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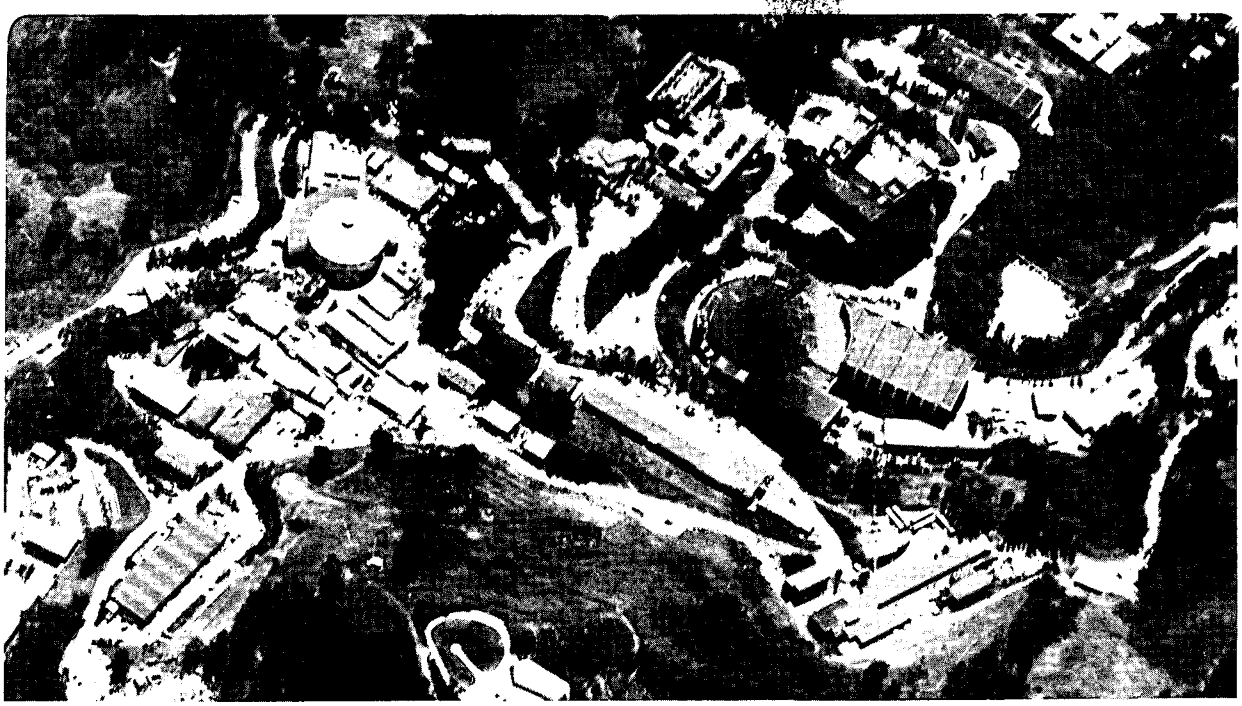

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A MECHANISM FOR BARYOGENESIS IN SUPERSYMMETRIC INFLATIONARY COSMOLOGIES*

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ABSTRACT

We present another mechanism to generate the baryon asymmetry of the universe within supersymmetric inflationary cosmologies. The gravitational coupling of the inflaton to the heavy fields in the theory is used to generate the baryon excess. We find that in models with an inflaton field and some heavy fields, there is generation of baryon number due to the transfer of energy from the inflaton to the heavy sector. We study this general mechanism for two simple models–one in which the inflaton does not break supersymmetry and one for which it does. We find that we can get the observed value of baryon to entropy ratio in these models. The thermal constraint [stabilization of the inflaton in the plateau region at high temperatures] is violated in both these models. We discuss the possibility of the introduction of direct couplings to satisfy this constraint.

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The new inflationary universe scenario provides an elegant solution to many cosmological problems of the hot big–bang model [1]. Supersymmetry on the other hand has been used to solve many serious problems in particle physics in a beautiful way [2]. In fact, inflationary scenarios employing local supersymmetry seem to be very attractive for providing "natural" solutions to many cosmological conundrums [3]. The success of these models is somewhat marred by one potentially serious problem – a low reheating temperature after the exit from the inflationary era. A low reheating temperature is undesirable because it is a potential blow to one of the most important achievements of the application of Grand Unified Theories to cosmology – the generation of baryon-antibaryon asymmetry from symmetric initial conditions [4]. This is so because in the standard scenario, in order to generate a baryon asymmetry after the de–Sitter expansion has diluted any primordial asymmetry, one needs to reheat the universe to at least a temperature of $O(10^9 - 10^{10} GeV)$ [5]. It could be argued that the standard out of equilibrium decay of the color-triplet Higgs is not the mechanism responsible for the generation of the asymmetry, but alternative mechanisms: decay of coherent Higgs field oscillations which are very far from equilibrium [6], low temperature baryon generation scenarios [7] etc. could be operative. While this may be reasonable, it still seems fruitful to us to investigate alternate origins for baryon number generation , since this feature is potentially the most restrictive on model building.

