Rapid anomalous baryon number violating ‘sphaleron’ interactions present at high temperatures in the standard model certainly play an important role in determining the baryon asymmetry of the universe. Despite the large effort invested [1] however, we still lack a convincing theory of baryogenesis at the electroweak phase transition (EPT). What is much better established, is that baryon number was badly broken at temperatures above the EPT [2,3], which is known to greatly alter the more conventional pictures of baryon number production, e.g. in GUTs [4]. Moreover, the necessity of sphaleron reprocessing of particle asymmetries has made it possible to devise new schemes of baryogenesis [5–7]. For example in leptogenesis [5] one first generates an asymmetry in the lepton number $L$, which then becomes (partly) reprocessed into a baryon asymmetry by sphalerons.

Loosely speaking the effect of sphaleron interactions is to break $B + L$ while conserving $B - L$. Then the equilibrium conditions due to various ‘ordinary’ and anomalous interactions imply [3,8], to lowest order, that the $B$ and $L$ asymmetries are proportional to $B - L$. Thus a universe that initially has $B - L = 0$, as is the case in the simplest GUTs, should have zero baryon asymmetry. This realization added motivation for studying the generation of a baryon number at the EPT. Also, any source of new exotic baryon or lepton number violation together with the sphalerons has the potential to destroy the baryon asymmetry, even with a nonzero primordial component of $B - L$. To avoid this outcome, it appeared to be necessary to place stringent limits on various effective operators [8–11]. These limits however, were shown to be weakened significantly due to the effective conservation of $e_R$-number at temperatures $T > \sim 10$ TeV, and to only apply to flavour independent couplings [7,10,12,13].

In more accurate treatment of the equilibrium conditions [13,14], it was shown that even in the case $B - L = 0$ – assuming that the EPT is weakly enough first order – the vacuum mass effects in the broken phase (re)generate a small but nonvanishing equilibrium value for $B$. Still it was maintained that above the EPT temperature $T_C$, $B$ should be zero. In that case, should the EPT be of first order, as might yet be the case even in the minimal standard model [15], the baryon number would again be zero. (Unless of course, it was generated at the phase transition)

In this letter we will show that the baryon and lepton numbers remain nonvanishing throughout the evolution of the universe, even in the case $B - L = 0$, if there are lepton flavour asymmetries. Above $T_C$ a small nonzero $B$ is preserved due to the nonidentical
dispersion relations of particles in a plasma. We will first find the effective thermal masses of particles in a primeval plasma, and then show how their presence alters the chemical equilibrium conditions so as to give rise to a nonzero $B$. We discuss the implications of our results to the Affleck-Dine [6] type baryogenesis, where large initial asymmetries may be generated [16].

When a particle propagates in a plasma, it continuously interacts with other particles in the background. Due to these interactions, for example the Dirac equation for a massless fermion becomes

$$[(1 + a)b + b\bar{u}]\Psi = 0,$$

where $a$ and $b$ are some perturbatively calculable momentum and energy dependent functions [17] that vanish at the zero temperature limit, and $u$ is the four velocity of the plasma. The operator multiplying $\Psi$ in (1) is just the inverse propagator, the zeros of which define the dispersion relation $E = E(k)$. The functional form of the dispersion relation is in general rather complicated, but one can introduce effective ‘thermal’ masses $M(T)$ such that to very good accuracy [17]

$$E(k) \approx \sqrt{k^2 + M^2(T)}.$$  

Thermal masses get contributions from both gauge and Higgs interactions and in general are different for different chiralities and families. Defining $x_i \equiv M_i(T)/T$, one finds the following expressions in the standard model:

$$x_{\ell_i L}^2 = \frac{3}{32}g^2 + \frac{1}{32}g'^2 + \frac{1}{16}h_i^e,$$

$$x_{e_i R}^2 = \frac{1}{8}g^2 + \frac{1}{8}h_i^e,$$

$$x_{q_j L}^2 = \frac{3}{32}g^2 + \frac{1}{288}g'^2 + \frac{1}{6}g_s^2 + \frac{1}{16}(h_{j u}^2 + h_{j d}^2),$$

