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Stability of strange quark matter at $T\neq 0$

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The possibility of stable strange quark matter (both bulk and quasibulk) at finite temperature and some of its gross properties have been investigated with the dynamical density-dependent quark mass model of confinement. The possibility of metastable and/or stable strangelet formation in ultrarelativistic heavy-ion collisions has also been discussed.

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I. INTRODUCTION

It was argued by Witten [1] that flavor-symmetric bulk strange quark matter (SQM) may be an absolute ground state of matter near nuclear density at zero temperature and zero pressure. This speculation was investigated [2–4] using the most popular phenomenological bag model results [5]. With the same model, the stability of strange droplets and their gross properties have also been studied [6]. Recently the stability of SQM at finite temperature has got some special interest [7,8].

It is generally expected that if the density of a supernova remnant (the neutron star) is sufficiently high, neutrons may overlap and form a large cluster of quark matter. Since SQM is energetically favorable over nonstrange quark matter (QM), the whole star will be converted to a strange star by the weak decays (and also reactions) of u and d quarks. The temperature of the star may be a few MeV [9–11].

During the quark-hadron phase transition in the early Universe, which is assumed to be first order (or at least weak first order) in nature, quark matter may be trapped within the infinite network of hadron bubbles. This trapped quark matter finally collapses to form a large cluster of quark matter or quark nugget and subsequently by the weak decay of light quarks a bulk structure of SQM will be produced [1].

Recently it has also been argued that if a baryon-rich quark-gluon plasma is formed in ultrarelativistic heavy-ion collisions [12], strangelets might play an important role. In this particular case an equal number of s and $\bar{s}$ is created by the reactions $q\bar{q}\rightarrow s\bar{s}$ and $gg\rightarrow s\bar{s}$ where $(\bar{q})$ q stands for light (anti)quarks and g stands for gluon. These are strong interaction processes and further creation of u and d quarks (by pair production of gluons) are suppressed or forbidden by the Pauli exclusion principle. The produced antistrange quarks will combine with u and d quarks and give rise to K^+ and K^0 mesons. They will come out of the hot and baryon-rich quark-gluon plasma (QGP) bubble. Since the abundances of $\bar{u}$ and $\bar{d}$ quarks are very low (due to Pauli suppression), the emission of K^- and $\bar{K}^0$ will also be very rare. As a result the strangeness fraction in the QGP bubble will increase continuously with time and finally, depending on the density, temperature and strangeness content of the system, the QGP bubble (which is now strangelet or strange droplet) may become a metastable or stable system. It is therefore interesting to study the stability of both bulk as well as quasibulk SQM at finite temperature both from the points of view of cosmology or/and astrophysics and ultrarelativistic heavy-ion collisions.

In a series of papers [11,13–16], we have investigated the stabilities of both bulk and quasibulk SQM at zero temperature and the results obtained differ significantly from those of the previous authors [2–4]. In our studies, a dynamical density-dependent quark mass model of confinement was introduced, which is one of the best models of quark confinement. According to this model the mass of a quark is extremely small inside a hadron and is infinitely large outside [13]. The predictions of this model are more or less consistent with the lattice results. On the other hand, the phenomenological bag model *a priori* assumes that, within the bag, quarks are asymptotically free. But the recent results [17] from lattice calculations show that quark matter does not become asymptotically free immediately after the phase transition, even if it is first order in nature; it approaches the free gas equation of state rather slowly; some hadronic degrees of freedom remain within quark matter immediately after the phase transition; in this context the bag model is thus an inadequate description of confinement. In our previous studies we have obtained a stable configuration of noninteracting bulk SQM at densities around 7–$8n_0$ ($n_0\sim 0.17$ fm^{-3}, the normal nuclear density) [13]. When an interaction is switched on among the quarks (which was treated by the relativistic version of the Landau theory of a Fermi liquid [15,18]) the stability point was shifted to $5n_0$ [15]; but in both the cases the internal kinetic pressures remain nonzero, which again contradicts Witten's speculation. This model also predicts the stability of strangelets at zero temperature and the stable density is $\sim 2n_0$ [14]. The stability point (density) is shifted due to the surface effect. We have also studied, with the same model of confinement, the gravitational stability and some gross properties of strange stars [15]. We have reproduced some of the expected parameters for submillisecond pulsars. The density of a strange star (from core to surface) is again very close to the density range within which bulk

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SQM is stable, which is not the case with the bag model of confinement.

In this paper we shall investigate the conditions of stability and some global properties of bulk and quasibulk SQM at nonzero temperature. A modified form of the dynamical density-dependent quark mass model of confinement at $T\neq 0$ will be used [11].

The paper is organized in the following manner. In Secs. II and III we have investigated the stability and some global properties of bulk and quasibulk strange quark matter. Section IV deals with the possibility of strangelet formation in ultrarelativistic heavy-ion collisions, and Sec. V contains the discussions and conclusions.

II. STABILITY OF BULK STRANGE QUARK MATTER AT $T\neq 0$

Let n_B be the baryon number density of strange quark matter, given by

$$n_B=\tfrac{1}{3}(\Delta n_u+\Delta n_d+\Delta n_s)\ , \qquad (1)$$

where

$$\begin{aligned}\Delta n_i&=n_i-\bar{n}_i\\ &=\frac{g_i}{(2\pi)^3}\int d^3p\left[\frac{1}{\exp[\beta(\epsilon_i-\mu_i)]+1}-\frac{1}{\exp[\beta(\epsilon_i+\mu_i)]+1}\right], \end{aligned}\qquad (2)$$

$(\bar{n}_i)$ n_i is the number density of the (anti)flavor i ($i=u$, d, or s), $g_i=6$, the degeneracy factor, μ_i is the chemical potential (for antiparticle $\bar{\mu}_i=-\mu_i$), and $\epsilon_i=\sqrt{p^2+m_i^2}$ is the single-particle energy of the same flavor.