In this paper, we will investigate the possibility of generating a satisfactory baryon excess within the framework of locally supersymmetric inflationary models. More specifically, we will use the hidden sector models [8], since they seem to be the most attractive phenomenologically. ("no–scale " models [9] will not be considered here.)

These models have a very weakly coupled scalar field, the inflaton which is responsible for the de–Sitter expansion and the subsequent reheating. The very weak interactions of the inflaton imply the reheating temperature is low because the lifetime is large and there is a significant redshifting of energy [5,10,12]. This causes problems for baryosynthesis.

We investigate the possibility of remedying this situation by using other heavy fields in the theory (eg. the adjoint Higgs in SU(5)). Due to the gravitational couplings between these heavy fields and the hidden sector, energy is transferred from the inflaton to these fields. Since these fields have gauge interactions and hence a short lifetime, their decays occur before any significant redshifting has taken place, giving rise to a significant baryon excess.

After establishing a general framework in Section I, we investigate two representative models in Section II and III. Supersymmetry is unbroken in the first model, which is simpler to analyse while in the second model it is broken. We compute the baryon to entropy ratio in both these models and show that with reasonable values of various model–dependent parameters we obtain a satisfactory baryon excess. Both the models, in spite of giving a satisfactory cosmology, do not however, satisfy the thermal constraint. We find that even with the incorporation of heavy fields, the situation does not change. Finally, we comment on the finite temperature corrections and the use of direct couplings between the heavy fields and the inflation in solving the thermal constraint and its effect on our results.

1. GENERAL FRAMEWORK:

Consider a set of scalar fields ϕ_i in a locally supersymmetric theory with a superpotential $W(\phi_i)$. Then the corresponding scalar potential is given by (assuming a flat Kähler metric)[11]

$$V(\phi_i) = \exp\left(\sum_i |\phi_i|^2/M^2\right)\left[\sum_i |D_{\phi_i}W(\phi_i)|^2 - \frac{3}{M^2}|W(\phi_i)|^2\right] \quad (1)$$

where $D_{\phi_i}W(\phi_i)$ is the Kähler covariant derivative

$$D_{\phi_i}W(\phi_i) = \frac{\partial W}{\partial \phi_i} + \frac{\phi_i^* W(\phi_i)}{M^2} \quad (2)$$

and $M = \frac{M_P}{\sqrt{8\pi}} \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck Mass.

We consider the superpotential W to be a function of two fields ϕ and Σ. ϕ is the field which causes inflation, the inflaton and Σ is some heavy field in the theory. Throughout we assume that ϕ is a gauge singlet while Σ can have non-trivial transformation properties under the gauge group. We will for our purposes take Σ to be the adjoint Higgs of SU(5) but most of the results will be independent of this choice.

As a first step, we assume that the superpotential $W(\phi,\Sigma)$ be written as the sum of two superpotentials $f(\phi)$ and $g(\Sigma)$. This implies that the two fields only interact gravitationally (we will comment on the effect of direct coupling later). Then,

$$W(\phi,\Sigma) = f(\phi) + g(\Sigma) \quad (3)$$

Next we demand that at the true minimum, ϕ_0, Σ_0, the cosmological constant is zero and supersymmetry is unbroken. It is easy to show that these conditions imply

$$\left.\frac{\partial f}{\partial \phi}\right|_{\phi_0} = 0 \quad (4a)$$

$$f(\phi_0) + g(\Sigma_0) = \left.\frac{\partial g}{\partial \Sigma}\right|_{\Sigma_0} = 0 \quad (4b)$$

The most general gauge invariant and renormalizable superpotential for Σ is given by

$$g(\Sigma) = \frac{b_1}{2}Tr\Sigma^2 + \frac{b_2}{3}Tr\Sigma^3 + b_0 \quad (5)$$

where the constants b_0, b_1, b_2 will be fixed by condition (4b). It is convenient to work with dimensionless variables x and y defined as

$$x \equiv \phi/M \qquad y \equiv \Sigma/M \quad (6)$$

Then

$$g(y) = \frac{b_1 M^2}{2}Tr\ y^2 + \frac{b_2 M^3}{3}Tr\ y^3 + b_0 \quad (7)$$

Furthermore, we want the true minimum in the Σ direction to break SU(5) $\rightarrow$ SU(3) $\times$ SU(2) $\times$ U(1) which implies that

$$y_0 = \frac{\Sigma_0}{M} = \frac{\Delta}{M}\begin{pmatrix} 2 & & & & 0 \\ & 2 & & & \\ & & 2 & & \\ & & & -3 & \\ 0 & & & & -3 \end{pmatrix} \tag{8}$$

