THERMODYNAMIC GENERATION OF THE BARYON ASYMMETRY

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A new technique for generating the baryon asymmetry of the universe is discussed. By using the natural *CPT* non-invariance of the universe during its early history, we discuss how a baryon asymmetry can develop while baryon violating interactions are still in thermal equilibrium with respect to an effective hamiltonian. Furthermore, both the ground state and fundamental interactions in these theories can be *CP* conserving.

1. Introduction and conclusions. One of the most peculiar physical parameters of the universe today is the baryon to photon ratio:

$$\sigma \equiv N_B/N_\gamma \simeq 10^{-9} . \tag{1.1}$$

One would like to explain this number on the basis of microphysics rather than on initial conditions. Since a non-zero value for σ breaks baryon number and *CP*, we expect *CP* and baryon violation to play an important role in its explanation. However, *B* and *CP* violating interactions are not sufficient for producing the observed asymmetry; as pointed out by Sakharov [1] and elaborated by Dimopoulus and Susskind [2], *CPT* invariance requires that the baryon number must be generated out of thermal equilibrium: *CPT* ensures that the energy spectra for baryons and anti-baryons will be identical, and consequently baryon and anti-baryon thermal distributions will be identical as well. In a finite volume there will still be a baryon asymmetry, but only due to random fluctuations, yielding an asymmetry $N_B \simeq \sqrt{N_\gamma}$. Since $N_\gamma \simeq 10^{80}$ in the visible universe, these fluctuations cannot account for the relatively large value of σ in eq. (1.1). Making use of this reasoning, many authors have followed Yoshimura's lead [3,4], offering explanations for the value of σ in terms of out-of-equilibrium decay of sundry relic particles.

In this paper we show that the above conclusion is avoidable – that baryogenesis can in fact occur while baryon violating interactions are still in thermal equilibrium. This can occur because *CPT* need not be realized as a good symmetry in the early universe. After all, we know that an expanding universe at finite temperature violates both Lorentz invariance and time reversal. We will show that it is possible for this lack of symmetry to lead to *effective CPT* violating interactions among the baryons. Of course, this can only happen in an evolving universe, and the value of the Hubble parameter plays an important role in determining the size of the *CPT* violation; at zero temperature all interactions must be *CPT* invariant if derivable from a relativistic field theory which has a Lorentz invariant ground state.

To illustrate the general phenomenon we have in mind, consider a theory in which a neutral scalar field is derivatively coupled to the baryon current:

$$\mathcal{L}_{\mathcal{I}} = (1/\Lambda)\partial_\mu \phi j^\mu_B . \tag{1.2}$$

This is a dimension-five effective operator, assumed to result from some unspecified dynamics at a scale of order $\Lambda \gtrsim T$, where T is the temperature. We will refer to the ϕ particle as the thermion. Suppose that the thermion field can be given cause to develop a slowly varying time derivative as the universe cools: $\dot\phi/\Lambda \equiv \mu$. Then replacing $\partial_\mu\phi$ in eq. (1.2) by its classical value leads to the effective interaction

$$\mathcal{L}_{\mathcal{I}} = \mu n_B , \tag{1.3}$$

where $n_B \equiv j^0_B$ is the baryon density. Note that this is a *CPT* violating interaction, which shifts the baryon

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and anti-baryon spectra relative to each other. If baryon number is conserved, this shift in the energy spectra will not allow a baryon asymmetry to develop since the total charge, B, cannot change. However, if there are baryon violating interactions in thermal equilibrium, the baryons and anti-baryons will equilibrate with different thermal distributions. Note that the parameter μ affects the thermal distributions exactly as would a chemical potential for baryon number. Occasionally we will refer to μ as an "effective chemical potential". The net baryon density is roughly given by

$$n_B(T) \simeq \mu T^2 + O(\mu/T) . \tag{1.4}$$

The true ground state of the universe must have $\langle \partial_\mu \phi \rangle = 0$. Thus μ must approach zero as the universe cools. If baryon violation remained in thermal equilibrium as $\mu \to 0$, the baryon asymmetry would disappear. If, however, B violation drops out of thermal equilibrium before μ relaxes to zero, the baryon number (1.4) becomes frozen in place. When μ subsequently returns to zero, one is left with a net baryon asymmetry.

