Primordial nuggets survival and QCD pairing

G. Lugones* and J. E. Horvath†
Instituto de Astronomia, Geofísica e Ciências Atmosféricas/Universidade de São Paulo, Rua do Matão 1226, (05508-900) São Paulo SP, Brazil
(Received 2 October 2003; published 17 March 2004)

We reexamine the problem of boiling and surface evaporation of quark nuggets in the cosmological quark-hadron transition with the explicit consideration of pairing between quarks in a color-flavor locked state. Assuming that primordial quark nuggets are actually formed, we analyze the consequences of pairing on the rates of boiling and surface evaporation in order to determine whether they could have survived with substantial mass. We find a substantial quenching of the evaporation+boiling processes, which suggests the survival of primordial nuggets for the currently considered range of the pairing gap Δ. Boiling is shown to depend on the competition of an increased stability window and the suppression of the rate, and is not likely to dominate the destruction of the nuggets. If surface evaporation dominates, the fate of the nuggets depends on the features of the initial mass spectrum of the nuggets, their evaporation rate, and the value of the pairing gap, as shown and discussed in the text.

DOI: 10.1103/PhysRevD.69.063509 PACS number(s): 98.80.Cq, 12.38.Mh, 95.35.+d

I. INTRODUCTION

Approximately 10^{-5} s after the big bang, the early Universe was filled with a hot and expanding mixture of elementary particles. The Universe was composed mainly of photons, charged leptons, neutrinos, quarks, and gluons (and the corresponding antiparticles) coexisting in thermal and chemical equilibrium through electroweak interactions. As the Universe expanded, this mixture cooled down to a critical temperature at which the plasma of free quarks and gluons converted into hadrons. Early studies of this transition started in the 1980s [1–3] and gave a broad-brush picture of the physics involved (for a more complete reference list, see [4–9]).

A very important question is whether the transition is actually first order, second order, or just a crossover. Dramatic effects are expected in the first case, while a second order or crossover would be much less spectacular. Lattice numerical simulation is the best approach currently available for the study of QCD near the finite-temperature transition point. While it has been known for a long time that the transition is first order in the case of pure gluonic calculations (corresponding to infinitely heavy quarks) and in the case of four light quarks, the actual physical case is elusive. At present, there are well established nonperturbative lattice techniques to study this transition at $\mu=0$ and $T\neq 0$. The order of the transition is known as a function of the quark masses showing that the physical point is probably in the crossover region unless the s quark mass is small (in which case it should be first order). For recent reviews, see [10,11] and references therein.

Interesting baryon fluctuations would have been produced by a first-order transition. The two phases need to coexist long enough for baryon transport to shuffle the baryon number across the phase boundary. As pointed out in early studies, the onset of the supposedly first-order transition requires some degree of supercooling [3]. If the transition is not first order, no supercooling could possibly occur (even if the equation of state gave rise to a very rapid change in the energy density) due to the extremely slow expansion of the Universe.

The generation of primordial isothermal baryon number inhomogeneities can be understood within a scenario of cosmic separation of phases [4,8]. When the Universe cools to the critical temperature T_{QCD}, nucleation of bubbles of the hadron phase could begin. However, it is a general feature of the nucleation theory that the nucleation probability is not large enough at the critical temperature but for temperatures below it. Therefore, the Universe supercools below T_{QCD} being still in the quark phase until the nucleation rate becomes sufficiently large. After a brief stage of nucleation during which the hadron bubbles grow and reheat the Universe back to T_{QCD}, nucleation is again inhibited due to its low probability and the expansion of the Universe causes hadron bubbles to grow slowly at the expense of the quark phase. Once hadron bubbles occupy roughly half of the total volume, they are able to collide and merge, leaving the Universe with shrinking droplets of quark-gluon plasma immersed in a hadron matter medium.

