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Strange Matter Lumps in the Early Universe

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We study the stability of hot strange matter lumps formed in the Early Universe. As the universe cools below the hadronization temperature, these lumps surround themselves with a hadronic crust to prevent boiling. Our analysis improves on previous finite temperature studies by correctly treating the constraints from chemical equilibrium of baryonic and electric charge between crust and core. Since the expansion rate of the Early Universe is much faster than the cooling rate of the strange matter lumps, their internal structure is characterized by constant entropy per baryon. We show that such lumps are stable against boiling only if their mass is much larger than the total mass inside the horizon. Hence stable strange stars cannot be created during hadronization in the Early Universe.

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As a possible consequence of a first-order phase transition from the quark-gluon to the hadron phase in the Early Universe, strange matter [1, 2] has drawn strong interest. Within a certain range of parameters allowed by QCD phenomenology strange matter, which consists of roughly equal numbers of up, down and strange quarks, is actually the true ground state of QCD. It was thus speculated that, if the nature of the quark-hadron transition is of first order, much of the baryon number of the early Universe could condense into lumps of strange matter during the quark-hadron phase transition, which might explain the dark matter problem.

As the Early Universe expands in the quark-gluon phase and cools below the transition temperature T_c , hadron bubbles form and grow. Due to the assumed first-order nature of the transition a large amount of latent heat is released which reheats the Universe back to T_c . Upon further expansion of the universe hadron bubbles continue to form and grow at T_c . As the volume in the quark phase decreases and the hadronic volume grows, quark bubbles remain inside the hadron matter. The quark bubbles keep shrinking until the latent heat cannot keep up the temperature decrease by expansion. Due to the rapid decrease of the temperature at this stage, quark bubbles with large baryon and strangeness number may decouple from the expansion and gravitationally condense into lumps of strange matter. If strange matter is indeed energetically more stable than normal nuclear matter [2], these strange matter lumps might survive and thus be candidates for the dark matter in our universe.

Two mechanisms can dissolve the strange matter: surface evaporation [3] and boiling [4]. It was claimed [4] that only strange matter with a baryon number larger than a critical value A_{crit} can survive the dissolution by surface evaporation and boiling. This critical value depends, however, strongly on poorly known parameters, such as the bag constant B , the surface tension of quark matter σ , and the temperature T of the Universe. Whether or not the minimum mass of strange matter which can survive boiling exceeds the causality limit (i.e. the total mass inside the horizon at the time of lump formation) is still a matter of debate [5], and we will discuss this point below.

Another interesting possibility pointed out in [6] is that the strange matter lump surrounds itself by a hadronic crust which is gravitationally bound to the quark matter core. This prevents evaporation from the surface of the strange matter lump directly into the surrounding vacuum. Once chemical and gravitational equilibrium between the core and the crust is reached, also the boiling process stops, since the gravitational pressure from the surrounding hadron shell balances the high pressure of the hot strange matter. The authors of [6] found gravitationally stable strange matter lumps with a hadronic crust, but the total baryon number of these objects turned out to be larger than the causality limit already mentioned. They concluded that strange matter lumps with a hadronic crust are unlikely to form in the Early Universe.

Unfortunately, in [6] the phase boundary between the quark core and the hadronic

crust was not treated correctly. As has been known now for several years from studies of the hadronization process in heavy-ion collisions [7, 8] and was recently stressed again in the context of cold hybrid stars (neutron stars with a quark core) [9], the treatment of the phase boundary and the mixed phase between core and crust requires special care if the hadronization process is constrained by more than one conserved charge. In the case of nuclear collisions the relevant conservation laws are those of baryon number and strangeness, while charge conservation plays energetically a minor role and can be neglected. In the case of strange stars and the dynamics of the Early Universe the typical time scales are much longer, such that strangeness-violating weak interaction processes can equilibrate and strangeness is no longer conserved; however, compared to the weak gravitational interaction electric forces play now a crucial role, and the conservation of electric charge across the phase boundary must be correctly taken into account. This was not properly done in Ref. [6].

