On Preserving a $B + L$ Asymmetry Produced in the Early Universe

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Abstract
One of the most efficient mechanisms for producing the baryon asymmetry of the Universe is the decay of scalar condensates in a SUSY GUT as was first suggested by Affleck and Dine. We show that given a large enough asymmetry, the baryon number will be preserved down to low temperatures even if $B - L = 0$, because the baryon number carrying scalars form bose condensates that give the $W$ a mass. We derive the conditions on the condensate needed to suppress electroweak sphaleron interactions which would otherwise drive the baryon asymmetry to zero when $B - L = 0$.

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*This work was supported in part by DoE grants DE-FG02-94ER-40823 and DE-AC03-76SF00098, by NSF grant AST-91-20005 and by NSERC.

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Models for baryogenesis at the weak scale and above have been under considerable scrutiny since the realization that baryon number violation due to non-perturbative electroweak interactions is rapid at high temperatures \(^1\). Baryon number violating interactions mediated by sphalerons in fact violate \(B + L\) while conserving \(B - L\). Thus any model of baryogenesis above the weak scale which also conserves \(B - L\) is likely to produce a negligible baryon asymmetry once sphaleron interactions have been incorporated \(^2\). This conclusion is clearly important as the simplest models of baryogenesis typically conserve \(B - L\). There are, however, several modifications to these simplest scenarios which can overcome this problem. One possibility is that the baryon asymmetry is generated during the electroweak phase transition \(^3\), though it is unlikely that this can be done within the standard model. It is possible that a more complicated grand unified theory which violates \(B - L\), (such as \(SO(10)\) rather than \(SU(5)\)) produces a baryon asymmetry which is only reshuffled by sphaleron interactions. Or, an extension of the standard model which violates lepton number and hence \(B - L\) (for instance the inclusion of right handed neutrinos used to generate neutrino masses via the see-saw mechanism \(^4\)), can produce a lepton asymmetry which is transformed into a baryon asymmetry at the weak scale \(^5\), \(^6\), \(^7\).

Alternatively, the conditions under which sphalerons are able to destroy a prior baryon asymmetry have come under consideration. For example, lepton mass effects in the presence of lepton flavour asymmetries and \(B + L\) violating electroweak interactions can generate a baryon asymmetry even though \(B + L = B - L = 0\) \(^8\), \(^9\). It has also been shown that because of the very small value of the electron Yukawa coupling constant, the equilibrium condition that leads to \(B = L = 0\) is only valid at late times (temperatures \(T \lesssim 1\ \text{TeV}\) close to the electroweak phase transition \(^10\), \(^11\)). At higher temperatures the decoupling of the right-handed electron in some sense safe-guards the baryon asymmetry. Though in the standard model sphalerons win in the end, the destruction of the baryon asymmetry is exponentially sensitive to parameters of the model. In an extension of the standard model the baryon asymmetry may yet survive. There are also attempts to protect baryon asymmetry introducing new fields \(^12\).

In this letter, we explore another possibility. We consider the effect of scalar Bose–Einstein (B-E) condensates on the sphaleron interaction rate in a \(B - L\) conserving supersymmetric grand unified theory. The presence of a B-E condensate gives rise to gauge boson masses, which, if they persist down to temperatures close to the electroweak phase transi-
tion, can in principle suppress sphaleron interactions sufficiently throughout the history of the universe to protect the baryon asymmetry.

Let us first recall some facts about baryogenesis in grand unified theories. The simplest mechanism is the out-of-equilibrium decay of heavy gauge or Higgs fields \[13\]. The same is generally true in a supersymmetric gut\[\ast\]. However, in the framework of inflationary cosmology \[12\], where generic models of inflation are constrained by the observed anisotropies in the microwave background \[13\], these out-of-equilibrium decay scenarios are also constrained \[1\]. In what follows below, we consider only a generic inflationary model in which the duration of inflation and the magnitude of density perturbations are determined by a single mass scale \(\mu\). We will assume that the inflaton potential can be written as 

\[ V(\psi) = \mu^4 P(\psi) \]

where \(\psi\) is the scalar field driving inflation, the inflaton, and \(P(\psi)\) is a function of \(\psi\) which possesses the features necessary for inflation, but contains no small parameters. The COBE data fixes the mass scale \(\mu\) relative to the Planck scale to be roughly

\[ \frac{\mu^2}{M_P^2} \simeq \text{few} \times 10^{-8} \tag{1} \]

