Cosmic $\Delta B$ from Lepton Violating Interactions at the Electroweak Phase Transition

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Abstract

We propose a new mechanism for late cosmological baryon asymmetry in models with first order electroweak phase transition. Lepton asymmetry arises through the decay of particles produced out of equilibrium in bubble collisions and is converted into baryon asymmetry by sphalerons. Supersymmetric models with explicitly broken $R$–parity may provide a suitable framework for the implementation of this mechanism.
The realization that the baryon (\(B\)) and the lepton (\(L\)) violating quantum effects in the standard electroweak theory are efficient at high temperature [1] has sparked a widespread interest on the issue of late cosmological \(B\) asymmetry production. The problem is that any matter–antimatter asymmetry created at some superheavy scale [2] can be easily wiped out by the \(B\)–violating quantum effects, which are abundantly in equilibrium throughout all the history of the early Universe until its temperature drops down to the Fermi scale range [1]. If this is indeed the case, then we are faced with the vital problem of originating a new cosmological \(B\) asymmetry at the scale of the Electro Weak Phase Transition (EWPT).

A key ingredient for the development of a sizeable late cosmological \(B\) asymmetry is that the EPTW be of first order. This seems necessary for the implementation of the so–called out of equilibrium condition which, together with \(B\) and \(CP\) violation, constitutes one of the essential requirements to generate a net \(\Delta B\) [3]. There already exist several studies of \(\Delta B\) production at the electroweak scale in the Standard Model (SM) or in minimal extensions of it [4]. In SM the \(B\)–violating quantum effects play a twofold role [5]: on one hand they are responsible for the washing of any preexisting \(B\) asymmetry, but, on the other hand, they are invoked for the production of the new \(\Delta B\) at the Fermi scale. In general, in order for these scenarios at the electroweak scale to work, the \(B\)–violating quantum effects must be sufficiently suppressed after the accomplishment of the EWPT, or else they can wipe out the \(B\) asymmetry once again. This condition translates into an upper bound on the mass of the light neutral physical Higgs [5,6]. This limit is in contrast with or dangerously close to the present LEP lower bound on Higgs masses [7], depending on the details of the different analyses and whether the SM [5] or the Minimal Supersymmetric Standard Model (MSSM)
are considered [6].

In this paper we focus on a different situation, where at the Fermi scale some new $B$– and/or $L$–violating phenomenon is effective in addition to the $B$– and $L$–violating quantum effects of the SM. In particular, we consider the case in which new $L$–violating interactions are responsible for producing a cosmological $\Delta L$ at the EWPT, while the $B$– and $L$–violating quantum effects are subsequently called into play in order to transform this $\Delta L$ into the cosmological $\Delta B$ \footnote{Baryogenesis via Leptogenesis was first considered in \cite{8}, however, in our model, the lepton asymmetry and the out of equilibrium condition are obtained in completely different manners.}. Notice that in our approach the requirement concerning the quantum effects just after the of the EWPT is opposite to that previously imposed in SM and MSSM: namely, they must be still efficient in order to convert $\Delta L$ into $\Delta B$. Hence, the dangerous upper bound on the physical Higgs mass is removed in our case. We provide an illustration of these ideas in the context of supersymmetric (SUSY) models [9] where $R$–parity is explicitly broken in the leptonic sector [10]. Unless the violation is extremely tiny, the $L$–violating transitions originated by $R$–parity breaking are in equilibrium down to temperatures at which the EWPT occurs. We show that the $R$–violating interactions can originate an adequate $\Delta L$ in theories with a first order EWPT and such $\Delta L$ can be efficiently converted into $\Delta B$ by $B$– and $L$–violating quantum effects still operative at the end of the phase transition.

Whether or not the EWPT in SM is of first order is still an open issue much debated in the present literature [11]. The new twist in this discussion is represented by the inclusion of infrared effects which contribute non–negligibly to the definition of the order of EWPT. It seems that, in any case, the transition may remain first order even after their inclusion, although possibly weaker than what appeared in previous analyses [12]. Even a weak first order EWPT can constitute
a suitable framework for our proposal.