The interesting phenomenon which occurs in this case if the transition is of first order is the “distillation” of strangeness in the case of nuclear collisions [7, 8, 10] and of electric charge in the case of neutron stars [9]. The core and the crust are separated by a mixed phase in which a fraction of the volume is in one phase with the remainder being in the other phase. When averaged over volumes large compared to the typical strong interaction scale of 1 fm, the overall baryon density in the mixed phase is homogeneous (unless the surface tension of the quark matter is too large [11]) while its charge density vanishes. The degree of hadronization at any given point in the mixed phase is parametrized in terms of the equilibrium hadronic volume fraction $f = V_h/V_{\text{tot}}$. Charge “distillation” means that, while at each point the mixed phase is on average charge neutral, the net charge in each subphase does not vanish, but that on overall charge neutrality is obtained by balancing the negative charge density in one of the phases (whose value depends on f) against the positive charge density in the other subphase. This phenomenon leads to a continuous variation of the chemical potentials for the conserved quantum numbers across the mixed phase, and as a result all other thermodynamic quantities also change smoothly across the phase transition region.

If instead charge neutrality is (wrongly) postulated in each subphase separately (as done in [6] and most other prior work on quark stars (e.g. [12–14])), there is a discontinuity in the charge chemical potential (in violation of the Gibbs equilibrium condition), resulting in discontinuities of most other thermodynamic quantities except for the pressure which is forced to be the same in both subphases. The consequence is discontinuous energy and baryon density profiles of the star. A correct treatment of chemical equilibrium in the mixed phase leads to continuous star profiles with a mixed phase of finite thickness, across which the pressure is no longer constant, but rises continuously towards the interior of the star. While this has been studied carefully in [9] for cold hybrid stars, we extend that work here for strange hybrid stars at high interior temperature, thus correcting the work of [6] and checking their conclusions.

Strange matter lumps with a hadronic crust which had survived until today would

be cold (due to heat conduction via neutrino emission), with a homogeneous temperature equal to that of the cosmic background radiation. Their stability against gravitational collapse has been studied in [9]. It is, however, interesting to ask whether the same stability criteria apply at higher temperatures, and whether it is at all possible to create such a cold strange hybrid star by cooling a primordial hot strange quark matter lump. In this dynamical situation constancy of the temperature across the star can in general not be maintained since, at least shortly after the assumed formation during the cosmic hadronization phase transition, the rate of expansion and cooling of the universe is presumably faster than the cooling rate of the star by neutrino emission [15]. In this case, the star stays hotter than its surrounding medium, and its internal temperature profile is determined by the laws of entropy and baryon number conservation during hydrodynamic gravitational contraction. It thus appears reasonable to assume that the structure of the hot strange quark matter lump is determined by the condition of constant specific entropy (i.e. entropy per baryon) across the star [16]. We will thus study strange stars with a quark core at finite temperature and their change in stability as they cool down, with special emphasis on a correct treatment of the mixed phase.

The paper is organized as follows: In Section II A possible scenarios for the formation of strange matter lumps will be reviewed. Simulation of the phase boundary through the mixed phase will be discussed in Section II B. In Section II C we study the influence of the equation of state on the temperature profile in a star with non-zero specific entropy. In Section II D we discuss the solution of the Oppenheimer-Volkoff equations at finite temperature. We present our results in Section III, and give a short summary in Section IV.

II. STRANGE MATTER LUMPS WITH HADRONIC CRUST

In this section we will investigate various aspects of strange matter lumps with a hadronic crust, such as their formation, the proper equilibrium treatment of the mixed phase layer, and their temperature profile.

Prior to that, let us make a few comments on the equation of state (EoS) employed. At finite temperature, where a large number of resonances are excited in the hadronic phase, evaluating the EoS is the critical factor in the solution of the Oppenheimer-Volkoff equations for the star’s structure. We have used a non-interacting gas of quarks, gluons, photons and leptons in thermal and chemical equilibrium above the phase transition. Photons and leptons are common elements of all the phases of the star as well as of the background radiation from the universe surrounding the star. As usual, deconfinement is parametrized in terms of a bag constant B . For the hadronic phase we assume a mixture of the known mesonic and baryonic resonances with hard core repulsion [17, 8] which are in thermodynamic equilibrium with the photons and leptons (for details see [6]). It is known (see [9] and references therein) that this is not sufficient to ensure the proper saturation and binding properties of cold nuclear matter in its ground state. To improve this one should

add other types of interactions in the hadronic phase, for example in the framework of relativistic mean field theory [9]. However, since we are mostly interested in the structure of stars at very high temperature (near T_c), these fine details in the binding properties should not be too important; they should be kept in mind, however, when comparing our results at low temperatures with other work on cold strange stars.

