Abstract

The possibility of baryogenesis at the electroweak phase transition is considered within the context of a minimal supersymmetric standard model with spontaneous R-parity violation. Provided that at least one of the sneutrino fields acquires a large enough vacuum expectation value, a sufficient baryon asymmetry can be created. Compared to R-parity conserving models the choice of soft supersymmetry breaking parameters is less restricted. The observed baryon asymmetry, $n_B/s \sim 10^{-10}$, can be explained by this scenario and the produced baryon-to-entropy ratio may easily be as high as $n_B/s \sim 10^{-9}$. 

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One of the major open questions in theoretical particle physics is to explain the observed baryon asymmetry of the universe. The well-known Sakharov conditions for baryogenesis, namely T, B, C and CP-violation [1], can be fulfilled in the early universe during a phase transition for suitable models of particle physics. In particular, it has been shown that the Standard Model (SM) includes all the requirements for electroweak baryogenesis [2]. However, the electroweak phase transition is too weakly first order to preserve the generated baryon asymmetry in the SM [3]. Furthermore, the CP-violation due to the phase of the Cabibbo-Maskawa-matrix is too small for a sufficient generation of baryons [4]. The realization of electroweak baryogenesis therefore requires physics beyond the SM.

The Minimal Supersymmetric Standard Model (MSSM) has appeared to be a promising candidate for the explanation of the origin of the observed baryon asymmetry in the universe. It has been shown that with an appropriate choice of parameters a large enough baryon asymmetry, $n_B/s \sim 10^{-10}$, could have been created at the electroweak phase transition [4, 5]. The necessary choice of parameters involves, however, quite stringent bounds. Furthermore the choice of $|\sin \phi_\mu|$, where $\phi_\mu$ is the phase of one of the soft supersymmetry breaking parameters $\mu$, has been constrained even more in a more recent analysis of the bubble wall profile [6]. It is therefore useful and interesting to expand the idea of electroweak baryogenesis to less constrained supersymmetric models. In this paper we shall consider the creation of baryon asymmetry in the context of a model where R-parity is spontaneously violated [8].

The Minimal Supersymmetric Standard Model (MSSM) assumes the conservation of a discrete symmetry, R-parity [9, 10]. R-parity is related to the spin (S), total lepton (L) and baryon (B) number of a particle by $R_p = (-1)^{3B+L+2S}$ so that all the standard model particles have an even R-parity whereas their supersymmetric partners have an odd R-parity. There is, however, no a priori reason to expect that R-parity is conserved. Therefore it can be a spontaneously or even explicitly broken symmetry. A model that spontaneously violates R-parity proposed in ref. [11] has a superpotential of the form (suppressing all $SU(2)$ and generation indices)

$$h_u \hat{Q} \hat{H}_u \hat{u}^c + h_d \hat{H}_d \hat{Q} \hat{d}^c + h_e \hat{H}_d \hat{e}^c + (h_d \hat{H}_d \hat{H}_d - \mu^2) \hat{\Phi} + h_v \hat{H}_u \hat{v}^c + h \hat{\Phi} \hat{S} \hat{v}^c + h.c., \quad (1)$$

which conserves both total lepton number and R-parity (hat denotes a superfield). The new superfields ($\hat{\Phi}, \hat{v}^c, \hat{S}_i$) are $SU(2) \otimes U(1)$ -singlets and carry conserved lepton numbers $(0, -1, -1)$ respectively.

In this paper we shall focus on the $h_v \hat{H}_d \hat{e}^c$ term. Using standard methods [3], one can show that the following terms are present in the Lagrangean:

$$\mathcal{L}_R = \overline{H}[h_{\nu ij} \tilde{\nu}^e_1 R_i + h_{ekj} \tilde{\nu}^e_2 L_i] e_j, \quad (2)$$
where
\[ \tilde{H} = \begin{pmatrix} \tilde{H}_+^d \\ \tilde{H}_-^u \end{pmatrix} \] (3)
and \( R, L \) are the chirality projection operators.