The masses of u, d (or $\bar{u}$ and $\bar{d}$), and s (or $\bar{s}$) quarks are parametrized in the following manner [13]:

$$m_{u,d,\bar{u},\bar{d}}=B/3n_B\ , \qquad (3a)$$

$$m_{s,\bar{s}}=m_s^0+B/3n_B\ , \qquad (3b)$$

where B is the vacuum ($n_B\rightarrow 0$) energy density, m_s^0 is the current mass of s quark (and also $\bar{s}$), whereas for the light quarks and antiquarks the current mass is assumed to be zero, which is in agreement with the asymptotic freedom and the restoration of chiral symmetry at very high density.

Inside SQM, s (and also $\bar{s}$) quarks are produced through the weak processes

$$\begin{aligned}&u+d\leftrightarrow u+s,\quad s\rightarrow u+e^-+\bar{\nu}_e\ ,\\ &d\rightarrow u+e^-+\bar{\nu}_e,\quad u+e^-\rightarrow d+\nu_e\ ,\end{aligned}\qquad (4)$$

and similarly for the antiquarks. The dynamical chemical equilibrium among the participants yields

$$\mu_d=\mu_s=\mu(\text{for example})\text{ and }\mu_u=\mu-\mu_e\ , \qquad (5)$$

where u, d, s, and e stand for the u quark, d quark, s quark, and electron, respectively (for antiparticles $\bar{\mu}_i=-\mu_i$, where $i=u$, d, s, and e). The created neutrinos are assumed to leave the system freely (no Pauli blocking), so that $\mu_\nu=-\bar{\mu}_\nu=0$.

Overall electrical charge neutrality of the SQM system gives

$$2\Delta n_u=\Delta n_d+\Delta n_s+3\Delta n_e\ , \qquad (6)$$

where Δn_e is given by Eq. (2). Then the free energy density ϵ of the system can be written as

$$\epsilon=\epsilon_q+\epsilon_e\ , \qquad (7a)$$

where

$$\epsilon_q=-P_q+\sum_i\mu_i n_i-T\sum_i\frac{\partial\Omega_i}{\partial T} \qquad (7b)$$

is the free energy density of the quark sector, $i=u$, d, and s (also their antiparticles), and

$$\epsilon_e=-P_e+\sum_i\mu_i n_i-T\sum_i\frac{\partial\Omega_i}{\partial T} \qquad (7c)$$

is the electronic free energy density; here $i=e$ and $\bar{e}$, P_q, and P_e are the kinetic pressures of quark part and electronic part, respectively, and are given by

$$P_i=\sum_i\frac{1}{3}\frac{g_i}{(2\pi)^3}\int d^3p\frac{p^2}{\sqrt{p^2+m_i^2}}\times\left[\frac{1}{\exp[\beta(\epsilon_i-\mu_i)]+1}+\frac{1}{\exp[\beta(\epsilon_i+\mu_i)]+1}\right], \qquad (7d)$$

where $i=u$, d, s, or e. The thermodynamic potential density Ω_i is given by

$$\Omega_i=-\frac{g_iT}{(2\pi)^3}\int d^3p\ \ln(1+e^{-\beta(\epsilon_i-\mu_i)})=-P_i\ , \qquad (7e)$$

where $i=u$, d, s (or $\bar{u}$, $\bar{d}$, $\bar{s}$), and e (or $\bar{e}$).

For the stable configuration of SQM, the free energy per baryon (ϵ_B) should be less than M_N, the nucleon mass. One can make the energy per baryon for SQM even less than the corresponding quantity for iron by properly adjusting the numerical value of the parameter B. The value of this parameter has to be reduced slightly. Here, we have not considered the reference energy as the energy needed to emit a neutron. We have studied numerically the emission of neutrons from the surface of SQM, using a Monte Carlo coalescence model. At the surface we have chosen three quarks randomly with momenta very close to each other, and spins, colors, and flavors are adjusted properly to construct neutrons. We have seen that the emission rate of neutrons is negligibly small.

To obtain the chemical potentials for the constituents of SQM we have solved Eqs. (1) and (6) numerically with the chemical equilibrium conditions (5), for a given baryon number density n_B and temperature T. The parameter B is adjusted in such a way that bulk QM system becomes just unstable (i.e., ϵ_B for QM becomes just greater than M_N).

Knowing the values of chemical potentials for all the constituents we have evaluated the electron and strangeness densities, given by $n_i = -\partial\Omega_i/\partial\mu_i$, where $i = e$ and s, for various values of baryon density and temperature. We have also calculated energy/baryon for SQM for the same values of density and temperature.

In Fig. 1, we have plotted the variation of energy/baryon for SQM with baryon number density, for two different temperatures [30 MeV (curve a) and 10 MeV (curve b)]. The values of the parameter B are respectively $(148\ \text{MeV})^4$ and $(180\ \text{MeV})^4$ for these two temperatures. It is to be noted that beyond $T=35$ MeV there is no stable configuration of SQM for an acceptable value of B $[\geq(140\ \text{MeV})^4?]$. For both these cases, the minima of ϵ_B occur near $2.2n_0$.

In Fig. 2 the temperature dependence of ϵ_B's are plotted for two different baryon number densities ($2n_0$ and $4n_0$). Both these curves show that ϵ_B is a monotonically increasing function of T and SQM becomes unstable for $T>35$ MeV.

