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# Supersymmetric CP-violating Currents and Electroweak Baryogenesis

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**Abstract**

In this work we compute the CP-violating currents of the right-handed stops and Higgsinos, induced by the presence of non-trivial vacuum expectation values of the Higgs fields within the context of the minimal supersymmetric extension of the Standard Model (MSSM) with explicit CP-violating phases. Using the Keldysh formalism, we perform the computation of the currents at finite temperature, in an expansion of derivatives of the Higgs fields. Contrary to previous works, we implement a resummation of the Higgs mass insertion effects to all orders in perturbation theory. While the components of the right-handed stop current $j^{\mu}_{\tilde{t}_R}$ become proportional to the difference $H_2\partial^\mu H_1 - H_1\partial^\mu H_2$ (suppressed by $\Delta\beta$), the Higgsino currents, $j^{\mu}_{\widetilde{H}_i}$, present contributions proportional to both $H_2\partial^\mu H_1 \pm H_1\partial^\mu H_2$. For large values of the charged Higgs mass and moderate values of $\tan\beta$ the contribution to the source proportional to $H_2\partial^\mu H_1 + H_1\partial^\mu H_2$ in the diffusion equations become sizeable, although it is suppressed by the Higgsino number violating interaction rate $\Gamma_\mu^{-1/2}$. For small values of the wall velocity, $0.04 \lesssim v_\omega \lesssim 0.1$, the total contribution leads to acceptable values of the baryon asymmetry for values of the CP-violating phases $\varphi_{CP}$ in the range $0.04 \lesssim |\sin\varphi_{CP}| \lesssim 1$. Finally, we comment on the relevance of the latest results of Higgs searches at LEP2 for the mechanism of electroweak baryogenesis within the MSSM.

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# 1 Introduction

The origin of the baryon asymmetry of the Universe is one of the most important open questions in cosmology and particle physics. It has been long understood that, in order to generate the observed baryon asymmetry, three requirements [1] need to be fulfilled: the non-conservation of baryon number, CP-violation and the existence of non-equilibrium processes [2]. Interestingly enough, at temperatures above the electroweak phase transition temperature, $T_c$, the Standard Model fulfills these requirements. Baryon number violation is induced by anomalous [3] sphaleron processes [4], which are suppressed at zero temperature, but whose rate grows linearly with the temperature above $T_c$ [5]. The non-conservation of CP is an essential property of the Standard Model, and non-equilibrium processes may be obtained through the expansion of bubbles of true vacuum, which occurs after the electroweak phase transition.

In spite of fulfilling all the desired properties, the rate of the CP-violating processes in the Standard Model (SM) is too small to induce the required baryon asymmetry [6, 7]. Moreover, the preservation of the generated baryon asymmetry after the electroweak phase transition requires a strongly first order phase transition, with $v(T_c)/T_c \gtrsim 1$, where $v(T_c)$ is the Higgs vacuum expectation value at the critical temperature $T_c$. For the experimentally allowed values of the Higgs mass, this requirement is not fulfilled in the Standard Model [8].

Supersymmetric particles lead to new radiative corrections to the Higgs effective potential at finite temperature [9]-[11]. Light boson fields with relevant couplings to the Higgs field may induce a stronger first order electroweak phase transition [12]-[26]. The supersymmetric partners of the top quark are the only new bosons which couple in a relevant way to the Higgs boson which acquire vacuum expectation value and hence play a relevant role in defining the strength of the phase transition [1]. For sufficiently small values of the stop masses the strength of the phase transition is enhanced [12, 20]. In order to get values of $v(T_c)/T_c \geq 1$, however, the right handed stop soft supersymmetry breaking squared mass parameter, $m_U^2$, should be small or even slightly negative and the stop mixing mass parameter $X_t = |A_t - \mu_c/\tan\beta|$ must be smaller than $\sim 0.6\ m_Q$, with $m_Q$ the left-handed stop supersymmetry breaking mass parameter. Under these conditions, and for $m_Q \lesssim$1–3 TeV, a strongly first order phase transition may be obtained up to values of the lightest CP-even Higgs boson mass as high as $\sim$ 110–115 GeV [25, 20].

Moreover, supersymmetric particles lead to new, relevant CP-violating sources for the generation of the baryon asymmetry [27]. Several computations have been performed [28]-[40] in recent years, showing that if the CP-violating phases associated with the chargino mass parameters are not too small, these sources may lead to acceptable values of the baryon asymmetry. In this work, we shall perform a computation of these new CP-violating sources in an expansion in derivatives of the Higgs background fields. Similarly to Ref. [29], we shall use the Keldysh formalism [41] for the computation of the CP-violating currents at finite temperature. We improve the computation of Ref. [29] in

[1] Although bottom and tau Yukawa couplings become large for large values of $\tan\beta$, the bottom and tau superfield couplings to the Higgs boson combination which acquires vacuum expectation value, $\Phi = H_1^0 \cos\beta + H_2^0 \sin\beta$, remains small, apart from an enhancement of the $\Phi$-trilinear coupling to left and right sbottoms and staus, which increases the corresponding mixings, but does not lead to an enhancement of the phase transition strength.

two main aspects. On the one hand, instead of computing the temporal component of the current in the lowest order of Higgs background insertions, we compute all current components by performing a resummation of the Higgs mass insertion contributions to all order in perturbation theory. The resummation is essential since it leads to a proper regularization of the resonant contribution to the temporal component of the current found in Ref. [29] and leads to contributions which are not suppressed for large values of the charged Higgs mass. On the other hand, we consider, in the diffusion equations, the contribution of Higgsino number violating interaction rate [40] from the Higgsino $\mu$ term in the lagrangian, $\Gamma_\mu$, that was considered in our previous calculations in the limit $\Gamma_\mu/T \to \infty$.

This article is organized as follows. In sections 2 and 3 we present the detailed derivation of the CP-violating currents for the cases of right-handed top squarks ($j^\mu_{\tilde{t}_R}$) and charginos ($j^\mu_{\widetilde{H}}$), respectively, by making use of the Keldysh formalism and resumming to all order in Higgs background insertions. These two sections deal with all the technical details of the computation, with the main results given in Eqs. (2.17), (3.16) and (3.18). In section 4 we present explicit, analytical, solutions to the diffusion equations and an explicit expression for the baryon asymmetry in the broken phase after the phase transition in the MSSM. In section 5 we exhibit the results of a numerical analysis of our solutions. A discussion of present Higgs mass constraints is made in section 6, and in section 7 we present our conclusions and outlook.

# 2 The squark sector

Our aim in this section is to compute the Green functions for left-handed ($\widetilde{t}_L(x)$) and right-handed ($\widetilde{t}_R(x)$) stop fields, describing the propagation of these scalars in the presence of a bubble wall. The bubble wall is assumed to be located at the space-time point $z$, where there is a non-trivial background of the MSSM Higgs fields, $H_i(z)$, which carries dimensionful CP-violating couplings to the left- and right-handed stops. We shall use these Green functions to compute the right-handed and left-handed stop currents at the point $z$. The starting point is the lagrangian for the stop system:

$$\mathcal{L}(x) = \left|\partial_\mu \widetilde{t}_L(x)\right|^2 + \left|\partial_\mu \widetilde{t}_R(x)\right|^2 + \left( \begin{array}{cc} \widetilde{t}^*_L(x) & \widetilde{t}^*_R(x) \end{array} \right) \mathcal{M}(x) \left( \begin{array}{c} \widetilde{t}_L(x) \\ \widetilde{t}_R(x) \end{array} \right) , \qquad (2.1)$$

where $\mathcal{M}$ is the stop squared mass matrix which depends, through the Higgs background, on the space-time point.

Clearly this is not a free lagrangian, since the mass matrix depends on the space-time coordinates, and we must identify the free and perturbative parts out of it. In order to make such a selection we will expand the mass matrix around the point $z^\mu \equiv (\vec{r}, t)$ (the point where we are calculating the currents in the plasma frame) up to first order in derivatives as,

$$\mathcal{M}(x) = \mathcal{M}(z) + (x-z)^\mu \mathcal{M}_\mu(z) , \qquad (2.2)$$

where we use the notation $\mathcal{M}_\mu(z) \equiv \partial \mathcal{M}(z)/\partial z^\mu$, and we can split the initial Lagrangian as:

$$\begin{aligned}
\mathcal{L}_0(x) &= \left|\partial_\mu \widetilde{t}_L(x)\right|^2 + \left|\partial_\mu \widetilde{t}_R(x)\right|^2 + \left( \begin{array}{cc} \widetilde{t}_L^*(x) & \widetilde{t}_R^*(x) \end{array} \right) \mathcal{M}(z) \left( \begin{array}{c} \widetilde{t}_L(x) \\ \widetilde{t}_R(x) \end{array} \right) \\
\mathcal{L}_{\rm int} &= (x-z)^\mu \left( \begin{array}{cc} \widetilde{t}_L^*(x) & \widetilde{t}_R^*(x) \end{array} \right) \mathcal{M}_\mu(z) \left( \begin{array}{c} \widetilde{t}_L(x) \\ \widetilde{t}_R(x) \end{array} \right) . \qquad (2.3)
\end{aligned}$$

Let $\mathcal{U}(z) \in SU(2)$ be the matrix that diagonalizes $\mathcal{M}(z)$. We can then rewrite $\mathcal{L}_0$ and $\mathcal{L}_{\rm int}$ as:

$$\begin{aligned}
\mathcal{L}_0 &= \sum_{i=1}^{2} \left\{ |\partial_\mu \chi_i(x)|^2 + m_i^2(z)\, |\chi_i(x)|^2 \right\} , \\
\mathcal{L}_{\rm int} &= (x-z)^\mu \left( \begin{array}{cc} \widetilde{t}_L^*(x) & \widetilde{t}_R^*(x) \end{array} \right) \mathcal{U}(z) \mathcal{M}_\mu(z)\, \mathcal{U}^\dagger(z) \left( \begin{array}{c} \widetilde{t}_L(x) \\ \widetilde{t}_R(x) \end{array} \right) , \qquad (2.4)
\end{aligned}$$

where $m_i^2(z)$, $\chi_i(z)$ $(i = 1, 2)$ are the eigenvalues and eigenvectors of $M(z)$. Note that the description in terms of the mass eigenstates $\chi_i(z)$ is useful so far the Higgs field variations are small for propagation lengths of the order of the inverse of the width of the stop fields, $\Gamma^{-1}$. Under these conditions, namely $L_w \Gamma / v_w \lesssim 1$, with $L_w$ and $v_w$ being the bubble wall width and velocity, respectively, an expansion in derivatives is justified [29].

