SCIPP 95-13
MIT-CTP-2421
SLAC-PUB-95-6776

Supersymmetry Breaking in the Early Universe

Michael Dine*
Santa Cruz Institute for Particle Physics
University of California, Santa Cruz, CA 95064

Lisa Randall†
Laboratory for Nuclear Science and Department of Physics
Massachusetts Institute of Technology
Cambridge, MA 02139

Scott Thomas‡
Stanford Linear Accelerator Center
Stanford University
Stanford, CA 94309

Abstract

Supersymmetry breaking in the early universe induces scalar soft potentials with curvature of order the Hubble constant. This has a dramatic effect on the coherent production of scalar fields along flat directions. For the moduli problem it generically gives a concrete realization of the problem by determining the field value subsequent to inflation. However it might suggest a solution if the minimum of the induced potential coincides with the true minimum. The induced Hubble scale mass also has important implications for the Affleck-Dine mechanism of baryogenesis. This mechanism requires large squark or slepton expectation values to develop along flat directions in the early universe. This is generally not the case if the induced mass squared is positive, but does occur if it is negative. The resulting baryon to entropy ratio depends mainly on the dimension of the nonrenormalizable operator in the superpotential which stabilizes the flat direction, and the reheat temperature after inflation. Unlike the original scenario, it is possible to obtain an acceptable baryon asymmetry without subsequent entropy releases.

⋆ Work supported in part by the Department of Energy.
† NSF Young Investigator Award, Alfred P. Sloan Foundation Fellowship, DOE Outstanding Junior Investigator Award. Supported in part by DOE contract DE-AC02-76ER03069 and by NSF grant PHY89-04035.
‡ Work supported by the Department of Energy under contract DE-AC03-76SF00515.

1. Introduction

Low energy supersymmetry, if it exists in nature, is likely to have dramatic consequences for the early universe. One of the most striking stems from the existence of flat directions in the scalar potential. Such directions are a generic feature of supersymmetric theories, unfamiliar in conventional field theories. In string theory, for example, there are often moduli which label degenerate classical vacuum states of the string. These states remain degenerate to all orders in perturbation theory. In the minimal supersymmetric standard model (MSSM) there exist, at the level of renormalizable operators and ignoring supersymmetry breaking, a large number of flat directions along which some combination of squark, slepton, and Higgs fields have expectation values. In the early universe if the fields parameterizing a flat direction start displaced from the true minimum, coherent oscillations result when the Hubble constant becomes smaller than the effective mass. The energy stored in these oscillations amounts to a condensate of nonrelativistic particles. The production of such condensates should be a generic feature of supersymmetric theories. In this paper we discuss the effect of supersymmetry breaking in the early universe on coherent field production, with emphasis on the cosmological moduli problem [1-3] and Affleck-Dine scenario for baryogenesis [4].

Most discussions of the coherent production of scalar fields assume the potential along flat directions arises from the same SUSY breaking responsible for the mass splitting among the standard model fields in the present universe. The curvature of the potential would then be set by the gravitino mass, $V'' \sim m_{3/2}^2$. If this where the case the field would be highly overdamped for $H \gg m_{3/2}$, and only begin to oscillate when $H \sim m_{3/2}$. Here we observe that the finite energy density in the early universe induces a soft potential with curvature of order the Hubble constant, $V'' \sim H^2$. The flat directions are then always parametrically near critically damped, and efficiently evolve to an instantaneous minimum of the potential. For both the moduli problem and AD mechanism, this leads to a precise way of understanding the "initial conditions" for the amplitude of the fields when they begin to oscillate freely at $H \sim m_{3/2}$. In the case of the moduli problem, this suggests a possible solution if the minimum of the induced potential coincides with the true minimum. For the AD mechanism, it gives a much more complete understanding of the conditions for baryogenesis, namely a negative mass squared from the finite energy breaking. This permits an estimate of the asymmetry which systematically includes the effects of nonrenormalizable terms in the superpotential. The resulting asymmetry is largely independent of any assumptions about initial conditions.

2. Supersymmetry Breaking

The finite energy density in the early universe breaks supersymmetry. In a thermal phase this is manifest through the disparate occupation numbers for bosons and fermions. In an inflationary phase in which a positive vacuum energy dominates, the inflaton F or D component is necessarily nonzero, implying supersymmetry breaking. The same is true in the post-inflationary phase before reheating, when the inflaton oscillations dominate, and the time averaged vacuum energy is positive. In principle, SUSY breaking can be transmitted to flat directions by both renormalizable and nonrenormalizable interactions. However, for large field values all fields which couple through renormalizable interactions gain a mass larger than any relevant scale of excitation. These states then effectively decouple and do not lift the flat directions.

