Supersymmetric Baryogenesis

Submitted to Physics Letters B

A. G. Cohen $^a$

Physics Department
Boston University
Boston, MA 02215

A. E. Nelson$^b$

Department of Physics
9500 Gilman Drive 0319
University of California, San Diego
La Jolla, CA 92093-0319

Abstract

Requiring that the baryon number of the universe be generated by anomalous electroweak interactions places strong constraints on the minimal supersymmetric standard model. In particular, the electric dipole moment of the neutron must be greater than $10^{-27}$e-cm. Improvement of the current experimental bound on the neutron’s electric dipole moment by one order of magnitude would constrain the lightest chargino to be lighter than 88 GeV, and the the lightest neutralino to be lighter than 44 GeV. In extensions of this model with gauge singlet superfields all of these bounds are eliminated.

UCSD/PTH 92-32, BU-HEP-92-20

Email: cohen@andy.bu.edu, anelson@ucsd.edu

$^a$ DOE Outstanding Junior Investigator

$^b$ Sloan fellow, SSC fellow
Despite the success of the standard model of weak and strong interactions, we still remain ignorant of the mechanism of electroweak symmetry breaking; we are just beginning to probe this sector directly through Higgs searches at LEP. We are also still ignorant of the origins of the CP violation observed in the kaon system. Unfortunately, the minimal standard model accounts for all solid experimental results observed to date, and we have few experimental constraints on the symmetry breaking or CP violating aspects of the theory. There is a strong cosmological argument that the minimal standard model cannot be the whole particle physics story. The baryon to entropy ratio of the universe is \((0.4 - 1.0)10^{-10}\) \([1]\), and explaining this observable requires baryon number violation coupled with out-of-equilibrium CP violation in the early universe \([2]\). While the standard model does have CP violation, the effects of the CP violating phase in the Kobayashi-Maskawa matrix are too suppressed by small masses and mixing angles in order to account for the observed baryonic asymmetry of the universe (BAU) \([1]\). The standard model also contains anomalous baryon number violating interactions \([3]\), which should be rapid enough at high temperatures to affect cosmology \([7-9]\). Furthermore, the standard model can satisfy the out of equilibrium condition for baryogenesis if the phase transition is first order, proceeding via nucleation and expansion of bubbles of the broken phase \([8]\). Unfortunately it is necessary for the vev of the Higgs field after the transition to be large in order to avoid washing out any baryon number created during the transition \([3,10-14]\). In the minimal standard model this requirement cannot be satisfied unless the Higgs is lighter than \(\sim 35\) GeV \([12]\), which conflicts with current experimental bounds. Thus we should look beyond the standard model in order to discover the origin of the BAU. However we may not need to look very far, as possible mechanisms for baryogenesis have been suggested in several reasonable extensions such as axion models with additional light scalar doublets \([15]\), the singlet Majoron model \([16]\), the two Higgs model \([17,19]\), the supersymmetric standard model \([18,19]\), extended supersymmetric models \([20]\) and left-right symmetric models \([21]\).

Requiring that sufficient BAU be generated during the weak transition, and that it not be washed out afterwards can give us new information about the CP violating and symmetry breaking sectors of the weak interactions, allowing us to rule out some models (such as

\[1]\ Shaposhnikov has suggested two conceivable ways to enhance the CP violation in the standard model at high temperature \([3]\); the first mechanism, dynamical high temperature spontaneous CP violation, is contradicted by non-perturbative computation \([1]\), and the second mechanism, reflection of baryon number from expanding bubble walls, according to our estimates cannot provide a large enough asymmetry \([3]\).
the minimal standard model), and to constrain parameters in others. It is the aim of this letter to use baryogenesis to find new constraints on supersymmetric models. We are able to severely constrain almost all unknown parameters of the minimal model, while considerable freedom remains for models with additional gauge singlets.