Now we can write down the two point Green function for the field $(\chi_1(x)\ \chi_2(x))^T$ in matrix form:

$$G^\chi(x, y; z) = G(x, y; z) + \int d^4w\ (w-z)^\mu G(x, w; z)\, \mathcal{U}(z)\, \mathcal{M}_\mu(z) \mathcal{U}^\dagger(z)\ G(w, y; z) + \dots \qquad (2.5)$$

where $x$ and $y$ are assumed to be close to $z$, the point at which the current is being evaluated and around which the expansion is being performed ($|x-z|, |y-z| \ll \Gamma^{-1}$), $G(x, y; z)$ is the two by two diagonal free Green function of the stop mass eigenstates with masses $m_i(z)$, the trace over internal $(a = 1, 2)$ indices being understood in Eq. (2.5). Explicitly, the free Green functions for each of the two stop eigenstates can be written as [41]:

$$\begin{aligned}
G_i^{11} &= P_i^+ + f_B \left( P_i^+ - P_i^- \right) \\
G_i^{12} &= \left[ \theta(p^0) + f_B \right] \left( P_i^+ - P_i^- \right) \\
G_i^{21} &= \left[ \theta(-p^0) + f_B \right] \left( P_i^+ - P_i^- \right) \\
G_i^{22} &= -P_i^- + f_B \left( P_i^+ - P_i^- \right) , \qquad (2.6)
\end{aligned}$$

where $f_B \equiv n_B(|p^0|)$ is the Bose-Einstein distribution function, which contains the dependence on the temperature $T$,

$$P_i^\pm = \frac{1}{p_0^2 - \vec{p}^2 - m_i^2(z) \pm 2i\Gamma_{\widetilde{t}} |p^0|} , \qquad (2.7)$$

and $\Gamma_{\widetilde{t}}$ is the stop width which can be taken to be $\Gamma_{\widetilde{t}} \sim \alpha_s T$ independently of the stop mass eigenstate.

Since we need to calculate the CP-violating currents induced by the right-handed stop states, we should first go to the weak eigenstate basis. The Green functions in the weak eigenstate basis can be obtained from the ones given above, which were computed in the basis of mass eigenstates, by the following expression

$$G^{\widetilde{t}}(x,y;z) = \mathcal{U}^\dagger(z) G^\chi(x,y;z) \mathcal{U}(z) \ .$$

Therefore, the current for right-handed stops takes the form:

$$j^\mu_{\tilde{t}_R}(z) = \lim_{x,y\to z} \mathrm{Tr}\left[P_2 \frac{\partial G^{\widetilde{t}}(x,y;z)}{\partial(x-y)_\mu}\right] \ , \tag{2.8}$$

where $P_2 = (\sigma_0 - \sigma_3)/2$, $\sigma_i$ being the two by two Pauli matrices and $\sigma_0$ the two by two identity matrix, is a projection matrix which allows to separate the current induced by the right-handed stops from the one induced by the left-handed stops. Nevertheless, since baryon number is conserved at this point, the total CP-violating currents induced by left- and right-handed top squarks must be zero, $\mathrm{Tr}[\partial^\mu G^{\widetilde{t}}(x,y)] = 0$, and therefore

$$j^\mu_{\tilde{t}_R}(z) = -\frac{1}{2} \lim_{x,y\to z} \mathrm{Tr}\left[\sigma_3 \frac{\partial G^{\widetilde{t}}(x,y;z)}{\partial(x-y)_\mu}\right] \ . \tag{2.9}$$

After integrating over the $w$ space-time variable, and going to momentum space, we can write the current in terms of free Green functions of the mass eigenstates at the point $z$:

$$\begin{aligned} j^\mu_{\tilde{t}_R}(z) \; &= \; \frac{1}{2}\int \frac{d^4p}{(2\pi)^4} p^\mu \mathrm{Tr}\left[\sigma_3 \mathcal{U}^\dagger(z)\, G^\nu(p;z)\, \mathcal{U}(z)\, \mathcal{M}_\nu(z)\, \mathcal{U}^\dagger(z)\, G(p;z) \mathcal{U}(z)\right. \\ &- \; \left. \sigma_3 \mathcal{U}^\dagger(z)\, G(p;z)\, \mathcal{U}(z)\, \mathcal{M}_\nu(z)\, \mathcal{U}^\dagger(z)\, G^\nu(p;z)\, \mathcal{U}(z)\right] \end{aligned} \tag{2.10}$$

since the contribution induced by the linear term in $z$ in Eq. (2.5) trivially vanishes because $G(p;z)$ only depends on $|\mathbf{p}|$ and $p^0$. We are using the notation $G^\nu(p;z) = \partial G(p;z)/\partial p_\nu$. Note that in the above expression only off-diagonal terms of the derivatives of the mass matrix $\mathcal{M}_\nu(z)$ at the point $z$ give a non-vanishing contribution. We shall denote by $\widetilde{\mathcal{M}}_\nu(z)$, the matrix containing only the derivative of the off-diagonal terms of the matrix $\mathcal{M}(z)$.

The current could be simplified a little bit more by using:

$$\mathcal{U}^\dagger(z) D\, \mathcal{U}(z) = \sigma_1 D \sigma_1 + \frac{1}{2}\mathrm{Tr}[\mathcal{U}(z)]\mathrm{Tr}[D\sigma_3]\mathcal{U}^\dagger(z)\sigma_3 \tag{2.11}$$

where $D$ is a diagonal matrix. Then $j^\mu_{\tilde{t}_R}(z)$ can be written as:

$$\begin{aligned} j^\mu_{\tilde{t}_R}(z) =& -\frac{i}{4}\mathrm{Tr}[\mathcal{U}(z)]\mathrm{Tr}\left[\widetilde{\mathcal{M}}_\nu(z)\mathcal{U}(z)\right]\int \frac{d^4p}{(2\pi)^4} p^\mu \mathrm{Tr}\left[\sigma_1 G(p;z)\sigma_2 G^\nu(p;z)\right] \\ =& \frac{1}{4}\mathrm{Tr}[\mathcal{U}(z)]\mathrm{Tr}\left[\widetilde{\mathcal{M}}_\nu(z)\mathcal{U}(z)\right]\int \frac{d^4p}{(2\pi)^4} p^\mu \epsilon^{ij} G_i(p;z)\ G^\nu_j(p;z) \ . \end{aligned} \tag{2.12}$$

Expanding $G_i$ in terms of $P_i^{\pm}$ one gets:

$$
\begin{aligned}
j^i_{\tilde{t}_R}(z) =& \frac{8\,C^i}{3\pi}\mathrm{Im}\left\{\int_{-\infty}^{\infty}\frac{dp^0}{2\pi}(1+2f_B)\int_0^{\infty}\frac{d\mathbf{p}}{2\pi}\;\mathbf{p}^4(P_1^+(p;z)P_2^+(p;z))^2\right\}\\
j^0_{\tilde{t}_R}(z) =& -\frac{2\,C^0}{\pi}\mathrm{Im}\left\{\int_{-\infty}^{\infty}\frac{dp^0}{2\pi}|p^0|\left[(1+2f_B)\left(|p^0|+i\Gamma_{\tilde{t}}\right)\int_0^{\infty}\frac{d\mathbf{p}}{2\pi}\left(P_1^+(p;z)P_2^+(p;z)\right)^2\right.\right.\\
&\left.\left.-\int_0^{\infty}\frac{d\mathbf{p}}{2\pi}\frac{f_B'}{m_1^2(z)-m_2^2(z)}\,P_1^+(p;z)P_2^-(p;z)\right]\right\}
\end{aligned}
\tag{2.13}
$$

where $f_B'$ is the derivative of $f_B$ with respect to its argument and $C_\mu$ is given by

$$
C_\mu = (m_1^2(z)-m_2^2(z))\,\mathrm{Tr}[\mathcal{U}(z)]\mathrm{Tr}\left[\widetilde{\mathcal{M}}_\mu(z)\mathcal{U}(z)\right]\;. \tag{2.14}
$$

Using now the particular value of the squared mass matrix $\mathcal{M}$ for the stop system,

$$
\mathcal{M}(z) = \begin{pmatrix} m_Q^2 + h_t^2\,H_2^2(z) & h_t\,(A_t H_2(z) - \mu_c^* H_1(z)) \\ h_t\,(A_t^* H_2(z) - \mu_c H_1(z)) & m_U^2 + h_t^2\,H_2^2(z) \end{pmatrix}\;, \tag{2.15}
$$

where $h_t$ is the top-quark Yukawa coupling, $A_t$ the left-right stop mixing parameter, and $\mu_c$ the complex Higgsino mass parameter, defined as $\mu_c \equiv \mu\exp(i\varphi_\mu)$, with $\mu$ real (positive or negative). In the above, we have neglected corrections $\mathcal{O}(g^2)$. In this approximation, the above constant vector $C_\mu$, Eq. (2.14), can be written as:

$$
C^\mu = 2\,h_t^2\;\mathrm{Im}(A_t\,\mu_c)\;\{H_2(z)H_1^\mu(z) - H_1(z)H_2^\mu(z)\}\;. \tag{2.16}
$$

Hence, in order to compute the CP-violating currents induced by the stop fields, the momentum integrals should be performed. Due to the form of the free Green functions, Eq. (2.6), the integral over the temporal component of the momentum cannot be performed by standard integration methods in the complex plane. It is therefore better to perform the integration over the spatial components of the momentum and express the results as an integral function over $p^0$, which admits a simple physical interpretation. In order to perform the integrals of the spatial components of the momentum, one should note that all functions depend only on $|\mathbf{p}|^2$. Therefore, the angular integration can be trivially performed and the integral over the modulus $|p|$ of a function $\mathcal{F}(|p|)$ can be written as half the integral on the whole real plane of the function $\mathcal{F}(x)$, with $\mathcal{F}(x)=\mathcal{F}(-x)$. Doing this, we can perform the spatial momentum integrals in Eq. (2.13) by means of standard techniques of integration in the complex plane and the residues theorem, and we can cast the resulting currents as:

$$
j^\mu_{\tilde{t}_R}(z) = h_t^2\;\mathrm{Im}(A_t\,\mu_c)\;\{H_2(z)H_1^\mu(z) - H_1(z)H_2^\mu(z)\}\left\{\mathcal{F}_B(z) + \delta^{\mu\,0}\mathcal{G}_B(z)\right\} \tag{2.17}
$$

where

$$
\begin{aligned}
\mathcal{F}_B(z) =& \frac{1}{6\pi^2}\mathrm{Re}\int_0^{\infty} dp^0\,(1+2\,f_B)\left(\frac{1}{z_1+z_2}\right)^3\\
\mathcal{G}_B(z) =& \frac{1}{3\pi^2}\mathrm{Re}\int_0^{\infty} dp^0\,p^0 f_B'\left\{\left(\frac{1}{z_1+z_2}\right)^3\right.\\
&\left.-\frac{3}{m_1^2(z)-m_2^2(z)}\left[\frac{z_1}{m_1^2-m_2^2-4i\Gamma_{\tilde{t}}\,p^0}+\frac{z_2}{m_1^2-m_2^2+4i\Gamma_{\tilde{t}}\,p^0}\right]\right\}
\end{aligned}
\tag{2.18}
$$

and $z_i$ is defined as the pole of $P_i^+$, i.e.

$$z_i(p^0) = \sqrt{p^0\,(p^0 + 2\,i\Gamma_{\widetilde{t}}) - m_i^2(z)} \tag{2.19}$$

with positive real and imaginary parts satisfying $\mathrm{Re}(z_i) = \Gamma_{\widetilde{t}}\, p^0/\mathrm{Im}(z_i)$.

