THE ELECTROWEAK THEORY AND BARYOGENESIS*

MICHAEL DINE
Santa Cruz Institute for Particle Physics
University of California, Santa Cruz, CA 95064

ABSTRACT

The electroweak interactions violate baryon number, as a consequence of anomalies. At high temperatures, this violation is enhanced. This raises the prospect of producing the observed matter-antimatter asymmetry at the electroweak phase transition. In this lecture, we briefly review the various aspects of this problem. After explaining the nature of the high-temperature enhancement, we discuss the phase transition, sources of CP violation in extensions of the Standard Model, and various mechanisms which have been considered for producing the asymmetry. Limitations of our present understanding and computational abilities are stressed.

The past few years have witnessed the triumph of the standard electroweak theory. At this conference, we have heard much about precision tests of the theory. Over the past few years, aside from these stunning experimental results (and their theoretical evaluation and interpretation), there has been an important realization on the theoretical side: the electroweak interactions badly violate baryon number at high temperatures,[1-3] and thus have the potential to explain the density of matter in the universe.[1,4-11] Today I will review the basic ideas which underlie this possibility. We will see that, if weak interactions are the origin of the asymmetry, something more than the minimal Higgs sector is required. Thus there is the real potential for interplay between particle physics experiments and our understanding of the behavior of the universe at extremely early times.

It will be helpful to review a few highlights of the conventional big bang theory. Standard big bang cosmology assumes that the universe is homogeneous and isotropic on large distance scales. From observations of the cosmic microwave radiation background, we know that this is true to an extraordinary degree.[†] Solving Einstein's equations with these assumptions gives a universe whose scale,

⋆ Work supported in part by the U.S. Department of Energy.

† The recent COBE observations measure extremely small deviations from isotropy. They also show that the radiation is extremely isotropic. See the talk by Smoot at this meeting.

R, grows with time as $1/\sqrt{t}$ during the radiation dominated era, or $t^{-2/3}$ in a matter dominated era. The temperature of the early cosmic soup behaves as $T \propto 1/R$. We can think of the temperature as a sort of cosmic clock. Some high points in the history of the universe include:

a. $t \sim 100$ sec. $T \sim$ MeV synthesis of light elements,

b. $t \sim 10^5$ yrs. $T \sim$ eV electrons, nuclei combine to form atoms,

c. $t \sim 10^9$ yrs. $T \sim 10^{-2}$ eV stars, galaxies form; light elements are "cooked" to form heavier elements, and

d. $t \sim 2 \times 10^{10}$ yrs. $T = 2.7$ K the present epoch.

Support for this picture comes from observation of the Hubble expansion, observation of the cosmic microwave radiation background (a relic of the time of recombination), and the successful prediction of the abundance of light elements.

About earlier times we have no direct evidence. Interesting moments include:[12]

e. $T \sim 100$ MeV QCD phase transition (system changes to one described by free quarks and gluons).

f. $T \sim 100$ GeV Electroweak phase transition. The W and Z masses decrease with increasing temperature; above the transition temperature, the W and Z behave as if they were (virtually) massless.

g. $T \sim 10^{16}$ GeV (?) Inflation(?):[12,13] gives rise to the observed homogeneity and isotropy of the universe; sweeps away magnetic monopoles and other topological defects; gives rise to primordial fluctuations [responsible for galaxy formation, recent COBE observations(?)].

Of course, for each of these, the degree of speculation increases with the temperature.

The question we would like to examine here is: why does the universe contain more matter than antimatter? Today, the ratio of baryons to photons is

$$\frac{n_B}{n_\gamma} = 10^{-9} - 10^{-10} . \qquad (1)$$

One could imagine that this is simply an initial condition, but this is rather perverse. For example, at $T \sim 1$ GeV, there are of order 10^{39} quarks and antiquarks per cc. At this time, the tiny asymmetry looks almost absurd, and is completely irrelevant to the dynamics of the quark-gluon plasma.

Sakharov was the first to recognize that one could scientifically address the question of the matter-antimatter asymmetry.[14] He enumerated three critical ingredients to any understanding:

1. Baryon Number Violation in the fundamental, microscopic laws. This is necessary if one is to understand how, starting with a $B = 0$ universe one obtains a $B \neq 0$ universe.
2. C and $CP(T)$ violation in these laws. Otherwise, for every interaction which produces baryons, there will be one proceeding at exactly the same rate which produces antibaryons.
3. An arrow of time (departure from thermal equilibrium). If the system is always at equilibrium, the baryon number will vanish, even if it is not conserved – the equilibrium state of a system is determined solely by conserved quantities. Alternatively, we can understand this by noting that CPT guarantees that particles and antiparticles possess equal mass, so their equilibrium distributions are identical.