In Fig. 3 we have presented the variation of the strangeness fraction for a stable SQM system with the baryon number density for the temperatures and parameter B as discussed above. Both the curves show that with the increase of baryon number density the strangeness fraction increases and saturates to a constant value $\frac{1}{3}$, which is the flavor symmetric situation. Since, at high baryon density regime, the production cross section of s quarks and the absorption cross section of electrons are very high (mean free path of electron becomes extremely small), the system will try to become globally charged neutral by properly adjusting the numbers of u, d, and s quarks and electrons (and also their antiparticles). As a consequence, the system will become almost flavor symmetric ($\Delta n_u \approx \Delta n_u \approx \Delta n_u$) at very high density.

Figure 4 represents the variation of the strangeness fraction with temperature for $n_B=2n_0$ and $4n_0$. These curves show that the strangeness fraction is an increasing function of temperature. But beyond some baryon number density ($\geq 5n_0$), it saturates to a maximum value $\frac{1}{3}$ and becomes independent of temperature.

The variation of electron density with baryon density is plotted in Fig. 5. The behavior of electron density is just opposite to that of the strangeness fraction; it decreases with the increase of baryon number density. Unlike the strange quark density, the electron density is extremely small for all densities and temperatures.

Figure 6 shows the variation of electron number density with temperature and is an increasing function of tem-

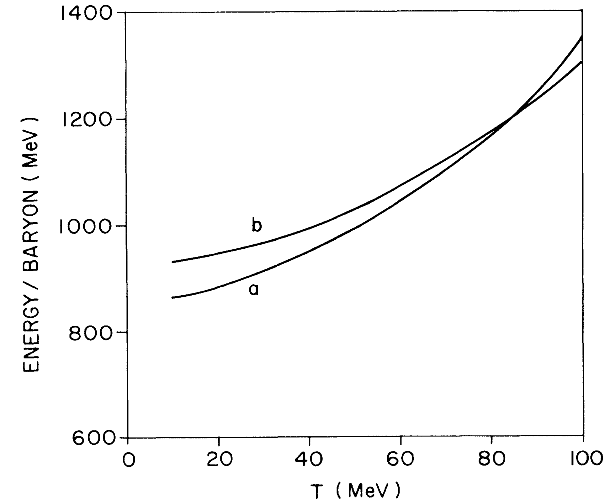

FIG. 2. Temperature dependence of energy per baryon; curves (a) $n_B=2n_0$ and (b) $n_B=4n_0$.

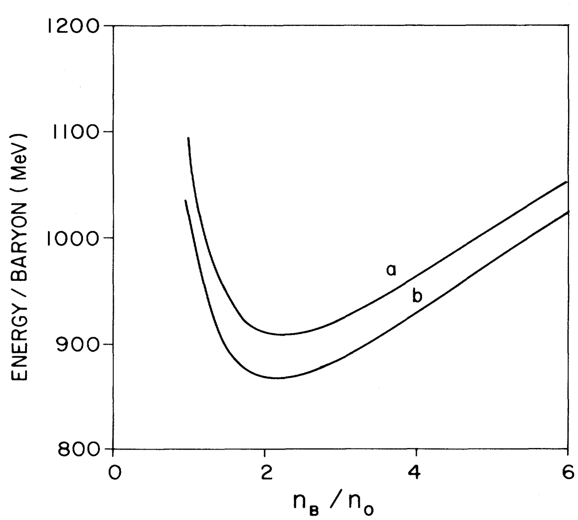

FIG. 1. Variation of energy per baryon with the baryon number density (expressed in terms of n_0) for $T=10$ MeV (curve b) and 30 MeV (curve a).

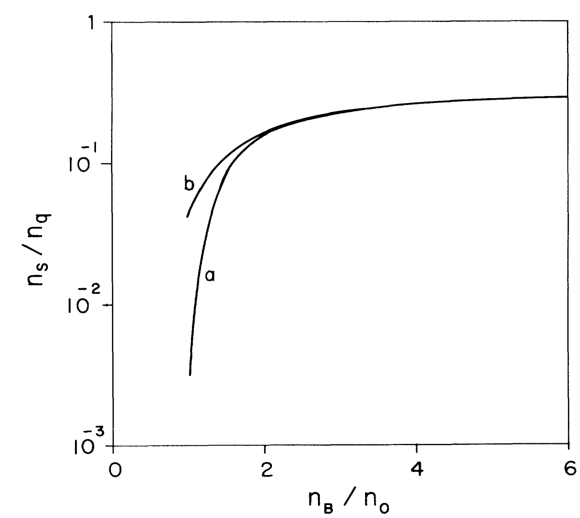

FIG. 3. Variation of strangeness fraction (n_s/n_q) with the baryon number density; curves (a) $T=10$ MeV and (b) $T=30$ MeV.

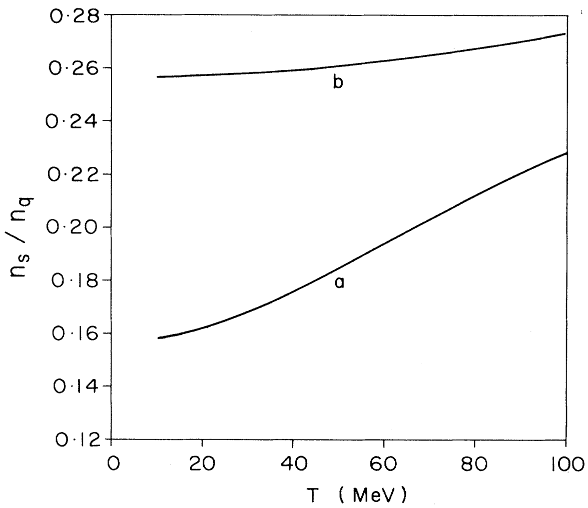

FIG. 4. Temperature dependence of strangeness fraction; curves (a) $n_B = 2n_0$ and (b) $n_B = 4n_0$.

