$$

where $(g_\sigma/m_\sigma)^2=11.79\ \mathrm{fm}^{-2}$, $(g_\omega/m_\omega)^2=7.149\ \mathrm{fm}^{-2}$, $(g_\rho/m_\rho)^2=4.411\ \mathrm{fm}^{-2}$, $b=0.002947$, $c=-0.001070$ [13].

Here P_i and ρ_i are the expressions for a Fermi gas of relativistic, noninteracting particles:

$$P_i=\frac{1}{3}\frac{g_i}{(2\pi)^3}\int d^3p\ \frac{p^2}{(p^2+m_i^{*2})^{1/2}} \times(f_i(T)+\overline{f}_i(T)), \tag{3}$$

$$\rho_i = \frac{g_i}{(2\pi)^3} \int d^3p (p^2 + m_i^{*2})^{1/2} \times (f_i(T) + \overline{f}_i(T)), \quad (4)$$

where $f_i(T)$ and $\overline{f}_i(T)$ are the Fermi-Dirac distribution functions for particles and antiparticles respectively:

$$f_i(T) = (\exp([(p^2 + m_i^{*2})^{1/2} - \nu_i]/T) + 1)^{-1}, \quad (5)$$

$$\overline{f}_i(T) = (\exp([(p^2 + m_i^{*2})^{1/2} + \nu_i]/T) + 1)^{-1}. \quad (6)$$

Note that for baryons we must use, instead of masses m_i and chemical potentials μ_i, ''effective'' masses m_i^* and chemical potentials ν_i given by

$$m_i^* = m_i + x_{\sigma i}(g_\sigma \sigma), \quad (7)$$

$$\nu_i = \mu_i - x_{\omega i}\left(\frac{g_\omega}{m_\omega}\right)^2 \rho'_B - x_{\rho i} I_{3i} \left(\frac{g_\rho}{m_\rho}\right)^2 \rho'_{I_3}, \quad (8)$$

where we assume for the relative coupling strengths the values $x_{\sigma i} = x_{\rho i} = 0.6$ and $x_{\omega i} = 0.653$ [13]. I_{3i} is the third component of the isospin of each baryon.

The weighted isospin density ρ'_{I_3} and the weighted baryon density ρ'_B are given by

$$\rho'_{I_3} = \sum_{i=B} x_{\rho i} I_{3i} n_i, \quad (9)$$

$$\rho'_B = \sum_{i=B} x_{\omega i} n_i, \quad (10)$$

where n_i is the particle number density of each baryon:

$$n_i = \frac{g_i}{(2\pi)^3} \int d^3p (f_i(T) - \overline{f}_i(T)). \quad (11)$$

The mean field $g_\sigma \sigma$ satisfies

$$\left(\frac{g_\sigma}{m_\sigma}\right)^{-2} (g_\sigma \sigma) + b m_n (g_\sigma \sigma)^2 + c (g_\sigma \sigma)^3 = \sum_{i=B} x_{\sigma i} n_i^s, \quad (12)$$

where n_i^s is the scalar density:

$$n_i^s = \frac{g_i}{(2\pi)^3} \int d^3p \frac{m_i^*}{(p^2 + m_i^{*2})^{1/2}} (f_i(T) + \overline{f}_i(T)). \quad (13)$$

Note that in the last equation particle and antiparticle contributions are added. The scalar field density does not distinguish between particles and antiparticles.

The hadron phase is assumed to be charge neutral and in weak equilibrium. Electric charge neutrality states:

$$n_p + n_{\Sigma^+} - n_{\Sigma^-} - n_{\Xi^-} - n_\mu - n_e = 0. \quad (14)$$

Chemical weak equilibrium in the presence of trapped electron neutrinos implies that the chemical potential μ_i of each baryon in the hadron phase is given by

$$\mu_i = q_B \mu_n - q_e (\mu_e - \mu_{\nu_e}), \quad (15)$$

where q_B is its baryonic charge and q_e is its electric charge. In most of the following study we shall assume (if not explicitly stated) that muon and tau neutrinos are not present; so, we have $\mu_\mu = \mu_e - \mu_{\nu_e}$. In the case in which muon neutrinos are present, the above relation modifies to $\mu_\mu = \mu_e - \mu_{\nu_e} + \mu_{\nu_\mu}$.

All the above equations can be solved numerically by giving three quantities (four if ν_μ is present); for example, the values of the temperature T, the mean field $g_\sigma \sigma$ and the chemical potential(s) of the electron neutrino μ_{ν_e} (and muon neutrino μ_{ν_μ}).