We formulate all thermodynamic quantities as functions of the temperature T , the baryon chemical potential μ_b , and the charge chemical potential μ_c , which is defined as the negative of electron chemical potential, $\mu_c = -\mu_e$. Two chemical potentials are needed to parametrize baryon number conservation and charge neutrality. The chemical potentials for all the particles in each phase can be expressed in terms of these two chemical potentials via their baryonic and electric charges, b_i and q_i , respectively:

$$\mu_i = b_i \mu_b + q_i \mu_c. \quad (1)$$

A. The formation of strange matter lumps

Above T_{c0} , thermally equilibrated photons, gluons, electrons, neutrinos, quarks, and their antiparticles form the constituents of the Early Universe. We will distinguish T_{c0} , which is the hadronization temperature of the (essentially) baryon-free Universe, from the generic transition temperature at finite baryon density which we denote by T_c , and which decreases with baryon chemical potential. Since T_{c0} is of the same order or larger than the strange current quark mass (~ 150 MeV), the numbers of up, down and strange quarks are roughly equal.

The mechanism of hadronic bubble formation by supercooling in a first order phase transition has already been described in the Introduction. As more and more matter converts to the hadronic phase the picture reverses into one where quark bubbles are immersed in the hadron gas. It is a well-known fact [1, 18] that (due to the specific form of the entropy balance between the two phases in the early universe where the photons and leptons play an important role) the baryon density contrast gets strongly enhanced by this process, i.e. the baryon density in the remaining quark bubbles is much larger than in the hadronic environment. A large surface tension or gravitational collapse may further increase the baryon density and thus reduce the critical temperature for hadronization. Such regions of large baryon density would thus hadronize much later than the rest of the universe. Since the expansion of the universe and the velocity with which hadron bubbles grow and quark bubbles shrink are all slow compared to the weak interaction time scale, strangeness can always adjust to its equilibrium abundance (roughly equal to the up and down abundance), thus providing an entrance channel to stable strange quark matter droplets which will no longer hadronize at all.

Due to the density inhomogeneities caused by the violent phenomena of growing and colliding bubbles it is quite likely that some of these remaining quark bubbles

undergo gravitational contraction. They will then accumulate a crust of hadrons on their surface which is bound by gravitation. At this point their temperature will still be of the order of the transition temperature. One of the questions which we will address concerns the conditions of stability for such a selfbound very hot object, and how these conditions change as the universe and thus (at a slower rate) also the strange quark star cools down.

B. The mixed phase layer

As one heats a medium with a first order phase transition at constant pressure, or as one increases the pressure at constant temperature, the medium goes through a mixed phase where part of the volume (say, a fraction f) is in one and the rest of the volume (a fraction $1-f$) is in the other phase. Everyday examples, like a glass of water at 100° C covered with a tight lid, suggest that under the influence of gravitation phase separation occurs, i.e. the denser phase settles at the bottom and the more dilute phase rises to the top, leaving a sharp surface in the middle across which the density is discontinuous.

This is in fact what happens if there is only one conserved quantum number (like the number of water molecules) characterizing the system. In this case at fixed temperature the pressure remains constant across the phase transition, i.e. for all volume fractions f the thermodynamic pressure is the same, and thus the gravitational pressure determines the spatial configuration of the mixed system. If there are several conserved quantum numbers, even at fixed temperature the pressure is a function of the volume fraction f . This was first discussed in the context of relativistic nuclear collisions, where the hadronization phase transition is constrained by the conservation of baryon number and strangeness [6], and recently stressed again in the context of neutron stars with quark core [9] where the two conserved quantum numbers are baryon number and electric charge. Since the interior structure of a star is determined by the condition of mechanical stability, i.e. the balance between the thermodynamic and gravitational pressure, and since the latter is a function of distance r from the center (as it is caused by the total mass $m(r)$ inside the radius r), different points of the mixed phase with different volume fractions f will be able to satisfy the pressure balance condition at different values of r . This means that no longer all the dense phase settles to the bottom with the dilute phase on top and a discontinuity at r_c in between (which would imply $f = 0$ for all $r < r_c$ and $f = 1$ for all $r > r_c$). Rather there will be no density discontinuity at all, instead the density varies smoothly from its lower value n_H in the pure dilute (hadronic) phase at large r to its higher value n_Q in the pure dense (quark) phase at small r , with

$$n(r) = f(r) n_H + (1 - f(r)) n_Q \quad (2)$$

in between. Here $f(r)$ is the fractional volume of the hadronic phase as a function of r in the mixed phase region. (It should be noted that there is in addition an implicit

r -dependence through the densities $n_{H/Q}(T(\tau), \mu_b(\tau), \mu_c(\tau))$ via their dependence on the temperature and chemical potentials whose radial profiles are dictated by the Oppenheimer-Volkoff equations, see below.)