# 3 The chargino sector

For the case of the charged gaugino-Higgsino system we will follow similar steps as the ones we performed before to compute the stop current. In this case the role of right-handed stops is played by the (left- and right-handed) Higgsinos. The starting point is the lagrangian:

$$\mathcal{L}(x) = \overline{\widetilde{h}}_c(x)\partial_\mu\gamma^\mu\widetilde{h}_c(x) + \overline{\widetilde{W}}_c(x)\partial_\mu\gamma^\mu\widetilde{W}_c(x) + \begin{pmatrix}\overline{\widetilde{W}}_c(x) & \overline{\widetilde{h}}_c(x)\end{pmatrix} M(x) \begin{pmatrix}\widetilde{W}_c(x) \\ \widetilde{h}_c(x)\end{pmatrix} \tag{3.1}$$

where

$$\widetilde{h}_c = \begin{pmatrix}\widetilde{h}_2^+ \\ \widetilde{h}_1^{-\,*}\end{pmatrix}, \qquad \widetilde{W}_c = \begin{pmatrix}\widetilde{W}^+ \\ \widetilde{W}^{-\,*}\end{pmatrix} .$$

From the structure of the chargino mass matrix we can write the lagrangian in the following form:

$$\begin{aligned}\mathcal{L}(x) =& \psi_R(x)^\dagger\sigma_\mu\partial^\mu\psi_R(x) + \psi_L(x)^\dagger\overline{\sigma}_\mu\partial^\mu\psi_L(x) \\ &+ \psi_R(x)^\dagger M(x)\psi_L(x) + \psi_L(x)^\dagger M^\dagger(x)\psi_R(x)\end{aligned} \tag{3.2}$$

where in this expression we have used

$$\psi_R(x) = \begin{pmatrix}\widetilde{W}^+ \\ \widetilde{h}_2^+\end{pmatrix}, \qquad \psi_L(x) = \begin{pmatrix}\widetilde{W}^- \\ \widetilde{h}_1^-\end{pmatrix} .$$

Expanding the masses around the point $z$ and splitting the lagrangian into a free and a perturbative part, we can write, to first order in derivatives:

$$\begin{aligned}\mathcal{L}_0(x) =& \psi_R(x)^\dagger\sigma_\mu\partial^\mu\psi_R(x) + \psi_L(x)^\dagger\overline{\sigma}_\mu\partial^\mu\psi_L(x) \\ &+ \psi_R(x)^\dagger M(z)\psi_L(x) + \psi_L(x)^\dagger M^\dagger(z)\psi_R(x) \\ \\ \mathcal{L}_{int}(x) =& (x-z)^\mu \left\{\psi_R(x)^\dagger M_\mu(z)\psi_L(x) + \psi_L(x)^\dagger M_\mu^\dagger(z)\psi_R(x)\right\} .\end{aligned} \tag{3.3}$$

Like for the scalar case we will diagonalize $M(z)$ by means of the matrices $\mathcal{U}(z)$, $\mathcal{V}(z) \in SU(2)$. Additional phase redefinition can be performed in order to bring the mass eigenstates to be real and positive. In general, the lagrangian can be written as:

$$\mathcal{L}_0(x) = \varphi_R(x)^\dagger \sigma_\mu \partial^\mu \varphi_R(x) + \varphi_L(x)^\dagger \overline{\sigma}_\mu \partial^\mu \varphi_L(x) + \varphi_R(x)^\dagger \begin{pmatrix} m_1(z) & 0 \\ 0 & m_2(z) \end{pmatrix} \varphi_L(x) + \varphi_L(x)^\dagger \begin{pmatrix} m_1^*(z) & 0 \\ 0 & m_2^*(z) \end{pmatrix} \varphi_R(x)$$

$$\mathcal{L}_{int}(x) = (x-z)^\mu \left\{ \varphi_R(x)^\dagger \mathcal{U}(z) M_\mu(z) \mathcal{V}^\dagger(z) \varphi_L(x) + \varphi_L(x)^\dagger \mathcal{V}(z) M_\mu^\dagger(z) \mathcal{U}^\dagger(z) \varphi_R(x) \right\} \quad (3.4)$$

where $m_i(z)$ are the eigenvalues of $M(z)$ and

$$\varphi_R(x) = \mathcal{U}(z)\psi_R(x), \qquad \varphi_L(x) = \mathcal{V}(z)\psi_L(x)$$

are the mass eigenstates at the point $z$.

At this point we can write the Green functions describing the propagation of the right- and left-handed fermion $\varphi$ fields, $S_\varphi^{RR}$ and $S_\varphi^{LL}$, respectively, as

$$\begin{aligned} S_\varphi^{RR}(x,y;z) = & S^{RR}(x,y;z) \\ & + \int d^4w (w-z)^\mu S^{RR}(x,w;z) \mathcal{U}(z)\, M_\mu(z) \mathcal{V}^\dagger(z)\, S^{LR}(w,y;z) \\ & + \int d^4w (w-z)^\mu S^{RL}(x,w;z) \mathcal{V}(z)\, M_\mu^\dagger(z) \mathcal{U}^\dagger(z) S^{RR}(w,y;z) \\ S_\varphi^{LL}(x,y;z) = & S^{LL}(x,y;z) \\ & + \int d^4w (w-z)^\mu S^{LL}(x,w;z) \mathcal{V}(z)\, M_\mu^\dagger(z) \mathcal{U}^\dagger(z) S^{RL}(w,y;z) \\ & + \int d^4w (w-z)^\mu S^{LR}(x,w;z) \mathcal{U}(z)\, M_\mu(z) \mathcal{V}^\dagger(z)\, S^{LL}(w,y;z) \end{aligned} \quad (3.5)$$

where $S^{LL}$, $S^{RR}$, $S^{LR}$ and $S^{RL}$ denote the left-left, right-right, left-right and right-left Green functions of free fermions with mass $m_i(z)$. In the approximation where both fermionic widths are equal, we can rewrite the free fermionic Green functions in terms of bosonic ones as:

$$\begin{aligned} S^{RR}(p;z) &= \sigma_\mu p^\mu G(p;z) & S^{RL}(p;z) &= \begin{pmatrix} m_1(z) & 0 \\ 0 & m_2(z) \end{pmatrix} G(p;z) \\ S^{LR}(p;z) &= \begin{pmatrix} m_1^*(z) & 0 \\ 0 & m_2^*(z) \end{pmatrix} G(p;z) & S^{LL}(p;z) &= \overline{\sigma}_\mu p^\mu G(p;z) \end{aligned} \quad (3.6)$$

where the free Green functions $G(p;z)$ are given by (2.6) with $f_B \to f_F \equiv -n_F(|p^0|)$, $n_F$ being the Fermi-Dirac distribution function, $m_i(z) \to |m_i(z)|$, and $\Gamma_{\widetilde{t}} \to \Gamma_{\widetilde{H}} \sim \alpha_W T$. Using the relations between Green functions in the mass and weak eigenstate basis, as we did in the stop case, we obtain in the weak eigenstates basis,

$$\begin{aligned} S_\psi^{RR}(p;z) &= \mathcal{U}^\dagger(z) S_\varphi^{RR}(p;z)\, \mathcal{U}(z) & S_\psi^{RL}(p;z) &= \mathcal{U}^\dagger(z) S_\varphi^{RL}(p;z)\, \mathcal{V}(z) \\ S_\psi^{LR}(p;z) &= \mathcal{V}^\dagger(z) S_\varphi^{LR}(p;z)\, \mathcal{U}(z) & S_\psi^{LL}(p;z) &= \mathcal{V}^\dagger(z) S_\varphi^{LL}(p;z)\, \mathcal{V}(z)\ . \end{aligned} \quad (3.7)$$

The Higgsino currents can now be defined as:

$$j^{\mu}_{\widetilde{\mathcal{H}}_{\pm}}(z) = \lim_{x,y\to z} \left\{ \mathrm{Tr}\left[P_2 \sigma^{\mu} S^{RR}_{\psi}(x,y;z)\right] \pm \mathrm{Tr}\left[P_2 \overline{\sigma}^{\mu} S^{LL}_{\psi}(x,y;z)\right]\right\} \tag{3.8}$$

where $P_2$ is the projection operator used in Eq. (2.8).

By replacing (3.6) and (3.7) in these currents, and taking into account that, as happened in the stop case, the contribution of the linear term in $z$ in Eq. (3.5) is zero by symmetry reasons, one gets [2]

$$\begin{aligned} j^{\mu}_{\widetilde{\mathcal{H}}_{\pm}}(z) =& \frac{1}{2}\int \frac{d^4p}{(2\pi)^4} \left\{ p^{\mu}\, \mathrm{Tr}\left[\sigma_3 \left( \mathcal{U}^{\dagger}(z) G(p;z) \mathcal{U}(z) M_{\rho}(z) M^{\dagger}(z) \mathcal{U}^{\dagger}(z) G^{\rho}(p;z)\mathcal{U}(z)\right.\right.\right. \\ & - \mathcal{U}^{\dagger}(z) G^{\rho}(p;z) \mathcal{U}(z) M(z) M^{\dagger}_{\rho}(z) \mathcal{U}^{\dagger}(z) G(p;z) \mathcal{U}(z) \\ & \pm \mathcal{V}^{\dagger}(z) G(p;z) \mathcal{V}(z) M^{\dagger}_{\rho}(z) M(z) \mathcal{V}^{\dagger}(z) G^{\rho}(p;z) \mathcal{V}(z)) \\ & \left.\left.\mp\, \mathcal{V}^{\dagger}(z) G^{\rho}(p;z) \mathcal{V}(z) M^{\dagger}(z) M_{\rho}(z) \mathcal{V}^{\dagger}(z) G(p;z) \mathcal{V}(z)\right)\right] \\ & + \mathrm{Tr}\left[\mathcal{U}^{\dagger}(z) G(p;z) \mathcal{U}(z) (M^{\mu}(z) M^{\dagger}(z) - M(z) M^{\mu\,\dagger}(z)) \mathcal{U}^{\dagger}(z) G(p;z) \mathcal{U}(z)\right. \\ & \left.\left. \pm\, \mathcal{V}^{\dagger}(z) G(p;z) \mathcal{V}(z) (M^{\mu\,\dagger}(z) M(z) - M^{\dagger}(z) M^{\mu}(z)) \mathcal{V}^{\dagger}(z) G(p;z) \mathcal{V}(z)\right]\right\} \; . \end{aligned} \tag{3.9}$$