Nonrenormalizable interactions can have important effects though. To illustrate this consider a term in the Kahler potential of the form

$$\delta K = \frac{1}{M_p^2} \chi^\dagger \chi \phi^\dagger \phi \tag{1}$$

where χ is a field which dominates the energy density of the universe, ϕ is a canonically normalized flat direction, and $M_p = m_p/\sqrt{8\pi}$ is the reduced Planck mass. No symmetry prevents such a term, which can be present directly at the Planck scale, or be generated by running to a lower scale. If χ dominates the energy density, then $\rho \simeq \langle \int d^2\theta \chi^\dagger \chi \rangle$. In a thermal phase the expectation value arises from kinetic terms over the χ component thermal distributions. In an inflaton dominated era it is given by the inflaton F components and kinetic energy. The interaction (1) gives an effective mass for ϕ of $\delta \mathcal{L} = (\rho/M_p^2)\phi^\dagger\phi$ (note that a positive contribution in the Kahler potential gives a negative contribution to m^2). In a flat expanding background, $\rho = 3H^2 M_p^2$, so that $m^2 \sim H^2$. This is a generic result, independent of what specifically dominates the energy density. For $H \gtrsim m_{3/2}$, this source for the soft mass is more important than any hidden sector breaking.

In order to be concrete about the evolution along flat directions, we will assume an inflationary anzatz. In most models the correct magnitude of density and temperature fluctuations in the present universe is obtained for $H \sim 10^{13-14}$ GeV during inflation. In order to avoid the gravitino problem the reheat temperature after inflation can not (conservatively) be larger than about 10^9 GeV [5]. This implies that by the era of reheating $H \ll m_{3/2}$. With this restriction the induced potential discussed above is only important (ignoring an pre-inflationary evolution) during inflation and in the pre-reheating era dominated by inflaton oscillations. We therefore only need to consider the couplings of the inflaton to the flat directions.

Since the important couplings between the inflaton and flat directions arise from Planck scale operators, supergravity interactions should be included. The supergravity scalar potential is

$$V = e^{K/M_p^2} \left(D_i W K^{i\bar{j}} D_{\bar{j}} W^* - \frac{1}{3M_p^2}|W|^2 \right) + \frac{1}{8} f_{ab}^{-1} D^a D^b \tag{2}$$

where $D_i W \equiv W_i + K_i W/M_p^2$, $W_i \equiv \partial W/\partial \varphi_i$, $K^{i\bar{j}} \equiv (K_{i\bar{j}})^{-1}$, f_{ab} is the gauge kinetic function. $W(\varphi)$ and $K(\varphi^\dagger, \varphi)$ and the superpotential and Kahler potential, $D^a \equiv K_\varphi T^a \varphi$, where φ includes in general the flat directions, inflaton(s), and hidden sector. If the inflaton potential arises from F terms, the term in parenthesis has positive expectation value and a nontrivial potential along flat directions arises. Even if D terms dominate the inflaton potential, with nontrivial Kahler potential couplings (such as (1)) a potential results. The general form for the induced potential from (2) along an exact flat direction is of the form

$$V(\phi) = H^2 M_p^2 \ f(\phi/M_p) \tag{3}$$

where f is some function. Notice that the curvature is set by the Hubble constant, $V'' \sim H^2$, and the scale for variations in the potential is M_p. The general lesson is that in the early universe, when $H \gg m_{3/2}$, the characteristic scale for soft parameters is of order the Hubble constant [6].

In the rest of this letter we describe some of the consequences of this observation for the moduli problem and AD mechanism of baryogenesis. In a forthcoming paper, we will present a much more detailed discussion of these issues, with particular attention to the computation of the baryon asymmetry [7].

3. Moduli

The coherent production of string moduli leads to the string version [2,3] of the Polonyi problem [1]. The late decay of such a condensate can lead to a number of cosmological problems, including modification of the light element abundances. During inflation the moduli evolve in the potential (3) with $H \sim$ constant. Since the fields are parametrically close to critically damped, within a few e-foldings they are driven to a local minimum of the potential (up to quantum deSitter fluctuations) [8,9]. This is in contrast to the usual assumption that "scalars are not diluted during inflation." However, the form of the potential does not necessarily coincide with that after inflation, or from hidden sector SUSY breaking. In general the minima are separated by $\mathcal{O}(M_p)$. Once $H \sim m_{3/2}$ the moduli then start to oscillate freely about a true minimum with amplitude of $\mathcal{O}(M_p)$ [7]. This just gives a concrete realization of the initial conditions for the moduli problem by specifying the field for $H \gtrsim m_{3/2}$.

