Baryon isocurvature fluctuations at small scales and baryonic dark matter

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A model of baryogenesis is described which gives rise to fluctuations in baryon number that are large on small scales but of low amplitude on large scales. This provides a mechanism for primordial black-hole formation, and allows the possibility of a critical density of dark matter in baryonic (and antibaryonic) form. Since high-density regions naturally possess both signs of baryonic excess, our model also predicts a small fraction of the mass density of the Universe to be in the form of compact antibaryonic regions. These objects may be observable via a redshifted annihilation feature in the diffuse extragalactic γ-ray background, as both steady sources of γ-ray radiation and annihilation-powered γ-ray bursters at cosmological distances, by a distortion of the spectrum of the 3K background radiation, or by the presence of antihelium nuclei as a rare component of high-energy cosmic rays.

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I. INTRODUCTION

The large-scale structure in the Universe is assumed to be formed by gravitational instability from initially small primordial density inhomogeneities. The origin of these inhomogeneities remained unknown even a decade ago. A breakthrough was achieved with the development of the inflationary universe scenario, in which the exponentially rapidly rising scale factor results in the transformation of small-scale (both in amplitude and wave length) fluctuations of quantum fields into cosmologically interesting density perturbations [1–4]. Energy density fluctuations generated by the inflaton field are of sufficiently large amplitude to account for structure formation. (In fact, they are too large with a natural strength of the inflation coupling and special care must be taken to reduce them to a sufficiently low value.) This mechanism generates so-called *adiabatic* perturbations when inhomogeneities in the matter energy density are about the same magnitude as those in the radiation density. The perturbation spectrum is of the Harrison-Zel'dovich form that is (practically) scale independent.

Unfortunately, these two features make the simplest models questionable. Adiabatic perturbations are strongly bounded from above by the strict limits on angular fluctuations of the background radiation temperature [5] and the scale-independent spectrum poorly describes the observed structure on large ($\gtrsim 10$ Mpc) scales [6], at least when normalized in the conventional way to fluctuations in the galaxy counts. These shortcomings of an otherwise elegant model have stimulated consideration of other possibilities which could give rise to a nonflat spectrum and/or to isocurvature (or isothermal) fluctuations. The latter are fluctuations in the chemical content of the primeval plasma [e.g., fluctuations in baryonic charge (B) density] with a constant total energy density. Baryon number fluctuations are transformed into energy density fluctuations during the later stages of the evolution of the Universe when baryons (or, generally speaking, any other particles) become nonrelativistic.

A number of models for creation of isocurvature fluctuations have been proposed during the last decade [7–22] with a variety of spectra ranging from flat to those possessing a prominent peak at a particular wavelength. We will show in what follows that a specific spectrum of (isocurvature) baryonic charge-density perturbations can be generated in some scenarios of baryogenesis. We show, in particular, that there can exist very large fluctuations in baryonic charge density at sufficiently low spatial scale that their existence does not contradict the bounds on $\Delta T/T$. Moreover, the sign of baryon asymmetry can be different in different regions of high baryonic density, so that we can expect both high density and small size baryonic and antibaryonic regions. There is no observational difference between the two if the density is so high that those regions collapsed at some early epoch into compact stellar remnants and black holes. In fact, the model described here presents a new mechanism for early black-hole formation from large amplitude isothermal fluctuations at small spatial scales. In this case, dark matter in the Universe would be in the form of baryonic (and antibaryonic) black holes. Smaller uncollapsed bubbles of antibaryonic matter would be observable either as pointlike sources of γ radiation or, if they annihilated earlier, as some bright spots in the otherwise isotropic background radiation. If the number density of these objects were sufficiently high, early $p\bar{p}$ annihilation could result in the distortion of the spectrum of background radiation. Unfortunately, there is too much freedom in the model to make any specific predictions. The amount of uncollapsed antimatter may vary from an unnoticeable amount to an amount in contradiction with existing data.

We will consider two possible mechanisms for the realization of such scenarios. One is based on spontaneous charge symmetry breaking in a first-order phase transition. This mechanism proves to be considerably less effective in creating a baryon asymmetry than a second one that we give, but we still present it since the features may be of more general interest.

Another possible mechanism for generation of a large baryon asymmetry in a small fraction of space is based on the model of baryogenesis by the baryonic charge condensate $\langle \chi \rangle$ [32]. Our only modification of the original scenario is an introduction of a rather generic coupling of χ to the inflaton field Φ. This coupling permits generation of a large value of $\langle \chi \rangle$ in a reasonably small part of space, while in the rest of the space $\langle \chi \rangle$ would have a much smaller value. Except for this difference, baryogenesis proceeds along the same lines in both phases. The net result for the baryon asymmetry can differ by several orders of magnitude.

For the realization of both mechanisms, it is necessary that the bubbles of the broken phase neither coalesce nor disappear by the epoch of baryogenesis. Since the size of the baryon asymmetry of the Universe is proportional to the amplitude of *C*, *CP* violation, it should be different inside and outside the bubbles. In fact, if spontaneous symmetry breaking is the only source of *C*, *CP* violation the baryon asymmetry could only be generated inside the bubbles while outside of them the Universe would be charge symmetric. If, however, there is another source of *C*, *CP* violation, namely an explicit one, the asymmetry would be generated in both phases while the size of the asymmetry might be very different. After baryogenesis is over, the bubble walls can, and in fact must, disappear so as to avoid excessive inhomogeneities in the cosmological energy density. In other words, after baryogenesis, the phase transition should proceed in the reverse direction restoring the charge symmetry of the vacuum state. The temporary existence of the broken phase remains imprinted, however, in inhomogeneities in the baryon number density. Thus we end up with a homogeneous background baryon asymmetry coming from explicit *C*, *CP* violation and strong but presumably short-scale inhomogeneities inside the former bubbles of the broken phase. To avoid possible confusion, we mention that we use the term "bubble" both for the bubbles of the broken symmetry phase separated from the symmetric (false) vacuum state by a domain wall and for the high baryonic number density remnants of such bubbles. The characteristic wavelength of the baryon density perturbations is determined by the onset of the phase transition. To make the scale cosmologically significant, the phase transition should take place during the inflationary stage. Such a scenario has been considered in Refs. [14,15].

There are many features in common between these mechanisms of generation of large baryonic charge-density fluctuations. Both are based on a specific phase transition going back and forth once in the early universe. The same kind of coupling of the underlying scalar field to the inflaton field is necessary in both cases. As a result, the bubble evolution is essentially the same in both cases.

The paper is organized as follows. In the next section the first model based on spontaneous breaking of charge symmetry is discussed. The kinetics of the bubble formation is considered in Sec. III. The results of this section are applicable both to the first model as well as to the case of the baryonic charge condensate. The latter is described in Sec. IV. The evolution of the bubbles with high baryon number density, after the bubble walls that separate two different vacuum states have disappeared, is considered in Sec. V. Cosmological implications and conclusions are presented in Sec. VI.

II. CHARGE SYMMETRY BREAKING

A mechanism which gives rise to the bubbles of high baryon number density can be realized as follows. Assume that there exist two sources of charge symmetry breaking, spontaneous and explicit. If the spontaneous symmetry breaking proceeded through the first-order phase transition which started at the inflationary stage and did not terminate during baryogenesis, one would expect that *CP*-odd amplitudes were not universal space independent quantities but had different values inside and outside bubbles of broken phase. Correspondingly, the size of the baryon asymmetry is different inside and outside the bubbles. The resulting picture depends upon the characteristic bubble size l_B and upon the relation between the amplitude of explicit *CP* breaking, ϵ_e, and that of spontaneous *CP* breaking, ϵ_s. Onset of the phase transition during the inflationary stage permits l_B to be on a scale that is astronomically large. The value of l_B is very much parameter dependent and can be larger or comparable with the horizon size l_h or much smaller, down to stellar size or even below. The first case has been considered in Refs. [14,15] where the island universe model was discussed. Here we concentrate mostly on the case $l_B \ll l_h$, so that the baryonic charge inhomogeneities are observable inside our universe.

If the amplitude of spontaneous symmetry breaking is small in comparison with the explicit one, $\epsilon_s < \epsilon_e$, there would be relatively small fluctuations in the baryonic number density over the uniform baryonic background. In the other case, $\epsilon_e < \epsilon_s$, there would be denser baryonic or antibaryonic regions in the low-density baryonic background. While we cannot make specific predictions, we can use observations to constrain the model. The absence of visible $p\bar{p}$ annihilation sets strong limits on bubble size and bubble separation [23].