The first compelling framework in which one could hope to understand the baryon asymmetry was provided by Grand Unified Theories (GUT's).[15] Because quarks and leptons are unified in multiplets of a single gauge group, baryon number is violated. In GUT's, the couplings of superheavy gauge bosons and color-carrying Higgs bosons violate B and L. These fields typically have masses of order the unification scale, M_{GUT}. In GUT scenarios for baryogenesis, it is assumed that for temperatures larger than M_{GUT}, these particles (and B-violating interactions) are in thermal equilibrium; as the temperature drops below this, they drop out of equilibrium. GUT's generally contain many sources of CP-violation; thus the rates for decays, for example, into different channels of opposite baryon number need not be the same. By carefully studying rate equations, one can compute the resulting asymmetry. Typically, assuming weak coupling at the unification scale, one obtains small asymmetries which can easily be consistent with what is observed.

GUT baryogenesis might well be the complete story. If this is true, it is rather disappointing, since we have little hope of studying experimentally, even indirectly, the details of any GUT. There are, however, some troubling things about this picture. First, inflation[12,13] if it occurs, dilutes any pre-existing baryon number. In order to produce baryons, it is necessary after inflation, that the universe reheat to M_{GUT}. In typical models of inflation, this is hard to achieve and requires fine-tuning of parameters. Second, many particle physics models lead to entropy generation at temperatures below M_{GUT}. Since $n_\gamma \propto s$ (the specific entropy), this means that the baryon number, typically rather small to begin with, is diluted below the observed level. Variants of these problems occur, for example, in supersymmetric models, where inflation is required to get rid of gravitinos, but any significant reheating gives back an unacceptable gravitino density.[16] For all of these reasons and more, mechanisms which produce baryons at lower temperature are quite attractive.

While there have been a number of proposals for low temperature baryon number generation,[17] perhaps the most appealing and exciting possibility is that the baryon asymmetry might be generated at the temperatures of order the

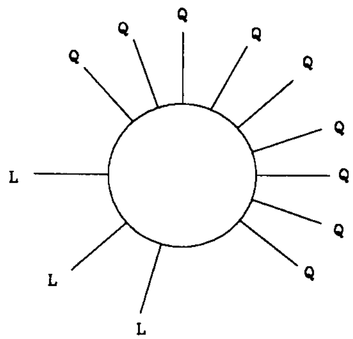

Fig. 1 Effective baryon and lepton number-violating vertex in the Standard Model.

electroweak scale, by electroweak interactions. At first sight, this appears absurd. After all, the fundamental vertices of the Standard Model preserve baryon number, *i.e.*, baryon number is a good symmetry to all orders of perturbation theory. However, it has been known since the work of 't Hooft[18] that at low energies baryon and lepton number are violated by a tiny amount in the electroweak theory. This is associated with the fact that the baryon and lepton currents in the electroweak theory are anomalous. As a result, the potential of the gauge fields has local minima which differ in baryon number. These minima are separated by a barrier; the barrier height is referred to as the "sphaleron energy," E_{sp}.[19] In general,

$$E_{sp} = AM_W/\alpha_W \tag{2}$$

where A is of order 1–2; the barrier height is thus of order 10's of TeV. At low energies (compared to M_W), the system can tunnel between these minima. The tunneling rates are computed by means of "instantons," solutions of the classical Euclidean equations of motion. The results can be described by effective B-violating interactions; these can be represented as in Fig. 1. The interaction strength is proportional to $e^{-2\pi/\alpha_W}$. This is smaller than any power of α_W, *i.e.*, it is consistent with the statement that baryon number is conserved in perturbation theory; numerically, it is so small that even before worrying about prefactors, one can say that no proton has decayed as a result in the age of the universe.

