# Baryogenesis in extended inflation. I. Baryogenesis via production and decay of supermassive bosons

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(Received 21 May 1990)

We consider baryogenesis occurring during the thermalization stage at the end of extended inflation. In extended inflation, the Universe passes through a first-order phase transition via bubble nucleation; inflation comes to an end when bubbles collide and their collisions convert energy stored in the bubble walls into particles. This naturally provides conditions well out of thermal equilibrium in which baryon-number-violating processes may proceed; we estimate the amount of baryon asymmetry which may be produced this way. The avoidance of a monopole or domain-wall problem can also be ensured and isothermal density perturbations may arise as a remnant of spatial variation in the baryon asymmetry.

## I. INTRODUCTION

Recently the spirit of the original inflationary cosmology[1] has been revived in the context of "extended" inflationary models by La and Steinhardt.[2] In such models, the Universe is trapped in a false-vacuum state as it cools from high temperatures; the energy density of this false vacuum then drives a rapid expansion in the scale factor of the Universe, solving a variety of cosmological conundrums. Inflation is ended by the quantum-mechanical process of formation of bubbles of the true vacuum via tunneling; bubbles form with a characteristic size determined by microphysics[3] (provided gravitational corrections are small). These bubbles then expand at the speed of light, eventually colliding with adjacent bubbles. The percolation of these bubbles then brings the inflationary era to an end. The original "old inflation" scenario of Guth[1] was flawed by what became known as the "graceful exit" problem: regions trapped in the false-vacuum state expand exponentially; the expansion generically overcomes the decay to the true-vacuum state and percolation of the Universe by true-vacuum bubbles never occurs.[4] Extended inflation circumvents this obstacle by considering modified gravitational theories (such as the Jordan-Brans-Dicke theory) in which the gravitational constant may vary. In such theories the inflationary expansion is a rapid power law rather than exponential, and the exponential bubble nucleation rate will always eventually overcome the expansion and bring the inflationary era to a satisfactory end.

As pointed out by Weinberg[5] and by La, Steinhardt, and Bertschinger,[6] the original extended inflation model based on a Jordan-Brans-Dicke theory fails because bubbles nucleated early in inflation have time to grow to large sizes. Such bubbles do not have time to thermalize before radiation decoupling (a lower bound on the thermalization time being easily obtained simply from causality) and would cause unacceptably large distortions in the microwave background. To resolve this conflict, several more involved models have been proposed,[7–10] with the common theme of arranging that the production of bubbles early in inflation is suppressed. This appears to be a necessary ingredient for a successful extended inflation model, and here we shall assume, without tying ourselves down to a particular model, that the vast majority of bubbles are produced in a rapid burst right at the end of inflation. These bubbles have little time to grow before the inflationary era is brought to an end by percolation. A detailed examination of the dynamics of extended inflationary models is given in Ref. 11. We note also that it is simply the falling Hubble expansion rate that enables the phase transition to proceed to completion in extended inflationary models, and similar conclusions could be drawn in any power-law inflationary model[12] in which a first-order transition occurs. Thus the picture we shall present is more general than the "extended" inflationary universe model which we use to provide a context.

In this paper we will not address problems in the dynamics of the bubble nucleation rate, and only assume that some satisfactory explanation will result in an acceptable bubble distribution at the end of extended inflation. Rather, we will concentrate on the inflaton sector of the theory, and investigate whether an acceptable baryon asymmetry can be produced after extended inflation.

One of the most important results in particle astrophysics is the development of a framework that provides a dynamical mechanism for the generation of the baryon asymmetry. Before reviewing the basic ingredients necessary, it is useful to quantify exactly what is meant by the baryon asymmetry. The baryon number density is defined as the number density of baryons, minus the number density of antibaryons: $n_B \equiv n_b - n_{\bar{b}}$. Today, $n_B = n_b = 1.13\times 10^{-5}(\Omega_B h^2)$ cm$^{-3}$, where $h$ is Hubble's constant in units of 100 km s$^{-1}$ Mpc.$^{-1}$ Of course, the baryon-number density changes with expansion, so it is most useful to define a quantity $B$, called the *baryon number of the Universe,* which is the ratio of the baryon-number density to the entropy density $s$. Assuming three species of light neutrinos, the present entropy density is $s = 2970$ cm$^{-3}$, and the baryon number is