$$x_{u_j R}^2 = \frac{1}{18}g^2 + \frac{1}{6}g_s^2 + \frac{1}{8}h_j^u$$

$$x_{d_j R}^2 = \frac{1}{72}g^2 + \frac{1}{6}g_s^2 + \frac{1}{8}h_j^d,$$

where $g, g'$ and $g_s$ are the usual $SU(2), U(1)_Y$ and $SU(3)$ couplings and the Yukawa couplings $h^I$ derive from the lagrangean $L_Y = \sum_{ij} (h_{ij}^e \bar{L}_i \Phi_{e_j R} + h_{ij}^u \bar{Q}_i \Phi_{u_j R} +$ 


\[ h^d i \bar{Q}_i \Phi d_{jR} h.c., \] with \( \Phi = i\sigma_2\Phi \). We choose the RH fermion basis such that \( h^f \dagger h^f \) is diagonal, the LH lepton basis such that \( h^e h^e \dagger \) is diagonal, and the LH quarks such that \( h^u h^u \dagger + h^d h^d \dagger \) is diagonal. In this basis the eigenvalues of \( h^\dagger h \) and \( hh \dagger \) are the same, and equal to \( 2m^2_f/v^2 \), where \( v \) is the Higgs vacuum expectation value (\( \approx 246 \text{ GeV} \)). Note that we do not know the eigenvalues of the matrix \( h^u h^u \dagger + h^d h^d \dagger \). However, since the CKM mixing angles are small, we can assume that the eigenvalues \( \sim m^2_u + m^2_d \). In any case, this is not important for our calculation. Numerically one finds\(^1\) for example \( x_{\ell,L} \simeq 0.21, x_{\ell,R} \simeq 0.13 \) and \( x_{tL} \simeq 0.58 \) (see also table 1 below).

Dispersion relations for bosons can be similarly derived and approximated by thermal masses. The thermal mass of the Higgs boson turns out to give \( x_h \simeq 0.59 \), assuming \( T \gg T_C, m_H = 60\text{GeV} \) and \( m_{\text{top}} = 174\text{GeV} \); for a complete expression see e.g. ref. [12], but the actual value of \( x_h \) is not important here. The thermal masses of gauge bosons will turn out to be irrelevant for our purposes.

Thermal masses can have an effect on baryon number through the boundary conditions, such as charge neutrality \( Q_{em} = 0 \) (cf. equation (7) below). Total charges are proportional to asymmetries in particle densities, whereas the equilibrium conditions are relations between chemical potentials. It is then clear that finite mass effects can modify the solutions to equilibrium equations subject to boundary conditions such that nontrivial solutions may exist [14].

Indeed, the occupation number of a given particle species with momentum \( k \) (in the rest frame of the heat bath) is given by

\[ n(k, T, \mu_i) = g(\exp((E(k) - \mu_i)/T) \pm 1)^{-1}, \quad (4) \]

where \( \pm \) refers to fermions (bosons), and \( g \) is the number of internal degrees of freedom of the particle. Then the asymmetry in that particle species is

\[ \Delta n_i = \int \frac{d^3k}{(2\pi)^3} (n(k, T, \mu_i) - n(k, T, -\mu_i)) \]
\[ \equiv \frac{g}{6} (\mu_i T^2) \times (a_+ - a_-), \quad (5) \]

where we have expanded to first order in \( \frac{\mu_i}{T} \), and \( a_+ = 1 \) and \( a_- = 2 \). Assuming that the dispersion relation is of the form (2), the function \( \delta \) becomes (In the notation of

\(^1\) We assumed that the coupling constants take their values measured at the weak scale. More accurately we could write, e.g. \( x_{tL}^2 \simeq 0.23\chi + 0.10 \), with \( \chi = \ln(M_Z/\Lambda_{QCD})/\ln(T/\Lambda_{QCD}) \).
ref. [13] \( \delta = a_\pm - \alpha. \)

\[
\delta_\pm = a_\pm - \frac{6}{\pi^2} \int_x^\infty dy \, y \sqrt{y^2 - x^2} \frac{e^y}{(e^y \pm 1)^2}
\]

(6)

