FLAT DIRECTIONS AND BARYOGENESIS IN SUPERSYMMETRIC THEORIES

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Abstract

Flat directions are a generic feature of supersymmetric theories. They are of cosmological interest because they can lead to coherent production of scalars. In the early universe such flat directions could be dangerous due to the potentially large energy density and the late decay of the associated scalars when they have only $1/M_p$ couplings (Polonyi problem). On the other hand, flat directions among the standard model fields can carry baryon number and lead to a possible mechanism for baryogenesis (Affleck Dine baryogenesis). When considering the cosmological consequences of the flat directions, it is important to take into account the soft potential with curvature of order the Hubble constant due to supersymmetry breaking in the early universe. In this talk, we discuss flat directions, their potential cosmological implications focusing on Affleck-Dine baryogenesis, and how the standard picture of their evolution must be modified in the presence of the large supersymmetry breaking in the early universe.

to appear in the Proceedings of the XXXth Rencontres de Moriond
"Electroweak Interactions and Unified Theories"
Les Arcs, France, March 1995

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[1]NSF Young Investigator Award, Alfred P. Sloan Foundation Fellowship, DOE Outstanding Junior Investigator Award. This work is supported in part by funds provided by the U.S. Department of Energy (D.O.E.) under contract #DE-FC02-94ER40818 and by N.S.F. grant PHY89-04035.

[*]Based on work done in collaboration with Michael Dine and Scott Thomas.

Flat directions are a common feature of supersymmetric theories and have many important implications. Here we will focus on some of the cosmological aspects of the existence of flat directions which can lead to the coherent production of a scalar condensate. Before we proceed however, we address the question of why supersymmetric cosmology is of interest. One well known reason is that one can potentially constrain particle physics models. Chief among the requisite constraints are by that the universe is not overclosed and that nucleosynthesis can proceed successfully. This latter constraint gives rise to the cosmological moduli (Polonyi) problem [1, 2, 3] for example. For a more extensive discussion of the potential problems, see ref. [10]

A related reason for studying supersymmetric cosmology is the difficulty of experimentally probing the high scales associated with supersymmetric physics. Cosmology offers an alternative probe to high physics scales like that associated with supersymmetry breaking, flavor, or GUTS. In particular, operators which are suppressed by high mass scales might nonetheless be relevant at early times when fields take large expectation values. Unfortunately, this also means that if a mechanism of baryogenesis exists which exploits CP violating high dimension operators, this nonstandard model CP violation is experimentally inaccessible, so that there is not necessarily a connection between CP violation which will be explored in a laboratory and that which was important for baryon number creation.

A third reason is that we know there is a baryon asymmetry in the universe and it is worthwhile to understand its origin. GUT scale baryogenesis must contend with $B+L$ violation at late times and a relatively low reheat temperature after inflation. Weak scale baryogenesis is a nice possibility, but there are many difficult questions regarding the nature of the phase transition. It is useful to study in detail alternative mechanisms for baryogenesis. The Affleck Dine (AD) mechanism [4] in particular is a beautiful way to utilize the flat directions which are in any case present in the supersymmetric standard model for the creation of baryons. Although there this mechanism has been studied in the past, we show that the more likely picture of how the AD mechanism works is substantially different from what has been studied, and yields qualitatively and quantitatively different conclusions.

In this talk, we will first review the flat directions of supersymmetric theories, and the old picture for coherent production in the early universe. We will argue that this picture is modified because of supersymmetry breaking in the early universe, which effectively generates a soft superysymmetry breaking scale of order H, where H is the Hubble constant. When the Hubble constant is bigger than $m_{3/2}$ which is of order of the weak scale, the supersymmetry breaking potential determined by H scale terms is dominant, and changes significantly the picture of the evolution of the fields at early times.

We will then discuss two implications of this revised picture of the early universe field evolution. We will show that it changes the standard scenario for the Polonyi problem, and potentially presents a solution. The major focus of this talk however will be the implications for baryogenesis through the Affleck-Dine [6] mechanism. We will discuss the evolution of the Affleck-Dine field in the presence of the Hubble constant scale potential and also include higher dimension operators which can generate a potential for the "flat" directions, even in the supersymmetric limit. We will see that we have a relatively simple and predictive scenario for baryogenesis. We find that depending on the identity of the AD field, one can naturally obtain $n_b/s \geq 10^{-10}$. This is in contrast to the previous picture, according to which additional entropy deposition in the late universe was required.