The chargino mass matrix is given by

$$M(z) = \begin{pmatrix} M_2 & u_2(z) \\ u_1(z) & \mu_c \end{pmatrix} \tag{3.10}$$

where we have defined $u_i(z) \equiv g H_i(z)$. The diagonalizing matrices are [30]

$$\begin{aligned} \mathcal{U} =& \frac{1}{\sqrt{2\,\Lambda(\Delta+\Lambda)}} \begin{pmatrix} \Delta+\Lambda & M_2\, u_1 + \mu_c^*\, u_2 \\ -\,(M_2\, u_1 + \mu_c\, u_2) & \Delta+\Lambda \end{pmatrix} \\ \mathcal{V} =& \frac{1}{\sqrt{2\,\Lambda(\bar\Delta+\Lambda)}} \begin{pmatrix} \bar\Delta+\Lambda & M_2\, u_2 + \mu_c\, u_1 \\ -\,(M_2\, u_2 + \mu_c^*\, u_1) & \bar\Delta+\Lambda \end{pmatrix} \; , \end{aligned} \tag{3.11}$$

where field redefinitions have been made in order to make the Higgs vacuum expectation values, as well as the weak gaugino mass $M_2$, real,

$$\begin{aligned} \Delta =& (M_2^2 - |\mu_c|^2 - u_1^2 + u_2^2)/2 \\ \bar\Delta =& (M_2^2 - |\mu_c|^2 - u_2^2 + u_1^2)/2 \\ \Lambda =& \left(\Delta^2 + |M_2\, u_1 + \mu_c^*\, u_2|^2\right)^{1/2} \; , \end{aligned} \tag{3.12}$$

and the mass eigenvalues are given by

$$\begin{aligned} m_1(z) =& \frac{(\Delta + \Lambda + u_1^2(z))\, M_2 + u_1(z) u_2(z) \mu_c^*}{\sqrt{(\Delta+\Lambda)(\bar\Delta+\Lambda)}} \\ m_2(z) =& \frac{(\Delta + \Lambda - u_2^2(z))\, \mu_c - u_1(z) u_2(z) M_2}{\sqrt{(\Delta+\Lambda)(\bar\Delta+\Lambda)}} \; . \end{aligned} \tag{3.13}$$

[2]Notice that the phases $\varphi_i$ of the mass eigenvalues, $m_i(z) = |m_i(z)| \exp\{i\varphi_i(z)\}$ can be absorbed in a redefinition of the matrix $\mathcal{V}(z)$, as $\mathcal{V}(z) \to \mathrm{diag}(\exp\{i\varphi_1(z)\}, \exp\{i\varphi_2(z)\})\mathcal{V}(z)$. As required, the currents (3.9) do not depend on this phase redefinition.

Using these expressions, and the property $2\,\Lambda = |m_1(z)|^2 - |m_2(z)|^2$, we can cast the Higgsino currents in the following general form:

$$j^\mu_{\widetilde{\mathcal{H}}_+} = \frac{\mathrm{Im}(M_2\,\mu_c)}{\Lambda} \left\{ [u_2(z)u_1^\nu(z) - u_1(z)u_2^\nu(z)] \int \frac{d^4p}{(2\pi)^4}\, p^\mu\, G_i(p;z)\epsilon^{ij} G_j^\nu(p;z) \right. \tag{3.14}$$
$$\left. + \frac{u_2^2(z)-u_1^2(z)}{2\,\Lambda}\,[u_2(z)u_1^\nu(z) + u_1(z)u_2^\nu(z)] \int \frac{d^4p}{(2\pi)^4}\, p^\mu\, G_i(p;z)\left(\delta^{ij} - \sigma_1^{ij}\right) G_j^\nu(p;z)\right\}$$

where $\epsilon^{12} = +1$, and

$$j^\mu_{\widetilde{\mathcal{H}}_-} = \frac{\mathrm{Im}(M_2\,\mu_c)}{2\,\Lambda}\,[u_2(z)u_1^\nu(z) + u_1(z)u_2^\nu(z)] \int \frac{d^4p}{(2\pi)^4}\, p^\mu$$
$$\times \left\{ [(G_2+G_1)(G_2-G_1)]^\nu + \left(\frac{\Delta+\bar{\Delta}}{\Lambda}\right)(G_2-G_1)\,(G_2^\nu - G_1^\nu)\right\}\ . \tag{3.15}$$

Notice that while the first term in $j^\mu_{\widetilde{\mathcal{H}}_+}$ is similar to the squark current $j^\mu_{\tilde{t}_R}$ (it is proportional to $u_2(z)u_1^\nu(z) - u_1(z)u_2^\nu(z)$), the second term in $j^\mu_{\widetilde{\mathcal{H}}_+}$ and the current $j^\mu_{\widetilde{\mathcal{H}}_-}$ have no counterpart in the scalar sector. The contribution proportional to $u_2(z)u_1^\nu(z) - u_1(z)u_2^\nu(z)$ is proportional to the variation $\Delta\beta$ of the angle $\beta$ along the bubble wall. Since $\Delta\beta \lesssim 10^{-2}$, the corresponding contribution is suppressed. The contribution proportional to $u_2(z)u_1^\nu(z) + u_1(z)u_2^\nu(z)$, instead, is not affected by this suppression factor, although it is suppressed, for large values of $\tan\beta$, as $1/\tan\beta$.

Now the integration over the spatial components of the momentum can be performed as in the previous section and the final currents can be cast as follows. For the current $j^\mu_{\widetilde{\mathcal{H}}_+}$ one obtains,

$$j^\mu_{\widetilde{\mathcal{H}}_+}(z) = 2\,\mathrm{Im}(M_2\,\mu_c)\left\{[u_2(z)u_1^\mu(z) - u_1(z)u_2^\mu(z)]\left\{\mathcal{F}_F(z) + \delta^{\mu\,0}\mathcal{G}_F(z)\right\}\right.$$
$$\left. + \left[u_2^2(z) - u_1^2(z)\right][u_2(z)u_1^\mu(z) + u_1(z)u_2^\mu(z)]\,\mathcal{H}_F(z)\right\} \tag{3.16}$$

where the functions $\mathcal{F}_F$, $\mathcal{G}_F$ are defined in (2.18) after changing $f_B \to f_F$, $m_i(z) \to |m_i(z)|$ and $\Gamma_{\tilde{t}} \to \Gamma_{\widetilde{H}}$, and

$$\mathcal{H}_F(z) = \frac{1}{8\pi^2}\mathrm{Re}\int_0^\infty dp^0\,(1+2\,f_F)\frac{1}{z_1\,z_2}\left(\frac{1}{z_1+z_2}\right)^3 \tag{3.17}$$

with $z_i(z)$ defined in (2.19), after changing $m_i(z) \to |m_i(z)|$ and $\Gamma_{\tilde{t}} \to \Gamma_{\widetilde{H}}$. Note that, being proportional to $u_2^2(z) - u_1^2(z) \equiv -u^2(z)\cos 2\beta(z)$, the second term of $j^\mu_{\widetilde{\mathcal{H}}_+}$ vanishes at the lowest order in the Higgs field insertions, in agreement with our previous results [29], and it also vanishes in the case $\tan\beta = 1$.

For the current $j^\mu_{\widetilde{\mathcal{H}}_-}$ one obtains,

$$j^\mu_{\widetilde{\mathcal{H}}_-}(z) = 2\,\mathrm{Im}(M_2\,\mu_c)\,[u_2(z)u_1^\mu(z) + u_1(z)u_2^\mu(z)]\left\{\mathcal{K}_F(z) + 2\,\left[\Delta+\bar{\Delta}\right]\,\mathcal{H}_F(z)\right\} \tag{3.18}$$

where the function $\mathcal{K}_F$ is defined as,

$$\mathcal{K}_F(z) = -\frac{1}{4\,\pi^2}\mathrm{Re}\int_0^\infty dp^0\,(1+2\,f_F)\frac{1}{z_1\,z_2}\left(\frac{1}{z_1+z_2}\right)\ . \tag{3.19}$$

Note that the current (3.18) appears to leading order in the Higgs mass insertion and it is not suppressed by $\Delta\beta$. However, $j^{\mu}_{\widetilde{\mathcal{H}}_-}$ is suppressed for large values of $\tan\beta$ (which are needed to push the Higgs mass beyond the most recent LEP bounds, as we will discuss in section 6) and, moreover, its effects on the corresponding Higgs density are damped by the presence of the Higgsino number violating interaction rate $\Gamma_\mu$. Accordingly, its contribution to the BAU is small.

A similar calculation for the neutral gaugino-Higgsino system would involve diagonalization of the four-by-four neutralino mass matrix, making the analytic resummation treatment much more involved than for the chargino case. The analysis performed in Ref. [29], to lowest order in the Higgs mass insertions, showed, as expected from a naive counting of degrees of freedom, that the neutralinos contribute to the Higgsino current as half the chargino contribution, with a total effect given by 3/2 that of the chargino. After resummation of the neutralino sector it would be reasonable to expect a total contribution to the Higgsino current equal to $\sim 3/2$ that of the chargino sector. However, and to be as conservative as possible in our calculation of the baryon asymmetry, we would just consider as source of baryon number the chargino current (remember that left-handed squarks are assumed very heavy and decouple from the thermal bath) that was computed in this section, keeping in mind that an enhancement factor $\sim 3/2$ might appear after a rigorous calculation of the currents in the neutralino sector.

# 4 The baryon asymmetry

To evaluate the baryon asymmetry generated in the broken phase we need to first compute the density of left-handed quarks and leptons, $n_L$, in front of the bubble wall (in the symmetric phase). These chiral densities are the ones that induce weak sphalerons to produce a net baryon number. Since, in the present scenario, there is essentially no lepton asymmetry, the density to be computed in the symmetric phase [3] is $n_L = n_Q + \sum_{i=1}^{2} n_{Q_i}$ where the density of a chiral supermultiplet $Q \equiv (q, \tilde{q})$ is understood as the sum of densities of particle components, assuming the supergauge interactions to be in thermal equilibrium, $n_Q = n_q + n_{\tilde{q}}$. If the system is near thermal equilibrium, particle densities, $n_i$, are related to the local chemical potential, $\mu_i$ by the relation $n_i = k_i \mu_i T^2/6$, where $k_i$ are statistical factors equal to 2 (1) for bosons (fermions) and exponentially suppressed for particle masses $m_i$ much larger than $T$. For the calculation of the density $n_L$ we will use the formalism described in Refs. [28, 29].

We will consider those particle species that participate in fast particle number changing transitions, neglecting all Yukawa couplings except those corresponding to the top quark. In this approximation, there is no left-handed lepton number contribution to $n_L$. By introducing strong sphaleron effects, first and second family quark number is generated. Assuming that all quarks have nearly the same diffusion constant it turns out that [28], $n_{Q_1} = n_{Q_2} = 2(n_Q + n_T)$, and then,

$$n_L = 5\, n_Q + \, 4\, n_T \ . \tag{4.1}$$

[3] We use, for the third family, the notation $Q \equiv Q_3,\ T \equiv T_3$.

In general we will relate particle number changing, or fermion number violating, rates $\Gamma_X$ with the corresponding rates per unit volume $\gamma_X$, as,

$$\Gamma_X = \frac{6\,\gamma_X}{T^3}\ . \tag{4.2}$$

The involved weak and strong sphaleron rates are:

$$\Gamma_{ws} = \ 6\,\kappa_{ws}\,\alpha_w^5 T, \quad \Gamma_{ss} = \ 6\,\kappa_{ss}\,\frac{8}{3}\,\alpha_s^4 T\ , \tag{4.3}$$

respectively, where $\kappa_{ws} = 20 \pm 2$ [42] and $\kappa_{ss} = \mathcal{O}(1)$. The particle number changing rates that will be considered both in the symmetric and in the broken phase are: $\Gamma_{Y_2}$, corresponding to all supersymmetric and soft breaking trilinear interactions arising from the $h_t H_2 Q T$ term in the superpotential, $\Gamma_{Y_1}$, which corresponds to the supersymmetric trilinear scalar interaction in the Lagrangian involving the third generation squarks and the Higgs $H_1$, and $\Gamma_\mu$, which corresponds to the $\mu_c \tilde{H}_1 \tilde{H}_2$ term in the Lagrangian. There are also the Higgs number violating and axial top number violation processes, induced by the Higgs self interactions and by top quark mass effects, with rates $\Gamma_h$ and $\Gamma_m$, respectively, that are only active in the broken phase.