where Δ is a scale characteristic of Σ (typically M_{GUT}). Now the condition $g(y_0) = 0$ implies

$$15b_1\Delta^2 - 10b_2\Delta^3 + b_0 = 0 \tag{9a}$$

and

$$\left.\frac{\partial g}{\partial y_{ab}}\right|_{y=y_0} = 0 (\text{with the constraint Tr y} = 0)$$

implies

$$b_0 = -5\Delta^3 \tag{9b}$$

$$b_1 = \Delta b_2 \tag{9c}$$

With the choice $b_2 = 1$, we have

$$g(y) = \frac{\Delta M^2}{2} Tr(y^2) + \frac{M^3}{3} Tr(y^3) - 5\Delta^3 \tag{9d}$$

For our case

$$W(x,y) = f(x) + g(y)$$

and

$$V(x,y) = \frac{e^{x^2+y^2}}{M^2}\left[\left(\frac{\partial W}{\partial x} + xW\right)^2 + \left(\frac{\partial W}{\partial y_{ab}} + y_{ab}W\right)^2 - 3W^2\right] \tag{10}$$

assuming x and y to be real.

From this expression, it is straightforward but tedious to compute the derivatives of the potential in the two directions. We only display $\frac{\partial V}{\partial x}$ since the others are messy and not particularly illuminating

$$\begin{aligned}\frac{\partial V}{\partial x} &= 2xV + \frac{2e^{x^2+y^2}}{M^2}[(f' + xW)(f'' + W + xf') \\ &+ Wf'Tr\, y^2 + \Delta M^2 f'Tr\, y^2 \\ &+ M^3 f'Tr\, y^3 - 3Wf']\end{aligned} \tag{11}$$

where primes denote $\frac{\partial}{\partial x}$. Using these expressions, one can determine what the value of the Σ field is when $\phi = 0$ i.e. at the beginning of inflation.

In the Appendix we show that it is impossible to simultaneously satisfy $\frac{\partial V}{\partial x} = \frac{\partial^2 V}{\partial x^2} = \frac{\partial V}{\partial y} = 0, V > 0$ and V $\sim 0(\mu^4)$ at $\phi = 0$ if the Σ field is sitting at its true mimimum i.e. in the 3-2-1 phase. Since all the above conditions are necessary for a successful inflationary model, the Σ field must start its evolution away from the true minimum. If the Σ field is at its true minimum when $\phi = 0$ then it will be less likely that Σ oscillations will be generated as ϕ evolves from $\phi = 0$ to $\phi = \phi_0$.

We now estimate the baryon to entropy ratio in two representative models.

2. Model I

The superpotential for the inflaton field is [12]

$$f(x) = \mu^2 M(x-1)^2 \qquad x \equiv \phi/M \tag{12}$$

where the scale μ is fixed at $(10^{-3} - 10^{-4})M$ by demanding that the model gives the correct order of magnitude of density fluctuations which lead to galaxy formation. [12,13]

This superpotential leads to an absolute minimum at $x = 1$ with zero cosmological constant and unbroken supersymmetry.

The evolution equations for x and y can be solved numerically and the energy stored in the Σ field can be determined. However, this is not particularly illuminating. We find that a more physically transparent strategy is to solve the evolution equations analytically using various physically reasonable approximations. This is the approach we chose in the following analysis.

There are two natural scales in this model: the scale μ associated with the inflation sector $\left(\frac{\mu}{M} \sim 0(10^{-3} - 10^{-4})\right)$ and the scale Δ associated with the Σ sector which has a typical value $\sim 10^{-2}M$ [18]. Thus a reasonable parameter to use is μ/Δ. We will throughout keep only the lowest order terms in μ/Δ.

At $\phi = 0$, we need to determine the value of the Σ field. Assuming that the value at $\phi = 0$ is a small perturbation from the true minimum, we write

$$y = \frac{\Delta}{M}\begin{pmatrix} 2+a\mu/\Delta & & & & \\ & 2+a\mu/\Delta & & & \\ & & 2+a\mu/\Delta & & \\ & & & -3-\frac{3}{2}a\mu/\Delta & \\ & & & & -3-\frac{3}{2}a\mu/\Delta \end{pmatrix} \tag{13}$$

Using the derivatives $\frac{\partial V}{\partial y}$ we can solve for a to get

$$a = \frac{5}{21}\frac{\mu}{M} \sim 10^{-5}$$

which confirms our expectations of keeping only the lowest order terms in μ/Δ.

Next we need to trace the evolution of the ϕ and Σ system in the $\phi - \Sigma$ plane as ϕ evolves from $\phi = 0$ to $\phi = \phi_0 = M$. Once again we need to solve the evolution equations numerically, but we can simplify matters. Since the position of $< y >$ at $\phi = 0$ is not very different from that at $\phi = \phi_0$, it is reasonable to assume that the evolution of ϕ is unaltered.