B. Quark matter

The quark plasma phase is composed by u, d, and s quarks, electrons, muons, electron neutrinos and muon neutrinos. We describe this phase by means of the MIT bag model at finite temperature with zero strong coupling constant, zero u and d quark masses and strange quark mass m_s in the range 100–200 MeV. The pressure is given by

$$P = \sum_i P_i - B, \quad (16)$$

where the sum goes on $i = u, d, s, e, \mu, \nu_e, \nu_\mu$ and P_i is given by a formula like Eq. (3). The degeneracy of quarks is $g = 6$, for electrons and muons $g = 2$, and for neutrinos $g = 1$. The particle densities are those of a free Fermi relativistic gas and are given by formulas like Eq. (11). Unfortunately, there exists a large uncertainty in the actual value of B. We assume different values for the bag constant B, but we neglect any dependence of B with density or temperature.

III. DECONFINEMENT PHASE TRANSITION FROM HADRON MATTER TO QUARK MATTER IN PRESENCE OF TRAPPED NEUTRINOS

A. A simple model at zero temperature

We develop now a simple model for the phase transition at zero temperature. The hadron phase is assumed to be composed by ideal non-relativistic protons and neutrons and relativistic massless electrons and electron neutrinos. So, as it is well known [14], the pressure and mass-energy density are given by

$$P_h(\mu_n, \mu_p, \mu_e, \mu_{\nu_e}) = P_n + P_p + P_e + P_{\nu_e} = \frac{(\mu_n^2 - m_n^2)^{5/2}}{15\pi^2 m_n} + \frac{(\mu_p^2 - m_p^2)^{5/2}}{15\pi^2 m_p} + \frac{\mu_e^4}{12\pi^2} + \frac{\mu_{\nu_e}^4}{24\pi^2}, \quad (17)$$

$$\rho_h(\mu_n,\mu_p,\mu_e,\mu_{\nu_e})=\rho_n+\rho_p+\rho_e+\rho_{\nu_e}$$
$$=\frac{m_n(\mu_n^2-m_n^2)^{3/2}}{3\pi^2}+\frac{m_p(\mu_p^2-m_p^2)^{3/2}}{3\pi^2}+\frac{\mu_e^4}{4\pi^2}+\frac{\mu_{\nu_e}^4}{8\pi^2}, \tag{18}$$

and particle densities are

$$n_n=\frac{(\mu_n^2-m_n^2)^{3/2}}{3\pi^2}, \tag{19}$$

$$n_p=\frac{(\mu_p^2-m_p^2)^{3/2}}{3\pi^2}, \tag{20}$$

$$n_e=\frac{\mu_e^3}{3\pi^2}, \tag{21}$$

$$n_{\nu_e}=\frac{\mu_{\nu_e}^3}{6\pi^2}. \tag{22}$$

The baryon number density is given by $n_B=n_p+n_n$. We can also define the abundance of each particle as $Y_i=n_i/n_B$. The nuclear phase is in β-equilibrium so, we have the following chemical equilibrium condition:

$$\mu_n=\mu_e+\mu_p-\mu_{\nu_e}. \tag{23}$$

The system is charge neutral, so we also have $n_e=n_p$, or

$$\mu_e=(\mu_p^2-m_p^2)^{1/2}. \tag{24}$$

With these two last conditions, the state of hadron matter depends only on two chemical potentials, for example μ_p and μ_{ν_e}.

Neutrons and protons are composed only by u and d quarks so, after deconfinement, the quark phase contains only u, d, e and ν_e. For simplicity, we assume here that all these particles are massless so, the pressure is

$$P_q(\mu_u,\mu_d,\mu_e^q,\mu_{\nu_e}^q)=P_u+P_d+P_e+P_{\nu_e}-B$$
$$=\frac{\mu_u^4}{4\pi^2}+\frac{\mu_d^4}{4\pi^2}+\frac{\mu_e^{q4}}{12\pi^2}+\frac{\mu_{\nu_e}^{q4}}{24\pi^2}-B. \tag{25}$$

Particle densities are

$$n_u=\frac{\mu_u^3}{\pi^2}, \tag{26}$$

$$n_d=\frac{\mu_d^3}{\pi^2}, \tag{27}$$

$$n_e^q=\frac{\mu_e^{q3}}{3\pi^2}, \tag{28}$$

$$n_{\nu_e}^q=\frac{\mu_{\nu_e}^{q3}}{6\pi^2}, \tag{29}$$

and baryon number density is $n_B^q=(n_u+n_d)/3$.

Phase transition conditions between the two phases are given by the following.

