Electroweak Baryogenesis

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ABSTRACT: We present a summary of the physical framework necessary for electroweak baryogenesis. We discuss in detail the strong constraints on the physical parameters of the models from the preservation of the baryon asymmetry of the Universe for the Standard Model and the Minimal Supersymmetric Standard Model. For the Constrained MSSM we discuss the region in which also cosmological constraints from dark matter are simultaneously satisfied. We briefly discuss alternative scenarios which can increase the available parameter space and in particular we comment on models with low reheating temperatures.

1. Introduction

A compelling and consistent theory which can explain the observed baryon asymmetry of the Universe (BAU) is one of the most challenging theoretical aspects of the interplay between particle physics and cosmology. Many mechanisms for the production of the baryonic asymmetry have been discussed for different periods of the evolution of the early Universe which include GUT-baryogenesis, leptogenesis, etc. For reviews on this subject see for instance [1]. The electroweak scale is the last instance in the evolution of the Universe in which the baryon asymmetry could have been produced. Moreover, the scenario of electroweak baryogenesis places constraints on the physical parameters of specific models which are testable at current and future accelerators.

In 1967, Sakharov [2] enunciated the necessary ingredients for the production of the baryon asymmetry, which are: baryon number violation, C and CP violation, and a departure from thermal equilibrium. The Standard Model contains all necessary physical aspects, and thus was considered that solely within this framework baryogenesis could be explained. As far as the required CP-violation, present from the CKM matrix in the Standard Model, we refer the reader to the literature [3, 4].

Let us quickly summarize the current understanding related to the other two requisites for baryogenesis. Baryon number violation occurs in the Standard Model through anomalous processes. At low temperatures this anomalous baryon number violation only proceeds via a tunnelling process which is exponentially suppressed, at a rate $\Gamma \sim e^{(-\frac{4\pi}{\alpha_W})}$. However, at finite temperature these interactions are in equilibrium above the electroweak scale with a rate propotional to $\Gamma \sim \alpha_W^5 T^4$ and may erase any previously produced B-asymmetry [5][1].

At finite temperature T the rate Γ_s per unit time and unit volume for fluctuations between neighboring minima with different baryon number is [6]

$$\Gamma_s \sim 10^5\, T^4 \left(\frac{\alpha_W}{4\pi}\right)^4 \kappa \frac{\zeta^7}{B^7} e^{-\zeta}, \tag{1.1}$$

where we have used $\zeta(T) = E_s(T)/T$ and $E_s(T) = [2m_W(T)/\alpha_W]\, B(\lambda/g^2)$ is the sphaleron energy, $m_W(T) = \frac{1}{2} g\langle\phi(T)\rangle$, $B \simeq 1.9$ is a function which depends weakly on the gauge and the Higgs quartic couplings g and λ, $\alpha_W = g^2/4\pi = 0.033$.

[1]The loophole is that sphaleron transitions conserve $B - L$, so a net $B - L$ generated previously through interactions in a specific model will not be erased.

Thus, below the electroweak phase transition the sphaleron transition rate changes as the $SU(2)$ gauge field acquires a mass $\Gamma_s/T^3 \sim e^{-\frac{4\pi\phi}{gT}} \equiv e^{-\frac{E_{sph}}{T}}$.

In 1985, Kuzmin, Rubakov and Shaposhnikov[7] suggested that if the electroweak phase transition was first order it provided a natural way for the Universe to depart from equilibrium. Eventually bubbles of the broken phase nucleate and grow until they fill the Universe. Local departure of thermal equilibrium takes places near the walls of the expanding bubbles and if asymmetries of some local charges are produced which then diffuse into the unbroken phase where baryon number violation is active this converts the asymmetries into a baryon asymmetry. Finally, the baryon number flows into the broken phase. Now it is necessary that the sphaleron transitions are turned off after the phase transition (in the broken phase) so as to *not* wash out the asymmetry produced during the transition. This is a crucial test for any model of electroweak baryogenesis. The condition which guarantees the non-erasure of the produced asymmetry occurs when the transition rate of the sphaleron interactions is small (out-of-equilibrium) compared to the Hubble expansion rate. In a radiation dominated Universe the Hubble parameter is $H \sim \frac{T^2}{M_{Pl}}$.