Any neutral scalar field will generically acquire effective couplings to the baryon current as in eq. (1.2) if not forbidden by some discrete symmetry [#1]. It remains to be shown how the thermion can develop an expectation value for its velocity. In this paper we consider two general scenarios:

(1) *Symmetry breaking.* During a second order phase transition, a scalar order parameter must evolve from one field value to another as the true minimum of its effective potential changes. Then in general there will be a period during which the velocity of the field develops an expectation value. Thus the thermion could be associated with the spontaneous breaking of a discrete symmetry. Alternatively, one could consider a charged scalar Φ which develops a vacuum expectation value (VEV). Φ cannot itself couple to the baryon current since it is charged, but its phase θ (the resultant Goldstone boson) can. If the U(1) symmetry is only approximate, immediately after Φ develops a VEV, θ will be at some initial value θ_i which will not in general be at the true minimum of its potential. It will then slowly roll to its true minimum. We discuss this model in detail in the next section and consider the possibility that Φ could be the inflaton.

(2) *Non-minimal gravity couplings.* If the ϕ field couples directly to the Ricci scalar R in the lagrangian, then the trace of the stress–energy tensor serves as a source for ϕ. Examples of such scalars are the Brans–Dicke scalar, and the dilaton which arises in string theories [#2]. In such models ϕ will acquire a time derivative directly from the expansion of the universe, with $\dot{\phi}/\phi \simeq H$ typically. In models like these, one finds the simple relation $\sigma \simeq T_D/M_{Pl}$, where T_D is the decoupling temperature for baryon violating interactions. With this simple relation, $\sigma \simeq 10^{-9}$ requires $T_D \simeq 10^{10}$ GeV.

There are several problems that we must address. Firstly, we must consider what happens to the baryon number we have created after the baryon violating interactions have dropped out of thermal equilibrium, and the derivative coupling stops looking like an effective chemical potential. It is possible that the evolution of the universe will wash out the baryon number that we have generated. In fact, as we shall see in section 3, this is not the case – the subsequent dynamics will, in most cases, *increase* the baryon to photon ratio. In fact, the development of a baryon asymmetry below the decoupling temperature will turn out to be a possible source of baryogenesis even without the thermal generation above decoupling.

Secondly, at least for the symmetry breaking alternative, we must worry about spatial homogeneity of the effective chemical potential. In particular, how are the initial conditions which govern the evolution of the scalar field to be set uniformly, so that the baryon asymmetry averaged over the observable universe today is not zero? In section 4 we describe several alternatives for insuring this spatial homogeneity.

In summary, what we demonstrate in this paper through the example of the simple interaction in eq. (1.2) is that it is possible for CPT violation to be fed into the baryon–anti-baryon spectrum. This allows for generation of the baryon asymmetry during an era when baryon violating interactions are still in

[#1] The current need not be the baryon current itself; any current not orthogonal to B would work, such as $B-L$ for example.

[#2] The dilaton may or may not be much lighter than the Planck scale.

thermal equilibrium. This asymmetry gets frozen in place at the decoupling temperature, T_D, when baryon violation goes out of equilibrium. Since this effect is due to the expansion of the universe, the Hubble constant, H, plays an important role. In fact, there is a natural value for the baryon to photon ratio, which is

$$\sigma \simeq H(T_D)/T_D \simeq T_D/M_{Pl} . \tag{1.5}$$

In these scenarios the size of the baryon asymmetry is determined by the hierarchy of the interactions; σ is on the order of 10^{-9} because the baryon violating interactions decouple at a temperature of order 10^{10} GeV.

We intend the model we discuss below to serve as a paradigm for thermodynamic baryogenesis; the essential feature is the ability of the natural CPT violation in the early universe to alter the baryon spectrum. There are undoubtedly many ways for this to occur, of which the model below is just a simple example.