Flat directions are peculiar to supersymmetric theories. They correspond to fields with no classical potential. In the absence of supersymmetry, they would be highly unnatural. However, supersymmetric nonrenormalization theorems protect massless fields to keep them massless, even with radiative corrections. There are several contexts in which flat directions are relevant. One is the moduli space (of string theory for example) which parameterizes a large vacuum degeneracy of physically inequivalent theories not related by a symmetry. These flat directions might have no perturbative superpotential couplings. However, even if they do couple in the superpotential, one can often find combinations of fields for which there is no potential due to accidental degeneracies. This happens even in the supersymmetric standard model, where there are a large number of such directions.

A simple example of such a flat direction is when the H_u and L fields have equal expectation values so that the D term contribution to the potential is cancelled. Another example of a flat direction of the supersymmetric standard model is $Q_1 L_1 \bar{d}_2$, where the numerical index labels generation number. There are many such flat directions for which both the D and F type contributions to the potential vanish.

It is not true however that the potential for flat directions vanishes identically. There are two ways in which one expects the flat directions to be lifted. One possibility is nonrenormalizable operators in the superpotential. These are not necessarily present, but in many cases, if they are consistent with all existing symmetries, one would expect such operators to occur, suppressed by a high dimension scale which might be M_p or M_G or some other high scale of the theory. These operators turn out to be very important to the AD mechanism because they are the source of B (or L) and CP violation.

The other source of the potential for the flat direction fields is soft supersymmetry breaking. In the absence of additional symmetry, one would expect all fields to get a mass of order $m_{3/2}$ when supersymmetry is broken. Of course, in the early universe, when $m_{3/2} \ll H$, this is negligible.

Why are these flat direction fields of cosmological relevance? Assume a flat direction field has “no” potential, so the initial field value is undertermined. In this case, one would expect some random initial displacement of the field from its zero temperature minimum. Consider now the classical evolution of the zero mode,

$$\ddot{\phi} + 3H\dot{\phi} + V'(\phi) = 0 \tag{1}$$

where H is the damping term in an expanding background. It is clear that if $H^2 \gg V''$ we are in the overdamped case and$\phi \approx$ constant, whereas if $H^2 \ll V''$, the field is underdamped and ϕ evolves to minimum and oscillates freely. These oscillations imply a coherent production of nonrelativistic particles This coherent production of scalar particles in the early universe is generic in supersymmetric theories with flat directions.

Now what are the implications of these particles? Clearly it will depend on the couplings and quantum numbers of the field. If there are only gravitational strength $(1/M_p)$ interactions, one finds the Polonyi problem. The standard statement of the problem is as follows. Suppose the ϕ field has only $1/M_p$ couplings and $m_\phi \approx O(m_{3/2})$. Suppose also that the initial field value $\phi_0 \sim M_p$. Then the field is frozen until $H \sim m_{3/2}$, after which it begins to oscillate freely. At this point, $\rho_\phi \approx m^2 M_p^2 \Rightarrow \rho \approx \rho_{universe}$. Subsequently, ρ_ϕ dominates the energy density of the universe. When $H \approx \Gamma \approx m^3/M_p^2$, ϕ decays and the associated reheat temperature $T_R \approx \sqrt{\Gamma M_p} \approx 10\text{keV}$. This is too low for successful nucleosynthesis. Moreover, even if the density of the condensate is somewhat lower, late decays of the condensate would destroy the nuclei which have already been created, so the problem is even more severe than it naively appears.

An alternative scenario for the coupling of the flat direction is that the field is a flat direction of the standard model carrying net $B-L$. In general, the AD mechanism requires that $B-L$ is violated at large field value, but conserved for small field value. It is assumed that the field at early times is displaced from its true minimum. Eventually, the field is driven towards the origin through the equations of motion. The mechanism works when baryon number violation is important as the field moves in towards the origin, so that baryon number is stored in the coherent condensate. Once the field is oscillating about the minimum, the baryon number violating operators are no longer significant and baryon number is conserved. Notice that the three conditions for baryogenesis [5] can be satisfied. There is CP violation through the phase difference between the initial phase of the field and that of the baryon number violating operator. There is B violation by assumption (although we have not yet specified the source). Finally, the large initial value for the field ϕ is a nonequilibrium situation.