(a) Pressure equilibrium:

$$P_h(\mu_p,\mu_{\nu_e})=P_q(\mu_u,\mu_d,\mu_e^q,\mu_{\nu_e}^q). \tag{30}$$

(b) Chemical equilibrium, that is equality of the Gibbs energy per baryon in both phases:

$$g_h(\mu_p,\mu_{\nu_e})=g_q(\mu_u,\mu_d,\mu_e^q,\mu_{\nu_e}^q), \tag{31}$$

where we have

$$g_h(\mu_p,\mu_{\nu_e})=Y_n\mu_n+Y_p\mu_p+Y_e\mu_e+Y_{\nu_e}\mu_{\nu_e} \tag{32}$$

and

$$g_q(\mu_u,\mu_d,\mu_e^q,\mu_{\nu_e}^q)=Y_u^q\mu_u+Y_d^q\mu_d+Y_e^q\mu_e^q+Y_{\nu_e}^q\mu_{\nu_e}^q \tag{33}$$

for hadrons and quarks respectively.

(c) Deconfinement condition: The deconfinement of hadron matter is a process that occurs in a very fast timescale typical of strong interactions, so, weak interactions do not have time to operate and the number per baryon of u, d, e and ν_e must be the same in the hadron phase as well as in the quark phase:

$$Y_i^q=Y_i^h \quad i=u,d,e,\nu_e. \tag{34}$$

Note that this condition automatically makes the quark phase to be charge neutral.

Protons are composed by two quarks u and one d, and neutrons by one u and two d so: $Y_u^h=2Y_p+Y_n$ and $Y_d^h=Y_p+2Y_n$.

Also, we have by definition: $Y_u^q=n_u/n_B^q$, $Y_d^q=n_d/n_B^q$, $Y_e^q=n_e^q/n_B^q$ and $Y_{\nu_e}^q=n_{\nu_e}^q/n_B^q$. So, eliminating n_B^q we have $Y_d^q/Y_u^q=n_d/n_u$, $Y_e^q/Y_u^q=n_e^q/n_u$ and $Y_{\nu_e}^q/Y_u^q=n_{\nu_e}^q/n_u$, or equivalently:

$$\mu_d=\left(\frac{Y_d^q}{Y_u^q}\right)^{1/3}\mu_u \tag{35}$$

$$\mu_e^q=\left(\frac{3Y_e^q}{Y_u^q}\right)^{1/3}\mu_u \tag{36}$$

$$\mu_{\nu_e}^q=\left(\frac{6Y_{\nu_e}^q}{Y_u^q}\right)^{1/3}\mu_u. \tag{37}$$

We can easily solve the problem as follows. As we have already seen, the state of the nuclear phase depends only on μ_p and μ_{ν_e}. With μ_p and μ_{ν_e} we can calculate n_p, n_n, n_e, n_{ν_e}, and so Y_p, Y_n, Y_e, Y_{ν_e}, $Y_u^h=2Y_p+Y_n$, and $Y_d^h=Y_p$

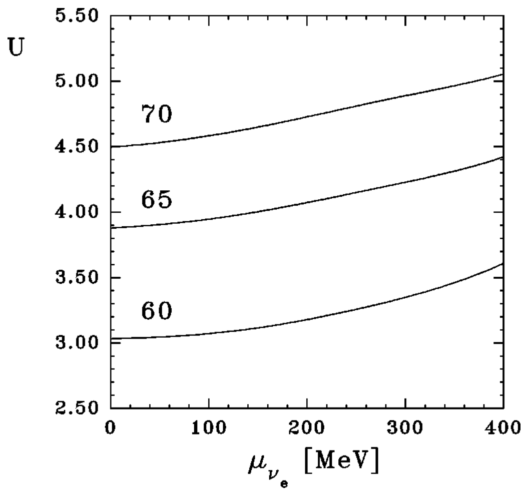

FIG. 1. The mass-energy density of hadron matter at which deconfinement phase transition occurs versus the chemical potential of the electron neutrinos present in hadron matter at $T=0$ as given by the simple model developed in Sec. III A. Density is given in units of the nuclear saturation density: $U=\rho_h/\rho_0$ being $\rho_0=2.7\times10^{14}$ g cm^{-3}. The curves are shown for different values of the bag constant B. We see that the mass-energy density of phase transition is an increasing function of the neutrino chemical potential. Of course, for densities over those of phase transition the preferred phase is deconfined quark matter and for densities below hadron matter is preferred. For values of $B\gtrsim 73$ MeV fm^{-3} no transition is found in this simple model.