We will write now a set of diffusion equations involving $n_Q$, $n_T$, $n_{H_1}$ (the density of $H_1 \equiv (h_1, \tilde{h}_1)$) and $n_{H_2}$ (the density of $\bar{H}_2 \equiv (\bar{h}_2, \bar{\tilde{h}}_2)$), and the particle number changing rates and CP-violating source terms discussed above. In the bubble wall frame, and ignoring the curvature of the bubble wall, all quantities become functions of $z \equiv r + v_\omega t$, where $v_\omega$ is the bubble wall velocity. The diffusion equations are:

$$\begin{aligned} v_\omega n_Q' =& D_q n_Q'' - \Gamma_Y \left[\frac{n_Q}{k_Q} - \frac{n_T}{k_T} - \frac{n_H + \rho\, n_h}{k_H}\right] - \Gamma_m \left[\frac{n_Q}{k_Q} - \frac{n_T}{k_T}\right] \\ & -6\Gamma_{ss}\left[2\,\frac{n_Q}{k_Q} - \frac{n_T}{k_T} + 9\,\frac{n_Q + n_T}{k_B}\right] + \tilde{\gamma}_Q \end{aligned} \tag{4.4}$$

$$\begin{aligned} v_\omega n_T' =& D_q n_T'' + \Gamma_Y \left[\frac{n_Q}{k_Q} - \frac{n_T}{k_T} - \frac{n_H + \rho\, n_h}{k_H}\right] + \Gamma_m \left[\frac{n_Q}{k_Q} - \frac{n_T}{k_T}\right] \\ & +3\Gamma_{ss}\left[2\,\frac{n_Q}{k_Q} - \frac{n_T}{k_T} + 9\,\frac{n_Q + n_T}{k_B}\right] - \tilde{\gamma}_Q \end{aligned} \tag{4.5}$$

$$v_\omega n_H' = D_h n_H'' + \Gamma_Y \left[\frac{n_Q}{k_Q} - \frac{n_T}{k_T} - \frac{n_H + \rho\, n_h}{k_H}\right] - \Gamma_h\,\frac{n_H}{k_H} + \tilde{\gamma}_{\tilde{H}_+} \tag{4.6}$$

$$v_\omega n_h' = D_h n_h'' + \rho\,\Gamma_Y \left[\frac{n_Q}{k_Q} - \frac{n_T}{k_T} - \frac{n_H + n_h/\rho}{k_H}\right] - (\Gamma_h + 4\,\Gamma_\mu)\,\frac{n_h}{k_H} + \tilde{\gamma}_{\tilde{H}_-} \tag{4.7}$$

where all derivatives are with respect to $z$, $D_q \sim 6/T$ and $D_h \sim 110/T$ are the corresponding diffusion constants in the quark and Higgs sectors [44], $n_H \equiv n_{H_2} + n_{H_1}$, $n_h \equiv n_{H_2} - n_{H_1}$, $k_H \equiv k_{H_1} + k_{H_2}$, $\Gamma_Y \equiv \Gamma_{Y_2} + \Gamma_{Y_1}$ and $\rho\,\Gamma_Y \equiv \Gamma_{Y_2} - \Gamma_{Y_1}$. The parameter $\rho$ is in the range $0 \leq \rho \leq 1$. In previous analyses [28, 29, 38] the limit $\Gamma_\mu \to \infty$ was implicitly considered, leading to the solution $n_h \to 0$. However, as we will see, for finite values of $\Gamma_\mu$ we obtain non-vanishing values of the density $n_h$.

For the sources $\tilde{\gamma}_{Q,\widetilde{\mathcal{H}}_\pm}$ in Eqs. (4.4)-(4.7) we will follow the formalism of Refs. [28, 32] where $\tilde{\gamma}_X \simeq j^0_X/\tau_X$, $\tau_X$ being the corresponding typical thermalization time. Thus we will use as sources of our diffusion equations,

$$\begin{aligned}
\tilde{\gamma}_Q \simeq &- \, v_\omega \, h_t^2 \, \Gamma_{\tilde{t}} \, \mathrm{Im}(A_t \mu_c) \; H^2(z) \, \beta'(z) \; \{\mathcal{F}_B(z) + \mathcal{G}_B(z)\} \\
\tilde{\gamma}_{\widetilde{H}_+} \simeq &- \, 2 \, v_\omega \, g^2 \, \Gamma_{\widetilde{\mathcal{H}}} \, \mathrm{Im}(M_2 \mu_c) \left\{ H^2(z) \, \beta'(z) \; [\mathcal{F}_F(z) + \mathcal{G}_F(z)] \right. \\
&+ \left. g^2 \, H^2(z) \cos 2\beta(z) \left[ H(z) H'(z) \sin 2\beta(z) + H^2(z) \cos 2\beta(z) \beta'(z) \right] \mathcal{H}_F(z) \right\} \\
\tilde{\gamma}_{\widetilde{H}_-} \simeq & \; 2 \, v_\omega \, g^2 \, \Gamma_{\widetilde{\mathcal{H}}} \, \mathrm{Im}(M_2 \mu_c) \left[ H(z) H'(z) \sin 2\beta(z) + H^2(z) \cos 2\beta(z) \beta'(z) \right] \\
& \left\{ \mathcal{K}_F(z) + 2 \left( \Delta + \bar{\Delta} \right) \mathcal{H}_F(z) \right\} \; . \qquad (4.8)
\end{aligned}$$

Notice that our sources, Eq. (4.8), are proportional to the wall velocity $v_\omega$, and so die when the latter goes to zero, which is a physical requirement.

We can find an approximate solution for $n_Q$ and $n_T$ by assuming that $\Gamma_Y$ and $\Gamma_{ss}$ are fast so that $n_Q/k_Q - n_T/k_T - (n_H + \rho\, n_h)/k_H = \mathcal{O}(1/\Gamma_Y)$ and $2\, n_Q/k_Q - n_T/k_T + 9\,(n_Q + n_T)/k_B = \mathcal{O}(1/\Gamma_{ss})$. In this case we can write

$$\begin{aligned}
n_Q &= \frac{k_Q \, (9 k_T - k_B)}{k_H \, (k_B + 9 k_Q + 9 k_T)} \; (n_H + \, \rho \, n_h) + \mathcal{O}\left( \frac{1}{\Gamma_{ss}}, \frac{1}{\Gamma_Y} \right) \\
n_T &= - \; \frac{k_T \, (9 k_Q + 2 k_B)}{k_H \, (k_B + 9 k_Q + 9 k_T)} \; (n_H + \, \rho \, n_h) + \mathcal{O}\left( \frac{1}{\Gamma_{ss}}, \frac{1}{\Gamma_Y} \right) \; . \qquad (4.9)
\end{aligned}$$

If the left-handed third generation squarks were light ($m_Q \sim T$) we could expect that all supersymmetric and supersymmetry breaking interactions arising from the $h_t\, H_2\, Q\, T$ term in the superpotential are in thermal equilibrium and similar in size, so that $\Gamma_{Y_1} \simeq \Gamma_{Y_2}$, or $\rho \ll 1$. In such case, which was considered in Ref. [40], the influence of $n_h$ in the quark densities $n_Q$ and $n_T$, through Eqs. (4.9), is $\rho$-suppressed although this suppression can be arguably mild depending on the particularly chosen value of $\rho$. However, in the case where left-handed squarks are heavy ($m_Q \gg T$), as preferred to get a good agreement of the MSSM with electroweak precision measurements, their corresponding interactions decouple, $\Gamma_{Y_1} \simeq 0$ and $\rho \simeq 1$. This is the case we will consider from here on.

We now take (for $\rho = 1$) the linear combinations of Eqs. (4.4), (4.5), (4.6) and (4.7) which are independent of $\Gamma_Y$ and $\Gamma_{ss}$. They are given by,

$$\begin{aligned}
v_\omega \left[ n'_Q + 2\, n'_T - n'_H \right] =& D_q \left[ n''_Q + 2\, n''_T \right] - D_h \, n''_H + \Gamma_m \left[ \frac{n_Q}{k_Q} - \frac{n_T}{k_T} \right] \\
&+ \Gamma_h \frac{n_H}{k_H} - \left( \tilde{\gamma}_Q + \tilde{\gamma}_{\widetilde{H}_+} \right) \qquad (4.10) \\
v_\omega \left[ n'_Q + 2\, n'_T - n'_h \right] =& D_q \left[ n''_Q + 2\, n''_T \right] - D_h \, n''_h + \Gamma_m \left[ \frac{n_Q}{k_Q} - \frac{n_T}{k_T} \right] \\
&+ \left[ \Gamma_h + 4\, \Gamma_\mu \right] \frac{n_h}{k_H} - \left( \tilde{\gamma}_Q + \tilde{\gamma}_{\widetilde{H}_-} \right) \; . \qquad (4.11)
\end{aligned}$$

When $n_Q$ and $n_T$ are replaced by the explicit solutions of Eqs. (4.9), as functions of $n_H$

and $n_h$, Eqs. (4.10) and (4.11) yield the system of coupled equations for $n_H$ and $n_h$:

$$v_\omega\,\mathcal{A}\begin{pmatrix} n'_H \\ n'_h \end{pmatrix} = \mathcal{D}\begin{pmatrix} n''_H \\ n''_h \end{pmatrix} - \mathcal{G}\begin{pmatrix} n_H \\ n_h \end{pmatrix} + \begin{pmatrix} f_+ \\ f_- \end{pmatrix} \tag{4.12}$$

where the sources are

$$f_\pm = \frac{G}{F+G}\left(\tilde{\gamma}_Q + \tilde{\gamma}_{\widetilde{H}_\pm}\right) \ , \tag{4.13}$$

with

$$\begin{aligned} F &\equiv 9k_Q k_T + k_Q k_B + 4k_T k_B \\ G &\equiv k_H(9k_Q + 9k_T + k_B) \ , \end{aligned} \tag{4.14}$$

and $\mathcal{A}$, $\mathcal{D}$ and $\mathcal{G}$ are the $2\times 2$ matrices,

$$\mathcal{A} = \begin{pmatrix} 1 & \frac{F}{F+G} \\ \frac{F}{F+G} & 1 \end{pmatrix} \ , \ \mathcal{D} = \begin{pmatrix} \overline{D}_q + \overline{D}_h & \overline{D}_q \\ \overline{D}_q & \overline{D}_q + \overline{D}_h \end{pmatrix}$$

$$\mathcal{G} = \begin{pmatrix} \overline{\Gamma}_m + \overline{\Gamma}_h & \overline{\Gamma}_h \\ \overline{\Gamma}_m & \overline{\Gamma}_m + \overline{\Gamma}_h + 4\,\overline{\Gamma}_\mu \end{pmatrix} \ , \tag{4.15}$$

with

$$\overline{D}_q \equiv \frac{F}{F+G}\,D_q, \quad \overline{D}_h \equiv \frac{G}{F+G}\,D_h$$

$$\overline{\Gamma}_i \equiv \frac{G}{F+G}\,\frac{\Gamma_i}{k_H}, \quad (i = m,\ h,\ \mu)\ . \tag{4.16}$$