The fate of these baryon number inhomogeneities depends on how heat and baryon number are transported across the transition front [4,8]. Latent heat (or entropy) could be carried out by neutrinos, surface evaporation of hadrons (mostly pions), and by the motion of the transition front which converts the volume of one vacuum into another. The baryon number transport across the conversion front depends on the bulk properties of both phases and on the penetrability of the interface (which quantifies the chance of the baryon number to pass from one side of the boundary to the other). Estimations of baryon number penetrability have been made within the frame of the chromoelectric flux tube model [8,12–14]. If the baryon penetrability indeed happens to be small, it may be possible to accumulate almost all the baryon number density in the quark-gluon phase (see below). In

*Email address: glugones@astro.iag.usp.br
†Email address: foton@astro.iag.usp.br

such a case, and depending on the parameters, the inhomogeneities may be large enough to produce strange quark matter (SQM). This results in a universe in thermal equilibrium but with an inhomogeneous baryon distribution (i.e., out of chemical equilibrium).

The study of the extreme case in which quark nuggets form has been undertaken by a number of authors [15–22]. An absolute upper limit to the baryon number contained in the nuggets is determined by the size of the cosmological horizon evaluated at the critical temperature. The simplest estimate yields the well-known value

$$A_{\max}=10^{49}\left(\frac{100\ \mathrm{MeV}}{T_{\mathrm{QCD}}}\right)^2. \quad (1)$$

Actually, the details of the dynamics will determine an initial mass function at the end of the transition. This is a quite complicated problem and has not been solved in detail, although Bhattacharyya *et al.* [23] presented a series of calculations showing that the maximum baryon number of the nuggets is $\sim 10^{43}$ for $T_{\mathrm{QCD}}=150$ MeV, which fit comfortably within the horizon size.

After the QCD phase transition, the temperature in the primordial Universe is still high enough to allow for evaporation of hadrons from the surface of the nuggets, and in principle to allow for boiling (nucleation of hadronic bubbles inside its volume). This is a consequence of the presence of the $-TS$ term in the free energy, which disfavors the strange quark matter phase at intermediate temperatures. It is only at low temperatures ($T\approx 2$ MeV) that nuggets begin to be preferred to free hadrons. Previous work found that boiling is not possible for reasonable values of the bag constant B since the time scale is too short for bubble nucleation to take place [20,21]. If so, surface evaporation seems to be the only mechanism able to destroy the primordial nuggets, although the very survival of these entities may be considered as still subject to uncertainties. Since it is likely that quarks inside the nuggets settle in paired states at a relatively high temperature (see Fig. 1 for a qualitative sketch), we shall examine in the remainder of this work the effects of quark pairing on the evaporation/boiling at intermediates temperatures, thus revisiting the question of nugget survival.

II. BOILING OF CFL NUGGETS

As stated above, quark nuggets are born hot and therefore nucleation of hadronic bubbles could occur inside them. Nevertheless, as has been shown previously [21], boiling is unlikely to destroy primordial nuggets for reasonable values of the bag constant B. However, we must note that pairing must occur at temperatures below the critical temperature $T_\Delta=0.57\Delta$. Therefore, the analysis made in [21] holds only in the temperature regime between T_{QCD} and T_Δ, while below T_Δ pairing effects must be taken into account. A remarkable consequence of QCD pairing is that the stability window for strange matter is considerably enlarged, allowing a wider range for B [24]. Thus, although pairing should make boiling difficult because more energy is necessary to produce a hadron lump, it is not clear *a priori* to what extent the modification of the stability characteristics of SQM counteracts this effect.