2. Symmetry breaking and thermal baryon number. We begin by realizing the symmetry breaking scheme outlined in the introduction. We will derivatively couple the Goldstone boson produced from the spontaneous breakdown of a global U(1) symmetry to a current composed of light quarks and leptons, which we assume is not orthogonal to B. In order to obtain an expectation value for the time-derivative of the Goldstone field, we will also include a potential term in the lagrangian – that is, we will softly break the U(1) symmetry explicitly, as well as spontaneously. We will refer to the resulting pseudo-Goldstone boson as the thermion. Since we assume that the thermion arises from the spontaneous breakdown (at a temperature $T_0=f$) of some approximate U(1) symmetry, we can write the lagrangian in terms of an angular variable θ which runs from 0 to 2π. In the notation of the introduction, $\theta=\phi/f$. For simplicity, we will assume in this paper that the broken U(1) is proportional to baryon number itself, although this is not necessary. Then

$$\mathscr{L}_\theta = f^2 \partial_\mu\theta\partial^\mu\theta - V(\theta) - q\partial_\mu\theta j^\mu_B , \tag{2.1}$$

where q is a dimensionless constant. $V(\theta)$ is some arbitrary periodic potential, which we will approximate by

$$V(\theta) \simeq \tfrac{1}{2}m^2f^2\theta^2 , \quad m \ll f . \tag{2.2}$$

We will also assume that the full theory has baryon violating interactions which go out of thermal equilibrium at a temperature T_D during the evolution of the universe. It is not necessary for our purpose to know the form of these interactions, nor to specify the particle content of the theory at high temperatures.

The equation of motion that we get from the above lagrangian for a spatially constant θ field in a Robertson–Walker universe is

$$\ddot\theta + 3H\dot\theta + m^2\theta = -(q/f^2)(\dot n_B + 3Hn_B) , \tag{2.3}$$

where n_B is the baryon density and H is the Hubble constant. We will assume that the universe is flat and radiation dominated so that

$$H = 1/2t = \kappa T^2/M_{Pl} . \tag{2.4}$$

In the standard model $\kappa \simeq 17$ above the SU(2) $\times$U(1) symmetry breaking scale.

We wish to consider the thermodynamics of this model as the universe cools down from the symmetry breaking temperature, f. Above the decoupling temperature, the effect of the interaction in eq. (2.1) is as discussed in section 1; provided that the rate of change of $\dot\theta$ is sufficiently slow, this interaction will shift the baryon and anti-baryon energy levels like a chemical potential for baryon number:

$$\mu_B = -q\dot\theta . \tag{2.5}$$

"Sufficiently slow" simply means that the typical time scale of baryon violating interactions, $\tau_{\Delta B}$, must be fast enough to maintain thermal equilibrium:

$$\tau_{\Delta B}(T) < \dot\mu_B/\mu_B = \ddot\theta/\dot\theta . \tag{2.6}$$

If this condition is satisfied, these interactions will be able to track along with the time dependent effective chemical potential μ_B. Then it is a simple exercise in statistical mechanics to determine the baryon density, which to first order in μ/T is given by [#3]

$$n_B = -\tfrac{1}{12}Bq\dot\theta T^2 , \tag{2.7}$$

where B is the absolute value of the baryon number

[#3] This is an extremely good approximation for our purposes, since we will find $\mu_B \simeq H \ll T$.

summed over each spin degree of freedom. Thus when the conditions discussed above are satisfied, the equation of motion (2.3) can be written as:

$$\ddot{\theta}+3H\dot{\theta}+m^2\theta$$
$$=-\tfrac{1}{12}q^2B[(\partial_t+3H)(\dot{\theta}H/H_0)]\ , \qquad (2.8)$$

where H_0 is the initial value of the Hubble constant, at temperature $T_0=f$.

We will first consider this equation in the limit $q\to 0$, i.e., without the damping due to baryon creation. The solution to (2.3) is then given by

$$\theta(z)=(2z)^{-1/4}[C_-J_{-1/4}(z)+C_+J_{1/4}(z)]\ ,$$
$$z\equiv mt=m/2H=mM_{\rm Pl}/2\kappa T^2\ . \qquad (2.9)$$

We impose the initial conditions $\theta=\theta_0$, $\dot{\theta}=0$ at $z=z_0\ll 1$. Then we find, to leading order in z_0

$$C_-=-\theta_0\tfrac{5}{4}\Gamma(\tfrac{3}{4})z_0^{5/2}\ ,$$
$$C_+=(\theta_0/\sqrt{8})\Gamma(\tfrac{1}{4})\gg C_-\ . \qquad (2.10)$$

This solution for θ is graphed in fig. 1 as a function of the logarithmic of the temperature. During the period when $H\gg m$ the θ field changes slowly, decreasing like a power of T. Eventually the Hubble constant becomes comparable with m, and θ begins oscillating with a frequency m, with an amplitude that damps out as $T^{3/2}$.