The system (4.12) amounts to equations for $n_H$ and $n_h$, with sources induced by $\tilde{\gamma}_{\widetilde{Q}}$ and $\tilde{\gamma}_{\widetilde{H}_\pm}$, and by the same densities $n_{H,h}$ and their derivatives. It can be re-written as,

$$v_\omega\, n'_H = \overline{D}\, n''_H - \overline{\Gamma}\, n_H + f_+ + \Delta f_+ \tag{4.17}$$

$$v_\omega\, n'_h = D_h\, n''_h - [\Gamma_h + 4\Gamma_\mu]\,\frac{n_h}{k_H} + v_\omega\, n'_H - D_h\, n''_H + \Gamma_h \frac{n_H}{k_H} + \tilde{\gamma}_{\widetilde{H}_-} - \tilde{\gamma}_{\widetilde{H}_+} \tag{4.18}$$

where

$$\overline{D} = \overline{D}_q + \overline{D}_h, \quad \overline{\Gamma} = \overline{\Gamma}_m + \overline{\Gamma}_h \tag{4.19}$$

$$\Delta f_+ = -\frac{F}{F+G}\,v_\omega n'_h + \overline{D}_q\, n''_h - \overline{\Gamma}_m\, n_h\ . \tag{4.20}$$

We have solved the system (4.12) numerically and the results are presented in section 5. However a very useful analytical approximation can be worked out as follows. Using Eq. (4.6) and the approximate relations (4.9) we can write for $n_h$ the following equation,

$$v_\omega\, n'_h = D_h\, n''_h - [\Gamma_h + 4\Gamma_\mu]\,\frac{n_h}{k_H} + \tilde{\gamma}_{\widetilde{H}_-}\ . \tag{4.21}$$

In this way the equation for $n_h$ has been decoupled from the other equations and can be easily solved. On the other hand, from (4.17) and the expression for $\Delta f_+$, Eq. (4.20), we see that $n_h$ acts as a source for $n_H$, and the equations (4.17) and (4.21) can be solved analytically.

We will only quote the solutions in the symmetric phase ($z<0$) since that would be needed to compute the baryon asymmetry from $n_L(z)$, as we will see. Finally we will impose boundary conditions $n_h(\pm\infty)$ and $n_H(\pm\infty)$ and continuity of the functions and first derivatives at $z=0$.

From Eq. (4.21) we obtain the solution for $n_h(z)$, for $z\leq 0$ as,

$$n_h(z) = \mathcal{A}_h\, e^{z\alpha_+} \tag{4.22}$$

where

$$\mathcal{A}_h = \frac{2}{\sqrt{v_\omega^2+4\Gamma_1 D_h}+\sqrt{v_\omega^2+4\Gamma_2 D_h}} \int_0^\infty d\zeta\; \tilde{\gamma}_{\widetilde{H}_-}(\zeta)\, e^{-\zeta\beta_+} \tag{4.23}$$

and

$$\begin{aligned}
\alpha_\pm &= \frac{1}{2D_h}\left\{v_\omega \pm \sqrt{v_\omega^2+4\Gamma_1 D_h}\right\}\\
\beta_\pm &= \frac{1}{2D_h}\left\{v_\omega \pm \sqrt{v_\omega^2+4\Gamma_2 D_h}\right\}\\
\Gamma_2 &= \frac{\Gamma_h+4\Gamma_\mu}{k_H}\\
\Gamma_1 &= \frac{4\Gamma_\mu}{k_H}\ .
\end{aligned} \tag{4.24}$$

Note that, from expression (4.24), the coefficient $\mathcal{A}_h$ behaves as $\Gamma_\mu^{-1/2}$, in the limit of large $\Gamma_\mu$, and so the $n_h$ density tends to zero when $\Gamma_\mu$ tends to infinity, as anticipated.

From Eq. (4.17), the solution for $n_H(z)$, for $z\leq 0$ is given by

$$n_H(z) = \mathcal{A}_H\, e^{z\,\alpha_+} + \mathcal{B}_H\, e^{z\,v_\omega/\overline{D}} \tag{4.25}$$

where

$$\begin{aligned}
\mathcal{B}_H =& \mathcal{A}_0 + \mathcal{A}_h \frac{F}{F+G}\left\{\frac{D_q\alpha_+-v_\omega}{\overline{D}}\left[\frac{1}{\lambda_+}+\frac{1}{\alpha_+-v_\omega/\overline{D}}\right]\right.\\
&+\left.\frac{1}{\overline{D}\lambda_+}\left[v_\omega - D_q(\alpha_++\lambda_+)+\frac{\lambda_+-\alpha_-}{\lambda_+-\beta_+}\;\frac{F\lambda_+(-v_\omega+D_q\lambda_+)-G\,\Gamma_m/k_H}{F(\lambda_+-\beta_-)}\right]\right\}\\
&-\frac{1}{\overline{D}\lambda_+}\,\mathcal{A}_\lambda\,\frac{1}{F+G}\;\frac{\alpha_+-\beta_-}{(\lambda_+-\beta_+)(\lambda_+-\beta_-)}\;\left[F\lambda_+(-v_\omega+D_q\lambda_+)-G\Gamma_m/k_H\right]
\end{aligned} \tag{4.26}$$

and

$$\mathcal{A}_H = -\,\mathcal{A}_h\,\frac{F}{F+G}\;\frac{D_q\,\alpha_+-v_\omega}{\overline{D}\alpha_+-v_\omega} \tag{4.27}$$

with

$$
\begin{aligned}
\mathcal{A}_\lambda =& \frac{2}{\sqrt{v_\omega^2+4\Gamma_1 D_h}+\sqrt{v_\omega^2+4\Gamma_2 D_h}} \int_0^\infty d\zeta\ \tilde{\gamma}_{\widetilde{H}_-}(\zeta)\, e^{-\zeta\lambda_+} \\
\mathcal{A}_0 =& \frac{1}{\overline{D}\lambda_+} \int_0^\infty d\zeta\ f_+(z)\ e^{-\zeta\lambda_+}
\end{aligned}
\tag{4.28}
$$

and

$$
\lambda_\pm = \frac{1}{2\overline{D}} \left\{ v_\omega \pm \sqrt{v_\omega^2 + 4\overline{\Gamma}\ \overline{D}} \right\} \ . \tag{4.29}
$$

Since we assume the sphalerons are inactive inside the bubbles, the baryon density is constant in the broken phase and satisfies, in the symmetric phase, an equation where $n_L$ acts as a source [28] and there is an explicit sphaleron-induced relaxation term [43, 40]

$$
v_\omega n_B'(z) = -\theta(-z)\left[n_F \Gamma_{ws} n_L(z) + \mathcal{R} n_B(z)\right] \tag{4.30}
$$

where $n_F = 3$ is the number of families and $\mathcal{R}$ is the relaxation coefficient [43],

$$
\mathcal{R} = \frac{5}{4} n_F\, \Gamma_{ws} \ . \tag{4.31}
$$

Eq. (4.30) can be solved analytically and gives, in the broken phase $z \geq 0$, a constant baryon asymmetry,

$$
n_B = -\frac{n_F \Gamma_{ws}}{v_\omega} \int_{-\infty}^0 dz\, n_L(z)\ e^{z\mathcal{R}/v_\omega} \ . \tag{4.32}
$$

Using now the explicit solutions for $n_H$ and $n_h$ given in Eqs. (4.25) and (4.22), we can cast the explicit solution for the baryon asymmetry as,

$$
n_B = n_F\, \Gamma_{ws}\ \frac{5k_Q k_B + 8k_T k_B - 9k_Q k_T}{k_H\,(k_B + 9k_Q + 9k_T)} \left\{ \frac{\mathcal{A}_H + \mathcal{A}_h}{\mathcal{R} + v_\omega \alpha_+} + \frac{\overline{D}\mathcal{B}_H}{\overline{D}\mathcal{R} + v_\omega^2} \right\} \tag{4.33}
$$

where all symbols used in Eq. (4.33) have been previously defined.

The validity of our analytical approximation is guaranteed by the dominance of $n_H$ over $n_h$, which in turn is related to the $\tan\beta$ suppression of $\tilde{\gamma}_{\widetilde{H}_-}$ and the presence of $\Gamma_\mu$. In fact were we working in the limit $\Gamma_\mu \to \infty$ we would find that the density $n_h$ is negligible. On the other hand, in the limit $\Gamma_\mu \to 0$ and $\tan\beta \simeq 1$ we would really expect $n_h > n_H$, due to the dominance of $\tilde{\gamma}_{\widetilde{H}_-}$ over $\tilde{\gamma}_{\widetilde{H}_+}$, at least for large values of $m_A$ where the $\Delta\beta$ suppression of $\tilde{\gamma}_{\widetilde{H}_+}$ is more severe. However small values of $\tan\beta$, as we noticed earlier in this paper, are strongly disfavored in our scenario by recent LEP bounds on the Higgs mass. Hence, we have found that the analytical approximation is accurate with an error which depends on the chosen values of the supersymmetric parameters, but it is always much smaller than the other uncertainties involved in the final calculation. In section 5 we will provide explicit comparison with the numerical result, while all plots will be done using the numerical solution of system (4.12).

# 5 Numerical results

In this section we present the numerical results for the baryon asymmetry computed in section 4 and, in particular, of the baryon-to-entropy ratio $\eta \equiv n_B/s$, where the entropy density is given by,

$$s = \frac{2\pi^2}{45} g_{eff}\, T^3 \tag{5.1}$$

with $g_{eff}$ being the effective number of relativistic degrees of freedom. The profiles $H(z)$, $\beta(z)$ have been accurately computed in the literature [14, 26]. For the sake of simplicity, in this paper we will use a kink approximation [29]

$$\begin{aligned} H(z) =& \frac{1}{2}\, v(T) \left(1 - \tanh\left[\alpha\left(1 - \frac{2\,z}{L_\omega}\right)\right]\right) \\ \beta(z) =& \beta - \frac{1}{2}\, \Delta\beta \left(1 + \tanh\left[\alpha\left(1 - \frac{2\,z}{L_\omega}\right)\right]\right) \ . \end{aligned} \tag{5.2}$$

This approximation has been checked to reproduce the exact calculation of the Higgs profiles within a few percent accuracy [22], provided that we borrow from the exact calculation the values of the thickness $L_\omega/2\alpha$ and the variation of the angle $\beta(z)$ along the bubble wall, $\Delta\beta$, as we will do. In particular we will take $\alpha \simeq 3/2$, $L_\omega = 20/T$, and we have checked that the result varies only very slowly with those parameters, while we are taking the values of $\Delta\beta$ which are obtained from the two-loop effective potential used in our calculation.