We consider the one–loop effective potential coming from the SUSY model whose Higgs sector contains, in addition to the usual two Higgs doublet superfields $\hat{H}_1^0$ and $\hat{H}_2^0$, one or more singlet superfields $\hat{N}$'s [13], with a typical coupling in the superpotential $\xi_i\hat{H}_1^0\hat{H}_2^0\hat{N}_i$. The theory simplifies in the limit in which the SUSY breaking mass scale $\tilde{m}$ is much larger than the weak scale. Indeed, only one combination of the two neutral Higgs bosons $H_1^0$ and $H_2^0$ remains light. For instance, in the case of one singlet $\hat{N}$

$$h = \cos \beta \text{Re}H_1^0 + \sin \beta \text{Re}H_2^0 + \mathcal{O}\left(\frac{v}{x}\right)\text{Re}N, \quad (1)$$

where

$$\text{Re}\langle H_1^0 \rangle \equiv v_1, \quad \text{Re}\langle H_2^0 \rangle \equiv v_2, \quad \tan \beta \equiv \frac{v_2}{v_1}, \quad (2)$$

and

$$\text{Re}\langle N \rangle \equiv x, \quad v \equiv \sqrt{v_1^2 + v_2^2}, \quad (3)$$

Here we have assumed $x \gg v$.

Hence, at energy scales lower than $\tilde{m}$, the theory approximately reduces to the standard model with only one Higgs doublet. The one–loop potential at finite temperature $V(h, T)$ can be written as the sum of the classical potential, the Coleman–Weinberg quantum corrections [14] and a temperature–dependent part [15]. Calling $T_0$ and $v(T_0)$ the temperature at which the potential is flat at the origin and the vacuum expectation value of $h$ at $T = T_0$, respectively, we obtain

$$\frac{v(T_0)}{T_0} = \frac{\alpha}{\lambda_T}, \quad (4)$$

where $\alpha$ and $\lambda_T$ are the coefficients of the cubic and quartic term in $h$, respectively. $\lambda_T$ can be written as $\lambda_T = \lambda - \frac{1}{2}K$, where $\lambda$ is the coefficient of the $h^4$ term in SM and $K$ denotes the correction coming from the supersymmetric contribution.
to the effective potential. Given our assumption that $\tilde{m} \gg m_W$, we can expand the mass squared of each SUSY particle in power of $\tilde{m}$ [6]

$$m_i^2 = \tilde{m}^2 + g_i h^2 + \mathcal{O}\left(\frac{1}{\tilde{m}^2}\right).$$

(5)

Then $K$ is found to be [6]

$$K = \sum_i \pm \frac{N_i g_i^3}{8\pi^2} \frac{v^2}{\tilde{m}^2}.\quad (6)$$

The sum in (6) extends over all bosons (+) and fermions (−) present in the theory with number of degrees of freedom $N_i$.

A non zero value of $v (T_0) / T_0$ signals a first order phase transition. Here comes the abovementioned crucial difference between baryogenesis schemes where $\Delta B$ is generated before the end of the EWPT and our scenario. In the former case, in order to avoid the cosmic $\Delta B$ be washed out soon after the phase transition, a lower bound on $v (T_0) / T_0$ must be imposed, $v (T_0) / T_0 \gtrsim 1.3$ [5,16]. From eq. (4) it is easy to realize that this bound implies an upper bound on the Higgs boson mass in SM where, roughly, $m_h$ is proportional to $\lambda^{1/2}$. In our scenario the anomalous electroweak processes must be still effective soon after the accomplishment of the EWPT and hence, no upper bound on $m_h$ has to be imposed. However, the presence of a singlet $\hat{N}$ is still relevant also in our treatment. It was recently observed [17] that in SM the EWPT may proceed via percolation of subcritical bubbles, originated by thermal fluctuation, unless the Higgs boson mass, and consequently $\lambda$, are small enough, $m_h \lesssim 100$ GeV, according to the recent estimate of Dine et al. [11]. Thus, if one decreases the value of $\lambda_T$, the first order phase transition is made easier. Taking the result for $K$ from ref. [6] for the case of the SUSY model with two singlets $\hat{N}_1$ and $\hat{N}_2$, one can see that it is indeed possible to lower $\lambda_T$ considerably with respect to its value in SM [8] without decreasing

\footnote{In the presence of only one singlet $\hat{N}$, the fermionic contribution increases $\lambda_T$. We have}
the value of $m_h$. In fact, one must be careful not to decrease $\lambda_T$ too much, or else $v(T_0)/T_0$ becomes too large and the anomalous electroweak phenomena are no longer operative after the EWPT. Taking $m_h \approx m_t \approx 100$ GeV, a value of $\lambda_T \approx 2 \times 10^{-2}$ generates a suitable phase transition for the baryogenesis scenario.