The Higgsino current associated with charged higgsinos can be written as
\[ J^\mu_H = \bar{\tilde{H}} \gamma^\mu H. \] (4)

The Higgsino current (4) is associated with a triangle diagram where Higgsinos interact with electron-like leptons, similar to the diagram in ref. [6]. The CP-violating source in the diffusion equation is in this case created by interactions arising from the term
\[ \mathcal{L}_R = h_{eij} \bar{H} v_i^L e_j, \] (5)
where \( v_i \equiv< \tilde{\nu}_{Ri} > \). It is worth noting that a non-vanishing source requires only that one of the sneutrino fields acquires a vev. This is due to a complex vev so that two degrees of freedom are present. In a spontaneously R-parity breaking model the global B-L symmetry is spontaneously broken by assuming non-zero vev’s for \( v_R \equiv< \tilde{\nu}_{R3} >, v_S \equiv< \tilde{S}_3 > \). In general we may also choose a non-zero value for \( v_L \equiv< \tilde{\nu}_{L3} > \). However, spontaneous R-parity breaking generates a majoron that are produced in a stellar environment and to suppress the stellar energy loss one must require \( \frac{v_L^2}{v_R} m_W \lesssim 10^{-7} \). This is easily achieved for \( v_R = \mathcal{O}(1 \text{ TeV}) \) provided that \( v_L \leq \mathcal{O}(100 \text{ MeV}) \). For generality we consider here the case where all left- and right-handed sneutrinos have a non-zero vev. In numerical estimations we make the usual assumption that only \( v_{R3} \) is large.

If one chooses interactions similar to ref. [6] where both left- and right-handed electrons are included the amount of CP-violation is proportional to \( \tilde{\nu}_R \tilde{\nu}_L \) -term. However, if we consider only interactions like (5), the total CP-violation has a factor \( \tilde{\nu}_R^2 \), which, assuming that \( |v_R| \gg |v_L| \), can give a much larger contribution.

The scenario considered here sets some limitation to the hierarchy of the electroweak phase transition, namely that the higgsino and sneutrino field(s) acquire their respective vev’s simultaneously. Otherwise sphaleron transitions wash out the previously created baryon asymmetry.

The CP-violating source may now be computed using CTP-formalism as described in ref. [12]. The interactions of the Higgsinos lead to a contribution to the self-energy of the form
\[ \Sigma_{CP}(x, y) = g^L_{CP}(x, y) R G_\nu^o <(x, y)L + g^L_{CP}(x, y) L G_\nu^o <(x, y)R, \] (6)
where
\[ g^L_{CP} = h_{eij} v_i(x) h_{ekj} v_k^*(y). \] (7)
Similarly for the other component $\Sigma_{CP}$. High temperature corrections change the dispersion relation of charginos and neutralinos. The spectral function of Higgsinos may, in the approximation $\Gamma_{\tilde{H}} \ll m_{\tilde{H}}$, be written as

$$\rho_{\tilde{H}}(k, k^0) = i(\not{k} + m_{\tilde{H}})[((k^0 + i\epsilon + i\Gamma_{\tilde{H}})^2 - \omega_{\tilde{H}}^2(k))^{-1} - ((k^0 - i\epsilon - i\Gamma_{\tilde{H}})^2 - \omega_{\tilde{H}}^2(k))^{-1}],$$

(8)

where $\omega_{\tilde{H}}^2(k) = k^2 + m_{\tilde{H}}^2(T)$. The effective squared Higgsino plasma mass $m_{\tilde{H}}^2(T)$ may be approximated by $m_{\tilde{H}}^2(T) \approx |\mu|^2$. Similarly, $|\mu|$ should be replaced by the electron plasma mass $\frac{3}{2}g_2^2 T^2$ for $\rho_{e_i}(k, k^0)$.

The source term

$$S_{\tilde{H}} = -\int d^3r_3 \int_{-\infty}^{T} \text{tr} \left[ \Sigma_{CP}^{>}(X, x_3) G_{\tilde{H}}^{<}(x_3, X) - G_{\tilde{H}}^{<}(X, x_3) \Sigma_{CP}^{>}(x_3, X) \right] + G_{\tilde{H}}^{<}(X, x_3) \Sigma_{CP}(x_3, X) - \Sigma_{CP}^{<}(X, x_3) G_{\tilde{H}}^{>}(x_3, X)$$

(9)

now contains the following function

$$f_g = g_{CP}^L(X, x_3) - g_{CP}^R(x_3, X)$$

$$= i \text{Im} [(h_{eij}^* h_{ekj})(v_i(X) v_k^*(x_3) - v_i(x_3) v_k^*(X))],$$

(10)

which after the Higgs insertion expansion becomes

$$f_g = i \text{Im} [(h_{eij}^* h_{ekj})(v_i(X) \partial_X^k v_k^*(X) - v_i^*(X) \partial_X^k v_k(X))].$$

(11)

The possibility of complex Yukawa coupling constants, $h_e$, deserves further discussion. The $h_e$ matrix clearly has 18 degrees of freedom i.e. 9 complex phases are present.