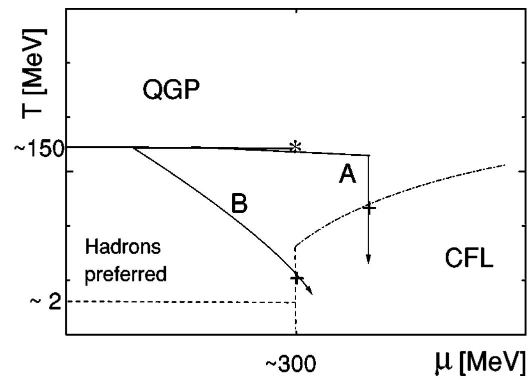

FIG. 1. Path of the strange quark matter nuggets in the T-μ plane. Nuggets are quickly formed starting at $T=T_{\mathrm{QCD}}$ ~ 150 MeV, and they are fragile to evaporation/boiling at intermediate temperatures, as discussed in the text. A transition to the CFL phase occurs at the points marked with crosses. The path labeled "A" assumes a very quick formation of the nugget (that is, $t_{\mathrm{formation}} \ll H^{-1}$, see Ref. [22]), which evolves at constant density afterwards following a vertical path. A (perhaps more realistic) path "B" has been also sketched, in which the formation is slower, $t_{\mathrm{formation}} \sim H^{-1}$. After crossing the CFL boundary, nuggets are "safe" because the pairing now protects them against evaporation/boiling, and attaining the dashed line temperature is no longer relevant for their fate. Thus, their masses freeze at a higher value when quarks become locked in CFL states.

Let us briefly examine the thermodynamic description of boiling of primordial nuggets including the effect of pairing between quarks. We assume, for simplicity, that the nucleated phase is in thermal and chemical equilibrium with the nugget, which itself evolves at fixed μ (that is, along a path of the type "A" in Fig. 1). The work done to form a bubble of radius r composed by the hadronic phase inside the quark phase is

$$W=-\frac{4}{3}\pi r^3\Delta P+4\pi\sigma r^2-2\pi\gamma r+\frac{4}{3}\pi r^3\times\frac{3}{2}n_B\Delta, \quad (2)$$

where $\Delta P=P_h-P_q$ is the pressure difference between both phases, $\sigma=\sigma_q+\sigma_h$ is the surface tension, $\gamma=\gamma_q-\gamma_h$ is the curvature coefficient, n_B is the baryon number density in the hadronic phase, and Δ is the gap of the CFL pair. The innovation here is the last term which is introduced by considering that to put three quarks in a hadron, an energy Δ must be expended for each CFL pair that is disassembled inside the quark phase. Further refinements should be considered if hadrons are assumed to be composed by a diquark plus a free quark; in this case, the net energy released (per baryon) should be $\sim\frac{1}{2}\Delta$. The effect of Δ on σ and γ themselves is unknown and will be neglected in this first approach. Note that the gap also enters the free energy through the pressure as a quadratic term. The critical radius r_c is found by extremizing W,

$$r_c=\frac{\sigma}{F}[1+\sqrt{1-b}] \quad (3)$$

where $b \equiv \gamma F / 2\sigma^2$ and

$$F \equiv \Delta P - \frac{3}{2} n_B \Delta. \tag{4}$$

Only those bubbles with a radius greater than r_c will be able to grow. The nucleation rate for the growing bubbles is given by

$$\mathcal{R}_{\mathrm{boil}} = T^4 \exp(-W_c / T) \tag{5}$$

where $W_c = 4\pi\sigma^3 / 3F^2 [2 + 2(1-b)^{3/2} - 3b]$. As we shall see below, the contribution of Δ tends to suppress the rate since it enters in such a way that F becomes smaller and W_c becomes larger. However, since the stability behavior is modified by pairing, allowing stability for a much wider range of the bag constant B, it is necessary to determine which is the leading effect on the boiling process, as we shall do in the following.

The effect of boiling on the quark nuggets can be analyzed by means of a slightly different condition (see [21] and references therein). Comparing the total area of the nugget and the bubbles

$$\sum_i A_i^{\mathrm{bubbles}} = A^{\mathrm{nugget}}, \tag{6}$$

it is found that boiling is less important than surface evaporation for a baryon number A below the value A_{boil} given by

$$A_{\mathrm{boil}} = 7.90 \times 10^{-61} \left(\frac{2F}{(1+\sqrt{1-b}) T\sigma} \right)^6 \times \exp\left(\frac{\pi\sigma^3 [2 + 2(1-b)^{3/2} - 3b]^2}{TF^2} \right). \tag{7}$$

For a given value of A_{boil}, the last equation gives the critical bag constant B and σ separating the boiling and the nonboiling regions. The value of the critical B and σ is almost insensitive to the value of A_{boil}. Spanning the range $1 < A_{\mathrm{boil}} < 10^{57}$ only changes the critical values of $B^{1/4}$ and $\sigma^{1/3}$ by a few MeV [19].

