PHYSICAL REVIEW D
PARTICLES AND FIELDS

THIRD SERIES, VOLUME 32, NUMBER 6 15 SEPTEMBER 1985

Evaporation of strange matter in the early Universe

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(Received 6 May 1985)

Strange matter, a stable form of quark matter containing a large fraction of strange quarks, may have been copiously produced when the Universe had a temperature of ~ 100 MeV. We study the evaporation of lumps of strange matter as the Universe cooled to 1 MeV. Only lumps with baryon number larger than $\sim 10^{52}$ could survive. This places a severe restriction on scenarios for strange-matter production.

INTRODUCTION

A new candidate for the dark matter of the Universe is strange matter.[1] This substance consists of roughly equal numbers of up, down, and strange quarks confined in a quark phase which is *conjectured* to have a lower energy per baryon number than ordinary nuclei. Strange matter is absolutely stable, has a density comparable to that of nuclei and can exist in lumps ranging in size from a few fermis to ~ 10 km. If it is distributed in space in lumps larger than ~ 1 cm, it could close the Universe without ever encountering the Earth and would be astronomically unobservable.[2]

To be a viable candidate for the dark matter of the Universe, a substance must at least pass the following three tests. First, there must be a reasonable theoretical expectation, independent of cosmology, that this form of matter exists. Second, there must be a reasonable scenario for copious production of the substance during some cosmic epoch. Third, the substance must be durable enough to survive until today.

Strange matter has never been seen in the laboratory, yet there are plausible arguments that this form of matter exists. A lump of strange matter contains ordinary quarks (up, down, and strange) and gluons plus a small component of electrons to guarantee charge neutrality. The hadronic material is in a "quark phase" in which nucleons and mesons do not exist and the quarks are free to roam within the lump.[3] Witten[1] suggested that this form of matter could be absolutely stable, i.e., the energy per baryon number (equal to one-third the number of quarks) could be less than 930 MeV. The inclusion of strange quarks is crucial to having a low energy since their presence increases the fermion degeneracy. (Higher-mass quarks, such as the charm quark, are not included since their masses are more than the chemical potentials involved.) A precise calculation of the energy is beyond our present abilities. However, a detailed study[4] has shown that within the uncertainties inherent in a strong-interaction calculation, the existence of stable strange matter is reasonable. Ordinary nuclei would then be unstable but the rate of decay into strange matter would be effectively zero since the decay could only proceed through a very-high-order weak interaction. It is possible that the true ground state of the strong interactions is not iron and that this possibility has been overlooked.

Witten[1] also outlined a scenario for the production of strange matter in the early Universe. The production occurred when the Universe cooled through the QCD phase transition at a temperature T_c (roughly 100—200 MeV), which is characteristic of the strong-interaction energy scale. Above this temperature quarks are unconfined, below it they are confined. As the Universe cooled to a temperature T_c, bubbles of the confined phase appeared and began to expand into the deconfined phase. (There is some reason to believe that the phase transition was first order.) The Universe stayed at the temperature T_c until the confined phase occupied nearly all of space. However, it is possible that a large fraction, say 95%, of the net baryon number condensed into the regions of the deconfined phase. If these regions lost heat primarily through processes such as neutrino emission, which did not deplete their baryon number, they would have settled into nuggets of strange matter. These nuggets would contain most of the baryon number and mass of the Universe.

Witten's scenario has been criticized by Applegate and Hogan,[5] but the physics involved is sufficiently complex that a clear determination of the outcome of the QCD phase transition is unlikely to appear soon. We examine one of the consequences of this phase-transition scenario. We work with the hypothesis that large amounts of strange matter were produced as the Universe cooled through T_c. Our goal in this paper is to examine the fate of a lump of strange matter as the Universe cooled below T_c, i.e., we will examine the *durability* of the lump. By working below T_c we avoid any subtle discussion of the QCD phase transition.