From (5) we note that 7 of the complex phases may be set zero by redefining the sneutrino, higgsino and electron-like lepton fields. Generally we then have two complex phases left. For simplicity, we shall from now on assume that all Yukawa coupling constants are real i.e. $h_{eij} \in \mathbb{R}$, $i, j = 1, 2, 3$. However, the sneutrino vev’s may be complex, so that we may write

$$v_i(X) = A_i(X) e^{i\theta _i(X)},$$

(12)

where $A_i(X), \theta_i(X) \in \mathbb{R}$. Function $f_g$ can now be written as

$$f_g = i(h_{eij} h_{ekj}) \left[ \sin(\theta_i - \theta_k)[A_i A_k'] - A_i A_k' \cos(\theta_i - \theta_k) A_i A_k(\theta_i + \theta_k)' \right].$$

(13)

The first term is of the form presented in (6) with an extra factor of $\sin(\theta_i - \theta_k)$. The second term is of a new form, which exists due to the non-vanishing complex phases of the sneutrino vev’s.

We can now write the expression for the CP-violating source,

$$S_{\tilde{H}} = 2i(h_{eij} h_{ekj}) \left[ \sin(\theta_i - \theta_k)[A_i A_k' - A_k A_i'] - \cos(\theta_i - \theta_k) A_i A_k(\theta_i + \theta_k)' \right] \mathcal{I}^{e_{ij}}_{\tilde{H}},$$

(14)
where $I^e_j$ is slightly modified from the result in [12],

$$I^e_j = \int_0^\infty \frac{dk}{2\pi^2} k^2 (\omega_{ej}^2 - k^2) \left[ (1 - 2\text{Re}(n_{ej})) F(\omega_{ej}^H, \Gamma_{ej}, \Gamma_{e_j}) + (1 - 2\text{Re}(n_{ej})) F(\omega_{ej}, \Gamma_{ej}, \omega_{ej}^H, \Gamma_{e_j}) + 2\left( \text{Im}(n_{ej}) + \text{Im}(n_{ej}) \right) G(\omega_{ej}^H, \Gamma_{ej}, \Gamma_{e_j}) \right]$$

where $n_{ej} = \left[ \exp \left( \frac{\omega_{ej}^H (e_{ej})}{T} + i\Gamma_{ej} \right) \right] + 1$ and

$$F(a, b, c, d) = \frac{1}{2} \left[ (a + c)^2 + (b + d)^2 \right]^{-1} \sin[2\arctan \frac{a + c}{b + d}],$$

$$G(a, b, c, d) = -\frac{1}{2} \left[ (a + c)^2 + (b + d)^2 \right]^{-1} \cos[2\arctan \frac{a + c}{b + d}].$$

The damping rate of Higgsinos and electron-like leptons is dominated by weak interactions so that we may take $\Gamma_{ej}^H \approx \Gamma_{ej} \sim 0.05 \ T$.[6].

We shall now review the basic ingredients of electroweak baryogenesis as presented in ref. [6]. In the paper the baryon to entropy ratio was shown to be

$$\frac{n_B}{s} = -g(k_i) \frac{\mathcal{A} \bar{D} \Gamma_{ws}}{v_w^2 s},$$

where $g(k_i)$ is a numerical coefficient depending on the degrees of freedom, $\bar{D}$ the effective diffusion rate, $\Gamma_{ws} = 6\kappa \alpha_w^4 T (\kappa = 1)$ [14] the weak sphaleron rate and $v_w$ the velocity of the bubble wall. The entropy density $s$ is given by

$$s = \frac{2\pi^2 g_{ss} T^3}{45},$$

where $g_{ss}$ is the effective number of the relativistic degrees of freedom. $\mathcal{A}$ was shown to be a weighted integral over the CP-violating source $\tilde{\gamma} = v_w f(k_i) \partial_u J^0(u)$:

$$\mathcal{A} = \frac{1}{\bar{D}^+} \int_0^\infty du \tilde{\gamma}(u) e^{-\lambda^+ u},$$

