Affleck-Dine baryogenesis after thermal inflation

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We argue that an extension of the minimal supersymmetric standard model (MSSM) that gives rise to viable thermal inflation, and so does not suffer from a Polonyi or moduli problem, should contain right-handed neutrinos which acquire their masses due to the vacuum expectation value of the flaton that drives thermal inflation. This strongly disfavors SO(10) grand unified theories. The μ term of the MSSM should also arise due to the VEV of the flaton. With the extra assumption that $m_L^2 - m_{H_u}^2 < 0$, but, of course, $m_L^2 - m_{H_u}^2 + |\mu|^2 > 0$, we show that a complicated Affleck-Dine-type baryogenesis employing an LH_u D-flat direction can naturally generate the baryon asymmetry of the Universe. [S0556-2821(96)00822-3]

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I. INTRODUCTION

Thermal inflation [1–3] provides the most compelling solution to the moduli (Polonyi) problem [4–6]. However, for a theory of the early Universe to be viable it must be capable of producing a baryon asymmetry [7]

$$\frac{n_B}{s} \sim 3 \times 10^{-11} \tag{1}$$

by the time of nucleosynthesis. Thermal inflation probably dilutes any preexisting baryon asymmetry to negligible amounts, and the final reheat temperature after thermal inflation ($T_f \sim$ few GeV) is probably too low even for electroweak baryogenesis. Thus, if thermal inflation really is the solution of the moduli problem, then it is likely also to be responsible for baryogenesis.

In Sec. II we explain why the flaton that gives rise to thermal inflation probably also generates the masses of right-handed neutrinos as well as the μ term of the minimal supersymmetric standard model (MSSM). We also note the various ways in which a potential domain wall problem can be avoided. In Sec. III we describe how a somewhat complicated Affleck-Dine-type mechanism can naturally generate the required baryon asymmetry after thermal inflation. In Sec. IV we give our conclusions.

II. THERMAL INFLATION, RIGHT-HANDED NEUTRINOS, AND THE μ TERM

A. Thermal inflation and right-handed neutrinos

The superpotential of the MSSM is [8][1]

$$W_{\mathrm{MSSM}} = \lambda_t Q H_u t + \lambda_b Q H_d b + \lambda_\tau L H_d \tau + \mu_H H_u H_d . \tag{2}$$

Thermal inflation [1] requires that there is, in addition, at least one flaton ϕ with vacuum expectation value $|\phi| = M$ in the range 10^{10} GeV$\lesssim M \lesssim 10^{12}$ GeV, the lower bound coming from the requirement that thermal inflation sufficiently dilutes the moduli, and the upper bound from the requirement that the final reheat temperature after thermal inflation T_f be high enough to thermalize the lightest supersymmetric particles (LSP's) produced in the flaton's decay and so avoid an excess of LSP's.

Also, in order for ϕ to be held sufficiently strongly[2] at $\phi = 0$ by the finite temperature during thermal inflation, ϕ must have unsuppressed couplings to at least one other field, say ψ, that is light when $\phi = 0$. We, therefore, either require a term $\lambda_\phi \phi \psi^2/2$ with $|\lambda_\phi| \sim 1$ in the superpotential, or require ϕ to spontaneously break a continuous gauge symmetry with gauge coupling $g_\phi \sim 1$ (with ψ being the gauge field in this case). One reason to prefer the Yukawa coupling over the gauge coupling is that the renormalization group [8] effect of the Yukawa coupling would be to drive the soft supersymmetry-breaking mass squared of ϕ negative at low energies as is required for a flaton, while the gauge coupling would have the opposite effect. Another reason is that the gauge symmetry has no independent motivation, while, as we shall see, the Yukawa coupling is very well motivated. We will, therefore, focus on the case of the Yukawa coupling.

After ϕ acquires its vacuum expectation value M, ψ will acquire a mass $|\lambda_\phi| M \sim 10^{10}$ GeV–10^{12} GeV and so is not a MSSM field. In order for ϕ to be coupled to the thermal bath, which is presumably composed of MSSM fields, we, therefore, require ψ to couple to the MSSM. We, therefore,[3]

[1]Here and throughout most of this paper, all indices (the usual gauge and generation indices, as well as any singlet indices) have been suppressed. We will for the most part be focusing on the third generation as is suggested by our notation.

[2]Note the strong dependence on T_C in Eq. (35) of Ref. [1].

[3]Assuming ψ is not charged with respect to the MSSM continuous gauge symmetries as this would, in general, destroy SUSY GUT gauge coupling unification. However, if ψ is a complete SU(5) multiplet the unification of the gauge couplings will be unaffected. For example, one could have $\lambda_\phi \phi \psi^2 = \lambda_\phi \phi \bar{5} 5$. Alternatively, appropriate choices of representations could shift the unification scale to the string scale [9].

require at least one[4] of the terms $LH_u\psi$ or $\psi H_u H_d$ in the superpotential, as these are the only possible renormalizable couplings of a singlet to the MSSM.

The former is the standard coupling $\lambda_\nu L H_u \nu$ of a right-handed neutrino $\psi=\nu$ to the MSSM, and furthermore, ν automatically acquires a mass $M_\nu = |\lambda_\phi| M \sim 10^{10}$ GeV–10^{12} GeV in the right range for the seesaw mechanism [10] to generate a left-handed τ neutrino mass

$$m_{\nu_L} = \frac{m_D^2}{M_\nu} = \frac{|\lambda_\nu|^2 (174\ \text{GeV})^2 \sin^2\beta}{|\lambda_\phi| M} \tag{3}$$

$$\sim 5\ \text{eV}\left(\frac{3\times 10^{12}\ \text{GeV}}{M/|\lambda_\nu|^2}\right)\left(\frac{1}{|\lambda_\phi|}\right)\left(\frac{\sin^2\beta}{0.5}\right) \tag{4}$$

suitable for the mixed dark matter scenario for the formation of the large-scale structure of the Universe [11].