For fermions one finds to first order \( \delta_+ \simeq \frac{3}{2\pi^2} x^2 \), which is a very good approximation for all quarks and leptons. In the bosonic case the integral expression for \( \delta_- \) is nonanalytic near \( x = 0 \), but one can show that for \( x \lesssim 1 \) to good accuracy \( \delta_- \simeq x \).

| Particle      | \( x \) | \( \delta \)   |
|---------------|--------|---------------|
| \( t_L (b_L) \) | 0.58   | 4.9 \times 10^{-2} |
| other \( q_L \) | 0.52   | 4.0 \times 10^{-2} |
| \( t_R \)      | 0.60   | 5.3 \times 10^{-2} |
| other \( q_{uR} \) | 0.49   | 3.6 \times 10^{-2} |
| \( q_{dR} \)    | 0.48   | 3.4 \times 10^{-2} |
| \( i_L \)       | 0.21   | 6.7 \times 10^{-3} |
| \( i_{R} \)     | 0.13   | 2.4 \times 10^{-3} |
| \( H \)         | 0.59   | 0.51          |

Table 1. Shown are the effective masses \( x \equiv M(T)/T \) and corresponding values of \( \delta \)-functions as given by equation (6). We took the QCD coupling corresponding to the scale \( M_Z \) (see footnote 1). Only the top quark Yukawa coupling is large enough to show up in these results. We have used \( m_{\text{top}} = 174 \text{ GeV} \) and \( \alpha_s = 0.11 \).

We will now solve for the baryon number in the presence of finite temperature mass effects in a universe with \( B - L = 0 \), but \( L_i \neq 0 \). At \( T \gtrsim T_C \), there are additional boundary conditions of vanishing total charge- and isospin densities, so that the complete set of boundary conditions to be used is:

\[
Q_{em} = 0 \quad Q_3 = 0 \quad B - L = 0.
\]

(7)

Rapid decays and inverse decays (when kinematically allowed) and various \( 2-2 \) scattering processes induce the usual set\(^2\) of ‘ordinary’ equilibrium relations between

\(^2\) We can assume that the temperature is less than \( \mathcal{O}(\text{few}) \text{TeV} \), so that also right handed electrons are in thermal equilibrium [12].
chemical potentials [8]:

\[ \mu_+ + \mu_0 = \mu_W \]
\[ \mu_{u_R} - \mu_{u_L} = \mu_0 \]
\[ \mu_{d_R} - \mu_{d_L} = -\mu_0 \]
\[ \mu_{i_R} - \mu_{i_L} = -\mu_0 \]
\[ \mu_{d_L} - \mu_{u_L} = \mu_W \]
\[ \mu_{i_L} - \mu_i = \mu_W, \]

where the subscripts indicate particle species — “0” and “−” being the Higgs doublet, \( i \) a neutrino species, and \( i_L \) and \( i_R \) are charged leptons of the \( i \)th generation.

Additionally, there is an equilibrium condition due to anomalous sphaleron processes which, with help of (8), can be written as

\[ 9\mu_{u_L} + 6\mu_W + \mu = 0, \]

where \( \mu \equiv \sum_i \mu_i \). We can use the equations (5) for particle asymmetries along with the equilibrium conditions (8) to write the global charges as

\[ Q = (6 - 2\Delta_u + \Delta_d) \mu_{u_L} - (18 - 4\delta_W - 2\delta_\pm - \Delta_d - \Delta_e) \mu_W \]
\[ + (14 - 2\Delta_{u_R} - \Delta_{d_R} - \Delta_{e_R} - 2\delta_\pm) \mu_0 - 2\mu + \Delta\mu_L + \Delta\mu_R \]
\[ Q_3 = -(11 - 4\delta_W - \delta_\pm - \frac{3}{2}\Delta_{d_L} - \frac{1}{2}\Delta_{e_L}) \mu_W \]
\[ + \frac{3}{2}(\Delta_{d_L} - \Delta_{u_L}) \mu_{u_L} + \frac{1}{2}(\Delta\mu_L - \Delta\mu_\nu) + \mu_0(\delta_0 - \delta_\pm) \]
\[ B = (12 - \Delta_q) \mu_{u_L} + (6 - \Delta_d) \mu_W + (\Delta_{d_R} - \Delta_{u_R}) \mu_0 \]
\[ L = 3\mu - \Delta\mu_\nu - \Delta\mu_L - \Delta\mu_R + (6 - \Delta_e) \mu_W - (3 - \Delta_{e_R}) \mu_0, \]

