Baryogenesis and the Thermalization Rate of Stop

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Abstract

We take the first steps towards the complete computation of the thermalization rate of the supersymmetric particles involved in electroweak baryogenesis by computing the thermalization rate of the right-handed stop from the imaginary part of the two-point Green function. We use improved propagators including resummation of hard thermal loops. The thermalization rate is computed at the one-loop level in the high temperature approximation as a function of $M_{\tilde{t}_R}(T)$. We also give an estimate for the magnitude of the two-loop contributions which dominate the rate for small $M_{\tilde{t}_R}(T)$. If the stop is non-relativistic with $M_{\tilde{t}_R}(T) \gg T$, thermalization takes place by decay and is very fast.

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It is now commonly accepted that the generation of the baryon asymmetry during the electroweak phase transition \cite{1} requires some new physics at the weak scale. Therefore electroweak baryogenesis in the framework of the Minimal Supersymmetric Standard Model (MSSM) has attracted much attention in the past years, with particular emphasis on the strength of the phase transition \cite{2} and the mechanism of baryon number generation \cite{3, 4, 5}. It has recently been shown both analytically \cite{6, 7} and by lattice simulations \cite{8} that the phase transition can be sufficiently strongly first order if the lightest stop is not much heavier than the top quark, the ratio of the vacuum expectation values of the two neutral Higgses tan $\beta$ is smaller than $\sim 4$ and the lightest Higgs is lighter than about 85 GeV.

Moreover, the MSSM may contain, besides the CKM matrix phase, new CP-violating phases in the soft supersymmetry breaking parameters associated with the stop mixing angle and with the gaugino and neutralino mass matrices. In the MSSM large values of the stop mixing angle are restricted in order to preserve a sufficiently strong first order electroweak phase transition. Therefore, an acceptable baryon asymmetry may only be generated through a delicate balance between the values of the different soft supersymmetry breaking parameters contributing to the stop mixing parameter, and their associated CP-violating phases \cite{4, 9}. In the MSSM it turns out that the main contribution to the baryon asymmetry comes from charginos and neutralinos and the phase of the parameter $\mu$ must be larger than about 0.1 to be responsible for the generation of the observed baryon asymmetry \cite{4, 9, 5}. If the strength of the electroweak phase transition is enhanced by the presence of some new degrees of freedom beyond the ones contained in the MSSM, e.g. some extra standard model gauge singlets, light stops (predominantly the right-handed ones) and charginos/neutralinos are expected to give quantitatively the same contribution to the final baryon asymmetry.

Extra and sizeable CP-violating phases in the MSSM are therefore a necessary ingredient for a successful electroweak baryogenesis scenario. If the particles involved in the process of baryon number generation thermalize rapidly, CP-violating sources however loose their coherence and are diminished. This phenomenon can be intuitively understood by means of the following example: let us focus on the right-handed stop current and imagine that a right-handed stop scatters off the advancing bubble wall and is transformed into a left-handed stop. If the latter scatters off the Higgs background once again, a right-handed stop reappears in the plasma; since in both interactions with the wall CP is violated at the vertices because of the explicit phase in $A_t\mu$, quantum interference may give rise to a final right-handed current. This, however, only takes place if, along their way, the stops do not interact with the surrounding plasma and disappear: large thermalization rates reduce the final baryon asymmetry.

Moreover, CP-violating currents in supersymmetric baryogenesis are more easily built up if the degrees of freedom in the stop and in the gaugino/neutralino sectors
are nearly degenerate in mass \([4]\). For instance, phases of \(\mu\) smaller than 0.1 are only consistent with the observed baryon asymmetry for values of \(|\mu|\) of the order of the gaugino mass parameters. This is due to a large enhancement of the computed baryon asymmetry for these values of the parameters. This resonant behaviour is associated with the possibility of absorption (or emission) of Higgs quanta by the propagating supersymmetric particles. For momenta of the order of the critical temperature, this can only take place when, for instance, the Higgsinos and gauginos do not differ too much in mass. By using the Uncertainty Principle, it is easy to understand that the width of this resonance is expected to be proportional to the thermalization rate of the particles giving rise to the baryon asymmetry \([4]\).