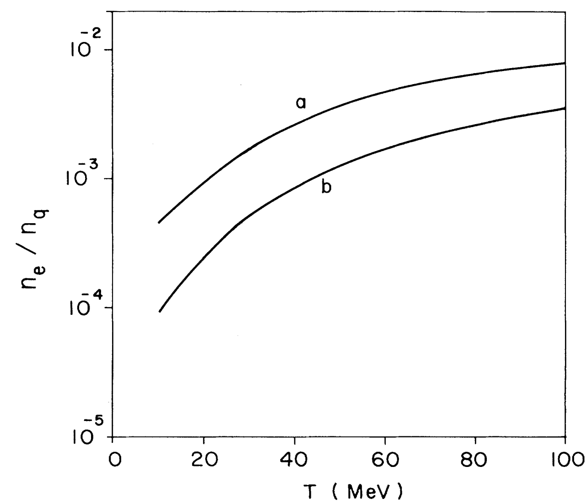

FIG. 6. Temperature dependence of electron number density; curves (a) $n_B = 2n_0$ and (b) $n_B = 4n_0$.

perature. An increase of temperature increases the mean free path of electrons inside SQM.

III. STRANGELETS AT $T \neq 0$

For the sake of simplicity we assume the shape of strangelets to be spherical (which is again a necessary condition to minimize the surface energy). Then the radius R can be parametrized in the form

$$R = r_0 A^{1/3} , \tag{8}$$

where A is the total number of baryons present in the strange droplet and r_0 is an unknown parameter and has to be determined. Assuming a uniform baryon density inside the strangelets, we have

$$n_B = \frac{A}{V} = \frac{3}{4\pi r_0^3} ; \tag{9}$$

hence

$$r_0 = \left[\frac{3}{4\pi n_B} \right]^{1/3} , \tag{10}$$

and then following (3a) and (3b), we have

$$m_{u,d,\bar{u},\bar{d}} = C_1 r_0^3 , \tag{11a}$$

$$m_{s,\bar{s}} = m_s^0 + C_1 r_0^3 , \tag{11b}$$

where $C_1 = 4\pi B/9$.

Unlike bulk SQM, in this case the total energy is given by the sum of volume and surface energies

$$E_{\text{tot}} = \epsilon_{\text{vol}} V + \epsilon_{\text{surf}} S , \tag{12}$$

where ϵ_{vol} is the volume energy density, V is the total volume, ϵ_{surf} is the surface energy density (energy/surface area), and S is the surface area surrounding the volume V.

Then following Berger [19] (see also Ref. [20]), we have

$$\epsilon_{\text{surf}} = \sigma = \frac{T}{32\pi^2} \sum_i g_i \int \frac{d^3k}{k} \left[1 - \frac{2}{\pi} \arctan \frac{k}{m_i} \right] \times \ln \left[1 + \exp \left[- \frac{\epsilon_i - \mu_i}{T} \right] \right] , \tag{13}$$

where σ is the surface tension and $i = u$, d, and s (and also their antiparticles). In the case of electrons (or positrons), $m_e = 0$ and, therefore, $\sigma = 0$ (no surface contribution from massless components). This expression for surface tension follows immediately from the work of Balian and Bloch [21] using the bag model boundary conditions (see also Ref. [22] for an excellent discussion). In our

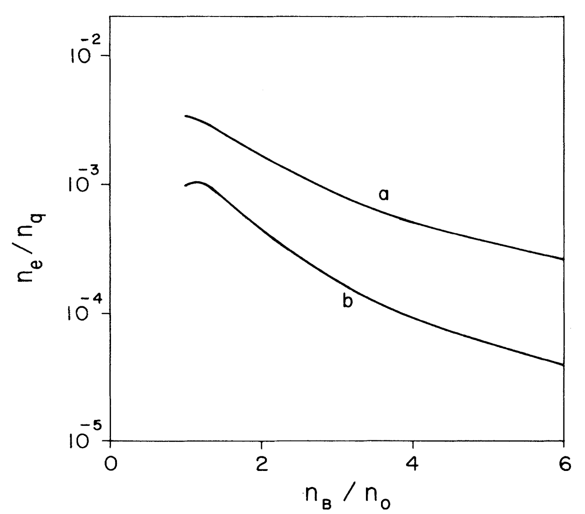

FIG. 5. Variation of electron number fraction (n_e/n_q) with the baryon number density; curves (a) $T = 30$ MeV and (b) $T = 10$ MeV.

confinement model, which is entirely different from the MIT bag model, the mass of the quark becomes infinity just outside the hadrons. The scenario is physically analogous to the confinement of quarks in an infinite potential well. In this case, the quark wave function vanishes at the boundary, which is the well-known Dirichlet boundary conditions. This condition uniquely specifies a solution at the boundary, which is zero in this particular model of confinement. Here we do not need derivatives of the wave function at the boundary. However, in this model, since the total energy density is finite in vacuum (which is B, for $n_q \rightarrow 0$), the second derivative of the wave function must become infinity at the surface and, as a consequence, its first derivative is discontinuous at the boundary. Then following Balian and Bloch [21], the modified form of Eq. (13) which is appropriate for our model, is given by

$$\epsilon_{\rm surf}=\sigma$$

$$=\frac{T}{32\pi^2}\sum_i g_i \int \frac{d^3k}{k} \ln\left[1+\exp\left[-\frac{\epsilon_i-\mu_i}{T}\right]\right], \quad (13a)$$

where $i=u$, d, and s, and their antiparticles. Here the electrons can go out of the system and form a thin layer at the strangelet surface. This electronic layer will be bounded to the positively charged strange nucleus by a strong attractive Coulomb potential ($V \approx 20$ MeV). In this calculation we have neglected the surface energy contribution from the electronic part.