Fixing \(\mu^2/M_P^2\) has immediate general consequences for inflation \[17, 6\]. For example, the Hubble parameter during inflation, \(H^2 \simeq (8\pi/3)(\mu^4/M_P^2)\) so that \(H \sim 10^{-7} M_P\). The duration of inflation is \(\tau \simeq M_P^3/\mu^4\), and the number of e-foldings of expansion is \(H \tau \sim 8\pi(M_P^2/\mu^2) \sim 10^9\). If the inflaton decay rate goes as \(\Gamma \sim m_\psi^3/M_P^2 \sim \mu^6/M_P^5\), the universe recovers at a temperature \(T_R \sim (\Gamma M_P)^{1/2} \sim \mu^3/M_P^2 \sim 10^{-11} M_P \sim 10^8 \text{GeV}\).

The relatively low inflaton mass \(m_\psi \sim \mu^2/M_P \lesssim 10^{12} \text{GeV}\) is problematic for the out-of-equilibrium decay in a susy gut. The low mass scale would require at least one baryon number violating gauge or Higgs boson with a smaller mass causing proton decay at experimentally disallowed rates. There is a natural alternative to the out-of-equilibrium decay scenario in a supersymmetric GUT, which is the decay of sfermion condensates as first proposed by Affleck and Dine \[18\]. In this scenario, squark and slepton fields obtain large vacuum expectation values along flat directions of the scalar potential. Effective baryon number violating operators induce a baryon asymmetry in coherent flat direction oscillations. The decay of these flat directions, due to supersymmetry breaking effects characterized by a scale \(\tilde{m}\), will produce a potentially large baryon asymmetry if \(C\) and \(CP\) are violated (explicitly or spontaneously) in that sector. If the oscillations of the flat direction, \(\phi\), come to dominate

\[\ast\text{See Ref. 14 for a recent analysis in minimal SUSY SU}(5).\]
the energy density of the Universe, indeed a large baryon to entropy ratio, \( n_B/s \sim O(1) \) is expected \([18, 19]\). It was already pointed out in the original paper that B-E condensates of the squarks and sleptons may form, and Dolgov and Kirilova \([20]\) showed that the B-E condensate may persist down to relatively low temperature in a Universe where the energy density is dominated by the flat direction. We will show in this letter that in an inflationary Universe, the B-E condensate can persist down to the electroweak phase transition for \( n_B/s \sim O(10^{-2})–O(1) \). Then the sphaleron transitions are effectively killed, and the original \( B + L \) asymmetry generated by the decay of the sfermions along the flat direction is preserved.

Note that this relatively large \( n_B/s \) can be naturally obtained within the inflationary cosmology and Affleck-Dine scenario. In the context of inflation, it is quite likely that the Universe is dominated by the inflaton and the radiation products of inflaton decay rather than the sfermion oscillations. Inflation, in fact, offers a natural explanation for the large initial value of the sfermion fields. During inflation, quantum fluctuations drive massless scalar fields as \( \langle \phi^2 \rangle = H^3 t/4\pi^2 \) \([21]\), which gives \( \phi^2_o \sim \mu^2 \sim m_\psi M_P \) at the end of inflation. In this case the baryon asymmetry produced is somewhat lower as it is diluted by the entropy produced by inflaton decay. The result is \([17]\)

\[
\frac{n_B}{s} \simeq \frac{\epsilon \phi^4_o m_\psi^{3/2}}{M_X^2 M_P^{5/2} \tilde{m}}
\]

where \( \epsilon \) is a combination of coupling constants whose value parametrizes the CP-violation, \( \phi_o \) is the initial sfermion value, which is determined by quantum fluctuations during inflation, and \( M_X \simeq 10^{-4}–10^{-3} M_P \) is the unification scale. For \( \tilde{m} \sim 10^2–10^3 \) GeV, and \( m_\psi \sim 10^{11}–10^{12} \) GeV required from COBE data \([17, 6]\), \( n_B/s \) ranges from \( 10^{-6} \epsilon \) to \( \epsilon \), making this scenario very attractive. Note further that \( \epsilon \) can be \( O(1) \) since the flat direction vacuum expectation value (vev) can break CP spontaneously and the vev is the same over the whole universe due to inflation. So a relatively large value of \( n_B/s \) can be naturally obtained within this scenario.