$+2Y_n$. Therefore we obtain, using Eq. (34) the abundance of each particle in the quark phase. Using Eqs. (35)–(37) the state of the quark phase depends only on μ_u. Fixing the value of μ_{ν_e}, Eqs. (30) and (31) allow us to obtain μ_p and μ_u. So, the mass-energy density of hadron matter at which it deconfines depends only on one variable at $T=0$, that is, μ_{ν_e}.

The results are given in Fig. 1 for different values of μ_{ν_e} and the bag constant B. We find that the phase transition density is an increasing function of the neutrino chemical potential. This is the result we shall find in the following subsection, that has been correctly described with the above given treatment in spite of its simplicity. However, we warn the reader that, as consequence of the large differences between the detailed treatment of the nuclear equation of state and the free gas approximation here employed, the estimation performed in this subsection should not be taken as a quantitatively correct one.

B. Full treatment at finite temperature

For a more detailed study of the transition we describe hadron matter by means of the mean field model equation of state given in Sec. II A and quark matter by the equation of state given in Sec. II B. We analyze two different situations: one in which we only consider the presence of ν_e's and other in which both ν_e's and ν_μ's are present in matter.

In order to compute the transition conditions, we apply, as in the previous subsection, the Gibbs criteria, i.e. equality of pressure, temperature, and Gibbs energy per baryon in both phases:

$$P_q=P_h, \quad T_q=T_h, \quad g_q=g_h. \tag{38}$$

The Gibbs energy per baryon is

$$g=\sum_i Y_i\mu_i \tag{39}$$

where $Y_i=n_i/n_B$ is the abundance of each particle and the sum goes over all particles composing each phase, i.e. for nuclear matter $i=n,\ p,\ \Lambda,\ \Sigma^+,\ \Sigma^0,\ \Sigma^-,\ \Xi^-,\ \Xi^0,\ \mu,\ e,\ \nu_e$ (and ν_μ if present) and for quark matter $i=u,\ d,\ s,\ \mu,\ e,\ \nu_e$ (and ν_μ if present). Deconfinement condition give us six (seven) equations:

$$Y_i^q=Y_i^h, \tag{40}$$

with $i=u,\ d,\ s,\ \mu,\ e,\ \nu_e$ (and ν_μ if present). Keeping in mind the quark content of the nuclear phase particles we have

$$Y_u^h=2Y_p+Y_n+Y_\Lambda+2Y_{\Sigma^+}+Y_{\Sigma^0}+Y_{\Xi^0}, \tag{41}$$

$$Y_d^h=Y_p+2Y_n+Y_\Lambda+Y_{\Sigma^0}+2Y_{\Sigma^-}+Y_{\Xi^-}, \tag{42}$$

and

$$Y_s^h=Y_\Lambda+Y_{\Sigma^+}+Y_{\Sigma^0}+Y_{\Sigma^-}+2Y_{\Xi^0}+2Y_{\Xi^-}. \tag{43}$$

With all these equations the transition is univocally determined at a given temperature by the value(s) of μ_{ν_e} (and μ_{ν_μ} if ν_μ if present). The results of such computations are given in Figs. 2–13.

In Figs. 2–5 we show the density of hadron matter deconfinement for different temperatures and electron neutrino chemical potentials, assuming muon neutrinos to be absent. It is clearly noticed that, for a given value of B, as consequence of the presence of trapped neutrinos, the deconfinement is pushed to higher densities.

For the sake of completeness, we show in Figs. 6–9 the conditions for the transition in presence of a Fermi muon neutrino gas. In such plots we assumed a fixed value for μ_{ν_μ} and the same values of μ_{ν_e} considered in Figs. 2–5. For quark matter it is assumed $B=80$ MeV fm^{-3} and $m_s=150$ MeV. It is noticeable that the presence of muon neutrinos affects the transition density pushing it up to even larger values (compare Figs. 6–9 with dashed lines of Fig. 3). This effect is clear when μ_{ν_μ} is of the order of μ_{ν_e}, otherwise, the transition conditions are dominated by the type of neutrino with the largest chemical potential.

Note that, due to the presence of strange hadrons in the nuclear phase, the deconfinement leaves quark matter with finite strangeness, and not pure two flavor quark matter as in the previous subsection. This can be seen in Figs. 10–13 where we plot the number densities of particles in the quark phase.