With these assumptions, we now obtain the position of the Σ field at the end of inflation. The inflationary epoch is characterized by a slow rollover in the ϕ direction and in terms of the potential this implies,[5]

$$V''(\phi) \leq \frac{3}{M^2}|V(\phi)| \tag{14a}$$

$$V'(\phi) \leq \frac{\sqrt{6}}{M}|V(\phi)| \tag{14b}$$

For the potential we consider, the first equation breaks down first at a value

$$x_e \sim 0.2425 \tag{15}$$

Using this value of x_e, we once again solve $\frac{\partial V}{\partial y}$ to get the value of Σ at this point (to lowest order in μ/Δ). Assuming the form of y to be as in (13) we get

$$y(x = x_e) = \frac{\Delta}{M}\begin{pmatrix} 2+1.15\frac{\mu^2}{\Delta M} & & & & \\ & 2+1.15\frac{\mu^2}{\Delta M} & & & \\ & & 2+1.15\frac{\mu^2}{\Delta M} & & \\ & & & -3-1.725\frac{\mu^2}{\Delta M} & \\ & & & & -3-1.725\frac{\mu^2}{\Delta M} \end{pmatrix} \tag{16}$$

The evolution of the ϕ and Σ fields is governed by the evolution equations which are [6]

$$\ddot{x} + 3H\dot{x} + \Gamma_x \dot{x} = -\frac{1}{M^2}\frac{\partial V}{\partial x}$$
$$\ddot{y}_{ab} + 3H\dot{y}_{ab} + \Gamma_y \dot{y}_{ab} = -\frac{1}{M^2}\frac{\partial V}{\partial y_{ab}} \quad (17)$$

where

$$H^2 = \frac{1}{3M^2}[V(\phi,\Sigma) + \frac{1}{2}\dot{\phi}^2 + \frac{1}{2}\dot{\Sigma}^2 + \rho_\gamma] \quad (18)$$

Here Γ_x and Γ_y are the decay rates of the ϕ and Σ fields respectively and ρ_γ is the energy density in radiation. The equation for y can be rewritten as an equation for a using (13)

$$\ddot{a} + 3H\dot{a} + \Gamma_y \dot{a} = -\frac{1}{\mu M}\frac{\partial V}{\partial y} \quad (19)$$

We can get a sensible approximation scheme for these quantitites by comparing the order of magnitude. Since the ϕ field has only gravitational couplings, its decay rate is

$$\Gamma_\phi \sim \frac{m_\phi^3}{M^2} \quad (20)$$

On the other hand, Σ is a gauge nonsinglet and its decay rate is

$$\Gamma_\Sigma \sim \alpha m_\Sigma \sim \alpha\Delta(\text{assuming } m_\Sigma \sim \Delta) \quad (21)$$

where α is the GUT gauge coupling constant.

At the origin in the ϕ direction, the value of the Hubble parameter H is $\sim \frac{\mu^2}{M}$. Assuming $m_\phi \sim \frac{\mu^2}{M}$ [12] and and $\alpha \sim \frac{1}{20}$ [13], we obtain

$$\Gamma_\phi << 3H << \alpha\Delta \quad (22)$$

Furthermore, the time taken for slow rollover, t_e, is given by [12]

$$t_e \sim \frac{M}{\mu^2} >> \Gamma_\Sigma^{-1} \quad (23)$$

The physical picture which emerges from this is as follows: at $t = 0$, the ϕ field is at its origin while the Σ field is displaced from its true minimum at a value given by (13). From $t = 0$ to $t = t_e$, the ϕ field evolves slowly from $\phi = 0$ to $\phi = \phi_e$, giving rise to the de-Sitter expansion of the scale factor. Since this time is much longer than the lifetime of the $\Sigma's$, all the primordial $\Sigma's$ decay and the density of the decay products is exponentially diluted. However, at $t = t_e$, Σ is not at its true minimum but is displaced to a value given by (16).

Taking into account the inequalities given by (22), we can approximately solve the evolution equations for ϕ and Σ. These equations give us essentially the same result as if the Σ field was moving in a pure quadratic potential around the true minimum. Thus for our purposes, we take the motion in the Σ direction to be governed by

$$V = \frac{1}{2}M^2 m_\Sigma^2 (y - y_0)^2$$
$$= \frac{1}{2}M^2\Delta^2 a^2 \frac{\mu}{\Delta} \quad (24)$$

At time $t = t_e$, the value of Σ is given by (16) and the total energy in the Σ direction is at least

$$\rho_\Sigma(t = t_e) \sim \Delta^2\frac{\mu^4}{M^2} \quad (25)$$

The field is oscillating in a pure quadratic potential with a frequency given by its mass. Since this frequency is comparable to the decay rate of Σ, this energy rapidly goes into decay products before redshifting decreases it significantly. On the other hand, the ϕ field has a very long lifetime and it continues to oscillate near $\phi = \phi_0$ for a long time, with its energy redshifting significantly before decay into radiation. So we need to study the evolution of the energies associated with the ϕ and Σ directions from time $t = t_e$ to $t = t_\phi \equiv \Gamma_\phi^{-1}$ and compute the ratio $\frac{n_B}{S}$ at $t = t_\phi$.