First let us consider baryogenesis in the minimal supersymmetric standard model (MSSM) [22]. Several authors have claimed that this model is ruled out for baryogenesis because the bound on the lightest scalar mass in this model is the same as in the minimal standard model [11,13]. However Myint [14] finds that the bounds on the scalar mass are relaxed somewhat, to 64 GeV, due to top quark and squark corrections to the high temperature effective potential when the top quark is heavy and the squark masses are not too heavy. The bound is relaxed because top squarks can play the role envisioned by Anderson and Hall for a gauge singlet scalar, whose coupling to the Higgs doublet increases the strength of the transition [23]. The upper bound on the Higgs mass was computed in one loop perturbation theory and higher order corrections to the gauge propagator will reduce this number by a factor of about \(\sqrt{2/3}\) [12], to \(\sim 50\) GeV. Furthermore this bound will receive corrections proportional to \((m_H/m_Z)^2\) from two loop corrections. Thus avoiding baryon number washout is difficult in the MSSM, but not impossible. To increase the upper bound on the Higgs mass as much as possible, one should take the top quark to be heavy (\(> 150\) GeV), the squark masses not much heavier than 150 GeV, the trilinear soft supersymmetry breaking terms ("A-terms") small, and the parameter \(\tan \beta\) (the ratio of the two Higgs vevs) less than 1.7. Then, in improved one loop perturbation theory, avoiding baryon number washout requires the lightest scalar mass to be lighter than \(\sim 50\) GeV, which is not in conflict with current bounds provided that higher order corrections do not decrease the mass bound.

The MSSM has two possible sources of CP violation which are absent in the minimal standard model [24]; a combination of these will be constrained by the BAU. The interactions in this model are given by

\[
\begin{align*}
[U\lambda_U QH + D\lambda_D QH' + E\lambda_E LH' + |\mu| \ e^{-i\phi_B} HH']_F \\
+ m_{3/2} \ [|A| \ e^{i\phi_A} (U\xi_U QH + D\xi_D QH' + E\xi_E LH') + |\mu_B| \ HH']_A.
\end{align*}
\]

Note that unlike the generic two Higgs doublet model discussed in refs. [17,19], the Higgs potential does not contain any CP violating phases at tree level—the phases \(\phi_A\) and \(\phi_B\) occur only in interactions involving the super-partners of the ordinary particles. We find that the contribution of the phase \(\phi_A\) to the BAU is small, and so we focus on the effects
of the phase $\phi_B$, which appears in the mass matrices of the supersymmetric partners of the gauge bosons and the Higgs scalars (the “inos”).

Our strategy for calculating the BAU in the MSSM is as follows. We will fix the top mass at 170 GeV, the lightest scalar mass at 48 GeV, the soft supersymmetry breaking top squark masses at 150 GeV, the A-terms at zero, and $\beta$ at 0.85 ($\tan\beta = 1.14$), since these parameters are already constrained to be near these values [1, and in any case allowing them to vary will only affect the BAU by $O(1)$. Using these values we then calculate the critical temperature, the effective potential at the critical temperature, and the shape of the bubble walls. (Note that these phase transition quantities are not very sensitive to other supersymmetric parameters such as ino masses.) We use the improved one loop approximation for the effective potential, including the order $T^2$ corrections to the gauge propagators, and neglecting the contribution of Higgs doublet loops. When the mass of the lightest scalar is far below the gauge boson masses this approximation is reliable at the critical temperature around the symmetry breaking minimum and in the vicinity of the symmetric minimum for scalar field expectation values larger than $O(g_{w_k}T)$. Although perturbation theory is not valid for a calculation of the effective potential between the two minima, which will affect the width and shape of the phase boundary, fortunately our calculation of the BAU will turn out to be insensitive to the detailed shape of this boundary, provided it is much thicker than $O(1)/T$. The one loop estimate gives the width of the phase boundary to be $11/T$ at the temperature where the two minima are degenerate, so we will assume the phase boundary is thick. The one loop effective potential indicates that the ratio of the expectation values of the two Higgs remains constant during the transition, so we take

$$H' = H \left( \frac{\langle H' \rangle}{\langle H \rangle} \right),$$

where $\langle H' \rangle$ and $\langle H \rangle$ are the expectation values at the critical temperature in the symmetry breaking minimum. The temperature of the transition\(^2\) is 59 GeV, and at this temperature the minimum of the effective potential occurs at $H = 63$ GeV and $H' = 53$ GeV.