The latter coupling $\psi H_u H_d$ does not have any good independent motivation, apart from the fact that it is possible. We will, therefore, focus on the former case. Later though, we shall make use of this possibility for coupling to the MSSM in a different context.

B. The final reheat temperature after thermal inflation and the μ term

Now that we have a more precise picture of the couplings of the flaton,

$$W_{\text{so far}} = W_{\text{MSSM}} + \lambda_\nu L H_u \nu + \frac{1}{2}\lambda_\phi \phi \nu^2, \tag{5}$$

we can hope to make a more precise estimate of the decay rate of the flaton and so the final reheat temperature after thermal inflation T_f.

We first note that the effective superpotential coupling,

$$W_{\text{seesaw}} = -\frac{(\lambda_\nu L H_u)^2}{2\lambda_\phi \phi}, \tag{6}$$

obtained by integrating out ν, i.e., eliminating ν via the constraint $\partial W/\partial \nu = 0$, will give a decay rate of order $\Gamma \sim m^5/M^4$ which is negligible.

The ϕ dependence of the low energy renormalized coupling constants will give larger decay rates. To estimate these we first need to know the contributions of the right-handed neutrinos to the renormalization group equations. They are [12]

$$16\pi^2 \frac{d}{dt}\lambda_t = |\lambda_\nu|^2 \lambda_t + \cdots, \tag{7}$$

$$16\pi^2 \frac{d}{dt}\lambda_\tau = |\lambda_\nu|^2 \lambda_\tau + \cdots, \tag{8}$$

$$16\pi^2 \frac{d}{dt}\mu_H = |\lambda_\nu|^2 \mu_H + \cdots, \tag{9}$$

$$16\pi^2 \frac{d}{dt} m_L^2 = 2|\lambda_\nu|^2 (m_L^2 + m_\nu^2 + m_{H_u}^2 + |A_{LH_u\nu}|^2) + \cdots, \tag{10}$$

$$16\pi^2 \frac{d}{dt} m_{H_u}^2 = 2|\lambda_\nu|^2 (m_L^2 + m_\nu^2 + m_{H_u}^2 + |A_{LH_u\nu}|^2) + \cdots, \tag{11}$$

$$16\pi^2 \frac{d}{dt} A_{QH_u t} = 2|\lambda_\nu|^2 A_{LH_u\nu} + \cdots, \tag{12}$$

$$16\pi^2 \frac{d}{dt} A_{LH_d\tau} = 2|\lambda_\nu|^2 A_{LH_u\nu} + \cdots, \tag{13}$$

$$16\pi^2 \frac{d}{dt} A_{\mu_H H_u H_d} = 2|\lambda_\nu|^2 A_{LH_u\nu} + \cdots, \tag{14}$$

where t is the logarithm of the renormalization scale, the m's are the soft supersymmetry-breaking masses, the A's are the soft supersymmetry-breaking parameters in the terms in the scalar potential of the form AW+c.c., the m's and the magnitudes of the A's are of the order of the soft supersymmetry-breaking scale $m_s \sim 10^2$–10^3 GeV, and the ellipsis stands for other terms independent of the right-handed neutrinos. $|\phi|$ sets the threshold for the right-handed neutrinos, and so writing $|\phi| = M + \delta\phi_r/\sqrt{2}$ we get the effective couplings

$$W_{\text{eff}} = \frac{|\lambda_\nu|^2}{16\sqrt{2}\pi^2 M}(\lambda_t Q H_u t + \lambda_\tau L H_d \tau + \mu_H H_u H_d)\,\delta\phi_r \tag{15}$$

and

$$V_{\text{soft eff}} = \frac{|\lambda_\nu|^2}{8\sqrt{2}\pi^2 M}(m_L^2 + m_\nu^2 + m_{H_u}^2 + |A_{LH_u\nu}|^2) \times (|L|^2 + |H_u|^2)\,\delta\phi_r + \frac{|\lambda_\nu|^2}{8\sqrt{2}\pi^2 M} \times A_{LH_u\nu}(\lambda_t Q H_u t + \lambda_\tau L H_d \tau + \mu_H H_u H_d)\,\delta\phi_r. \tag{16}$$

From these couplings we estimate the total decay rate to be

$$\Gamma \sim \frac{|\lambda_\nu|^4 m_s^3}{10^4 M^2}. \tag{17}$$

This would give a final reheat temperature after thermal inflation of

[4] If we have both then we need two different ψ's, one charged and the other neutral under R parity.

$$T_f \simeq g_*^{-1/4} \Gamma^{1/2} M_{\rm Pl}^{1/2}, \tag{18}$$

$$\sim 10\ {\rm MeV} \left(\frac{3\times 10^{12}\ {\rm GeV}}{M/|\lambda_\nu|^2} \right) \left(\frac{m_{\rm s}}{300\ {\rm GeV}} \right)^{3/2}, \tag{19}$$

where the first pair of brackets is constrained to be of order one, or perhaps less, by Eq. (4). However, in order not to overproduce LSP's we require [13]

$$T_f \gtrsim 1\ {\rm GeV} \left(\frac{m_{\rm s}}{300\ {\rm GeV}} \right)^{1.5\text{–}2}. \tag{20}$$

Thus, our model seems to be in trouble unless we can add some extra coupling that gives a stronger decay rate. The only possibility is to couple ϕ to $H_u H_d$ in the superpotential. A term $\phi H_u H_d$ would require a very small coupling constant to avoid generating too large a μ term, but a term[5]

$$W_{\rm decay} = \frac{\lambda_\mu \phi^2 H_u H_d}{M_{\rm Pl}} \tag{21}$$

is not only allowed but could naturally generate a μ term

$$\mu_\phi = \frac{\lambda_\mu \langle \phi \rangle_{\rm vac}^2}{M_{\rm Pl}} \tag{22}$$

of the required size [14]. For example, for $M = 10^{11}$ GeV and $|\lambda_\mu| = 0.1$ we get $|\mu_\phi| = 400$ GeV. From now on we will assume that the μ term is generated in this way so that $\mu_H = 0$, or at least $|\mu_H| \lesssim |\mu_\phi|$, and $|\mu_\phi| \sim m_s$.