where we have dropped the common multiplicative factor \( \frac{15}{4\pi^2 g_\ast T} \), (we define \( L \equiv \Delta n/s \), where \( s \) is the entropy density so that \( g_\ast \) is the number of relativistic degrees of freedom taken to be 106.75 from now on) and defined

\[ \Delta\mu_\nu \equiv \sum_i \mu_i \delta_i \]
\[ \Delta\mu_X \equiv \sum_i \mu_i \delta_{iX}, \quad X = L, R \]
\[ \Delta f_X \equiv \sum_i \delta_{f_{iX}}, \quad X = L, R \quad f_i = u_i, d_i, e_i \]

and moreover \( \Delta f \equiv \Delta f_L + \Delta f_R \) and \( \Delta q \equiv \Delta_u + \Delta_d \). Notice that the chemical potential \( \mu_i \) in definitions (11) is always that of neutrinos. We must now apply the boundary
conditions (7) to the expressions (10). At first sight this looks rather messy, but fortunately some of the coefficients in (10) turn out to be zero, so that an analytic solution of decent length can be obtained for $B$.

First of all, dispersion relations of particles within the same weak doublet are identical, so that, $\delta_{fuL} - \delta_{fdL} = 0$, $\delta_0 - \delta_- = 0$, and $\Delta \mu_L = \Delta \mu_\nu$. Then only the term proportional to $\mu_W$ remains in the expression for $Q_3$ and therefore, even with thermal corrections we get $\mu_W = 0$. After this simplification it is rather straightforward to proceed: we use (9) to solve $\mu_{uL} = -\mu/9$ and then solve $\mu_0$ from the charge equation:

$$\mu_0 = \frac{1}{14 - \Delta_1} \left\{ \left( \frac{8}{3} - \Delta_2 \right) \mu - \Delta \mu_L - \Delta \mu_R \right\}, \quad (12)$$

where we used the shorthand notations $\Delta_1 \equiv 2\Delta_{uR} + \Delta_{dR} + \Delta_{eL} + 2\delta_0$ and $\Delta_2 \equiv \frac{1}{9}(2\Delta_u - \Delta_d)$. Finally, the last constraint $B - L = 0$ becomes a consistency condition for $\mu$:

$$\left\{ \frac{13}{3} - \frac{\Delta q}{9} - \left( \frac{8}{3} - \Delta_2 \right) \frac{3 + \Delta_3}{14 - \Delta_1} \right\} \mu = \Delta \mu_L + (1 - \frac{3 + \Delta_3}{14 - \Delta_1})(\Delta \mu_L + \Delta \mu_R), \quad (13)$$

where $\Delta_3 \equiv \Delta_{dR} - \Delta_{uR} - \Delta_{eR}$. If all thermal corrections vanish, equation (13) has only the trivial solution $\mu = 0$, which would immediately lead to $B = L = 0$, as indeed was earlier thought to be the case [8,13,14].

At first sight, (13) might be thought to be giving rise to relatively large asymmetries, because the gauge contributions to effective masses are not particularly small. However, the gauge contributions are the same for all families and their added contribution is proportional to $\mu$. Indeed, we can write the $\Delta \mu_X$-terms in the R.H.S of equation (13) as

$$\Delta \mu_X = \delta^q_X \mu + \sum \delta^Y_{iX} \mu_i, \quad X = L, R \quad (14)$$

where $\delta^q$ and $\delta^Y$ respectively are the gauge and Yukawa mass corrections (see equation (3)). Thus the equation (13) is of the form $A\mu = a\mu + c$, where $c$ is nonzero, and hence a nontrivial solution exists, only because lepton family number is conserved and the lepton Yukawa interactions differ from one family to another. This is in close analogy to the situation when $T \lesssim T_C$ [13,14]. Using (14) one can rewrite (13) as

$$\left\{ \frac{13}{3} - \frac{\Delta q}{9} - \left( \frac{8}{3} - \Delta_2 \right) \frac{3 + \Delta_3}{14 - \Delta_1} - \delta^q_L - (1 - \frac{3 + \Delta_3}{14 - \Delta_1})(\delta^q_L + \delta^q_R) \right\} \mu = (4 - \frac{3(3 + \Delta_3)}{14 - \Delta_1})\Delta \mu_Y, \quad (15)$$

