PHYSICAL REVIEW D VOLUME 43, NUMBER 4 15 FEBRUARY 1991

# Boiling of strange-quark matter

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(Received 3 October 1990)

Boiling of strange-quark matter is studied assuming baryon chemical equilibrium between a quark phase described by the bag model, and a hadron phase described by Walecka's mean-field theory. Boiling of quark nuggets at high temperatures is shown to be much less efficient than previously believed. Thus quark nuggets with large baryon numbers may survive from the early Universe after all. Similar arguments may be important for other kinds of nontopological solitons as well.

## I. INTRODUCTION

The possible existence of an absolutely stable phase of quark matter, called strange matter because it consists of roughly equal numbers of up, down, and strange quarks, has created much interest in the last couple of years.[1–6] It has been suggested that this sort of matter might form in neutron stars or during the cosmic quark-hadron phase transition at temperature $T \approx 100$ MeV.[1] Limits on the cosmic abundance of strange-matter lumps (quark nuggets) have been derived from big-bang nucleosynthesis[2] and neutron-star glitches,[3] but heavy lumps (with baryon number $A$ in excess of $10^{35}$) are not sensitive to these arguments.

Whether heavy quark nuggets are formed in the early Universe is still debated in the literature, but if they can form it is obviously of great interest to know whether they can survive in the hot environment. Alcock and Farhi[4] showed that surface evaporation of neutrons and protons is a severe threat to nugget survival. The authors argued that nuggets less massive than $10^{-5} M_\odot$ (baryon number $A \lesssim 10^{52}$) evaporated in this way, and since the causality limit (the baryon number within the horizon) is $10^{49}\,(T/100\text{ MeV})^{-3}$, they concluded that quark nuggets were unlikely to survive.

This result was modified by taking account of flavor equilibration near the nugget surface.[5] Emission of nucleons leads to an enrichment of strangeness in the surface layers, and this in turn reduces the emission rate significantly. The evaporation turns out to be governed by an equilibrium between nucleon and kaon emission, with an effective nucleon binding energy of several hundred MeV. Nuggets with $A \gtrsim 10^{46}$ can survive this evaporation, and it was argued that even smaller nuggets might survive due to reabsorption because of insufficient hadron transport away from the surface.

Recently Alcock and Olinto[6,7] have discussed *boiling* as a more efficient mechanism for nugget destruction. The idea is based on the fact that a free nucleon gas is thermodynamically favored for temperatures above a few MeV, because it has higher entropy and hence lower free energy than strange matter. Analogous to boiling of superheated water, bubbles of hadronic gas spontaneously nucleate throughout the nugget volume. Those bubbles larger than a critical radius given by the surface tension and pressure difference between the phases can grow, converting the quark phase into nucleons. As long as the total surface area of the hadronic bubbles is much larger than the surface area of the nugget, this boiling is less sensitive to flavor disequilibration than surface evaporation.

Based on this idea, the authors conclude, that quark nuggets within the causality limit are unable to survive boiling at temperatures near 100 MeV unless the surface tension $\sigma$ exceeds $(178\text{ MeV})^3$. Such a high value seems excluded; e.g., Berger[8] concludes that $\sigma$ is certainly smaller than $(105\text{ MeV})^3$.

The present paper shows that nugget boiling is significantly less efficient than calculated in Ref. 6, and that the evaporation results in Ref. 5 may give the relevant survival parameters after all. The rate of hadron bubble formation is recalculated for a degenerate, interacting nucleon gas, rather than a gas obeying the ideal-gas equation, including a self-consistent calculation of the surface tension and bag constant $B$ in the quark phase. At a given temperature boiling turns out to be extremely sensitive to the values of $B$ and $\sigma$, typically with boiling being unimportant for $B^{1/4} \lesssim 135-160$ MeV, corresponding to $\sigma^{1/3} \lesssim 40-60$ MeV, the exact values depending on temperature. (We thus find that there is a *lower* bound on $\sigma$ for boiling to be important at a given $T$, rather than an upper bound, as found in Ref. 6.)