Spontaneous *CP* violation [24] is achieved by a condensate of complex scalar field ϕ. We assume that its potential has the usual form

$$U(\phi)=m_{\rm eff}^2|\phi|^2+\lambda|\phi|^4\ln\frac{|\phi|^2}{\sigma^2} \ . \tag{1}$$

The logarithmic factor in this expression comes as a result of summing higher loop corrections [25]. There should be, of course, some other terms either in $U(\phi)$ or in Yukawa coupling of ϕ to fermions which are not invariant with respect to phase rotation $\phi\to\phi\exp(i\alpha)$ so that *CP* breaking becomes operative. The effective mass factor $m_{\rm eff}^2$ should generically include the contribution from different kinds of interaction of ϕ with other fields:

$$m_{\rm eff}^2 = m_0^2 + \xi R + \beta T^2 + \lambda_1(\Phi - \Phi_1)^2 . \quad (2)$$

Here the first term m_0^2 is the vacuum mass of ϕ (barring the contribution from $\lambda_1\Phi_1^2$). The term ξR comes from the possible nonminimal coupling of ϕ to the curvature scalar. During the inflationary stage $R = 12H_I^2$, in the radiation-dominated stage $R \approx 0$, and in the matter-dominated stage $R \approx t^{-2}$. The third term is a result of temperature corrections to the effective potential which became effective when the Universe was heated as a result of inflaton decay (the Gibbons-Hawking temperature effects are hidden in a redefinition of ξ). There should be some other temperature corrections to $U(\phi)$ but they are not essential for our purpose. The last term comes from the general renormalizable coupling of ϕ to the inflaton field

$$L_{\rm int} = \lambda_1|\phi|^2\Phi^2 + g|\phi|^2\Phi .$$

There can also be terms of the form ($\tilde{g}\phi\Phi^2$+H.c.) which are not considered here. Here Φ_1 is a constant. We assume for definiteness that the chaotic inflationary scenario is valid [26] (though our model is compatible with other inflationary scenarios), so that Φ evolved from large values $\Phi > m_{\rm Pl}$ and $\Phi > \Phi_1$ down to zero.

With all these terms taken into account, $m_{\rm eff}^2$ can possess the rather peculiar oscillatory behavior (shown in Fig. 1) which is necessary for realization of our model. To get this nonmonotonic time dependence, some not particularly strong fine tuning might be necessary. All the parameters in Eq. (1) may have either sign. We assume, however, that λ_1 and Φ_1 are positive to ensure a minimum of $m_{\rm eff}^2$ during inflation. Numerically, Φ is very large in the chaotic inflation scenario, $\Phi \gg M_{\rm Pl}$. We assume that Φ_1 is also large, $\Phi_1 > m_{\rm Pl}$, so that Φ becomes equal to Φ_1 in the course of inflation. On the contrary, the coupling constant λ_1 should be very small, $\lambda_1 < 10^{-6}$. Otherwise, loop correction would induce effective $\lambda_2\Phi^4$ interactions with $\lambda_2 \approx \lambda_1^2 > 10^{-12}$. This in turn would give rise to unacceptably large energy density fluctuations [1–4]. As we see below, λ_1 should be even smaller ($\lambda_1 < 10^{-10}$) to ensure a successful realization of our second scenario. This fits the notion of a very weak coupling of the inflaton field.

Little is known about numerical values of ξ. Conformal invariance (which is absent in our model) implies $\xi = \frac{1}{6}$, while for Goldstone bosons (also absent) $\xi = 0$ [27]. Radiative corrections to minimal gravitational coupling of ϕ give rise to generically nonzero ξ proportional to a power of essential coupling constants which gives a rather small result $\xi \lesssim 10^{-2}$, although $\xi = O(1)$ is not excluded. The temperature correction term βT^2 arises from interactions of ϕ with the thermal bath and is proportional to the interaction strength. β is of the order of λ for the $\lambda\phi^4$ interaction and of the order of the coupling constant squared for Yukawa or gauge couplings. Usually $\beta > 0$, but it is also possible for β to be negative [28]. A reasonable value for λ is about 10^{-2} as is typical for supersymmetric models.

To ensure the behavior shown in Fig. 1, the following conditions should be satisfied:

$$m_0^2 + \xi R < 0 , \quad (3a)$$

$$m_0^2 + \lambda_1\Phi_1^2 + \xi R > 0 , \quad (3b)$$

$$m_0^2 + \lambda_1\Phi_1^2 + \beta T^2 > 0 , \quad (3c)$$

$$m_0^2 + \lambda_1\Phi_1^2 < 0 . \quad (3d)$$

Note that R might be different in Eqs. (3a) and (3b). We assume that m_0^2 is of the order of H_I^2 and, with $H_I = (10^{-5} - 10^{-6})m_{\rm Pl}$ and $\lambda_1 = 10^{-10} - 10^{-12}$, all the terms in Eqs. (3) are more or less of the same order of magnitude. A naive estimate of the reheating temperature is $T_{\rm rh} \approx \sqrt{m_{\rm Pl}H_I}$ but a weak coupling of the inflaton to matter makes $T_{\rm rh}$ proportional to the Yukawa coupling constant $g \sim \lambda_2^{1/4} \leq 10^{-3}$. Hence, $\beta T_{\rm rh}^2$ is similar to other terms in (3).

At the beginning of inflation, Φ is large and so $m_{\rm eff}^2 > 0$. The potential $U(\phi)$ (1) has a minimum at $\phi = 0$ and charge symmetry is unbroken. With diminishing $m_{\rm eff}^2$,

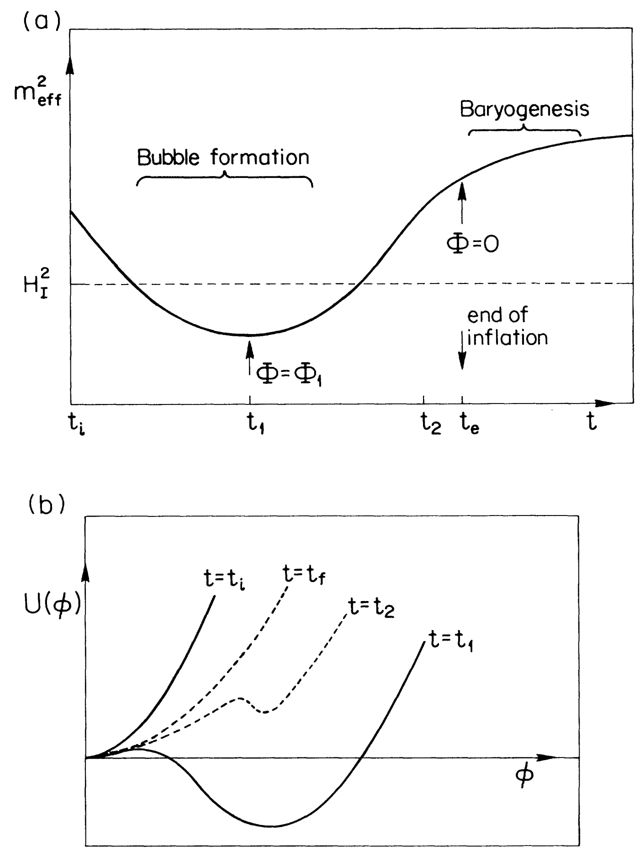

FIG. 1. (a) The evolution of the effective mass of the baryon number-generating field ϕ driven by the inflaton field Φ. (b) The shape of the potential $U(\phi)$ for different values of the effective mass of ϕ. In the first scenario with spontaneous charge symmetry breaking, baryogenesis should proceed while $U(\phi)$ still has the second minimum at $\phi \neq 0$ while baryogenesis by the scalar field condensate starts when this minimum disappears.

another minimum at $\phi \neq 0$ is formed which gradually becomes deeper than the one at $\phi = 0$. Still, the transition probability between the two minima is exponentially suppressed. Even for $m_{\rm eff} = 0$, when the barrier between two extremes of $U(\phi)$ is absent, the transition probability is very small. For the case of the $-\lambda|\phi|^4$ potential, it is determined by the so-called Fubini instanton and is proportional to $\exp(-8\pi/3\lambda)$. For the Coleman-Weinberg potential [Eq. (1)] the transition probability is slightly larger but still negligible [29] for small λ.