Requiring $\Gamma_s/T^3 \lesssim H$ at the bubble nucleation temperature T_b leads to the condition on $\zeta(T_b)$. It can be shown that to avoid erasing any generated baryon asymmetry via sphaleron interactions we require $\zeta(T_c) = \frac{E_{sph}}{T_c} \gtrsim 45$. Given that E_{sph} is defined in terms of ϕ, this in turn implies that we must have a sufficiently strong phase transition such that $\frac{\phi}{T_c} \gtrsim 1$. Where ϕ is the order parameter of the transition.

One of the main analytic tools for studying the thermodynamics of the phase transition is the finite temperature effective potential. At one-loop, the non-interacting bosonic or fermionic particles, whose mass depends on the background field ϕ, contribute to the finite temperature effective potential.

The contributions from bosons in the high-temperature expansion is

$$\frac{m_i^2(\phi)T^2}{24} - \frac{m_i^3(\phi)T}{12\pi} + \frac{m_i^4(\phi)}{64\pi^2}(\ln\frac{T^2}{\mu^2} + C_i), \quad (1.2)$$

and from fermions

$$\frac{m_i^2(\phi)T^2}{48} - \frac{m_i^4(\phi)}{64\pi^2}(\ln\frac{T^2}{\mu^2} + C_i). \quad (1.3)$$

We can easily see that the terms of the form $m_i^2(\phi)T^2 \sim \phi^2T^2$, are responsible for symmetry restoration at very high temperature, and that the terms $-m_i^3(\phi)T \sim -\phi^3T$ produce a barrier in the potential as the temperature decreases and thus allows a first order phase transition to occur. In many extensions of the Standard Model, finite temperature computations are complicated by a hierarchy of mass scales, indeed many mass scales can appear for which the high-temperature expansion is not adequate. In some cases, however a low-temperature expansion can be employed to a very good approximation. Recently it was shown how a perturbative calculation to two-loops can be computed without any temperature expansions [14].

In general, we can parametrize the effective potential at one-loop in the following way,

$$V_{eff}^{1-loop} = (\gamma T^2 - m^2)\phi^2 - ET\phi^3 + \lambda\phi^4. \quad (1.4)$$

Thus, the phase transition will occur when

$$\frac{\phi(T_c)}{T_c} \sim \frac{E}{\lambda} \sim \frac{g^3}{\lambda} \sim g\frac{m_W^2}{m_h^2}, \quad (1.5)$$

where we have used the Standard Model leading contributions to E. At zero-temperature, λ is related to the mass of the Higgs particle. At tree level in the SM, $m_h^2 = 4\lambda\phi^2$, so small λ corresponds to small Higgs mass. When quantum corrections are included, the relationship becomes less simple, but the erasure condition still translates into an upper bound on the Higgs mass. Thus, requiring $\frac{\phi(T_c)}{T_c} \gtrsim 1$ shows that we need a small value of m_h, which is not possible due to current experimental bounds.

At two-loops there are new contributions which could modify our previous conclusions, as the leading correction is of the form $\Delta V_{eff}^{2-loop} \sim g^2T^2m_i^2(\phi)\ln(\frac{m_i(\phi)}{T}) \sim \gamma BT^2\phi^2\ln(\frac{\phi}{T})$. Thus we can now parametrize

$$\frac{\phi(T_c)}{T_c} \sim \frac{E}{\lambda} + \left(\frac{E^2}{\lambda^2} + \frac{2\gamma B}{\lambda}\right)^{1/2}. \qquad (1.6)$$

For a given model, the parameters γ, E, λ and B are calculable functions of the zero temperature coupling constants. This shows that two-loop effects can strengthen the phase transition if B is large, and that we need a large E and a small quartic coupling λ. This does not happen in the Standard Model.