Strange matter has an energy per baryon number of less than 930 MeV but most likely more than 850 MeV.[4] Therefore, the energy I required to liberate one neutron from a lump of strange matter is certainly less than 100 MeV. In *equilibrium,* a lump of strange matter could not exist at a temperature I and at a density comparable to that of the Universe at that temperature. Entropy considerations would simply favor the breakup of strange matter into neutrons. However, a large lump maintained at the temperature I for a brief enough period would evaporate only its outer layer. The Universe cools from 100 to 1 MeV in 1 sec. We find that this is enough time to evaporate all but the largest lumps of strange matter. A detailed analysis gives the minimum baryon number of any lump that survives as roughly 10^{52}. This restriction severely limits the viability of strange matter as a candidate for dark matter.

PROPERTIES OF STRANGE MATTER

In order to calculate the evaporation of strange matter in the early Universe we need to establish some of its properties.[4] A lump of strange matter with baryon number A is a collection of $3A$ quarks in a quark-matter phase which is absolutely stable at zero temperature and external pressure. The mass of a lump of strange matter is ϵA where ϵ is between 850 and 930 MeV. The radius grows as $A^{1/3}$. The baryon number density is roughly $(125\text{ MeV})^3$ or $2.54\times10^{38}/\text{cm}^3$. The minimum value of A for which strange matter is stable is around 10^2. For lower A, shell effects raise the energy per baryon number above the mass of a nucleon and the substance can decay. The maximum A is determined by requiring stability against gravitational collapse into a black hole,[6] $A\sim2.5\times10^{57}$.

Quark matter consists of up, down, and strange quarks and electrons with chemical equilibrium between the flavors maintained by the weak reactions: $d\leftrightarrow u+e+\bar{\nu}$, $s\leftrightarrow u+e+\bar{\nu}$, and $s+u\leftrightarrow d+u$. Electrons are included to guarantee electric-charge neutrality. The chemical potentials of the quarks inside strange matter are typically 300 MeV. A single quark can never be removed from strange matter because of confinement. The baryon number of strange matter can change only through the emission or absorption of baryons such as neutrons, protons, and Λ's.

We define the binding energy I to be the energy needed to liberate one neutron from a zero-temperature lump of strange matter,

$$(A+1)+I\rightarrow(A)+(n)\,, \tag{1}$$

and in the Introduction we argued that I is certainly less than 100 MeV. We focus on the neutron since it is one of the two lightest baryons and unlike the proton it does not face a Coulomb barrier. If a slow neutron strikes a large lump of strange matter, it will fall in and break apart releasing an energy I. We expect the neutron to enter unimpeded for several reasons. The quarks inside the neutron typically have energies above the top of the Fermi sea of the strange matter so there is no Pauli blocking. The neutron interior is in the same vacuum configuration as the strange matter and since we assume a positive surface tension, combining material to make a larger lump is energetically favored. (If the surface tension were negative, strange matter would form tiny balls which would presumably evaporate very easily.) We expect, therefore, a neutron absorption cross section which is geometric:

$$\sigma=f_n 4\pi r_s{}^2=f_n\sigma_0 A^{2/3}\,, \tag{2}$$

where f_n is the absorption efficiency and $f_n\lesssim1$. We will discuss the sensitivity of our results to f_n later. We do not expect $f_n\ll1$ under any circumstances of interest in this calculation. $\sigma_0=3.1\times10^{-4}\text{ MeV}^{-2}$ is obtained by assuming a density of $(125\text{ MeV})^3$.

We can use detailed balance to compute the neutron emission rate for strange matter at a temperature T. First, consider the *equilibrium* situation, at a temperature T, where strange lumps of all baryon numbers are coexisting with neutrons. The reaction $(A+1)\leftrightarrow(A)+n$ will establish the equilibrium relative number densities. Because of (1), the chemical potentials obey

$$\mu(A+1)=\mu(A)+\mu(n)+I\ . \tag{3}$$

For nonrelativistic, nondegenerate particles (we always work at a temperature much less than the neutron mass),

$$\mu=T\ln\left[\left[\frac{2\pi}{mT}\right]^{3/2}\frac{N}{g}\right]\,, \tag{4}$$

where N is the number density, m is the mass of the particle, and g is the internal partition function. For neutrons, $g=2$. For strange matter, the internal partition function $g(A,T)\sim\text{const}\times\exp(aAT/\mu_q)$ where a is of order unity and μ_q is the quark chemical potential inside the lump of order 300 MeV. The ratio $g(A+1,T)/g(A,T)=1$ up to corrections of order T/μ_q. Then (3) and (4) imply

$$\frac{N_nN_A}{N_{A+1}}=2\left[\frac{m_nT}{2\pi}\right]^{3/2}e^{-I/T} \tag{5}$$

for the number densities in equilibrium.