Therefore stars with two or more conserved charges possess an extended mixed phase layer [9] in which the quark and hadronic phases coexist. The microscopic structure of this mixed phase (bubbles, rods, sheets, etc.) depends on the surface tension between the two phases and requires more detailed considerations [11]. In [11] it was pointed out that if the surface tension is too large and begins to dominate the free energy near the phase transition (and thus the Gibbs equilibrium conditions) the mixed phase may disappear again and the density discontinuity between the quark core and hadron crust may be recovered. Recent results from lattice QCD [19], however, give quite small values for the quark-hadron surface tension and thus do not support such a picture.

In the mixed phase all extensive thermodynamic quantities (energy, entropy, baryon, and charge density) vary continuously and are given by expressions analogous to Eq. (2). For each value of f the critical pressure can be calculated from the Gibbs conditions for two-phase equilibrium, namely $P_Q = P_H$, $T_Q = T_H$, $\mu_b Q = \mu_b H$ and $\mu_c Q = \mu_c H$, or

$$P_Q(T, \mu_b, \mu_c) = P_H(T, \mu_b, \mu_c). \quad (3)$$

The charge neutrality constraint takes the form

$$n_c(T, \mu_b, \mu_c) = f n_{c,H}(T, \mu_b, \mu_c) + (1 - f) n_{c,Q}(T, \mu_b, \mu_c) = 0. \quad (4)$$

These two conditions allow to determine for each value of f the two chemical potentials as a function of T . Please note that Eq. (4) in general requires nonvanishing, but compensating charge densities in the two subphases; in the context of heavy-ion collisions this phenomenon applies to the strangeness quantum number and is known as “strangeness separation” in the mixed phase [7, 8]; together with kaon evaporation from the hadronic subphase, it may give rise to “strangeness distillation” and subsequent strangelet formation [10] in those collisions.

C. Specific entropy and the temperature profile

As described above we assume that the strange matter lumps with hadronic crust are formed by gravitational contraction of a density perturbation in the early universe near T_c . The matter inside the region which is about to contract is characterized by a certain net baryon number A and entropy content S and thus by a certain specific entropy S/A . It is reasonable to assume that initially the specific entropy is nearly homogeneous, i.e. that the density fluctuation was created adiabatically. Since the microscopic processes leading to local thermal and chemical equilibrium are governed by the strong, electromagnetic and weak interactions and thus occur with a much faster rate than the gravitational contraction, the matter will always remain very

close to local thermodynamic equilibrium. Thus the contraction is governed by the laws of ideal hydrodynamics, and the specific entropy will thus remain constant and homogeneous.

For a system of ultrarelativistic particles the specific entropy S/A is a function only of the ratio μ_b/T between the baryon chemical potential and the temperature. Constant S/A thus implies constant μ_b/T . Since μ_b tracks the density profile, $n_b \sim \mu_b^3(1 + c(\pi T/\mu_b)^2)$, so will the temperature. Thus the initial configuration will be characterized by a dense and hot interior of quark matter, surrounded by a colder, less dense hadronic crust with a temperature roughly equal to that of the surrounding universe, with constant S/A throughout the interior of the “star”.

As the universe continues to expand and cool down, the “star” will try to follow by radiating heat. The most efficient cooling mechanism is neutrino radiation. Since the hot interior will radiate with the largest rate, this mechanism tends to reduce the temperature gradients across the star. However, initially neutrino heat conductivity is very small: Appelgate and Hogan estimated the mean free path of neutrinos as $\lambda_\nu \simeq 10$ cm $(T/100 \text{ MeV})^{-5}$ [15]. In the hot core region the mean free path of neutrinos and thus the temperature difference between layers of distance λ_ν is small. Hence heat conduction through neutrino emission cannot be very efficient in reducing the temperature gradients there. In the surface region the neutrino mean free path is much larger: for $T = 20$ MeV, $\lambda_\nu = 310$ m, and for $T = 10$ MeV, $\lambda_\nu = 10$ km. But there the energy density of neutrinos and thus their total emission rate (which drops like T^4) is small. We conclude that the thermal evolution of the star must decouple from that of the expanding universe. We will therefore in our calculations of the star’s structure assume that to good accuracy the outside temperature (i.e. the temperature of the heat bath surrounding the star) is zero.