The physics of boiling is described in Sec. II. Section III introduces the relevant physical quantities to be calculated for the quark phase, and Sec. IV gives the similar prescriptions for the hadron phase, which is assumed to be described by Walecka's mean-field theory for an interacting neutron-proton-electron ($npe$) gas. Section V contains the results concerning boiling of strange matter to hadron matter (assuming equal baryon chemical potentials in the two phases). Finally Sec. VI is devoted to a discussion of the results and conclusions.

## II. NUCLEATION OF HADRONIC BUBBLES

According to classical nucleation theory[9,6] boiling is governed by the formation rate of critical bubbles, i.e., bubbles of the minimum free energy state with radius $r_c$ maximizing the thermodynamic work $W$ expended to create the bubble. In general the thermodynamic work is given by

 

$$W=\frac{4\pi}{3}r^3(P_q-P_h)+4\pi\sigma r^2 , \tag{1}$$

where $r$ is the bubble radius, $P_q$ is the pressure of strange matter (equal to the mean pressure in the Universe), $P_h$ the pressure in the hadronic bubble, and $\sigma$ the surface tension. The thermodynamic work is maximized for bubbles with a critical radius

$$r_c=\frac{2\sigma}{P_h-P_q} , \tag{2}$$

corresponding to

$$W_c=\frac{16\pi}{3}\frac{\sigma^3}{(P_h-P_q)^2} . \tag{3}$$

Critical bubbles nucleate with a rate given by

$$p\approx\epsilon^4\exp(-W_c/T) , \tag{4}$$

where $\epsilon$ is a characteristic energy expected to be comparable to $T$.

Alcock and Olinto[6] argued that the two phases in question had equal pressure contributions from the thermal spectrum of light particles (photons, neutrinos and electrons), so that $\Delta P\equiv P_h-P_q$ was given entirely by the pressure due to neutrons and protons. They described this pressure difference in terms of the ideal-gas equation of state for the nucleons, so that the bubble nucleation depended on the baryon chemical potential in the strange matter through the neutron binding energy.

However, the density inside a quark nugget is comparable to densities in nuclei, so an ideal-gas equation of state is not appropriate for the problem at hand. Instead one has to use an equation of state that incorporates the effects of degeneracy and interactions in the hadron phase. As described in Sec. IV we have chosen Walecka's mean-field theory for nuclear matter, extended to an *npe* gas. The strategy is then to calculate the pressure difference between the phases from the equation

$$\mu_q(T,P_q)=\mu_h(T,P_h) , \tag{5}$$

where $\mu_q$ and $\mu_h$ are the baryon chemical potentials in the quark and hadron phase, respectively. The surface tension $\sigma$ is calculated self-consistently, i.e., at the same values of parameters.

Boiling is efficient provided that the total surface area of bubbles $\alpha_b$ exceeds the nugget surface area $\alpha_n$. If $V$ denotes the nugget volume ($V\approx 6.31\times10^{-7}A$ MeV$^{-3}$) and $n_b$ the number density of bubbles [$n_b\approx 0.1M_P T^2\exp(-W_c/T)$, where $M_P$ is the Planck mass, because the expansion time is $0.1M_P/T^2$], then $\alpha_b\approx Vn_b 4\pi r_c^2$. The nugget surface area is $\alpha_n\approx 3.58\times10^{-4}A^{2/3}$ MeV$^{-2}$. Comparing the two areas one finds that boiling is efficient for nuggets with baryon number exceeding

$$A_{\rm boil}\approx 7.90\times10^{-61}\frac{\Delta P^6}{T^6\sigma^6}\exp\left[16\pi\frac{\sigma^3}{T\Delta P^2}\right] . \tag{6}$$

[The expressions used for the nugget volume and surface area in deriving Eq. (6) assumed a constant baryon density in the nuggets. Actually the baryon densities vary with temperature and bag constant, so that $A_{\rm boil}$ given above can be wrong by a few orders of magnitude at high $T$. The results given below are based on a more exact numerical calculation of $A_{\rm boil}$, but the difference in terms of limits on $B$ and $\sigma$ are negligible because of the strong dependence of $A_{\rm boil}$ on the relevant physical quantities].