In equilibrium, lumps of strange matter are absorbing and emitting neutrons. The rate per unit volume at which lumps with baryon number $A+1$ are being created is proportional to the density of those with baryon number A times the density of neutrons times the cross section,

$$R[(A)+n\rightarrow(A+1)]=N_AN_nf_n\sigma_0A^{2/3}(T/2\pi m_n)^{1/2}\,, \tag{6}$$

where the last factor is the mean neutron velocity. (This form for the velocity is exact only for large negative chemical potentials. However, it is still approximately correct in our regime.) The rate per unit volume at which lumps of baryon number $A+1$ are disappearing is equal to the density of those lumps times their decay rate r:

$$R[(A+1)\rightarrow(A)+n]=N_{A+1}r \ . \tag{7}$$

The principle of detailed balance tells us that in equilibrium these two rates, (6) and (7), must be equal. Using the equilibrium relation (5) we obtain

$$r=\frac{1}{2\pi^2}m_n T^2 e^{-I/T} f_n \sigma_0 A^{2/3} \ . \tag{8}$$

This evaporation rate has been evaluated by considering an equilibrium situation. However, it is simply the evaporation rate, into neutrons, of a lump of strange matter at a temperature T and Eq. (8) can be used in nonequilibrium situations.

STRANGE MATTER IN THE EARLY UNIVERSE

An object with an evaporation rate dA/dt given by Eq. (8) will evaporate completely in a time

$$\tau(A)=\frac{6\pi^2 e^{I/T}A^{1/3}}{m_n T^2 f_n \sigma_0} \ . \tag{9}$$

To see if a lump survives we should compare this lifetime to the age of the Universe at a temperature T_u. In the radiation-dominated era, the age of the Universe is

$$\tau_u=\left[\frac{3}{32\pi G\rho}\right]^{1/2} , \tag{10}$$

where G is Newton's constant and ρ is the energy density. The energy density is mostly in photons, electrons, and three types of neutrinos for which $\rho=43\pi^2 T_u^4/120$ and

$$\tau_u=\left[\frac{45}{172\pi^3 G}\right]^{1/2}\frac{1}{T_u^2} \ . \tag{11}$$

For a lump to survive until the temperature is roughly its binding energy I, it must have a baryon number A greater than $\sim 10^{55} f_n^3$.

This simple analysis indicates that the minimum baryon number required for survival is extremely large. To confirm this, we present a calculation of the evaporation in the environment of the early Universe which includes the following effects.

(i) *Energy balance:* The heat capacity of the strange matter is very low, and the lump cools rapidly as it evaporates. The lump will cease evaporating unless energy is continuously supplied. The source of this energy is predominantly incoming neutrinos.

(ii) *Diffusion away from the lump:* The emitted neutrons and protons scatter off ambient particles and the baryon-number flow, away from the lump, is constricted. The baryons must push through the surrounding medium.

(iii) *Absorption:* A high concentration of baryons near the surface of the lump will lead to absorption. If the concentration near the surface is too high, the net evaporation could be zero.

(iv) *Obscuration of neutrinos:* The emitted baryons form a blanket around the lump and shield it from incoming neutrinos. This limits the energy supply and the subsequent evaporation.

We will calculate the net evaporation rate dA/dt of a lump of strange matter when the Universe had a temperature T_u. Using Eq. (11) we then get

$$\frac{dA}{dT_u}=\frac{dA}{dt}\frac{dt}{dT_u}=\frac{dA}{dt}\left[\frac{45}{172\pi^3 G}\right]^{+1/2}\left[\frac{-2}{T_u^3}\right] . \tag{12}$$

Given a lump with a certain initial baryon number, this equation can be integrated to trace the evolution of the lump as the Universe cools. We now turn to the calculation of dA/dt.