The further dynamics of entropy loss by neutrino emission is complicated, since it requires knowledge of the mechanisms which transport the loss of specific entropy from the surface (which is transparent to neutrinos) into the neutrino-opaque interior. Heat conduction by weak interactions and convection are expected to play the dominant roles. We will not enter into a detailed discussion of these issues here, but simply assume that the star is able to maintain a constant specific entropy profile while cooling down.

From these considerations, it is clear that the density profile of the star presents a cut through the hadronic equation of state and through the phase diagram along a trajectory of fixed S/A . Such isentropes are shown in Fig. 1 for various values of S/A for $B^{1/4} = 150$ MeV. Please note that for small values of S/A the latent heat from the first order phase transition reheats the system, while for $S/A > 10$ the temperature decreases during the hadronization phase transition. Correspondingly the baryon chemical potential μ_b increases during hadronization, if S/A is large, and decreases if S/A is small. The difference relative to the situation in heavy-ion collisions, where μ_b always decreases along isentropes (i.e. the system reheats during hadronization) is caused [18] by the contribution of leptons and photons to the entropy balance; due to

their small interaction cross sections, these particles are decoupled from the relatively small nuclear fireballs in heavy-ion collisions, but remain in thermal equilibrium and thus contribute to the pressure balance in the case of a quark-hadron star.

The characteristic features of the isentropic trajectories through the phase diagram shown in Fig. 1 will be reflected in the density and temperature profiles of the star. This will be further discussed in Section III. One crucial observation, however, needs to be made already at this point. The temperature profiles at the edge of the star are determined by the low-temperature limit of the isentropes. We see that for small values of S/A the isentropes in the T - μ -plane approach a value $\mu_b(T=0) > m_N$ from above, while for large S/A the baryon chemical potential approaches the nucleon mass from below; there is a critical value of S/A where in the mixed and hadronic phase μ_b stays more or less constant and close to the nucleon mass.

Since in the region $\mu_b < m_N$ no nucleons (and thus no hadronic matter at all) can exist at $T=0$, at low but non-zero temperatures this domain is characterized by the approximate absence of heavy baryons altogether and by the dominance of leptons, photons and light mesons. Thus in the region $\mu_b < m_N$, T small, the equation of state is essentially that of radiation of massless particles. For $\mu_b > m_N$ baryons can exist at low temperatures in appreciable densities, and the equation of state at low temperatures is that of a gas of massive particles.

For very small values of the bag constant, $B < 58 \text{ MeV/fm}^3$ ($B^{1/4} < 145 \text{ MeV}$), both strange and non-strange cold quark matter are stable against hadronization since in this case the baryon chemical potential $\mu_b(T=0) < m_N$ at the phase boundary. The finite temperature hadronic “crust” around the quark core of a strange quark star will then always be approximately of radiation type.

This observation has an important consequence for hot strange matter lumps with a hadronic crust: As is well known [20], radiation dominated pressure results in supermassive stars with very large radii. In the limit of pure massless particle radiation the mass and radius of the star diverges because the Oppenheimer-Volkoff solution for the density profile at large r decays like a power rather than exponentially. Thus, if S/A is large and/or B is small such that the hadronic skin of a hot strange star possesses a radiation-like equation of state, this skin will become very thick and spread to very large radii, leading to superheavy stars with very large masses and radii. In fact, we will see that already for rather modest values of S/A the radii of such stars are much larger than the horizon radius at the hadronization phase transition.

D. Integration of the Oppenheimer-Volkoff equations

The profile and its stability of any “star”-like object bound by gravitation can be obtained by solving the Oppenheimer-Volkoff equation:

$$\frac{dP(r)}{dr} = -\frac{GM(r)\rho(r)}{r^2} \left(1 + \frac{P(r)}{\rho(r)}\right) \left(1 + \frac{4\pi r^3 P(r)}{M(r)}\right) \left(1 - \frac{2GM(r)}{r}\right)^{-1}, \quad (5)$$

where

$$M(r) = 4\pi \int_0^r \rho(r') r'^2 dr'. \quad (6)$$

The energy density ρ is related to the pressure P through the EoS, which in our representation are functions of the three parameters, T , μ_b and μ_c . Since the constraints from charge neutrality and constant entropy per baryon throughout the star,