## III. THE QUARK PHASE

We shall assume that the weak interaction maintains equilibrium between the different quark flavors through the processes

$$\begin{aligned} &d\rightleftarrows u+e^-+\bar{v}_e , \\ &s\rightleftarrows u+e^-+\bar{v}_e , \\ &d+u\rightleftarrows u+s . \end{aligned} \tag{7}$$

This gives the following relations between the chemical potentials:

$$\mu_d=\mu_s\equiv\mu, \quad \mu_u=\mu-\mu_e . \tag{8}$$

Local charge neutrality gives a relation between the number densities $n_i$:

$$\tfrac{2}{3}(n_u-n_{\bar{u}})-\tfrac{1}{3}(n_d-n_{\bar{d}})-\tfrac{1}{3}(n_s-n_{\bar{s}})-(n_{e^-}-n_{e^+})=0 \tag{9}$$

which leaves only one independent chemical potential $\mu$. The chemical potential per baryon is then given as $\mu_q=3\mu-\mu_e$. The number densities are given as the integral over a Fermi-Dirac distribution

$$\begin{aligned} n_i&=\frac{g}{(2\pi)^3}\int d^3\mathbf{k}\left[1+\exp\left[\frac{\epsilon_i-\mu_i}{T}\right]\right]^{-1} , \\ n_{\bar{i}}&=\frac{g}{(2\pi)^3}\int d^3\mathbf{k}\left[1+\exp\left[\frac{\epsilon_i+\mu_i}{T}\right]\right]^{-1} , \\ &\qquad i=u,d,s,e, \end{aligned} \tag{10}$$

with $\epsilon_i=(k^2+m_i^2)^{1/2}$, where $k$ and $m_i$ are momentum and mass, respectively (in the calculations $m_u=m_d=0$). The statistical weight $g=2$ for electrons and 6 for quarks.

The independent chemical potential $\mu$ can be found from the constraint that the fermion pressures in the nugget must be balanced by the bag constant $B$ and the external pressure $P_{\rm ext}$, i.e.,

$$\Omega_u+\Omega_{\bar{u}}+\Omega_d+\Omega_{\bar{d}}+\Omega_s+\Omega_{\bar{s}}+\Omega_{e^-}+\Omega_{e^+}=-(B+P_{\rm ext}) , \tag{11}$$

where $\Omega_i$ is the thermodynamic potential:

$$\Omega_i=-\frac{gT}{(2\pi)^3}\int d^3\mathbf{k}\ln\left[1+\exp\left[-\frac{\epsilon_i-\mu_i}{T}\right]\right] . \tag{12}$$

The bag model was originally proposed in order to describe the quark-quark interaction at zero temperature. At any finite temperature, things get complicated by the presence of thermal radiation and particle-antiparticle pairs and it is necessary to find a way of dealing with the extra pressures introduced here in a self-consistent manner. Outside the nuggets, the pressure is produced by thermal photons, neutrinos, electrons, and positrons plus a small hadronic contribution which will be neglected here. Inside there is a nonthermal pressure from $u$, $d$, and $s$ quarks and electrons plus a thermal pressure from photons, neutrinos and $u$, $d$, $s$, and $e$ pairs. The thermal photons and neutrinos cancel out and the relevant external pressure is thus equal to the pressure from thermally produced electrons and positrons. Note then that the bag constant in Eq. (11) exactly balances the fermion pressures in the quark matter *minus* the thermal electron and positron pressures. (For the parameters studied below, the thermal $e^+e^-$ contribution to $P_{\rm ext}$ is less than 1% of $B^{1/4}$.)