To study the evolution, note that the energy associated with the oscillations in the ϕ direction is $0(\mu^4)$ and that in the Σ oscillations is $0(\mu^4\frac{\Delta^2}{M^2})$. Since $\Delta \sim 10^{-2}M$, we can safely ignore the contribution of ρ_Σ to the evolution of the scale factors.

We assume that the dominant mechanism for the production of baryon asymmetry is the decay of color triplet Higgs which is produced in the decay of Σ. This will give us a lower limit on the magnitude of $\frac{n_B}{S}$.

Let n_H be the number density of the Higgs triplets of mass m_H produced by the decay of the $\Sigma's$. Then the energy density ρ_H is given by, since the Higgs' are non relativistic,

$$n_H = \rho_H / m_H \tag{26}$$

Further let a fraction f of the Σ energy before decay go into the triplets and for simplicity the rest into photons. Then

$$\rho_H = f \rho_\Sigma \tag{27}$$

and the reheat temperature is

$$T_{RH}^{(1)} = \left[\frac{30}{\pi^2 g_*} (1 - f) \rho_\Sigma \right]^{1/4} \tag{28}$$

where g_* is the effective relativistic degrees of freedom.

The potential in the ϕ direction is given by

$$V = e^{x^2} \mu^4 [x^6 - 4x^5 + 7x^4 - 4x^3 - x^2 + 1] \tag{29}$$

and near $x = x_0$ by

$$V = \mu^4 e^{x^2} [4(x - x_0)^2 + 12(x - x_0)^3 + \cdots]$$

Thus near $x = x_0$, the dominant term is the quadratic term and the expansion is matter dominated [10]. The energies at $t = t_e$ and $t = t_\phi$ are related by

$$\rho_H(t = t_\phi) = \rho_H(t = t_e) \left[\frac{R(t = t_\phi)}{R(t = t_e)} \right]^{-3} \tag{30}$$

where R is the cosmic scale factor. But

$$\frac{R(t_\phi)}{R(t_e)} = \left[1 + \frac{3}{2} H_{t=t_e} (t_\phi - t_e) \right]^{2/3} \tag{31}$$

where $H_{t=t_e}$ is the Hubble parameter at $t = t_e$.

From (22), (23) we obtain

$$\frac{R(t_\phi)}{R(te)} \sim \left(1 + \frac{3}{2} H_{t=t_e} \Gamma_\phi^{-1} \right)^{2/3} \tag{32}$$

Also from (18) and the fact that $\rho_\phi(t_e) \sim \mu^4$ we get

$$\rho_H(t_\phi) = \frac{4}{3} \rho_H(t_e) \frac{\mu^8}{M^8} \tag{33}$$

Using (33) and $\rho_\Sigma(t_e) \sim \frac{\Delta^2}{M^2} \mu^4$ we obtain the number density of the triplets at the time of ϕ decay as

$$n_H(t_\phi) = \rho_{H/m_H} \sim \frac{f}{m_H} \Delta^2 \frac{\mu^{12}}{M^{10}} \tag{34}$$

Assuming that ϵ_B is the baryon excess produced per triplet decay we obtain the number density of excess baryons as

$$n_B \sim \frac{\epsilon_B f}{m_H} \Delta^2 \frac{\mu^{12}}{M^{10}} \tag{35}$$

From Ref. 12, we know the reheat temperature for this model,

$$T_{RH} \sim \sqrt{M \Gamma_\phi} \sim \mu^3 / M^2 \tag{36}$$

Note that this is the final reheat temperature, produced by the decay of the inflaton. There might be some intermediate reheating associated with the decay of other particles, for eg. $T_{RH}^{(1)}$ associated with the decay of $\Sigma's$. This produces a negligible amount of entropy because the small amount of energy gets redshifted significantly between t_e and t_ϕ. Thus the baryon to entropy ratio at $t = t_\phi$ is given by

$$\frac{n_B}{S} \sim \frac{45}{2\pi^2 g_*} \frac{\epsilon_B f}{m_H} \frac{\Delta^2 \mu^3}{M^4} \tag{37}$$

Using eq. 37, we can estimate the numerical value of $\frac{n_B}{S}$ and compare it to the observed value of $\sim 10^{-10}$. There are however, ambiguities in the values of the parameters entering eq. 37. The values of $\frac{\mu}{M}$ and $\frac{\Delta}{M}$ can be fixed, as already indicated at $10^{-3} - 10^{-4}$ and 10^{-2} respectively [12, 14]. g_* can be assumed to be $0(2 \times 10^2)$ at these scales. ϵ_B, f and m_H are more uncertain and model dependent.