From an inspection of Eq. (7) it is easily recognized that the main effect of pairing on boiling enters only through a boost in the bag constant B. This can be understood by comparing the boiling of CFL strange matter with the boiling of unpaired SQM. If we assume that the strange quark mass m_s is zero, the pressure of SQM and CFL strange matter differs only by a term proportional to Δ^2, i.e., $P_{\mathrm{SQM}} = P_{\mathrm{CFL}} - (3\Delta^2\mu^2)/\pi^2$ [24]. In both the CFL and the unpaired SQM phases, there exists a symmetric flavor composition with $n_u = n_d = n_s$. Note that the last will not be true when considering that the strange quark has a finite mass; in this case, the CFL phase will still have a symmetric composition but SQM will not. However, we do not expect strong departures from this simple analysis. The nucleation of hadron bubbles occurs in a very fast time scale $\tau_s \sim 10^{-24}$ s, typical of strong interactions, therefore the transition happens out of weak equilibrium irrespective of the pairing of quark matter. This means that flavor will be conserved during the process of nucleation (only after $\tau_w \sim 10^{-8}$ s will β decays lead the just formed hadron phase to its equilibrium configuration). The conservation of flavor during the formation of bubbles, and the fact of both SQM and CFL strange matter having the same composition (in a first approximation), guarantee that exactly the same gas of hadrons will be produced both beginning with a CFL or with an unpaired SQM composition. Therefore, the pressure difference ΔP_0 between the SQM phase and the hadron phase, and the difference ΔP between the CFL phase and the hadron one, are related by

$$\Delta P = \Delta P_0 - \frac{3\Delta^2\mu^2}{\pi^2}. \tag{8}$$

The simple relation given by Eq. (8) allow us to gauge straightforwardly the impact of pairing in the boiling process. The difference F can be written, in the case of massless quarks, as

$$F = \Delta P - \frac{3}{2} n_B \Delta \approx \Delta P_0 - \frac{3}{\pi^2}\Delta^2\mu^2 - \frac{3}{2} n_B \Delta = P_h - P_{\mathrm{free}} - B_{\mathrm{eff}}. \tag{9}$$

The effect of pairing has been included in the bag constant B by defining an "effective bag constant"

$$B_{\mathrm{eff}} = B - \frac{3}{\pi^2}\Delta^2\mu^2 - \frac{3}{2} n_B \Delta, \tag{10}$$

where P_{free} is the pressure of a flavor-symmetric mixture of free quarks. Note that pairing enters through a leading contribution $\sim \mu^3\Delta$ associated with the condensation work and a second-order contribution $\sim \Delta^2\mu^2$ associated with the modification of the pressure in the CFL phase. Although our analysis does not include the important effects of the finite mass of the strange quark and of the chemical composition of the phases, it seems clear that the tendency shown here should be qualitatively the same in a more realistic study. Also, as stated above, some refinements would need to be considered due to the effect of finite temperature and the pairing gap in the surface tension and curvature terms. In summary, the only difference with the boiling of quark nuggets made up of unpaired massless quarks is that here we must use the effective bag constant defined by Eq. (10).

The likelihood of boiling can be analyzed in the B versus T plane (see Fig. 2). For an unpaired quark mixture, the critical B above which boiling is allowed (for any baryon number A) is always greater than the maximum B that allows stable SQM for a transition out of the equilibrium [21]. As we have shown in Eq. (10), the pairing shifts up the critical curve by a magnitude $(3\Delta^2\mu^2)/\pi^2 + (3n_B\Delta)/2$. On the other hand, as shown in Ref. [24], the maximum B that allows stable CFL strange matter also shifts up a magnitude $m_n^2\Delta^2/(3\pi^2)$. Therefore, the net shift is