It is therefore clear that the computation of the thermalization rate of the particles responsible for supersymmetric baryogenesis represents a necessary step towards the final computation of the baryon number. Despite its relevance, no substantial effort has been devoted to a detailed computation of the decay width of SUSY particles in the thermal bath. The goal of this paper is to take the first step towards a complete evaluation of the thermalization rate of the supersymmetric particles involved in the generation of the baryon asymmetry, right-handed stops, charginos and neutralinos. At present, we restrict ourselves to the computation of the thermalization rate \(\Gamma_R\) of light right-handed stops from the imaginary part of the two-point Green function in the unbroken phase of the MSSM. We use improved propagators including resummation of hard thermal loop; the thermalization rate is computed at the one-loop level in the high temperature approximation and an estimate is given for two-loop contributions. Finite mass effects as well as the computation of the thermalization rate for charginos and neutralinos will be presented in a longer publication \([11]\).

Let us first present some technical details needed for the computation of the right-handed stop thermalization rate. A great deal of the formalism may be found in \([12]\), but we briefly summarise here some details and also present some new necessary technical tools.

The thermalization rate \(\Gamma_R\) depends upon the imaginary part of the two point Green function \(\Sigma_R\) via the relation

\[
\Gamma_R = -\frac{\text{Im} \Sigma_R}{\omega},
\]

where \(\omega = \omega(k)\) is the energy of a given mode. In the present paper we shall focus on the long wave-length stops with \(\omega \simeq M_{\tilde{t}_R}\). At the one-loop level the quantity \(\text{Im} \Sigma_R\) receives different contributions from diagrams involving SM fermions, charginos, neutralinos, scalar and gauge bosons\([1]\). At two loop level there are several potentially relevant sources that contribute to the thermalization rate. The strong interactions

\[\text{We assume R-parity conservation so that one of the two particles running in the loop is always a superpartner.}\]
contribute, due to diagrams including gluon interactions being proportional to the square of the strong coupling constant. Also, there are possible interactions of two right-handed stops appearing from $D$-terms of the superpotential. They are, however, proportional to the fourth power of the electroweak gauge couplings (and so are two-loop diagrams containing electroweak gauge bosons) and therefore can be neglected. Possible interactions arise also from the $F$-terms of the superpotential. Of particular importance is the one proportional to the top-quark Yukawa coupling given by $h^2 |H_2|^2 |\tilde{t}_R|^2$. The diagram including (only) this interaction is not suppressed by the coupling, but is expected to be small due to phase space suppression unless the right-handed stop is light. The very same argument can be used in the case of gluon interactions, too.

The imaginary parts are cuts across the relevant diagrams and correspond to the differences between the absorption and the emission rates, or between decay and inverse decay rates. For brevity we denote these respectively by “absorption” and “decay”. For all the fermions and bosons unbroken phase thermal corrections must be taken into account. They merely change the pole structure of the propagators and, to leading order in the temperature $T$ and for the massless SM fermions, they are the hard thermal loops. Accordingly, there appear energy thresholds which, in particular for the fermionic loops, are rather complicated because of the complicated nature of fermionic dispersion relations, which have two branches called particles and holes. However, the vertex corrections relevant in the present discussion contains no hard thermal loops and can therefore be omitted.

The plasma corrected left-handed Dirac fermion propagator reads generically ($P_\mu = (p_0, \mathbf{p})$, and $p = |\mathbf{p}|$)

$$S(P) = P_L \frac{F_\mu \gamma^\mu + \mu}{F^2 - \mu^2} P_R,$$

where $P_L$ and $P_R$ are the left- and right-handed projections and $\mu$ is the bare mass, and

$$\gamma_\mu F^\mu = [1 + a(P)] P_\mu \gamma^\mu + b(P) \gamma^0.$$  

In the case of a left-handed field, at high temperature $T \gg \mu$ the functions $a$ and $b$ are given by

$$a_L(P) = \frac{m_L^2}{p^2} \left( 1 - \frac{p_0}{2p} \ln \left| \frac{p_0 + p}{p_0 - p} \right| \right)$$