We come now to a short description of the dynamics of bubble nucleation and collision in a first order phase transition. At high temperatures, $T \gg T_0$, the Universe is in the symmetric phase and there exists a unique minimum of $V(h, T)$ at $\langle h \rangle = 0$. As temperatures lowers, a second minimum appears and becomes degenerate with the minimum at $\langle h \rangle = 0$ at

$$T_C = T_0 \left(1 - \frac{2 \alpha^2}{9 \omega \lambda_T} \right)^{-1/2},$$

(7)

where $\omega$ is the coefficient of the $h^2 T^2$ term in the potential.

The Universe decays from its false vacuum state at $\langle h \rangle = 0$ by bubble nucleation. Nucleation takes place at $T_{NUC} \sim T_0 \sim 150$ GeV, if $\lambda_T \approx 2 \times 10^{-2}$. The difference between the false and the true vacuum energy densities at $T \sim T_0$ reads

$$\rho_{VAC} \approx \frac{4 \alpha^2 \omega}{9 \lambda_T^2} T_0^4,$$

(8)

and is transformed into potential energy in the bubble walls whose energy density is (see Enqvist et al. [18])

$$\eta \approx \frac{2^{3/2}}{3^4} \frac{\alpha^3}{\lambda_T^{5/2}} T_0^3.$$  

(9)

At $T < T_{NUC}$ bubbles keep growing. If they expand by a factor $\gamma$ since $T_{NUC}$, their kinetic energy becomes, roughly speaking, $O(\gamma)$ larger than the potential energy [19]. The total energy of the bubble reads

$$E_{TOT} \approx 4\pi R^2 \eta \gamma,$$

(10)

checked that the inclusion of the bosonic contribution may decrease $\lambda_T$ enough only if one performs a fine-tuning.
where $R$ is the radius of the expanding bubble.

When bubbles collide the energy (10) can be released with two different mechanisms: i) direct particle production due to quantum effects [20] and ii) generation of coherent scalar waves [19,20] which decay into light particles (such a production appears to be less useful for our scenario). In both cases the resulting particles are obviously produced with a distribution far from the equilibrium distribution and this represents the crucial ingredient in our scenario to implement the out of equilibrium condition necessary to originate a net $\Delta L$.

Our aim now is to estimate the number of particles of each species generated through the abovementioned direct particle production mechanism. The mean energy of a particle produced in the bubble collision is of order of the inverse thickness of the wall,

$$\langle E \rangle \sim \Delta^{-1} \equiv \Delta^{-1} \gamma,$$

where $\Delta \approx 6\sqrt{2\lambda T}/(\alpha T_0)$ is the size of the wall when it is formed and we are assuming that $\gamma \approx 1/\sqrt{1 - v_w^2}$, $v_w$ being the wall velocity [19].

The number density of particles of species $\tilde{\chi}$ produced in the collision is of order

$$n_{\tilde{\chi}} \approx f_{\tilde{\chi}} \rho_{VAC} \Delta \gamma \tag{11}$$

where $f_{\tilde{\chi}}$ parametrizes the fraction of particles of species $\tilde{\chi}$ produced in the direct bubble collision. We proceed to give an estimate of this factor $f_{\tilde{\chi}}$. Following ref. [20], one computes the probability of producing particles $\tilde{\chi}$’s in the decay of off–shell Higgs fields describing the dynamics of the bubble collision. Such probability reads

$$P \approx \int \frac{d^4k}{(2\pi)^4} |\mathcal{H}(k)|^2 \mathcal{F}(k^2) \Theta \left(k^2 - 4m_{\tilde{\chi}}^2\right), \tag{12}$$

where $k = (k_0, k_x, k_y, k_z)$, $\mathcal{H}(k)$ denotes the Fourier transform of the suitable Higgs field configuration, $\mathcal{F}(k^2)$ is the imaginary part of the two points one–particle irreducible Green function constructed from the $h-\tilde{\chi}-\tilde{\chi}$ coupling and $m_{\tilde{\chi}}$.
represents the mass of the particle $\tilde{\chi}$. The important message of (12) is that particles $\tilde{\chi}$'s with masses larger than $m_h/2$ can be produced through the decays of virtual Higgs fields. A simple ansatz for the field configuration describing the two bubble collision is

$$h(t, z) \approx \begin{cases} 
0 & \text{if } v_w t < z < -v_w t, \quad t < 0 \\
0 & \text{if } -v_w t < z < v_w t, \quad t > 0 \\
v(T_0) & \text{otherwise,}
\end{cases}$$