We must recall that, although this quark phase is not stable under weak interactions (and will decay to strange

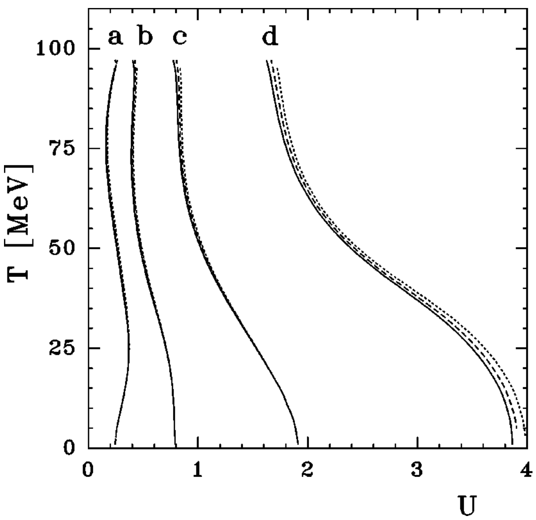

FIG. 2. The hadron matter mass-energy density of deconfinement phase transition versus the temperature T as given by the full calculation of Sec. III B. Density is given in units of the nuclear saturation density: $U=\rho_h/\rho_0$ being $\rho_0=2.7\times10^{14}$ g cm^{-3}. The bag constant is assumed to be $B=60$ MeV fm^{-3}. Full lines correspond to a value of the strange quark mass of $m_s=100$ MeV, dashed lines to $m_s=150$ MeV, and dotted lines to $m_s=200$ MeV. The labels a, b, c, and d correspond to different values of the chemical potential of the electron neutrinos in hadron matter: 0, 100, 200, and 300 MeV respectively.

quark matter in a typical timescale of $\sim10^{-8}$ sec) it is an unavoidable intermediate step in the transition to strange quark matter.

IV. DISCUSSION

We have shown that the presence of a gas of trapped neutrinos in hadron matter makes the mass-energy density of deconfinement transition to quark matter to be appreciably higher than in the case in which neutrinos are not present.

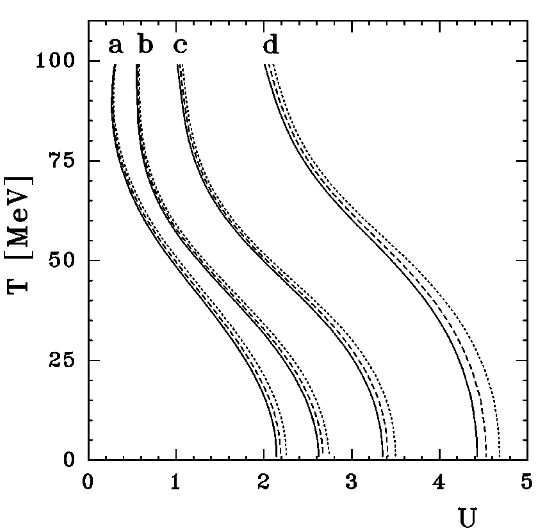

FIG. 3. The same as Fig. 2 but for $B=80$ MeV fm^{-3}.

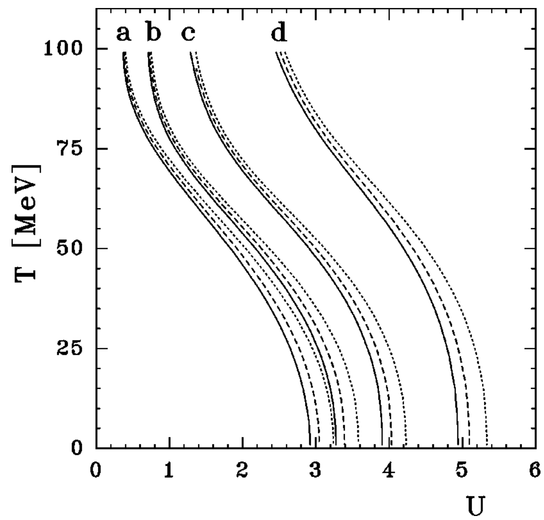

FIG. 4. The same as Fig. 2 but for $B=100$ MeV fm^{-3}.

We can qualitatively explore the consequences of this result for supernovae and proto-neutron stars (PNSs).

In order to discuss the possibility of the occurrence of such transition we must employ accurate computations of the evolution of PNSs. Such task was carried out e.g. in Refs. [4] and [5]. In this work we shall use the results of Ref. [5] on the evolution of a PNS of 1.68 $M_\odot$ (where $M_\odot$ is the solar mass) made up of hadronic matter. Let us remark that in such work the effects of muon neutrinos were neglected and also that the thermal effects on the equation of state of hadron matter were considered in an approximate way.