In a minimal SUSY GUT where \( B - L \) is conserved, this asymmetry can be destroyed by sphaleron interactions. (However, the Affleck-Dine scenario itself does extend to incorporate \( B - L \) violation \([22, 23]\). Here we are interested in a \( B - L \) conserving theory, and we will see whether the sphaleron interactions are indeed present when the asymmetry takes near
maximal values. The sphaleron interaction rate is roughly
\[ \Gamma_{sph} \sim \frac{10^{-2} m_W^7}{\alpha_W^3 T^6} e^{-E_{sph}/T}, \] (3)
where \( E_{sph} \sim 4m_W/\alpha_W \) and the numerical value of the prefactor depends on the gauge and higgs couplings. (Here we took \( \lambda = g^2 \) since this is a supersymmetric theory.) In the presence of a condensate, \( m_W \neq 0 \) and the sphaleron rate is exponentially suppressed. We will see that a condensate is likely to be present at lower temperatures if the degeneracy due to a baryon asymmetry is high, hence we will assume \( n_B/s \sim O(1) \), and discuss later how large \( n_B/s \) needs to be.

In Ref. [17], parameters were chosen so that \( n_B/s \) was relatively small, i.e. so that no additional entropy generation was required (more on this below). In that case a relatively simple sequence of events was found to occur: Initially (after the exponential expansion due to inflation), the inflaton begins to oscillate at a value of the cosmological scale factor \( R = R_\psi \). Soon afterwards, the sfermions start oscillating along the flat direction at \( R = R_\phi \sim (m_\psi/\tilde{m})^{2/3} R_\psi \). The inflaton oscillations decayed first at \( R = R_d\psi = (M_P/m_\psi)^{4/3} R_\psi \) followed by the decay of the sfermion oscillations at \( R = R_d\phi = (m_\phi^{7/15} \tilde{\phi}^{2/5} M_P^{2/15} / \tilde{m}) R_\psi \). For a given value of \( R \), the number density of the asymmetry can be written as
\[ n_B = \frac{\epsilon \phi_0^4 m_\psi^2}{M_X^2 \tilde{m}} \left( \frac{R_\psi}{R} \right)^3. \] (4)
At this point the radiation from inflaton decays has not yet had enough time to thermalize, and equilibrium was established at the slightly later time corresponding to \( R = R_T \sim R_d\psi/\alpha^2 \), where \( \alpha \) is a gauge fine structure constant. The temperature of the Universe now is well defined and takes the value \( T \sim m_\psi^{3/2} \alpha^2/M_P^{1/2} N_T^{1/3} \lesssim 10^5 \text{ GeV} \), where \( N_T \sim 200 \) is the number of particle degrees of freedom at \( T \).

A large baryon asymmetry will make both quantitative as well as qualitative changes to this picture, because the presence of the flat direction or B-E condensate gives mass to the gauge and Higgs multiplets and the interactions are suppressed as \( \sigma \sim (\alpha^4/\pi^2)T^2/\tilde{\phi}^4 \) where \( \phi \) is the typical amplitude of either the flat direction or the B-E condensate. In particular,

\[ \text{We refer the readers to Refs. [24, 25, 26] on discussions of the symmetry breakdown in the presence of high number asymmetry.} \]
the history before thermalization may be drastically altered. However, it does not affect the baryon asymmetry, as we will see below.

Let us follow the evolution before thermalization for the sake of completeness. For simplicity, we will assume that the sfermion condensate decays directly only to fermions. Prior to thermalization, chemical equilibrium between sfermions and fermions is not reached, and the decay to fermions is Pauli blocked. Consider the decay rate of a scalar $\phi$ to two massless fermions,

$$\Gamma_{\phi} = \frac{1}{2\tilde{m}} \int \frac{d^3p_1}{(2\pi)^32E_1} \frac{d^3p_2}{(2\pi)^32E_2} (2\pi)^4 \delta^4(\Sigma p)|\mathcal{M}|^2(1 - f_1)(1 - f_2)$$

where $|\mathcal{M}|^2 = 2(\tilde{m}^2/\phi^2)F(g_i, h_j)(p_1 \cdot p_2)$; $p_1, p_2$ are the outgoing fermion four-momenta, and $F(g_i, h_j)$ is a function of gauge and Yukawa coupling constants in both the numerator and the denominator. We assume $F(g_i, h_j) \sim 1$. The factor $(\tilde{m}^2/\phi^2)$ accounts for the suppression due to masses of the fields directly coupled to $\phi$. $f = (\exp((p - \mu)/T) + 1)^{-1}$ is the fermion momentum distribution with chemical potential $\mu$. At $T = 0$, $f = 0$ for $E = p \geq \mu$ and $= 1$ otherwise. Thus,

$$\Gamma_{\phi} = \begin{cases} \frac{\tilde{m}^3}{16\pi^3\mu^2} & \mu < \frac{\tilde{m}}{2} \\ 0 & \mu \geq \frac{\tilde{m}}{2} \end{cases}$$

and the baryon density contained in fermions is

$$n^{(f)}_B = \frac{g_f}{3\pi^2} \mu^3$$

where $g_f$ is the number of fermion species sharing the baryon asymmetry. Clearly, $n^{(f)}_B \ll n_B$ for $\mu = \tilde{m}/2$ since $\tilde{m} \ll n_B^{1/3}$. The decay of the sfermion condensate is therefore delayed.