For negative $m^2_{\rm eff}$, the suppression is absent, and the phase transition is of second order without bubble nucleation and supercooling. If, however, $m^2_{\rm eff}$ is negative only during a finite time, the phase transition may not be accomplished, and only a few bubbles of new phase are formed. When $m^2_{\rm eff}$ becomes negative, the process starts as a second-order phase transition with ϕ moving to the bottom of the potential according to

$$\phi = \phi_i \exp\left\{ \left[\left[\frac{9}{4}H^2 + |m^2_{\rm eff}| \right]^{1/2} - \frac{3}{2H} \right] t \right\} , \tag{4}$$

where ϕ_i is the initial amplitude of (quantum) fluctuations with sufficiently long wavelength l_i. The latter scales during inflation as $l = l_i \exp(Ht)$. When $m^2_{\rm eff}$ once again becomes positive, the local maximum at $\phi = 0$ turns back into a local minimum and, if ϕ is smaller than the critical value $\phi_{\rm cr} \approx m/\lambda^{1/2}$, it should roll back to $\phi = 0$. If during the period Δt, when $m^2_{\rm eff} < 0$, ϕ rises above $\phi_{\rm cr}$, bubbles of new phase may form. It is evident that the probability of bubble formation must vary between 0 and 1 depending upon the behavior of $m_{\rm eff}$. The production probability as a function of time depends upon how fast the inflaton field Φ evolves near Φ_1.

Two limiting cases can consequently be considered. First, W is a sharp function of time. In this case, the distribution in bubble sizes is also narrow, strongly peaked near a particular value of l_B. In another limiting case, when the inflaton evolution is slow in comparison with the rate of rise of ϕ [Eq. (4)], the bubble production probability is a slowly varying function of time for a (model-dependent) time interval Δt. That case corresponds to the situation when the fluctuations of ϕ during this period reached and exceeded the critical value ϕ_c. The earlier a bubble formed, the larger would be its size and the larger would be the average distance between the bubbles. Correspondingly, their distribution in l is

$$\frac{dN}{dl} \sim l^{-4+\kappa} , \tag{5}$$

where κ corrects for the possibility of a larger probability of earlier bubble production.

When Φ evolves away from Φ_1, the effective mass squared becomes positive and the potential barrier reappears. This barrier separating the two minima should exist at least until baryogenesis so that the symmetry is not restored. It is achieved if $m^2_{\rm eff}$ at this stage [Eqs. (3b) and (3c)] is positive but not too large. In Fig. 1(b) the moment of baryogenesis in this scenario corresponds to $t = t_2$. Subsequent cooling would again lead to $m^2_{\rm eff} < 0$ if condition (3d) is valid. The symmetry would be broken over all space and the domain walls would disappear. Another possibility to get rid of the walls might be realized if $m^2_{\rm eff} = m_0^2 + \lambda\Phi_1^2 > 0$ and is so large that the second minimum at $\phi = 0$ vanishes and symmetry is restored over all space (but after the baryogenesis epoch). We show in the next section that quantum fluctuations of ϕ in the de Sitter background are essential for the dynamics of the bubble formation in the case of $|m_{\rm eff}| < H$, and we will calculate more accurately the probability of their production and the bubble spectrum.

To summarize, we see that it is possible (and perhaps even natural) to obtain at the end of inflation a peculiar space pattern of CP violation, when in relatively small and chaotically but uniformly distributed pieces of space, CP-odd amplitudes are rather large, while in the rest of space they are (much) smaller. The details of the picture such as bubble size, number density, the relative values of CP-odd terms, etc. are very much model dependent. Moreover, there is no theoretical input for the parameters of the model, and so we are unable to calculate the relevant quantities. Instead, we adopt a different approach, namely we assume that such a scenario was realized and discuss the observational consequences.

III. BUBBLE FORMATION AND SPECTRUM

We assume that initially $m^2_{\rm eff}$ is large and positive so that $U(\phi)$ [Eq. (1)] has only one minimum at $\phi = 0$ and there is no ϕ condensate. Initially, a large value of $m^2_{\rm eff}$ is naturally realized in the chaotic inflationary scenario due to the term $\lambda_1(\Phi - \Phi_1)^2$. In particular, in the chaotic inflation scenario, the inflaton field Φ evolves from a very large value $\Phi > m_{\rm Pl}$ down to zero. When $\Phi \approx \Phi_1$, the contribution from the term $\lambda_1(\Phi - \Phi_1)^2$ into $m^2_{\rm eff}$ vanishes, and it may happen that for some period of time $|m_{\rm eff}| < H_I$. We neglect for simplicity in what follows the contributions ξR and βT^2 into $m_{\rm eff}$. During this stage, the behavior of ϕ was governed by its quantum fluctuations which are known to arise during inflation [1–4]. The probability distribution of the fluctuations $P(\phi, t)$ satisfies the equation [3,30]:

$$\frac{\partial P}{\partial t} = \frac{H_I^3}{8\pi^2}\frac{\partial^2 P}{\partial \phi^2} + \frac{1}{3H_I}\frac{\partial}{\partial \phi}\left[P \frac{\partial U}{\partial \phi} \right] . \tag{6}$$

Here U is the potential energy of ϕ and P is normalized as

$$\int d\phi\, P(\phi, t) = 1 . \tag{7}$$

If the potential term is negligible, $\langle \phi^2 \rangle$ rises as $H^3 t$ until it reaches the limiting value of the order of H^4/m^2 (for $U = m^2\phi^2$) or $H^2/\sqrt{\lambda}$ (for $U = \lambda\phi^4$), but if the slope of the potential is negative (as in our case for large ϕ), the stochastic rise of ϕ turns into a classical rolling down. The change from quantum to classical behavior takes place when the variation of ϕ due to quantum fluctuations, $\delta\phi_q$, becomes smaller than $\delta\phi_c$ originating from classical rolling down the potential slope. Quantum fluctuations can be visualized as Brownian motion with steps H in the time interval H^{-1}, while classical evolution is governed by the equation of motion, so that

$\delta\phi_c \approx H^{-1}\partial^2 U/\partial\phi^2$ for the same time interval $\Delta t = H^{-1}$ (see, e.g., Ref. [26]). For the case of $U(\phi) = -\lambda|\phi|^4$ the regime is changed from the quantum into the classical one at

$$\phi_{qc}^2 \approx H^2/\lambda \ , \tag{8}$$

The result is slightly different for $U(\phi) = \lambda|\phi|^4 \ln|\phi|^2/\sigma^2$, but this change is not essential for our estimates.

Thus the following picture for the evolution of ϕ emerges. When $m_{\rm eff}$ becomes smaller than H, quantum fluctuations of ϕ start to develop in accordance with Eq. (6). The probability for ϕ to reach the value ϕ_{qc} becomes non-negligible. After reaching this value, ϕ starts to roll down to another potential minimum at $\phi \approx \sigma$. If $m_{\rm eff}$ was smaller than H for a sufficiently long time, the phase transition to the broken-symmetry phase would be accomplished over all space and the alternating regions with opposite sign of *CP*-odd amplitudes would occupy all of the Universe. If, however, the duration of low $m_{\rm eff}$ was short, the probability of bubble formation would be small and the bubbles of *CP*-odd phase would be spatially separated, filling only a small fraction of space. In contrast to the previous case, such a picture is not excluded by observations even if the size of the bubbles and the distance between them is smaller than the present-day horizon.

The phase transition would take place, roughly speaking, if ϕ exceeds the critical value $\phi_{\rm cr} = O(m/\sqrt{\lambda})$ which corresponds to the maximum of $U(\phi)$ when $\Phi \ll \Phi_1$ and $m_{\rm eff}$ once again becomes large. We assume that $\phi_{\rm cr}$ is larger than ϕ_{qc} or in other words that $m_{\rm eff} > H$ at the end of inflation. When ϕ reaches ϕ_{qc} it evolves classically as $(t-t_0)^{-1}$ with $t_0 \approx \lambda^{-1/4}H^{-1}$.

Thus the phase transition would be quickly accomplished practically everywhere once $\phi > \phi_{qc}$. The linear size of the fluctuation when it reaches the value ϕ_{qc}, which by assumption is considerably larger than $\langle\phi(t)\rangle$, is typically of the order of H^{-1}. Bubbles of smaller size are energetically unfavorable while the probability of a fluctuation with $l > H^{-1}$ and large $\phi > \langle\phi(t)\rangle$ is small.