where $\lambda^+ = (v_w + \sqrt{v_w^2 + 4\bar{D}^+ \bar{\mathcal{D}}})/(2\bar{\mathcal{D}})$ and the wall was defined to begin at $u = 0$, where $u$ denotes the co-moving coordinate $u = z + v_w t$ and the wall is assumed to move in the direction of the positive $z$-axis. $(f(k_i))$ is a coefficient depending on the number of degrees of freedom present in thermal path and related to the definition of the effective source [6, 15].) In the notation of the current paper and after performing
a partial integration (assuming the CP-source vanishes at \(u = 0\) and \(u \to \infty\)) the coefficient \(A\) is related to the shape of the bubble wall through

\[
A \propto I_1 = \int_0^\infty du S_H(u) e^{-\lambda+u} \\
= 2i(h_{eij} h_{ekj}) \left( \int_0^\infty du \left[ \sin(\theta_i - \theta_k) [A_i A'_k - A_k A'_i] - \cos(\theta_i - \theta_k) A_i A_k (\theta_i + \theta_k)' \right] e^{-\lambda+u} \right) \mathcal{I}_{ij}^e.
\]

(20)

In comparison, in the paper by Carena et al. the coefficient \(A\) is related to the shape of the bubble wall through

\[
A \propto I_2 = \int_0^\infty \frac{\partial}{\partial u} (H_1 H'_2 - H_2 H'_1) e^{-\lambda+u}, \\
\equiv \int_0^\infty du \frac{\partial}{\partial u} (H(u)^2 \frac{\partial \beta(u)}{\partial u}) e^{-\lambda+u},
\]

(21)

where \(H = (H^2_1 + H^2_2)^{1/2}\), \(\tan \beta = H_1/H_2\) and \(H_i's\) are the real parts of the neutral components of the Higgs doublets. In the paper by Riotto, the CP-violating source associated with the charginos was considered and in that case \(I_2\) can be written as

\[
A \propto I_2 = \text{Im}(\mu) \left[ \int_0^\infty du H(u)^2 \frac{\partial \beta(u)}{\partial u} e^{-\lambda+u} \right] [3M_2 g_2^2 I_{ij}^{W} + M_1 g_1^2 I_{ij}^{B}],
\]

(22)

where \(M_i\) are the soft supersymmetry breaking parameters, \(I_{ij}^{W,B}\) are corresponding integrals for the winos and the bino similar to \(I_{ij}^{e}\).

(20) can be written as (writing out the sums explicitly)

\[
I_1 = 4i \sum_j \left[ - \sum_i h^2_{ij} \int_0^\infty A^2_i \theta'_i e^{-\lambda+u} du \\
+ \sum_{k<i} h_{ij} h_{kj} \int_0^\infty \left[ \sin(\theta_i - \theta_k) [A_i A'_k - A_k A'_i] - \cos(\theta_i - \theta_k) A_i A_k (\theta_i + \theta_k)' \right] e^{-\lambda+u} du \right] \mathcal{I}_{ij}^e
\]

\[
\equiv 4i \sum_j (B_1 + B_2 + B_3) \mathcal{I}_{ij}^e.
\]

(23)

In the wall shape was assumed to take a simple sinusoidal form where the field \(H(u)\) can be given by

\[
H(u) = \frac{v}{2} \left[ 1 - \cos \left( \frac{u \pi}{L_w} \right) \right] [\theta(u) - \theta(u - L_w)] + v \theta(u - L_w)
\]

(24)

and the angle \(\beta(u)\) by

\[
\beta(u) = \frac{\Delta \beta}{2} \left[ 1 - \cos \left( \frac{u \pi}{L_w} \right) \right] [\theta(u) - \theta(u - L_w)] + \Delta \beta \theta(u - L_w),
\]

(25)
where $\Delta \beta$ is given by $\Delta \beta = \beta(T_0) - \arctan(m_1(T_0)/m_2(T_0))$, calculated at the temperature where the curvature of the one-loop effective potential vanishes at the origin. (It is thus the angle between the flat direction and vacuum direction.) We shall assume that both the modulus and the phase of the sneutrino vev’s have a similar sinusoidal form (24). In this approximation, one can write an analytical expression for the integral in (22),

$$I_2 \propto \frac{\pi^2}{4} \frac{2 \lambda_+^4 L_w^4 + 20 \lambda_+^2 L_w^2 \pi^2 + 3 (11 + 5 \lambda_+ L_w) \pi^4}{(\lambda_+^6 L_w^6 + 14 \lambda_+^4 L_w^4 \pi^2 + 49 \lambda_+^2 L_w^2 \pi^4 + 36 \pi^6)} \Delta \beta v^2$$

where $L_w$ is the width of the bubble wall and $v^2 = <H_1>^2 + <H_2>^2$.