Finally, the contribution of each fermion species to the surface tension is given by[8]

$$\sigma_i=\frac{gT}{64\pi^2}\int\frac{d^3\mathbf{k}}{k}\left[1-\frac{2}{\pi}\arctan\frac{k}{m_i}\right]\ln\left[1+\exp\left[-\frac{\epsilon_i-\mu_i}{T}\right]\right]. \tag{13}$$

Because $\arctan(\infty)=\pi/2$ only the massive $s$ quarks contribute to the surface tension. The surface tensions found in our calculations agree with those published by Berger.[8]

## IV. THE HADRON PHASE

The hadron phase is assumed to consist of neutrons, protons, and electrons in weak equilibrium. Thus the chemical potentials follow the relation

$$\mu_n=\mu_p+\mu_e\ . \tag{14}$$

The requirement of local charge neutrality then gives $\mu_e$ through

$$n_p=n_e\ , \tag{15}$$

where the densities are found from the distribution functions [Eqs. (18)–(21)]. The baryon chemical potential in the hadron phase is equal to $\mu_n$.

In Walecka's mean-field theory,[10] extended to include protons and electrons,[11] the strong interactions are described by introducing a scalar and a vector field with coupling constant and mass $g_{SF}, m_{SF}$, and $g_{VF}, m_{VF}$, respectively. The baryon-number density $n_B$ and the pressure, $P$, are given by

$$n_B=\sum_{i=n,p}\frac{g}{(2\pi)^3}\int d^3\mathbf{k}[f_i(k,T)-f_{\bar{i}}(k,T)] \tag{16}$$

and

$$P=\frac{g_{VF}^2}{2m_{VF}^2}n_B^2-\frac{m_{SF}^2}{2g_{SF}^2}(m_n-M^*)^2+\frac{1}{3}\sum_{i=n,p,e}\frac{g}{(2\pi)^3}\int d^3\mathbf{k}\frac{k^2}{(k^2+M^{*2})^{1/2}}[f_i(k,T)+f_{\bar{i}}(k,T)]\ , \tag{17}$$

where $g=2$ is the statistical weight, $f_i(k,T)$ and $f_{\bar{i}}(k,T)$ are the Fermi-Dirac distribution functions,

$$f_i(k,T)=\left[1+\exp\left[\frac{E^*-\nu_i}{T}\right]\right]^{-1}, \tag{18}$$

$$f_{\bar{i}}(k,T)=\left[1+\exp\left[\frac{E^*+\nu_i}{T}\right]\right]^{-1}, \tag{19}$$

$E^*=(k^2+M^{*2})^{1/2}$, and

$$\nu_i=\mu_i-\frac{g_{VF}^2}{m_{VF}^2}n_B,\quad i=n,p\ , \tag{20}$$

$$\nu_i=\mu_i,\quad i=e\ . \tag{21}$$

For electrons, $M^*$ is equal to $m_e$, but for the nucleons the coupling to the scalar field reduces the mass, and $M^*$ is given by

$$M^*=m_n-\frac{g_{SF}^2}{m_{SF}^2}\frac{g}{(2\pi)^3}\int d^3\mathbf{k}\frac{M^*}{(k^2+M^{*2})^{1/2}}\times[f_n(k,T)+f_{\bar{n}}(k,T)+f_p(k,T)+f_{\bar{p}}(k,T)]\ . \tag{22}$$

We have neglected the small mass difference between neutrons and protons and thus $M^*$ is the same for both species.

Equations (16) and (18)–(22) have to be solved self-consistently to find $M^*$, $\nu_n$, and $n_B$ [from (14) and (20) one has $\nu_p=\nu_n-\mu_e$] and the pressure of the nucleon gas for a given $\mu_n$ is then given by Eq. (17). Using $3\mu-\mu_e=\mu_q=\mu_h=\mu_n$ one then has the wanted $\Delta P$ as a function of $B$, $m_s$, and $T$, apart from corrections for thermal electrons and positrons, the pressure of which should be subtracted from $P$ for reasons similar to those described in Sec. III for the quark phase. Thus

$$\Delta P=P-\frac{1}{3}\frac{g}{(2\pi)^3}\int d^3\mathbf{k}\frac{k^2}{(k^2+m_e^2)^{1/2}}\times[f_{e^-}(k,T)+f_{e^+}(k,T)]_{\mu_e=0}\ . \tag{23}$$