We expect the Universe to thermalize at a similar scale factor in the presence of the condensate as in the model considered in [17]. Subsequent to thermalization, sfermion decays occur in a thermal bath, and the decay to fermions proceeds rapidly.

During the decay, equilibrium between fermions and scalars will be achieved and the equilibrium condition $\mu_F = \mu_B = \tilde{m}$ will insure that the baryon number flows back to scalars. However, for a highly degenerate boson gas at low temperature (low with respect to thermalization), the sphaleron transitions are not present prior to thermalization, since collective excitations like the sphaleron can be formed only in a dense plasma, while there is only a dilute gas of high energy particles before thermalization.

This may happen if the flat direction $\phi$ corresponds to the lightest scalar field, making the decays to scalars kinematically forbidden. As we will see below, our results are independent of this assumption.

This is true for the flat direction in Ref. [23], where all gauge interactions are killed by the flat direction.
a critical temperature we derive shortly), a B-E condensate will develop. The total baryon asymmetry may then be written as

\[ n_B = n_B^{(f)} + n_B^{(b)} + n_B^{(c)} \]  

(8)

where the superscripts denote contributions in fermions \( f \), \( p \neq 0 \) bosons \( b \) and a B-E condensate \( c \). For \( T \gg \mu = \tilde{m} \) it is straightforward to expand in \( \mu/T \), neglecting \( \tilde{m} \), to compute \( n_B^{(f)} \) and \( n_B^{(b)} \)

\[ n_B^{(f)} = \frac{g_f}{6} \mu_f T^2 \]  

(9)

\[ n_B^{(b)} = \frac{g_b}{3} \mu_b T^2. \]  

(10)

We can define a (scale factor dependent) critical temperature \( T_c \) below which a B-E condensate will be present by

\[ n_B^{(b)} + n_B^{(c)} = \frac{g_b}{3} \tilde{m} T_c^2 \]  

(11)

If the temperature of the Universe is below \( T_c \), the baryon number in the B-E condensate will be

\[ n_B^{(c)} = \frac{g_b}{3} \tilde{m} \left( 1 - \left( \frac{T}{T_c} \right)^2 \right) T_c^2 \]  

(12)

So far we have followed the history until thermalization. Now we follow the evolution after the thermalization. At \( R > R_T \), when the decay is complete, and assuming equilibration, \( n_B^{(f)} \ll n_B \) (since \( \mu T^2 \ll T^3 \sim n_B \)) and \( n_B^{(b)} = 2n_B^{(f)} \) so the baryon number is now almost entirely in the form of the B-E condensate. One sees now that the assumption of the sfermion decay to fermions is not crucial as decay into \( (p \neq 0) \) bosons would not be possible if \( T \) is below \( T_c \). The critical temperature evaluated at \( R_T \) is (from eqns (11) and (4) with \( \epsilon = 1 \))

\[ T_c \simeq (3n_B/g_b\tilde{m})^{1/2} \sim (3/g_b)10^7 \text{ GeV}. \]

Notice that the presence of the B-E condensate at this late stage (\( T \sim 10^5 \text{ GeV} \)) depends crucially on a large baryon asymmetry. From eq. (11) one sees that while \( T_c \sim R^{-3/2} \) (because \( n_B \) scales like \( R^{-3} \)), \( T \sim R^{-1} \), so as the Universe expands and cools, it will eventually be at a temperature above the critical temperature, and the B-E condensate will evaporate. The phase transition will occur at a scale \( R_F \sim 10^3 R_T \) or at a temperature \( T_F \sim 100 \text{ GeV}, below the normal electroweak phase transition temperature \( \equiv T_{EPT} \). (The presence of the B-E condensate could affect the electroweak phase transition, \footnote{We treat the scalar particles as free bose gas. Then \( \mu = \tilde{m} \) is the maximum possible chemical potential.}
because $SU(2)$ is broken above the EPT, and the scalar vacuum expectation value (vev) will give masses to most particles. However, if the B-E condensate evaporation temperature $T_c$ is below $T_{EPT}$, it should be safe to assume that there is always a vev present.)