Writing $\phi = \mathcal{M} + \delta\phi$, where $|\mathcal{M}| = M$ and $\mu_\phi = \lambda_\mu \mathcal{M}^2 / M_{\rm Pl}$, we get the relatively unsuppressed couplings

$$\begin{aligned} \mathcal{L}_{\rm decay} = {} & 2\mu_\phi \widetilde{H}_u \widetilde{H}_d \frac{\delta\phi}{\mathcal{M}} - 2|\mu_\phi|^2 (|H_u|^2 + |H_d|^2) \frac{\delta\phi}{\mathcal{M}} \\ & - 2\mu_\phi (\overline{\lambda_t Q t} H_d + \overline{\lambda_b Q b} H_u + \overline{\lambda_\tau L \tau} H_u \\ & + A_\mu H_u H_d) \frac{\delta\phi}{\mathcal{M}} + \text{c.c.}, \end{aligned} \tag{23}$$

where a tilde denotes the fermionic component of the superfield, a bar denotes the Hermitian conjugate, and A_μ is the soft supersymmetry-breaking parameter in $V_{\rm soft} = A_\mu \mu_\phi H_u H_d + \text{c.c.}$ and so $|A_\mu| \sim m_s$. The decay rate is then estimated to be

$$\Gamma \sim \frac{m_s^3}{10^2 M^2}, \tag{24}$$

where we have roughly assumed $m_\phi \sim |\mu_\phi| \sim |A_\mu| \sim m_s$. We, therefore, get a reheat temperature

$$T_f \sim (1\text{–}10\ {\rm GeV}) \left(\frac{10^{11}\ {\rm GeV}}{M} \right) \left(\frac{m_s}{300\ {\rm GeV}} \right)^{3/2}, \tag{25}$$

which is sufficiently high.

We thus expect the flaton to have the following couplings to the MSSM[6]:

$$W_{\rm couplings} = \lambda_\nu L H_u \nu + \frac{1}{2} \lambda_\phi \phi \nu^2 + \frac{\lambda_\mu \phi^2 H_u H_d}{M_{\rm Pl}}. \tag{26}$$

C. The flaton potential and domain walls

In this section we consider the self-couplings of the flaton. The flaton (or, in the case of a multicomponent flaton, at least one component of the flaton) should have a negative soft supersymmetry-breaking mass squared $-m_\phi^2 |\phi|^2$ to drive it away from $\phi = 0$ after thermal inflation. It will also need a term to stabilize its potential at its vacuum expectation value (VEV) $|\phi| = M \sim 10^{10}$–10^{12} GeV. The simplest possibility is[7]

$$W_{\rm VEV} = \frac{\lambda_M \phi^4}{4 M_{\rm Pl}} \tag{27}$$

and this is what we shall assume. In the case of a multicomponent flaton this could be interpreted as, for example, $\lambda_M \phi_1^4 / 4 M_{\rm Pl}$, $\lambda_M \phi_1 \phi_2^3 / M_{\rm Pl}$, or a sum of such terms. See Ref. [15] for an explicit multicomponent example. Here, for simplicity, we will focus on the case of a single-component flaton, though it should be borne in mind that a multicomponent flaton might be preferable from the model-building point of view.

We then get the following scalar potential:

$$V(\phi) = V_0 - m_\phi^2 |\phi|^2 + \left(\frac{A_M \lambda_M \phi^4}{M_{\rm Pl}} + \text{c.c.} \right) + \frac{|\lambda_M|^2 |\phi|^6}{M_{\rm Pl}^2}, \tag{28}$$

where $m_\phi \sim |A_M| \sim m_s$. This potential has four degenerate minima with $|\phi| = M$, where

$$M^2 = \frac{2 m_\phi M_{\rm Pl}}{3|\lambda_M|} \left(\frac{|A_M|}{m_\phi} + \sqrt{\frac{|A_M|^2}{m_\phi^2} + \frac{3}{4}} \right). \tag{29}$$

For example, for $m_\phi = |A_M| = 300$ GeV and $|\lambda_M| = 0.1$, we get $M = 10^{11}$ GeV. The eigenvalues of the mass squared matrix at the minima are

$$m_{\delta\phi_\theta}^2 = \frac{16 m_\phi |A_M|}{3} \left(\frac{|A_M|}{m_\phi} + \sqrt{\frac{|A_M|^2}{m_\phi^2} + \frac{3}{4}} \right) \tag{30}$$

and

[5] A term $\lambda_\mu \phi^3 H_u H_d / M_{\rm GUT}^2$ would be an alternative if $M \gtrsim 3\times 10^{11}$ GeV. We assume the displayed case for simplicity.

[6] As mentioned before, $\lambda_\mu \phi^3 H_u H_d / M_{\rm GUT}^2$ is a possible alternative to $\lambda_\mu \phi^2 H_u H_d / M_{\rm Pl}$.

[7] $W_{\rm VEV} = \lambda_M \phi^5 / M_{\rm GUT}^2$ would be an alternative if $M \gtrsim 3\times 10^{11}$ GeV. Again, we assume the displayed case for simplicity. One might even be able to use the renormalization group running of m_ϕ to stabilize the potential but then one would not automatically get a value for M in the correct range.

$$m_{\delta\phi_r}^2 = 4m_\phi^2 + m_{\delta\phi_\theta}^2 = 4m_\phi^2\left[1 + \frac{4|A_M|}{3m_\phi}\times\left(\frac{|A_M|}{m_\phi} + \sqrt{\frac{|A_M|^2}{m_\phi^2} + \frac{3}{4}}\right)\right] \tag{31}$$

and the flatino mass squared is

$$m_{\tilde{\delta\phi}}^2 = 3m_\phi^2 + \frac{3}{2}m_{\delta\phi_\theta}^2 = 4m_\phi^2\left(\frac{|A_M|}{m_\phi} + \sqrt{\frac{|A_M|^2}{m_\phi^2} + \frac{3}{4}}\right)^2. \tag{32}$$