and

$$b_L(P) = \frac{m_L^2}{p^2} \left[ -\frac{p_0}{p} + \frac{1}{2} \left( \frac{p_0^2}{p^2} - 1 \right) \ln \left| \frac{p_0 + p}{p_0 - p} \right| \right],$$

where $m_L$ is the plasma mass of the left-handed particle. An identical formula can be written for right-handed fermion component with $m_L$ replaced by $m_R$. 
Eqs. (4, 5) apply to all particles with masses \( \ll T \), that is, to all the SM fermions in the unbroken phase. For Dirac fermions with non-zero bare mass \( m \), the situation is more complicated because the left- and right-handed components couple. The Majorana propagators for massive neutralinos would be even more complicated [15]. In the present paper we shall focus on the limit \( T \gg m \), which we believe cover most of the cases of interest. In present context the only Dirac particles with non-zero bare mass are the top quarks and the charginos, whose right- and left-handed plasma masses are equal. The bare mass of the charged gauginos \( \tilde{W}^\pm \) is the \( SU(2)_L \) soft supersymmetry breaking gaugino mass \( M_2 \) and the bare mass \( \mu \) of charged Higgsinos \( \tilde{H}^\pm \) emerges from the bilinear term \( \mu \tilde{H}_1 \tilde{H}_2 \) in the superpotential.

We assume that the only particles light enough to be in equilibrium with the thermal bath at temperature \( T \) are the SM particles, the right-handed stop \( \tilde{t}_R \), the charginos \( \tilde{W}^\pm \) and \( \tilde{H}^\pm \), the neutralinos \( \tilde{B}, \tilde{W}_3 \) and \( \tilde{H}_1^0, \tilde{H}_2^0 \), and the neutral and charged scalar fields in the two Higgs doublets. This amounts to assuming that all the other supersymmetric particles, e.g. left-handed stops and gluinos, are much heavier than \( T \). This choice is motivated by considerations about the strength of the phase transition [7, 6, 8] and the mechanism for baryon number generation [3, 4, 5] in the MSSM. For instance, a strongly first order electroweak phase transition can be achieved in the presence of a top squark not much heavier than the top quark [7, 6]. In order to naturally suppress its contribution to the parameter \( \Delta \rho \) and hence preserve a good agreement with the precision measurements at LEP, it should be mainly right-handed. This can be achieved if the soft supersymmetry breaking mass \( m_Q \) of \( \tilde{t}_L \) is much larger than \( M_Z \) so that \( \tilde{t}_L \) is decoupled from the thermal bath. It is important to keep in mind that this mass hierarchy between the left- and right-handed stops may be relaxed if the strength of the phase transition is enhanced by light boson degrees of freedom other than the right-handed stops themselves. This happens, for instance, if the MSSM content is increased by adding a standard model gauge singlet.

The relevant high \( T \) plasma masses are given in the Table 1. We have separated the SM contribution \( m_{SM} \) from the MSSM contribution \( m_{MSSM} \), so that the plasma mass is actually the sum of the two. We also denote by \( g_1 = 0.247 \), \( g_2 = 0.640 \) and \( g_3 = 1.243 \) respectively the values of \( U(1)_Y \), \( SU(2)_L \) and \( SU(3)_c \) couplings at \( T \simeq M_W \), and we take the (top-quark) Yukawa coupling to be \( h_t = 1 \). Other couplings and interactions are not included because of their relative smallness. The relevant fermionic one-loop diagrams are those, where the right-handed stop couples to one of the pairs \( \tilde{H}_2^0 t_L, \tilde{H}_1^0 b_L \) or \( \tilde{B}_L t_R \). The possible bosonic one-loop diagrams consist of those including \( \tilde{t}_R \) and \( U(1) \) gauge boson \( B \) or \( \tilde{t}_R \) and gluon.