We would like to estimate how large a baryon to entropy ratio is required to make this scenario work. Although the Affleck-Dine flat direction consists of vevs for only a small number of fields, we expect the equilibrium B-E condensate at low temperatures to share the baryon number equally among all the squarks (assuming that they all have similar masses), and the lepton number $L_i$ equally among the $i$th generation sleptons. For convenience, let us assume that all the lepton number is in muons (as it would be in Affleck and Dine’s flat direction), in which case the slepton vevs will be slightly larger than those of the squarks, because there are fewer of them among whom to share the asymmetry. We require that the three B-E-condensed scalars carrying muon number remain as large as $\langle \phi_i \rangle \sim T/3$ at $T_{EPT} \sim 300$ GeV (more later on this choice of $\langle \phi_i \rangle$), so that

$$n_\mu \simeq \tilde{m} T_{EPT}^2 + \frac{1}{2} \tilde{m} T_{EPT}^2 + 3 \tilde{m} T_{EPT}^2 \frac{T}{9} \tag{13}$$

Taking $s = 2 \pi^2 N T^2 / 45$ gives

$$\frac{n_B}{s} \simeq \frac{n_\mu}{s} \gtrsim \frac{45 \tilde{m}}{\pi^2 N T_{EPT} T_{EPT}} \simeq 0.01 \tag{14}$$

where we have used $\tilde{m} = 100$ GeV, $T_{EPT} = 300$ GeV, and $N_{T_{EPT}} = 200$. If we assume that there is no entropy generation between $R_{EPT}$ and $R_T$, then to preserve a B-E condensate down to below the EPT we need $n_B / s \gtrsim 0.01$. As we discussed previously, this should be possible.

We have shown that given a large enough baryon asymmetry, a scalar condensate will persist down to very low temperatures. This same conclusion was found in another context by Dolgov and Kirilova \[20\]. However in order to protect the baryon asymmetry against wash-out from sphaleron interactions, we need to still show that the sphaleron interaction rate (eq. 3) is supressed. If we define $v = \sqrt{2 \sum_i |\phi_i|^2}$, where $\phi_i$ are the B-E condensate vevs carrying $SU(2)$ quantum numbers, thus $m_W = g v / 2$, and

$$\frac{\Gamma_{sph}}{H} \sim \frac{g v^7 M_p}{12 N T^{1/2} T^8} \exp \left[ -\frac{8 \pi v}{g T} \right] \tag{15}$$

**$\tilde{m}$ here is the slepton mass, and not necessarily equal to $\tilde{m}$ of equation (2), which is the overall supersymmetry-breaking mass.**
which is out-of-equilibrium at $T \simeq 300$ GeV for $v \gtrsim 0.8T$. So if we assume that the vevs corresponding to the lepton condensates are larger than $T/3$ near the electroweak phase transition, as we required in the previous paragraph, the sphalerons are clearly out of equilibrium.

Finally, we wish to comment on the problem of entropy production. Typically this is usually a problem concerning excess entropy. Here, we are in a perhaps more enviable situation requiring some entropy generation at or around the weak scale. There has been a considerable amount of discussion of excess entropy production in supergravity and superstring-inspired models [27]. This question was addressed specifically in the context of the AD mechanism for baryogenesis in [28]. There a constraint on intermediate scale models was derived to avoid the overproduction of entropy. A relaxation of that bound could possibly yield the entropy needed here.

In summary, we have shown that if the baryon to entropy ratio at thermalisation is large ($\gtrsim 0.01$) in a theory with $B - L = 0$, $B + L$ carrying bose condensates may keep the sphalerons out of equilibrium until the electroweak phase transition, and thereby prevent them from washing out the asymmetry.

Acknowledgements

We would like to thank Bruce Campbell and John March-Russell for useful discussions. This work was supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contracts DE-FG02-94ER-40823 and DE-AC03-76SF00098, by NSF grant AST-91-20005 and by NSERC. The work of KAO was in addition supported by a Presidential Young Investigator Award.

††Using the dilution factor $\Delta \sim 10^2 m_I^3/\tilde{m}^{5/2} M_\nu^{1/2}$ (Eq. (14) in [28]), one indeed obtains the desired $\Delta \sim 10^{7}$ for $m_I \sim 10^7$ GeV, corresponding to $n = 1$ in their Eq. (6).
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