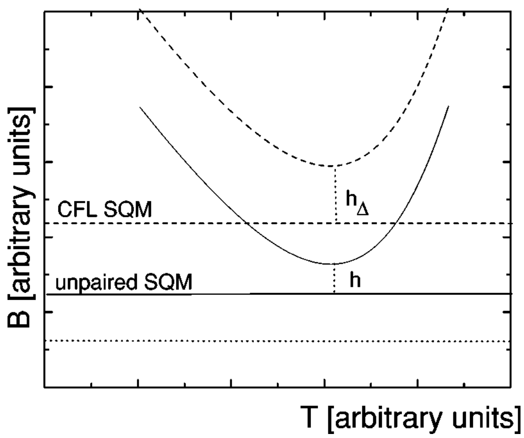

FIG. 2. Sketch of the effects of CFL pairing on the boiling of nuggets. The stability window of SQM is realized between the lower dotted line and the solid horizontal line. Analogously, the wider stability window of CFL SQM holds between the lower dotted line and the short-dashed horizontal line, as marked [24]. For a given temperature, boiling of primordial quark nuggets is allowed only above the parabolic-like curves (solid line for unpaired SQM and dashed for CFL SQM). In spite of the rising of the upper stability limit for CFL SQM, the boiling curve also rises, and since both curves lie above the corresponding upper limits for matter stability (horizontal lines), quark matter nuggets must survive boiling (with or without pairing).

$$h_{\Delta}-h=\frac{(9\mu^2-m_n^2)\Delta^2}{3\pi^2}+\frac{3}{2}n_B\Delta, \qquad (11)$$

which is clearly positive in the range of interest since the leading term is the second one and scales as $\mu^3\Delta$. This means that nucleation is impossible in the temperature regime where QCD pairing operates ($T<T_{\Delta}$), even in the most favorable situation in which the nucleated phase is in equilibrium.

III. SURFACE EVAPORATION

Surface evaporation of hadrons has been addressed as a first mechanism for nugget destruction by Alcock and Farhi [15], as explained above. Using simple detailed balance arguments, they concluded that the flux of baryons from the surface was more than enough to evaporate the nuggets for all but the highest (unphysical) masses. Further work revisited the issue by employing chromoelectric flux tube expressions [13], which happened to be much smaller than the naive flux employed originally. It was found that nuggets having a baryon number larger than 10^{39} could survive evaporation. Another detailed study by Madsen, Heiselberg, and Riisager [16] also considered explicitly the effect of flavor nonequilibrium at the surface of the nugget, and found that lumps with baryon number as low as $A\sim10^{46}$ could survive evaporation. These and other calculations [12] suggest that a low surface baryon flux allows the survival of large, but not extreme, nuggets, perhaps down to $A\sim10^{40}$ provided the evaporation flux is low enough. Therefore, it is of interest to understand what happens when quark pairing is introduced in this picture.

An evaporating lump is slightly cooler than the environment, which for temperatures ~100 MeV is composed mainly by photons, neutrinos, electrons, and their antiparticles. Heat flows from the surrounding medium into the lump, providing the energy to power the evaporation. Near the surface of the lump, there is a high concentration of baryons that have just evaporated. Due to their large mean free path, neutrinos are the most efficient energy carriers. We shall discuss the effects of pairing using a scaling of the simplest rate derived by Alcock and Farhi [15], since the latter provides a good description of all surface evaporation rates presented in the literature. The number of hadrons evaporated from the surface per second is written as

$$\mathcal{R}=\alpha\frac{\sigma_0 m_n}{2\pi^2}T^2A^{2/3}e^{-I/T} \qquad (12)$$

with $I=20$ MeV the binding energy, $\sigma_0=3.1\times10^{-4}$ MeV^{-2}, and m_n the neutron mass. The parameter α is introduced in order to reproduce approximately the behavior of the flux within very different models, which differ by several orders of magnitude (see, e.g., [13–15]).

Above the critical temperature for pair formation T_{Δ}, the evaporation rate would be given by Eq. (12). Below T_{Δ}, we use the same combinatorial criterion as in the previous section for the breakup of quark pairs to write down the rate of evaporation from CFL nuggets $\mathcal{R}_{\Delta}$ in a first approximation as

$$\mathcal{R}_{\Delta}=\mathcal{R}\times e^{-3\Delta/2T}. \qquad (13)$$

Since the energy cost of pair breakup is a general feature of the models, we expect this suppression to hold quite independently of the detailed physics. The important feature is that surface evaporation rates get effectively quenched as soon as the CFL phase sets in, at a temperature $T_{\Delta}=0.57\Delta$, which is certainly much higher than the ~2 MeV necessary to stabilize the nuggets. Therefore, it may be said that CFL states freeze out the mass of the nuggets once they cool down to T_{Δ}.