It is known [15], that in supersymmetric models, apart from the usual dimension 6 operators responsible for proton decay, there can also exist dimension 5 operators which could give a disastrously small proton lifetime. If these operators are present, we have a lower bound on the mass of the superpartners of the triplets given by [16].

$$m_{\tilde{H}_3} \geq 10^{16} GeV. \tag{38}$$

However, one can invoke certain symmetries, for example a Peccei Quinn symmetry or a discrete symmetry, which forbid proton decay by dimension 5 operators. In these cases the limit is much smaller. For example Ref. 17 shows that it is possible to reconcile a low mass Higgs triplet with the experimental bounds on proton lifetime. The lower bound is considerably reduced to

$$m_H \geq 2.85 \times 10^{10} \mathrm{GeV} \tag{39}$$

The value of ϵ_B, or the net baryon number produced by the decay of a particle-antiparticle pair is also very model dependent. At tree level, $\epsilon_B = 0$ and $\epsilon_B \neq 0$ comes from loop diagrams. For supersymmetric guts, no "surprising" cancellations occur at one loop level and so $\epsilon_B \leq 0(\alpha/4\pi)$[18]. The quantity f is to be determined by looking at the decay modes of the $\Sigma's$. The $\Sigma's$ can decay into anything lighter–triplet, doublet Higgs, gluons etc. A value of 1/10 is not an unreasonable value for this parameter. Using $\epsilon_B \sim 10^{-3}$ [17], we obtain from (37).

$$\frac{n_B}{S} \sim 10^{-10} 10^{(-9-12)} \frac{M}{m_H} \tag{40}$$

If we use $m_H \sim 10^{+10} GeV$ and $\mu \sim 10^{-3} M$, we obtain a value of $\frac{n_B}{S}$ which almost agrees with that observed. However, if the higher bound on

m_H is taken from models where dimension 5 operators are not suppressed by some symmetry, then this mechanism gives us a much smaller value of $\frac{n_B}{S}$ in disagreement with observations.

3. MODEL II:

Having computed $n_{B/S}$ for this simple model with no supersymmetry breaking, we go on to consider a model with supersymmetry breaking in the inflaton sector.

Consider the inflaton superpotential [19],

$$f(x) = \mu^2 M[\beta + \epsilon + x + \beta x^2 - \frac{1}{12}\beta x^4] \tag{41}$$

where $\beta = -\frac{3}{8}\sqrt{2} - \frac{3}{8}\epsilon + 0(\epsilon^2)$. The minimum is supersymmetry breaking and is at

$$x = \sqrt{2} + \left(\frac{2\sqrt{2}}{3}\epsilon\right)^{1/2} \tag{42}$$

and the gravitino mass is

$$m_{3/2} = \frac{2}{9}\sqrt{3\sqrt{2}e}\frac{\mu}{M}^2 \epsilon^{3/2} \tag{43}$$

In this model, supersymmetry breaking is associated with a non-zero value of ϵ. However, for the first part of our analysis we will assume $\epsilon = 0$ since this does not change our conclusions. We start with a superpotential

$$f(x) = \sqrt{2}\mu^2 M\left[-\frac{3}{8} + \frac{x}{\sqrt{2}} - \frac{3}{8}x^2 + \frac{x^4}{32}\right] \tag{44}$$

Coupling the Σ field to ϕ and carrying out the same analysis as for model I, we obtain the value of Σ at the end of inflation. The slow rollover or the inflationary epoch ends at a time $t = t_e$ when the inequalities in eq. 14 are no longer satisfied. It turns out that the second inequality breaks first when the ϕ field is at ϕ_e given by [19],

$$\phi_e \sim \cdot 71M \tag{45}$$

Using this value of ϕ_e, we solve $\frac{\partial V}{\partial y} = 0$ to obtain

$$y(x = x_e) = \frac{\Delta}{M}$$

$$\begin{pmatrix} 2 + \cdot 41\frac{\mu^2}{\Delta M} & & & & \\ & 2 + \cdot 41\frac{\mu^2}{\Delta M} & & & \\ & & 2 + \cdot 42\frac{\mu^2}{\Delta M} & & \\ & & & -3 - \cdot 62\frac{\mu^2}{\Delta M} & \\ & & & & -3 - \cdot 62\frac{\mu^2}{\Delta M} \end{pmatrix} \tag{46}$$

Once again, as for Model I, we use these initial conditions to approximately solve the evolution equations for Σ and ϕ. Not surprisingly, we find again that the motion in the Σ direction is governed by a pure quadratic potential. At time t_e, the Σ field sits away from its true minimum and has energy $\rho_\Sigma \sim \mu^4\frac{\Delta^2}{M^2}$ which rapidly goes into its decay products. In computing $\frac{n_B}{S}$, we need to trace the evolution of ρ_ϕ and ρ_Σ from t_e to t_ϕ. It is in this part that the difference from Model I comes in.