Including the right-hand side of (2.8) in the equation of motion has little effect on the solution for θ. Because of the factor H/H_0 in (2.8) baryon damping is apparently only significant at early times. In fact, numerical integration of eq. (2.8) shows that the term on the right side of the equation has only a negligible effect on the evolution of the θ field for *all* times, provided that $z_0\lesssim 0.1$ [#4].

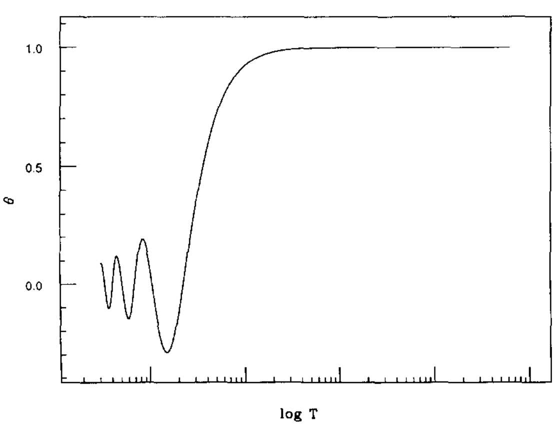

Fig. 1. The solution for θ as given by (2.9) and (2.10) versus the logarithm of the temperature. $\theta_0=1$, $z_0=10^{-3}$.

As the universe cools below $T_0=f$, the baryon density remains in thermal equilibrium an effective chemical potential μ_B, which is decreasing. Eventually we reach the decoupling temperature T_D, where the baryon violating interactions drop out of equilibrium. From then on, the baryon distribution functions will continue to look thermal, but with constant effective chemical potential given by the value of μ_B when the baryon violation freeze-out occurs. To calculate the baryon to photon ratio, we simply evaluate the effective chemical potential at the decoupling point, $z=z_D\equiv m/2H_D$.

The decoupling temperature is determined by equating baryon violating interaction time to the expansion rate

$$\tau_{\Delta B}(T_D)\simeq H_D^{-1}\ . \qquad (2.11)$$

However, we must also insure that the condition for slow variation of the effective chemical potential [eq. (2.6)] is always satisfied. From eq. (2.9) we see that at early times, $\ddot{\theta}/\dot{\theta}\simeq m/z\simeq H$. However, at somewhat later times, $\ddot{\theta}/\dot{\theta}$ drops to m, and then remains approximately constant. This crossover occurs when $H\simeq m$, that is when the θ field enters its oscillating phase. Then the condition for slow variation (2.6) may be rewritten as

$$\tau_{\Delta B}(T)<\min\{H^{-1},m^{-1}\}\ . \qquad (2.12)$$

This condition in conjunction with (2.11) guarantees that decoupling must occur before θ begins to oscillate; or equivalently, that

$$z_D<1\ . \qquad (2.13)$$

Consequently we can approximate the Bessel functions of (2.9) by the first few terms of their Taylor series expansions (note that $z\geqslant z_0$):

$$\dot{\theta}=-m\tfrac{2}{5}\theta_0 z[1-(z_0/z)^{5/2}]\ . \qquad (2.14)$$

[#4] We remind the reader that the condition $z_0\equiv m/2H_0\ll 1$ means that the thermion's mass is small enough so that its motion is strongly Hubble damped at the initial temperature $T_0=f$.

Thus our final result for the baryon number to photon number ratio, σ, is

$$\sigma=\frac{\pi^2}{60\zeta(3)}\left(\frac{\theta_0 qB}{2\kappa}\right)\frac{m^2 M_{\rm Pl}}{T_{\rm D}^3}\left[1-\left(\frac{T_{\rm D}}{f}\right)^5\right]. \qquad (2.15)$$

Note that σ depends only weakly on the symmetry breaking scale f, so that it is specified primarily by the remaining free parameters in our scenario, m and $T_{\rm D}$. These two parameters are constrained by the condition (2.13). Setting σ equal the observed value of 10^{-9}, this constraint forces the decoupling temperature $T_{\rm D}$ to be greater than 10^8 GeV.