$$n_b = 0, \quad S/A = \text{const.}, \quad (7)$$

reduce the number of parameters by two, there remains only one parameter to be determined from Eq. (5) as a function of the radial distance, r . Alternatively, one can solve Eq. (5) directly for $P(r)$ as a function of r , after tabulating ρ as a function of P from the EoS subject to the constraints Eq. (6), taking care that in the phase transition region the charge neutrality condition is solved according to Eq. (4). This tabulation of the EoS $P(\rho)$ turns out to be the most time consuming part of the calculation but it needs to be done only once for all stars with the same S/A .

We have solved Eq. 5 for $P(r)$ using the Runge-Kutta method. The mass $M(r)$ inside the radial distance r is evaluated from Eq. (6) by the Simpson method as we integrate eqs. (5, 6) inside out over r .

III. RESULTS

We have studied strange matter lumps at a variety of values for the specific entropy and investigated their mass-radius relation and their stability against gravitational collapse. For different masses in the gravitationally stable domain we have analyzed their density and temperature profiles and established whether or not they contain a quark or mixed phase core. By following lumps of a given mass from high to low specific entropy we can study how the conditions of stability change as the star cools down (we neglect matter loss during cooling). When solving the Oppenheimer-Volkoff equations, we only consider the bulk energy of the matter, neglecting possible microstructures and the surface and Coulomb energies connected with them [11].

In Fig. 2 we plot the mass of hot strange matter lumps against their central energy density, for various values of the specific entropy S/A , choosing for the bag constant the value $B = 150 \text{ MeV/fm}^3$. The background temperature outside the star is taken to be zero. Stability against gravitational collapse requires the lump to sit on those parts of the curve where the slope is positive [20]. From the figure we see that for this bag constant stable configurations correspond to stars which have either very small specific entropy S/A (low temperature) or are just at the verge of forming a quark core. (The region of central densities which corresponds to a mixed phase at the origin is clearly recognizable by the nearly horizontal sections in the mass-density-plot.) When S/A becomes much larger than 10, even the latter configurations become unstable. For subcritical central densities, i.e. purely hadronic strange hot stars, at large S/A all solutions are unstable: they correspond to unstable supermassive stars

with an approximately radiation-type equation of state (see the discussion of Fig. 1 above). Only for small S/A the equation of state of the hadronic crust is matter dominated ($\mu_b > m_N$); in this case we find stable solutions with and without quark cores, and (for this value of B) their maximum mass is rather conventional, roughly 1.4 times the solar mass.

It is interesting to observe that stars with a mixed phase at the origin appear to be generically unstable: such stars correspond to the rather flat section in the $M - \rho$ diagram, Fig. 2, but a closer look shows that this section actually has a slightly negative slope. This point will be further discussed below. As the central density exceeds the threshold for a pure quark phase, the slope becomes even more negative (i.e. the solution is even less stable); only for central densities well above the pure quark matter threshold the lumps become stable again. Thus the two stable branches for small S/A in Fig. 2 correspond to stars which either consist only of hadronic matter, or have a large quark matter core followed by a comparatively thin mixed phase layer and a hadronic crust on the outside.

These characteristics of stability of hot strange matter stars do not change qualitatively as the bag constant is changed. This is shown in Fig. 3, where we plot the mass against the central energy density for stars with $S/A=1$ and $S/A=2$, for various values of the bag constant. The maximum mass of hot stars with a quark matter core decreases as the bag constant increases, while it does not change appreciably with S/A in the considered range. [We notice that the maximum mass of purely hadronic stars without a quark core appears to stabilize as B is increased: this behaviour reflects our approximate treatment of the hadronic proper volume in the hadronic equation of state (by multiplying all thermodynamic quantities by a factor $1/(1 + p^*/4B)$, see [17]), and should thus be taken with a grain of salt.]