$$B = 3.81\times 10^{-9}(\Omega_B h^2) \ . \tag{1.1}$$

Primordial nucleosynthesis provides the constraint $0.010 \le \Omega_B h^2 \le 0.017$,[13] which implies $B = (3.81-6.48)\times 10^{-11}$. So long as baryon-number-violating processes are slow compared to the expansion rate and no entropy is created in the expansion, $B$ is constant.

A key feature of inflation is the creation of a large amount of entropy in a volume that was at one point in causal contact. The creation of entropy in inflation would dilute any pre-existing baryon asymmetry, so it is necessary to create the asymmetry after, or very near the end of, inflation. In order for the baryon number to arise after inflation in the usual picture, where *CPT* invariance and unitarity hold, it is necessary for three criteria to be satisfied: baryon-number- ($B$-) violating reactions must occur, $C$ and $CP$ invariance must be broken, and nonequilibrium conditions must be obtained. There are two standard scenarios for baryogenesis:[14] In the first picture the baryon asymmetry is produced by the "out of equilibrium" $B$-, $C$-, and $CP$-violating decays of some massive particle, while the second scenario involves the evaporation of black holes.[15] We shall discuss the role of the latter mechanism in a second paper on this subject.

In the out of equilibrium decay scenario, the most likely candidate for the decaying particle is a massive boson that arises in grand unified theories (GUT's). In the simplest models, the degree of $C$ and $CP$ violation is larger for Higgs scalars than for the gauge vector bosons, so we will assume that the relevant boson is a massive Higgs particle. This Higgs particle is also taken to be different from the inflaton. The Higgs particle of GUT's naturally violate $B$. The origin of the $C$ and $CP$ violation necessary for baryogenesis is uncertain. It is practical simply to parametrize the degree of $C$ and $CP$ violation in the decay of the particle. To illustrate such a parametrization, imagine that some Higgs scalar $H$ has two possible decay channels: to final states $f_1$, with baryon number $B_1$, and $f_2$, with baryon number $B_2$. Consider the initial condition of an equal number of $H$ and its antiparticle $\bar{H}$. The $H$'s decay to final states $f_1$ and $f_2$ with decay widths $\Gamma(H\to f_1)$ and $\Gamma(H\to f_2)$, while the $\bar{H}$'s decay to final states $\bar{f}_1$ and $\bar{f}_2$ with decay widths $\Gamma(\bar{H}\to\bar{f}_1)$ and $\Gamma(\bar{H}\to\bar{f}_2)$. The decays produce a net baryon asymmetry per $H$-$\bar{H}$ given by

$$\epsilon \equiv \sum_{i=1,2} B_i \frac{\Gamma(H\to f_i) - \Gamma(\bar{H}\to\bar{f}_i)}{\Gamma_H} \ , \tag{1.2}$$

where $\Gamma_H$ is the total decay width. Of course $\epsilon$ can be calculated if one knows the masses and couplings of the relevant particles. Reasonable upper bounds for $\epsilon$ are in the range of $10^{-2}$ to $10^{-3}$, but it could be much smaller. For more details, the reader is referred to Ref. 16.

The nonequilibrium condition is most easily realized if the particle interacts weakly enough so that by the time it decays when the age of the Universe is equal to its lifetime, the particle is nonrelativistic. Then the decay products will be rapidly thermalized, and the "back reactions" that would destroy the baryon asymmetry produced in the decay will be suppressed.

In most successful models of new inflation the reheat temperature is constrained to be rather low. This is due to the fact that new inflation requires flat scalar potentials in order for inflation to occur during the "slow roll" of the scalar field toward its minimum. In order to maintain the flatness of the potential, the inflaton field must be very weakly coupled to all fields so that one-loop corrections to the scalar potential do not interfere with the desired flatness of the potential. The feeble coupling of the inflaton to other fields means that the process of converting the energy stored in the scalar field to radiation ("re" heating) is inherently inefficient. Although it is possible to overcome this difficulty in several ways, it remains a concern for new inflation.