The present discussion suggests a possible solution of the moduli problem [6]. If the minima coincide at early and late times the moduli are driven to the true minimum during inflation. One possibility under which the minima can coincide occurs if there are no Kahler potential couplings between the moduli and either the inflaton or hidden sectors. If there is no SUSY preserving nonperturbative superpotential generated on moduli space, then the potential arises from supergravity interactions coupling the moduli to F components in the inflaton or hidden sectors. Every minimum of the moduli Kahler potential then coincides with a minimum of the potential at both early and late times [10].

A less restrictive circumstance under which the minima might coincide is if there is a point of enhanced symmetry on moduli space [11]. The potential is necessarily an extremum about such points since the moduli transform under some symmetry. Such enhanced symmetry points are familiar in string theory. In many string compactifications, there are points in the moduli space where all of the moduli, with the notable exception of the dilaton, transform non-trivially under some enhanced symmetry.

An example of this phenomenon is provided by the Z orbifold [12]. This orbifold is usually described by taking a product of three two dimensional tori. In this construction, the resulting theory has a variety of symmetries including an $SU(3)$

gauge symmetry and two Z_3 R symmetries. All the moduli are charged under some of these symmetries except those which describe the three two-dimensional tori. At special points in the moduli space there are further enhanced symmetries under which these remaining moduli, with the exception of the dilaton, are charged.

It might be that the true ground state of string theory is near such a point of enhanced symmetry. Alternately, some or all of these symmetries might be broken by small $\mathcal{O}(m_{3/2})$ vev's of other fields. The main problem with this idea is the dilaton. It is not known if such an enhanced symmetry exists for this field, and even if it does, it is likely to lie at a point where the gauge coupling is extremely large. So if symmetries are the solution of the moduli problem, the dilaton must be on a different footing than the other moduli. For example, the dilaton mass might arise from dynamics which does not break supersymmetry. The serious difficulties which such an idea must face have been discussed in ref. [3]. The possibility also exists to solve the moduli problem with a late inflation [10]. However unless $H \ll m_{3/2}$ the minimum may be shifted as for standard inflation.

4. Baryogenesis

In the MSSM, at the level of renormalizable operators, there are numerous flat directions in the space of scalar fields. Most of these involve squarks or sleptons and carry B and/or L. A simple example is the direction [13]

$$H_u = \begin{pmatrix} \phi \\ 0 \end{pmatrix} \qquad L = \begin{pmatrix} 0 \\ \phi \end{pmatrix} \tag{4}$$

where ϕ parameterizes the flat direction. In the original discussion of ref. [4] it was assumed that directions such as this were exactly flat in the supersymmetric limit [14] and that ϕ was initially $\mathcal{O}(M_{GUT})$ or $\mathcal{O}(M_p)$. For $H \lesssim m_{3/2}$ the field would begin to oscillate about the true minimum at $\phi = 0$. In addition to the B and L conserving terms, the soft potential was assumed to contain B and/or L violating dimension four terms suppressed by $m_{3/2}^2/M_p^2$. As a result, the coherently oscillating field develops a large baryon number. The subsequent decay of the condensate then gives a substantial (even enormous) baryon asymmetry [4,13,15].

With the inclusion of nonrenormalizable terms in the superpotential, and the induced soft potential, the scenario for AD baryogenesis is very different. Nonrenormalizable terms in the superpotential, if present, will lift flat directions even in the supersymmetric limit. These can take the form

$$\delta W = \frac{\lambda}{nM^{n-3}}\phi^n \tag{5}$$

where M is some large mass scale such as the GUT or Planck scale. For the LH_u example given above the lowest order term of this form, assuming R parity, is $\frac{\lambda}{M}(LH_u)^2$. The power law growth in the potential from these terms limits the fields to be parametrically less than M_p (even for $M \sim M_p$). In addition, A terms, proportional to ϕ^n, can result from cross terms in (2) and higher order terms in the Kahler potential. In light of our discussion of early universe SUSY breaking, the scalar potential for $H \gg m_{3/2}$ then has the form [6]

$$V(\phi) \simeq cH^2|\phi|^2 + \frac{a\lambda H\phi^n}{nM^{n-3}} + |\lambda|^2\frac{|\phi|^{2n-2}}{M^{2n-6}} \quad . \tag{6}$$

where c and a are constants of $\mathcal{O}(1)$. The A term has the important effect of violating B or L and has a definite CP violating phase relative to ϕ.