3. Oscillating baryon number. In this section, we will consider the subsequent development of the baryon asymmetry after baryon violating interactions have fallen out of equilibrium. Below $T_{\rm D}$, the interaction of the thermion with the baryon current cannot be interpreted as an effective chemical potential for baryon number, and we must solve the equations of motion directly. We will begin by considering the lagrangian of the previous section (2.1) without the expansion of the universe.

We can of course integrate by parts in the interaction term to obtain

$$q\theta\partial_\mu j^\mu_{\rm B}\,. \qquad (3.1)$$

If the baryon current were conserved, the divergence would vanish. However, since baryon number is assumed violated, the divergence may be replaced by the operator that violates baryon number. We will not specify this operator, but will assume that it gives rise to a decay of the thermion field with a width Γ. (For example, if the operator is of dimension six with a coupling of Λ^{-2}, then $\Gamma\simeq q^2m^5/\Lambda^4$.) Then we may approximate the effect of the decay of the motion of the thermion field due to its baryon violating interaction by including an extra term in the equations of motion:

$$\ddot\theta+m^2\theta+\Gamma\dot\theta=0\,. \qquad (3.2)$$

This is the equation of motion for a damped harmonic oscillator; the θ field oscillates with a frequency m with an amplitude that decreases exponentially with time.

By comparing (3.2) with the equation of motion derived from (2.1) we see that the baryon density as a function of time is given by

$$\dot n_{\rm B}=-\Gamma f^2\dot\theta q^{-1}\,. \qquad (3.3)$$

Integrating this equation we obtain

$$n_{\rm B}(t)=n_{\rm B}(t_0)-\Gamma f^2q^{-1}[\theta(t)-\theta(t_0)]\,. \qquad (3.4)$$

As the θ field oscillates, so does the baryon number. As $t\to\infty$, θ damps out to zero. (We have assumed that the true minimum of the potential is at $\theta=0$.) However, the net baryon density left is not zero, but rather depends on how far from the true minimum of the potential we started the θ field. This result is more transparent if we consider a graph of $\dot\theta$ in fig. 2. When the velocity of the field is greater than zero, the decay of the thermion produces mostly baryons; when the velocity is negative, the decay produces mostly anti-baryons. Since the amplitude of the θ field is damped, the regions where the velocity is positive are not equivalent to the regions where the velocity is negative. Hence as $t\to\infty$ the net baryon density is determined by the net asymmetry, which is in turn only a function of the difference in initial and final positions of θ.

The important point to remember is that we are now below the decoupling temperature, which means that all baryon violating interactions are out of thermal equilibrium, and are no longer able to affect the baryon density we have previously calculated. The only important baryon violating effect is the conversion of energy stored in the θ field oscillations into baryons.

To finish our computation we must now include

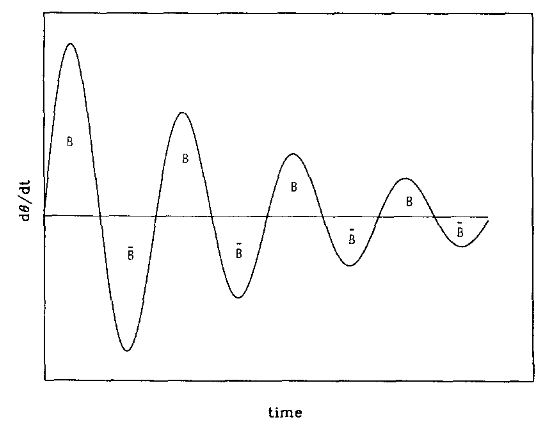

Fig. 2. $\dot\theta$ versus time below decoupling. $\dot\theta>0$ implies baryon production, while $\dot\theta<0$ implies antibaryon production.

the effects of the expanding universe. All of our previous equations may be directly transcribed by replacing any volume factors by a comoving volume proportional to R^3; equivalently, we may simply include the appropriate metric factors in the lagrangian. In either case, we see that, provided $H \ll \dot{\theta} \simeq m$, and the expansion serves only to dilute the amount we have previously calculated by the ratio of the scale factor cubed.