The fraction of the volume of the Universe occupied by bubbles of broken charge symmetry phase can be evaluated as

$$\epsilon(t) = \int_{t_i}^{t} dt' H \int_{\phi_{qc}}^{\infty} d\phi\, P(\phi,t) \ . \tag{9}$$

The factor H enters since it is the inverse characteristic time scale of the development of the fluctuations. This expression is valid if $\epsilon \ll 1$. Though the distance between the bubbles is exponentially increasing, their size is rising with practically the same rate, and so ϵ rather weakly depends on the expansion. This is precisely true if the bubble walls are at rest in the comoving frame. However, even if they move with the speed of light, the relative increase of a bubble with initial size l is given by the factor $(1+2H/l)$. For bubbles with typical size $l = H^{-1}$, this yields a factor of 3 and the order of magnitude estimate (9) remains valid.

There are a few conditions which should be satisfied by the parameters so that this picture is realized. These conditions are seen in the final expression (20) but it is worthwhile to analyze them before giving the formal solution of the equations. Firstly, the duration of the inflation Δt_1 after the inflaton field reached the value Φ_1 should not be too long so that the size of the bubbles would be much smaller than the present-day horizon: $H_I \Delta t_1 \le 50$. Secondly, the period Δt_2, when Φ is close to Φ_1 so that $|m_{\rm eff}| < H_I$, should be bounded from above and from below. The bound from above follows from the condition that the phase transition takes place only in a small fraction of space. Since $\langle\phi^2\rangle$ is known to rise as $H_I^3 t/4\pi^2$ (see, e.g., Ref. [2]), one should demand $H_I^3 t/4\pi^2 < \phi_{qc}^2 = H_I^2/\lambda$. Hence $H_I \Delta t_2 < 4\pi^2/\lambda$. On the other hand, Δt_2 should be large enough so that there was sufficient time to reach the value ϕ_{qc}. This gives the condition $H_I \Delta t_2 > \phi_{qc}/H_I = 1/\sqrt{\lambda}$. These conditions are satisfied for reasonable values of λ: $10^{-2} < \lambda < 1$.

The condition $|m_{\rm eff}| \le H_I$ means that $\lambda_1(\Phi-\Phi_1)^2 < H_I$ or in other words $\lambda_1 \dot{\Phi}_1^2 \Delta t_2^2 \approx H_I^2$. Correspondingly, the variation of Φ after it passed the value Φ_1 would be approximately

$$\Delta\Phi = \dot{\Phi}_1 \Delta t_1 \approx (H_I/\sqrt{\lambda_1})(\Delta t_1/\Delta t_2) \ .$$

To get a consistent picture (in the framework of the chaotic inflationary scenario), one needs $\Delta\Phi = O(m_{\rm Pl})$. The above condition can be realized if $\lambda_1 < 10^{-10}$ even for H_I as small as $10^{-5} m_{\rm Pl}$ (the latter is necessary for sufficiently small density fluctuations). Such a small coupling of the inflaton field to ϕ seems unnatural but this is of the same order of magnitude as the self-coupling of the inflaton, so that this may reflect the universal weakness of the inflaton interactions. In other inflationary scenarios, $\Delta\Phi$ can be much smaller than $m_{\rm Pl}$ and the restriction on λ_1 may be relaxed.

The potential $U(\phi)$ near the minimum $\phi = 0$ can be approximated by the harmonic term $U \approx m_{\rm eff}^2|\phi^2|$. In this case, Eq. (6) can be solved analytically. The initial distribution of ϕ was determined by the effective mass $m_{\rm eff}^2$ at an early inflationary stage when the amplitude of the inflaton field was large and $m_{\rm eff}^2$ was also large [see Eq. (2)]. In this case, ϕ was initially concentrated near the origin, and so the initial distribution of ϕ can be well described by a δ function $\delta(\phi)$. The solution of Eq. (6) with such an initial condition is known to be Gaussian:

$$P(\phi,t) = C(t)\exp[-\phi^2/2\langle\phi(t)^2\rangle] \ , \tag{10}$$

where

$$\langle\phi^2(t)\rangle = \frac{H_I^3}{4\pi^2}\int_0^t dt' \exp\left[-\frac{2}{3H}\int_{t'}^{t} d\tilde{t}\, m_{\rm eff}^2(\tilde{t})\right] \tag{11}$$

and $C(t)$ can be found from the normalization condition:

$$\int_{-\infty}^{\infty} P\, d\phi = 1 \ .$$

This gives

$$C = (2\pi\langle\phi^2\rangle)^{-1/2} \ . \tag{12}$$

The number density n and the size distribution of bubbles, dn/dl, can be found as follows. The rate of the increase of the relative volume occupied by the broken-symmetry phase at the moment t is

$$\dot{\epsilon}=HW\ , \quad (13)$$

where

$$W=\int_{\phi_{qc}}^{\infty} d\phi\, P(\phi,t)$$

[see Eq. (9)]. It can be expressed through n as

$$\dot{\epsilon}=\int dl\, l^3 \frac{dn}{dt\, dl}\ .$$

As we have argued, the bubble size distribution at the moment of their formation is strongly peaked near H^{-1},

$$\frac{dn}{dl\, dt}=\dot{n}\,\delta(l-H^{-1})\ ,$$

where n satisfies the equation

$$\dot{n}=H_I^4 W(t)-3Hn \quad (14)$$

and W is defined by Eq. (13).

Since the bubble size is completely determined by the epoch of their production t_p,

$$l(t)=H_I^{-1}\exp[H_I(t-t_p)]\ , \quad (15)$$

the size distribution of bubbles at the end of inflation t_e is given by

$$\frac{dn}{dl}=\frac{1}{l^4}W\left[t_e-\frac{\ln H_I l}{H_I}\right]$$

$$=\frac{1}{l^4}\int_{\phi_{qc}}^{\infty} d\phi\, P\left[\phi,t_e-\frac{\ln H_I l}{H_I}\right]\ . \quad (16)$$

The distribution of this shape remains after inflation at least until the bubbles reenter the horizon.

P is given by Eqs. (10)–(12). The probability has the maximum value when the inflaton field Φ is near Φ_1. The effective mass of ϕ in this region can be represented as

$$m_{\rm eff}^2=m_0^2+\mu^4(t-t_1)^2\ , \quad (17)$$

and hence $\langle\phi^2(t)\rangle$ is equal to

$$\langle\phi^2(t)\rangle=\frac{H_I^3}{4\pi^2}\int_0^{\infty} d\eta\, \exp\left[-\frac{2m_0^2}{3H}\eta-\frac{2\mu^4}{9H}\eta[(t-t_1)^2+(t-t_1-\eta)^2+(t-t_1)(t-t_1-\eta)]\right]\ . \quad (18)$$

This integral can be easily evaluated in the particularly interesting case $\mu^2/3m_0H_I \ll 1$. Then only the linear (in η) terms in the exponent in (18) are essential and

$$\langle\phi^2(t)\rangle=\frac{3H_I^4}{8\pi^2[m_0^2+\mu^4(t-t_1)^2]}\ . \quad (19)$$

Substituting all the factors into expression (15), we obtain the following distribution of the bubbles in the post-inflationary stage

$$\frac{dn}{dl}=\frac{1}{\sqrt{2\pi}l^4}\frac{\exp[-\delta-\gamma\ln^2 l/L(z)]}{[\delta+\gamma\ln^2 l/L(z)]^{1/2}}\ , \quad (20)$$

where

$$\delta=\frac{4\pi^2 m_0^2}{3\lambda H_I^2}\ ,$$

$$\gamma=\frac{4\pi^2\mu^4}{3\lambda H_I^4}\ ,$$

$$L(z)=H_I^{-1}z\exp[H_I(t_e-t_1)]\ ,$$

and $z=T_{\rm rh}/T$ is the ratio of the reheating temperature at the end of inflation to the running temperature. The distribution (20) remains valid until the bubble size l is larger than the horizon. After this, dynamical effects become essential which depend upon the density contrast and equation of state. We will consider them in Sec. V.