With minimal Kahler potential, the coefficient c arising from (2) is positive ($c = 3$ during inflation for F type inflaton breaking). The flat direction is then driven exponentially quickly to the origin during inflation. Quantum deSitter fluctuations give $\langle\delta\phi^2\rangle \sim H^2$, but with a correlation length of $\mathcal{O}(H^{-1})$ [17]. Any resulting baryon number then averages to zero over the present universe [18]. In addition the relative magnitude of the B violating term in (6) is small for $H \ll M$.

A non-negligible baryon number can result if the B violating term in (6) has the same magnitude as the B conserving terms. This will occur if $c < 0$. This is perfectly possible for suitable choices of the Kahler potential; no fine-tuning is required. In this case the minimum of the potential, ignoring for the moment the A term, is given by

$$|\phi_0| = \left(\frac{\sqrt{-c}HM^{n-3}}{(n-1)\lambda}\right)^{\frac{1}{n-2}} \quad . \tag{7}$$

Inclusion of the contribution of the A term does not substantially change the magnitude of the minimum, but does give n discrete minima for the phase of ϕ. During inflation if $|c|$ is not too small, the system quickly settles into one of the minima. The observable universe is then left with a single value of the intial phase of ϕ. After inflation, H changes with time as in a matter dominated universe and ϕ_0 decreases. A straightforward analysis of the equations of motion in this era indicates that for $n \geq 4$, the field oscillates about a point slightly larger than $\phi_0(t)$ where $V''(\phi) \sim H^2$. Thus when $H \sim m_{3/2}$, $\langle\phi\rangle \sim \phi_0$. At this time the soft potential from hidden sector SUSY breaking becomes important. The A term from this source is comparable in magnitude to the other terms in the potential (as may be seen by simply plugging ϕ_0 into eq. (6)) and in general has a different phase than any arising from coupling to the inflaton. The B or L violation is therefore maximal during the epoch at which the field begins to oscillate freely, thereby imparting a substantial asymmetry to the condensate. The resulting baryon number per condensate particle is near maximal, $n_b/n_\phi \sim \mathcal{O}(10^{-1})$ (if the relative phases are $\mathcal{O}(1)$). Notice that this is *independent* of λ/M. Once $H \ll m_{3/2}$ the field value decreases and the relative importance of the A term is reduced. The baryon number imparted to the condensate is therefore conserved in this epoch. This scenario has been checked by numerical integration of the equations of motion [7].

The baryon to entropy ratio depends on the total density in the condensate, and the inflaton reheat temperature , T_R. The flat direction ϕ begins to oscillate freely when the coherent oscillations of the inflaton still dominate the total energy density, ρ_I. Since $\rho_\phi \sim m_{3/2}^2 \phi_0^2$, the fractional energy in the condensate is

$$\frac{\rho_\phi}{\rho_I} \approx \left(\frac{m_{3/2} M^{n-3}}{\lambda M_p^{n-2}} \right)^{2/(n-2)} . \tag{8}$$

Notice that ρ_ϕ is larger for smaller (λ/M^{n-3}). After the inflaton decays the baryon to entropy ratio is then

$$\frac{n_b}{s} \approx \frac{n_b}{n_\phi} \frac{T_R}{m_\phi} \frac{\rho_\phi}{\rho_I}. \tag{9}$$

This estimate is insensitive to the details of the decay of the AD flat direction, as long as it has nonzero $B - L$. The baryon to entropy ratio depends mainly on T_R and the order at which the flat direction is lifted. For T_R just below the gravitino bound and $M \sim M_p$, n_b/s is too large for $n \geq 6$. However for the LH_u direction with $n = 4$, after sphaleron processing of the resulting lepton number, $n_b/s \sim 10^{-10}(T_R/10^9\mathrm{GeV})(M/\lambda M_p)$. This is a quite reasonable range. At low energies the operator $\frac{\lambda}{M}(LH_u)^2$ which lifts this flat direction gives rise to neutrino masses. 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, $n_b/s \sim 10^{-10}(T_R/10^9\,\mathrm{GeV})(10^{-5}\mathrm{eV}/m_\nu)$. The total baryon density in the condensate grows rapidly with n; only the LH_u direction gives a reasonable result without additional entropy releases after inflaton decay.