EVAPORATION IN AN OPTICALLY THIN ENVIRONMENT

An evaporating lump has a lower temperature than the environment. Heat flows from the surrounding medium to the lump, supplying energy for evaporation. At the temperature of interest, $T_u < 100$ MeV, the environment consists primarily of photons, electrons, neutrinos, and their antiparticles. Near the lump there may also be a high concentration of emitted baryons. The neutrinos, which only interact weakly, have the longest mean free paths and are the best energy transporters through the medium. However, in order to heat the object the neutrinos must be absorbed. In this section we study evaporation under conditions for which neutrinos dominate the heating and in addition the environment is transparent to neutrinos. In the next section we will determine the range of values of T_u and A which imply these conditions.

If the medium near the lump is optically thin to neutrinos, then neutrinos from far away can heat the lump. These neutrinos have a thermal distribution at a temperature T_u which is greater than the temperature at the surface of the lump T_s. The rate at which energy flows into a lump of radius r_s is $4\pi r_s^2(7\pi^2/160)T_u^4$ where we have included three generations of neutrinos, i.e., $\nu_e, \bar{\nu}_e, \nu_\mu, \bar{\nu}_\mu, \nu_\tau, \bar{\nu}_\tau$. A neutrino is assured of being absorbed only if its mean free path is smaller than r_s. The incoming neutrinos have energies of order T_u. They scatter off a Fermi gas of quarks with chemical potentials μ_q, much larger than T_u. The mean free path of these neutrinos is roughly $(G_F^2\mu_q^2 T_u^3)^{-1}$ where G_F is the weak-interaction constant. The chemical potentials in strange matter are typically 300 MeV so at a temperature T_u of 10 MeV the mean free path is around 16 m. At this temperature an object of this radius or larger, corresponding to $A\sim 5\times 10^{48}$, will absorb all incident neutrinos.

If the lump is smaller than the neutrino mean free path $l(T_u)$, the probability $p(r_s,T_u)$ of absorption will vary linearly with the radius. If the lump is much larger than $l(T_u)$, then $p(r_s,T_u)=1$. We will use

$$p(r_s,T)=\begin{cases}1, & r_s>\frac{3}{4}l(T)\\ \dfrac{4r_s}{3l(T)}, & r_s\le\frac{3}{4}l(T),\end{cases} \tag{13}$$

where $l(T)=(G_F^2\mu^2T^3)^{-1}$. The factor of $\frac{4}{3}$ in Eq. (13)

arises from an average over solid angle. Equation (13) overestimates $p(r_s,T)$ when $r_s \sim l(T)$. The rate at which heat is absorbed is then

$$L^{\mathrm{abs}} = 4\pi r_s{}^2 \frac{7\pi^2}{160} T_u{}^4 p(r_s,T_u) \ . \tag{14}$$

For small lumps the energy absorbed is proportional to the volume of the lump.

If the lump is larger than a neutrino mean free path, it is black to neutrinos and will also *emit* neutrinos with a thermal spectrum. The rate at which energy is emitted is then $4\pi r_s{}^2(7\pi^2/160)T_s{}^4$ since the lump emits at its temperature T_s. If the lump is smaller than the neutrino mean free path we should include an emission probability factor of $p(r_s,T_s)$. The net heating rate is

$$L = 4\pi r_s{}^2 \left[\frac{7\pi^2}{160} \right] [T_u{}^4 p(r_s,T_u) - T_s{}^4 p(r_s,T_s)] \ . \tag{15}$$

We have not included photons or electrons because we are working in a domain where electromagnetic heating is highly diffusion limited.

The lump of strange matter is losing its baryon number through the emission of nucleons, i.e., neutrons and protons. The neutron evaporation rate is just Eq. (8):

$$r_n = \frac{m_n T_s{}^2}{2\pi^2} e^{-I/T_s} f_n \sigma_0 A^{2/3} \ . \tag{16}$$

The proton evaporation rate can be determined in much the same way as the neutron evaporation rate. However, the proton faces a Coulomb barrier which is equal to the electron chemical potential μ_e in strange matter. This is typically a few tens of MeV. A proton incident on strange matter must have a kinetic energy larger than μ_e or it will not be absorbed (except by penetrating the barrier, a small effect). Thus, the absorption efficiency for protons f_p is energy dependent and the evaporation rate r_p has a somewhat different form than r_n. We will make the approximation

$$r_p = \frac{m_n T_s{}^2}{2\pi^2} e^{-I/T_s} f_p \sigma_0 A^{2/3} \ , \tag{17}$$

where f_p is a constant absorption efficiency less than f_n, i.e., $0 < f_p \lesssim f_n \lesssim 1$. Our conclusions will turn out to be insensitive to f_p and f_n so this approximation is valid.