Similarly we can now estimate the first term in (23),

$$\mathcal{B}_1 = -\sum_i h_{eij}^2 \theta_f f_1(L_w),$$

where $\theta_f \equiv <\theta_i>$. Here it is assumed that the phases of sneutrino fields acquire their vev’s sinusoidally from $<\theta_i> = 0$ to $<\theta_i> = \theta_f$. $\mathcal{B}_{2,3}$ cannot be solved analytically. However, before numerical analysis, some special cases are worth a further study. Let us define new phase variables, $\psi_{ik}(X) \equiv \theta_i(X) - \theta_k(X)$ and $\phi(X)_{ik} \equiv \theta_i(X) + \theta_k(X)$. Now, if $\psi_{ik}(X) = \psi_{ik}$, $\mathcal{B}_{1,3}$ clearly vanish and we are left with an expression similar to (26) with an additional $\sin \psi_{ik}$ factor,

$$I_1^+ = +4i \sum_j \sum_{k<i} h_{eij} h_{ekj} f_1(L_w) \sin \psi_{ik}(v_i^2 + v_k^2) \Delta \beta_{ik} \tilde{T}_{ij}^{ik},$$

where $\Delta \beta_{ik}$ is the angle between the flat and the vacuum direction at the origin. If, on the other hand, $\psi_{ik}(X) = 0$ i.e. all sneutrino fields acquire the vev of their phases in uniform, we are left with

$$I_1^- = -4i \sum_j \left[ \sum_i h_{eij}^2 \theta_f f_1(L_w) + 2 \sum_{k<i} h_{eik} h_{ijk} v_i v_k \theta_f f_1(L_w) \right] \tilde{T}_{ij}^{ik},$$

where $\theta_f \equiv <\theta_i>$ (all phases have the same vev). Clearly since $\mathcal{B}_2$ is the only positive term and $\sin \psi_{ik} \leq 1$, $I_1^+$ gives the maximum value of $I_1$. Similarly $I_1^-$ gives the minimum value. However, the overall sign of $I_1$ is not significant so that we are only concerned with the absolute value of $I_1$. Since $I_1^+$ and $I_1^-$ have opposite signs, one of them give the maximum value of $|I_1|$ depending on the choice of parameters. Generally $I_1^-$ is greater in magnitude than $I_1^+$ due to the $\Delta \beta_{ik}$ factor in $I_1^+$. To compare the magnitude of the CP-violating source considered in the present paper with previous results, numerical estimates are in order. In choosing parameters we shall follow closely ref. [3]. The bubble wall width is chosen to be $L = 25/T$ (our results are quite insensitive to the exact choice of the bubble wall width).
First we shall consider (22). We choose \( M_2 \approx |\mu| \approx M_1 \approx T \approx 100 \text{ GeV} \) and \( \mathcal{I}_{H}^{W} \approx \mathcal{I}_{H}^{B} \) to give us an estimate of the effect of the CP-violating source. With these values \( \sin \theta_\mu \) is roughly of the order 0.1 \[6\] and \( M_2 |\mu| \mathcal{I}_{H}^{W} \) can numerically be shown to have a value of about 10. Putting all these values into (22), we get an estimate \( I_2 \approx 0.35 v^2 \Delta \beta \). Estimating \( v = 246 \text{ GeV} \) and \( \Delta \beta \approx 0.015 \) \[6\] we finally arrive at an estimate \( I_2 \sim 300 \). Recent calculations using a more realistic bubble wall shape have lead to an additional suppression factor of \( \sim 0.3 \) \[7\]. This effect is ignored here since the analytical approximation is used in both cases.