The parametrization of the surface evaporation rates given in Eqs. (12) and (13) allows a simple analytic solution of the evolution equation $dA/dt=\mathcal{R}$ of the baryon number of the nugget. This solution is

$$A^{1/3}(T)=A_0^{1/3}-C\int_T^{T_0}\frac{e^{-I/T}}{T}\,dT, \qquad (14)$$

where we have identified the initial baryon number $A_0\equiv A(T_0)$ and $C=\alpha(-2\sigma_0 m_n)/(6\pi^2)\times[45/(172\pi^3G)]^{1/2}$ (with G the Newton constant).

As discussed in Refs. [15,18], the beginning of the evaporation is possible when the nuggets are surrounded by an optically thin environment. The neutrino influx is then capable of powering the baryon evaporation at the nugget surface. The transition from the diffusive to transparent regime satisfies the condition $G_F^2T^4A^{1/3}\approx1$ (with G_F the Fermi con-

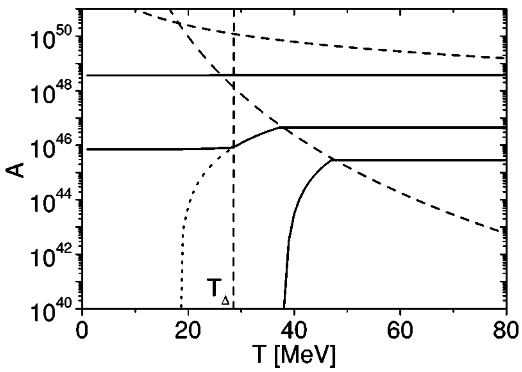

FIG. 3. The baryon number evolution of quark nuggets as a function of the temperature of the Universe for $\Delta=50$ MeV. The upper dashed line is the baryon number included within the horizon. The lower dashed line (given by $G_F^2 T^4 A^{1/3} \approx 1$) divides the regions of diffusive neutrino heating and neutrino transparency of the environment of the nuggets. There is almost no evaporation in the diffusive neutrino heating regime, simply because not enough energy is supplied to power the evaporation. The vertical line shows the temperature T_Δ below which pairing operates. The solid lines show the evolution of nuggets with three different initial baryon numbers. The heavier nugget reaches T_Δ even before leaving the diffusive neutrino heating regime, and therefore freezes out, retaining its initial baryon number. The lighter of these nuggets evaporates completely as soon as it enters the neutrino transparent regime. The intermediate mass nugget maintains its initial baryon number until it enters the neutrino transparent regime. Thereafter, it evaporates substantially until it reaches the critical temperature T_Δ, where it freezes out with a smaller mass. This nugget would have evaporated in the absence of pairing, as is apparent from the dotted curve corresponding to $\Delta=0$. These calculations assume $\alpha=10^{-3}$, which is much higher than the values given by chromoelectric flux tube models but smaller than the extreme value given by detailed balance.

stant). Evaporation is small before the nugget crosses this boundary curve. Thus, the temperature T_0 at which each nugget begins to evaporate is a function of the initial baryon content A_0, assumed to be the value it had at the formation temperature (see Figs. 3 and 4).

It can be checked that, depending on the value of the pairing gap Δ, a subset of nuggets that proceed directly from

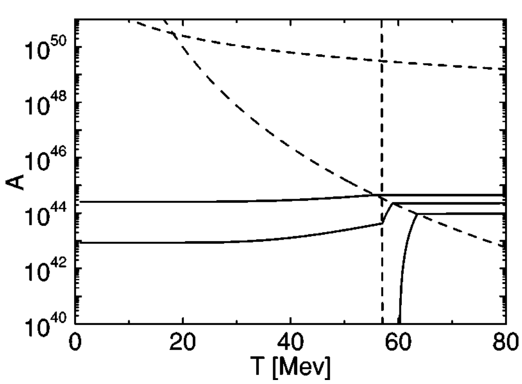

FIG. 4. The same as the previous figure but for $\Delta=100$ MeV. Lower masses can now reach the pairing temperature and freeze out.