The thermalization process of bubble-wall collision at the end of extended inflation provides a natural arena for baryogenesis in the early Universe, as it automatically creates conditions far from thermal equilibrium, exactly as required for $B$-, $C$-, and $CP$-violating GUT processes to produce an asymmetry. The aim of this paper is to investigate how the baryon asymmetry produced at the end of extended inflation can be estimated.

In this, the first paper of two, we consider the production in bubble-wall collisions of supermassive baryon-number-violating bosons whose decays generate the baryon asymmetry. In the second paper we consider the further possibility that the bubble-wall collisions may produce a significant density of black holes, which then decay via the emission of Hawking radiation. These decays may lead to the radiation of more baryons than antibaryons, providing an alternative mechanism for the generation of the baryon asymmetry.

In the next section we will describe the Universe at the end of extended inflation. In particular we will derive the physical parameters that describe the true-vacuum bubbles. In Sec. III we will discuss baryogenesis from the decay of Higgs particles produced in the bubble-wall collisions. The final section discusses our results.

## II. THE END OF EXTENDED INFLATION

Most of the work on extended inflation has concerned the gravitation sector of the theory, which will not concern us. Our only assumption about the gravitational sector is that the parameter that determines the efficiency of bubble nucleation, $\epsilon_N(t)=\Gamma_N(t)/H^4(t)$, where $\Gamma_N$ is the nucleation rate per volume and $H$ is the expansion rate of the Universe, has a time dependence that suppresses bubble nucleation early in inflation, then rapidly increases so inflation is brought to a successful conclusion in a burst of bubble nucleation. This will be the case so long as $H(t)$ falls as $t$ increases. Hence we see that it could occur in any power-law inflationary universe[12] driven by an appropriate phase transition. In fact, our experience with "new" and "chaotic" inflation, as well as inflation driven by higher-order curvature terms in the gravitational Lagrangian,[17,18,19] indicates that we can have "intermediate" and hyperinflation where the scale factor of the Universe increases as $\exp(At^n)$, with $A$ constant, and $n<1$ or $n>1$ respectively. When $n>1$, as is possible in some quadratic Lagrangian inflationary scenarios,[20] we will have $\dot{H}>0$ and a phase transition could not complete even if the effective potential allowed one to occur. However, when $0<n<1$, as considered in Ref. 18, the phase transition could proceed to completion just as in the power-law and extended inflationary models.

Here we shall refer to the extended inflationary model for the definiteness, and we shall be concerned with the inflaton sector of the theory. So far the only restriction on the inflaton sector has been that it must result in a first-order phase transition. Here, we examine the results of requiring that it must also produce a baryon asymmetry.

In order to keep our discussion as general as possible, we will not specify any particular inflaton model, but rather describe the salient features of the potential in terms of a few parameters that can be easily identified with any scalar potential that undergoes spontaneous symmetry breaking. We denote the inflaton field throughout to be $\sigma$, which has a potential of the general form suitable to provide a first-order phase transition necessary for extended inflation. The parameters of the potential are assumed to be the following.

(1) $\sigma_0$, the energy scale for spontaneous symmetry breaking (SSB), i.e., the vacuum expectation value (VEV) of the scalar field.

(2) $\lambda$, a dimensionless coupling constant of the inflaton potential. We will assume that the potential is proportional to $\lambda$.

(3) $\xi$, a dimensionless number that measures the difference between the false- and the true-vacuum energy density via $\rho_V=\xi\lambda\sigma_0^4$. $\xi$ must be less than unity for sufficient inflation to occur; this is also precisely the condition that allows the thin-wall approximation (discussed below) to be made.

From these few parameters we can find all the information we require about the bubbles formed in the phase transition.[3] For instance, an important parameter is the size of bubbles nucleated in the tunneling to the true vacuum. In the thin-wall approximation, the size of a nucleated bubble is given by[14]

$$R_C \sim 3(\xi\lambda^{1/2}\sigma_0)^{-1} . \tag{2.1}$$

Bubbles smaller than this critical size will not grow, and it is exponentially unlikely to nucleate bubbles larger than this critical size. We will assume that all the true-vacuum bubbles are initially created with size $R=R_C$.