Because baryon number violation is a tunneling process, one can hope that it is enhanced at high temperatures[20] or energies.* One can give rather convincing arguments that the rate is enhanced at high temperatures.[1,3] Below the weak phase transition, the rate behaves as

$$\Gamma \sim e^{-AM_W(T)/\alpha_W T}\ , \tag{3}$$

where A is typically 1–2, and $M_W(T)$ is the temperature-dependent W mass, which, as we have mentioned earlier, vanishes at some critical temperature. Above the critical temperature, the rate is substantial; it behaves as

$$\Gamma = \kappa(\alpha_W T)^4\ . \tag{4}$$

Thus the electroweak theory gives precisely what we would like for baryogenesis: at low temperatures baryon number violation is negligible, while at high temperatures it is rapid.

Before asking whether one can produce the asymmetry through electroweak interactions, it is interesting to consider the consequences for GUT baryogenesis. Since B is violated, and B-violating interactions are in equilibrium down to rather low temperature, there is the danger that any asymmetry produced early on may be lost. Fortunately, the minimal standard model, while it violates B and L, conserves $B - L$. So, if the grand unified theory violates $B - L$ and produces some net value of this quantity, electroweak effects will leave a net B and L. However, if lepton number is violated by neutrino masses, all B and L may disappear.[23] This can be used to give an upper limit on the lightest Majorana neutrino mass:

$$M_\nu < \frac{4\ \text{eV}}{(T_{B-L}/10^{10}\ \text{GeV})^{1/2}}\ . \tag{5}$$

Returning to the question of electroweak baryogenesis, the rapid violation of B at high temperatures means that the first of Sakharov's conditions is satisfied. As for the second, CP violation exists already in the minimal Standard Model; most extensions of the model produce more. (We will see shortly that there is too little CP violation in the minimal model to produce the asymmetry). What about departure from equilibrium? This will arise if the electroweak phase transition is first order.[1] Even in the minimal theory, the phase transition turns out to be first order, for light enough Higgs particles. In fact, over the last year, the problem of electroweak baryogenesis has stimulated a great deal of work on the electroweak phase transition. We turn to this now.

In this lecture we will confine our attention to the minimal theory. Most of the analysis we will describe here is readily extended to models with more scalars. (Moreover, in many cases, the minimal theory emerges as an effective theory at the phase transition.) To determine the nature of the phase transition, one needs

---

* The question of whether the baryon number violation rate is enhanced at high energies is a controversial one. In ref. 21, it is argued that there is no enhancement. Other viewpoints on this question are reviewed in ref. 22.

to calculate the effective potential for the Higgs particle, *i.e.*, one must find the free energy cost (or gain) associated with a given value of the Higgs expectation value. To lowest order, this means one has to calculate the free energy of a gas of quarks, leptons, and gauge bosons, with mass corresponding to a given value of ϕ, *i.e.*,

$$V_T(\phi) = \pm \sum_i \int \frac{d^3p}{2\pi^3} \ln\left(1 \mp e^{-\beta\sqrt{p^2+M_i^2(\phi)}}\right) \tag{6}$$

where the sum is over all particle species (physical helicity states), the minus sign is for bosons and the plus sign for fermions. For example, for W bosons,

$$M_W^2(\phi) = g_2^2\phi/4 \ . \tag{7}$$

Note here, and in what follows, we are normalizing ϕ as if it were a real field. (For better or worse, this is standard in this subject.)

The integrals are straightforward to evaluate. We will take the Higgs mass smaller than the W mass, and ignore the scalar contribution to simplify the writing. One then finds

$$V(\phi,T) = D(T^2 - T_o^2)\phi^2 - ET\phi^3 + \frac{\lambda_T}{4}\phi^4 \ . \tag{8}$$

Here

$$D = \frac{1}{8v_o^2}(2M_W^2 + M_Z^2 + 2M_t^2) \ , \tag{9}$$

$$E = \frac{1}{4\pi v_o^3}(2M_W^3 + M_Z^3) \sim 10^{-2} \ , \tag{10}$$

$$T_o^2 = \frac{1}{2D}(\mu^2 - 4Bv_o^2) = \frac{1}{4D}(M_H^2 - 8Bv_o^2) \ , \tag{11}$$

$$\lambda_T = \lambda - \frac{3}{16\pi^2 v_o^4}\left(2M_W^4 \ln\frac{M_W^2}{a_B T^2} + M_Z^4 \ln\frac{M_Z^2}{a_B T^2} - 4M_t^4 \ln\frac{M_t^2}{a_F T^2}\right) \ , \tag{12}$$

where $\ln a_B = 2\ln 4\pi - 2\gamma \simeq 3.91$, $\ln a_F = 2\ln\pi - 2\gamma \simeq 1.14$.