To summarize, although the ingredients for the production of the baryon asymmetry exist in the Standard Model as far as having a mechanism which produces a departure of thermal equilibrium and baryon number violating interactions, we have shown using purely perturbative arguments that the simpler constraint on the preservation of the produced asymmetry cannot be fulfilled.

A drawback to the perturbative evaluation of the effective potential is that it has infra-red divergences as the Higgs vacuum expectation value ϕ approaches zero. These are due to bosonic modes whose only mass in the perturbative calculation is proportional to ϕ. Resummation is then employed to deal with these divergences. The net effect in the case of the Standard Model is to reduce the contribution from the gauge bosons to the cubic term, thus weakening the strength of the phase transition. An alternative way to compute the ratio (1.6) is using Monte Carlo 3-d lattice simulations via dimensional reduction [11]. The important point is that this will include also all relevant non-perturbative effects. In dimensional reduction an effective 3-d theory is constructed perturbatively integrating out all the massive modes. The characteristics of the phase transition can then be determined by simulating on the lattice the remaining 3-dimensional theory of massless bosonic modes. One of the strong points of the lattice simulations is that they can constrain a whole class of models which are described by the same effective 3-d theory containing a single light scalar at the phase transition[12].

2. Electroweak Phase Transition in the MSSM

We have shown that the relevant contributions to the strength of the first order phase transition arise from bosonic particles. Extensions of the Standard Model with new bosons could modify the strength of the phase transition. In addition, we would also like to have new CP-violating interactions.

The Minimal Supersymmetric Standard Model (MSSM) provides an interesting scenario in which there are many new bosonic particles. In particular, the stops couple strongly to the Higgs field producing the largest modification to the strength of the phase transition. Analytically, to determine the contribution from the stops we use the finite temperature stop mass matrix which is given by

$$M_{stops}(\phi, T) = \begin{pmatrix} M_{11} & M_{12} \\ M_{21} & M_{22} \end{pmatrix}, \qquad (2.1)$$

where

$$\begin{aligned} M_{11} &= m_Q^2 + \frac{1}{2}h_t^2\sin^2\beta\phi^2 + a_1T^2, \\ M_{12} &= M_{21} = -\frac{1}{\sqrt{2}}h_t\sin\beta\phi(A+\mu\cot\beta), \\ M_{22} &= m_U^2 + \frac{1}{2}h_t^2\sin^2\beta\phi^2 + a_2T^2, \qquad (2.2) \end{aligned}$$

where a_1 and a_2 are combinations of the relevant couplings constants in the theory.

In the previous expressions we have neglected the contributions from the gauge couplings, this will suffice to illustrate which are the most relevant contributions from the stop sector. Note that for the case of zero stop-mixing, the contribution from the right-handed stop to the cubic term of the effective potential is

$$(m_{\tilde{t}_R}^2)^{3/2} = \Big(m_U^2 + +\frac{1}{2}h_t^2\sin^2\beta\phi^2 + [\frac{4}{9}g_s^2$$

$$+ \frac{1}{6}(1+\sin^2\beta)h_t^2]T^2\Big)^{3/2}, \qquad (2.3)$$

so, if $m_U^2 \lesssim 0$ it can cancel the finite temperature contributions such that $m_{\tilde{t}_R}^3 \sim \delta E \phi^3$. This is called the light-stop mechanism [13]. The crucial effect of the parameters of the phase transition is that now at one-loop,

$$\frac{\phi(T_c)}{T_c} \sim \frac{(E+\delta E)}{\lambda} \qquad (2.4)$$

where $\delta E \sim h_t^3 \sin^3\beta(1 - \frac{\tilde{A}_t}{m_Q^2})^{3/2}$, and $\tilde{A}_t = A + \mu \cot\beta$. Something analogous would occur for the left-handed stop $m_{\tilde{t}_L}$, which could also strengthen the phase transition. However, this is not possible for two reasons: the zero temperature one-loop corrections to the Higgs mass from the stop sector are large enough to satisfy the experimental constraints only when one of the stop fields is very heavy; secondly, precision electroweak constraints require $m_{\tilde{t}_L}$ to be heavy. It is also worth mentioning here that the two-loop QCD effects from the stop sector are sizeable and they strengthen the phase transition [16, 18, 15].