Solving numerically for chemical potentials of the constituents as before, we have evaluated the surface tension of the strangelet for various values of baryon density, temperature, and parameter B. Also in this case the parameter B is adjusted to make the nonstrangelet just unstable. The values of $B^{1/4}$ are ≈ 155 MeV, ≈ 162 MeV, and ≈ 184 MeV for the temperatures 10, 35, and 50 MeV, respectively. In this case $T=55$ MeV is the upper limit for the strangelet temperature, beyond which there cannot be any stable configuration.

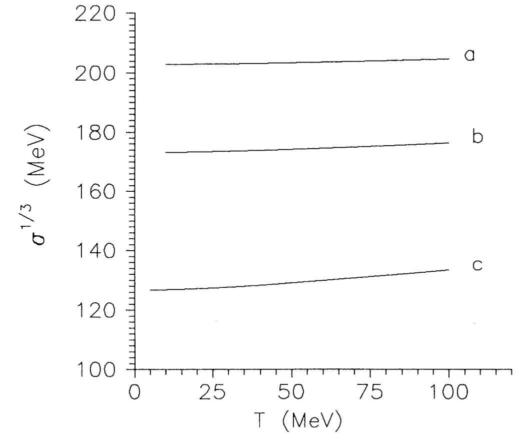

FIG. 8. Temperature dependence of strangelet surface tension; curves (a) $n_B=6n_0$, (b) $n_B=4n_0$, and (c) $n_B=2n_0$.

In Fig. 7, we have plotted density dependence of $\sigma^{1/3}$ for three different temperatures. These curves show that the surface tension of the strangelets is an increasing function of baryon number density.

Figure 8 shows the variation of surface tension with the temperature of the system and is a decreasing function of temperature, but the variation is very slow. Figure 9 shows the variation of $\sigma^{1/3}$ with the parameter B for three different temperatures and fixed baryon number density $n_B=2n_0$. The surface tension first increases with B, reaches a maximum value, and then decreases continuously. The maximum value is a function of temperature. Within the values of B, for which strangelets are predicted to be stable, the surface tension is an increasing function of B.

Now for the stability of strangelets, the total energy per baryon should be less than M_N, i.e.,

$$\frac{\epsilon_{\rm tot}}{A}=\frac{1}{A}(\epsilon_{\rm vol}V+\epsilon_{\rm surf}S)<M_N \ . \quad (14)$$

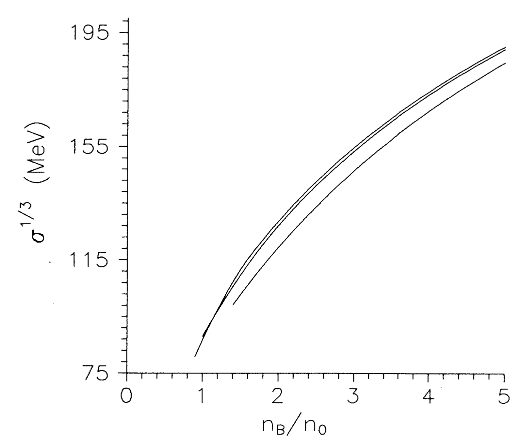

FIG. 7. Baryon density dependence of strangelet surface tension; curves (a) $T=50$ MeV, (b) $T=30$ MeV, and (c) $T=10$ MeV.

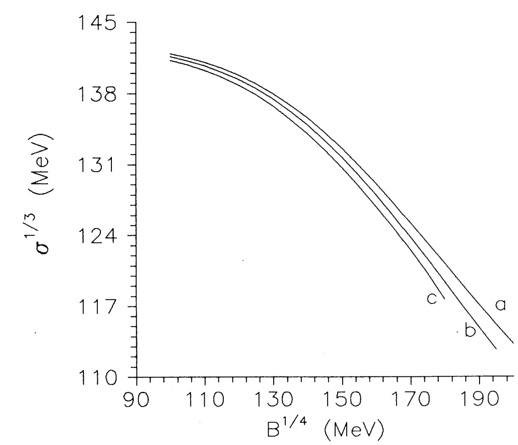

FIG. 9. Variation of surface tension with the parameter B; curves (a) $T=50$ MeV, (b) $T=30$ MeV, and (c) $T=10$ MeV; $n_B=2n_0$.

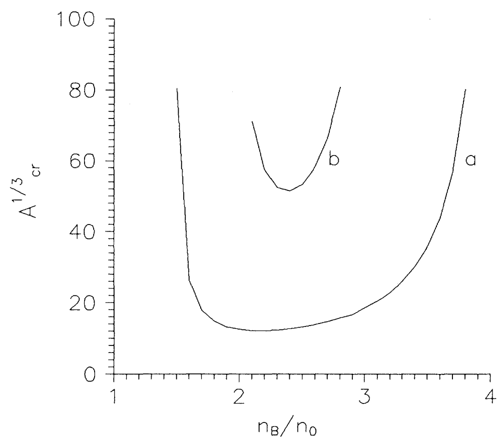

FIG. 10. Variation of critical mass number with baryon number density; curves (a) $T = 10$ MeV and (b) $T = 30$ MeV.

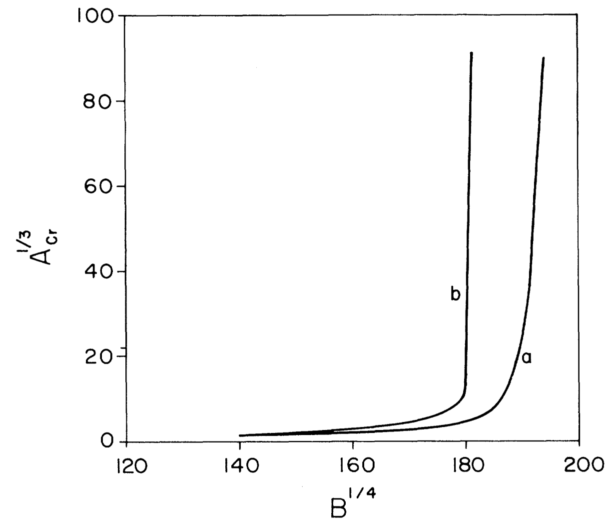

FIG. 12. Variation of critical mass number with the parameter B; curves (a) $T = 10$ MeV and (b) $T = 30$ MeV; $n_B = 2n_0$.