Outside the lump the neutrons and protons are strongly interacting amongst themselves and the protons are electromagnetically interacting with the electrons, positrons, and photons. The nucleons are in local thermal equilibrium with the electromagnetic quanta. We can establish the local densities of neutrons and protons using pressure equilibrium. The contribution to the pressure from electrons, positrons, and photons is $(11\pi^2/180)T^4$ where T is the local temperature, $T_s \le T \le T_u$. The neutrons and protons contribute $(N_n+N_p)T$ to the pressure where N_n and N_p are the local densities. Far from the lump there are no nucleons so

$$(N_n+N_p)T + \frac{11\pi^2}{180} T^4 = \frac{11\pi^2}{180} T_u{}^4 \ . \tag{18}$$

We have not included neutrinos in Eq. (18) because the material near the lump is assumed to be thin to neutrinos. For quickly evaporating lumps the weak interactions do not have time to establish the equilibrium relationship $N_n = N_p$. However, at temperatures as low as 20 MeV, there are enough pions around to establish equilibrium through strong interactions like $\pi^+ + n \to p + \pi^0$. Therefore, at the surface of the lump the densities of nucleons are

$$N_n(T_s) = N_p(T_s) = \frac{11\pi^2}{360 T_s} (T_u{}^4 - T_s{}^4) \ . \tag{19}$$

These nucleons, near the surface, can be reabsorbed by the lump. From Eq. (2) the absorption rates are

$$\begin{aligned} r_n^{\mathrm{abs}} &= f_n N_n(T_s) \sigma_0 A^{2/3} (T_s/2\pi m_n)^{1/2} \ , \\ r_p^{\mathrm{abs}} &= f_p N_p(T_s) \sigma_0 A^{2/3} (T_s/2\pi m_n)^{1/2} \ , \end{aligned} \tag{20}$$

where the last factor is the mean thermal velocity of the nucleons near the surface. The net evaporation rate, emission (16) and (17) minus absorption (20), is

$$\frac{dA}{dt} = \left[\frac{m_n T_s{}^2}{2\pi^2} e^{-I/T_s} - \frac{11\pi^2}{360 T_s} (T_u{}^4 - T_s{}^4)(T_s/2\pi m_n)^{1/2} \right] \times \sigma_0 A^{2/3} (f_n + f_p) \ . \tag{21}$$

The mean kinetic energy of an emitted nucleon is $2T_s$ (this form for the kinetic energy is exact only for large negative chemical potentials; however, it is still approximately correct in our regime) and the energy required to release each nucleon is I. Therefore, energy is being expended at the rate

$$L = \frac{dA}{dt}(I + 2T_s) \ . \tag{22}$$

By energy balance this equals the incoming power given by Eq. (15). This equality allows us to determine the surface temperature T_s of a lump with baryon number A given the temperature of the Universe T_u and the binding energy I.

Once we know T_s as a function of A and T_u we know dA/dt, by Eq. (21), as a function of A and T_u for a given I. Since dA/dT_u is simply determined by Eq. (12), we can integrate, starting at some initial high temperature and find the fate of a lump as the Universe cools.

DIFFUSIVE HEATING

In the preceding section we assumed that the material surrounding the evaporating lump forms a layer which is transparent to neutrinos. We now want to see under what conditions that assumption is valid. A lump which starts significantly evaporating at some temperature T_i when it has a baryon number A_i can at most emit A_i nucleons and will typically be surrounded by roughly A_i nucleons. The number density of nucleons goes like $(T_u{}^4 - T^4)/T$

where T is the local temperature. T is not very different from T_u or T_s so the typical nucleon density is T_u^3. These nucleons will extend out to a radius of order $A_i^{1/3}/T_u$. Incoming neutrinos scattering off these nucleons have a mean free path of $\sim(G_F^2T_u^5)^{-1}$. For the surrounding medium to be thin, we require $(G_F^2T_u^5)^{-1}$ to be greater than $A_i^{1/3}/T_u$.