The effects of non-zero doublet-singlet mixing reduce the strength of the phase transition, as shown in eq. (2.4), as it reduces the contribution of the light stop to δE, it also makes the Higgs heavier and increases the value of the critical temperature. These undesirable effects of the trilinears on the survival of the baryon asymmetry can be partially compensated by decreasing the value of m_U^2.

Simulations have also been done for the case where there is a colored $SU(2)$ singlet among the light/massless modes at $\phi = 0$ [17]. This corresponds to the light right-handed stop scenario that could make the phase transition strong enough in the MSSM. The area of MSSM parameter space where the electroweak phase transition is strong enough was found in the lattice analysis to be larger than the area found in perturbation theory[17]. This means that calculating $\phi(T_c)/T_c$ from the perturbative effective potential is conservative, so this is why we employ the 2-loop effective potential. We construct the effective potential via dimensional reduction. The effective 3D theory constructed with dimensional reduction reproduces the perturbative 4D effective potential results. The 3-D theory naturally incorporates the effects of resummation and some higher order corrections. We use the results given in ref. [15] for the two-loop finite-temperature effective potential of the MSSM with a light stop. It assumes that the b-quark Yukawa coupling is small (this restricts the value of $\tan\beta \lesssim 12$), and is calculated in the limit where all the supersymmetric particles are heavy ($\sim$ TeV) except for the stops.

Allowing $m_U^2 \lesssim 0$ is constrained at zero temperature from the limits on the stop mass; for the current limits it is still possible for the light stop scenario to be available. We will return to this below when we discuss experimental constraints. Another constraint on the negative values of m_U^2 appears as colour (and charge)-breaking minima can develop in the stop direction both at zero and finite-temperature. In fact, the strongest constraint arises at finite-temperature, when a two-stage phase transition can appear in certain regions of parameter space. The fact that the stop is light allows the possibility of tunneling into a color and charge breaking minimum [18, 13, 15, 19], from which the Universe would subsequently undergo a transition to the $SU(2)$ broken minimum. The analysis of ref. [19] shows that the second phase transition may not take place, thus giving stronger constraints on the allowed parameter space. This gives a lower bound on the stop mass for every set of values of the mixing parameter and $\tan\beta$. This lower bound is larger than the direct experimental search limit on $m_{\tilde{t}_2}$, as shown in fig. 1.

The conclusions of the two-loop perturbative analysis for the MSSM [18, 13, 15] are that the sphaleron transitions are suppressed when the stop and Higgs bosons are light enough. In the MSSM at two loops, the light singlet stop contribution to E is sufficient to satisfy $\frac{\phi}{T_c} \gtrsim 1$. However, the maximum possible values of $m_{\tilde{t}_2}$ (the light stop mass), and m_h for which the sphaleron transitions are suppressed depend also on the trilinear terms and on $\tan\beta$, as discussed above.

In fig. 1 we display the full allowed region

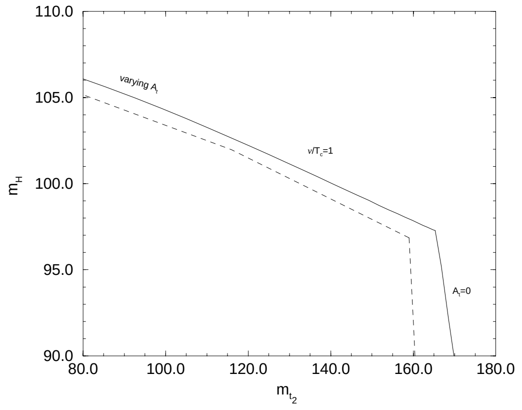

Figure 1: *Available region in parameter space in the $m_h - m_{\tilde{t}_2}$ plane for $m_Q = 300$ GeV, varying $\tilde{A}_t$ and $\tan\beta$ defined by the contour of $\frac{\phi}{T} = 1$. The dashed line is defined when the critical temperatures in the $\phi-$ and $\chi-$ directions are equal for the same variations of $\tilde{A}_t$ and $\tan\beta$.*