The coupling strengths $g_{SF}^2/m_{SF}^2$ and $g_{VF}^2/m_{VF}^2$ for the scalar and vector field, respectively, can be found from fitting to ordinary nuclei:[10]

$$\frac{g_{SF}^2}{m_{SF}^2}=3.026\times10^{-4}\ \mathrm{MeV}^{-2}\ , \tag{24}$$

$$\frac{g_{VF}^2}{m_{VF}^2} = 2.219 \times 10^{-4}\ \text{MeV}^{-2} . \quad (25)$$

## V. BOILING

With the ingredients described in Secs. III and IV, the minimum baryon number $A_{\text{boil}}$ above which boiling is efficient can be calculated from Eq. (6) [or rather numerically as described after Eq. (6)]. It turns out that $A_{\text{boil}}$ is an extremely rapidly decreasing function of the bag constant, so that typically $A_{\text{boil}}$ drops from essentially infinity to essentially zero over an interval of a few MeV for $B^{1/4}$. This rapid dependence on parameters, illustrated in Fig. 1, is due to the high powers and exponential term in Eq. (6), which ensures that the number density of hadron bubbles becomes comparable to the baryon density almost as soon as boiling sets in. It means that the results of the investigation can be presented in terms of a critical bag constant $B_{\text{boil}}(T)$ below which boiling is unimportant.

Figure 2 shows $B_{\text{boil}}$ (defined as the bag constant corresponding to $A_{\text{boil}} = 10^{50}$) as a function of temperature for two values of the strange-quark mass. Also shown is the coexistence temperature $T_c$ at which the quark nuggets are assumed to be formed ($T_c \approx 0.72 B^{1/4}$)[12], i.e., the highest temperature cosmologically produced quark nuggets are likely to experience.

It is seen that boiling is unimportant for $B \leq (158\ \text{MeV})^4$ for $T \lesssim 25$ MeV, dropping to $B \leq (135\ \text{MeV})^4$ at $T \approx 80$ MeV, and then increasing again for higher temperatures, tracing the coexistence curve very closely. It is not immediately clear that this should be the case, since $T_c$ is found[12] assuming pressure equilibrium between the quark phase (thermal quarks and gluons minus the bag constant) and the hadron phase (mainly thermal pions), while $B_{\text{boil}}(T)$ is found from a pressure *difference* between the two phases for equal chemical potentials. (It should be noted that the pressure difference, at least for high temperatures, is relatively small and thus unimportant for the determination of $T_c$.)

However, at high tempertures the thermal pressure in the quark phase dominates, and for a given bag constant this results in a maximum temperature above which the quarks are unbound: $T_{\max} \approx 0.75 B^{1/4}$. This corresponds almost exactly to $T_c$ and thus one needs not worry about what happens at still higher temperatures. When $T = T_{\max}$, the nonthermal pressure must be zero and thus $\mu_q = 0$ (at fixed $B$), but as $T$ decreases below $T_{\max}$, one gets a nonthermal component, and $\mu_q$ achieves a small but nonzero value. Because of chemical equilibrium, this