We will compute the heating rate in this regime by assuming a stationary state has been achieved for each T_u as the Universe cools. This is not strictly correct, and probably results in a slight underestimate of the heating rate. However, for reasons that will become clear below, this error is not important.

Since the material is optically thick to the neutrinos, the pressure due to neutrinos must be included in calculating the pressure balance. Equation (18) becomes

$$(N_n+N_p)T+\frac{43\pi^2}{360}T^4=\frac{43\pi^2}{360}T_u^4 . \tag{23}$$

The energy flux due to a neutrino of type i ($i=\nu_e,\bar{\nu}_e,\nu_\mu,\bar{\nu}_\mu,\nu_\tau,\bar{\nu}_\tau$) is proportional to the temperature gradient and to the mean free paths l_i. The rate $R_{ia}(E)$ at which neutrinos of type i and energy E scatter off particles of type a ($a=n,p,e,\bar{e}$) can be calculated using the standard electroweak theory. The mean free path $l_i(E)$ is then given by

$$l_i(E)^{-1}=\sum_a R_{ia}(E) . \tag{24}$$

The cross section for neutrino nucleon scattering is given by $C_{in}G_F^2E^2$ or $C_{ip}G_F^2E^2$ where C_{in} and C_{ip} are calculable coefficients. The nucleons are taken to be at rest since $T \ll m$. The neutrino nucleon reaction rate is then

$$R_{in}+R_{ip}=(C_{in}+C_{ip})G_F^2E^2N_n . \tag{25}$$

The neutrino lepton cross sections depend on the lepton energy so these reaction rates must be computed by integrating over the Fermi distributions of lepton energies. This gives

$$\sum_{a=\text{leptons}} R_{ia}=D_iG_F^2T^4E , \tag{26}$$

where D_i as well as $C_{in}+C_{ip}$ are given by the standard weak-interaction theory. These calculations have all neglected the Pauli blocking associated with finding the final states occupied. This reduces the rates by roughly 10% and we include this factor at later stages.

Neutrino-neutrino scattering has a different effect on the calculation. It is clear that scattering between neutrinos of the same type i has no effect on the flux carried by neutrinos of type i; this is a consequence of momentum conservation. Furthermore, scattering between neutrinos of different type but identical $l_i(E)$ has no effect on their fluxes. The effect of collisions between the neutrinos of different type and different $l_i(E)$ is to increase the flux carried by the type with smaller $l_i(E)$ and decrease the flux carried by the other type. This effect may be approximately modeled by averaging the collision rates over the neutrino types. The average mean free path is then

$$l_A^{-1}(E)=\langle l_i^{-1}(E)\rangle=\tfrac{1}{6}\sum_i\sum_a R_{ia}(E) . \tag{27}$$

The flux carried by all six neutrino types is

$$F=-2\frac{dT}{dr}\int_0^\infty l_A(E)\frac{\partial}{\partial T}\left[\frac{1}{e^{E/T}+1}\right]\frac{4\pi E^3dE}{(2\pi)^3} . \tag{28}$$

This formula is the fermionic counterpart of the Rosseland approximation.[7] The total flux can be written as

$$F=-\psi(T)\frac{dT}{dr} , \tag{29}$$

where $\psi(T)$ is determined by Eq. (28). $\psi(T)$ depends on T_u through the neutron and proton number densities given by Eq. (23).

The assumption of stationarity implies that

$$F=-\frac{L}{4\pi r^2} , \tag{30}$$

where L is the total neutrino heating rate. L is determined by integrating (29) from the surface of the lump to infinity, using (30). The result is

$$L=\frac{r_s}{G_F^2T_u}\Phi\left[\frac{T_s}{T_u}\right] , \tag{31}$$

where Φ is a dimensionless function given by

$$\Phi\left[\frac{T_s}{T_u}\right]=\frac{4T_u}{\pi}\int_0^\infty dx\int_{T_s}^{T_u}dT\frac{x^3e^x}{(e^x+1)^2}\times\frac{T^2}{a(T_u^4-T^4)x+bT^4} . \tag{32}$$

The constants a and b are determined by Eqs. (23) and (25)–(27).