(13)

where we have approximated the two colliding bubbles as plane–symmetric walls moving along the $z$ direction with velocity $v_w$. The Fourier transform of (13) yields

$$H(k) = \frac{4v_w v(T_0)}{k_0^2 - k_z^2 v_w^2}. $$

(14)

From (12) we can estimate the number of particles $\tilde{\chi}$'s produced per unit area

$$\frac{N}{A} \approx \int \frac{dk_z dk_0}{(2\pi)^2} |H(k_0, k_z)|^2 \mathcal{F}(k_0^2 - k_z^2) \Theta \left( k_0^2 - k_z^2 - 4m_{\tilde{\chi}}^2 \right).$$

(15)

The computation simplifies assuming the step–function configuration (13) and the scattering to be elastic [20]. These approximations give accurate results for $k_z \lesssim \Delta_\gamma^{-1}$ and $k_0 \lesssim \left( \Delta_\gamma^{-1} v_w \right)$. Note that the slow power law–fall at large $k_0^2 - k_z^2 v_w^2$ in expression (14) is due to the fact that $h(t, z)$ is discontinuous. For more realistic thick walls, $H(k_0, k_z)$ would cut off exponentially for $k_z \gtrsim \Delta_\gamma^{-1}$ and $k_0 \gtrsim (\Delta_\gamma v_w)^{-1}$ [20].

If we consider a typical Yukawa coupling of the Higgs field to fermions $\tilde{\chi}$'s, $\mathcal{L} = gh \bar{\tilde{\chi}} \tilde{\chi}$, from eq. (15) we obtain

$$\frac{N}{A} \approx g^2 v(T_0)^2 \ln \left( \frac{\gamma \Delta^{-1}}{2m_{\tilde{\chi}}} \right).$$

(16)

From the last expression we can argue that there is the possibility of producing particles with mass approximately up to $\gamma \Delta^{-1}$. Hence, the value of $\gamma$ becomes crucial in our discussion. For our mechanism for $\Delta L$ production to work it
is important to be able to generate particles with mass larger than \( m_h \) through bubble collisions and, thus, we need a large value of \( \gamma \). The value of this parameter has been the focus of interest of a few recent analyses. In ref. [21] it is claimed that bubbles can expand with \( v_w \) close to the speed of light. However Dine et al. [11] have identified relevant mechanisms which may slow down \( v_w \). Their result depend on microsopic conditions for thermalization processes in the vicinity of the wall and a more complete analysis of the Boltzmann equation to determine the departure from equilibrium in the particle phase–space density across the wall is needed [21]. Moreover, even a small gradient in temperature across the wall can be crucial for accelerating the wall itself, see Enqvist et al. [18] and ref. [22].

The estimate of \( f_{\tilde{\chi}} \) comes from the direct comparison of \( \mathcal{N}/A \) given in (16) and that calculated from (10)

\[
\frac{\mathcal{N}}{A} \approx \frac{E_{\text{TOT}}}{\langle E \rangle A} f_{\tilde{\chi}} \approx \eta \Delta \gamma \ f_{\tilde{\chi}} \approx v(T_0)^2 f_{\tilde{\chi}},
\]

where the last equality is readily obtained using eqs. (4) and (9). We obtain

\[
f_{\tilde{\chi}} \approx g^2 \ln \left( \frac{\gamma \Delta^{-1}}{2m_{\tilde{\chi}}} \right).
\]

We have now all the factors entering the expression (11) for the number density of particles \( \tilde{\chi} \)'s. Then, \( \Delta L \) produced in \( \tilde{\chi} \) decays is

\[
\Delta L \equiv \frac{n_l - n_{\bar{l}}}{s} = \frac{n_L}{s} \approx \frac{\varepsilon_L f_{\tilde{\chi}} \rho_{\text{VAC}} \Delta \gamma^{-1}}{45 g_s s(T_0) T_0^3},
\]

where

\[
\varepsilon_L = \sum_i B_i \frac{\Gamma(\tilde{\chi} \to l_i) - \Gamma(\bar{\tilde{\chi}} \to \bar{l}_i)}{\Gamma(\tilde{\chi} \to \text{all})},
\]