The generic instability of stars with a mixed phase at the origin is true even for stars at zero temperature, for all values of the bag constant; this is shown in Fig. 4. This observation seems to contradict the zero temperature results obtained by Glendenning [9] using a relativistic mean field equation of state for the hadron phase. He found stars with a well-developed mixed phase center extending to radii of several kilometers. A characteristic feature of his calculations is a very strong increase of the pressure during the mixed phase (i.e. a strong dependence of P on the volume fraction f), which we are unable to reproduce with our equation of state for any set of parameters. This feature appears to be crucially influenced by the symmetry and Coulomb energy contributions to the equation of state, which in our case are only due to the increase in Fermi energy for unequal quark flavor occupancies (as encountered in the electrically charged subphases of the mixed phase), whereas in the case of Ref. [9] the symmetry energy is strongly enhanced by virtual vector meson exchange whose coupling was adjusted to reproduce the proper saturation properties of cold nuclear matter in its ground state [21]. Since these effects become less crucial at high temperature, our conclusion that the window of stability for stars with a mixed phase core closes rapidly for finite values of S/A appears to be safe.

We can also study the change in the stability of strange matter lumps under the influence of a nonvanishing background pressure due to the temperature of the surrounding universe. In Fig. 5 we show the analogue of Fig. 2, for a lump with specific entropy $S/A=100$, for various values of the temperature of the environment. In these calculations the surface of the star, i.e. its radius R , is determined by pressure equilibrium between the hot hadronic crust matter and the surrounding radiation, i.e. we stop the integration of the Oppenheimer-Volkoff equations when the pressure has dropped to the outside radiation pressure. Since the surrounding radiation has vanishing baryon density, S/A and n_b are discontinuous at $r = R$. While this crude approximation could, of course, be improved, it is sufficient for our qualitative stability analysis.

As the universe expands and cools and thus the background temperature decreases, the star may readjust its density profile to the changing external pressure conditions. How will this change its gravitational stability? While its total mass will in general change as a consequence of the relaxation of the density profile (even if S/A does not change on the cooling time scale of the expanding universe), its baryon number will remain constant. Thus the fate of a stable hot strange matter lump in the once hot universe under the influence of a decreasing external radiation pressure should be obtained from Fig. 5 by following points with equal total baryon number. It turns out, however, that the changes in gravitational binding due to relaxation of the external pressure are small, and thus for a qualitative discussion it is sufficient to follow the points with equal total mass in Fig. 5. In a very hot universe with $T > 100$ MeV stable hot stars with a quark matter core have a maximum mass of the order of one solar mass; as the temperature of the surrounding universe drops, the typical mass of a hot strange star with a quark core at the same S/A is nearly a factor 10 larger. Thus primordial quark stars will not keep their quark core as the universe cools down; as the external radiation pressure decreases, they expand and hadronize, first moving to the steeply rising hadronic branch on the left of the curves in Fig. 5 and then completely evaporating. They could remain stable gravitationally bound objects only if they found some mechanism to cool down very rapidly, reducing their specific entropy S/A at a similar rate as the expansion rate of the universe.

In Fig. 6a-d we show the internal structure of a typical strange matter lump with $S/A = 2$ for $B^{1/4} = 150$ MeV, with a mass of $1.4 M_\odot$ and a radius of 10.5 km. Please note the rather thin mixed phase layer of only 0.5 km thickness. As discussed in Section II B, the mixed phase would shrink to zero width and the density would be discontinuous at this point if we did not correctly include the simultaneous conservation of baryonic and electric charge near the interface. The occurrence of a finite width mixed phase is entirely due to the Gibbs requirement of continuity for the chemical potentials across the phase transition.

The spike in temperature distribution in Fig. 6d is due to the release of latent heat in the hadronization transition, as discussed in Section II C. As follows also from this discussion, for stars with much larger values of S/A this temperature spike

in the mixed phase goes away and turns into a steep drop of T in the mixed phase layer.

Supermassive stars of very large specific entropy ($S/A > 100$, i.e. very hot quark stars), have a core of quark matter and a mixed phase layer of similar size as in Fig. 6, but are surrounded by an enormously thick hadronic skin with a nearly radiation-type equation of state, which also contains most of their total mass. Their radii are typically hundreds to thousands of kilometers, i.e. much larger than the horizon radius of the universe at the temperature when they were supposed to be formed. This constitutes an unsurmountable obstacle for their creation by gravitational collapse of a density fluctuation in the early universe near the hadronization temperature, because the gravitational instability signal leading to collapse cannot reach all the matter needed to form the stable star.