If we ignore the small ϕ^3 term, the transition is second order, with the mass of the scalar vanishing at temperature T_o. Above this temperature, the expectation value of ϕ vanishes and the gauge bosons are massless; below it they are massive. Including the ϕ^3 term modifies the picture, rendering the transition first order. The behavior of the system as the temperature is lowered is shown in Fig. 2. For temperatures well above T_o, the potential has a unique minimum at $\phi = 0$. At a temperature T_1, slightly above T_o, a second minimum appears in the potential. T_1 is given by

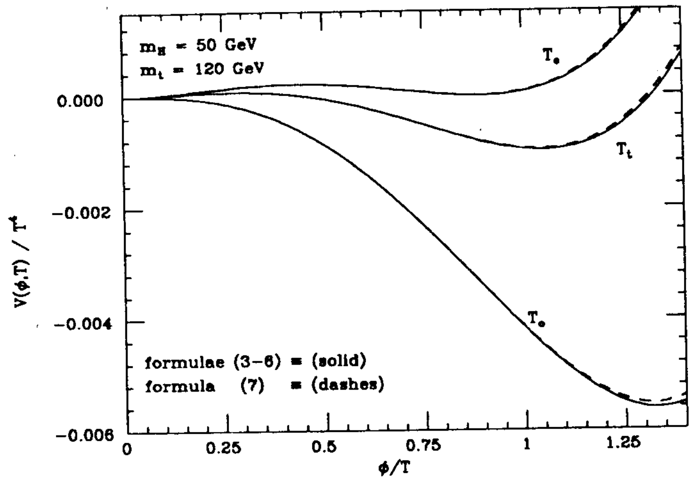

Fig. 2. Behavior of the potential as the temperature is lowered.

$$T_1 = \frac{T_o}{1 - 9E^2/8\lambda_T D} . \tag{13}$$

At a temperature T_c, this new minimum becomes degenerate with the minimum at the origin; T_c and the corresponding minimum of the potential ϕ_c are given by

$$T_c^2 = \frac{T_o^2}{1 - E^2/\lambda D} \qquad \phi_c = \frac{2E\,T_c}{\lambda_{T_c}} \ . \tag{14}$$

Thus in lowest order the transition is first order. Such a transition should proceed through formation of bubbles, in which the Higgs field is non-vanishing; these bubbles should then grow until they collide. Near the surface of these bubbles, the system is far from equilibrium, and it is here than one might expect that baryon number is generated.

If we take the above picture seriously, before asking how the phase transition actually proceeds, or how baryon number is produced, we can already set an upper limit on the value of the Higgs mass for which baryogenesis can possibly occur.[24] The point is that, once the phase transition is over, and ϕ has settled to its minimum everywhere, it is necessary that the baryon violating rate, Γ, be small compared to the expansion rate of the universe. Otherwise, whatever B

was produced during the transition will disappear. The rate, however, depends on the value of the ϕ after the transition, essentially as

$$\Gamma(\phi) = \gamma \alpha_2^{-3} M_W^7 e^{-Ag\phi/2\alpha_W T} . \quad (15)$$

The expectation value of the Higgs field after the transition is found from minimizing the effective potential above. One finds that the Higgs mass must be less that 45 ± 10 GeV in order that the baryon number not be washed out.[25] This is already below present limits. Once certain corrections to the potential discussed below are taken into account, the limit is lowered by about 12 GeV.[26] Of course, the true theory need not be the minimal Standard Model, and the limits can be significantly relaxed in theories with additional, relatively light scalar fields.[27] A particularly simple model, with a light singlet, is described in ref. 28.

To compute the baryon number produced at the transition, we need some understanding of how it actually takes place. We might expect such a first order transition to proceed through the formation of bubbles. Outside the bubble the system is in the "wrong" (*i.e.*, non-equilibrium) state; inside is in the "true" state. In the present case, this means that outside the bubble, the mean value of the scalar field vanishes, while inside it is non-zero, as indicated in Fig. 3. There is a standard procedure for calculating the nucleation rate for such bubbles.[29] In vacuum, the bubbles would expand at nearly the speed of light, filling all space. In the plasma of the early universe, they grow more slowly. It is somewhat more difficult – as we will discuss below – to determine the final velocity of the bubble walls. However, one expects that this velocity will be at least some modest fraction of the speed of light. With this assumption, one finds that bubbles fill the universe at a temperature only slightly below T_c (typically about 98% of T_c).[8,9]