Let us first consider the fermionic contribution to the one-loop imaginary part of the right-handed stop self-energy. The right-handed stop can absorb a right-handed Higgsino hole (left-handed quark hole) to produce a left-handed quark particle (right-
defined the appropriate plasma mass. The solutions are depicted in Fig. 1 for the case of
where the upper and lower signs refer to particles and holes, respectively, and
\( \omega \)
by
Here \((R, h)\) and \((L, p)\) refer to right-handed holes and left-handed particles, respectively, and \(\omega_{L,p}(k)\) and \(\omega_{R,h}\) are solutions to the fermion dispersion relations, given by

\[
\omega_{p,h} = -a(\omega_{p,h}, k)\omega_{p,h} - b(\omega_{p,h}, k) \pm (1 + a(\omega_{p,h}, k))k
\]

\[
= \pm \left[ k + \frac{m^2}{k} \left[ 1 + \frac{1}{2} \left( \pm 1 - \frac{\omega_{p,h}}{k} \right) \ln \left( \frac{\omega_{p,h} + k}{\omega_{p,h} - k} \right) \right] \right]
\]

where the upper and lower signs refer to particles and holes, respectively, and \(m\) is the appropriate plasma mass. The solutions are depicted in Fig. 1 for the case of \(t\) and \(\tilde{H}\). The coupling factor \(e^2 = h^2\) for the loops involving \(t_L\) and \(b_L\) (at high \(T\) their contributions are equal), and \(e^2 = g^2/2\) for the case of \(\tilde{B}t_R\)-loop. In Eq. (6) we have defined

\[
n^\pm(x) \equiv \frac{1}{e^x \pm 1}.
\]
Taking $\tilde{t}_R$ to be at rest, the momentum $k$ is determined by the energy conservation condition
\[
M_{\tilde{t}_R} + \omega_{R,h} = \omega_{L,p} \quad (M_{\tilde{t}_R} + \omega_{L,h} = \omega_{R,p}).
\] (9)

At high $T$ the direct decay channel contributes an imaginary part \[\text{Im } \Sigma_{\text{dec}} = 4e^2 k^2 (\omega_{L,p}^2 - k^2)(\omega_{R,p}^2 - k^2) \frac{16m_L^2 m_R^2}{16m_L^2 m_R^2} \left[1 - n^+(\beta \omega_{R,p}) - n^+(\beta \omega_{L,p})\right]
\] (10)

where $k$ is determined by
\[
M_{\tilde{t}_R} = \omega_{R,a} + \omega_{L,a}
\] (11)

with $a = p, h$. At high temperatures the plasma mass of $\tilde{t}_R$ is $M_{\tilde{t}_R}(T) = 1.020T$, whereas $M_{\tilde{H}_2}(T) = 0.329T$, $M_{\tilde{B}}(T) = 0.116T$, $M_{b_L}(T) = M_{t_L}(T) = 0.649T$, and $M_{t_R}(T) = 0.625T$ so that decay is clearly kinematically possible to both $\tilde{B}t_R$ and $\tilde{H}_2 t_L$ or $\tilde{H}_2 b_L$. The full fermionic high $T$ contribution to Im $\Sigma_R$ is thus given by the sum $2 \text{Im } \Sigma_{\text{abs}}(\tilde{H}_2 t_L) + \text{Im } \Sigma_{\text{abs}}(\tilde{B}t_R) + 2 \text{Im } \Sigma_{\text{decay}}(\tilde{H}_2 t_L) + \text{Im } \Sigma_{\text{decay}}(\tilde{B}t_R)$, and it can only be computed numerically by solving $k$ from Eqs. (9, 11) for a fixed $M_{\tilde{t}_R}(T)$, using the fermionic dispersion relations for holes and particles as given by Eq. (7). (The factor 2 comes about because the equality of $\tilde{H}_2 t_L$ and $\tilde{H}_2 b_L$ contributions.)