Just after the bounce of the supernova core, the just-formed PNS has a temperature $T\gtrsim30$ MeV in most of its structure, and more importantly, the chemical potential of electron neutrinos μ_{ν_e} is ~300 MeV in the center of the star and drops monotonically to the surface of the PNS. The timescale of neutrino diffusion (i.e. the time interval in which μ_{ν_e} drops to zero) is about 10 sec. In such stage, T is still very high (up to 50 MeV in the center), whereas the central

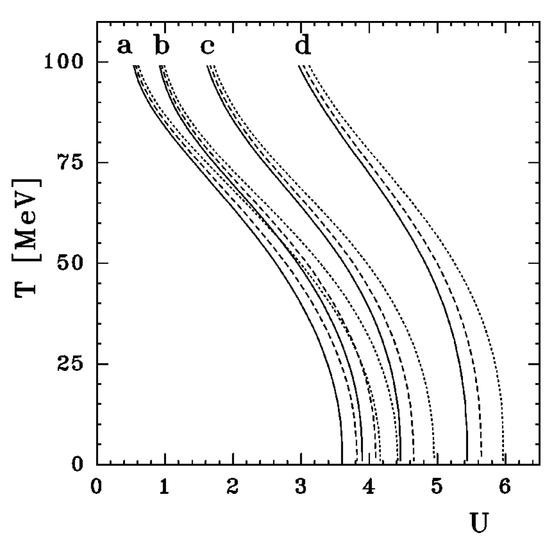

FIG. 5. The same as Fig. 2 but for $B=120$ MeV fm^{-3}.

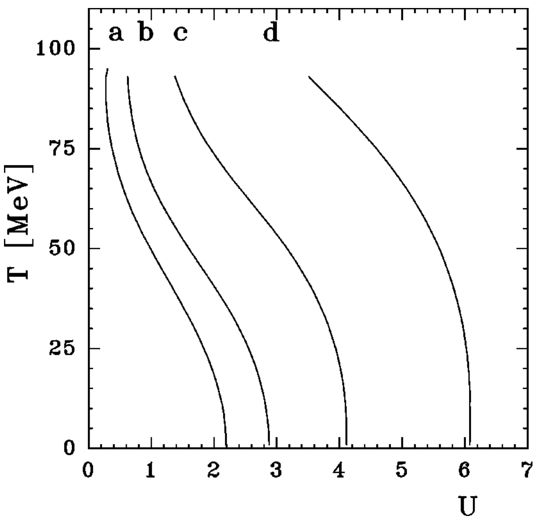

FIG. 6. The same as Fig. 2 with $\mu_{\nu_e}=0$ but in the presence of a muon neutrino gas. The labels a, b, c, and d correspond to different values of the chemical potential of the muon neutrinos in hadron matter: 0, 100, 200, and 300 MeV respectively. Here we assumed $B=80$ MeV fm^{-3} and $m_s=150$ MeV. Note that muon neutrinos have the same effect as electron neutrinos in pushing the transition density upwards.

density shows a modest increase from $U=3.11$ to $U=3.75$ in the same time interval ($U=\rho_h/\rho_0$ being $\rho_0=2.7\times10^{14}$ g cm^{-3}). Thus, due to the high values for μ_{ν_e}, if we want to study the actual occurrence of a phase transition from hadron matter to quark matter in PNSs, we must consider the effect of neutrinos, and so, apply the results of the preceding section.

These results strongly suggest that neutrino trapping should preclude the transition to quark matter in the first stages of evolution of the PNS formed in the center of a supernova, because the density of hadron matter will be below the very high density at which deconfinement is possible. However, as neutrinos are radiated away, neutrino chemical potential drops essentially to zero in ~ 10 sec making the transition density to drop appreciably. Also, the density of matter increases during the PNS evolution. These two effects modify the conditions of hadron matter helping the occurrence of the transition somewhere in the star. So, if transition to quark matter actually occurs in supernovae, it should *not* happen in the first stages of PNS evolution but after (or at least after a fraction of the) neutrino diffusion timescale. A detailed study of the astrophysical conditions at which a transition is expected to occur will be presented elsewhere.

We should remark that first order phase transitions are mediated by bubble nucleation. Such process has been studied related with the transition to quark matter without considering the effects of neutrino trapping [1]. At present, due to the large uncertainties in the value of the bag constant B and on the surface tension of quark matter is not possible to

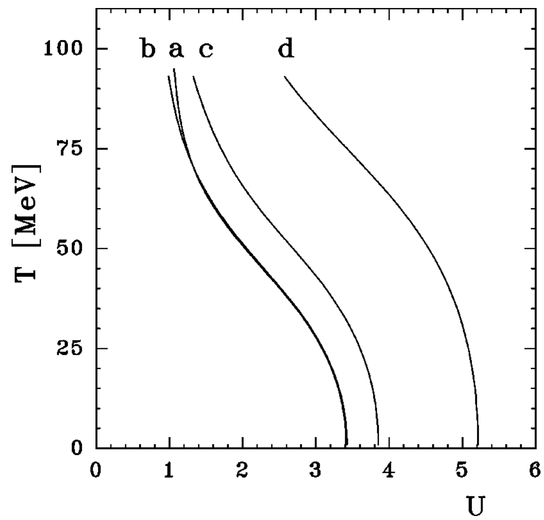

FIG. 8. The same as Fig. 6 but with $\mu_{\nu_e}=200$ MeV.