In the various scenarios which have been considered, baryogenesis occurs near the bubble walls, since this is the region in which there is the largest departure from equilibrium. Thus it is important to understand the features of bubble propagation in the medium, particularly the size and terminal velocity of the wall. As we noted above, these turn out to be difficult problems. Naively,[29] one expects that the wall is slowed by the scattering of particles, particularly top quarks and W and Z bosons, which gain mass as they pass through the walls. Recent analyses are somewhat more sophisticated, but many potentially important processes have not yet been included.[26,30] These calculations give, for the minimal Standard Model with Higgs particles of relatively low mass (where the transition is reasonably first order)

$$\gamma v \sim 0.1 - 0.8, \quad (16)$$

$$l \sim 10 - 30 T^{-1} \quad (17)$$

where v and l are the wall thickness and velocity.

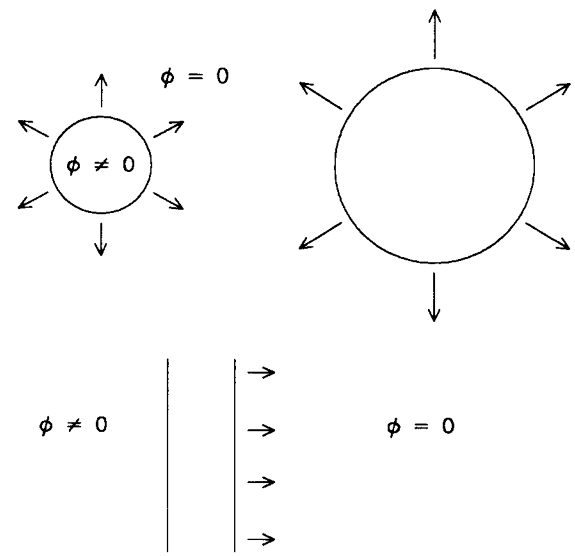

Fig. 3. Formation and growth of bubble during the phase transition.

We have seen, the lowest order analysis leads to a first order phase transition for sufficiently light Higgs in the minimal model; the analysis is easily extended to more complicated models. In all of these calculations, the ϕ^3 term in the potential is crucial. Yet, even in the minimal model, there is a great deal of confusion and controversy. A number of papers have appeared over the last year disputing virtually all aspects of the picture described above.

1. Order of the phase transition? The ϕ^3 term is an infrared effect. It is necessary to check that perturbation theory is under control, and important contributions have not been missed. In refs. 31 and 32, it was argued that in addition to the cubic terms, there are linear terms, which might weaken or strengthen the transition. In ref. 26, it was shown that for small enough λ, there is in fact a systematic expansion in powers of λ/g^2, and that there is no large linear term. However, the coefficient of the cubic term actually is 2/3 of its lowest order value, resulting in the stronger limit on the Higgs mass mentioned above.

2. How does the phase transition proceed? There has been debate over whether the transition, for the range of Higgs masses we have been discussing, is genuinely first order. In refs. 33 and 34 it is argued that the transition is effectively second order; counterarguments have been given in

ref. 26 and 35.

3. How do bubbles propagate? Here, a variety of results have been obtained, including the possibility that the wall is ultrarelativistic.[36] As of this writing, while there is perhaps not complete agreement, a more or less consistent set of analyses have emerged, giving the range of velocities mentioned above.[26,30]

For the rest of this talk, we will assume the validity of the picture which I have outlined up to now. Then all of the conditions required to produce an asymmetry are satisfied. The question which remains is: under what circumstances is an adequate asymmetry produced? Two types of scenarios have been considered:

1. Baryons are produced in the bubble wall due to time-variation of scalar fields.[7,8,9] As we will see shortly, this mechanism may just barely be able to produce the observed asymmetry. The problem is that ϕ itself suppresses B-violating processes (recall that $\Gamma \sim e^{-A\phi/2\alpha_W T}$).[8,25] Thus, assuming that baryon-number violating processes are nearly in equilibrium, baryons are only produced in a rather small layer in the wall.