\( B_i \) being the branching ratios of the decay mode \( \tilde{\chi} \to l_i \) and \( g_s s(T_0) = 106.75 \) counts for the relativistic degrees of freedom contributing to \( s \), the entropy density. In (19) we have taken into account that the "reheating" temperature is \( \sim T_0 \).
Using eq. (8) and taking $\alpha \sim 10^{-2}$, $\omega \sim 0.2$, and $\lambda_T \sim 2 \times 10^{-2}$, we obtain from eq. (19)

$$\Delta L \approx 10^{-2} \varepsilon_L f_{\chi} \gamma^{-1}. \quad (21)$$

As we previously mentioned, we consider $R$–parity explicitly broken in the leptonic sector [10] as the concrete framework for the illustration of our ideas. The relevant $L$–violating terms in the superpotential are

$$W_{\Delta L \neq 0} = \frac{1}{2} \lambda_{ijk} \left[ \hat{L}_i, \hat{L}_j \right] \hat{e}^c_k + \lambda'_{ijk} \hat{L}_i \hat{Q}_j \hat{d}^c_k, \quad (22)$$

where $i, j, k$ are generation indices, $\hat{L}$ and $\hat{Q}$ are the left–handed lepton and quark doublet superfields, and $\hat{e}^c$ and $\hat{d}^c$ denote (the charge conjugate of) the right–handed lepton and charge $-1/3$ quark singlet superfields.

There already exist several experimental constraints on the $\lambda$ and $\lambda'$ couplings in (22) [23]. The strongest bounds by far are those obtained from the cosmological considerations on the survival of the cosmic $\Delta B$ [24]. If we ask for the $L$–violating interactions due to the couplings (22) to be constantly out of equilibrium until the weak scale, so that they cannot wash out any $(B - L)$ asymmetry previously generated, a severe upper limit $\lambda, \lambda' \lesssim 10^{-7}$, independently from the generation indices is obtained\footnote{This is true only if one allows for the violation of all partial lepton number symmetries in (22). It was pointed out that one can evade these bounds if one of the lepton numbers $L_e, L_\mu, L_\tau$, say $L_e$, is preserved in (22), so that $\frac{1}{2}B - L_e$ is not washed out [25].}. In our case these limits are no longer valid since we ask for the $L$–violating interactions in (22) to be still operative at the weak scale in order to give rise to the late $\Delta L$. Other phenomenological and astophysical limits [26] exist, though they are substantially less tight, and we shall take them into account.

The scenarios we have in mind for the generation of $\Delta L$ goes as follows. We assume the lightest supersymmetric particle (LSP) to be the lightest neutralino
\( \tilde{\chi} \). Differently from what occurs in MSSM, the LSP can now decay through the \( R \)-violating couplings in (22). Since \( \tilde{\chi} \) possesses decay channels with different \( L \)-numbers, a net \( \Delta L \) can develop in its decay if \( CP \) is violated and the out of equilibrium condition is implemented. \( CP \) violation can be easily provided by the complexity of the \( \lambda \) and \( \lambda' \) couplings in (22). This is a remarkable advantage with respect to late \( \Delta B \) production in SM where \( CP \) violation provided by the CKM phase seems to be too small. Concerning the out of equilibrium requirement, this is automatically satisfied if the \( \tilde{\chi} \)'s are produced in the bubble collision through their couplings to \( h \).

\( \tilde{\chi} \) is generally given by the superposition of the neutralinos of the model (notice, in particular, the presence of the singlet fermion \( \tilde{N} \), in the case of one singlet superfield \( \hat{N} \))

\[
\tilde{\chi} = Z_{11} \tilde{W}_3 + Z_{12} \tilde{B} + Z_{13} \tilde{H}_1^0 + Z_{14} \tilde{H}_2^0 + Z_{15} \tilde{N}.
\tag{23}
\]

As shown in ref. [27], when the LSP is \( \tilde{\chi} \approx \frac{1}{2} \left( \tilde{H}_1^0 + \tilde{H}_2^0 \right) - \frac{1}{\sqrt{2}} \tilde{N} \), it can be very massive, \( m_{\tilde{\chi}} = \mathcal{O}(500) \text{ GeV} \). In this limit, the coupling of \( \tilde{\chi} \) to \( h \) reads

\[
\mathcal{L} = Z_{15} (Z_{14} \cos \beta + Z_{13} \sin \beta) \ h \tilde{\chi}^\dagger \frac{1 - \gamma_5}{2} \tilde{\chi} + \text{h.c.}.
\tag{24}
\]

The coefficient \( g \) in eq. (18) is then

\[
g = \sqrt{2} Z_{15} (Z_{14} \cos \beta + Z_{13} \sin \beta),
\tag{25}
\]

where \( \sqrt{2} \) takes into account that \( \tilde{\chi} \) is a Majorana fermion.

