Fermi National Accelerator Laboratory

SUSSEX-AST 90/5-1
FNAL-PUB-90/62-A
(May, 1990)

Baryogenesis in Extended Inflation. I. Baryogenesis via Production and Decay of Supermassive Bosons

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Abstract

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.

Operated by Universities Research Association Inc. under contract with the United States Department of Energy

I. INTRODUCTION

Recently the spirit of the orginal 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 tunnelling; 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,[7,8,9,10] 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, $n_B = n_b = 1.13\times10^{-5}(\Omega_B h^2)$ 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 $s \simeq 2970$ 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 $0.010 \leq \Omega_B h^2 \leq 0.017$,[13] which implies $B = (3.81 \text{ to } 6.48) \times 10^{-11}$. 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 non-equilibrium conditions must obtain. 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 (GUTs). 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 is also taken to be different from the inflaton. The Higgs of GUTs naturally violate B. The origin of the C and CP violation necessary for baryogenesis is uncertain. It is practical simply to parameterize the degree of C and CP violation in the decay of the particle. To illustrate such a parameterization, imagine that some Higgs scalar H has two possible decay chanels, 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 Γ_H is the total decay width. Of course ϵ can be calculated if one knows the masses and couplings of the relevant particles. Reasonable upper bounds for ϵ 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 non-equilibrium 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 Section 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 gravitation sector is that the parameter that determines the efficiency of bubble nucleation, $\epsilon_N(t) = \Gamma_N(t)/H^4(t)$, where Γ_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 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 σ, 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:

1. σ_0, the energy scale for SSB, i.e., the VEV of the scalar field.

2. λ, a dimensionless coupling constant of the inflaton potential. We will assume that the potential is proportional to λ.

3. ξ, a dimensionless number that measures the difference between the false and the true vacuum energy density via $\rho_V = \xi\lambda\sigma_0^4$. ξ 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 tunnelling to the true vacuum. In the thin-wall approximation, the size of a nucleated bubble is given by

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

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 crucial 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.

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 the bubble at the end of extended inflation is still approximately R_C.

Consider first the growth of the bubble in comoving 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

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

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

$$R(t_{\rm END}) \equiv R = r(t_{\rm NUC})a(t_{\rm END}) = R_C[a(t_{\rm END})/a(t_{\rm NUC})]. \tag{2.5}$$

With the assumption that successful extended inflation will have a burst of bubble nucleation at the end of inflation, and that there is not much growth in the size of a typical bubble between the time of nucleation and when it collides with other bubbles at the end of inflation, for a first approximation we will assume that at the end of extended inflation all bubbles have the same size, $R = \alpha R_C$, where $\alpha \equiv a(t_{\rm END})/a(t_{\rm NUC}) \sim \mathcal{O}(1)$. In the concluding section we will discuss possible implications of relaxing this assumption.

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 Universe is first percolated by bubbles of true vacuum. The "filling fraction" of the Universe at percolation is some fraction f_V. Based upon Monte-Carlo simulations, it is expected that $f_V \sim 0.3$.[21] The final step is the release of this energy into radiation via bubble-wall collisions.

III. BARYOGENESIS BY DIRECT PRODUCTION OF SUPERMASSIVE BOSONS

Let us concentrate on a single bubble of radius R. The total mass of the bubble is

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

The collisions of the bubble walls produces some spectrum of particles, which are subsequently thermalized. The mean energy of a particle produced in the collision is of order of the thickness of the wall, $\langle E\rangle \sim \Delta^{-1}$, and hence the mean number of particles produced in the collisions is

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

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 mass exceeds Δ^{-1} by a significant amount, we can expect some suppression, presumably exponential, in the number of Higgs formed. This possibility will be discussed at the end of this section. For now, we simply parameterize the fraction of the primary annihilation products that are supermassive Higgs by a fraction f_H. 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 is from the primary particles produced in the bubble-wall collisions. This will be true if the reheat temperature, T_{RH}, is below the Higgs 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 ϵ per decay, where ϵ 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

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

We must now calculate the entropy generated in bubble-wall collisions. As stated above, the mass of a bubble is $M = 4\pi\sigma_0^3\lambda^{1/2}R^2$. Thermalization of the mass in the bubble walls will redistribute this energy throughout the bubble, resulting in a radiation energy density (recall that at percolation, a fraction f_V of the Universe is taken up by bubbles of radius R)

$$\rho_R \sim f_V M/(4\pi R^3/3) \sim 3 f_V \lambda^{1/2}\sigma_0^3/R. \tag{3.6}$$

We may now express the reheat temperature to the radiation energy density via

$$\rho_R = \frac{g_* \pi^2}{30} T_{RH}^4, \tag{3.7}$$

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:

$$s = \frac{2\pi^2}{45} g_* T_{RH}^3 \sim 2.3\; g_*^{1/4} f_V^{3/4} \lambda^{3/8} \sigma_0^{9/4} R^{-3/4}. \tag{3.8}$$

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

$$B \equiv n_B/s \sim \epsilon f_V^{1/4} f_H g_*^{-1/4} \lambda^{-3/8} \left(\sigma_0 R\right)^{-1/4}. \tag{3.9}$$