2. Baryons are produced in front of the wall in a region where $\phi = 0$. Potentially, this is much more efficient.[11]

All of these scenarios require extensions of the Standard Model with more CP violation, such as multi-Higgs models, supersymmetry, or technicolor. The two cases which have been examined most extensively are multi-Higgs theories and supersymmetric theories, and it is perhaps worthwhile to comment on CP-violation in both cases. In a two-Higgs model, the most general potential contains several CP-violating phases. In order to suppress flavor changing neutral currents in this model, however, one often imposes a discrete symmetry. This symmetry, eliminates all of the CP-violating angles, so it is then necessary to suppose that it is at least softly broken if one wants CP violation. Supersymmetric theories contain several new phases. For example, the minimal supersymmetric Standard Model has potential CP-violation in the gaugino and squark mass matrices. Unless the superpartners of ordinary fields are very heavy, however, the neutron and electron electric dipole moments set stringent limits on these phases.

Let us now consider some specific scenarios of the two types listed above. First, we consider the case where the asymmetry is produced in the bubble wall. The simplest such case to describe (though also, as we will see, the least effective) arises when one has some new fields at a scale $M \gg T_c$ (here the new fields might be additional Higgs fields, fields connected with supersymmetry, or something else). At scales below M, suppose as well that there is a single Higgs doublet, ϕ. Integrating out physics at or above M, the leading CP-violating bosonic operator is

$$\mathcal{L}_{CP} = \frac{g^4 \sin(\delta)}{64\pi^2 M^2}|\phi|^2 F\tilde{F}\,, \tag{18}$$

where $\sin(\delta)$ is a measure of CP violation, and M represents the scale of new physics. Using the anomaly equation, this may be rewritten

$$\mathcal{L}_{CP} = \frac{g^2 \sin(\delta)}{12M^2}|\phi|^2 \partial_\mu j^{\mu\ B}. \tag{19}$$

Now during the phase transition, ϕ is changing in time near the bubble walls. Suppose ϕ changes slowly in time, and can be viewed as nearly homogeneous in space. Then we may rewrite the interaction term as

$$\mathcal{L}_{CP} = \frac{g^2 \sin(\delta)}{12M^2}\partial_0|\phi|^2 n_B.$$

If ϕ is sufficiently slowly varying, this is a chemical potential for baryon number, n_B, and the minimum of the free energy is shifted away from zero. A simple calculation gives for the location of the new minimum

$$n_B^o = \mathcal{L}_{CP} = \frac{g^2 \sin(\delta)}{48M^2}\partial_0|\phi|^2. \tag{20}$$

Since baryon number is violated, the system tries to get to this minimum. A simple detailed balance argument gives for the rate of approach to this quasi-equilibrium state

$$\frac{dn_B}{dt} = -18\Gamma(t)T^{-3}[n_B - n_B^o(t)]. \tag{21}$$

Note that the result is sensible: the rate vanishes at the minimum; it is proportional to Γ and the number of quark doublets.

We assume that the change in the field is adiabatic, in the sense that everything is in equilibrium except n_B. In this limit, $\Gamma(t)$ is determined by the instantaneous value of ϕ, *i.e.*,

$$\Gamma(t) = \Gamma[\phi(t)] \sim e^{-AM_W(\phi)/(\alpha_W T)}. \tag{22}$$

This turns off for $M_W \sim \alpha_W T$. Because of the two factors of ϕ in $\mathcal{L}_{CP}$, this is potentially a large suppression. To make a rough estimate of the asymmetry, take

$$\begin{aligned}\Gamma &= \kappa(\alpha_W T)^4 \quad M_W < \sigma\alpha_W T \\ &= 0 \quad M_W > \sigma\alpha_W T\end{aligned} \tag{23}$$

where σ is an unknown numerical constant, about which we will comment below. It is now a trivial matter to solve the rate equation, since both sides involve total derivatives

$$\frac{n_B}{n_\gamma} = \alpha_W^6 \frac{3\sigma^2 T^2 \sin(\delta)\kappa}{g^* M^2}. \tag{24}$$

Here g^* is the number of light degrees-of-freedom at the phase transition (of order 100). Of the six powers of α_W, four come from the rate, while two come from the factors of ϕ in the operators. α_W^6 is already of order 10^{-9}, so to have any chance of obtaining the observed asymmetry in such a model, one needs $T \sim M$, maximal CP violation, and good luck with κ and σ.