The high $T$ fermionic contributions to $\Gamma_R$ are depicted in Fig. 1 plotted against $M_{\tilde{t}_R}(T)$, which is related to the zero temperature stop mass by $M_{\tilde{t}_R}(T) = M_{\tilde{t}_R}(0) + \left(\frac{1}{3}g_1^2 + \frac{1}{3}h_1^2 + \frac{36}{108}g_2^2\right)T^2$. Note that for the fermionic contributions the high $T$ approximation does not concern $\tilde{t}_R$; the results apply equally to the case where $\tilde{t}_R$ is no longer relativistic ($T \ll M_{\tilde{t}_R}$) as well as close to the phase transition, where cancellation between the zero temperature and high temperature mass terms is possible so that $T \gg M_{\tilde{t}_R}$. Thus for a fixed stop mass, from the figures one can read the various contributions to $\Gamma_R$ at a fixed temperature.
The one-loop, high \( T \) gauge boson contribution to \( \text{Im} \, \Sigma_R \) can be cast in the form

\[
\text{Im} \, \Sigma_{gb} = -\frac{e^2}{\pi} \frac{\omega_g M_{tR}^2 k^4}{\Pi(M_g^2 - \Pi)\omega_i} \left[ n^-(\beta \omega_i) - n^-(\beta \omega_g) \right],
\]

(12)

where \( \Pi(\omega_g, k) \) is the longitudinal gauge boson self energy, given by

\[
\Pi(\omega_g, k) = 3M_g^2 \left( 1 - \frac{\omega_g^2}{k^2} \right) \left[ 1 - \frac{\omega_g}{2k} \ln \left( \frac{\omega_g + k}{\omega_g - k} \right) \right].
\]

(13)

\( \omega_g \) is solved implicitly through \( \omega_g^2 = k^2 + \Pi \), and \( \omega_i^2 = k^2 + M_{tR}^2 \). The coupling factor \( e^2 \) are \( 4g_3^2/3 \) and \( 4g_1^2/9 \) for gluons and \( B \), respectively. The energy conservation condition in this case is

\[
M_{tR} + \sqrt{k^2 + M_{tR}^2} = \omega_g .
\]

(14)

Physically Eq. (12) corresponds to absorption; for kinematic reasons decay is not possible for gauge loops. Here we have implicitly assumed that \( T \gg |M_{tR}(0)| \) so that \( M_{tR}(T) \leq 1.02T \). (If we assume \( M_{tR}^2(0) \) to be negative, it is possible that \( M_{tR}(T) \ll T \).) This is necessary because of the implementation of the high \( T \) approximation in the loop, which now also concerns \( \bar{t}_R \) as it is circulating in the loop. Hence the gauge boson contribution Eq. (12) is not correct for non-relativistic \( \bar{t}_R \), but it is nevertheless valid even if there is a cancellation between the stop bare mass and plasma mass terms.

The gauge boson contributions to \( \Gamma_R \) at high \( T \), plotted against \( M_{tR}(T) \), are shown in Fig. 3. One sees that gauge bosons do not contribute to \( \Gamma_R \) when \( \bar{t}_R \) is in full equilibrium, but may be important if \( M_{tR}(T) \lesssim T \). In Fig. 3 we also show the total high \( T \) thermalization rate \( \Gamma_R \). At high temperature, far from the the critical point, one finds that at one loop \( \Gamma_R \simeq 8 \times 10^{-4}T \).
Figure 3: (a) The gauge boson contributions to $\Gamma_R$ at high $T$; (b) the total $\Gamma_R$. Both the decay rate and the mass are given in the units of temperature in a logarithmic scale.

The total one-loop result presented in Fig. 3 has a range in which $\Gamma_R$ is vanishingly small. This reflects only the kinematics of absorption and decay and will disappear at higher loops. For instance, at two loops the cuts through various diagrams give rise in addition to absorption and decay (involving more than two particles in the initial or final state) to processes which correspond to ordinary scattering. An example of the two-loop contributions are the so-called sunrise-diagrams appearing both from gluon coupling and quartic Yukawa coupling (the electroweak gauge boson can be neglected as discussed earlier). Neglecting the mass differences between different loop particles, the imaginary parts can be estimated as in \cite{12}. We obtain $\Gamma_{R,h} \simeq h_4^2(T^2/128\pi M_{\tilde{t}_R}(T)) = 0.0024T^2/(M_{\tilde{t}_R}(T))$ and $\Gamma_{R,gluon} \simeq 7g_4^2T^2/(576\pi M_{\tilde{t}_R}(T)) = 0.0092(T^2/M_{\tilde{t}_R}(T))$. There appear also t-channel processes from gluon interactions. These may be also important whenever the right-handed stop is light. By making a rough approximation of the t-channel gluon processes in the spirit of ref. \cite{12}, we conclude that the t-channel contribution is comparable to the sunrise-diagrams; we estimate that $\Gamma_{R,scatt} \simeq 3 \times 10^{-3}T$ when $M_R(T) \sim T$. Thus two-loop diagrams are expected to dominate the thermalization rate when $M_{\tilde{t}_R}(T) \lesssim 0.1T$, or in the range $0.5T \lesssim M_{\tilde{t}_R}(T) \lesssim T$, with $\Gamma_R \simeq 10^{-3}T^2/M_{\tilde{t}_R}(T)$. In the range $0.1T \lesssim M_{\tilde{t}_R}(T) \lesssim 0.5T$ absorption of gauge bosons dominates the thermalization rate, and it can be as high as $0.1T$.