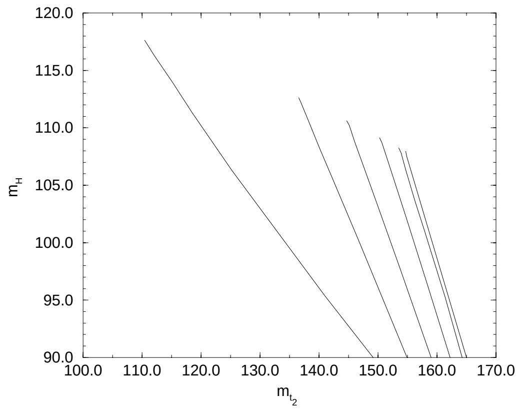

Figure 2: *Contours of $\frac{\phi}{T_c} = 1$ in the $m_h - m_{\tilde{t}_2}$ plane for $m_Q = 1$ TeV, varying $\tilde{A}_t$. The contour which corresponds to no-mixing corresponds to the first curve on the right. For increasing values of the mixing parameter the curves shift to the left on the plane.*

in parameter space coming from the requirement of having a strong enough phase transition. The solid vertical line is determined by varying $\tan\beta$ for $\tilde{A}_t = 0$ and the solid diagonal line corresponds to the variation of $\tilde{A}_t$ for $\tan\beta = 12$. In the region to the left and below these solid lines sphaleron transitions are suppressed for values of $2 \lesssim \tan\beta \lesssim 12$ and $0 \lesssim \tilde{A}_t \lesssim 280$ GeV. The dashed line shown defines when the critical temperatures in the ϕ- and χ-directions are equal,. For the same variations of $\tan\beta$ and $\tilde{A}_t$, according to the analysis of ref. [19] the region below this dashed line would be excluded. For this value of m_Q the maximum allowed masses are $m_h \lesssim 105$GeV and $m_{\tilde{t}_2} \lesssim 170$GeV.

In figure 2 we plot in the $m_h - m_{\tilde{t}_2}$ plane contours of $\frac{\phi}{T} = 1$ for $m_Q = 1$ TeV. The parameters that are varied are $\tan\beta, m_{\tilde{t}_R}$ and $\tilde{A}_t$. The effect of increasing $\tan\beta$ increases the value of m_h, thus weakening the phase transition and so a smaller value of m_{t_R} is necessary to maintain $\frac{\phi}{T} = 1$. It is clear that the least stringent bound on m_{t_2} arises when $\tilde{A}_t = 0$, but the largest possible value of the Higgs mass m_h is determined by the largest allowed value of the mixing parameter. The most important modification in the analysis provided by the light-stop mechanism at 2-loops is that there is no longer an upper bound on m_h, from requiring the phase transition to be sufficiently strong, it is really a zero temperature effect which gives the largest possible Higgs mass, and as long as the stop is light enough there is a sufficiently strong first order phase transition which preserves the produced baryon asymmetry.

The current experimental bounds on the Higgs masss, are for a Standard Model-like Higgs $m_h \gtrsim 106$GeV; we emphasize that the results presented here are valid for large values of the pseudoscalar mass m_A, and a maximum value of $\tan\beta = 12$, as corrections from the bottom- sbottom sector have not been included. The experimental bound for the Higgs for smaller values of m_A is about 10 GeV smaller for large $\tan\beta$. From experiments the stop mass is required to be from the Tevatron $m_{\tilde{t}} \gtrsim 90$ GeV for $m_{\chi^o} \leq 50$GeV. As can be seen from the results in figure 2 there is only a small window in which electroweak baryogenesis in the MSSM is possible. Before we discuss alternative scenarios, which could increase the allowed region in parameter space, let us for the moment make a small digression.