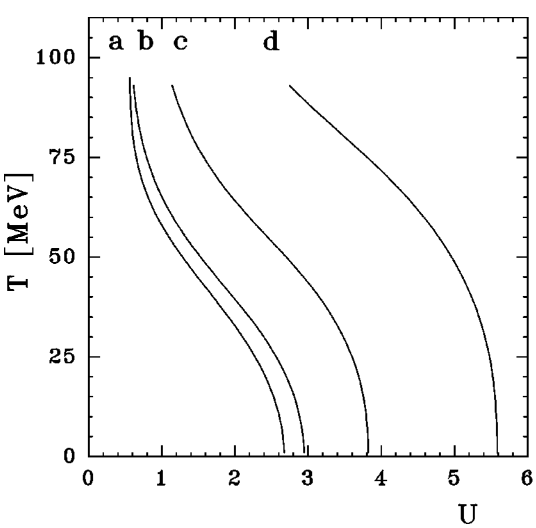

FIG. 7. The same as Fig. 6 but with $\mu_{\nu_e}=100$ MeV.

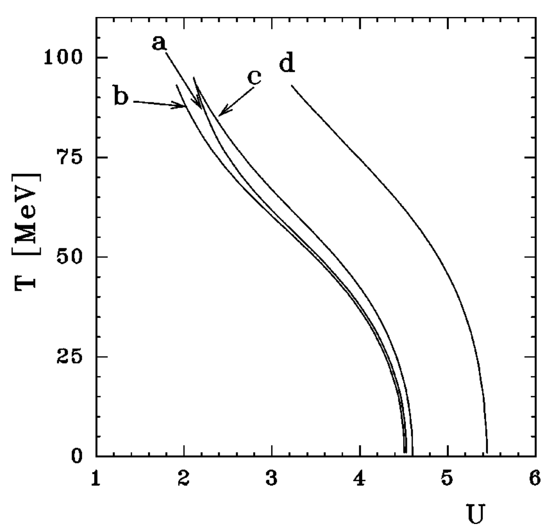

FIG. 9. The same as Fig. 6 but with $\mu_{\nu_e}=300$ MeV.

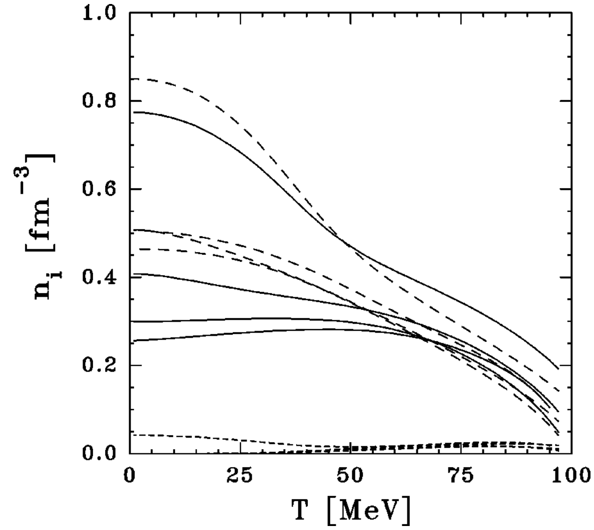

FIG. 10. The particle densities of quarks u (full lines), d (long dashed lines), and s (short dashed lines) in the just deconfined quark phase versus the temperature T as given by the full calculation of Sec. III B. We consider that ν_μ's are absent. The bag constant is assumed to be $B=60$ MeV fm^{-3} and strange quark mass $m_s=150$ MeV. For each type of line, from bottom to top and at low temperatures, each curve corresponds to a value of the chemical potential of the electron neutrinos in hadron matter of 0, 100, 200, and 300 MeV (notice that at high temperatures the dependence on μ_{ν_e} is not monotonous).