As stated above, a successful extended inflation model is expected to have the feature that a typical bubble does not grow significantly between nucleation and percolation, so we can assume that the bubble size at percolation is $R = \alpha R_C = \alpha(\xi\lambda^{1/2}\sigma_0)^{-1}$. The baryon asymmetry B is then given by

$$B = \epsilon f_V^{1/4} f_H \, \alpha^{-1/4} g_*^{-1/4} \lambda^{-1/4} \xi^{1/4}. \tag{3.10}$$

Provided the mass of the Higgs is less than T_{RH}, f_H is just g_H/g_*, where g_H is the number of Higgs degrees of freedom. Substituting this in gives the final result

$$B = \epsilon f_V^{1/4} g_H \, \alpha^{-1/4} g_*^{-5/4} \lambda^{-1/4} \xi^{1/4}. \tag{3.11}$$

This allows us to make numerical estimates of B based on sample values of these parameters. Notice that the dependence on both λ and ξ, which are the two parameters on which the inflaton's potential depends, is very weak, as is that on the filling factors f_V and α. The important contributions are the degree of asymmetry in CP violating Higgs decays and the number of particle species available for production in the wall collisions. Numerical estimates for B based upon this expression will be made in the concluding section.

We now elaborate upon the implications of two assumptions of our scenario. The first assumption is that the mass of the Higgs 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 mass is of order $\lambda_H^{1/2}\sigma_0$, where λ_H is the coupling constant of the quartic term in the Higgs potential coupling σ and H. Clearly $\lambda_H^{1/2}$ must not be too much larger than $\lambda^{1/2}$, or there will be a large supression in f_H.

The second assumption is that the reheat temperature is less than the mass of the Higgs, so that thermal production of H is not important. This implies that $m_H > \lambda^{1/4} g_*^{-1/4} f_V^{1/4} \alpha^{-1/4} \xi^{1/4} \sigma_0$. Again assuming that $m_H = \lambda_H^{1/2}\sigma_0$, the requirement becomes $\lambda_H > \lambda^{1/2} g_*^{-1/2} f_V^{1/2} \alpha^{-1/2} \xi^{1/2}$. 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.

IV. DISCUSSION AND CONCLUSIONS

Here we examine some typical numbers for the baryon asymmetry which may be obtained from Eq. (3.11), in the light of the experimental limits discussed in the introductory section setting B at around 10^{-10}. The two filling factors f_V and α both appear to the quarter power, and hence as they are both of order one, they can be dropped. The number of Higgs degrees of freedom, g_H, is expected also to be of order one, with simply one degree of freedom for each polarization in the case of a single Higgs 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

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

The remaining parameters ϵ, λ and ξ are less certain, with some dependence on the particular unified theory under examination, though it is reassuring that both λ and ξ also enter only to the quarter power and hence the dependence on these quantities is weak. This does however have the further implication that ϵ should be very small, as we shall shortly see. That a suitable baryon asymmetry can be produced with such a small ϵ indicates that the bubble wall collisions are very efficient in producing a baryon asymmetry. A reasonable estimate of ξ 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 λ should probably be less than of order 1, though nothing in principle prevents it from being much smaller; note that a smaller λ increases the baryon asymmetry as it leads to a less efficient production of entropy, though λ must also be sufficently large that Higgs particles can be produced in the wall collisions. From these arguments, it seems likely that the ratio ξ/λ will be within a few orders of magnitude of unity, implying that if this mechanism is to generate the appropriate baryon asymmetry ϵ must be of order 10^{-8}. 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, Δ^{-1}, should exceed the Higgs mass. (If this does not hold, then f_H will have an extra suppression. While this may allow a larger ϵ 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 mass in order to avoid the Higgs produced in wall collisions reaching a state of thermal equilibrium. These two constraints translate into an upper and lower bound for λ_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 λ_H are only possible provided

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

although to allow a range of λ_H this bound should be stronger. Although this provides a non-trivial constraint on ξ and λ that $\xi < 100\lambda$, it is not a particularly strong one and it leads only to a weak lower bound on ϵ of around 10^{-8}. 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 breaking does not lead to an unacceptable density of relic monopoles (monopoles being the inevitable outcome of any symmetry breaking from a semi-simple group to the standard model). In several theories this can be arranged by creating the monopoles in a pre-inflationary 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

$$\begin{aligned} SO(10) \quad &\longrightarrow \quad SU(3) \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}, \\ &\qquad\qquad \longrightarrow SU(3) \otimes SU(2)_L \otimes U(1)_Y \otimes Z_2, \end{aligned} \tag{4.4}$$

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 anomolous currents in the non-perturbative sector of the electroweak theory. It has been conjectured[27] that this anomolous 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 non-conserving 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 should also comment here on 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 phase 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.

Acknowledgements

EJC and ARL are supported by the SERC, and EWK was an SERC-supported Visiting Research Fellow at the Astronomy Centre, University of Sussex. EWK was also supported by the DoE and NASA (grant #NAGW–1340) at Fermilab. EWK would also like to thank John Barrow for his hospitality at the University of Sussex where part of this work was done.

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