the diffusive neutrino heating regime to the CFL paired state may exist. Therefore, these nuggets remain essentially frozen with the same baryon number they had at their formation. The minimum baryon number $A_{\rm freeze}$ that satisfies this condition is found inserting the relation $T_\Delta = 0.57\Delta$ in the transition condition above, and is given by

$$A_{\rm freeze} = 3.38\times 10^{44}\left(\frac{100\ \ \rm MeV}{\Delta}\right)^{12}. \tag{15}$$

Nuggets smaller than $A_{\rm freeze}$ will evaporate substantially once they enter the optically thin neutrino regime, but may survive if they manage to cool down to T_Δ with some finite mass.

While quarks remain unpaired, the evaporation rate will be given by Eq. (12). Therefore, for $T_\Delta < T < T_0$ the baryon number density as a function of temperature is given by

$$A^{1/3}(T) = A_0^{1/3} - C[\mathrm{Ei}(-I/T) - \mathrm{Ei}(-I/T_0)], \tag{16}$$

where Ei(x) is the exponential integral.

Once $T < T_\Delta$, pairing reduces the evaporation rate to the expression Eq. (13), and the baryon number density as a function of temperature follows:

$$A^{1/3}(T) = A_0^{1/3} - C\left[\mathrm{Ei}\left(\frac{-I-\frac{3}{2}\Delta}{T}\right) - \mathrm{Ei}\left(\frac{-I-\frac{3}{2}\Delta}{T_\Delta}\right) + \mathrm{Ei}(-I/T_\Delta) - \mathrm{Ei}(-I/T_0)\right]. \tag{17}$$

A simple approximation for Ei(x), which is good within a few percent in the range of interest, is the following:

$$\mathrm{Ei}(x) = \frac{\exp(x)}{x^2 - 2x}, \tag{18}$$

which is useful for understanding the relative weight of each term in the corresponding temperature regimes. The complete results are depicted in Figs. 3 and 4. The effects of CFL pairing are apparent when nuggets reach T_Δ, since many of them are able to survive while they would have been evaporated in the absence of this pairing. As a corollary, we may state quite generally that a given initial mass function of nuggets would be stretched towards the smallest masses because of CFL pairing. Detailed calculations of these features will be presented in a future publication. We finally stress that this evaporating population may not exist at all, depending on the form of the initial mass function.

IV. CONCLUSIONS

We have discussed in this work the effects of QCD pairing on the evaporation/boiling rates of quark nuggets assumed to be formed during the cosmological quark-hadron phase transition. These nuggets would be produced at $T_{\rm QCD} \sim 150$ MeV with maximum baryon numbers $A_{\rm max}$

$\sim 10^{49}(100\ \mathrm{MeV}/T_{\mathrm{QCD}})^2$ corresponding to the horizon scale at that epoch. After formation, the nuggets are fragile because of the hot environment and may boil and/or evaporate into hadrons. The nuggets may survive if their destruction is not complete when the Universe cools down to a sufficiently low temperature.

We have shown in this work that the consideration of pairing brings an additional twist to the problem of nugget survival at intermediate temperatures. Specifically, we have shown that both the boiling and the surface evaporation get suppressed because of the presence of the gap Δ in the respective rates.

Boiling of nuggets has already been discussed in the literature and was found to be unlikely in the most realistic calculations. When CFL pairing is included, the boiling is also unlikely because, in spite of the increase of the stability window, the rate is suppressed by Δ and the net effect produces $h_{\Delta}>0$ in realistic cases.

In the case of surface evaporation, the fate of the nuggets depends mainly on the (unknown) characteristics of the initial mass spectrum of the nuggets, their evaporation rate, and the value of the pairing gap. However, and independently of these uncertainties, many general trends can be noticed. If the value of the pairing gap Δ is sufficiently high, the nuggets perhaps as small as $\sim 10^{42}$ and up to $A_{\max}$ enter the CFL phase before leaving the regime that is opaque to neutrino transport. Since pairing quenches the rate by a large factor, all these nuggets freeze out with essentially the same baryon number they had at formation. In general, the net result is that many nuggets survive with smaller masses, which could not have otherwise survived if pairing had not operated. Therefore, any initial mass function of nuggets will be *stretched* towards the low-mass region after being partially evaporated. Note that this behavior is obtained for evaporating fluxes that may be many orders of magnitude larger than the very small values indicated by the chromoelectric flux tube models.