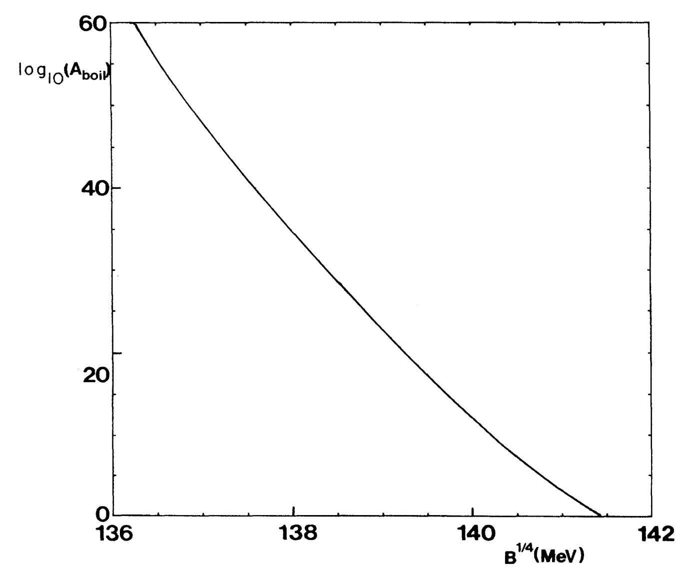


FIG. 1. The baryon number $A_{\text{boil}}$ above which boiling takes place is shown as a function of $B^{1/4}$ for $T$=70 MeV and $m_s = 150$ MeV. The range of possible $A$ values for cosmological quark nuggets ($10^2 \lesssim A \lesssim 10^{50}$) is spanned by only a few MeV in $B^{1/4}$. For lower *and* higher temperatures the range in $B$ is even smaller.

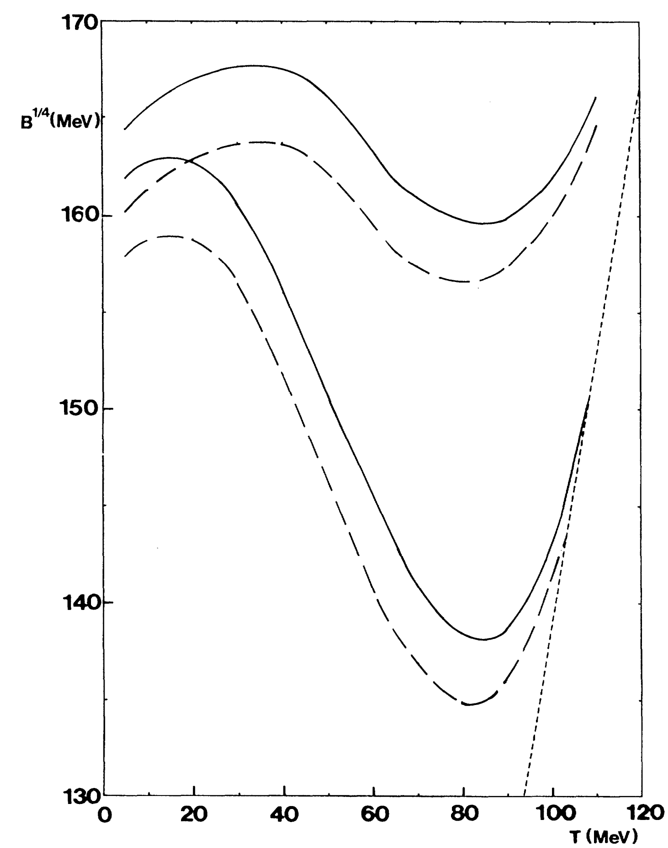


FIG. 2. The critical value of the bag constant $B_{\text{boil}}$ above which boiling of strange-quark matter takes place for baryon numbers below $10^{50}$, is shown as a function of temperature for strange-quark mass $m_s = 100$ MeV (lower full line) and 150 MeV (lower dashed line). Also indicated is the temperature $T_c$ at which the cosmic quark-hadron phase transition takes place (short-dashed line), and the values of $B_{\text{boil}}$ below which the strangeness of nucleated bubbles exceeds 10 (upper full line for $m_s = 100$ MeV and upper dashed line for $m_s = 150$ MeV).

corresponds to a small $\mu_h$ and thus to a small $\Delta P$. From Eq. (6) it is seen that this will result in a large $A_{\rm boil}$ (since the exponential term is clearly dominating), and thus just at the low-$T$ side of the line $T=0.75B^{1/4}$, boiling will not occur. However, it appears from Fig. 2 that this interval becomes smaller with increasing $T$ and $B$.