The evaporation rate is determined as in Eq. (21), except that neutrino pressure is taken into account,

$$\frac{dA}{dt}=\left[\frac{m_nT_s^2}{2\pi^2}e^{-I/T_s}-\frac{43\pi^2}{720T_s}(T_u^4-T_s^4)\left[\frac{T_s}{2\pi m_n}\right]^{1/2}\right]\times\sigma_0A^{2/3}(f_n+f_p) . \tag{33}$$

The heat lost in evaporation is given by Eq. (22), and as before, Eqs. (22) and (31)–(33) determine the temperature T_s, L, and dA/dt.

An important property of Eq. (31) is that the heating rate varies inversely with T_u. This may be understood as follows. The flux is proportional to the gradient of T^4 ($\sim T^4/r$), inversely proportional to the number density of scatterers ($\sim T^3$), and inversely proportional to a typical cross section ($\sim T^2$). The luminosity is $4\pi r^2$ times the flux. The product of these terms is $\sim r/T$.

Since the heating rate varies as T_u^{-1}, and the age of the Universe varies as T_u^{-2}, it is clear that, in this regime, most of the evaporation occurs at cooler temperatures. This means that the total evaporative loss of a strange lump that starts evaporating in this regime is insensitive to the initial T_u.

We can check the stationary solution to see if the region over which the temperature changes from T_s to $\approx T_u$ is optically thick to neutrinos. This condition is $G_F{}^2 r_s T_u{}^5 > 1$, a more stringent condition than what we estimated earlier by requiring that enough baryons could have been emitted to make the surroundings optically thick. This line is shown in Fig. 1.

ELECTROMAGNETIC HEATING

For sufficiently small lumps, the neutrino heating rate will fall below the heating rate due to photons, electrons, and positrons. This heating rate is difficult to calculate precisely because it depends on the number of baryons that have been previously emitted. By neglecting the electromagnetic heating we have underestimated the heating rate and therefore the evaporation rate particularly for small lumps. This means that the low-A part of the curves in Fig. 1 should fall even faster, but this has no effect on our conclusions, so we can safely neglect electromagnetic heating.

FLAVOR EQUILIBRATION

High-temperature lumps of strange matter emit neutrons and protons from the surface. Neutrons and protons are made of up and down quarks whereas the equilibrium configuration of strange matter has comparable numbers of up, down, and strange quarks. The lightest strange baryon, the Λ, is too heavy (1116 MeV) to be emitted. For the evaporation to proceed, up and down quarks must be supplied to the surface. As the number of up and down quarks drops, the strange quarks convert to ups and downs via the weak interactions. The rate at which the strange quarks convert places an upper limit on the total evaporation rate.

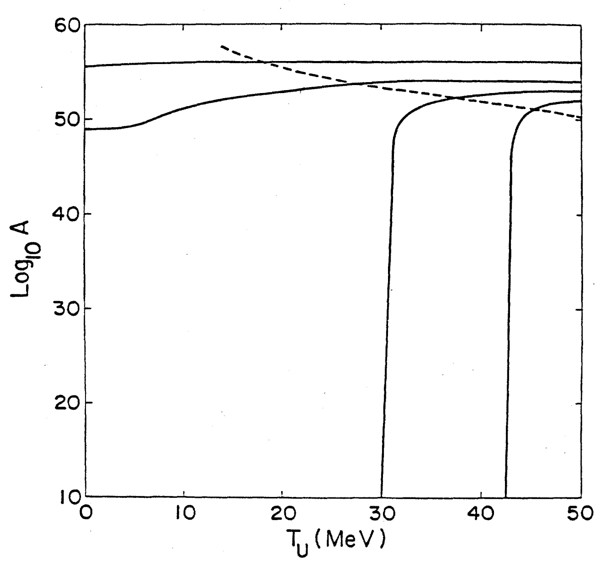

FIG. 1. The variation of A with T_u for four lumps of strange matter. The transition from diffusive neutrino heating to optically thin neutrino heating is shown by the dashed line. The binding energy is 20 MeV for all four lumps.