6
which shows explicitly the dependence on Yukawa terms \( \Delta \mu_L^Y \equiv \sum_i \delta_{iL}^Y \mu_i \). In principle it is easy to trace backwards the steps from (15) to (12) to (10) and solve for \( B \) exactly. However, glancing at the size of the various thermal correction terms (see also table 1),

\[
\Delta_q \simeq 0.49 \quad \Delta_1 \simeq 1.39 \quad \Delta_2 \simeq 0.03 \quad \Delta_3 \simeq -0.03 \quad \delta^3_L \simeq 0.007 \quad \delta^3_R \simeq 0.002,
\]

one can see that all of the quantities in (16) are small enough to be neglected. Then one finds a simple expression for \( B \propto \mu_u L \simeq -\frac{4}{3} \mu \),

\[
B \simeq -\frac{94}{79} \cdot \frac{15}{4\pi^2 g_*} (\Delta \mu^Y_L/T).
\]

Including all \( \Delta \)-terms (with \( x_h = 0.59 \)) one obtains \( B \simeq -4.15 \times 10^{-3} \Delta \mu^Y_L/T \), which is less than 3 percent off from the value given by (17)! The Yukawa couplings of leptons are rather small and obey the hierarchy \( h_e \ll h_\mu \ll h_\tau \) (see table (2)), so that to very high accuracy \( \Delta \mu_L^Y \simeq \delta^Y \mu_\tau \propto \delta^Y L_\tau \).

| Particle | \( h_\ell \) | \( x^2_\ell \) | \( \delta^Y \) |
|----------|----------------|-----------------|----------------|
| \( \tau_L, (\nu_\tau) \) | \( 1.03 \times 10^{-2} \) | \( 6.6 \times 10^{-6} \) | \( 1.0 \times 10^{-6} \) |
| \( \mu_L, (\nu_\mu) \) | \( 6.09 \times 10^{-4} \) | \( 2.3 \times 10^{-8} \) | \( 3.5 \times 10^{-9} \) |
| \( e_L, (\nu_e) \) | \( 2.94 \times 10^{-6} \) | \( 5.4 \times 10^{-13} \) | \( 8.2 \times 10^{-14} \) |

**Table 2.** Shown are the Yukawa contributions to the lepton thermal masses and the corresponding values of \( \delta \) functions. For right handed charged leptons use \( x^2_R = 2x^2_L \) and \( \delta_R = 2\delta_L \).

We still need to express \( B \) in terms of primordial quantities. It is easy to show that if \( B - L = 0 \) and \( \frac{1}{3}B - L_i \) are conserved, then

\[
L_\tau(t) = \frac{1}{6} (B + L)(t) + 2\Delta L_{\tau e} + 2\Delta L_{\tau \mu}).
\]

where \( L_\tau \equiv L_{\tau L} + L_{\tau R} + L_{\nu_\tau} \) and \( \Delta L_{ij} \equiv L_i - L_j \). The sphalerons force \( B + L \) and the Higgs chemical potential to be zero, up to the small mass corrections that we are calculating, so that \( L_\tau \simeq \frac{45}{4\pi^2 g_*} (\mu_\tau) \), and

\[
B \simeq -\frac{94}{79} \frac{3}{2\pi^2} x^2_\tau (\Delta L_{\tau e} + \Delta L_{\mu \tau}) \simeq 1.3 \times 10^{-7} (\Delta L_{e \tau} + \Delta L_{\mu \tau}),
\]

where \( x_\tau = \frac{1}{2} x_{\tau L} \).
This is our final result. It should be noted that mass effects alone are not enough to provide a nonzero $B$, although they set the scale of $B$ in terms of primordial asymmetries. Instead, primordial values of leptonic asymmetries must be different, i.e. $\Delta L_{ij} \neq 0$!