The bubble mass is given by the expression $M(l)=m_{\rm Pl}^2 l^3/8t^2$, and when the bubble reenters the horizon at the moment t_h, it is equal to $M(l=t_h)=m_{\rm Pl}^2 t_h/8$. The distribution of bubbles in mass can be obtained from Eq. (20) and, up to a nonessential preexponential factor logarithmically depending on mass, is given by

$$\frac{dn}{dM}=\frac{m_{\rm Pl}^6}{M_0^4}\exp\left[-\alpha-\frac{\gamma}{4}\ln^2(M/M_0)\right]\ , \quad (21)$$

where $\alpha=\delta+16/\gamma$ and

$$M_0=\tfrac{1}{4}m_{\rm Pl}^2 H_I^{-1}\exp[2H_I(t_e-t_1)-8/\gamma]\ .$$

In fact, the omitted common factor results in a redefinition of α but since the numerical value of the latter is not known, we will not take account of it. In what follows, we assume that δ, γ, and M_0 are free parameters, and consider the cosmological implications of the possible existence of objects with high baryon (antibaryon) number density possessing the mass distribution (21). The calculations presented above allow us to express δ, γ, and M_0 through "fundamental" parameters H_I, m_0, μ, and λ, and show that it is possible to get M_0 in the range $(1-10^6)M_{\odot}$, $\alpha \gg 1$, and $\gamma/4 \le 1$ with natural values of the fundamental parameters. The result (21) was obtained formally for a fixed value of the baryon asymmetry β, but one can see that the mass distribution remains essentially of the same form for uniformly distributed β. Indeed, in the latter case, the distribution would be modified by a power-law factor and the latter can be reabsorbed into expression (21) by a redefinition of the parameters.

The potential $U(\phi)$ is not harmonic of course for large ϕ, especially near the top of the barrier, and the expressions (20) or (21) should be considered only as approximations. Nevertheless, the exponential dependence of both the scale factor on time and the probability distribution on ϕ^2 makes one believe that the exact expressions do not drastically differ from the approximate ones presented here.

In the more realistic case, $U(\phi)$ deviates from harmonic behavior when ϕ is far from the origin. This would modify the simple results presented here but the qualitative behavior should remain unaffected. In particular, since the probability of quantum fluctuations of ϕ with large amplitude is largest when the potential is the most shallow or, in other words, when $m_{\rm eff}^2$ is smallest, the mass distribution of the bubbles should be peaked at some particular value M_0 similar to that given by the approximate equation (21).

IV. BARYOGENESIS SCENARIOS

As we shall see in what follows, the model considered here is especially interesting if the magnitude of the baryon asymmetry inside the bubbles β_B is much larger than that observed in the Universe, $\beta_0 \approx 3\times 10^{-10}$. It is not easy to get β considerably larger than β_0 in most baryogenesis scenarios (for a review see, e.g., Ref. [31]). In fact, the general trend is $\beta < \beta_0$ and the ingenuity of the model builders is aimed towards getting the largest possible β. The only exception is the model of Ref. [32] where the authors obtained $\beta > 1$ and attempted to diminish its value by various mechanisms. This model proves to be very appropriate here as we see in what follows.

The simplest way to have a larger value of β inside the bubbles is to assume that the nonzero ϕ generates CP nonconservation which is much larger than that outside the bubbles. Though it is easy to construct such models, it is difficult to get large β in any natural version of them. If, for example, the baryon asymmetry is generated in out-of-equilibrium decays of heavy particles, there are several suppression factors which make it hard to get β larger than 10^{-4}. In fact, natural values of β are of the order of or below β_0. The origin of the suppression lies, first, in the entropy dilution which is usually by several orders of magnitude and, second, in the inherent smallness of CP-nonconserving effects in particle reactions. Such effects require an interference of CP-odd and -even amplitudes, so that CP violation manifests itself only in higher orders of perturbation theory. Strictly speaking, it is not formally excluded that the CP-violating amplitudes are of the same order of magnitude as the CP-conserving ones and thus the difference in branching ratios for charge-conjugated channels can be large. This would permit generation of a large $\beta(\beta > 10^{-3})$ inside the bubbles because of possible large CP-violating effects induced by a nonzero ϕ. However, concrete realizations of this mechanism need specific fine-tuning and so do not look particularly appealing. Thus it is worthwhile to consider other possibilities for strong baryogenesis inside the bubbles.

In this connection, the model of Ref. [32] is a promising alternative. The basic features are the following. The baryon asymmetry is first accumulated in the form of the condensate of the scalar superpartner of a colorless and electrically neutral combination of quark and lepton fields $\langle \chi \rangle$. This condensate might be formed during the inflationary stage if baryonic and leptonic charges were not conserved and the potential $U(\chi)$ has flat directions. When inflation was over, the decay of this condensate could give rise to a large, $\beta = O(1)$, baryon asymmetry. It is easy to accommodate this scenario with our model. It can be done by a simple substitution of χ instead of ϕ in all of the previous expressions. Thus the condensate of χ would be formed only in relatively small bubbles in the same way as the condensate of ϕ has previously been formed. This condensate could give a large baryon asymmetry inside the bubbles while outside, where $\langle \chi \rangle = 0$, the baryon asymmetry is small.

For the realization of this scenario, the effective potential $U(\chi)$ during inflation should, as above, have two minima: at $\chi = 0$ and at $\chi = \sigma \neq 0$. If the barrier between them becomes small, e.g., due to the mechanism considered in the previous section, the field χ could penetrate the barrier and move to the other minimum at nonzero χ. This minimum might disappear at a later stage, forcing χ to return back to the origin at $\chi = 0$. It can be done with a very simple behavior of the effective mass

$$m_{\rm eff}^2 = m_0^2 + \lambda_1(\Phi - \Phi_1)^2 \tag{22}$$

[compare with Eq. (2)]. For large Φ, $m_{\rm eff}^2$ is large and $U(\chi)$ has the only minimum at $\chi = 0$. If Φ is near Φ_1, another minimum at $\chi = \sigma$ might exist and be deeper than that at $\chi = 0$. Thus it would be possible for χ to roll down to the second deeper minimum in some part of space by the mechanism described in Sec. III. Inside the bubbles, the amplitude of χ is large, $\chi \geq H_I/\sqrt{\lambda} >> H_I$ even if χ did not evolve down the potential minimum. When Φ passes through the point Φ_1 and becomes small, $m_{\rm eff}^2$ once again became large so that the minimum at $\chi = \sigma$ might vanish. As a result, χ has to return back to the original minimum but now it could have a large baryonic charge which would be revealed after χ decay into fermions.

For a better understanding of the dynamics, the following mechanical analogy is instructive. The baryonic current of χ is given by the expression

$$j_\mu^B = iB(\chi^* \partial_\mu \chi - \partial_\mu \chi^* \chi) , \tag{23}$$

where B is the baryonic charge of χ particles. For spatially homogeneous χ, the baryonic charge density can be visualized as the angular momentum of a pointlike body moving in the two-dimensional potential $U(\mathrm{Re}\chi, \mathrm{Im}\chi)$. This is conserved for a spherically symmetric potential $U(|\chi|)$.

It is essential for the model that j_μ^B is not conserved, otherwise the charge density would die out during inflation and not accumulate as we originally assumed. The necessary asphericity is induced by the quartic terms like, e.g., $\chi^4 + \chi^{*4}$, while near the origin at small $|\chi|$, the potential is B conserving, $U(\chi) \approx m_\chi^2 |\chi|^2$. Those χ's which succeed in passing over the potential barrier might acquire a nonzero angular momentum by the same mechanism as in the original model [32] if the slope of the potential of the phase of χ is not large in comparison with H_I. As a result the phase of χ may be apart from the equilibrium point of the potential. So when inflation ends and the Hubble friction becomes small, a nonzero angular momentum may arise. This angular momentum would remain when χ returns to the origin after the second minimum disappears. The subsequent B-conserving decay of χ might provide a large baryon asymmetry in small spatial regions.

We will not dwell further on the details here because they are essentially the same as in the original version of the scenario [32], the only difference being that the condensate of χ is not formed uniformly in space (see, however, Ref. [19]) but only inside a small fraction of the total volume of the Universe. Since the baryonic charge of χ is created by the quantum fluctuations of the latter, its value and sign can be different in different bubbles. This is in contrast to the model of spontaneous CP violation considered in the beginning of this section when the amplitude of β is the same for all the bubbles and only the sign is different.

Another promising possibility for creation of large β_B is given by the model of spontaneous baryogenesis [33] when a Goldstone-type field is coupled by baryonic current J_μ^B as

$$L_{\rm int} = \partial_\mu \psi J_\mu^B / \Lambda \ , \tag{24}$$

where Λ is a constant with dimension of mass. The baryon asymmetry generated in this model is proportional to the time derivative $\dot\psi$. The latter can be large for those ψ which penetrate the potential barrier, thus permitting a large baryon asymmetry inside the bubbles.