Similarly we can now estimate (23). Choosing \( L_w = 25/T \) we first find that \( f_1(L_w) \approx 0.25 \) for \( T \sim 100 \text{ GeV} \). Values of \( \mathcal{I}_{H}^{\tau} \) are shown in fig. 1. The values are very indifferent to the lepton type considered so that we may very accurately take \( \mathcal{I}_{H}^{\tau} = \mathcal{I}_{H}^{\tau} \equiv \mathcal{I}_{H} \). As from fig. 1 can be seen, if we choose \( \mu \approx T \approx 100 \text{ GeV} \), \( \mathcal{I}_{H} \) is about −2. From the figure it is clear that there exists a resonance at value \( \mu \approx T \). However, this resonance is not as pronounced as in ref. \[6\].

Numerical examinations show that the integral in \( \mathcal{B}_2 \) is approximately linear in \( \psi_{ik} \) and we may estimate it with \( \sim 0.18 \psi_{ik} \) (no significant \( T \) dependence). Furthermore the integral in \( \mathcal{B}_3 \) is approximately constant and we may take it to be \( \sim 0.05 \). Substituting these values into (23), we then have

\[
I_1 \approx -i8 \sum_{j} \left[ -0.25 \sum_{i} h_{eij}^2 v_i^2 \theta_{f_1} + 0.18 \sum_{k<i} h_{eij} h_{ekj} (v_i^2 + v_k^2) \Delta \beta_{ik} \psi_{ik} - 0.05 \sum_{k<i} v_i v_k \phi_{ik} \right].
\]

(30)

Here we have several unknown parameters. To estimate the numerical value of (30), let us first assume that \( v_3 \gg v_1, v_2 \) and \( \theta_{f_1} = \theta_{f_2} \approx 0 \) i.e. only the tau-sneutrino (we
may have as well chosen any other generation) has a significantly large vev. In this approximation
\[ I_1 \approx -i8\epsilon v_3^2\theta_{f_3} \left[ -0.25 \sum_j h_{e3j}^2 + 0.18(\Delta\beta_{31} \sum_j h_{e3j}h_{e1j} + \Delta\beta_{32} \sum_j h_{e3j}h_{e2j}) \right]. \] (31)

If we now choose \( v_3 = 1000 \) GeV, estimate \( h_{eij} \sim 10^{-1} \) and take a conservative estimate for \( \Delta\beta_{ij} \sim 0.1 \), we get an estimate \( I_1 \approx -50000 \theta_{f_3} \), which is significantly greater than the previous estimate provided that \( \theta_{f_3} \gtrsim 0.01 \). The values of the different contributions are in this case \( I_1^+ \sim 10000 \sin\theta_{f_3} \) and \( I_1^- \sim -60000 \theta_{f_3} \). Even if \( h_{eij} \sim \mathcal{O}(10^{-2}) \) and \( \theta_{f_3} \sim \mathcal{O}(1) \), \( I_1 \) is still large enough to explain the origin of the baryon asymmetry of the universe (even when accounting for the extra suppression factor due to a more realistic bubble wall profile).

We can now easily estimate the possible baryon to entropy ratio resulting from the spontaneously R-parity violating model. In [6] it was estimated that with the same choice of parameters used in this paper, \( \frac{n_B}{s} \sim 10^{-11} \). Since \( I_1/I_2 \) can easily be \( \sim 10^2 \) the baryon asymmetry created within the context of the scenario considered here can be as high as \( \sim 10^{-9} \). So the required value, \( \frac{n_B}{s} \sim 4 \times 10^{-11} \), is accessible with a non-restrictive choice of parameters.

In the present paper we have estimated the contribution to CP-violation in the bubble wall at the electroweak phase transition arising from the terms present in minimally supersymmetric spontaneously R-parity violating models. It has been shown that the amount of CP-violation due to the higgsino-lepton -interaction can quite easily be larger than the contribution due to the higgsino-gaugino -interaction by a factor of \( 10^2 \). This contribution is due to the possible existence of a large (complex) sneutrino vev and a Yukawa coupling constant of order \( \mathcal{O}(10^{-1}) \). It is notable that this result is not dependent on the soft supersymmetry breaking parameters \( M_{1,2} \) and furthermore allows freedom in the choice of the value of \( \mu \)-parameter which may also be chosen real. Also the phase transition temperature dependence is quite weak so it can be concluded that this mechanism allows for the creation of baryon asymmetry at the electroweak phase transition with quite a flexible choice of parameters. However, it should also be noted that there are several sources of uncertainty present due to the approximate nature of the chosen bubble wall profile and how higher corrections to the CP-violating source behave. Furthermore, the weak sphaleron rate is under discussion [12] and changes on that may significantly change the results on electroweak baryogenesis.
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