A model with two Higgs doublets has been extensively studied as a possible framework for baryogenesis. [7,9-11] The analysis of this model is more complicated, since the t quark is massless at the phase transition. Its dynamics must thus be carefully considered. There is a potential problem in these models associated with the suppression described earlier. If one only considers quadratic terms in the potential, the model conserves CP. On the other hand, in order that baryon production be appreciable, it is necessary that ϕ be quite small. This leads to even further suppression over that described in the scenario we described earlier. Other questions have been raised about the behavior of these models, as well.[10]

The analyses which lead to these conclusions are all based on an adiabatic picture, in which the baryon-violation rate depends on the instantaneous value of ϕ. In their seminal paper on this subject, Turok and Zadrozny[7] assumed that there was no suppression of the type we have been describing. They argued that the changing Higgs field itself could force the system over the barrier; CP-violation could bias this process. Thus there is no obvious suppression as ϕ grows. Recently, Grigoriev, Shaposhnikov and Turok[37] have provided what they argue is numerical support for these ideas from numerical simulations of $1+1$ dimensional models. One particularly puzzling feature of their results is that the asymmetries which one finally produces are largest at low temperatures.

Even if these numerical results are correct, it is hard to see how this can be relevant to phase transitions of the type considered here. The problem is that there is a characteristic energy scale associated with baryon number violating transitions, of order the "sphaleron energy," $E_{sp} \sim M_W/\alpha_W$. There is also a corresponding length scale, M_W^{-1}. A simple estimate, however, shows that the energy deposited by the moving wall in a sphaleron volume is roughly a factor 30 to 100 times smaller than the sphaleron energy. This suggests that one must consider biasing of transitions "in process." In the wall, this is likely to be ineffective.[8]

Let us turn, finally, to scenarios in which baryon number is produced in front of the wall, as suggested by Cohen, Kaplan and Nelson (CKN).[6,11] These authors have considered this possibility in a number of models. In the two Higgs model which we have introduced above, it is easy to describe these analyses in words. CKN assume some form for the Higgs field as a function of x. They then solve the Dirac equation for top quarks striking the wall, and obtain the reflection and transmission coefficients for left and right handed top quarks (and antiquarks) striking the wall. Because CP is violated, there is an asymmetry in the scattering of top quarks of different handedness; this gives rise to an asymmetry of what we might call "axial top quark number" in front of the wall. CKN then simulate the motion of the top quarks through the plasma, allowing, for example, for elastic scattering with other quarks and with gluons. They obtain, in this way, a density profile as a function of distance from the wall. The results depend on the assumed form for the wall shape and on the assumed value of the wall velocity. The asymmetries one finds, however, can be substantial and extend out significant distances in front of the wall (*e.g.*, $100T^{-1}$). For baryon production, the important point is that this asymmetry *biases* the baryon number-violating process in the region where it is unsuppressed. To see this, one minimizes the free energy with non-zero density of axial top quark number, and zero density for other (approximately) conserved quantum numbers. One then obtains a rate equation quite similar to that described earlier,

$$\frac{dn_B}{dt} = c\Gamma T^3 (n_B - n_B^o) . \tag{25}$$

For maximal CP violation, CKN find that the baryon to photon ratio which results can be as large as 10^{-5}. This holds if the wall has thickness of order T^{-1} and velocity of order one. The asymmetry decreases rapidly as the wall thickness grows, and as the velocity decreases.

From all of this I believe that the baryon asymmetry may well have been created at the electroweak phase transition. Ideally, we would like to be in a position where, given a model, we could turn the crank and determine the resulting baryon to photon ratio. Before we reach that stage, we will need:

1. Improved calculations (simulations) of B-violating rates, both with and without a background ϕ.
2. Better understanding of the phase transition, including exploration of models with more strongly first order transitions, as well as simulations in models where the weak coupling expansion is poor.
3. Better understanding of the various processes which can generate the asymmetry.