We conclude that the survival of the nuggets (if formed) is quite likely if they settle in a CFL state at a temperature $T_{\Delta}=57\times(\Delta/100\ \mathrm{MeV})$ MeV, which may be true for the whole population. Thus, CFL prevents further evaporation/boiling and effectively freezes out the masses of the nuggets. A detailed numerical study of the whole evolution of the nuggets is desirable to address this issue.

ACKNOWLEDGMENTS

The authors wish to thank the Instituto de Astronomia, Geofísica e Ciências Atmosféricas de São Paulo, the São Paulo State Agency FAPESP for financial support through grants, and to acknowledge the partial support of the CNPq (Brazil).

[1] K.A. Olive, Nucl. Phys. **B190**, 483 (1981).
[2] E. Suhonen, Phys. Lett. **119B**, 81 (1982).
[3] C.J. Hogan, Phys. Lett. **133B**, 172 (1983).
[4] E. Witten, Phys. Rev. D **30**, 272 (1984).
[5] M. Gyulassy, K. Kajantie, H. Kurki-Suonio, and L. McLerran, Nucl. Phys. **B237**, 477 (1984).
[6] J.H. Applegate and C.J. Hogan, Phys. Rev. D **31**, 3037 (1985).
[7] K. Kajantie and H. Kurki-Suonio, Phys. Rev. D **34**, 1719 (1986).
[8] G.M. Fuller, G.J. Mathews, and C.R. Alcock, Phys. Rev. D **37**, 1380 (1988).
[9] S.A. Bonometto and O. Pantano, Phys. Rep. **228**, 175 (1993).
[10] Z. Fodor and S.D. Katz, J. High Energy Phys. **03**, 014 (2002); Nucl. Phys. B (Proc. Suppl.) **106**, 441 (2002).
[11] K. Kanaya, Nucl. Phys. **A715**, 233 (2003); Z. Fodor, *ibid.* **A715**, 319 (2003).
[12] K. Jedamzik and G.M. Fuller, Nucl. Phys. **B441**, 215 (1995).
[13] K. Sumiyoshi, T. Kusaka, T. Kamio, and T. Kajino, Phys. Lett. B **225**, 10 (1989); K. Sumiyoshi, T. Kajino, G.J. Mathews, and C.R. Alcock, Phys. Rev. D **42**, 3963 (1990).
[14] B.K. Patra, V.J. Menon, and C.P. Singh, Nucl. Phys. **B564**, 145 (2000).
[15] C. Alcock and E. Farhi, Phys. Rev. D **32**, 1273 (1985).
[16] J. Madsen, H. Heiselberg, and K. Riisager, Phys. Rev. D **34**, 2947 (1986).
[17] C. Alcock and A. Olinto, Phys. Rev. D **39**, 1233 (1989).
[18] K. Sumiyoshi and T. Kajino, in *Proceedings of the International Workshop on Strange Quark Matter in Physics and Astrophysics*, edited by J. Madsen and P. Haensel [Nucl. Phys. B (Proc. Suppl.) **24**, 80 (1991)].
[19] J. Madsen, [18], p. 85.
[20] J. Madsen and M. Olesen, Phys. Rev. D **43**, 1069 (1991).
[21] M. Olesen and J. Madsen, Phys. Rev. D **47**, 2313 (1993).
[22] L. Masperi and M. Orsaria, astro-ph/0307347.
[23] A. Bhattacharyya, Jan-e Alam, S. Sarkar, P. Roy, B. Sinha, S. Raha, and Pijushpani Bhattacharjee, Phys. Rev. D **61**, 083509 (2000).
[24] G. Lugones and J. Horvath, Phys. Rev. D **66**, 074017 (2002).