At large times, $\theta(t)$ approaches the true minimum, and defining $\Delta\theta \equiv \theta(t_0) - \theta(\infty)$, our expression for the baryon to photon ratio becomes

$$\sigma = \frac{\pi^2}{2\zeta(3)T^3}[n_B(t_0) + \Gamma f^2 \Delta\theta q^{-1}] \times \left(\frac{R(t_0)T(t_0)}{R(t)T(t)}\right)^3 T^3(t_0)$$

$$= \sigma_0 + \frac{\Gamma f^2 \Delta\theta}{T^3(t_0)} \left(\frac{R(t_0)T(t_0)}{R(t)T(t)}\right)^3 \frac{q^{-1}\pi^2}{2\zeta(3)}. \tag{3.5}$$

Since the evolution of the universe is approximately isentropic, we may approximate the factor in brackets by 1, that is we may treat RT as a constant [#5]. Since the oscillations of θ are assumed to begin when the baryon violating interactions fall out of thermal equilibrium, $T(t_0)$ is just T_D, the decoupling temperature, assuming that $T_D \lesssim f$. Our final result is

$$\sigma = \sigma_0 + \frac{q^{-1}\Gamma f^2 \Delta\theta}{T_D^3} \frac{\pi^2}{2\zeta(3)}. \tag{3.6}$$

Thus we conclude that, provided decoupling occurs before the thermion has rolled past its true minimum, the effect of the subsequent evolution of the thermion after decoupling is to *increase* the baryon to photon ratio.

Since the baryon asymmetry computed in this section just adds to the result from section 2, we see that the generation of an asymmetry from oscillations may occur quite independently of thermodynamic generation. Our result only depended on the initial value of θ at a temperature T_D, and so the result (3.10) remains true, even if $\sigma_0 = 0$. This phenomenon is similar to what is seen in the models of ref. [5], although in these papers spontaneous symmetry breaking plays no role.

[#5] This is not exactly true, since the entropy contained in the θ field is not present after the decay. However, since this is only a small fraction of the total entropy at the temperatures we are interested in, we may ignore its effect.

4. Spatial homogeneity. We have shown in the previous two sections that the baryon asymmetry may be generated quite naturally by the classical motion of a scalar field in its potential. In the first model of a PGB, we found two distinct regimes: the thermodynamic regime, corresponding to small acceleration of the scalar field ($\ddot{\theta}/\dot{\theta} \ll \tau_{\Delta B}^{-1}$), and the supercooled oscillating regime corresponding to relatively large accelerations of the scalar field ($\ddot{\theta}/\dot{\theta} \gg \tau_{\Delta B}^{-1}$). In both cases, whether baryons or anti-baryons are produced is an accident of the initial conditions; namely whether the field has to travel clockwise or counter clockwise to get from its initial value to its final position. Either possibility is equally likely. This arbitrariness follows from the fact that the model considered so far has no explicit CP violation.

It might seem that our scenarios have accomplished nothing: in a Friedmann universe the initial conditions for the θ field cannot possibly be homogeneous over distance scales greater than the horizon size at the initial temperature, $T_0 \simeq f$. Averaging the initial value of $\theta_0 \simeq \pm 1$ over the region that is to become today's horizon, one finds $\langle\theta_0\rangle_{\mathrm{RMS}} \simeq [10^{-4}\ \mathrm{eV}/f]^3$, where 10^{-4} eV is today's temperature. Even if f were as low as the $SU(2)\times U(1)$ breaking scale, one would find $\langle\theta_0\rangle_{\mathrm{RMS}} \simeq 10^{-45}$, certainly a number too small to explain why $\sigma = 10^{-9}$.

This difficulty is easily resolved by including CP violation in the original lagrangian. We could, for example, bias the fluctuations of θ_0 by including CP violating self-interactions for the thermion. This would destroy the $\theta \to -\theta$ symmetry of $V(\theta)$ and allow the spatially averaged value of the final velocity of the thermion field to be non-zero. The net baryon asymmetry of the universe can then be calculated by replacing the effective chemical potential in (2.6) by its spatially averaged value at decoupling, and replacing $\Delta\theta$ by its spatially averaged value in (3.10).