For spherical strangelet $S = 4\pi r_0^2 A^{2/3}$, which leads to the minimum value of the baryon number content of a stable strangelet A_{cr} (critical mass number) given by

$$A_{cr} = \left[\frac{4\pi r_0^2 \sigma}{M_N - \epsilon_B} \right]^3 , \tag{15}$$

where $\epsilon_B = \epsilon_{vol}/n_B$, the bulk energy per baryon.

Equation (15) shows that the critical mass number is a function of temperature, baryon number density, and parameter B.

In Fig. 10, the variation of the critical mass number with the baryon number density of the system is shown for two different temperatures. For both the temperatures, there are some allowed regions in the $A_{cr} - n_B$ space for which A_{cr} is finite, indicating the possibility of stable strangelet formation in this regime. Beyond these regions, the critical mass number becomes infinitely large. It is also seen from these figures that, with the increase of temperature, the allowed region gets narrower.

Figure 11 shows the temperature dependence of the critical mass number of a stable strangelet for two different baryon number densities ($n_B = 2n_0$ and $4n_0$). For low baryon number density, $A_{cr}^{1/3}$ remains almost constant up to a certain temperature and then suddenly increases to a very large value and finally goes to infinity, indicating the barrier and upper limit of the temperature beyond which there cannot be any stable strangelet. For high baryon densities, the critical temperatures are shifted towards the lower values.

In Fig. 12, the variation of A_{cr} with parameter B is plotted. In this case also, there is an upper limit of B beyond which there cannot be any stable strangelet.

Now the ratio of protonic charge (Z) and the total baryon number (A) of the strangelet is given by

$$\frac{Z}{A} = \frac{2n_u - n_d - n_s}{n_u + n_d + n_s} . \tag{16}$$

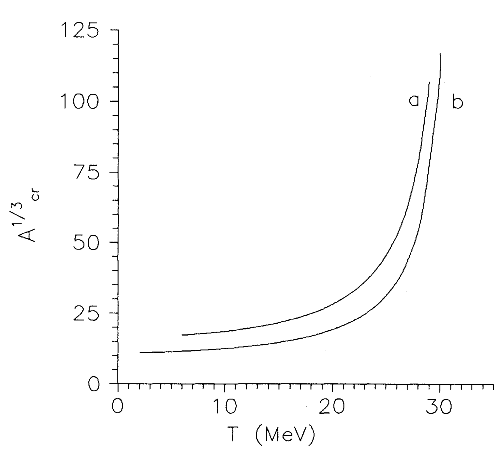

FIG. 11. Temperature dependence of critical mass number; curves (a) $n_B = 2n_0$ and (b) $n_B = 4n_0$.

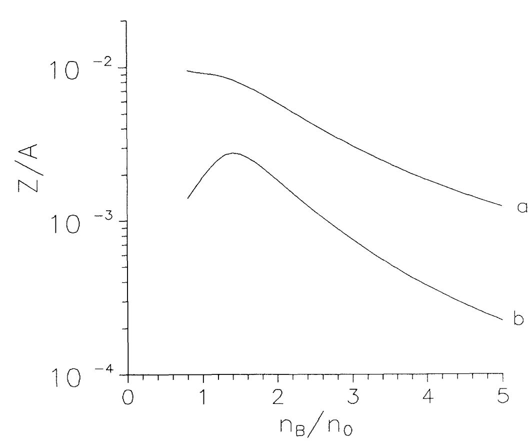

FIG. 13. Variation of Z/A with the baryon number density; curves (a) $T = 30$ MeV and (b) $T = 10$ MeV.

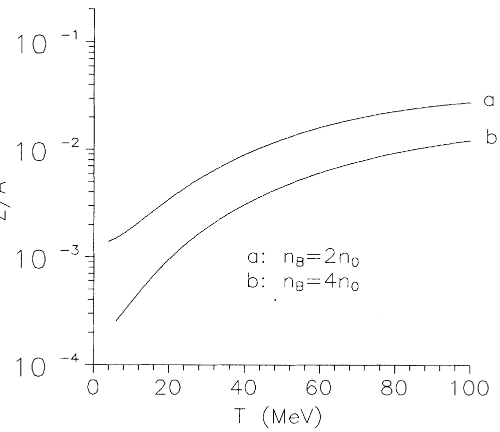

FIG. 14. Temperature dependence of Z/A.

This ratio has been evaluated for various values of n_B, T, and B.

In Fig. 13, the variation of Z/A with the baryon number density is plotted for two different temperatures. The ratio first increases, reaches a maximum, and then decreases continuously. Curve a is for $T=30$ MeV and curve b is for 10 MeV, respectively. Unlike a stable nucleus for which $Z/A\approx1/2$ here it is extremely small ($<10^{-2}$). This is one of the peculiarities of stable strangelets.

Figure 14 shows the variation of Z/A with temperature. The ratio remains much less than 10^{-2} up to the stability temperature 55 MeV, but the ratio is an increasing function of temperature.

On the other hand, Z/A is an increasing function of B, but the variation is very slow. In Fig. 15, we have plotted this variation for two different values of baryon number density ($=2n_0$ and $4n_0$) and constant temperature $T=10$ MeV.