Requiring zero cosmological constant at the minima gives

$$V_0 = \frac{2}{3}m_\phi^2 M^2\left[1 + \frac{2|A_M|}{3m_\phi}\left(\frac{|A_M|}{m_\phi} + \sqrt{\frac{|A_M|^2}{m_\phi^2} + \frac{3}{4}}\right)\right]. \tag{33}$$

With four degenerate minima we clearly have to worry about a potential domain wall problem. The simplest way to eliminate the domain walls is to add a small term which breaks the degeneracy of the vacua, the difference in pressure exerted on the walls causing the domains with greater vacuum energy to collapse. They collapse before the domain walls come to dominate the energy density if the difference in vacuum energies ϵ satisfies $\epsilon \gtrsim \sigma^2/M_{\rm Pl}^2$, where σ is the energy per unit area of the domain walls [16]. For flaton domain walls $\sigma \sim m_s M^2$ and so we require

$$\epsilon \gtrsim \frac{m_s^2 M^4}{M_{\rm Pl}^2}. \tag{34}$$

Therefore, a term in the superpotential of the form

$$W_{\rm walls} \gtrsim \frac{m_s M^{4-n}\phi^n}{M_{\rm Pl}^2} \tag{35}$$

with n odd would be sufficient to eliminate the domain walls. A term with $n = 2\,\mathrm{mod}\,4$, for example $W_{\rm walls} \sim \phi^6/M_{\rm Pl}^3$ would reduce the Z_4 domain walls to Z_2 domain walls. Note that $W_{\rm walls}$ can be extremely small, and hence have a negligible effect on the dynamics to be discussed in the next section, but still solve the domain wall problem.

Another way to avoid a domain wall problem is to gauge the discrete symmetry so that there is really only one vacuum. However, nontrivial anomaly cancellation conditions must be satisfied [17]. In the case of a single-component flaton with the superpotential of Eq. (36), and no extra light SU(3) multiplets, the mixed discrete-SU(3) anomaly cancellation condition requires ϕ^2 to be neutral under any unbroken[8] anomaly free discrete gauge symmetry. We, therefore, cannot use a discrete gauge symmetry to remove the Z_4 domain walls. However, the anomaly free Z_4 subgroup of $\mathrm{U}(1)_{B-L+4Y}$, under which ϕ has charge 2, can be used to gauge away the Z_2 domain walls left by $W_{\rm walls} \sim \phi^6/M_{\rm Pl}^3$ above. Furthermore, this symmetry is broken down to the standard matter parity of the MSSM (which is equivalent to the R parity of the MSSM) by the vacuum expectation value of ϕ.

In the case of a multicomponent flaton the discrete symmetry can be extended to a continuous symmetry which may or may not be gauged. In the case of a continuous global symmetry, for example, the Peccei-Quinn symmetry of Ref. [15], the Goldstone bosons may prove troublesome [18]. One also has more freedom to satisfy the anomaly cancellation conditions in the case of a multicomponent flaton.

[8] R symmetries are broken down to Z_2 by hidden sector supersymmetry breaking.

D. Summary

An extension of the MSSM that gives rise to viable thermal inflation, and so does not suffer from a moduli problem, should have the following terms in its superpotential:

$$W_{\rm ti} = \lambda_t Q H_u t + \lambda_b Q H_d b + \lambda_\tau L H_d \tau + \lambda_\nu L H_u \nu + \frac{1}{2}\lambda_\phi \phi \nu^2 + \frac{\lambda_\mu \phi^2 H_u H_d}{M_{\rm Pl}} + \frac{\lambda_M \phi^4}{4M_{\rm Pl}}. \tag{36}$$

III. LH_u AFFLECK-DINE BARYOGENESIS AFTER THERMAL INFLATION

A. Outline

To orientate the reader we will first sketch the basic idea we have in mind before plunging into the details.

The D-flat direction parametrized by LH_u provides an ideal Affleck-Dine field [19,20,5]. In order for it to behave as an Affleck-Dine field we must first get it away from zero. We, therefore, require its mass squared at the end of thermal inflation, $(m_L^2 - m_{H_u}^2 + |\mu_H|^2)/2$, to be negative. It is simplest to assume that it rolls away from zero before ϕ does. However, when $\phi=0$ the right-handed neutrinos are light, and so LH_u is not F flat (it has a quartic term in its potential coming from the superpotential). LH_u will thus be stabilized at a modest value.

Next, ϕ will roll away from zero. The right-handed neutrinos become heavy and so can be integrated out leaving the effective seesaw coupling $W_{\rm seesaw}$ given in Eq. (6). This is now the term that stabilizes the LH_u direction and we see that it gets smaller, and so the LH_u direction gets flatter, as ϕ gets larger. Thus, as ϕ rolls away from zero, LH_u will roll further away from zero. Furthermore, the soft supersymmetry-breaking term derived from $W_{\rm seesaw}$ will correlate the phases of ϕ and LH_u.

When ϕ becomes sufficiently large, it will start to feel the basin of attraction of one of the minima of its potential, and so will start curving in towards that minimum, i.e., its phase will be roughly determined modulo $\pi/2$. The phase of LH_u will then be roughly determined modulo $\pi/4$ by the soft supersymmetry-breaking term derived from $W_{\rm seesaw}$.