It was found by Farhi and Jaffe[1] that for $B<(145\ \mathrm{MeV})^4$ $ud$-quark matter is stable at $T=0$ (for the strong coupling constant $\alpha_c=0$). Such stability seems not to be allowed by experiments. Treating therefore $(145\ \mathrm{MeV})^4$ as a lower bound on $B$ one might be tempted to conclude that strange matter would boil for $50\ \mathrm{MeV} \lesssim T \lesssim 105$ MeV. However, as further discussed below, there is no clear-cut relation between the zero-temperature value of $B$ and the high-$T$ limits illustrated in Fig. 2. Furthermore one should keep in mind that all that is known is that $ud$-quark matter is not stable for $A \lesssim 200$, and in addition, the minimum bag constant depends strongly on the model chosen (Farhi and Jaffe, Ref. 1).

So far we have considered purely thermodynamic effects and disregarded the need to get rid of a lot of strange quarks in order to form the neutrons and protons. This can be properly done by allowing for $\Lambda$'s, kaons, and pions in the hadron phase, but it is not clear how to treat the interactions between these particles and the nucleons in Walecka's model. One can argue that if a protobubble in the quark phase contains more than $5-10$ $s$ quarks, the formation of nucleons will be strongly inhibited, since a high-order weak interaction is required. Alternatively, some of the baryon number must initially be incorporated in very massive $\Lambda$ particles, which are energetically less favorable to form. The effect of this is to make the bag constant at which boiling occurs somewhat higher, but how much is not clear. Figure 2 indicates the parameter regime where the strangeness in hadron bubbles exceeds 10 and the effects described above are likely to decrease the rate of boiling.

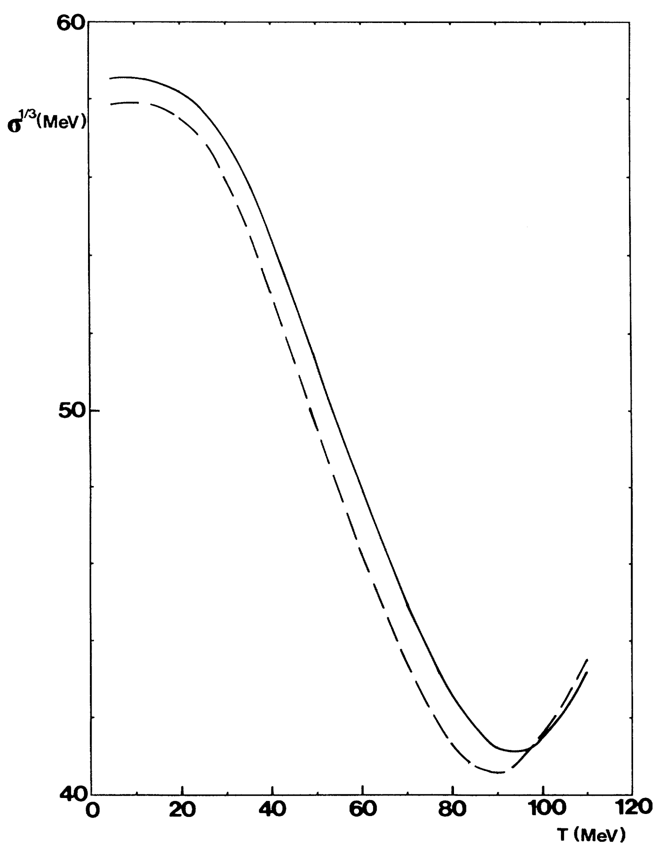


FIG. 3. The critical surface tension $\sigma_{\rm boil}$ above which boiling takes place is shown as a function of temperature for $m_s=100$ MeV (full line) and 150 MeV (dashed line).