The baryon asymmetry generated inside the bubbles should not be destroyed by later processes with baryonic charge nonconservation (if they took place). This is definitely true in the first scenario with spontaneous $C(CP)$ breaking, since all the features of the baryogenesis inside and outside the bubbles remain the same with the only difference arising in the amplitude of $C(CP)$ breaking. This is not obligatory in the second scenario. However, if this mechanism generates nonzero $(B-L)$ and the later processes conserve it, the asymmetry is evidently not destroyed. Another possibility is that the baryogenesis considered here goes at a low temperature T_1 (which can be due to a very weak inflaton coupling to the matter fields) and no other processes with baryon nonconservation were effective below T_1.

An interesting question is whether our model is compatible with the extensively discussed topic of electroweak baryogenesis (for a review see, e.g., Ref. [31]). Cosmologically interesting bubbles where the conditions for a generation of a large baryon asymmetry were created could only be formed during inflation and should not dissolve at least until baryogenesis. After that, the domain walls (if they existed) separating the bubbles from the rest of the Universe should disappear to avoid the well-known cosmological problems [34]. In the standard $SU(3)\otimes SU(2)\otimes U(1)$ model with minimal particle content, this scenario is hard to realize. It remains questionable if the observed baryon asymmetry can be generated in the standard model [35] or whether an extension of the latter is necessary by an introduction of additional Higgs fields, by supersymmetry or by some other means. It is not clear if our model of inhomogeneous baryogenesis can be realized in these limited frameworks or even in one of the many less restrictive "realistic" particle physics models. However, it is definitely possible to introduce in an *ad hoc* manner the necessary fields and interactions in a consistent way such that the proposed scenario is realized.

V. BUBBLE EVOLUTION

Initially the energy density contrast $\delta\rho$ between the bubbles of higher baryonic charge and the rest of the Universe is very small. During the inflationary stage, the difference of energy densities between the two phases is by assumption much smaller than the energy density of the inflaton. This energy difference completely vanishes after baryogenesis when the phase transition was completed everywhere and the bubble walls disappeared. We assume that the sizes of the bubbles prior to this were much larger than the horizon so that the surface effects were not essential. It is possible, however, that the initial density contrast was significant. In this case, the scenario would be more complicated and, in particular, small size black holes might form. If the initial $\delta\rho$ was not large enough for that, the density contrast should disappear in the course of the Universe evolution since the equation of state inside and outside the bubbles was the same.

The energy density contrast remains zero until the QCD phase transition when practically massless quarks transformed into nonrelativistic nucleons and the equation of state inside the bubbles started to deviate from the relativistic one, $p=\rho/3$. Depending on the value of β, these deviations could be either initially small nonrelativistic corrections or could almost instantly change the relativistic equation of state into the nonrelativistic one, $p\approx 0$. Because of the different equations of state inside and outside the bubbles, density inhomogeneities would develop. Small-density inhomogeneities might also appear before the electroweak phase transition when quarks, especially t quarks, acquire masses. These inhomogeneities soon disappear (as $1/t$) when the equation of state relaxes back to $p=\rho/3$. Although $\delta\rho\approx 0$, there was a temperature difference between the two phases due to the different chemical potentials of the baryons. The latter is negligible outside the bubbles but could be of the same order as the temperature inside the bubbles. The energy density of relativistic plasma with g effective degrees of freedom ($g=g_{\rm bosonic}+\frac{7}{8}g_{\rm fermiionic}$) and baryonic chemical potential $\mu=\xi T$ is equal to

$$\rho=\frac{\pi^2}{30}gT^4\left[1+\frac{15\xi^2}{g\pi^2}+\frac{30\xi^4}{8\pi^4 g}\right] . \tag{25}$$

This gives

$$T_e=T_i\left[1+\frac{15\xi^2}{\pi^2}+\frac{30\xi^4}{8\pi^4 g}\right]^{1/4} , \tag{26}$$

where T_i and T_e are the temperatures inside and outside the bubbles, respectively.

The baryon asymmetry of the Universe is defined as the ratio of the baryonic charge density to the entropy density, and can be expressed through ξ as

$$\beta=\frac{N_B}{S}=\frac{10\xi}{g}\left[1+\frac{\xi}{\pi^2}\right] . \qquad (27)$$

For β of order 1 and $g\approx 10^2$, the temperature inside the bubbles can be several percent below the surrounding temperature. The fate of a bubble depends upon its size and the magnitude of the baryon asymmetry β. The latter is not necessarily a universal constant for all the bubbles but may be distributed stochastically, as argued previously.

Very small bubbles with size smaller than the quark diffusion distance d should disappear by the time of the QCD phase transition. In comoving coordinates, d can be evaluated as $(tl_{\rm free})^{1/2}$ where $l_{\rm free}$ is the quark mean free path, $l_{\rm free}\approx(\sigma N)^{-1}\approx T^{-1}$. This gives

$$d\approx(t/T)^{1/2} . \qquad (28)$$

Bubbles of larger size survive to the QCD phase transition, and their evolution would depend upon the developing density contrast. One more factor might also be of importance; namely, the QCD phase transition should depend on the magnitude of the baryon asymmetry and proceed differently in the regions with different β. In particular, in bubbles with very large β, the phase transition is inhibited since the large energy density of the massive nucleons makes it energetically unfavorable (see, e.g., Ref. [36]). The gain in vacuum energy $\delta\rho_{\rm vac}$ should be larger than the loss due to nonzero nucleon mass $\delta\rho_B=m_Bn_B\approx\beta m_BT^3$, so that the phase transition could proceed only at sufficiently low temperatures

$$T<100\ \text{MeV}\left[\frac{100\ \text{MeV}}{m_B\beta}\right]^{1/3}\frac{\delta\rho_{\rm vac}}{(100\ \text{MeV})^4} . \qquad (29)$$

The large chemical potential of the baryons could also change the nature of the phase transition, making it strongly first order. This might open a way for creation of quark nuggets [37], identified as those pieces of space where the phase transition has never been completed. We assume here, however, that the phase transition has proceeded in the same way everywhere and postpone the above problems for a future investigation. We neglect also the temperature difference (26) which affects the onset of the phase transition.

The evolution of the bubbles with size l_B smaller than the horizon size h at the moment of the QCD phase transition depends upon the relation between l_B and the Jeans wavelength λ_J,

$$\lambda_J=c_s(\pi m_{\rm Pl}^2/\delta\rho_B)^{1/2} , \qquad (30)$$

where c_s is the speed of sound inside the bubble and $\delta\rho_B$ is the energy density contrast developed due to the different equations of state inside and outside the bubbles after the QCD phase transition. Initially $\delta\rho=0$ and the bubble temperature is smaller than the temperature of the surrounding plasma, firstly, due to the above-mentioned effect (26) and, secondly, because the latent heat of the phase transition heats the plasma of heavy particles to a smaller temperature than that of the light ones.

For the case of $l_B<\lambda_J$, one would expect that the density contrast is governed by sound wave propagation which ultimately smoothed it down to zero. This is not necessarily so in our case because the usually positive energy density contrast $\delta\rho>0$ implies simultaneously positive pressure differences $\delta p>0$ while we encounter the opposite situation when $\delta\rho>0$ and $\delta p<0$ because of the smaller plasma temperature inside the bubbles. Hence, although initially l_B was smaller than λ_J, this condition might be reversed in the course of the evolution. This is not, however, the case for adiabatic compression (or expansion) when $c_s=(5T/3m)^{1/2}$ and $\lambda_J\sim(Tl)^{1/2}l$. Since $Tl=$const, the ratio λ/l_B does not change. Thus we should expect that the condition $\lambda_J>l_B$ always remains valid for those small bubbles so that they do not collapse and would form diffuse clouds of baryons or antibaryons. In the latter case, there should be prominent annihilation processes on the boundary which would result either in disappearance of small clouds if there were sufficiently many baryons around or else in emptying the space of ordinary matter around the anticlouds. In the latter case, the clouds would survive to the present day in almost empty space, producing rare events of high-energy γ quanta. Note that these clouds are not necessarily spherically symmetric because the shape of the initial bubble could be arbitrary in contrast to the case when the bubbles are formed in a first-order phase transition.