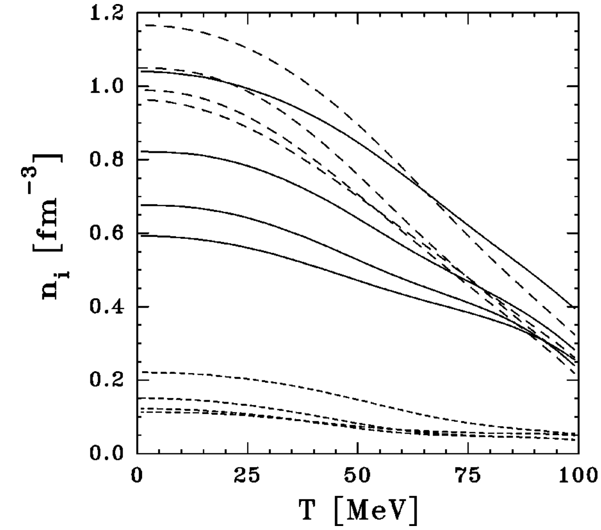

FIG. 12. The same as Fig. 10 but for $B=100$ MeV fm^{-3}.

make a safe estimate of the time involved in bubble growth. However, as temperatures here involved are very high, we should expect a very fast bubble growth, making the bulk treatment of the transition here presented as physically meaningful.

Once the first seed of quark matter is produced, it may grow (probably by means of a detonation wave [11,15]) and convert part of the PNS into quark matter. The just deconfined quark matter is not in weak equilibrium, so in a weak interaction timescale of $\sim 10^{-8}$ sec it will equilibrate and radiate a large amount of additional neutrinos. In this process, the temperature increases appreciably, so also the three kinds of neutrinos and antineutrinos are thermally generated. These additional neutrinos are radiated away in a diffusion time scale and may be crucial in the success of the delayed explosion mechanism because they could provide the necessary energy to revive the shock wave. Also, the time delay expected for the transition to quark matter should be observed by terrestrial detectors as a second neutrino signal superimposed to the standard one. Since the neutrino diffusion timescale is of the order of the time interval between the two neutrino peaks observed in SN1987A by the Kamiokande group [16] we believe (see [11]) that (although many researchers adjudicated them to a statistical fluctuation in the framework of a standard supernova explosion) such late Kamioka events are due to the neutrino-delayed quark deconfinement and subsequent decay to strange quark matter process.

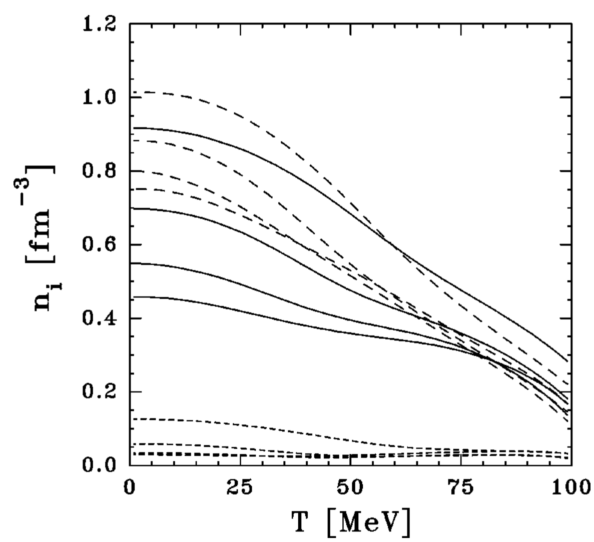

FIG. 11. The same as Fig. 10 but for $B=80$ MeV fm^{-3}.

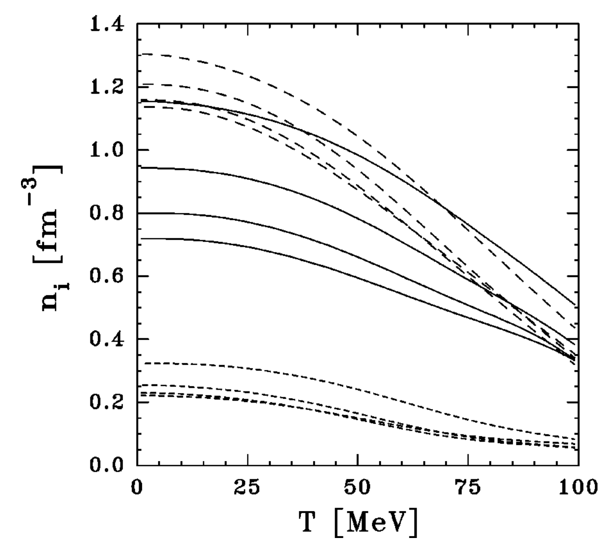

FIG. 13. The same as Fig. 10 but for $B=120$ MeV fm^{-3}.

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