---

\(^2\) Our definition of the transition temperature is the spinodal point where the local minimum at the origin vanishes. This temperature is slightly lower than the one loop estimates of the temperature at which the transition actually occurs [12], but these estimates require knowing the effective potential in a region where it is not calculable perturbatively. We use this definition in order to get a conservative upper bound on the BAU produced during the transition.
For thick bubble walls the relevant baryogenesis mechanism is known as “spontaneous 

baryogenesis”, reviewed below. This mechanism, first introduced for baryogenesis during 
a second order phase transition $^{[23]}$, involves a space-time dependent field, which evolves 
coherently during the transition. This time evolution produces a CPT-violating term in 
the effective Hamiltonian called a “charge potential”, which resembles a chemical potential$^{3}$. A charge potential will cause the free energy density inside the bubble walls to be 
minimized for nonzero baryon number, and hence the production of a net baryon number 
via anomalous weak interactions. Spontaneous electroweak baryogenesis during the 
weak phase transition has been suggested for the two Higgs model in refs. $^{[17,18]}$ and as a 
baryogenesis mechanism for supersymmetric models in refs. $^{[20,18]}$.

In the MSSM the $\textit{ino}$ mass matrices are space-time varying during the transition, and 
contain an irremovable CP violating phase, which leads to a charge potential for baryon 
number in the effective fermion Hamiltonian. If the $\textit{inos}$ are not too heavy and the phase is 
not too small, this charge potential is large enough to result in generation of an acceptable 
baryon number during the weak transition. We will find that the resultant BAU depends 
sensitively on the CP violating phase $\phi_B$, and on the masses of the $\textit{inos}$. We are able 
to use the BAU to place upper bounds on $\textit{ino}$ masses and lower bounds on $\phi_B$ and the 
electric dipole moment of the neutron.

Before launching into the specifics of the MSSM calculation, we show how to calculate 
the charge potentials resulting from space-time varying fermion mass matrices.

Consider a fermion mass term of the form 

$$\psi_i^T C m_{ij}(x_\mu) \psi_j + \text{h.c.},$$

(3)

where we take all fermions to be left-handed, and $C$ is the charge conjugation matrix. We 
can make a space-time dependent unitary change of basis on the fermions:

$$\psi_i \rightarrow U_{ij}(x_\mu) \psi_j.$$ 

(4)

$^{3}$ Recall that a chemical potential is a Lagrange multiplier introduced to implement a con-
straint, which appears in the effective Hamiltonian as an energy splitting between particles and 
anti-particles. Similarly, a charge potential results in different energy levels for particles and anti-
particles, but the energy difference is a real physical effect resulting from dynamical violation of 
CPT during a phase transition.
in order to make the fermion mass terms everywhere real, positive and diagonal; however the space-time dependence of $U$ requires that we replace the kinetic energy terms in the Lagrangian by

$$\mathcal{L}_{\text{K.E.}} \rightarrow \mathcal{L}_{\text{K.E.}} + \bar{\psi} \gamma^\mu (U^\dagger i \partial_\mu U) \psi .$$

(5)

Note that since $U$ is a unitary matrix, $U^\dagger \partial_\mu U$ may be written

$$U^\dagger \partial_\mu U = i \partial_\mu \sum_a \alpha_a(x_\mu) T_a ,$$

(6)

where for $n$ fermions the $T_a$s are generators of $U(n)$, and the functions $\alpha(x_\mu)$ are defined by eq. (3). Thus the Lagrangian with mass term (3) is equivalent to a Lagrangian with a real diagonal mass term but also containing a term

$$- \sum_a \partial_\mu \alpha_a(x_\mu) \bar{\psi} T_a \gamma^\mu \psi .$$

(7)

If the transformation (4) has a gauge anomaly there will also be a modification of the Lagrangian

$$\sum_\beta \theta F_\beta \tilde{F}_\beta \rightarrow \sum_\beta \left( \theta + \frac{g_\beta^2}{16\pi^2} \alpha_a \text{Tr} T_a t_\beta^2 \right) F_\beta \tilde{F}_\beta ,$$

(8)

where the $t_\beta$s are gauge generators in the left-handed fermion representation and the $F_\beta$s are the gauge field strengths.