However, at this point, we clearly do not have the whole story. We would like to better understand what are the initial field values, what provides the baryon number violation (for the AD mechanism), and whether the flat directions really are as flat as we have been assuming. In the rest of the talk, I argue that the picture we have been presenting is not the whole story. The most significant aspect which has been omitted is that supersymmetry is necessarily broken in the early universe. Furthermore the scale which acts as the soft supersymmetry breaking parameter is H, not $m_{3/2}$ [6, 7, 8]. This means that the picture of the fields as being "frozen" at early times is not correct. In the remainder of the talk, we explore the consequences.

We will briefly discuss the implication for the Polonyi problem. As for the AD mechanism, we will find that one can incorporate the B violation in higher dimension operators in the superpotential. We will find that the AD mechanism does not always work; what is required is that the effective mass squared in the early universe is negative in order to drive the AD field classically to large field value. However, at this point, all that is needed in order to derived n_b/s will be the dimension of the superpotential operator which stabilizes the potential and the reheat temperature after inflation. The picture of the evolution of the AD field is significantly altered, but a very elegant scenario emerges.

First let us understand why supersymmetry is necessarily broken in the early universe. The Hubble constant H is related to the matter energy density ρ by $H^2 = \rho/M_p^2$. Therefore, an expanding universe implies a finite positive energy density in the early universe, implying supersymmetry is broken. In the case the energy is carried by radiation, supersymmetry is broken by the different thermal occupation numbers of bosons and fermions. However this is well known and is not the effect which is of interest to us, since by the time the universe is radiation dominated after inflation, the zero temperature supersymmetry breaking is dominant. The case we will be most interested in is during and subsequent to inflation, when the Hubble constant is very large compared to $m_{3/2}$. Notice that this is in accord with the well known result that supersymmetry is broken in deSitter space.

We now consider the potential which is generated for the flat direction field. Suppose $\rho = F_I^\dagger F_I$ and $K \supset \phi^\dagger \phi I^\dagger I/M_p^2$. Then $m_\phi^2 \approx F_I^\dagger F_I/M_p^2 \approx H_I^2$. In fact, more generally if $\langle O \rangle = \rho$ and there is an operator in the potential of the form $O\phi^\dagger\phi/M_p^2$, then $m_\phi^2 \approx \rho/M_p^2 \approx H^2$. So we see there is generally a soft supersymmetry breaking mass in the early universe of order H. This Hubble scale mass is essential to the evolution when $H > m_{3/2}$.

One can ask whether this mass is necessarily present. The answer is yes, unless there is fine tuning. The point is that even with *minimal* Kahler potential, such a mass term occurs. Recall the supersymmetric potential takes the form

$$V = e^K \left(\left(W_i + \frac{K_i W}{M_p^2} \right) K^{i\bar{j}} \left(W_j + \frac{K_j W}{M_p^2} \right)^* - 3 \frac{W^* W}{M_p^2} \right) + \frac{1}{2} g^2 D^a D^a \tag{2}$$

Suppose there is minimal $K = \phi^\dagger \phi$. Then one can see that there is a mass of order H when the potential energy is finite and positive, and with the minimal Kahler potential the Hubble scale mass is positive.

With a nonminimal Kahler potential (or $I \approx M_p$) the mass formula is more complicated. For example, if $K \supset \phi^\dagger \phi I^\dagger I / M_p^2$, then $K_{I\bar{I}} = \phi^\dagger \phi / M_p^2 \Rightarrow m^2 \subset -F_I^\dagger F_{\bar{I}} / M_p^2 \approx -H^2$. Notice that a nonminimal K is to be expected. In fact, such higher dimension operators are necessarily present in the Kahler potential as counterterms for the running of the soft masses [9]. And as we will see later, the standard model running of the soft mass might be adequate to give negative mass squared to LH_u, the preferred AD field (as we will argue).

So we conclude there is certainly a mass of order H but its sign and magnitude is not determined.

We now consider the field evolution in the presence of the potential due to supersymmetry breaking in the early universe. The first observation is that generically, the fields are not frozen; that is they evolve to a local minimum as they are not overdamped, since $m \approx H$. For example, with minimal Kahler potential, the field will roll to the origin.