Dark Matter Abundance and Electroweak baryogenesis in the CMSSM

As we have shown above, the preservation of the baryon asymmetry is still possible in a small window of parameter space in the context of the MSSM. Another salient feature of this supersymmetric model is that it has a dark matter candidate, namely the lightest supersymmetric particle (LSP). If R-parity is conserved the LSP is stable and its mass density is large enough, $0.1 < \Omega h^2 < 0.3$, to be cosmologically interesting for a mostly gaugino-like neutralino [21]. In ref. [20] it was determined if there exists a region in which combined experimental, theoretical and cosmological constraints could be satisfied. The combined constraints that were used included: LEP bounds on particle masses, the branching ratio for $b \to s\gamma$, precision electroweak constraints to the ρ-parameter, electroweak phase transition bounds and the relic abundance constraint.

In the constrained MSSM (CMSSM), there are only four parameters $m_o, m_{1/2}, A_o, \tan\beta$ and the $sign(\mu)$ which define parameter space. Over much of parameter space of the CMSSM the LSP is a bino. For a gaugino-type neutralino annihilation occurs via sfermion exchange into fermion pairs (leptons). As m_o increases the slepton becomes more massive thus decreasing the neutralino annihilation rate. Thus, the neutralino relic abundance increases, until an upper bound for the scalar mass parameter m_o is obtained. A lot of work has been done constraining SUSY models using combined bounds from accelerator experiments and the cosmological relic abundance[21]. This defines a specific allowed region in the m_o- $-m_{1/2}$ plane. We are particularly interested in the overlap with the area of parameter space where the singlet stop soft SUSY-breaking mass is negative for the preservation of the baryon asymmetry. In order to do this the low energy parameters are obtained by running the couplings and gaugino masses with 2-loop renormalization group equations, and the 1-loop renormalization group equation for other masses. In fact, the light stop mass constraint from the electroweak phase transition is the hardest to satisfy. The one-loop renormalization group equation is given by

$$\frac{dm_U^2}{dt} = \frac{16}{3}\frac{\alpha_s}{4\pi}M_3^2 - 2\frac{h_t^2}{16\pi^2}(m_Q^2 + m_U^2 + m_h^2 + A_t^2) \tag{2.5}$$

where M_3 is the gluino mass and $t = \ln\frac{m_{GUT}^2}{Q^2}$. It can be shown that $m_U^2 \sim 0$ is obtained for $|A_o| \sim 11|m_{1/2}|$. The important point to notice is that for fixed $m_o, m_{1/2}$ we can independently vary A_o to produce a light stop. There is a similar term proportional to the tau-Yukawa couplings and for some ranges of medium to large $\tan\beta$ the light stau mass can be driven negative. Another important constraint from baryogenesis is the requirement of relatively small stop mixing, thus we choose the sign of μ to be opposite to A so as to cancel the off-diagonal elements of the stop mass matrix.

In figure 3 we plot the combined set of experimental and theoretical constraints in the $m_{1/2}-$ $-A_o$ plane. This figure is for values of $\tan\beta = 12$ and $m_o = 145$ GeV. The presence of a quasi-fixed point for A_t at low $\tan\beta$ tends to drive A_t to positive values in the vicinity of $2m_{1/2}$, which tends to be much larger than $|\mu|\cot\beta$. Smaller $|A_t|$ can be generated at moderate $\tan\beta$ by taking A_0 large and negative. Accordingly, to obtain essentially zero-mixing in the stop sector we consider negative A_0 and positive μ in fig. 3. The light shaded region is excluded because it contains either a tachyonic stop or stau. The dashed contour combines the experimental bounds from chargino, stop and Higgs searches with the condition that the LSP be the neutralino, to avoid an excess of charged dark matter. In Fig. 3, the vertical left side of the dashed contour is due to the chargino bound, the diagonal piece which parallels the light shaded region is due to the stop bound, and the horizontal piece is the line $m_\chi = m_{\tilde{\tau}_R}$ The experimental Higgs constraint, form LEP189 run, does not provide a useful bound for this value of $\tan\beta$. The solid line gives the BAU constraint: above the solid line, sphaleron processes wash out the baryon asymmetry, while below the solid line the baryon asymmetry is preserved. In the figures, the BAU boundary is typically set by the condition that the stop mass be sufficiently light. Lastly, the dark hatching marks the region where the neutralino relic density violates the cosmological upper bound $\Omega_{\tilde{\chi}} h^2 \leq 0.3$. The region which is allowed by all of the experimental and cosmological constraints is then highlighted by diagonal shading.