Bubbles of large size, $l_B>\lambda_J$, would form compact objects, either stars or black holes, at a very early stage of the evolution of the Universe. Stars of antimatter could emit considerable energy due to annihilation of the accreted matter. With a sufficiently large amount of surrounding matter, they should radiate at their Eddington limit,

$$L_{\rm Ed}=3\times10^4L_\odot\left[\frac{M}{M_\odot}\right] . \qquad (31)$$

The lifetime of such objects is of the order of 5×10^9 years. If the accretion rate is below the limiting one (e.g, due to the surrounding deficit of matter), the luminosities would be smaller and the lifetime would be larger. Those objects can be observed as γ-ray sources isotropically distributed over the sky. Infalling matter is likely to form an accretion disk heated by the annihilation radiation. We expect that the inner edge will continuously break away and annihilate in a series of discrete events. The duration of an individual event would be something like a few crossing times at $\sim 10-100$ Schwarzschild radii, or of order a second, and so may be identifiable as γ-ray burst events. If γ-ray bursters are cosmologically distant, as inferred from their isotropy and number-flux distribution [38], annihilation would be the most efficient, and hence preferred, form of energy release. At cosmological distances, the emitted photon energy in a γ-ray burst is $\sim10^{48-50}$ erg, corresponding to annihilation of $\sim10^{-6}$ to $10^{-4}M_\odot$. Particle acceleration is likely to occur in such energetic phenomena, and an interesting signature of such annihilation events would be via observation of rare light antinuclei and in particular of antihelium-4 in cosmic rays.

The bubbles with $l_B > h$ at the QCD phase transition have a strong possibility of forming black holes. The mass inside the horizon in the radiation-dominated Friedmann universe is equal to

$$M_h = m_{\rm Pl}^2 t/8 \approx 4\times 10^4 M_\odot (t/\text{sec}) . \quad (32)$$

If $\delta\rho/\rho$ is of the order of 1 at the moment of reentering the horizon—that is, when $l_B = t$—black holes with mass (32) should form. The energy density contrast between the bubbles and the surrounding matter developed when the nonrelativistic baryons began to dominate the bubble energy density. It took place at the moment

$$t_{\rm MD} \approx 10^{-6}\beta^{-2} \text{ sec} \quad (t > 10^{-4} \text{ sec}) . \quad (33)$$

Subsequently,

$$\delta\rho/\rho \rightarrow \tfrac{7}{9}(1 - 32 t_{\rm MD}/21t)$$

and the bubble size was close to its gravitational radius. Simultaneously the condition $l_B > \lambda_J$ was fulfilled, and the bubbles started to shrink, collapsing into black holes. The minimum mass of the black holes depends upon the value of β,

$$M_{\min} = \max(4M_\odot, 4\times 10^{-2}\beta^{-2} M_\odot) . \quad (34)$$

The number density of these black holes is given by Eq. (21) where α, γ, and M_0 can be chosen so that the observational constraints are satisfied.

VI. COSMOLOGICAL IMPLICATIONS

There is considerable freedom in our model since the essential parameters cannot be fixed theoretically. Hence, we cannot make rigorous theoretical predictions, but indicate how different physical effects arise as parameters are varied.

Varying the parameter γ, we can change the spectrum (21) from being sharply concentrated near M_0 to being broad, with abundant objects of mass several orders of magnitude different from M_0. The parameter α fixes the total number density of the bubbles.

The model considered here opens the interesting possibility that the dark matter in the universe is baryonic. To reconcile the model with existing observations, the bubbles of high-density baryonic or antibaryonic matter should predominantly form black holes. As reference values, we take $M_0 = 10^2 M_\odot$ and $\gamma = 0.4$. In this case, the bulk of the mass in the Universe would be contained in black holes with masses in the interval $(10^2-10^4)M_\odot$. With an appropriate α in Eq. (21), one can get the number density of these black holes to correspond to the amount of dark matter in the Universe. Larger black holes can also be formed. The relative number density of black holes with $M = 10^7 M_\odot$ is about 5×10^{-9} with respect to those with $M = 10^2 M_\odot$. These bigger black holes could serve as rare seeds for galaxy formation with smaller ones forming surrounding halos in the same way as in the more conventional cold dark matter model.

As a by-product, the model explains early quasar formation if the latter are black holes powered by matter accretion. The epoch of black-hole formation is determined by Eq. (33). This is much earlier than is allowed in the standard cold dark matter model when black holes are assumed to be formed by nonlinear growth of inhomogeneities in almost uniformly distributed matter. Note that in the present scenario, the usual objections against baryonic dark matter are avoided. There is no nucleosynthesis bound on the baryon-to-photon ratio because the regions with high β occupy a negligibly small fraction of space which ultimately collapsed into black holes or formed compact starlike objects. Hence, the observed abundances of light elements arise from regions of low β where nucleosynthesis proceeds in the standard way, while in regions with higher β, heavier elements might be produced. This would not be in conflict with observations if these heavier elements are not dispersed on the sky but confined inside the black holes and possibly inside primordial stars or quark nuggets.

Some fraction of high β bubbles which are below their Jeans wavelength would form dispersed clouds of matter or antimatter where heavy elements or antielements could be formed during primordial nucleosynthesis. These clouds as well as primordial massive stars that ejected matter could be the sources of heavy metals (and antimetals) in the early universe. One of the possible scenarios of strong baryogenesis based on a large *CP* violation inside the bubbles predicts a constant value of $|\beta|$ for all the bubbles. In that case, any contradiction with the observed element abundances can be easily avoided, since an arbitrary amount of baryonic matter can be hidden in black holes.

For the case of uniformly distributed β, there should be a considerable quantity of high-density baryonic matter outside black holes. This can result in formation of quark nuggets [37] and primordial stars, while a small amount may be dispersed into clouds. If we assume that the mass distribution of the bubbles is peaked near $(10^2-10^3)M_\odot$, then between 10^{-4} and 10^{-6} of the total mass of the Universe should be in the form of quark nuggets or primordial stars as well as any ejecta from the latter. This would result in the primordial enrichment of the Universe, presumably to an extreme population II level, with heavy elements. The ratios of heavy elements, such as the *r* and *s* process proportions, should differ from those in population II stars with, for example, $[\text{Fe/H}] \gtrsim -2$. Primordial enrichment of cosmic rays would yield spallation products, including Li, Be, and B, that is chemically decoupled from spallation products generated by conventional cosmic rays. A possible signature would arise in the dependence of, for example, Be and B on Fe/H that should be linear at very low Fe/H. One would also expect to observe antinuclei in cosmic rays, although their number density should be reduced in comparison with primordial heavy elements due to annihilation.

We would like to note that much more work is necessary here to refine these suggestions. Firstly, primordial nucleosynthesis with β of order 1 should deviate very much from the standard one. Secondly, the formation of compact starlike objects and their evolution might be very different from the usual case of star formation from cold matter. A detailed investigation of these phenomena could result in interesting bounds on the properties of the model and on the values of the essential parameters. We think, however, that this is outside the scope of the present paper. Our aim here has been to describe a model which could in principle produce the phenomena described above. We may be quite sure that the scenario can be made compatible with the existing observational data since it is possible to conceal high B regions in the form of black holes with arbitrarily good accuracy so that almost, or even absolutely, no trace of them remains, except for their mass. In the limiting case when all bubbles formed black holes at an early stage, there is no way to test the specific predictions of this model, although one is of course strongly constrained by the requirement that the ratio of mass in black holes not exceed that in ordinary matter and in radiation according to the observed parameter range. It is difficult to exclude this model but at the same time hard to prove its validity.

Another argument in favor of nonbaryonic dark matter is based on the standard model of inhomogeneity growth. The latter can reach the value

$$(\delta\rho/\rho)\sim(z_i+1)(\delta\rho/\rho)_i$$

(for $\Omega=1$) starting from initially small $(\delta\rho/\rho)_i$. The bounds on the angular variations of the cosmic microwave background radiation, $\Delta T/T$, restrict $(\delta\rho/\rho)_i$ to approximately 10^{-4} on galaxy cluster scales. Growth of inhomogeneities in baryonic matter could begin only after hydrogen recombination, that is after $z=10^3$. Hence there is not sufficient time for baryonic $(\delta\rho/\rho)$ to reach a value of the order of unity. The usual conclusion from this argument is support for the thesis that the inhomogeneities developed in dark nonbaryonic matter which does not interact with electromagnetic radiation and became unstable with respect to gravitational clustering after the epoch of first matter domination, $z\sim10^4$. It is evident that this argument is not applicable to our model. There is of course large $\Delta T/T$ in a baryon-dominated model, but at angular scales where relic fluctuations have been erased by the finite thickness of the last scattering surface.