Another interesting parameter is the thickness of the bubble wall separating the true-vacuum region inside from the false-vacuum region outside the bubble. For the potential described above, the bubble wall thickness is

$$\Delta \sim (\lambda^{1/2}\sigma_0)^{-1} . \tag{2.2}$$

Note that the ratio of the bubble thickness to its size is $\Delta/R_C \sim \xi$; as advertised, if $\xi \ll 1$, the thin-wall approximation is valid. We note here that our considerations for the rest of this paper are probably valid even in the absence of the thin-wall approximation.

Finally, the energy per unit area of the bubble wall is

$$\eta \sim \lambda^{1/2}\sigma_0^3 . \tag{2.3}$$

At the end of extended inflation all of the energy is in these bubble walls.

We must have some idea of the size of bubbles at the end of inflation, when bubbles of true vacuum percolate, collide, and release the energy density tied up in the bubble walls, so creating the entropy of the Universe. The bubbles of true vacuum are nucleated with size $R=R_C$. After nucleation the bubble will grow until it collides with other bubbles. We now show that the size of bubble at the end of extended inflation is still approximately $R_C$.

Consider first the growth of the bubble in co-moving coordinates. If a bubble is nucleated with *coordinate* radius $r$ at time $t_{\rm nuc}$, then at some later time $t$ the coordinate radius of the bubble will have grown by an amount $\Delta r(t,t_{\rm nuc})$, given by

$$\Delta r(t,t_{\rm nuc}) = \int_{\rm nuc}^{t} \frac{dt'}{a(t')} , \tag{2.4}$$

where $a(t)$ is the Robertson-Walker scale factor. Typically in extended inflation $a(t)$ grows as a power-law in time, say $a(t)\propto t^p$, $p\gg 1$. If this is true, then $\Delta r(t,t_{\rm nuc}) \sim t_{\rm nuc}^{1-p} - t^{1-p}$, which approaches an asymptotic value $\Delta r(\infty,t_{\rm nuc}) \sim t_{\rm nuc}^{1-p}$. Clearly bubbles nucleated at late time (large $t_{\rm nuc}$) will have little growth in coordinate radius, and any increase in the physical size of such a bubble is due solely to the growth in the scale factor between the time the bubble is nucleated and the end of inflation.

The physical size of a bubble nucleated at time $t_{\rm nuc}$ is related to its coordinate size by $R(t_{\rm nuc}) = r(t_{\rm nuc})a(t_{\rm nuc}) = R_C$. If there is negligible growth in the coordinate size of the bubble between the $t_{\rm nuc}$ and end of inflation $t_{\rm end}$, then at the end of inflation the bubble will have a physical size

$$R(t_{\rm end}) \equiv R = r(t_{\rm nuc})a(t_{\rm end}) = R_c[a(t_{\rm end})/a(t_{\rm nuc})] . \tag{2.5}$$

We will assume that the burst of bubble nucleation at the end of inflation leads to bubbles all of the same size, $R=\alpha R_C$, where $\alpha \equiv a(t_{\rm end})/a(t_{\rm nuc})$.

We conclude this section by a description of the Universe at the end of extended inflation. To a good approximation the Universe is percolated by bubbles of true vacuum of size $R=\alpha R_C$, with all the energy density residing in the bubble walls. We have spoken of the "end of inflation" as if it was a well-defined time, but in fact it is not. We simply define the end of inflation to be when the vast majority, say 99%, of the volume of the Universe is no longer in the false-vacuum phase. Our next step is to examine how the release of energy from the bubble walls into radiation via bubble-wall collisions takes place.