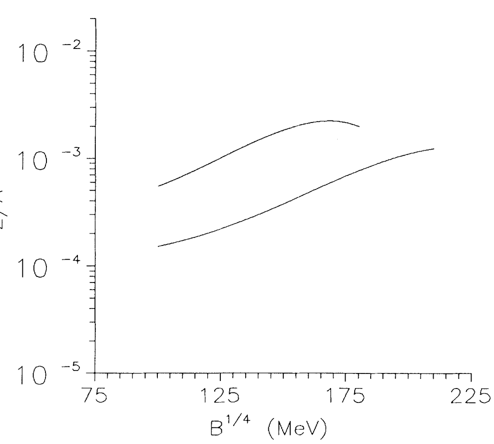

FIG. 15. Variation of Z/A with $B^{1/4}$ for two different values of n_B ($=2n_0$, upper curve and $4n_0$, lower curve) and $T=10$ MeV.

IV. STRANGELETS IN ULTRARELATIVISTIC COLLISIONS

If a baryon-rich QGP is formed in ultrarelativistic heavy-ion collisions, an equal number of s and $\bar{s}$ quarks will be created by the annihilation of light quarks and antiquarks and also by the pair production of gluon pairs. Since the chemical potentials of both u and d quarks (and their antiparticles) are nonzero, further creation of $u\bar{u}$ or $d\bar{d}$ pairs will be suppressed. Since initially (inside the colliding nuclei), there was no strange quark, the initial chemical potentials $\mu_s(t=0)=-\bar{\mu}_s(t=0)=0$. Now antistrange quarks produced in the system can combine with u or d quarks and give rise to $K^+(\bar{s}u)$ or $K^0(\bar{s}d)$ mesons and will be evaporated out through the QGP surface. Since the abundances of $\bar{u}$ and $\bar{d}$ quarks are extremely low, the production of $K^-(s\bar{u})$ and $\bar{K}^0(s\bar{d})$ mesons will consequently be very rare. As a result, the strangeness fraction in the QGP bubble will continuously increase with time and, as a consequence, the chemical potential of strange quarks will also increase (n_s or μ_s will be functions of time). The temperature of the system decreases due to the combined effects of expansion, pion evaporation (these two processes will not change the strangeness content of the system) and kaon evaporation. Since μ_s is a function of time, ultimately the strangeness fraction within QGP will saturate to a maximum value. At that moment the rate of creation of $s\bar{s}$ pairs and their annihilation rate plus evaporation rates of K^+ and K^0 become just equal in magnitude. Recently Greiner and coworkers [12] have done an exhaustive work on the distillation and survival of strangelets in ultrarelativistic heavy-ion collisions using the rate calculation method. In this section we are not going to repeat those calculations again. Our intention is to see whether the confinement model proposed by us [13] is able to predict the existence of stable strangelets, whose origin is the ultrarelativistic heavy-ion collisions.

In our treatment, strangeness (n_s), baryon density (n_B), and temperature (T) are variables and we shall try to find out the ranges of n_s, n_B, and temperature T within which such a system is stable. In our calculation, we are considering a situation where the strangeness fraction has already been saturated to a maximum value, and a strange droplet has been produced, i.e., the last stage of the evolution of the strange droplet. Since the saturated value of the strangeness density is a function of both baryon number density and temperature, we are treating all of them as parameters and we have investigated the stability of the strangelet in this parameter space.

Unlike cosmic or astrophysical SQM, where strange quarks are produced through the weak processes (4), here the underlying processes are strong in nature. Again, the system is not necessarily electrically neutral. The only condition one can impose is the baryon number conservation, provided no baryons are evaporated out from the QGP bubble (which is possibly true at the last stage of the evolution of the QGP bubble when the temperature becomes much less than M_N, the nucleon mass). Then writing $\mu_u=\mu_d=\mu_q$, one can solve for μ_q and, for a given strange quark density n_s, the strange quark chemical po-

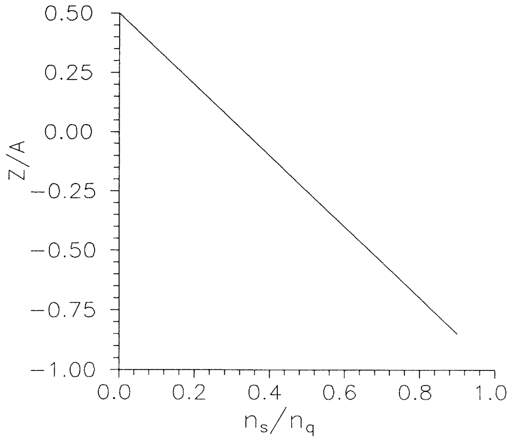

FIG. 16. Variation Z/A with the strangeness fraction n_s/n_q of a strangelet formed in ultrarelativistic heavy-ion collisions.

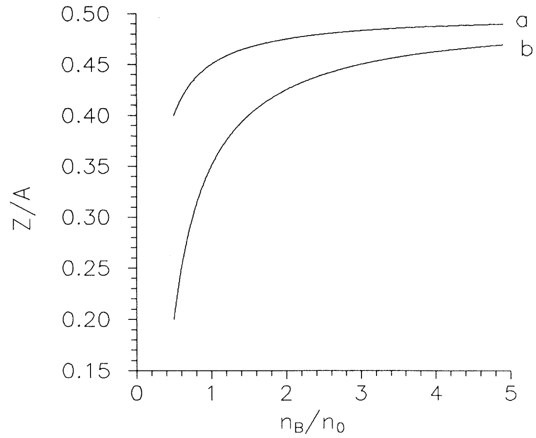

FIG. 17. Variation of Z/A with n_B/n_0, for $n_s=0.1n_q$ (upper curve) and $n_s=0.3n_q$ (lower curve).

tential μ_s can be evaluated, where n_B is given by Eq. (1) and

$$n_s=\frac{g_s}{(2\pi)^3}\int d^3p\left[\frac{1}{\exp[\beta(\epsilon-\mu_s)]+1}-\frac{1}{\exp[\beta(\epsilon+\mu_s)]+1}\right]. \quad (17)$$