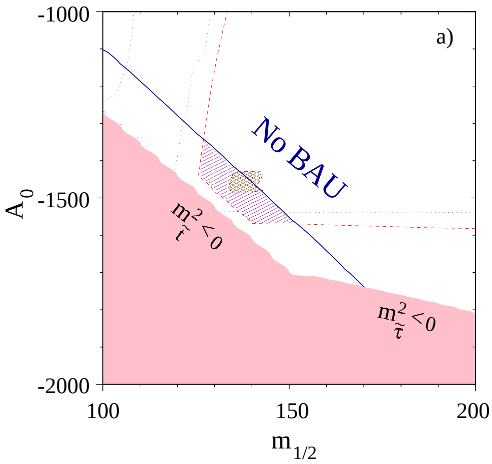

Figure 3: *The region of mSUGRA parameter space satisfying the BAU, experimental and cosmological constraints described in the text, for* $\tan\beta = 12, m_0 = 145$ *GeV and* $\mu > 0$*, is marked by diagonal shading. Above the solid line, sphaleron processes wipe out the baryon asymmetry. The light shaded region contains tachyons. The dark hatched region has* $\Omega_{\tilde{\chi}} h^2 > 0.3$*. The dotted lines are contours of constant* $\Omega_{\tilde{\chi}} h^2 = 0.1$.

For a more complete discussion of this scenario see ref. [20]. We have presented these results because although at present LEP2 bounds seem to exclude the combined region of allowed parameter space, a caveat may exist which can open up room for this region once more. The caveat is related to the inclusion of CP-violation in the MSSM which we will discuss a bit further in the next section.

3. Alternative Scenarios

Our results show that the allowed regions for baryogenesis in the (C)MSSM are quite small given the strong experimental constraints. Here we would like to mention some of the possible ways of enlarging parameter space.

- A straightforward procedure is to redo the electroweak phase transition analysis with values of the pseudoscalar mass $m_A \sim 100$ GeV with the light stop mechanism, as the LEP bound is weaker. Also we can explore larger values of $\tan\beta$ by including the contributions from the bottom-sbottom sector.
- The explicit inclusion of CP-violation into the analysis of the production and decay rate of the Higgs boson seem to hint that a relatively light Higgs boson could have escaped detection at LEP2[22, 23]. This clearly would enlarge parameter space by allowing a smaller value of m_h, which would also imply that a heavier stop would also be allowed.
- We could slightly modify the CMSSM by splitting m_Q (or m_U) from the other boson masses related to m_o. The effects of CP-violation mentioned above are also relevant here.
- One can consider other models such as the Two Higgs Doublet model, the Next to Minimal Supersymmetric Standard Model, or the MSSM with R-parity violation in which either the Higgs mass bounds are modified or we can have different couplings (not fixed by supersymmetry) which contribute to the ratio of $\frac{E}{\lambda}$.
- Alternative cosmologies can modify the constraints on the parameters related to the production and preservation of the baryon asymmetry of the Universe. We will discuss a particular scenario next.

Electroweak Baryogenesis in Low T_{reheat} *Models*

It is well known that the flatness and horizon problems of the standard big bang cosmology are solved if during the early Universe the energy density was dominated by some form of vacuum energy. The observed large scale density and temperature fluctuations can also be generated within this framework.