Isocurvature fluctuations of large but insufficient amplitude to have become bound and collapsed in the radiation-dominated era, namely, with amplitude of order unity, would survive as bound clouds during and after recombination provided their masses are at least $10^4-10^6M_\odot$ [39]. Annihilations from the rare antimatter regions that survive as discrete regions have observable signatures. These include spectral distortions of the cosmic microwave background, which restrict the amount of energy release in annihilations in the early universe through the chemical potential distortion that characterizes any energy input arising at [40] $10^4\lesssim z\lesssim10^8$, thereby probing isocurvature fluctuations of amplitude as large as $\sim10^4$. The current limit [41] $\mu/kT\lesssim5\times10^{-3}$ translates directly to an equivalent limit on $\Delta\rho_\gamma/\rho_\gamma$ in energy released via early annihilations. In the range $10^4\gtrsim z\gtrsim10^2$, the annihilation radiation is absorbed, and would result in a Comptonization distortion of the cosmic microwave background spectrum, the interpretation of which depends on the detailed thermal history of the absorbing medium. Any positive evidence for either type of distortion that is near the current limits would be a reasonably natural outcome of a model such as that advocated here with energy injection at very high redshift. At $z\lesssim10^2$ the diffuse γ-ray background directly restricts the antimatter annihilation rate: the ratio of the energy density in the γ-ray background to that in the cosmic microwave background is $\sim10^{-5}$. It has indeed been suggested [42] (see also Ref. [43]) that redshifted annihilation radiation may be responsible for a feature near 1 MeV in the isotropic γ-ray background spectrum.

Constraints on the amount of dark matter in black holes in the Galactic disk were presented in Ref. [44]. The arguments of this paper based on the observed metal abundance are not applicable here since in our case the formation mechanism of black holes is different from the traditional one by stellar evolution and so does not lead to an enhanced metallicity. X-ray emission generated by the gas accreted by black holes gives a rather strong bound on massive black holes in the Galactic disk. Any substantial component of dark mass in the disk probably does not consist of massive black holes. These arguments do not exclude black holes more or less uniformly distributed in the (extended) halo.

At the present time, the most promising way to detect the type of dark matter proposed here is via the ongoing searches [45] for gravitational microlensing in our galactic halo by monitoring several million Large Magellanic Cloud stars: these experiments are sensitive to compact halo objects with masses up to $\sim100M_\odot$.

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[1] A. H. Guth and S. Y. Pi, Phys. Rev. Lett. **42**, 1110 (1982).
[2] A. D. Linde, Phys. Lett. **116B**, 335 (1982).
[3] A. A. Starobinsky, Phys. Lett. **117B**, 175 (1982).
[4] For a review see S. K. Blau and A. H. Guth, in *300 Years of Gravitation,* edited by S. W. Hawking and W. Israel (Cambridge University Press, Cambridge, England, 1987).
[5] For a review see J. Silk, in *Current Topics in Astro-Fundamental Physics,* edited by N. Sanchez and A. Zichichi (World Scientific, Singapore, 1992), p. 183.
[6] J. A. Peacock, Mon. Not. R. Astron. Soc. **253**, 1P (1991).
[7] J. D. Barrow and M. S. Turner, Nature **291**, 469 (1981).
[8] J. R. Bond, E. W. Kolb, and J. Silk, Astrophys. J. **255**, 341 (1982).
[9] M. Yoshimura, Phys. Rev. Lett. **51**, 439 (1983).
[10] D. Seckel and M. S. Turner, Phys. Rev. D **32**, 178 (1985).
[11] A. D. Linde, Phys. Lett. **158B**, 375 (1985).
[12] L. A. Kofman, Phys. Lett. B **173**, 400 (1986).
[13] M. Fukugita and V. A. Rubakov, Phys. Rev. Lett. **56**, 988 (1986).
[14] A. D. Dolgov and N. S. Kardashov, Space Research Int.

Report No. 1190, 1986 (unpublished).
[15] A. D. Dolgov, A. F. Illarioniv, N. S. Kardashov, and I. D. Novikov, Zh. Eksp. Teor. Fiz. **94**, 1 (1987) [Sov. Phys. JETP **67**, 213 (1988)].
[16] L. A. Kofman and A. D. Linde, Nucl. Phys. **B282**, 555 (1987).
[17] L. A. Kofman and D. Yu. Pogosyan, Phys. Lett. B **214**, 508 (1988).
[18] M. S. Turner, A. G. Cohen, and D. B. Kaplan, Phys. Lett. B **216**, 20 (1989).
[19] A. D. Dolgov and D. P. Kirilova, J. Moscow Phys. Soc. **1**, 217 (1991).
[20] M. V. Chizhov and A. D. Dolgov, Nucl. Phys. **B327**, 527 (1992).
[21] J. Yokoyama and Y. Suto, Astrophys. J. **379**, 427 (1991).
[22] M. Sasaki and J. Yokoyama, Phys. Rev. D **44**, 970 (1991).
[23] G. Steigman, Annu. Rev. Astron. Astrophys. **14**, 336 (1976).
[24] T. D. Lee, Phys. Rev. D **8**, 1226 (1973).
[25] S. Coleman and E. Weinberg, Phys. Rev. D **7**, 1888 (1973).
[26] A. D. Linde, Phys. Lett. **129B**, 177 (1982).
[27] M. B. Voloshin and A. D. Dolgov, Yad. Fiz. **35**, 213 (1982) [Sov. J. Nucl. Phys. **35**, 373 (1982)].
[28] R. N. Mohapatra and G. Senjanovic, Phys. Rev. D **20**, 3390 (1979).
[29] A. D. Dolgov, in *Einshteinovski sbornik 1980-1981,* edited by I. Yu. Kobzarev (Nauka, Moscow, 1985), p. 111.
[30] A. A. Starobinsky, in *Field Theory, Quantum Gravity and Strings,* Proceedings of the Seminar Series, Meudon and Paris, France, 1984–1985, edited by H. J. de Vega and N. Sanchez, Lecture Notes in Physics Vol. 246 (Springer-Verlag, Berlin, 1986), p. 107.
[31] A. D. Dolgov, Phys. Rep. **222**, 309 (1992).
[32] I. Affleck and M. Dine, Nucl. Phys. **B249**, 361 (1985).
[33] A. Cohen and D. Kaplan, Phys. Lett. B **199**, 251 (1987); Nucl. Phys. **B308**, 913 (1988).
[34] Ya. B. Zel'dovich, I. Yu. Kobzarev, and L. B. Okun, Zh. Eksp. Teor. Fiz. **67**, 3 (1974) [Sov. Phys. JETP **40**, 1 (1975)].
[35] M. E. Shaposhnikov, Phys. Lett. B **277**, 324 (1992); **282**, 483(E) (1992).
[36] A. D. Linde, Phys. Rev. D **14**, 3345 (1976).
[37] E. Witten, Phys. Rev. D **30**, 272 (1984).
[38] S. Mao and B. Paczynski, Astrophys. J. **388**, L45 (1992).
[39] B. J. Carr and M. J. Rees, Mon. Not. R. Astron. Soc. **206**, 315 (1984).
[40] R. A. Sunyaev and Ya. B. Zel'dovich, Astrophys. Space Sci. **7**, 20 (1970).
[41] G. Smoot, in *Current Topics in Astrofundamental Physics* [5], p. 192.
[42] F. W. Stecker, D. L. Morgan, and J. Bredekamp, Phys. Rev. Lett. **27**, 1469 (1971); F. W. Stecker, in *High Resolution Gamma Ray Cosmology,* Proceedings of the Conference, Los Angeles, California, 1988, edited by D. B. Cline and E. J. Fenyves [Nucl. Phys. B (Proc. Suppl.) **10**, 93 (1989)].
[43] Y.-T. Gao and D. Cline, Mod. Phys. Lett. A **6**, 2669 (1991).
[44] D. J. Hegyi, E. W. Kolb, and K. A. Olive, Astrophys. J. **300**, 492 (1986).
[45] K. Griest *et al.*, Astrophys. J. **372**, L79 (1991).