REFERENCES

1. V.A. Kuzmin, V.A. Rubakov and M.E. Shaposhnikov, Phys. Lett. **B155** (1985) 36.
2. P. Arnold and L. McLerran, Phys. Rev. **D37** (1988) 1020.
3. E. Mottola and A. Wipf, Phys. Rev. **D39** (1989) 588; S. Khlebnikov and M. Shaposhnikov, Nucl. Phys. **B308** (1988) 885; M. Dine, O. Lechtenfeld, B. Sakita, W. Fischler and J. Polchinski, Nucl. Phys. **B342** (1990) 381.
4. M. Shaposhnikov, Nucl. Phys. **B287** (1987) 757.
5. L. McLerran, Phys. Rev. Lett. **62** (1989) 1075.
6. A. Cohen, D. Kaplan and A. Nelson, Phys. Lett. **245B** (1990) 561.
7. N. Turok and J. Zadrozny, Phys. Rev. Lett. **65** (1990) 2331.
8. M. Dine, P. Huet, R. Singleton and L. Susskind, Phys. Lett. **257B** (1991) 351.
9. L. McLerran, M. Shaposhnikov, N. Turok and M. Voloshin, Phys. Lett. **256B** (1991) 451.
10. A. Cohen, D. Kaplan and A. Nelson, Phys. Lett. **263B** (1991) 86.
11. A. Cohen, D. Kaplan and A. Nelson, "Baryogenesis at the Weak Phase Transition," ITP preprint NSF-ITP-90-85 (1990).
12. E.W. Kolb and M.S. Turner, *The Early Universe*, Addison-Wesley (1990), Redwood City.
13. A. Linde, *Particle Physics and Inflationary Cosmology*, Harwood (1990) Chur, Switzerland.
14. A.D. Sakharov, Pis'ma Zh. Eksp. Teor. Fix. **5** (1967) 32.
15. For a nice overview with references, see the textbook by Kolb and Turner (ref. 12).
16. See, *e.g.*, J. Ellis, A.D. Linde and D.V. Nanopoulos, Phys. Lett. **118B** (1982) 59.
17. For various suggestions for low temperature (but not electroweak) baryogenesis, see, for example, I. Affleck and M. Dine, Nucl. Phys. **B249** (1985) 361; J. Cline and S. Raby, Phys. Rev. **D43** (1991) 1381.
18. G. 't Hooft, Phys Rev. Lett. **37** (1976) 8; Phys. Rev. **D14** (1976) 3432.
19. N. Manton, Phys. Rev. **D28** (1983) 2019; F. Klinkhammer and N. Manton, Phys. Rev. **D30** (1984) 2212.
20. A. D. Linde, Phys. Lett. **70B** (1977) 306; S. Dimopoulos and L. Susskind, Phys. Rev. **D18** (1978) 4500.
21. M. Dine, Santa Cruz preprint SCIPP 90/27 (1990), in Proceedings of 18th SLAC Summer Institute (1990).
22. M. Mattis, Phys. Rep. **214** (1992) 139.
23. S. Barr and A. Nelson, Phys. Let. **B246** (1990) 141; J. Harvey and M. Turner, Phys. Rev. **D42** (1990) 3344; W. Fischler, G.F. Giudice, R.G. Leigh and S. Paban, Phys. Lett. **B258** (1991) 45.
24. M.E. Shaposhnikov, JETP Lett. **44** (1986) 465; Nucl. Phys. **287** (1987) 757; Nucl. Phys. **299** (1988) 797; A.I. Bochkarev, S. Yu. Khlebnikov and M.E. Shaposhnikov, Nucl. Phys. **B329** (1990) 490.
25. M. Dine, P. Huet, and R. Singleton, Jr., Nucl. Phys. **B375**(1992) 625.
26. M. Dine, P. Huet, R. Leigh, A. Linde and D. Linde, Phys. Lett. **B283** (1992) 319; Phys. Rev. **D46** (1992) 550.
27. A.I. Bochkharev, V.A. Kuzmin, M.E. Shaposhnikov, Phys. Lett. **244B** (1990) 27; N. Turok and J. Zadrozny, Nucl. Phys. **B369** (1992) 729.
28. G. Anderson and L. Hall, Phys. Rev. **D45** (1992) 2685.
29. A.D. Linde, Phys. Lett. **70B** (1977) 306; **100B** (1981) 37; Nucl. Phys. **B216** (1983) 421.
30. B.H. Liu, L. McLerran and N. Turok, Minn. preprint TPI-MINN-92-18-T.
31. D. Brahm and S. Hsu, Caltech preprints CALT-68-1705 and CALT-68-1762 (1991).
32. M.E. Shaposhnikov, Phys. Lett. **277B** (1992) 324.
33. M. Gleiser and E. Kolb, Fermilab preprint PUB-91-305-A (1991).
34. N. Tetradis, DESY preprint DESY-91-151 (1991).
35. G. Anderson, OSU preprint OHSTPY-HEP-T-92-009 (1992).
36. N. Turok, Princeton University preprint PUPT-91-1273.
37. D. Grigoriev, M.E. Shaposhnikov and N. Turok, Imperial preprint Imperial-TP-92-05 (1992).