IV. SUMMARY AND CONCLUSIONS

We studied the gravitational stability of hot strange matter lumps formed in the Early Universe. Proper treatment of the charge neutrality condition through the mixed phase results in continuous pressure, energy and baryon density profiles within the mixed phase which covers a finite but narrow (~ 0.5 km for $B = 150$ MeV/fm³) radial width. Assuming constant entropy per baryon S/A inside the star to fix the temperature profile, we obtain a large variety of gravitationally stable hot strange matter lumps with a quark matter core. The maximum mass of such stars depends critically on their specific entropy S/A (i.e. on their central temperature); for large values of S/A their hadronic skin is characterized by a nearly radiation-type equation of state, and this skin becomes very thick and massive, turning the star into a supermassive hot giant. For small values of S/A , say below $S/A = 5$, the maximum mass of the star depends only very weakly on S/A (i.e. upon whether it is cold or hot) and decreases from $1.7 M_\odot$ to $1.2 M_\odot$ as the bag constant is increased from $B = 100$ MeV/fm³ to $B = 200$ MeV/fm³.

Contrary to the case of cold hybrid neutron stars studied by others, we find that with our model EoS strange matter lumps with a mixed phase core are generically unstable. The difference seems to be due to the absence of Coulomb and symmetry energy contribution in our EoS. While at zero temperature these contributions to the EoS are appreciable and should be treated correctly (which we have not done here), their relative contribution is less crucial at high temperatures ($T > 10$ MeV) which are of interest here. We therefore think that our conclusions about the instability of stars with a mixed phase core are safe at high temperatures.

By varying the background temperature we see that stable primordial strange matter lumps with a quark matter core and a large specific entropy S/A needs to accumulate an enormous hadronic surface to remain stable as the universe cools. We do not see a reasonable mechanism how this can occur without violating causality, and thus we rather think that such primordial stars (if they were ever formed at all) hadronize and evaporate as the universe cools down.

As a further problem we find that the total baryon number of stable strange quark matter lumps at any temperature $T \leq T_c$ is in the range of 10^{46} to 10^{47} and thus larger than the total mass of the universe inside the horizon at T_c . At the point of hadronization the universe was a few microseconds old, with a horizon radius of only about 1 km. This is much smaller than the radius of the objects we have calculated, i.e. our stars cannot be causally connected, and the assumption that they are bound by gravitation (which requires gravitational signals to travel across the star) is inconsistent. The total net baryon number contained inside the horizon can be estimated in a homogeneous universe during the radiation dominated era as [15]

$$A_{horizon} = 10^{46} \left(\frac{100 \text{ MeV}}{T} \right)^3. \quad (8)$$

At $T = T_c$ this is about 8 orders of magnitude smaller than the baryon number needed to make one of the stable objects discussed in this paper. Even in the most violent imaginable hadronization scenarios it is very unlikely that a baryon density fluctuation of this order of magnitude could be created, allowing all of the baryons to communicate with each other and begin to gravitationally contract into a gravitationally stable hot strange quark matter star.

Strange quark matter of smaller size needs a nongravitational stabilization mechanism. Strange quark matter droplets without a hadronic crust, which are stabilized by strong interactions, are still under study [22].

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FIGURES

- FIG. 1. Phase diagrams for $B = 150 \text{ MeV/fm}^3$ with trajectories of $S/A=\text{const.}$: (a) phase diagram in (T, ρ_b) plane, (b) phase diagram in (T, μ_b) plane.
- FIG. 2. Mass vs. central energy density plot with $B = 150 \text{ MeV/fm}^3$ for various values of specific entropy, S/A . Strange matter lumps are stable only when they lie in the region of positive slopes in this figure.
- FIG. 3. Bag constant dependence of stability: same plot as in Fig. 2 for various values of bag constants by fixing $S/A = 2$.
- FIG. 4. Mass vs. central energy density plot of cold strange matter lumps at $T = 0$ for various bag constants.
- FIG. 5. Same as in Fig. 2 with different temperature environment. Note a change of stability of strange matter lumps as the universe expands.
- FIG. 6. Profiles of a strange matter lump with hadronic crust for $B = 150 \text{ MeV/fm}^3$ and $S/A=2$ as a function of radial distance from the core: (a) energy density, (b) pressure, (c) baryon density, and (d) temperature vs. r . The region of mixed phase spreads over a finite radial interval and energy and baryon densities are continuous functions of r . The temperature spike in the mixed phase region is due to the large latent heat from the first order phase transition. However, this peak exists only for stars with small S/A .

Fig. 1a

Fig. 1b

Fig. 2

Fig.4

Fig.3

Fig. 5

Fig. 6a

Fig. 6b