In Fig. 3 the *maximum* $\sigma$ for which boiling does not occur is given as a function of temperature, and it is seen that contrary to what was found by Alcock and Olinto,[6] nuggets do survive for small values of $\sigma$. (It should be noted that $\sigma$ is an increasing function of $B$ for fixed $T$, and thus a maximum bag constant automatically leads to a maximum in $\sigma$.) Again, the dependence on $m_s$ is very small. It is worth stressing, that the (perhaps counterintuitive) existence of a *lower* bound on $\sigma$ for boiling to be important results from the surface tension being calculated in terms of the same parameters as are the bulk properties. If the surface tension was varied independently of the bulk thermodynamic potential, boiling could be suppressed by increasing the surface tension.[6]

In Sec. III it was mentioned that it is not clear how to define the bag constant at finite temperatures. The calculations show that the difference between subtracting the thermal electron pressure or not is always less than 1% in $B^{1/4}$, and so this choice is not critical. However, exclusion of the thermal quarks can reduce the bag constant significantly, at least for high temperatures and low values of the bag constant [this can be understood from the fact that a low value of the bag constant corresponds to a low chemical potential, and thus at high temperatures ($T \gtrsim \mu$) the thermal pairs start dominating]. As argued in Sec. III it seems most natural to exclude the thermal electron pressure, but include the total quark-antiquark contributions when calculating $B$ at finite temperatures. But it is not clear how to relate bulk limits on $B(T)$ to low-baryon number measurements of $B(0)$. Therefore it is not clear which values of $B(T)$ can be considered as "allowed" by experiments.

## VI. DISCUSSION AND CONCLUSIONS

A self-consistent treatment of the boiling of strange quark matter into hadrons has been presented, assuming baryon chemical equilibrium across the phase-boundary. Contrary to previous studies treating the hadrons as an ideal gas, the use of Walecka's equation of state for the hadrons, and calculating the surface tension self-consistently led to the conclusion, that strange matter is unaffected by boiling for large parameter intervals for the bag constant $B$ and strange quark mass $m_s$.

The parameter intervals for survival are illustrated in Figs. 1–3. It is not clear how to compare the high-temperature bulk values of $B$ to the zero-temperature values determined from fits to hadron spectra. As discussed above, we have determined $B$ as the sum of pressures from the quarks and antiquarks in the strange matter (including thermal contributions), plus the nonthermal component of the electron pressure. A slightly higher value of $B$ results if one includes the thermal electrons and positrons, whereas subtraction of thermal quarks reduces $B$ significantly, at least at high temperatures and low $B$. The argument for subtracting thermal electrons and positrons, but not quarks and antiquarks is, that a similar thermal electron and positron, but not quark-antiquark, external pressure helps in stabilizing the quark nugget. But clearly, the meaning of $B$ at high temperatures compared to zero temperature is not well defined.

It should also be noted, that the strong-interaction coupling constant in the quark phase was set equal to zero in the calculations. A nonzero value shifts the relevant bag constants to lower values, but we have not performed a detailed study of these effects.

Future work should treat the strangeness problem self-consistently taking the presence of pions, kaons, and $\Lambda$'s in the hadron phase into account. However, it is not clear at present how to treat the high-density interactions properly under these circumstances. That is why we have so far chosen to treat the hadron phase in Walecka's model, only afterward estimating the likelihood of getting rid of the strange quarks. One should also consider nonequilibrium effects, as recently attempted for the formation of strange matter.[13] Nonequilibrium effects presumably suppress evaporation and boiling significantly, and also increase the chance of forming nuggets in the first place.

Independent work in progress along the same lines as discussed in this paper,[14] except treating the nucleons as noninteracting Fermi gases with finite-volume effects, seems to lead to similar conclusions, namely, that quark nuggets may survive boiling.

The evaporation and/or boiling scenario for quark nuggets has recently been generalized to other so-called nontopological solitons.[7,15] One should be aware that complications similar to those discussed in this paper (i.e., interactions, degeneracy, flavor disequilibration, and self-consistent calculation of the surface tension) could be important for some of these systems, depending on their detailed physical properties.

## ACKNOWLEDGMENTS

J.M. wishes to thank Charles Alcock for pointing out that an earlier attempt to calculate the boiling of strange matter for a nonideal hadron gas was in conflict with elementary thermodynamics. M.L.O. acknowledges support from the Danish Research Academy.

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