The dominant decay channel of the LSP is expected to be that of fig. 1 [\(^4\)], whose thermally averaged partial width is

\[
\langle \Gamma \left( \tilde{\chi} \to t l_i d_k^c \right) \rangle = \Gamma \left( \tilde{\chi} \to t l_i d_k^c \right) \frac{K_1(m_{\tilde{\chi}}/T)}{K_2(m_{\tilde{\chi}}/T)},
\tag{26}
\]

\(^4\)If the \( R \)-couplings are nonvanishing, LSP can decay also into two body final states via one-loop diagrams, for instance into \( W l_i \), but these channels are suppressed by a factor \( \mathcal{O}(100) \) in comparison to that of fig. 1., see Campbell et al. [24].
where $K_n$ denotes the modified Bessel functions and the prefactor in front of the usual (zero-temperature) decay rate can be interpreted as a time dilation factor. A straightforward calculation leads to

$$\langle \Gamma (\tilde{\chi} \rightarrow t l_i d_k^c) \rangle \approx \frac{\alpha W}{16} \left( \frac{m_t}{m_W} \right)^2 \frac{1}{768 \pi^3} \frac{m_{\tilde{\chi}}^5}{(m_{t_{L}}^2 + T_0^2)^2} |\lambda_{13k}'| \times$$

$$\times \left( 1 - \frac{3 T_0}{2 m_{\tilde{\chi}}} + \frac{15 T_0^2}{8 m_{\tilde{\chi}}^2} \right) Z_{14},$$

where, in our case, $Z_{14} = \frac{1}{2}$.

The expansion rate of the Universe is given by $H \simeq 17 T^2/M_P$, where $M_P \simeq 1.2 \times 10^{19}$ GeV is the Planck mass. Neutralinos can decay after their production if, roughly speaking, $\Gamma \gtrsim H$, which translates into the bound

$$|\lambda_{13k}'| \gtrsim 9.2 \times 10^{-5} \left( \frac{500 \text{ GeV}}{m_{\tilde{\chi}}} \right)^{5/2} \left( \frac{T_0}{150 \text{ GeV}} \right) \left( \frac{m_{t}}{1 \text{ TeV}} \right)^2 \left( \frac{100 \text{ GeV}}{m_{\tilde{\chi}}} \right).$$

The interference of the diagrams of figs. 1 and 2 gives rise to a non vanishing $\varepsilon_L$ of order

$$\varepsilon_L \simeq \frac{1}{16\pi} \frac{\sum_{i,k,m,n} (\lambda'_{1t_{m}}\lambda'_{s_{m}}\lambda'_{s_{3k}}\lambda'_{13k})}{\sum_{i,k} |\lambda'_{13k}|^2}.$$  

Actually, this expression must be multiplied by a suppression factor which takes into account that part of the neutralinos produced out of equilibrium by our mechanism may thermalize before decaying, e.g. through the $\tilde{\chi} t \rightarrow \tilde{\chi} t$ scattering, and, consequently, they cannot give contribution to $\varepsilon_L$. A rough estimate for such a suppression factor is given by $\langle \Gamma (\tilde{\chi} \rightarrow t l_i d_k^c) \rangle / \langle \Gamma (t \tilde{\chi} \rightarrow t \tilde{\chi}) \rangle \approx 10^3 |\lambda_{13k}'|^2$.