What does this imply in terms of the Polonyi problem? It should be clear that the "initial" field value is not random; the field evolves in the effective potential to a nearby local minimum. This might be a concrete realization of the problem, since in general, the minima in the early universe and today do not coincide. However it also suggests a solution to the problem in the cases where these minima are the same. This might be true for example when the minimum is a point of enhanced symmetry [6, 7]. It can also arise as a consequence of a factorization of the Kahler potential so that minima of K are local minima of the potential [10].

It should be noted however that an exception to this scenario is the case where additional symmetries protect the mass of the flat direction field, eg a pseudogoldstone boson. Then the associated symmetry breaking parameter is necessary to generate a mass, so that the mass will generically be suppressed relative to H. In this case, one expects the field to always be overdamped.

What are the implications of this Hubble scale mass for the Affleck Dine mechanism? First suppose the mass squared was positive. Then the field will be classically driven to the origin, in which case the amplitude of the AD field vanishes, and there is no baryon number creation! Previous authors suggested that quantum, not classical effects, drive the field away from the origin. But this will not work because when $m \approx H$, the coherence length is too small and baryon number would average to zero over the observable universe.

However, we have argued that the Hubble scale mass squared is not necessarily positive. In fact, we will now show that the negative mass squared scenario works perfectly. The AD field is driven classically to a large field value. When $H \approx m_{3/2}$, the low energy supersymmetry breaking gives a comparable contribution to the potential as the H dependent terms. About this time, the mass passes through zero to become positive, and the AD field rolls towards the true minimum at the origin. What makes this whole scenario work so efficiently is that the baryon number violating operators turn out to also be comparable to other terms in the potential at this time, so that baryon violation is essentially maximal. The condensate stores baryon number, and n_b/s is determined by the relative fraction of the AD field at this time.

Let us examine in more detail the salient features of the evolution. It will be useful to divide our analysis into three periods; during inflation, the post inflation-inflaton matter dominated era when $m_{3/2} \ll H$, and the post inflation-inflaton matter dominated era when $m_{3/2} \approx H$. In our analysis, we assume the AD potential arises from

- Nonrenormalizable terms in the superpotential
- Soft masses (and possibly soft A type terms) of order H due to supersymmetry breaking in the early universe
- Zero temperature supersymmetry breaking parameters whose scale is determined by $m_{3/2}$. These are negligible in the early stages of the evolution, but important when baryon number is established and subsequently, that is, for $H \leq m_{3/2}$.

Now let us consider the field evolution during the three epochs outlined above. First consider the potential during inflation.

$$V(\phi) = -cH_I^2|\phi|^2 + \left(\frac{a\lambda H_I \phi^n}{nM^{n-3}} + \ h.c.\ \right) + |\lambda|^2 \frac{|\phi|^{2n-2}}{M^{2n-6}} \tag{3}$$

where c and a are constants of $\mathcal{O}(1)$, and M is some large mass scale such as the GUT or Planck scale. For $H_I \gg m_{3/2}$ soft terms arising from the hidden sector are of negligible importance. The minimum of the potential (3), is given by

$$|\phi_0| = \left(\frac{\beta H_I M^{n-3}}{\lambda}\right)^{\frac{1}{n-2}} \tag{4}$$

where β is a numerical constant which depends on a, c, and n. Notice that ϕ_0 is parameterically between H_I and M.

- During inflation, the AD field evolves exponentially to the minimum of the potential, determined by the induced negative mass squared and nonrenormalizable term in the superpotential. This process may be thought of as establishing "initial conditions" for the subsequent evolution of the field. In the presence of a supersymmetry breaking "A" type term, the phase will also roll to its minimum. Otherwise a random value for the phase is taken. Either way, it is the difference between this phase and the real minimum for the phase which establishes baryon number.

  It is easy to determine that the field rolls efficiently to its minimum. Suppose it started far away since the field is rapidly oscillating with a slowly decreasing envelope. The time rate of change of the energy in ϕ can be found from the equations of motion,

$$\frac{dE}{dt} = \dot{\phi}\frac{d}{d\phi}(T+V) = -3H_I\dot{\phi}^2 = -6H_I(E-V) \tag{5}$$

  where T is the kinetic energy. Using the expression for $V(\phi)$ for large ϕ and averaging over a period gives $\dot{\phi}_m \simeq -6H_I/(2n-1)\phi_m$. We therefore conclude that in the large ϕ regime, ϕ decreases exponentially towards smaller values,

$$\phi_m \simeq e^{-6H_I t/(2n-1)}\phi_i \tag{6}$$

  where ϕ_i is the initial value of the field with espect to the origin. Thus after just a few e-foldings ϕ is near a minima. Once near a minimum, the field evolves like a damped harmonic oscillator.