## III. BARYOGENESIS BY DIRECT PRODUCTION OF SUPERMASSIVE BOSONS

Let us concentrate on a single bubble of radius $R=\alpha R_C$. The collisions of the bubble walls produce some spectrum of particles, which are subsequently thermalized. We need to estimate the typical energy of a particle produced in these collisions. When a bubble forms, the energy of the false vacuum has been entirely transformed into potential energy in the bubble walls, but as the bubbles expand, more and more of their energy becomes kinetic and the walls become highly relativistic. A simple calculation shows that if the bubble has expanded by a factor of $\alpha$ since nucleation, as discussed in the previous section, then only $1/\alpha$ of its energy remains as potential energy. The numerical simulations of bubble collisions by Hawking, Moss, and Stewart[21] demonstrate that during collisions the walls oscillate through each other, and it seems reasonable that the kinetic energy is dispersed at an energy related to the frequency of these oscillations (see their discussion of phase waves). The kinetic energy is presumably dispersed into lower-energy particles, and does not participate in baryogenesis. We are more interested in the fate of the potential energy. The bubble walls can be imagined as a coherent state of inflaton particles, so that the typical energy of the products of their decays is simply the mass of the inflaton. This energy scale is just equal to the inverse thickness of the wall. Note that by the time the walls actually disperse, most of the kinetic energy has been radiated away,[21] so the walls are probably no longer highly relativistic.

The probable first step in the reheating process is converting this coherent state of the Higgs boson into an incoherent state. The next step would be the conversion of the incoherent state of the Higgs boson into other particles either through decay of the Higgs boson, or through inelastic scattering. We are assuming that baryon-number-violating bosons $H$ will be produced in the process. The $\sigma$ field is typically in the adjoint representation of the gauge group, while $H$ is typically in the fundamental or some other representation. It is possible to envision some symmetry forbidding a direct $\sigma-H$ coupling, or that the coupling is very small compared to other couplings. If this is the case, production of $H$ relative to other particles will be suppressed by some power of the small coupling constant. However in the generic case where all couplings are of the same magnitude there will be no suppression. Of course the ultimate answer is model dependent but calculable.

As discussed earlier, we are assuming that the bubbles do not grow substantially before percolation in our idealized extended inflation model. Hence $\alpha$ remains not too far from 1, though as we shall see a growth by a factor of 1000 even will not necessarily rule out our model. The bubble-wall collisions yield a significant amount of the original false-vacuum energy in the form of potential energy, giving rise to high-energy particles. The potential energy in the bubble walls is given by

$$M_{\text{pot}}=4\pi\eta R^2\sim 4\pi\lambda^{1/2}\sigma_0^3 R^2 \ . \tag{3.1}$$

Taking the mean energy of a particle produced in the collisions to be of the order of the inverse thickness of the wall, $\langle E\rangle\sim\Delta^{-1}$, the mean number of particles produced in the collisions from the wall's potential energy is

$$\langle N\rangle\simeq M_{\text{pot}}/\langle E\rangle\sim 4\pi\Delta\lambda^{1/2}\sigma_0^3 R^2 \ . \tag{3.2}$$

In general, the bubble collisions will produce all species of particles, at least all species with masses not too large compared to $\langle E\rangle$. In the following we will assume that this is the case for the baryon-number-violating Higgs particles. If the Higgs-boson mass exceeds $\Delta^{-1}$ by a significant amount, we can expect some suppression, presumably exponential, in the number of Higgs bosons formed. This possibility will be discussed at the end of this section. For now, we simply parametrize the fraction of the primary annihilation products that are supermassive Higgs boson by a fraction $f_H$, which in general will depend on the masses and couplings of a particular theory in question. The typical number of Higgs particles produced per bubble is

$$\langle N_H\rangle\sim f_H\langle N\rangle\sim 4\pi f_H\Delta\lambda^{1/2}\sigma_0^3 R^2 \ . \tag{3.3}$$

We will now assume that the only source of the supermassive Higgs boson is from the primary particles produced in the bubble-wall collisions. This will be true if the reheat temperature $T_{\text{RH}}$ is below the Higgs-boson mass. (Note that throughout this paper we have set the Boltzmann constant equal to 1.) The validity of this approximation will also be discussed at the end of this section.