When $|\phi|$ becomes of order M, a cross term from the supersymmetric part of the potential becomes significant and changes the correlation between the phases of ϕ and LH_u, and so gives the phase of LH_u a kick. The direction of the kick is determined by the parameters in the Lagrangian (this is our CP violation) and so gives a nonzero net contribution when averaged over different spatial locations, unlike the rest of the angular momentum that is flying around. Further-

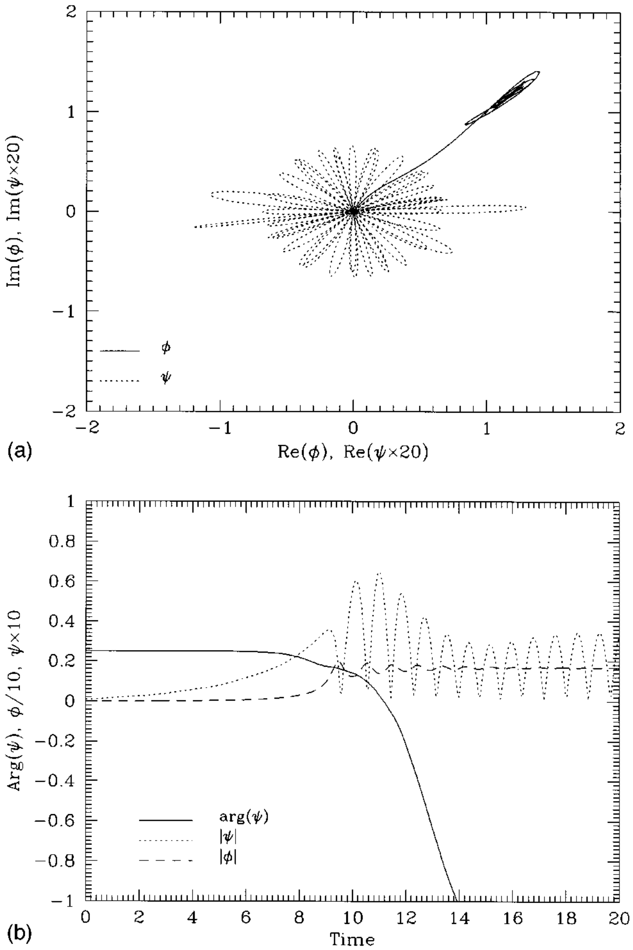

FIG. 1. Numerical simulation to illustrate the dynamics of ϕ and $LH_u=\psi^2/2$ after thermal inflation. The potential of Eq. (48) was used with the parameters $a_\phi=a_\psi=\alpha_\nu=\beta_M=1$, $|\alpha_M|=\beta_\mu=2$, $\gamma=10^{-3}$, and $\arg\alpha_M=3.1$. The initial conditions were $|\phi|=|\psi|=10^{-3}$, $\arg\phi=1$, and $\arg\psi=0.25$. A friction term $\Gamma\dot{\phi}$ with $\Gamma=0.75$ was added to the equation of motion of ϕ to crudely simulate the effects of parametric resonance. A friction term was not added for ψ because it would obscure the total lepton number generated which in reality is contained in both the homogeneous ψ field and its decay products (such as the inhomogeneous ψ modes produced by parametric resonance).

more, as this is happening $W_{\rm decay}$ [see Eq. (21)] starts to give a significant contribution to the mass of H_u and hence LH_u. For $m_L^2-m_{H_u}^2+|\mu_H+\mu_\phi|^2>0$ this gives the LH_u direction an overall positive mass squared (as it must because LH_u has a positive mass squared in the true vacuum), and so sends LH_u spiralling back in towards zero.

The effective friction on the motion of ϕ and LH_u coming from the Hubble expansion is negligible. However, the effective mass squareds of both ϕ and LH_u have been changing sign during the above dynamics and so one would expect them both to decay via broad parametric resonance [21]. This will lead to approximately critical damping, and so it seems reasonable to expect that both ϕ and LH_u will be trapped near their vacuum expectation values essentially immediately after the dynamics described above has occurred. Once they are trapped, parametric resonance becomes less efficient because the mass squareds are now always positive. LH_u's potential near $LH_u=0$ conserves angular momentum, or in other words, lepton number, and so LH_u's newly acquired lepton number is conserved. The dynamics outlined above is illustrated in Fig. 1.

The decay of the LH_u Affleck-Dine condensate will generate enough partial reheating to restore the electroweak symmetry, and so its lepton number can be converted to baryon number by the usual electroweak effects [22]. Note that the energy density is still dominated by the flaton and the reheating in the Affleck-Dine sector has a negligible effect on the now decoupled flaton. Finally, after the temperature has dropped to a few GeV, the flaton decay will complete, releasing substantial entropy.

B. Estimating the baryon asymmetry

Our basic model is

$$W_{\rm ti}=\lambda_t QH_u t+\lambda_b QH_d b+\lambda_\tau LH_d\tau+\lambda_\nu LH_u\nu+\frac{1}{2}\lambda_\phi\phi\nu^2+\frac{\lambda_\mu\phi^2H_uH_d}{M_{\rm Pl}}+\frac{\lambda_M\phi^4}{4M_{\rm Pl}}. \tag{37}$$

The squark fields have no linear terms in their potential and have positive mass squareds. They will, therefore, be held at zero apart from thermal fluctuations, and so can be ignored apart from their contribution to the finite temperature effective potential. The zero temperature potential for the other fields is

$$\begin{aligned}V=&|\lambda_\tau LH_d|^2+|\lambda_\nu LH_u+\lambda_\phi\phi\nu|^2+|\lambda_\tau H_d\tau+\lambda_\nu H_u\nu|^2\\&+\left|\lambda_\nu L\nu+\frac{\lambda_\mu\phi^2H_d}{M_{\rm Pl}}\right|^2+\left|\lambda_\tau L\tau+\frac{\lambda_\mu\phi^2H_u}{M_{\rm Pl}}\right|^2\\&+\left|\frac{1}{2}\lambda_\phi\nu^2+2\frac{\lambda_\mu\phi H_uH_d}{M_{\rm Pl}}+\frac{\lambda_M\phi^3}{M_{\rm Pl}}\right|^2+D\text{ terms}\\&+\Bigg(A_\tau\lambda_\tau LH_d\tau+A_\nu\lambda_\nu LH_u\nu+A_\phi\lambda_\phi\phi\nu^2\\&+\frac{A_\mu\lambda_\mu\phi^2H_uH_d}{M_{\rm Pl}}+\frac{A_M\lambda_M\phi^4}{M_{\rm Pl}}+\text{c.c.}\Bigg)\\&+m_\tau^2|\tau|^2+m_\nu^2|\nu|^2+m_L^2|L|^2\\&-m_{H_u}^2|H_u|^2+m_{H_d}^2|H_d|^2-m_\phi^2|\phi|^2,\end{aligned} \tag{38}$$