Note that right-handed stop plasma masses with $M_{\tilde{t}_R}(T) \ll T$ can occur only if $M_{\tilde{t}_R}^2(0)$ is negative. Even though present experimental bounds on the stop masses do not exclude such a possibility, one should keep in mind that $|M_{\tilde{t}_R}(0)|$ is bounded from above from considerations about color breaking at zero temperature \cite{6}. If we restrict ourselves to electroweak baryogenesis and therefore assume that the critical temperature is $T_c = \mathcal{O}(100)$ GeV, it is easy to show that $M_{\tilde{t}_R}(T_c)$ must be larger than
about $0.55 T_c$ for masses of the lightest Higgs boson around 80 GeV. (For lighter Higgs boson the bound for $M_R$ is even larger.)

Let us briefly discuss the implications of our findings. As we already mentioned in the introduction, the precise knowledge of the thermalization rate of the supersymmetric particles is a key ingredient for the computation of the final baryon asymmetry. Sizeable decay rates of the particles propagating in the plasma destroy the quantum interference out of which the the CP-violating sources are built up and therefore reduce the baryon asymmetry. Small decay rates, on the other side, are relevant when the particles reflecting off the advancing bubble wall have comparable masses and resonance effects show up. In such a case, the thermalization rates provide the natural width of these resonances and as the present calculation demonstrates, in supersymmetric theories these depend in a complicated way on the particles involved and their plasma masses.

Even though it is presently believed that the right-handed stops do not play a leading role in generating the baryon asymmetry, it is important to emphasize that this is only true in the context of the MSSM. There the phase transition is made strong by the infrared effects in the right-handed stop sector and the baryon number is mainly generated by charginos and neutralinos. Should the source of the strength of the phase transition reside somewhere else, i.e. in Standard Model gauge singlets, the role of right-handed stops will be comparable to the one played by charginos and higgsinos and the knowledge of the thermalization rate $\Gamma_R$ is important to obtain a precise estimate of the final baryon asymmetry.

If the stop is non-relativistic with $M_{\tilde{t}_R}(T) \gtrsim T$, thermalization is dictated by the one-loop thermal decay rate which can be larger than $T$. Thus in any case the thermalization of $\tilde{t}_R$ is rather fast, as can be seen in Fig. 3. Since the zero temperature limit on $|M_{\tilde{t}_R}(0)|$ from color breaking suggests that $M_{\tilde{t}_R}(T) \gtrsim 0.55 T$, we may conclude that during baryogenesis the thermalization rate of a relativistic right-handed stop is dominated by two-loop effects (i.e. scattering) with $\Gamma_R \simeq 10^{-3} T$. This is so in particular because the absorption channels close at $M_{\tilde{t}_R}(T) \sim 0.51 T$ before decay channels open at $M_{\tilde{t}_R}(T) \sim 0.66 T$. This means that the processes of quantum interference, necessary build up the axial stop number, may be damped by the incoherent nature of the plasma if $\tilde{t}_R$ is non-relativistic at baryogenesis. Even in the favorite case in which the left-handed and right-handed stops have comparable masses, the resonance effects will be washed out by the large right-handed stop thermalization rate.

We hope to present our results about the thermalization rate of charginos and neutralinos soon.
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