The presence of these anomalous terms complicates the discussion of the charge potentials, and consequently we will choose our unitary transformation to have no gauge anomaly; this is the strategy followed in [25].

Finally, if the transformation (4) does not correspond to a symmetry of the interactions the coupling constants will be affected by the change of basis; however the effect of these coupling constants on energy levels is higher order in perturbation theory and will not concern us here.

What effect does a charge potential have on a thermal system? There are two possibilities:

1) If there is a charge potential for an exactly conserved charge, (e.g. electric charge or B-L) we can ignore it. Although it looks like the system could lower its free energy by producing a net charge, charge conservation imposes a zero charge constraint. It can easily be seen by integration by parts that a charge potential for an exactly conserved
charge has no physical effect; equivalently, we can always redefine fields so such charge potentials never arise.

2) Charge potentials for non-conserved charges will lead to an asymmetry in the rates between processes which create and destroy the charges, until the system reaches thermal equilibrium. The thermal occupation numbers in general will be different for particles and their CP conjugates, due to CP violation and the dynamical CPT violation from the space-time varying scalar fields. For a charge which is approximately (but not exactly) conserved the system will take a long time to reach equilibrium. If the system is near thermal equilibrium the net rate of charge production can be computed using thermodynamic arguments. For instance if there is a small charge potential $\dot{\alpha}_B$ for the baryon number current and no other charge potentials, for a system starting with no net quantum numbers the constraint of zero net baryon number can be implemented by introducing a baryon chemical potential

$$\mu_B = -\dot{\alpha}_B .$$

The chemical potential $\mu_B$ is just the force of constraint on the system, i.e. the derivative of the free energy with respect to baryon number. Now the difference between the rates of anomalous processes which create and destroy baryon number is just proportional to the difference in the change in the free energy per event, leading to an anomalous baryon creation rate of

$$\dot{\rho}_B = -\frac{9\mu_B \Gamma_B}{T} = \frac{9\dot{\alpha}_B \Gamma_B}{T} ,$$

where $\Gamma_B$ is the rate of anomalous baryon violating events per unit volume. The factor of 9 comes about because each anomalous event changes the free energy by $3\mu_B$ and changes the baryon number by 3 units.

We can now calculate the BAU produced in the MSSM during the weak phase transition. We first compute the charged and neutral ino mass matrices, as functions of the scalar field expectation values, assuming equ. (2), and find the transformation on the ino fields $U_1(H)$ which renders their masses real, positive, and diagonal. $U_1$ will in general be anomalous and would, by itself, give rise to an effective $\tilde{W}\tilde{W}$ operator, as well as giving rise

---

4 Note that anomalous weak baryon number violating processes can be affected by a potential for any charge generator whose trace over left handed fermion weak doublets is nonzero \[18\].
to charge potentials for the *inos* via equ. (3). By further making a space-time dependent baryon number rotation of the light fermions we can remove this anomalous operator, at the cost of introducing a charge potential for baryon number 5:

\[
\frac{1}{3} \partial_\mu \left( \int_0^{(H(x))} dH \text{Tr} U_1^+(H) \frac{idU_1(H)}{dH} t_{w_k}^2 \right) j_B^\mu .
\]  

(11)

We then need to know how anomalous processes inside the bubble walls are affected by CP violating terms. Since the *inos* can quickly come into equilibrium with the *ino* charge potentials via ordinary non-anomalous interactions the *ino* charge potentials will not affect anomalous processes 6. The main effect on anomalous baryogenesis will come from the term (11). If the system is near thermal equilibrium then eqs. (10) and (11) can be used to find the total baryon number density produced during the transition 7:

\[
\rho_B = \int dt \dot{\rho}_B
\]

\[
= \int dt \frac{9 \Gamma_B(\langle H \rangle)}{T} \frac{d}{dt} \left( \frac{1}{3} \int_0^{(H)} dH \text{Tr} U_1^+(H) \frac{idU_1(H)}{dH} t_{w_k}^2 \right)
\]

\[
= \frac{3}{T} \int_0^{(H)} dH \text{Tr} U_1^+(H) \frac{idU_1(H)}{dH} t_{w_k}^2 \Gamma_B(H) .
\]  