- After inflation the universe enters a matter era dominated by the coherent oscillations of the inflaton. The minimum of the potential is time dependent (as it is tied to the instantaneous value of the Hubble parameter). The AD field oscillates near this time dependent minimum with decreasing amplitude.

  During a matter era the Hubble constant is related to the expansion time by $H = \frac{2}{3}t$. The equation of motion for ϕ is then

$$\ddot{\phi} + \frac{2}{t}\dot{\phi} + V'(\phi) = 0 \tag{7}$$

  where $V(\phi)$ is still given by (3), though the dimensionless constants c and a may be different, and H is now time dependent.

  We can obtain greater insight into the solutions of (7), and also obtain a form more suitable for numerical study by making changes of variables. We define

$$z = \log t.$$

  It is also useful to define the dimensionless field χ by scaling with respect to the instantaneous minimum of the effective potential.

$$\phi = \chi\phi_0(t) = \chi\left(\frac{\beta}{\lambda}M^{n-3}e^{-z}\right)^{\frac{1}{n-2}}$$

  where $\beta = \sqrt{c'/(n-1)}$ for $a = 0$, and $c' = \frac{4}{9}c$. The equation of motion in these rescaled variables is then

$$\ddot{\chi} + \left(\frac{n-4}{n-2}\right)\dot{\chi} - \left[c' + \frac{n-3}{(n-2)^2}\right]\chi + c'\chi^{2n-3} = 0 \tag{8}$$

  The rescaled problem is so simple because the effective mass term, Hubble damping term, and acceleration term are all homogeneous in z.

  The qualitative behavior of the solution is now much more apparent. We see that we have eliminated all large and small parameters from the differential equation. Unless the damping term is negative, we expect the field to track the minimum of the true potential. Notice that the effective potential for the rescaled variable has a new contribution to the effective mass but the AD field is of the same order of magnitude as if it were at the minimum of the potential.

  We now see that for $n > 4$, the field is driven towards the minimum, while for $n < 4$, it would be driven away. This latter case would however correspond to a field which was not flat, so it is not of interest to us. The case $n = 4$ is interesting in that the rescaled field is not driven closer to the minimum than its iniital value, although the true field is due to the scaling.

  In any case, it should be clear that in all cases of interest, the field amplitude essentially follows the effective minimum determined by balancing the time dependent (Hubble dependent) mass term and the nonrenormalizable term in the superpotential. Because the mass is decreasing with time, the field amplitude decreases with time.

- When $H \sim m_{3/2}$ the soft potential arising from hidden sector supersymmetry breaking becomes important and the sign of the mass squared becomes positive. At this time, the B-violating A term arising from the hidden sector is of comparable importance to the mass term, thereby imparting a substantial baryon number to the condensate. The fractional baryon number carried by the condensate is near maximal, more or less independent of the details of the flat direction. Subsequent to this time, the baryon number violating operators are negligible so the baryon number (in a comoving volume) is constant.

The potential takes the form

$$V(\phi) = m_\phi^2 |\phi|^2 - \frac{c'}{t^2}|\phi|^2 + \left(\frac{(Am_{3/2} + aH)\lambda\phi^n}{nM^{n-3}} + h.c. \right) + |\lambda|^2 \frac{|\phi|^{2n-2}}{M^{2n-6}} \tag{9}$$

where $m_\phi \sim m_{3/2}$. At early times the field tracks near the time dependent minimum as discussed in the last section. Therefore when $H \sim m_{3/2}$ all the terms in (9) have comparable magnitudes. Since the soft terms have magnitudes fixed by $m_{3/2}$ the field is no longer near critically damped, but becomes underdamped as H decreases beyond $m_{3/2}$. In addition the $m_\phi^2|\phi|^2$ term comes to dominate the $-cH^2\phi^2$ term as H decreases. The field therefore begins to oscillate freely about $\phi = 0$ when $H \sim m_{3/2}$, with "initial" condition given by $\phi_0(t)$ (eq. (4)) with $t \sim m_{3/2}^{-1}$.