The Higgs particles produced in the wall collisions decay, producing a net baryon asymmetry $\epsilon$ per decay, where $\epsilon$ is given in Eq. (1.2). Hence, the excess of baryons over antibaryons produced from a single bubble, $N_B=N_b-N_{\bar{b}}$. is given by

$$N_B=\epsilon\langle N_H\rangle\sim 4\pi\epsilon f_H\sigma_0^2 R^2 \ , \tag{3.4}$$

where we have substituted in for the bubble thickness from Eq. (2.2). This results in a baryon-number density of

$$n_B=N_B/(4\pi R^3/3)=3\epsilon f_H\sigma_0^2 R^{-1} \ . \tag{3.5}$$

We must now calculate the entropy generated in bubble-wall collisions. As stated above, the potential energy of a bubble is $M_{\text{pot}}=4\pi\sigma_0^3\lambda^{1/2}R^2$. Including the (possibly dominant) kinetic-energy contribution, the total mass of the bubble is $M=4\pi\sigma_0^3\lambda^{1/2}R^2\alpha$. Thermalization of the mass in the bubble walls will redistribute this energy throughout the bubble, resulting in a radiation energy density

$$\rho_R \sim M/(4\pi R^3/3) \sim 3\lambda^{1/2}\sigma_0^3\alpha/R = \xi\lambda\sigma_0^4 , \quad (3.6)$$

which is just the false-vacuum energy. The reheat temperature is related to the radiation energy density via

$$\rho_A = \frac{g_* \pi^2}{30} T_{\rm RH}^4 , \quad (3.7)$$

where $g_*$ is the effective number of degrees of freedom in all the species of particles which may be formed in the thermalization process. From this we obtain the entropy density $s$ produced by the thermalization of the debris from bubble-wall collisions:

$$s = \frac{2\pi^2}{45} g_* T_{\rm RH}^3 \sim g_s^{1/4}\xi^{3/4}\lambda^{3/4}\sigma_0^3 . \quad (3.8)$$

From Eqs. (3.5) and (3.8) we can calculate the baryon asymmetry $B$ as

$$B \equiv \frac{n_B}{s} = \epsilon f_H \alpha^{-1} g_*^{-1/4} \lambda^{-1/4} \xi^{1/4} . \quad (3.9)$$

Provided the mass of the Higgs boson is less than $T_{\rm RH}$, one might conjecture that $f_H$ is given simply by $g_H/g_*$, where $g_H$ is the number of Higgs degrees of freedom; that is, all suitably light particles are produced equally. In general the situation will be more complex, and the fraction of Higgs bosons produced will depend on the various couplings in the theory. This introduces a model dependence into the picture, though in fact one can always regard $\epsilon f_H$ as a single unknown parameter. For simplicity, we assume here that all particles are indeed produced equally. Substituting this in gives the final result

$$B = \epsilon g_H \alpha^{-1} g_*^{-5/4} \lambda^{-1/4} \xi^{1/4} . \quad (3.10)$$

This allows us to make numerical estimates of $B$ based on sample values of these parameters. Notice that the dependence of both $\lambda$ and $\xi$, which are the two parameters on which the inflaton's potential depends, is very weak. The important contributions are the degree of asymmetry in $CP$-violating Higgs-boson decays, the number of particle species available for production in the wall collisions and the factor $\alpha$ by which bubbles expand before colliding. Numerical estimates for $B$ based upon this expression will be made in the concluding section.

We now elaborate upon the implications of the two assumptions of our scenario. The first is that the mass of the Higgs boson is not much larger than the typical energy of particles produced in bubble-wall collisions, i.e., $m_H \lesssim \Delta^{-1} = \lambda^{1/2}\sigma_0$. If we take GUT theories as a guide, the Higgs-boson mass is of order $\lambda_H^{1/2}\sigma_0$, where $\lambda_H$ is the coupling constant of the quartic term in the Higgs potential coupling $\sigma$ and $H$. Clearly $\lambda_H^{1/2}$ must not be too much larger than $\lambda^{1/2}$, or there will be a large suppression in $f_H$.