(12)

We still need to know the rate of anomalous baryon density production $\Gamma_B$, which has been estimated in the symmetric phase to be 26

\[
\Gamma_B \sim \alpha_{w_k}^4 T^4 ,
\]  

(13)

while in the broken phase it is vanishingly small. Unfortunately, there is currently no reliable way of computing $\Gamma_B$ inside the wall where the scalar expectation values are changing. Furthermore it has been claimed, based on some 1 + 1 dimensional simulations

---

5 This could also be seen simply by using the anomaly equation to replace $F \tilde{F}$ with the divergence of the baryon current, and integrating by parts 20.

6 In the limit of vanishing gaugino mass or $\mu$ there is an additional approximate symmetry and the anomalous baryon production rate will be further suppressed. We always assume we are far from this limit, i.e. that the *ino* masses are not much smaller than the temperature.

7 The effect of the spatial component of the charge potential will be to also produce a baryon number current inside the walls. We find this current has no effect on the baryon density produced in the thermal frame in the limit where the walls are thick.
that anomalous baryon production inside the wall is an inherently non-equilibrium process. (Dine suggests this is unlikely to be the case for a transition with thick walls and small latent heat such as occurs in the MSSM [28].) We will use a conservative estimate for the maximum baryon number produced by simply computing the integral (12) using eq. (13) for $\Gamma_B$. We think it is more likely that the baryon production inside the wall will be suppressed when the Higgs vevs become large. For instance McLerran [29] estimates that baryon production is shut off when the Higgs vev reaches about $(1/2)$ its value inside the wall, while Dine et al. [11,20] estimate that this shutdown occurs for vevs of order $g_w k T/(4\pi)$. In Figure 1 we plot the baryon number produced as a function of the value of the Higgs vev where this shutdown occurs for a typical choice of parameters; note that the dependence is approximately quadratic. We conclude that it is most likely that the actual baryon number produced during the transition will be between $\sim 1/4$ and $\sim 10^{-2}$ times smaller than our most favorable estimate.

In order to find the allowed range of parameters in the MSSM, for each value of the gaugino and higgsino masses we adjust the CP violating phase $\phi_B$ to be as large as is compatible with current constraints on the EDMN [30]. We then compute the BAU, from equs. (12) and (13). The allowed range of ino masses are those for which the upper bound on the baryon to entropy ratio is greater than $0.4 \times 10^{-10}$. For masses in the allowed range, we compute the minimum value of $\phi_B$ which is consistent with the BAU, and use the calculation of Kizukuri and Oshimo in ref. [30] to find the lower bound on the EDMN.

In Figures 2a and 2b we plot the lower bound on the EDMN, as a function of $\mu$ and the wino mass parameter $m_2$. The first figure has a phase $\phi_B$ near 0, while the second figure has the phase near $\pi$. The bino mass parameter $m_1$ is assumed to satisfy the GUT relation $m_1 = m_2 (5/3) \tan^2 \theta_W$. The signs in fig. 2b indicate the sign of the EDMN, while in fig. 2a the EDMN is negative. The black excluded region corresponds to an EDMN greater than the current bound of $10^{-25}$ e-cm [31], the dark grey region is an EDMN greater than $10^{-26}$ e-cm, and the light grey region greater than $10^{-27}$ e-cm. Thus for any portion of the allowed parameter space the EDMN is greater than $10^{-27}$ e-cm. (Similarly, if we assume the selectron mass is equal to the squark mass, the electric dipole moment of the electron is greater than $O(3) \times 10^{-29}$ e-cm.) Throughout all of the region where the EDMN could be less than $10^{-26}$ e-cm the lightest chargino mass is lighter than 88 GeV.

\[8\] In computing the EDMN, we use the quark model and neglect the contribution of the phase $\phi_A$; we have assumed the A-terms are small.
and the lightest neutralino mass is lighter than 44 GeV. Also, if the EDMN is less than $10^{-26}$ e-cm the phase $\phi_B$ must be near $\pi$ in order to avoid a chargino mass lighter than $M_Z/2$.