5. Conclusions

In summary the large supersymmetry breaking in the early universe gives a precise realization of the "initial conditions" (when $H \sim m_{3/2}$) along flat directions. It seems quite difficult to solve the moduli problem unless there are symmetries which ensure that the high energy and low energy potentials possess the same minimum. We have seen that (much to the suprise of some of the authors) the AD mechanism is quite robust. Provided that the curvature of the induced ϕ potential at the origin is negative for $H \gg m_{3/2}$, a desirable value for n_b/s results for the LH_u direction when account is taken of higher dimension operators. More detail about the evolution of the fields, other standard model flat directions, the possible sources of supersymmetry breaking, and the decay of the condensate will be presented in ref. [7].

REFERENCES

1. C. Coughlan, W. Fischler, E. Kolb, S. Raby, and G. Ross, Phys. Lett. B **131** (1993) 59; J. Ellis, D.V. Nanopoulos and M. Quiros, Phys. Lett. B **174** (1986) 176.
2. R. de Carlos, D.V. Nanopoulos, and M. Quiros, Phys. Lett. B **318** (1993) 447; T. Banks, D. Kaplan, and A. Nelson, Phys. Rev. D **49** (1994) 779.
3. T. Banks, M. Berkooz, and P.J. Steinhardt, preprint RU-94-92, hep-th 9501053.
4. I. Affleck and M. Dine, Nucl. Phys. B **249** (1985) 361.
5. J. Ellis, A. Linde, and D. Nanopoulos, Phys. Lett. B **118** (1982) 59; D. Nanopoulos, K. Olive, and M. Srednicki, Phys. Lett. B **127** (1983) 30; W. Fischler, Phys. Lett. B **332** (277) 1994.
6. M. Dine, L. Randall, and S. Thomas, Talk presented at the *US-Polish Workshop on Physics from the Planck Scale to Electroweak Scale*, September 1994; and at the *DESY Theory Workshop on Supersymmetry*, October 1994.

7. M. Dine, L. Randall, and S. Thomas, to appear.
8. The fact that during inflation fields generically have masses of order H has been noted by many authors. In the context of the Polonyi problem, it was noted by M. Dine, W. Fischler, and D. Nemeschansky, Phys. Lett. B **136** (1984) 169. In another context, this was noted by O. Bertolami and G. Ross, Phys. Lett. B **183** (1987) 163.
9. G. Dvali, preprint IFUP-TH 09-95, hep-ph 9503259.
10. L. Randall and S. Thomas, preprint MIT-CTP-2331, SCIPP 94-16, hep-ph 9407248.
11. In the context of the Polonyi problem, M. Dine, W. Fischler, and D. Nemeschansky, Phys. Lett. B **136** (1984) 169 pointed out that symmetries could give minima which coincide at low and high energies.
12. L. Dixon, J. A. Harvey, C. Vafa, and E. Witten, Nucl. Phys. B **261** (1985) 678.
13. H. Murayama and T. Yanagida, Phys. Lett. B **322** (1994) 349.
14. This can in fact be guaranteed by a discrete Z_N R symmetry. Under such a symmetry the superpotential transforms as $W \to e^{4\pi i/N} W$. If, for example, ϕ transforms as $\phi \to \phi$, then no terms of the form ϕ^n are allowed. It is possible to forbid terms of the form $\chi\phi^n$ by symmetries as well; M. Dine and N. Seiberg, Nucl. Phys. B **160** (1985) 243.
15. A. Linde, Phys. Lett. B **160** (1985) 243; J. Ellis, D. V. Nanopoulos, Phys. Lett. B **184** (1987) 37; K. Enqvist, K.. W. Ng, and K. Olive, Phys. Rev. D **37** (1988) 2111; S. Davidson, H. Murayama, and K. Olive, Phys. Lett. B **328** (1994) 354.
16. J. Ellis, K. Enqvist, D. V. Nanopoulos, and K. Olive, Phys. Lett. B **191** (1987) 343; K. W. Wong, Nucl. Phys. B **321** (1989) 528.
17. T. Bunch and P. Davies, Proc. R. Soc. A **360** (1978) 117; A. Linde, Phys. Lett. B **116** (1982) 335; A. Starobinsky, Phys. Lett. B **117** (1982) 175; A. Linde, Phys. Lett. B **131** (1983) 330.
18. The Kahler potential might be tuned to give $m^2 \ll H^2$ but positive. In that case large fluctuations can develop with a correlation length of order $H^{-1} e^{3H^2/2m^2}$ [15,17]. To be coherent over a scale large compared to the current horizon size requires $(H/m)^2 \gtrsim 40$. However if $m^2 \sim H^2$ after

inflation, the field is driven toward the origin as a power law in time. So an entirely positive m^2 scenario must fine tune the mass even after inflation.