Recall that for Model I, the potential was predominantly quadratic in the ϕ direction and hence the universe expanded like a matter dominated one. In Model II however, there are two stages of expansion (once again $\rho_\phi >> \rho_\Sigma$ and the evolution is governed by ρ_ϕ). From time t_e to a time $t = t_t \sim 6 \text{ x } 10^{-2}\epsilon^{-1}\frac{M}{\mu^2}$, the ϕ^4 term dominates and the universe expands like radiation dominated [19].

Thus for $t_e \leq t \leq t_t$

$$\frac{R(t)}{R(te)} = [1 + 2H_{t=te}(t - te)]^{1/2} \tag{47}$$

From time t_t to $t_\phi \equiv \Gamma_\phi^{-1}$, the dominant term is quadratic and expansion is matter dominated.

$$t_t \leq t \leq t_\phi$$
$$\frac{R(t)}{R(t_t)} = \left[1 + \frac{3}{2}H_{t=t_t}(t - t_t)\right]^{2/3} \tag{48}$$

Now following the same steps as in Model I with the same notation, we find that

$$n_B = \frac{\epsilon_B \rho_H(t_\phi)}{m_H} = \frac{\epsilon_B f}{m_H} \frac{\Delta^2 \mu}{M^{1/2}} t_t^{1/2} \Gamma_\phi^2 \tag{49}$$

Since the energy density in Σ is much smaller than that in $\phi's$, one can easily check that the reheating temperature is the same as obtained in Ref. 19.

$$T_{RH} \sim \frac{\mu^3 \epsilon^{3/4}}{M^2} \tag{50}$$

Using eq. (49), (50) and $\Gamma_\phi \sim \frac{m^3 \phi}{M^2} \sim \frac{\mu^6 \epsilon^{3/2}}{M^5}$ [19] we obtain

$$\frac{n_B}{S} \sim \frac{\epsilon_B f}{g_*} \frac{\mu^3 \epsilon^{1/4} \Delta^2}{M^4 m_H} \tag{51}$$

From Ref. 19, we have $\mu/m \sim 10^{-3} - 10^{-4}$ and $\epsilon \sim 10^{-7\pm1.5}$. From the discussion for Model I, we know that the values of the other parameters are model–dependent. Taking $g_* \sim 210^2$, $\frac{\Delta}{m} \sim 10^{-2}$, $m_H \sim 10^{10} GeV$, $\epsilon_B \sim 10^{-3}$ and $f \sim 10^{-1}$ we obtain

$$\frac{n_B}{S} \sim 10^{-11} 10^{(-1.5-2)} \tag{52}$$

which is similar to that obtained in Model I apart from a factor of $\epsilon^{1/4}$. In fact the reheating temperature in this model is smaller by $\epsilon^{3/4}$ compared to Model I, and so one expects a larger $\frac{n_B}{S}$. This is not true however, because the inflaton field has a longer lifetime in Model II. Hence the energy in the triplets is redshifted more and the enhancement due to a lower reheating temperature is more than cancelled to give us $\frac{n_B}{S}$ in (52).

The two models we have considered, suffer from the same disease; they both violate the requirement that at high temperatures, a sufficient amount of energy is stored in the scalar field ϕ to give enough inflation – the thermal constraint. In other words, inflaton must start its evolution far away from its global minimum, slowly roll down and eventually settle in its global minimum. This is not surprising however, because of a general result given in Ref. 20. In a hidden sector with a single field and a flat Kähler metric, the temperature corrections do not stabilize the field at the origin.

The solution to this problem suggested in Ref. 12 and Ref. 19 is to allow for direct couplings between ϕ and another field ψ. For our case, we have until now, only considered the situation with the GUT sector and the inflaton sector are separate i.e. only coupled gravitationally. If direct couplings between the two sectors are allowed, the situation in the two models is somewhat different.

In Model I, the inflaton sector does not break supersymmetry and hence direct coupling of ϕ and Σ, will not have any danger of changing the breaking scale. In Model II however, the inflaton sector is also responsible for the breaking of supersymmetry (with $\epsilon \neq 0$). In this case we need to be careful because there is a danger that the supersymmetry breaking scale will be pushed up to m_{GUT} since the $\Sigma's$ now couple directly to the ϕ.