In order to prevent a washing out of the lepton asymmetry, we impose that $\Delta L = 1$ scatterings induced by $\lambda$ and/or $\lambda'$ couplings are out of equilibrium at $T_0$. The predominant process is shown in fig. 3. Since $T_0 < m_{\tilde{\chi}}$, this process has a thermally averaged rate
\[ \Gamma_{\Delta L=1} = n_{EQ}(T) \langle \sigma v \rangle \simeq \frac{27Z_{14}^2}{32\pi^3 \zeta(3)} |\lambda'_{i3k}|^2 \frac{T_0^3 m_{\tilde{\chi}}^2}{256} \alpha_W \times \left( \frac{m_t}{m_W} \right)^2 \frac{1}{(m_{t_L}^2 + T_0^2)^2} e^{-m_{\tilde{\chi}}/T_0}, \]  

(30)

where \( \zeta(3) \approx 1.202 \) is the Riemann zeta function and \( n_{EQ} \) denotes the equilibrium number density. Requiring \( \Gamma_{\Delta L=1} \lesssim H \) imposes

\[ |\lambda'_{i3k}| \lesssim 4.2 \times 10^{-4} \left( \frac{m_{t_L}}{1 \text{ TeV}} \right)^2 \left( \frac{500 \text{ GeV}}{m_{\tilde{\chi}}} \right) \left( \frac{150 \text{ GeV}}{T_0} \right)^{1/2} \left( \frac{100 \text{ GeV}}{m_t} \right). \]  

(31)

A similar upper bound is found requiring that inverse decays are out of equilibrium. The products of neutralino decays will be brought to thermal equilibrium at rates of order \( T_0 \). At this temperature, anomalous \((B+L)\)–violating processes are in equilibrium, since \( \Gamma_{\text{sphaleron}} \simeq \alpha_W^4 Te^{-\frac{4\pi}{g_{*W} S(T_0)}} \gtrsim H \) for \( \lambda_T \approx 2 \times 10^{-2} \) and the lepton asymmetry is converted into a baryon asymmetry \( \Delta B = n_B/s \), whose present value is \( (T_{\text{tod}} \approx 2.7^{0\text{K}}) [28] \)

\[ \Delta B \approx -\frac{21}{61} \frac{g_{*S}(T_{\text{tod}})}{g_{*S}(T_0)} \Delta L, \]  

(32)

where \( g_{*S}(T_{\text{tod}}) / g_{*S}(T_0) \) takes into account the increase of entropy since the EWPT up to now.

From eqs. (21), (29) and (32), we argue that a baryon asymmetry \( 4 \times 10^{-11} \lesssim \Delta B \lesssim 5.7 \times 10^{-11} \) [29] can be produced for resonable values of the \( R \)–violating couplings. For instance, taking \( m_{t_L} \approx 5 \text{ TeV}, \gamma \approx 10^2, \lambda_T \approx 2 \times 10^{-2} \) and all the \( \lambda' \) couplings of the same order, we obtain \( \lambda' \approx 8 \times 10^{-3} \). Note that this value is in agreement with the upper bounds given in [26].

Notice that our scheme makes use of the \( R \)–parity violating couplings in the leptonic sector. Previous proposals to use \( R \)–parity violation for late \( \Delta B \) production invoked \( R \)–breaking in the baryonic sector within contexts quite different to implement the out of equilibrium condition [30].
In conclusion, we have proposed a scheme for late baryon asymmetry in a first order EWPT through the interplay of lepton violating interactions and anomalous $B$– and $L$–violating quantum effects: $i$) $\Delta L$ is generated in the decay of the LSP produced out of equilibrium in bubble collision and $ii$) $\Delta L$ is converted into $\Delta B$ by sphalerons still operative after the EWPT. Obviously, a thorough discussion of several open issues of EWPT is eagerly awaited for to provide a more comprehensive and detailed scheme for late $\Delta B$ production along the lines suggested in this paper.

Acknowledgements

We should like to thank R. Barbieri, K. Enqvist, G. Giudice, E.W. Kolb, S. Matarrese, J. Miller, O. Pantano, S. Petcov and F. Zwirner for many helpful and stimulating discussions.
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**Figure Captions**

**Figure 1.** The tree–level Feynman diagram for the decay $\tilde{\chi} \to t l_i d_k$ induced by the $\lambda_{3ik}'$ couplings.

**Figure 2.** The one–loop Feynman diagram for the decay $\tilde{\chi} \to t l_i d_k$, whose interference with that of Figure 1 gives rise to a non vanishing lepton asymmetry.
Figure 3. The Feynman diagram for the $\Delta L = 1$ scattering $l_i d_k \rightarrow t \tilde{\chi}$, which can wash out any lepton asymmetry.