The calculation of the wall velocity $v_\omega$ is a very complicated phenomenon involving the hydrodynamics of the bubble interacting with the surrounding plasma. Some progress has been recently reported in this direction [45] indicating that, in the case of the MSSM, the wall is extremely non-relativistic, $v_\omega \ll 1$, and can be as slow as $v_\omega = 0.01$. Unless explicitly stated, in the numerical analysis of this section, we adopt the value $v_\omega = 0.05$, although the variation of the baryon asymmetry with respect to $v_\omega$ will also be analyzed. The widths, $\Gamma_m$, $\Gamma_h$ and $\Gamma_Y$ are as in Refs. [28, 29], while we are taking $\Gamma_\mu \simeq 0.1\,T$ and $\rho = 1$, in agreement with the large value we use for the left-handed third-generation squark masses, $m_Q \gtrsim 1$ TeV, which makes them decoupling from the thermal bath. On the other hand, and consistently with the latter assumption (which is required to render the MSSM in agreement with the Higgs mass bounds coming from LEP), the contribution to $n_B$ from the squark source, $\tilde{\gamma}_Q$, is negligible. The "observable" value for $\eta$ consistent with Big Bang Nucleosynthesis (BBN) has been considered to be $\eta_{\rm BBN} \sim 4\times 10^{-11}$ [46]. Finally we will consider the third generation squark mass and mixing parameters, $m_Q = 1.5$ TeV and $A_t = 0.5$ TeV, and $\tan\beta = 20$ and have checked that, for all plots in this section, the phase transition is strong enough first order, $v(T_c)/T_c \gtrsim 1$, and the Higgs mass is, within the accuracy of our calculations, $m_h \simeq 110$–$115$ GeV. These values are in rough agreement with present 95 % C.L. bounds on the Higgs mass coming from LEP, or even with the present excess of events observed at LEP, consistent with the detection of a SM-like Higgs at the runs with the highest center of mass energies, $\sqrt{s} > 206$ GeV. We will comment more about the LEP constraints in the next section.

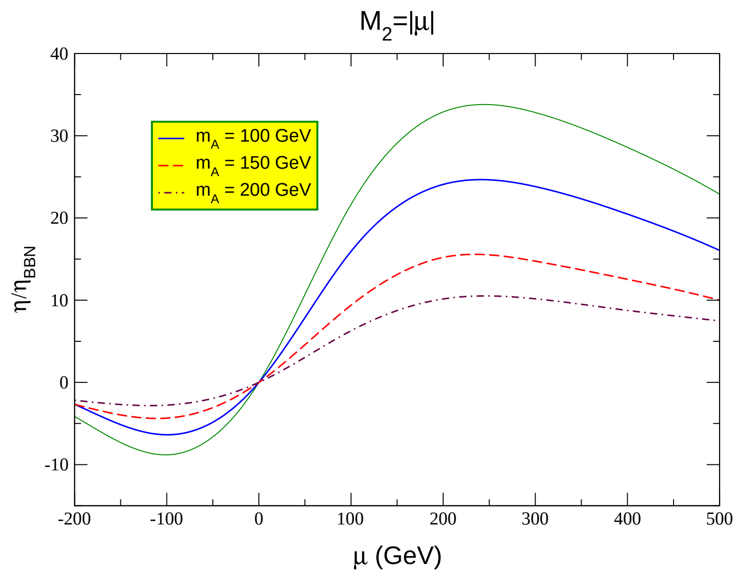


Figure 1: Plot of $\eta/\eta_{\rm BBN}$ as a function of $\mu$ for $M_2 = |\mu|$, $m_A = 100$ GeV (thick solid curve), $m_A = 150$ GeV (dashed curve) and $m_A = 200$ GeV (dash-dotted curve), and the rest of parameters as indicated in the text. The thin solid curve corresponds to the case $m_A = 100$ GeV when the approximate analytical solution in (4.33) is used.

In Fig. 1 we plot the ratio $\eta/\eta_{\rm BBN}$ for the values of the supersymmetric parameters that have just been described, $M_2 = |\mu|$, $\sin\varphi_\mu = 1$ and several values of the pseudoscalar Higgs mass $m_A$. Therefore, since $\eta$ is (almost) linear in $\sin\varphi_\mu$, one can read from Fig. 1 the value of $1/\sin\varphi_\mu$ that would reproduce $\eta_{\rm BBN}$. This observation applies to all plots presented in this section, where we have fixed $\sin\varphi_\mu = 1$. It follows that the region of parameters where we find $|\eta/\eta_{\rm BBN}| < 1$, is forbidden in all plots. For the value $m_A = 100$ GeV, we have presented both the exact result (thick solid curve), based on the numerical solution of Eqs. (4.12), and the approximate result (thin solid curve), based on the approximate analytical solution (4.33). We see that for values where $n_B/s$ is sizeable the discrepancy between the analytical and the numerical result is $\lesssim 30$ %. For the other curves in Fig. 1, as well as for the rest of plots in this paper, we will use the (exact) numerical solution of Eqs. (4.12).

We are, in Fig. 1, close to the resonance region discussed in Ref. [29], which is smoothed by the all order resummation in Higgs mass insertions. The departure from the resonance is exemplified in Fig. 2, where we plot $\eta/\eta_{\rm BBN}$ as a function of $\mu$ for $M_2 = 200$ GeV, $m_A = 150$ GeV and the other supersymmetric parameters as in Fig. 1.

In Fig. 3 we plot $\eta/\eta_{\rm BBN}$ as a function of $m_A$ for $M_2 = \mu = 200$ GeV (solid curve) and $M_2 = 200$ GeV, $\mu = 300$ GeV (dashed curve), and other supersymmetric parameters as in Fig. 1. Finally in Fig. 4 we plot $\eta/\eta_{\rm BBN}$ as a function of $v_w$ for $M_2 = \mu = 200$ GeV, $m_A = 150$ GeV and the other parameters as in Fig. 1. The maximum of this curve comes from the interplay between the relaxation and source terms in the equation for $n_B$, Eq. (4.30).

The numerical results exhibited in the plots of this section are an improvement of

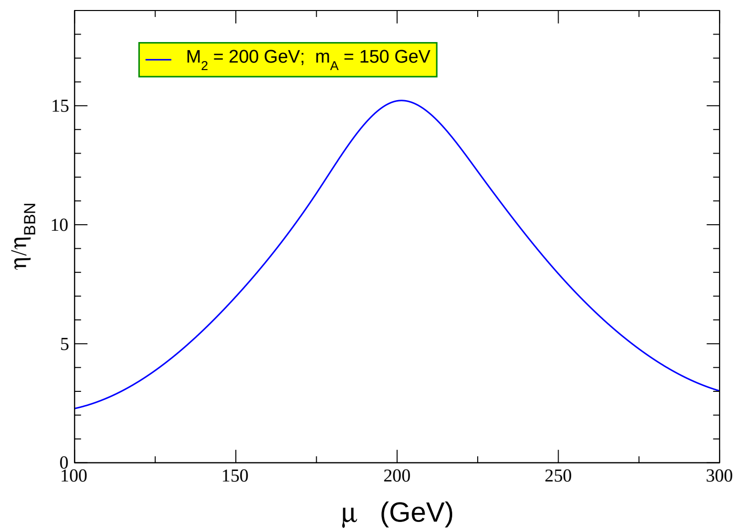


Figure 2: Plot of $\eta/\eta_{\rm BBN}$ as a function of $\mu$ for $M_2 = 200$ GeV and $m_A = 150$ GeV.

our previous results, Ref. [29], and include the all order resummation of the Higgs mass insertions in the current determination, as well as inclusion of finite $\Gamma_\mu$-effects in the diffusion equations. Since the first of these effects smooths out the resonant behaviour, which enhances the determination of $n_B$ for $M_2 = |\mu|$, while the second one slightly enhances $n_B$, our present numerical results are in rough agreement with those of Ref. [29]. On the other hand if we compare our numerical results with the recent ones of Ref. [40], that use WKB methods and values of $\rho < 1$ to deduce the source terms in the diffusion equations, we observe a discrepancy of a few orders of magnitude. However, we have been communicated [47] by the authors of Ref. [40] to have detected a problem in their numerical codes which enhances their numerical results by some orders of magnitude and that might explain part of this discrepancy. As explained above, large values of $m_Q$, implying $\rho = 1$, are necessary in order to fulfill the present experimental Higgs mass bounds.

# 6 Higgs mass constraints

In this section, we shall comment on the constraints coming from Higgs searches at LEP. The LEP experiments at CERN have collected data during the year 1999 at various energies between 192 GeV and 202 GeV, for a total integrated luminosity of about 900 $pb^{-1}$. A combined limit on the Standard Model Higgs mass of about 108 GeV at the 95 % C.L. was obtained, due to the absence of any significant Higgs signal in the LEP data [48]. Preliminary results of this year run [49] show that this limit moved up by a few GeV (up to about 113.2 GeV). More interesting, a slight excess of events, about 3 standard deviations above the SM predictions, has been observed, consistent with a SM like Higgs in the range of masses of about 113–116 GeV.

The present Higgs mass constraints become particularly relevant for small values of

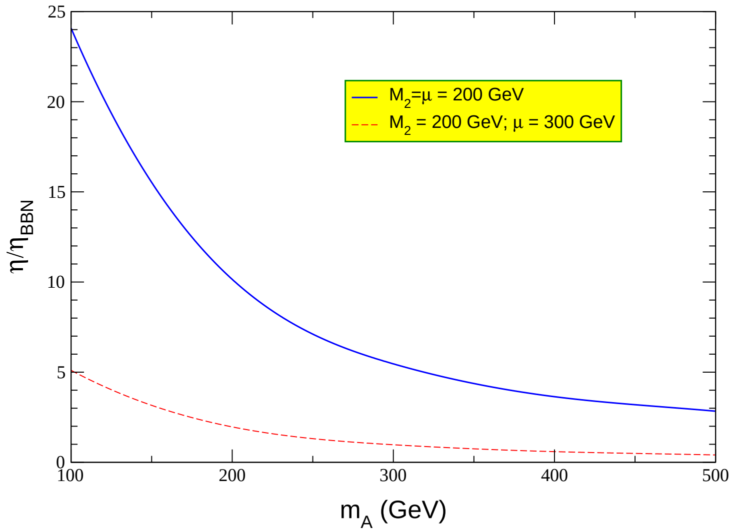


Figure 3: Plot of $\eta/\eta_{\rm BBN}$ as a function of $m_A$ for $M_2 = \mu = 200$ GeV (solid curve) and $M_2 = 200$, $\mu = 300$ GeV (dashed curve).

$\tan\beta$, $\tan\beta < 5$. In this case, due to the behaviour of the Higgs boson couplings to fermions and gauge bosons, the SM Higgs mass constraints translate with almost no variations into a bound on the lightest CP-even Higgs boson mass [4]. For values of $v(T_c)/T_c \gtrsim 1$ and $\tan\beta < 5$, and for left-handed stop masses smaller than $\sim 3$ TeV, the lightest CP-even Higgs mass never exceeds 105 GeV. Therefore, the mechanism of electroweak baryogenesis demands either values of $\tan\beta > 5$ or unnaturally large values of $m_Q$ [5].

Large values of $\tan\beta$ move the value of the Higgs boson mass, with relevant couplings to the gauge bosons, to larger values. However, if the values of the left-handed stop parameters are restricted to be below 3 TeV, for $v(T_c)/T_c \gtrsim 1$, the Higgs mass cannot exceed 115 GeV. Observe that these values are a few GeV higher than those obtained previously in Ref. [20], since in that reference we restricted ourselves to the case of left-handed stop masses below 1 TeV. The observed excess of events, with $b\bar{b}$ invariant masses of about 114 GeV, would be consistent with electroweak baryogenesis for large values of $\tan\beta$ and large values of the left-handed stop mass parameters $m_Q \gtrsim 1$ TeV, as the ones considered in the previous section.