where the m's and the magnitudes of the A's are of order m_s. We assume

$$m_{LH_u}^2\big|_{\phi=0}=\frac{1}{2}(m_L^2-m_{H_u}^2)<0 \tag{39}$$

so that the D-flat direction parametrized by LH_u is also unstable, in addition to the flaton ϕ. Note that after ϕ acquires its vacuum expectation value $M\sim\sqrt{m_sM_{\rm Pl}/|\lambda_M|}$, it will give an extra contribution $|\lambda_\mu|^2M^4/M_{\rm Pl}^2$ to H_u's mass squared. This will be of order m_s^2 if $|\lambda_\mu|\sim|\lambda_M|$. We assume

$$m_{LH_u}^2\big|_{|\phi|=M}=\frac{1}{2}\left(m_L^2-m_{H_u}^2+\frac{|\lambda_\mu|^2M^4}{M_{\rm Pl}^2}\right)>0 \tag{40}$$

so that the LH_u direction is stable in the true vacuum.

A rigorous study of the dynamics of this model is beyond the scope of this paper. Instead, we will make some simplifying assumptions in order to illustrate how the Affleck-Dine mechanism might be implemented after thermal inflation and to crudely estimate the resultant baryon asymmetry.

We assume that all fields are initially held at zero by the finite temperature during thermal inflation. We assume that the LH_u direction rolls away from zero first. It will be quickly stabilized by the term $|\lambda_\nu LH_u|^2$ at a value $|LH_u| \sim m_{LH_u}^2/|\lambda_\nu|^2$. The term $A_\nu \lambda_\nu LH_u \nu + \text{c.c.}$ then causes ν to roll away from zero. Then, the term $\overline{\lambda_\nu LH_u} \lambda_\phi \nu \phi + \text{c.c.}$ causes ϕ to roll away from zero[9] in the direction

$$\arg \phi \simeq \arg[\overline{\lambda_\phi} A_\nu (\lambda_\nu LH_u)^2]. \tag{41}$$

For simplicity, we will assume that τ and H_d remain at zero, or at least that any expectation values they acquire can be neglected. With this assumption, once ϕ and LH_u escape beyond the reach of the temperature, their dynamics will be governed by the zero temperature potential

$$\begin{aligned} V = {} & |\lambda_\nu LH_u + \lambda_\phi \phi \nu|^2 + |\lambda_\nu H_u \nu|^2 + |\lambda_\nu L \nu|^2 + \left| \frac{\lambda_\mu \phi^2 H_u}{M_{\rm Pl}} \right|^2 \\ & + \left| \frac{1}{2} \lambda_\phi \nu^2 + \frac{\lambda_M \phi^3}{M_{\rm Pl}} \right|^2 + D \text{ terms} \\ & + \left(A_\nu \lambda_\nu LH_u \nu + A_\phi \lambda_\phi \phi \nu^2 + \frac{A_M \lambda_M \phi^4}{M_{\rm Pl}} + \text{c.c.} \right) \\ & + m_\nu^2 |\nu|^2 + m_L^2 |L|^2 - m_{H_u}^2 |H_u|^2 - m_\phi^2 |\phi|^2. \end{aligned} \tag{42}$$

As $|\phi|$ increases, ν will quickly acquire a large mass $\sim |\lambda_\phi \phi|$, and so will be constrained to the minimum of its potential

$$\nu \simeq - \frac{\lambda_\nu LH_u}{\lambda_\phi \phi}. \tag{43}$$

The effective potential then becomes

$$\begin{aligned} V = {} & (|\lambda_\nu H_u|^2 + |\lambda_\nu L|^2) \left| \frac{\lambda_\nu LH_u}{\lambda_\phi \phi} \right|^2 + \left| \frac{\lambda_\mu \phi^2 H_u}{M_{\rm Pl}} \right|^2 + \left| \frac{\lambda_M \phi^3}{M_{\rm Pl}} \right|^2 \\ & + D \text{ terms} + \left[\left(A_\phi - A_\nu + \frac{\overline{\lambda_M} \overline{\phi^3}}{2\phi} \right) \frac{(\lambda_\nu LH_u)^2}{\lambda_\phi \phi} \right. \\ & \left. + \frac{A_M \lambda_M \phi^4}{M_{\rm Pl}} + \text{c.c.} \right] + m_L^2 |L|^2 - m_{H_u}^2 |H_u|^2 - m_\phi^2 |\phi|^2. \end{aligned} \tag{44}$$

We assume the D terms constrain L and H_u to the D-flat direction LH_u. Then, writing $LH_u = \psi^2/2$, we get

$$\begin{aligned} V = {} & \left| \frac{\lambda_\nu^2 \psi^3}{2\lambda_\phi \phi} \right|^2 + \frac{1}{2} \left| \frac{\lambda_\mu \phi^2 \psi}{M_{\rm Pl}} \right|^2 + \left| \frac{\lambda_M \phi^3}{M_{\rm Pl}} \right|^2 - m_\psi^2 |\psi|^2 - m_\phi^2 |\phi|^2 \\ & + \left[\left(A_\phi - A_\nu + \frac{\overline{\lambda_M} \overline{\phi^3}}{2\phi} \right) \frac{\lambda_\nu^2 \psi^4}{4\lambda_\phi \phi} + \frac{A_M \lambda_M \phi^4}{M_{\rm Pl}} + \text{c.c.} \right], \end{aligned} \tag{45}$$

where $m_\psi^2 = (m_{H_u}^2 - m_L^2)/2$. To make this potential more transparent, we make the following change of variables:

$$V = m_s^2 M^2 \tilde{V}, \quad \phi = M \tilde{\phi}, \quad \psi = M \tilde{\psi},$$

$$m_\phi = m_s a_\phi, \quad m_\psi = m_s a_\psi, \tag{46}$$

$$A_\nu - A_\phi = m_s \alpha_\nu, \quad A_M = m_s \alpha_M, \quad \lambda_\mu = \frac{m_s M_{\rm Pl}}{M^2} \beta_\mu,$$

$$\lambda_M = \frac{m_s M_{\rm Pl}}{M^2} \beta_M, \quad \frac{\lambda_\phi}{\lambda_\nu^2} = \frac{M}{m_s} \gamma, \tag{47}$$

where the a's and the magnitudes of the α's and β's are of order one and we assume $|\gamma| \ll 1$. We then get

$$\begin{aligned} \tilde{V} = {} & -a_\phi^2 |\tilde{\phi}|^2 + |\beta_M|^2 |\tilde{\phi}|^6 + \left(-a_\psi^2 + \left| \frac{\tilde{\psi}^2}{2\gamma\tilde{\phi}} \right|^2 \right. \\ & \left. + \tfrac{1}{2} |\beta_\mu|^2 |\tilde{\phi}|^4 \right) |\tilde{\psi}|^2 + \left[\alpha_M \beta_M \tilde{\phi}^4 \right. \\ & \left. - \left(\alpha_\nu - \alpha_M \frac{\overline{\alpha_M} \overline{\beta_M} \overline{\tilde{\phi}^4}}{2|\alpha_M \tilde{\phi}|^2} \right) \frac{\tilde{\psi}^4}{4\gamma\tilde{\phi}} + \text{c.c.} \right]. \end{aligned} \tag{48}$$

When $\tilde{\phi} \ll 1$, $\tilde{\psi}$'s potential is stabilized at

$$|\tilde{\psi}|^2 \sim |\gamma| |\tilde{\phi}| \tag{49}$$

while its phase is coupled to that of $\tilde{\phi}$ by the term[10]

$$- \left(\alpha_\nu - \alpha_M \frac{\overline{\alpha_M} \overline{\beta_M} \overline{\tilde{\phi}^4}}{2|\alpha_M \tilde{\phi}|^2} \right) \frac{\tilde{\psi}^4}{4\gamma\tilde{\phi}} + \text{c.c.}, \tag{50}$$

the second term in the brackets being negligible at this stage.

When $\tilde{\phi}^3 \gtrsim \gamma$, the potential for the phase of $\tilde{\phi}$ will be dominated by the term

$$\alpha_M \beta_M \tilde{\phi}^4 + \text{c.c.} \tag{51}$$

and so in some sense we can regard the phase of $\tilde{\phi}$ as being determined modulo $\pi/2$. Put in a different way, $\tilde{\phi}$ will be pulled towards one of the minima of its potential and so its phase will be strongly biased towards

[9]One might imagine that our Affleck-Dine-type mechanism could also be implemented using say the right-handed electron sneutrino, which could plausibly have a small quartic coupling $\lambda_{\phi e}$, instead of LH_u. However, unlike LH_u, if it was unstable it would roll away from zero at some early time because all its couplings would be small. The term $A_\phi \lambda_\phi \nu^2 \phi + \text{c.c.}$ would then cause ϕ to roll away from zero causing a premature end to thermal inflation.

[10]The correlation induced by this term is different from that of Eq. (41) and so the phase of ψ will get a kick in the direction $\sin(\arg A_\phi - \arg A_\nu)$ while the phase of ϕ will get a kick in the opposite direction. This may contribute to the net lepton number generated, in addition to the similar effect to be described below.

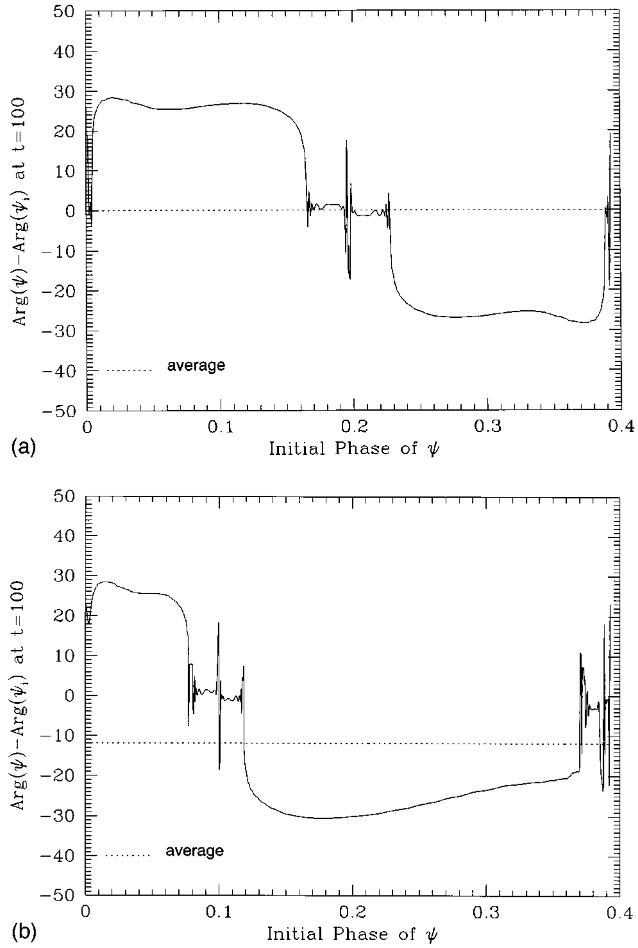

FIG. 2. Numerical simulation to show the nonzero net lepton number generated. The same parameters as in Fig. 1 were used except for the following. Motivated by Eq. (41), the initial phase of ψ was taken to be random while the initial phase of ϕ was taken to be given by $\arg(\phi)=4\arg(\psi)+C$. The lepton number produced, as measured by $\arg\psi(t=100)-\arg\psi(t=0)$, is plotted against the initial phase of ψ. The dotted line gives the average value. The plots correspond to the following values of the parameters (a) $\arg\alpha_M=\pi$ (which is CP conserving) and $C=0$, (b) $\arg\alpha_M=3.1$ and $C=0$.

$$\alpha_M\beta_M\tilde{\phi}^4=-|\alpha_M\beta_M\tilde{\phi}^4|. \tag{52}$$

The phase of $\tilde{\psi}$ is then determined modulo $\pi/8$ by the term in Eq. (50).