Let us compare our result (19) to that of refs. [13,14], where the situation at $T < T_C$ is considered. Equation (23) for $B$ in ref. [13] is similar to our equation (17) and it can also be brought into the form (19). Neglecting all mass corrections except the bare lepton Yukawa terms responsible for a nonzero result, one obtains from ref. [13] $B \simeq \text{few} \times 10^{-7}(\Delta L_{e\tau} + \Delta L_{\mu\tau})$. Of course, if the EPT is of second order or weakly enough first order, it is the latter result that is relevant. However, if the EPT is of first order, then sphaleron processes drop out of equilibrium instantaneously in the broken phase, the equilibrium conditions below $T_C$ never get realized and the value of $B$ given by (19) gets frozen into the system. The point we wish to make is that either way, a nonzero fraction of primordial baryon number survives.

In order for a component of primordial baryon number to survive at the level of the observed asymmetry $B_{\text{now}} = n_B/s \simeq 4 \times 10^{-11}$ [18], primordial asymmetries must be of order $\sim 10^{-4}$. This is perhaps at best marginally possible in the simplest GUT scenarios, but it is much more feasible in the Affleck-Dine mechanism [6], where a baryon asymmetry is generated during the decay of a scalar condensate, which can have a large component of baryon or lepton number stored within. It was previously found [16] that sufficiently large ($B \gtrsim 10^{-2}$) primordial $B + L$ asymmetries can survive in the Affleck-Dine scenario, because the squark condensate does not evaporate before the electroweak phase transition, giving the $W$ a large mass and keeping the sphalerons out of thermal equilibrium. However, a large amount of entropy production is then required after the phase transition. We have shown here that within the Standard Model, given a slightly smaller “large” primordial asymmetry $\sim 10^{-4}$, it would not be washed out, but rather be diluted by the sphalerons to give approximately the right baryon asymmetry today, so that no entropy production is required.

Let us finally discuss the apparent limitation that nonzero $\Delta L_{ij}'s$ are required for the mechanism to work. First of all, it is not at all unnatural to obtain differences in the asymmetries that are of the order of the asymmetries themselves. On the contrary, in e.g. A-D models it is natural to expect that the scalar condensate initially lies in some arbitrary direction in the slepton space, which in general is not symmetric in
flavours. In such case the subsequent evolution and decay of the condensate naturally gives rise to large differences in lepton asymmetries. Thus large $\Delta L_{ij}'$'s certainly can be produced. Secondly, in models with $B - L$ symmetry $\Delta L_{ij}$'s are always conserved quantities. Moreover, even in the case that the lepton number was broken at some higher temperature scale due to some new exotic interactions [9], $B$ typically would not vanish. This is due to the approximate $e_R$-number conservation, which protects the right handed electron asymmetry, and hence a nonzero $\Delta L_{\tau e}$ at any temperature $T \gtrsim 10 \text{ TeV}$ [12]. Thus, except in some very speculative cases (for details see ref. [12]) the survival of a nonzero $B$ is always guaranteed.

In conclusion, we have computed the equilibrium value of the baryon asymmetry $B$ in the early universe above the electroweak phase transition temperature $T_C$. We have found that due to thermal effects (nontrivial dispersion relations) a small, but nonzero value of $B$ persists throughout the evolution of the universe, even in the case $B - L = 0$. The final value of $B$ is roughly seven orders of magnitudes smaller than the primordial value, with the scale set by the Yukawa coupling of the tau lepton. Our results complement the analysis of refs. [13,14], where the vacuum mass effects at the broken phase at $T < T_C$ were shown to also secure a nonzero $B$. Combining these results one has the proof of persistence of a fraction of primordial baryon asymmetry regardless of the order of the electroweak phase transition. This result favours Affleck-Dine type baryogenesis models, which otherwise could produce too large a final baryon asymmetry; the scale of relaxation of $B$ due to modified equilibrium processes can make these theories work even for rather large initial asymmetries.

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Protecting the Baryon Asymmetry with Thermal Masses

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Abstract

We consider the evolution of baryon number $B$ in the early universe under the influence of rapid sphaleron interactions and show that $B$ will remain nonzero at all times even in the case of $B - L = 0$. This result arises due to thermal Yukawa interactions that cause nonidentical dispersion relations (thermal masses) for different lepton families. We point out the relevance of our result to the Affleck-Dine type baryogenesis.