These figures have assumed that the rate of baryon violation is given by equ. [26], and that this violation is occurring throughout the bubble wall. If this is not the case, the baryon number produced will be reduced and the resulting phase must then be larger to compensate, giving rise to larger dipole moments.

If the experimental bound on the EDMN is pushed down by an order of magnitude, or if numerical calculation finds the rate of anomalous baryon creation inside the phase boundary to be much less than it is in the symmetric phase, then there are strict upper bounds on $ino$ masses. If the MSSM is responsible for baryogenesis the prospects for discovering electric dipole moments and supersymmetric particles in the next few years are excellent. Should we fail to make these discoveries, there are several possible conclusions. The BAU may come from physics above the weak scale, from weak scale physics other than supersymmetry, or from a more complicated supersymmetric model. The simplest extension of the MSSM would be to add a gauge singlet superfield, which changes the allowed form of the Higgs potential and greatly relaxes all the constraints on scalar masses [32]. The tree level potential will generically contain large cubic terms, and gives a very strongly first order transition with a thin boundary between the two phases. Furthermore, with gauge singlets there is the possibility for CP violating phases in the Higgs potential. With thin bubble walls we expect baryogenesis to be dominated by the mechanism of refs. [16,19], in which CP violating particle scattering processes from the phase boundary leads to a transport of particle quantum numbers, biasing anomalous baryon production throughout the symmetric phase. The charge transport mechanism could produce the BAU for phases as small as $10^{-4}$ [19]. Unfortunately such extended supersymmetric models currently have too many free parameters to allow for calculation of the BAU.

In summary, we have shown that unlike the minimal standard model, the MSSM is still viable, but soon either it should be ruled out (for baryogenesis), or new experimental discoveries such as electric dipole moments will give us an important clue towards the explanation of why there are more baryons than anti-baryons. We are still a long way from testable predictions from baryogenesis in extensions of the MSSM.
Acknowledgments

We would like to thank David Kaplan and Stanley Myint for useful discussions. A. C. was supported in part by DOE contracts #DE-AC02-89ER40509, #DE-FG02-91ER40676, by the Texas National Research Laboratory Commission grant #RGFY91B6, and NSF contract #PHY-9057173; A. N. was supported in part by DOE contract #DE-FG03-90ER40546, by a fellowship from the Alfred P. Sloan Foundation, and by an SSC Fellowship from the Texas National Research Laboratory Commission.
References