When $\tilde{\phi}$ becomes of order one, two things happen. First, the second term in the brackets in Eq. (50) becomes of order one and gives the phase of $\tilde{\psi}$ a kick in the direction $\sin(\arg\alpha_\nu-\arg\alpha_M)$. Note that even before this, $\tilde{\psi}$ will have had some angular momentum about $\tilde{\psi}=0$, but it averages out to zero in the Universe as a whole, as is shown in Fig. 2(a). This new contribution has a direction determined by the parameters of the Lagrangian and so will give a nonzero net contribution, as is shown in Fig. 2(b). Put another way, the difference in phase between α_ν and α_M is our source of CP violation. Second, the last term in the brackets in

$$\left(-a_\psi^2+\left|\frac{\tilde{\psi}^2}{2\gamma\tilde{\phi}}\right|^2+\frac{1}{2}|\beta_\mu|^2|\tilde{\phi}|^4\right)|\tilde{\psi}|^2 \tag{53}$$

becomes of order one giving $\tilde{\psi}$ a net positive mass squared and so causing it to spiral back in to $\tilde{\psi}=0$.

Assuming the expected broad parametric resonance [21] provides enough damping, both ϕ and ψ should then become trapped near their vacuum expectation values, after which the parametric resonance becomes less efficient. ψ's potential near $\psi=0$ conserves angular momentum, or in other words lepton number, and so ψ's lepton number is conserved.

The Affleck-Dine condensate ψ will decay well before the Hubble expansion reduces its amplitude to the electroweak scale, and so will release enough thermal energy to restore the electroweak symmetry. The lepton asymmetry will then be converted into a baryon asymmetry

$$\frac{n_B}{s}\sim\frac{1}{3}\left(\frac{n_L}{s}\right) \tag{54}$$

by the usual electroweak effects [22,23]. Finally, after the temperature has dropped to a few GeV, the flaton decay will complete, releasing substantial entropy.

The baryon asymmetry generated in this way is roughly estimated to be

$$\frac{n_B}{s}\sim\frac{T_f n_L}{m_\phi n_\phi}\sim\frac{\theta T_f n_\psi}{m_\phi n_\phi}\sim\frac{\theta|\lambda_\phi|T_f}{|\lambda_\nu|^2 M}\sim\frac{\theta(100\ \mathrm{GeV})^2 T_f}{m_{\nu_L}M^2} \tag{55}$$

$$\sim 10^{-10}\theta\left(\frac{10\ \mathrm{eV}}{m_{\nu_L}}\right)\left(\frac{T_f}{\mathrm{GeV}}\right)\left(\frac{10^{11}\ \mathrm{GeV}}{M}\right)^2, \tag{56}$$

where we have used Eqs. (49) and (3). θ is defined by this equation and will depend on the phase difference between α_ν and α_M as well as the detailed dynamics.

As discussed in Sec. II, we expect $T_f\sim 1-10$ GeV and $M\sim 10^{10}-10^{12}$ GeV. Neutrino phenomenology [11,24] suggests $m_{\nu_{\tau L}}\sim 5$ eV, $m_{\nu_{\mu L}}\sim 10^{-2}-10^{-3}$ eV, and $m_{\nu_{eL}}\lesssim m_{\nu_{\mu L}}$. Therefore, in order for Eq. (56) to give the baryon asymmetry of Eq. (1), we require θ to be roughly[11]

$$\theta_\tau\sim 10^{-4}-10, \tag{57}$$

$$\theta_e\lesssim\theta_\mu\sim 10^{-7}-10^{-2}, \tag{58}$$

depending on which generations make up the Affleck-Dine LH_u direction. The eventual measurement of the Higgs boson and slepton masses should help to determine which of these ranges is the appropriate one (or rule out the whole scenario), and a measurement of $m_{\nu_{eL}}$ would narrow the uncertainty in θ_e.

[11] $\theta\gtrsim 1$ corresponds to the scenario being unviable.

IV. CONCLUSIONS

Right-handed neutrinos should acquire their masses due to the vacuum expectation value of the flaton that gives rise to thermal inflation, not some composite GUT operator. This will have important implications for GUT model building. In particular, SO(10) GUT's are strongly disfavored because the flaton would have to be in a **126** representation which is difficult to derive from superstrings and one would have a flaton-125 splitting problem, in addition to the usual doublet-triplet splitting problem.

The μ term of the MSSM should also be generated by the VEV of the flaton.

Our Affleck-Dine-type mechanism generates a baryon asymmetry which is roughly estimated to be

$$\frac{n_B}{s} \sim 10^{-10}\theta\left(\frac{10\ \mathrm{eV}}{m_{\nu_L}}\right)\left(\frac{T_f}{\mathrm{GeV}}\right)\left(\frac{10^{11}\ \mathrm{GeV}}{M}\right)^2, \qquad (59)$$

where θ is the lepton asymmetry per Affleck-Dine particle. θ depends on the difference in phase between the soft supersymmetry-breaking parameters of W_{seesaw} and W_{VEV} [c.f. Eqs. (6) and (27)], as well as the detailed dynamics.

We also make the prediction

$$m_L^2 < m_{H_u}^2 \qquad (60)$$

modulo renormalization effects, where $-m_{H_u}^2$ is the soft supersymmetry-breaking mass squared of H_u, and m_L^2 is the soft supersymmetry-breaking mass squared of a lepton doublet.

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