It might be thought that the introduction of CP violation would have been inevitable, following the arguments of Sakharov. However, in our symmetry breaking model this violation was necessary only to produce a net baryon asymmetry when averaged over space. If we prefer we can eliminate this explicit violation by using inflation [6].

Inhomogeneity is a problem inherent in Friedmann cosmologies, a problem that inflationary cosmologies were designed to cure. By introducing a period of exponential expansion into the thermal history of the universe, our entire horizon today can be fashioned out of a single correlated region at early time. For our purposes, if we asume that the symmetry breaking scale f is above the reheat temperature following inflation, $T_{\rm RH}$, then the visible universe today could be spawned from a region that had a homogeneous initial value θ_0, exactly as tacitly assumed in the previous sections. The only requirement is that the value for θ at the *end* of inflation be O(1). This is necessary since any baryons produced before or during inflation will be washed out by the expansion of the universe. This constraint is easily satisfied by having the mass of the θ field be much smaller than the Hubble constant during the inflationary phase.

It is important to note that one cannot achieve sufficient baryon number simply by assuming that our present horizon evolved from an inflated region of space that had a slight excess of baryons. On the contrary, since baryon number is conserved, the baryon density in such a scenario would be diluted by the ratio of initial to final volumes, a disastrous amount. The reason why baryon number has not suffered a similar fate in our model is that it is stored (in a loose sense) in the potential energy of the scalar field θ. Because this is a scalar quantity, it is not diluted during inflation. Instead one has, in effect, baryon creation driven by the expansion of the universe.

One might wonder if the baryon asymmetry might not be created by the inflaton itself. If the inflaton is identified with θ then thermodynamic baryogenesis is out of the question – that only occurs during a slow rolling period; for the inflaton, such a slow rolling only occurs during the period of exponential expansion, and, the resultant baryon number would be exponentially diluted. It is conceivable, however, that the inflaton could generate baryon number while it oscillates in its potential. This would correspond to the supercooled regime discussed in the previous section.

A more attractive possibility is that the inflaton is a complex field, and that θ should be identified with its phase. Then the thermion would be a pseudo-Goldstone boson as discussed earlier, but the scale f would be naturally related to inflation.

Finally we note that thermodynamic generation of the baryon asymmetry would resolve one of the problems of inflation. In most inflation models the inflaton potential must be adjusted to allow the universe to reheat to a temperature on the order of the GUT scale in order for conventional out-of-equilibrium baryogenesis to occur. In our scenario the baryon asymmetry may be generated at temperatures well below the GUT scale, eliminating the need for large reheating.

5. *Non-minimal gravity couplings.* We conclude by briefly considering the second possibility for thermodynamic baryogenesis outlined in the introduction, that of non-minimal gravity couplings. We will assume that the thermion is a scalar coupled to the trace of the energy–momentum tensor, such as a Brans–Dicke type field. Such a scalar may develop an expectation value on the order of the Planck mass, and, during the early part of its evolution, the rate of change of ϕ is set by the Hubble constant:

$$\dot{\phi} \simeq H \langle \phi \rangle \,. \tag{5.1}$$

Then the effective chemical potential is

$$\mu = (H/\Lambda) \langle \phi \rangle \,. \tag{5.2}$$

The scale for this coupling will be set by the Planck scale, so that $\langle \phi \rangle / \Lambda \equiv q$ is O(1). Then using our previous formula we get a baryon to photon ratio

$$\sigma = \frac{\pi^2}{2\zeta(3)} \frac{B}{12} q \frac{H}{T_{\rm D}} = \frac{\pi^2}{2\zeta(3)} \frac{B}{12} q\kappa \frac{T_{\rm D}}{M_{\rm Pl}} \,. \tag{5.3}$$

The only essentially free parameter in this formula is the decoupling temperature, $T_{\rm D}$. In order for σ to equal the observed baryon to photon ratio, the baryon violating interactions must decouple at a temperature of approximately 10^8 GeV.

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