The second assumption is that the reheat temperature is less than the mass of the Higgs boson, so that thermal production of $H$ is not important. This implies that $m_H > \lambda^{1/4} g_*^{-1/4} \xi^{1/4} \sigma_0$. Again assuming that $m_H = \lambda_H^{1/2}\sigma_0$, the requirement becomes $\lambda_H > \lambda^{1/2} g_*^{-1/2} \xi^{1/2}$. If this inequality is not satisfied, then $H$'s will be copiously produced in the thermalization process and baryogenesis will follow the standard out of equilibrium decay scenario rather than the mechanism we have outlined above.

The compatibility and naturalness of these two requirements will be discussed in the concluding section.

The factor of $\alpha^{-1}$ is also easy to understand. If the bubble has expanded by a factor of $\alpha$, only a fraction $\alpha^{-1}$ of the wall energy is "potential;" the rest is in the form of the kinetic energy of the wall. We have been conservative in assuming that only the potential energy of the wall leads to $H$ production.

## IV. DISCUSSION AND CONCLUSIONS

Here we examine some typical numbers for the baryon asymmetry which may be obtained from Eq. (3.10), in the light of the experimental limits discussed in the introductory section setting $B$ at around $10^{-10}$. The number of Higgs degrees of freedom $g_H$ is expected to be of order 1, with simple one degree of freedom for each polarization in the case of a single Higgs boson and further degrees in the case of a doublet or more of Higgs particles. The total number of degrees of freedom $g_*$ is expected to be of order 100–800 in a grand unified theory.[22] This implies, from Eq. (3.10)

$$B \sim 10^{-2}\epsilon \left[\frac{\xi}{\lambda}\right]^{1/4} \alpha^{-1} . \quad (4.1)$$

The remaining microphysical parameters $\epsilon$, $\lambda$, and $\xi$ are less certain, with some dependence on the particular unified theory under examination, though it is reassuring that both $\lambda$ and $\xi$ also enter only to the quarter power and hence the dependence on these quantities is weak. This does however have the further implication that $\epsilon$ should be very small, as we shall shortly see. That a suitable baryon asymmetry can be produced with such a small $\epsilon$ indicates that the bubble wall collisions are very efficient in producing a baryon asymmetry. In particular, this implies that the model can still comfortably work even if $\alpha$ is sizable, as there is plenty of scope for $\epsilon$ to be made larger. A reasonable estimate of $\xi$ is that it may be of order $10^{-2}$ (recalling $\xi \ll 1$ is the condition both for sufficient inflation and for the thin-wall approximation to be valid). The parameter $\lambda$ should probably be less than of order 1, though nothing in principle prevents it from being much smaller; note that a smaller $\lambda$ increases the baryon asymmetry as it leads to a less efficient production of entropy, though $\lambda$ must also be sufficiently large that Higgs particles can be produced in the wall collisions. From these arguments, it seems likely that the ratio $\xi/\lambda$ will be within a few orders of magnitude of unity, implying that if this mechanism is to generate the appropriate baryon asymmetry $\epsilon$ must be of order $10^{-8}\alpha$, emphasizing once more that a large $\alpha$ will not obstruct this generation mechanism. This argument will be made tighter below.

There are constraints that must be satisfied in order for this scenario to work. As mentioned at the end of the preceding section, the typical energy of particles produced in wall collisions, $\Delta^{-1}$, should exceed the Higgs-boson mass. (If this does not hold, then $f_H$ will have an extra suppression. While this may allow a larger $\epsilon$ it will most likely require some fine tuning of the amount of suppression.) A further constraint is that the reheat temperature be less than the Higgs-boson mass in order to avoid the Higgs bosons produced in wall collisions reaching a state of thermal equilibrium. These two constraints translate into an upper and lower bound for $\lambda_H$ (neglecting the volume factors and substituting for $g_*$ as before):

$$\lambda^{1/2} > \lambda_H^{1/2} > \frac{1}{\sqrt{10}} \left[ \frac{\xi}{\lambda} \right]^{1/4} \lambda^{1/2} . \tag{4.2}$$

Clearly suitable values of $\lambda_H$ are only possible provided

$$\left[ \frac{\xi}{\lambda} \right]^{1/4} < \sqrt{10} , \tag{4.3}$$

although to allow a range of $\lambda_H$ this bound should be stronger. Although this provides a nontrivial constraint on $\xi$ and $\lambda$ that $\xi < 100\lambda$, it is not a particularly strong one, and it leads only to a weak lower bound of $\epsilon$ of around $10^{-8}\alpha$. Hence this baryogenesis scenario appears satisfactory for a large set of possible model parameters.