What would happen if the excess of events present at LEP would not correspond to a Higgs signal, but would turn out to be a statistical fluctuation with the final outcome of an ultimate exclusion limit for a SM-like Higgs with mass below 115 GeV? To analyze this, let us stress that at large values of $\tan\beta$ the coupling of this Higgs boson to bottom

[4] In the presence of CP-violation, the Higgs mass eigenstates will not be CP-eigenstates. In our analysis we have used the CP conserving structure for the Higgs sector. This should lead to a good approximation if CP-violating effects in the Higgs potential are small, as happens when $\arg(\mu_c A_t) \simeq 0$. A more general analysis, similar to the one performed in Refs. [50, 51] would be appropriate to consider more general CP-violating effects.

[5] We have checked that, for the values of the stop mixing parameters consistent with electroweak baryogenesis, no significant modification of these bounds is obtained after considering CP-violating effects in the Higgs potential [50].

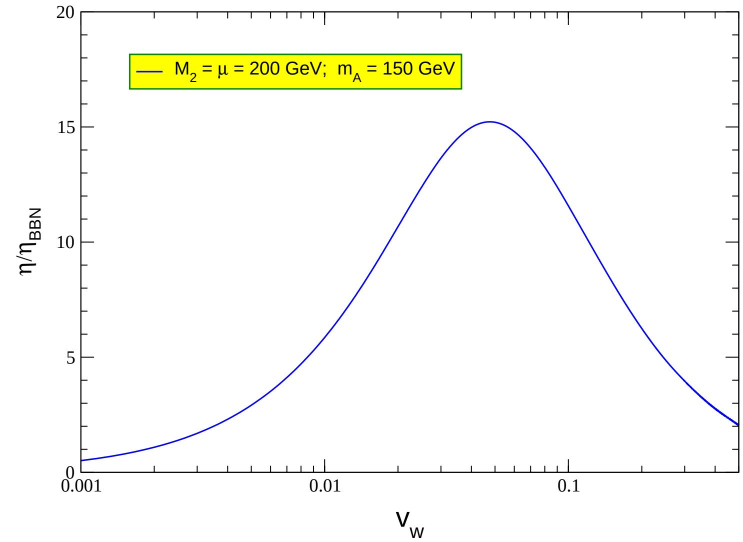


Figure 4: Plot of $\eta/\eta_{\rm BBN}$ as a function of $v_w$ for $M_2 = \mu = 200$ GeV and $m_A = 150$ GeV.

quarks can be significantly lower than in the SM [52, 50] with a corresponding reduction of the Higgs mass bound. These variations can only occur for small values of the CP-odd Higgs mass $m_A$, of order of the lightest CP-even Higgs boson mass. Unlike the case of small values of $\tan\beta$, for values of $\tan\beta > 10$, the values of $v(T_c)/T_c$ are only weakly dependent on the exact value of $m_A$. Intuitively, this can be understood by the fact that for large values of $\tan\beta$, the CP-odd Higgs can be approximately identified with the imaginary part of the neutral component of the Higgs doublet $H_1$, while the Higgs doublet which acquires vacuum expectation value is mainly $H_2$ ($v_2 \gg v_1$).

For the values of $A_t$ and $\mu$ consistent with electroweak baryogenesis, a reduction of the coupling of the CP-even Higgs boson to the bottom quark would demand not only small values of $m_A \simeq 100$–$150$ GeV, but also large values of $\tan\beta > 10$ and of $|\mu A_t|/m_Q^2 > 0.1$ (the larger $\tan\beta$, the easier suppressed values of the bottom quark coupling are obtained). We have checked that, assuming small CP-violating effects in the Higgs potential and in the Higgs-fermion couplings, and for values of $m_Q \simeq 1$ TeV, $v(T_c)/T_c \gtrsim 1$ and $|\mu| < 500$ GeV, a significant reduction of the coupling of the Higgs to bottom quarks only occur for $\tan\beta \gtrsim 30$. Therefore, if the excess of events observed at LEP is not associated with a Higgs signal, strong constraints on the electroweak baryogenesis scenario within the MSSM will be obtained.

# 7 Conclusions

In this article, we have performed a computation of the scalar- and fermion- CP-violating currents induced by the expansion of a true-vacuum bubble in the false vacuum plasma, within the framework of the minimal supersymmetric standard model. We made use of the Keldysh formalism and we have defined a systematic way of obtaining the currents in an expansion of derivatives of the Higgs fields, to all orders of the Higgs background insertions.

Although our method is similar to the one previously used by some of us in Ref. [29], our results differ from those presented in our previous work in several respects. First of all, they include a resummation of corrections associated with higher order of the Higgs background insertions. These corrections have two important effects. The first one is to dilute the resonant behaviour obtained in Ref. [29] for values of $|\mu| = M_2$. The second one is the appearance of a contribution proportional to $H_2\partial^\mu H_1 + H_1\partial^\mu H_2$ to the vector Higgsino current $j^\mu_{\widetilde{\mathcal{H}}_+}(z)$. This means that, as first observed in Refs. [28, 38], the vector Higgsino current does not vanish for large values of the CP-odd Higgs mass. Our method provides a way of obtaining the value of this non-vanishing contribution in a self-consistent way. In addition, we have also computed the axial Higgsino current $j^\mu_{\widetilde{\mathcal{H}}_-}(z)$, whose components are proportional to $H_2\partial^\mu H_1 + H_1\partial^\mu H_2$. Therefore, as first observed in Ref. [40], the chiral current is not suppressed for large values of the CP-odd Higgs mass and hence may become relevant in this regime.

The vector and axial Higgsino currents, $j^\mu_{\widetilde{\mathcal{H}}_\pm}(z)$, were used to determine the baryon asymmetry of the Universe, $n_B/s$. The computation of $n_B$ demands the solution of diffusion equations, with sources determined through $j^\mu_{\widetilde{\mathcal{H}}_\pm}(z)$. Following the method developed in Refs. [28, 32], we assumed that the sources are proportional to the temporal component of the currents, with a constant of proportionality given by the Higgsino width. Within this approximation, we computed the functional dependence of $n_B$ on the soft supersymmetry breaking parameters and on the bubble wall parameters. The most important parameters turn out to be the gaugino and Higgsino mass parameters, $|\mu_c|$ and $M_2$, their relative phase $\arg(\mu_c M_2)$, (equal to $\varphi_\mu$ in the basis in which $M_2$ is real) as well as the CP-odd Higgs mass $m_A$ and $\tan\beta$. We have also required that the condition of preservation of the baryon asymmetry $v(T_c)/T_c \gtrsim 1$ is fulfilled, what demands a light right-handed stop and, due to the present Higgs mass constraints coming from LEP (see section 6), also large values of the ratio of Higgs vacuum expectation values, $\tan\beta > 5$.

Under the above conditions, we have determined the value of $n_B$, compared to the value predicted by Big Bang Nucleosynthesis, for a value of $\sin\varphi_\mu = 1$. The ratio of the theoretically obtained to the BBN predicted baryon asymmetry can be reinterpreted as the inverse of the value of $\sin\varphi_\mu$ needed to obtain a value of $n_B$ in agreement with the BBN predictions. We conclude that, for small values of $m_A \simeq 100$ GeV and $|\mu| \simeq M_2$, values as low as $\varphi_\mu \simeq 0.04$ can lead to acceptable values of $n_B$. The predicted value of the phase $\varphi_\mu$ increases for larger values of $m_A$ and/or for $|\mu| \neq M_2$, but still there is a large fraction of parameter space in which the computed baryon number is in good agreement with BBN predictions, for phases such that $\sin\varphi_\mu \simeq 0.04$–1.

Values of $\varphi_\mu \gtrsim 0.04$ can lead to acceptable phenomenology if either peculiar cancellations in the squark and slepton contributions to the neutron and electron electric dipole moments (EDM) occur [53], and/or if the first and second generation of squarks are heavy [54]. This second possibility is quite appealing and, as has been recently demonstrated [55], leads to acceptable phenomenology, including the dark matter constraints [6].

[6]Third generation squarks would still contribute to the neutron and electron EDM, via two loop diagrams involving the would-be CP-odd Higgs boson [56]. These contributions can become sizeable at large values of $\tan\beta$, although they tend to be suppressed for small values of the mixing in the stop

Another important observable which, similarly to the value of the baryon number, depends on the precise value of the mass parameters in the gaugino, Higgsino and third generation squark sectors, as well as on the charged Higgs mass, is the rate of the rare decay $b \to s\gamma$ [57]. For small values of the charged Higgs and stop masses, and for moderate values of $A_t/m_Q$ and $|\mu|/m_Q$, the chargino-stop contribution, as well as the charged Higgs contribution, may become large for large values of $\tan\beta$ [58, 59, 60]. In scenarios with heavy first and second generation squarks, however, flavor violation couplings involving the third generation squarks could be non-negligible [54] and therefore the gluino-sbottom contributions to this rare decay rate may be enhanced [61]. Since these last contributions are strongly model dependent, and may be larger than the charged Higgs and chargino-stop ones, we have not imposed the $b \to s\gamma$ constraints in our analysis.

Finally, we have discussed the effect of the Higgs mass constraints coming from LEP. The preliminary data coming from the LEP experiments imply a lower bound on the mass of a SM-like Higgs boson of about 113 GeV. A small excess, consistent with a SM-like Higgs boson with a mass slightly above that value has also been observed. These relatively large values of the Higgs mass are consistent with electroweak baryogenesis within the MSSM if the value of $\tan\beta$ is large, $\tan\beta > 5$, if the left-handed stops are heavy $m_Q \gtrsim 1$ TeV, and if the stop mixing parameter is not small, $A_t \gtrsim 0.25\ m_Q$. On the other hand, for these values of the Higgs mass, values of $A_t \gtrsim 0.4\ m_Q$ make the phase transition weaker, leading to values of $v(T_c)/T_c$ that are in conflict with the condition of preservation of the baryon asymmetry. It is important to emphasize, however, that even if the CP-even Higgs boson coupling to the gauge boson is SM like, it can evade the LEP bounds if its coupling to the bottom quark is strongly suppressed, what can occur for very large values of $\tan\beta$, $\tan\beta \gtrsim 30$. More relevantly, if the excess of events at LEP has its origins in the presence of a SM-like Higgs boson of mass of about 113–115 GeV, one of the predictions of electroweak baryogenesis, namely the presence of a light neutral Higgs boson with SM-like couplings to the gauge bosons and a mass not larger than 115 GeV would have been fulfilled.

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sector, as the ones required for electroweak baryogenesis.

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T. Camporesi, for the DELPHI Collaboration,
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J.-J. Blaising, for the L3 Collaboration,
`http://l3www.cern.ch/analysis/latestresults.html`;
C. Rembser, for the OPAL Collaboration,
`http://opal.web.cern.ch/Opal/PPwelcome.html`;
C. Tully, for the LEP Working Group for Higgs boson searches,
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