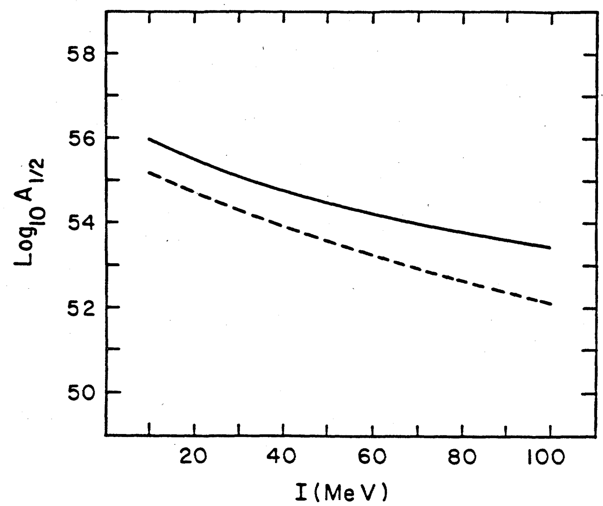

FIG. 2. The initial baryon number $A_{1/2}$, for a strange lump that loses 50% of its baryon number, plotted vs binging energy I. The solid curve is for $f_n + f_p = 1$, the dashed curve for $f_n + f_p = 0.1$.

The strange quarks have a chemical potential $\mu_q \approx 300$ MeV and the rate for a single quark to convert is roughly $G_F{}^2 \mu_q{}^5 \sin^2\theta_c$ where $\sin^2\theta_c \approx 0.04$. The total number of strange quarks is A, so we get the limit

$$\left| \frac{dA}{dt} \right| \le G_F{}^2 \mu_q{}^5 \sin^2\theta_c A \ . \tag{34}$$

Using Eq. (12) we can get a limit on the slope of the curves appearing in Fig. 2,

$$\frac{d\,\log_{10} A}{dT} \le 10^{10} \left[\frac{1\ \mathrm{MeV}}{T} \right]^2 \frac{1}{T} \ . \tag{35}$$

This slope is very large and places no significant restriction on any of our calculations.

THE EVAPORATION PROCESS

We show the evolution of several evaporating strange lumps in the A-T_u plane in Fig. 1. We assume a binding energy of 20 MeV in Fig. 1. The computations start at $T_u = 50$ MeV, with different initial A. The important features of the process are clear. Very little evaporation occurs in the optically thick region. Most of the evaporation occurs just as the surrounding nucleons become thin to neutrinos. There is a critical initial A_c below which the lump evaporates completely, a narrow range in A for which substantial but not total evaporation occurs, and large lumps suffer little evaporation.

We show the dependence of our results on the binding energy I in Fig. 2. We plot the initial A that results in 50% evaporation versus I. If the binding energy is large, the strange matter is more secure against evaporation. We also show the sensitivity of our results to $f_n + f_p$. The upper curve is for $f_n + f_p = 1$, the lower is for $f_n + f_p = 0.1$. Since we expect $f_n + f_p$ to be of order 1, the lower curve gives a reasonable lower bound on the minimum A which could survive. Even for the largest binding energies we consider, 100 MeV, the minimum size lump which could survive as the Universe cooled below 50 MeV is 10^{52}.

CONCLUSIONS

We have calculated the evaporation of strange matter *after* the QCD phase transition, without making assumptions about that transition, and we conclude that the minimum size lump that could survive has a baryon number $A \sim 10^{52}$. Such an object has a planetary mass.

These numbers are large in an important sense cosmologically. If we assume the Universe is closed by baryons in either the strange or normal phase, then the mean baryon number in the particle horizon at these early epochs is $\sim 10^{55}$ $(1\ \mathrm{MeV}/T_u)^3$. At $T_u = 50$ MeV this number is $\sim 8 \times 10^{49}$, much smaller than the minimum baryon number of a lump which could survive. This means that the process that leads to the formation of strange-matter lumps must involve large perturbations in the baryon number on the horizon scale. A mechanism for producing this perturbation is not known to us.

ACKNOWLEDGMENTS

We would like to thank R. L. Jaffe, J. Negele, and especially A. Olinto for valuable discussions, and the Alfred P. Sloan Foundation. This work was supported in part through funds provided by the U.S. Department of Energy (DOE) under Contract No. DE-AC02-76ER03069, and by NASA Grants Nos. NSG-7643 and NGL 22-009-638.

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