We should discuss more general requirements of the model in order for this type of extended inflation scenario to be considered as a sensible candidate for baryogenesis. One important requirement is that the symmetry breakings do not lead to an unacceptable density of relic monopoles (monopoles being the inevitable outcome of any symmetry breaking from a semisimple group to the standard model). In several theories this can be arranged by creating the monopoles in a preinflationary breaking. The monopoles are then subsequently diluted during the inflationary era and present no further problems. For example, in a specific model proposed by Olive and Turok,[23,24] SO(10) is broken in a two-step process:

$$\begin{aligned} SO(1) &\to SU(3)\otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L} , \\ &\to SU(3)\otimes SU(2)_L \otimes U(1)_Y \otimes Z_2 , \end{aligned} \tag{4.4}$$

where the first breaking is through the 45-dimensional representation of SO(10), and the second breaking is through the 126-dimensional representation.[25] In the symmetry-breaking scheme of Eq. (4.4), monopoles are produced at the first transition, but not at the second (existing monopoles are converted rather than new ones formed[24]). Hence, in such a picture at least two symmetry breakings are needed to reach the standard model, the latter causing both inflation and the out of equilibrium conditions required for baryogenesis. It is even possible for this later transition to produce cosmic strings as defects in the inflaton field.[26]

We should also comment on the possible role of the baryon-number-violating anomalous currents in the nonperturbative sector of the electroweak theory. It has been conjectured[27] that this anomalous current may cause a "wash-out" of any pre-existing baryon asymmetry at temperatures above the electroweak phase transition as the barrier height between sectors of different baryon number may only be around 10 TeV, and baryon-number-nonconserving interactions could proceed in thermal equilibrium. If these calculations prove correct, then this may destroy any baryon asymmetry generated in the wall collisions (this would wash out any primordial baryon asymmetry of course, and is not a problem specific to the scenario we are proposing here). It has also been proposed that the electroweak phase transition may actually create a suitable baryon asymmetry but so far these models have exhibited only limited success. The effect of sphalerons may be mitigated if a non-zero value of $B - L$ is generated, such as possible with the breaking scheme of Eq. (4.4).

We also note the possibility of isothermal perturbations arising from the thermalization process. While we have assumed throughout this paper that at percolation all the true-vacuum bubbles have the same size, the full picture is somewhat more complicated, as bubbles formed earlier in inflation will grow to larger sizes than those formed right at the end. While homogeneity of the microwave-background requires large bubbles to be suppressed,[28] one would still expect to see a range of sizes of small bubbles, and hence spatial variations in the ratio of baryon-number density to entropy density from point to point.

In conclusion then, we have seen that the result of the first-order transition bringing extended inflation to an end is an environment well out of thermal equilibrium. In such conditions baryogenesis via the decay of baryon-number-violating Higgs particles can proceed, and we have demonstrated a means by which the baryon number can be estimated. The mechanism has further been shown to work for a large range of model parameters and to have the capability of predicting a baryon asymmetry of the required magnitude. In a second paper, we shall consider a slightly different mechanism for baryogenesis in which primordial black holes formed during the bubble-wall collisions may play an important role.

## ACKNOWLEDGMENTS

E.J.C., A.R.L., and E.W.K. were supported by the United Kingdom Science and Engineering Research Council. E.W.K. was also supported by the Department of Energy and NASA (Grant No. NAGW-1340) at Fermilab. E.W.K. would also like to thank John Barrow for his hospitality at the University of Sussex where part of this work was done. J.D.B. and A.R.L. would like to thank Ian Moss for discussions on bubble-wall collisions. E.J.C. and E.W.K. would like to thank Erick Weinberg for his helpful comments.

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