Thus in both cases we see that if we include direct coupling of ϕ and Σ, then the thermal constraint can be satisfied. Furthermore, it is possible that with direct couplings, the value of $\frac{n_B}{S}$ will improve because more energy can be transferred now from the inflaton to the Σ. However, with the direct couplings, the analysis becomes very complicated. This is because firstly, one has to be careful that gauge radiative corrections do not spoil the nice features of the inflationary potential. Secondly, both the fields are now responsible for inflation and reheating [for an exception see Ref. 21]. We do not carry out this analysis since it is beyond the scope of the present work.

To conclude, we have studied a mechanism for the generation of baryon asymmetry which involves the use of the couplings of heavy fields with the hidden sector. This mechanism seems to be a very general one since in any model with an inflationary sector and a GUT sector which has heavy fields, there will exist the possibility of the transfer of energy from the inflaton to the heavy fields. We have obtained the value of $\frac{n_B}{S}$ in the case of two inflaton superpotentials (one with and one without supersymmetry breaking). The numerical value of $\frac{n_B}{S}$ however is seen to be dependent upon parameters which are model dependent. We saw that if we use the bound on m_H from supersymmetric GUTs where some symmetry prohibits dimensions 5 operators for baryon decay, then we obtain a value of $\frac{n_B}{S}$ which is almost in agreement with the observations. In both models we saw that the thermal constraint is violated unless one includes direct couplings between the inflaton and the Σ fields. We conclude then that there exists another possible mechanism for baryon number generation within the framework of supersymmetric inflationary cosmologies.

APPENDIX

In this appendix we show that under very general conditions it is impossible for the Σ field to sit at its absolute minimum when $\phi = 0$. The notation is that of the text. Let f(x) and g(y) be the superpotentials in the 2 sectors. Let

$$f(x) = \mu^2 M f_1(x) \qquad (A1)$$

$$g(y) = \Delta^3 g_1(y) \qquad (A2)$$

where $f_1(x)$ and $g_1(y)$ are dimensionless. Further assume that there is no direct coupling between the 2 fields. Then

$$W(x,y) = f(x) + g(y) \qquad (A3)$$

Now we impose the following conditions: at $x = x_0$, $y = y_0$ (the true minimum) we must have unbroken supersymmetry and zero cosmological constant. This implies

$$f_1(x_0) = f_1'(x_0) = 0 \qquad (A4)$$

$$g_1(y_0) = g_1'(y_0) = 0 \qquad (A5)$$

Assume that when $x = 0$, $y = y_0$ i.e. the field y starts off at its absolute minimum. Then demanding that the potential be flat means

$$\frac{\partial V}{\partial x} = \frac{\partial^2 V}{\partial x^2} = \frac{\partial V}{\partial y} = 0 \text{ at } x = 0\,,\ y = y_0 \qquad (A6)$$

These conditions imply

$$\mu^4 M^2 \left\{ f_1'[f_1'' + (y_0^2 - 2) f_1] \right\} = 0 \qquad (A7)$$

$$\mu^4 M^2 \left\{ f_1^2 (y_0^2 - 2) + f''^2_1 + f_1' f_1'' + y_0^2 f'^2_1 + (y_0^2 - 1) f_1 f_1'' \right\} = 0 \qquad (A8)$$

$$\mu^4 M^2 y_0 \left[f_1'^2 + (y_0^2 - 2) f_1^2 + f_1 g_1'' \frac{\Delta^3}{\mu^2 M} \right] = 0 \qquad (A9)$$

Futhermore at $x = 0$, $y = y_0$

$$V(0, y_0) = \frac{\mu^4 e^{y_0^2}}{M^2} [f'^2 + (y_0^2 - 3) f_1^2] \qquad (A10)$$

Using (A9) gives us

$$\begin{aligned} V(0, y_0) &= \mu^4 \frac{e^{y_0^2}}{M^2} \left[-f_1^2 - f_1 g_1'' \frac{\Delta^3}{\mu^2 M} \right] \\ &= -\frac{\mu^4 e^{y_0^2}}{M^2} f_1 \left[f_1 + g_1'' \frac{\Delta^3}{\mu^2 M} \right] \end{aligned} \qquad (A11)$$

But $g_1''(y_0) \sim 0(\frac{M^2}{\Delta^2})$ since $y_0 \sim 0(\frac{\Delta}{M})$ for the example in text which is quite general. Then (A11) immediately tells us that

$$V(0, y_0^2) \sim 0(\Delta^2 m^2).$$

This is unacceptable because we know that the potential at $\phi = 0$ must scale like μ^4 with $\mu \sim 0(10^{-4})$ to give us the correct density fluctuations! If the field Σ at $\phi = 0$ sits at its absolute minimum then the scale μ drops out of the potential.

Thus we assume that the field Σ starts at some other value at $\phi = 0$, i.e. we solve for $\frac{\partial V}{\partial y} = 0$ at $\phi = 0$ as in the text.

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