[1] E. Kolb and M. Turner, The Early Universe (Addison-Wesley, New York, (1990)
[2] A.D. Sakharov, JETP Lett. 6 (1967) 24
[3] M.E. Shaposhnikov, Nucl. Phys. B287 (1987) 757; Nucl. Phys. B299 (1988) 797; Phys. Lett. 277B (1992) 324, Erratum, Phys. Lett. 282B (1992) 483
[4] J. Ambjorn, K. Farakos, and M.E. Shaposhnikov, Niels Bohr Institute preprint NBI-92-20 (1992)
[5] Shaposhnikov’s original estimate of the BAU left out several important suppression factors, such as cancellations in CP violating phases, the slow rate of anomalous baryon number violation, and the large width of the bubble walls. See the Erratum in [3].
[6] G. t’Hooft, Phys. Rev. Lett. 37 (1976) 8; Phys. Rev. 14 (1976) 3432
[7] A. Linde, Phys. Lett. 70B (1977) 306; S. Dimopoulos and L. Susskind, Phys. Rev. D18 (1978) 4500; N. Christ, Phys. Rev. D21 (1980) 1591; N.S. Manton, Phys. Rev. D28 (1983) 2019; F.R. Klinkhammer and N.S. Manton, Phys. Rev. D30 (1984) 2212
[8] V.A. Kuzmin, V.A. Rubakov and M.E. Shaposhnikov, Phys. Lett. 155B (1985) 36
[9] P. Arnold and L. McLerran, Phys. Rev. D36 (1987) 581; Phys. Rev. D37 (1988) 1020
[10] M.E. Shaposhnikov, JETP Lett. 44 (1986) 465; A.I. Bochkarev, S. Yu. Khlebnikov and M.E. Shaposhnikov, Nucl. Phys. B329 (1990) 490.
[11] M. Dine, P. Huet, and R. Singleton, Nucl. Phys. B375 (1992) 625
[12] M. Dine, R. G. Leigh, P. Huet, A. Linde, and D. Linde, SLAC preprint SLAC-PUB-5741 (1992); B.H. Liu, L. McLerran, and N. Turok, Minnesota preprint TPI-MINN-92/18-T (1992)
[13] G.F. Giudice, Phys. Rev. D45 (1992) 3177
[14] S. Myint, Phys. Lett. 287B (1992) 325, and also work in progress
[15] L. McLerran, Phys. Rev. Lett. 62 (1989) 1075
[16] A.G. Cohen, D.B. Kaplan and A.E. Nelson, Phys. Lett. 245B (1990) 561; Nucl. Phys. B349 (1991) 727
[17] N. Turok and J. Zadrozny, Phys. Rev. Lett. 65 (1990) 2331; Nucl. Phys. B358 (1991) 471; L. McLerran, M. Shaposhnikov, N. Turok and M. Voloshin, Phys. Lett. 256B (1991) 451
[18] A.G. Cohen, D.B. Kaplan and A.E. Nelson, Phys. Lett. 263B (1991) 86
[19] A.E. Nelson, D.B. Kaplan and A.G. Cohen, Nucl. Phys. B373 (1992) 453
[20] M. Dine, P. Huet, R. Singleton and L. Susskind, Phys. Lett. 257B (1991) 351
[21] R.N. Mohapatra and X. Zhang, Maryland preprint UMDHEP 92-230 (1992)
[22] S. Dimopoulos and H. Georgi, Nucl. Phys. B193 (1981) 150; N. Sakai, Z. Phys. C11 (1981) 153
[23] G. Anderson and L. Hall, Phys. Rev. D45 (1992) 2685
[24] J. Ellis, S. Ferrara, and D.V. Nanopoulos, Phys. Lett. 114B (1982) 231; W. Buchmuller and D. Wyler, Phys. Lett. 121B (1983) 321; J. Polchinski and M.B. Wise, Phys. Lett. 125B (1983) 393; F. del Aguila, M. Gavela, J. Grifols, and A. Mendez, Phys. Lett. 126B (1983) 71; D.V. Nanopoulos and M. Srednicki, Phys. Lett. 128B (1983) 61; M. Dugan, B. Grinstein and L.J. Hall, Nucl. Phys. B255 (1985) 413
[25] A.G. Cohen and D.B. Kaplan, Phys. Lett. 199B (1988) 251; Nucl. Phys. B308 (1988) 913
[26] J. Ambjorn, T. Askgaard, H. Porter and M. Shaposhnikov, Phys. Lett. 244B (1990) 479
[27] D. Grigoriev, M. Shaposhnikov and N. Turok, Princeton preprint PUPT-91-1275 (1991)
[28] M. Dine, talk given at the Yale-Texas symposium on electroweak baryon violation, (March 1992) SCIPP 92/21
[29] L. McLerran, talk given at the ITP workshop on Cosmological Phase transitions, Santa Barbara, April, 1992
[30] R. Arnowitt, J.L. Lopez, and D.V. Nanopoulos, Phys. Rev. D42 (1990) 2423; Y. Kizukuri and N. Oshimo, Phys. Rev. D45 (1992) 1806
[31] K. Smith et al., Phys. Lett. 234B (1990) 191
[32] M. Pietroni, Padova preprint PFPD/92/TH/36 (1992)
**Figure Captions**

Fig. 1. Baryon number density as a function of the Higgs vev.

Fig. 2a. Lower bound on the EDMN for phase $\phi_B$ near 0. The black region is excluded, the grey region has an EDMN greater than $10^{-26}$ e-cm, and the white region has an EDMN greater than $10^{-27}$ e-cm.

Fig. 2b. Same as Figure 2a, except the phase $\phi_B$ is near $\pi$. The $\pm$ signs indicate the sign of the contribution of $\phi_B$ to the down quark electric dipole moment at one loop, when $\phi_B$ is chosen to produce a positive BAU.