Crucial for the generation of a baryon asymmetry are the B violating A terms in (9). However, when $H \sim m_{3/2}$ all the terms have comparable magnitude, including the A terms. Since $V_B \sim V_{\not B}$ when the field begins to oscillate freely a large fractional baryon number is generated in the "initial" motion of the field when m^2 becomes positive. Notice that in this negative mass squared scenario n_b/n_ϕ is roughly *independent* of λ/M. This is because the value of the field is determined precisely by a balance of (negative) soft mass squared term and nonrenormalizable supersymmetric term. That the B violating A term also has the same magnitude follows from supersymmetry since its magnitude is the root mean square of the soft mass term and nonrenormalizable supersymmetric term. In this scenario there is no need for ad hoc assumptions about the initial value of the field when it begins to oscillate freely. The expectation that $n_b/n_\phi \sim \mathcal{O}(1)$ falls out naturally.

We also did numerical simulations to confirm the above conclusions. Again it is useful to rescale variables. The field is rescaled as

$$\phi \rightarrow \left(\frac{m_{3/2}M^{n-3}}{\lambda} \right)^{\frac{1}{n-2}} \phi$$

From the arguments above , up to a numerical constant of order unity this is just the value of the field when $H \sim m_{3/2}$. All other mass scales and time are rescaled with respect to $m_{3/2}$. The equation of motion (9) with $a = 0$ and $\theta_A + \theta_\lambda = 0$ is then

$$\ddot{\phi} + \frac{2}{t}\dot{\phi} + \left(m_\phi^2 - \frac{c'}{t^2} \right) \phi + A\left(\phi^*\right)^{n-1} + (n-1)\left(\phi^*\phi\right)^{n-2}\phi = 0 \tag{10}$$

Notice that the independence from λ/M^{n-3} is manifest in this form. The equation of motion for the real and imaginary parts (appropriate for numerical integration) are

$$\ddot{\phi}_R + \frac{2}{t}\dot{\phi}_R + \left(m_\phi^2 - \frac{c'}{t^2} \right) \phi_R + A|\phi|^{n-1}\cos\left((n-1)\theta\right) + (n-1)|\phi|^{2n-4}\phi_R = 0$$

$$\ddot{\phi}_I + \frac{2}{t}\dot{\phi}_I + \left(m_\phi^2 - \frac{c'}{t^2}\right)\phi_I - A|\phi|^{n-1}\sin\left((n-1)\theta\right) + (n-1)|\phi|^{2n-4}\phi_I = 0 \tag{11}$$

where $\phi = \phi_R + i\phi_I$, and $\theta = \text{Arg}\ \phi$. It is straightforward to integrate these equations forward, assuming the field begins near the minimum at a time when the mass is still negative. When $t \sim 1$ ($H \sim m_{3/2}$) the field feels a "torque" from the A term, and spirals inward in the harmonic potential. The nonzero $\dot{\theta}$ in the trajectory gives rise to the baryon number. At late time, n_b/n_ϕ asymptotes to a constant value.

At very late stages of the evolution when $H \ll m_{3/2}$, the only potential term which is relevant in (9) is the soft mass term $m_\phi^2|\phi|^2$ which is of course B conserving. The baryon number created during the epoch $H \sim m_{3/2}$ is therefore conserved by the classical evolution of ϕ for $H \ll m_{3/2}$.

This establishes that $n_b/n_0 \approx 1$. What is important however is n_b/s. This is determined by both the baryon fraction of the condensate and also by the fractional density of matter carried by the condensate.