We have seen that the density range within which the system is stable is $1.6-3.5n_0$, and the upper limit of the temperature is ≈ 40 MeV. The lower and upper limits of the strangeness fraction (n_s/n_q) are 0.13 and 0.7, respectively (provided by the baryon density lying within the range stated above). In Fig. 16, we have plotted

$$Z/A=n_u/n_B-1=0.5(1-n_s/3n_q)$$

for a stable strangelet for different values of the strangeness fraction (n_s/n_q). Obviously, the ratio Z/A remains positive and less than 0.5 up to the strangeness fraction value $\frac{1}{3}$ and then becomes negative (when the system becomes *more than flavor symmetric,* which is an unphysical situation) and then decreases monotonically. We have also noted that for a given n_s, Z/A first increases and finally saturates to a constant value 0.5 for large n_q. In Fig. 17, we have shown the variation of Z/A with n_B/n_0 for two different values of n_s.

V. DISCUSSIONS AND CONCLUSIONS

The dynamical density-dependent quark mass approach to confinement is also successful in predicting the possibility of stable SQM at finite temperature. The density regime, within which bulk SQM is stable is lying between $1.8n_0$ and $5n_0$. The minimum value of ϵ_B occurs at $\approx 2.2n_0$. This density point is again independent of temperature, whereas the density range for which the strangelets are stable lies between $2n_0$ and $4n_0$.

The strangeness fraction increases with the increase of both the baryon number density as well as temperature of the system. Beyond $n_B=5n_0$, the strangeness fraction saturates to a constant maximum value $\frac{1}{3}$, which is the flavor-symmetric situation. In this regime, the rate of absorption of electrons by u quarks balances its rate of creation.

The ratio Z/A for strangelets is much smaller than that of stable nuclei. For cosmo or astrophysical strangelets, Z/A can be defined only after $n_B\geq 1.1n_0$, whereas for the strangelets, whose origin is ultrarelativistic heavy-ion collisions, this ratio is independent of T, and decreases linearly both with n_s and $1/n_B$ for given n_B and n_s, respectively. The ratio is positive and less than 0.5 up to a flavor-symmetric situation and then becomes negative and decreases continuously. On the other hand, for a given n_s, the ratio Z/A saturates to 0.5 for some large value of n_B and this limiting value increases with the increase of n_s. This is in good agreement with the recent results obtained by Greiner and co-workers [12].

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[1] E. Witten, Phys. Rev. D **30**, 272 (1984).
[2] C. A. Alcock and E. Farhi, Phys. Rev. D **32**, 1273 (1985).
[3] C. A. Alcock, E. Farhi, and A. Olinto, Astrophys. J. **310**, 261 (1986).
[4] P. Haensel, J. L. Zdunik, and R. Schaeffer, Astron. Astrophys. **160**, 251 (1986).
[5] A. Chodos *et al.*, Phys. Rev. D **9**, 3472 (1974); T. DeGrand *et al.*, *ibid.* **12**, 2060 (1975).
[6] M. S. Berger and R. L. Jaffe, Phys. Rev. C **35**, 213 (1987); E. G. Blackman and R. L. Jaffe, Nucl. Phys. **B324**, 205 (1989).
[7] H. Reinhardt and B. V. Dang, Phys. Lett. B **202**, 133

(1988); T. Chmaj and W. Slominski, Phys. Rev. D **40**, 165 (1989).

[8] T. Chmaj and W. Slominski, in *Strange Quark Matter in Physics and Astrophysics,* Proceedings of the International Workshop, Aarhus, Denmark, 1991, edited by J. Madsen and P. Haensel [Nucl. Phys. B (Proc. Suppl.) **24B**, 14 (1991)].

[9] M. L. Olesen and J. Madsen, in *Strange Quark Matter in Physics and Astrophysics* [8], p. 170, and references therein.

[10] H. Heiselberg, G. Baym, and C. J. Pethick, in *Strange Quark Matter in Physics and Astrophysics* [8], p. 144.

[11] S. Chakrabarty, Nuovo Cimento **106B**, 1023 (1991).

[12] C. Greiner, P. Koch, and H. Stöcker, Phys. Rev. Lett. **58**, 1825 (1987); C. Greiner, D. H. Rischke, H. Stöcker, and P. Koch, Phys. Rev. D **38**, 2797 (1988); C. Greiner, P. Koch, and H. Stöcker, *ibid.* **44**, 3517 (1991).

[13] S. Chakrabarty, S. Raha, and B. Sinha, Phys. Lett. B **229**, 112 (1989).

[14] S. Chakrabarty, Phys. Scr. **43**, 11 (1991).

[15] S. Chakrabarty, Phys. Rev. D **43**, 627 (1991).

[16] S. Chakrabarty, in *Strange Quark Matter in Physics and Astrophysics* [8], p. 148.

[17] A. Ukawa, in *Quark Matter '88,* Proceedings of the Seventh International Conference on Ultrarelativistic Nucleus-Nucleus Collisions, Lenox, Massachusetts, 1988, edited by G. Baym, P. Braun-Munzinger, and S. Nagamiya [Nucl. Phys. **A498**, 227c (1989)].

[18] G. Baym and S. A. Chin, Phys. Lett. **62B**, 241 (1976).

[19] M. S. Berger, Phys. Rev. D **40**, 2128 (1989); **43**, 4150(E) (1991).

[20] J. Madsen, in *Strange Quark Matter in Physics and Astrophysics* [8], p. 84.

[21] R. Balian and C. Bloch, Ann. Phys. (N.Y.) **60**, 401 (1970).

[22] R. L. Jaffe, in *Strange Quark Matter in Physics and Astrophysics* [8], p. 8.