- The inflaton decays when $H < m_{3/2}$ (consistent with the gravitino bound on the reheat temperature). The condensate will decay soon afterwards through scattering with the thermal bath (but see [6] for some caveats). The baryon to entropy ratio subsequent to inflation is determined from

$$\frac{n_b}{s} \approx \frac{n_b}{n_\phi}\frac{T_R}{m_\phi}\frac{\rho_\phi}{\rho_I} \tag{12}$$

where n_b and n_ϕ are baryon and AD field number densities, T_R is the reheat temperature, $m_\phi \sim m_{3/2}$ is the low energy mass for the AD field, and ρ_ϕ and ρ_I are the AD field and inflaton mass densities (both at the time of inflaton decay).

$$\frac{\rho_\phi}{\rho_I} \approx \left(\frac{m_{3/2}M^{n-3}}{\lambda M_p^{n-2}}\right)^{2/(n-2)}. \tag{13}$$

For $n = 4$, $\rho_\phi/\rho_I \sim 10^{-16}(M/\lambda M_p)$, while for $n = 6$, $\rho_\phi/\rho_I \sim 10^{-8}(M^3/\lambda M_p^3)^{1/2}$, Notice for smaller (λ/M^{n-3}) the direction is effectively flatter, and ϕ_0 and ρ_ϕ are larger. A greater total energy is therefore stored in the oscillating condensate for smaller λ or larger n.

Notice that for $n = 4$ and high reheat temperature (motivated by naturalness arguments), n_b/s turns out to be just about right. It is worth noting that all flat directions can be lifted by operators with $n \leq 6$, so that it is conceivable that any direction can give the correct n_b/s without additional entropy dump, although $n = 6$ requires a low reheat temperature, of order the weak scale.

The LH_u direction appears to be especially promising for natural production of the correct baryon to entropy ratio. The only directions which carry $B - L$ and can be lifted at $n = 4$ in the standard model are the LH_u directions. The nonrenormalizable operator is then

$$W = \frac{\lambda}{M}(LH_u)^2 \tag{14}$$

At low energies this is the operator which gives rise to neutrino masses. For baryogenesis along the LH_u direction in this scenario, n_b/s can therefore be related to the lightest neutrino mass since the field moves out furthest along the eigenvector of L_iL_j corresponding to the smallest eigenvalue of the neutrino mass matrix.

$$\frac{n_b}{s} \sim 10^{-10} \left(\frac{T_R}{10^9 \text{ GeV}}\right) \left(\frac{10^{-5} \text{ eV}}{m_\nu}\right) \tag{15}$$

If $T_R < 10^9$ GeV in order to satisfy the gravitino bound [11], then at least one must neutrino be lighter than roughly 10^{-5} eV.

Another reason this field is so interesting as a candidate for the AD field is that even with *minimal* Kahler potential, the scenario we have outlined might work. The point is that we know the H_u mass is driven negative by renormalization group running due to the large top quark Yukawa coupling. If the LH_u mass is driven negative at a sufficiently high scale, the scenario we have outlined works with no additional nonminimal Kahler terms present (although as we have emphasized these are there as counterterms in any case). Although it might seem dangerous to have the mass run so quickly negative, since the running also determines the low energy minimum, this is not necessarily the case. This is because of the indirect relation between the soft parameters at high and low temperature, and also because the "μ" term is not necessarily of order H, whereas there is certainly a μ term of order $m_{3/2}$ at zero temperature. This can allow for our scenario to work where the parameters are such as to give the correct low energy minimum. Diego Castano has verified this scenario numerically.

To conclude, we have seen that cosmology in a supersymmetric universe is intriguing and subtle. The existence of flat directions is very important; so is the fact that there is necessarily supersymmetry breaking in the early universe. This means that fields are *not* frozen; they evolve according to their classical field equations to their minimum. This qualitatively changes scenario for coherent field production. Whether or not they evolve to the low energy minimum determines whether or not there is a Polonyi problem.

Many of our interesting conclusions apply to the Affleck-Dine Mechanism. We have seen that the picture which has emerged is very different from the standard scenario in the literature. We have found a negative mass squared in the early universe is essential in order to drive the field to a large value (in the absence of a tuned small mass or a symmetry mechanism). We have also found a predictive scenario for the field evolution. The field value varies continuously subsequent to inflation, tracking the instantaneous minimum. This gives a definite motivation for the field magnitude at the time the AD field is driven to the origin. Happily, the A type terms have a magnitude such that baryon number violation is maximal. We have found a formula for n_b/s which is roughly independent of the details of the problem and depends primarily on T_R and n. The LH_u direction would be best, because it does not require an additional source of entropy; however, almost any direction could work.

Acknowledgements: I would like to thank Greg Anderson, Diego Castano, Dan Freedman, Steve Martin, Graham Ross, and Bob Singleton for useful conversations, and Greg Anderson for his comments on the manuscript.

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