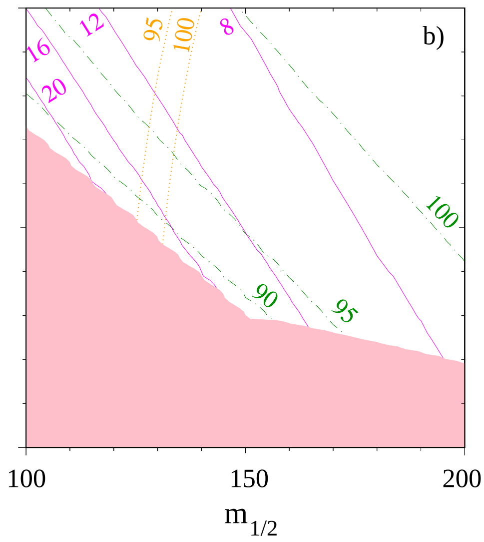

Figure 4: *Contours of constant Higgs mass (dot-dashed), chargino mass (dotted) and $\delta\rho \times 10^4$ (solid) are displayed.*

The inflationary stage can be parametrized by the evolution of some scalar field ϕ, the inflaton, which is initially displaced from the minimum of its potential. Inflation ends when the potential energy associated with the inflaton field becomes smaller that the kinetic energy of the field. The process which converts a low-entropy cold Universe dominated by the energy of the coherent motion of the ϕ-field into a high entropy hot Universe dominated by radiation is called reheating.

When the Universe becomes radiation dominated the energy density scales like $\rho \sim T^4_{reheat}$, this defines T_{reheat}. Recall that for a radiation dominated Universe the Hubble parameter is $H \sim \frac{T^2}{M_{Pl}}$. A common assumption is that the post-inflationary Universe contained a plasma in thermal equilibrium at a temperature $T_{reheat} \gg T_{EW}$ $\sim 10^2$ GeV. However, experimentally, all we know is that $T_{reheat} \geq$ *few* MeV, which is required from nucleosynthesis. Low reheat scenarios are particularly appealing as they avoid the overproduction of dangerous relics[24, 25]. A particularly interesting question is what happens to Sakharov's conditions of baryon number violation and disequilibrium in models with low reheat temperatures. In fact, baryon number violating interactions at low temperature are strongly constrained by laboratory bounds, also disequilibrium is hard to come by as the Universe is expanding very slowly, so it is close to equilibrium.

This would seem to rule out the possibility of baryogenesis within the context of these models. Nevertheless, let us inspect more carefully what are the implications for electroweak baryogenesis. For other baryogenesis models see refs. [27, 28]. The crucial point is to realize that reheating is **not** an instantaneous process and that the maximum temperature T_m during reheating can be much larger that T_{reheat}[26]. For a radiation dominated Universe the temperature scales as $T \sim a^{-1}$; during reheating the temperature is given by $T = T_m f(a)$, where $T_m = \frac{T_{reheat}}{\alpha_\phi^{1/2}}$, and α_ϕ is defined by the decay rate of the inflaton field, $\Gamma_\phi = \alpha_\phi M_\phi$. The function $f(a) = \kappa(a^{-3/2} - a^{-4})^{1/2}$, i such that $f(a_o) = 1$ has its maximum value for $a_o = (\frac{8}{3})^{2/5}$. The important thing to note is that the Hubble parameter is now modified to

$$H \sim H_{RD}\frac{T^2}{T^2_{reheat}}, \tag{3.1}$$

thus for smaller values of T_{reheat} the Hubble parameter increases with respect to its value in a radiation dominated Universe. Let us now suppose that the reheating temperature $T_{reheat} \ll T_c$. We have seen that during the reheating process the thermal bath may reach temperatures which are much larger than electroweak phase transition critical temperature $T_c \sim T_{EW}$. This means that the Universo can undergo a phase transition in which the electroweak symmetry is broken although the Universe is not yet radiation dominated. The main difference is that the phase transition occurs in a matter dominated Universe with an expansion rate given by eq. (3.1). This implies that electroweak baryogenesis may occur even when $T_{reheat} \ll T_c$. This is a nontrivial result and is the crucial fact that we want to point out. As discussed above, the generation of the baryon asymmetry is mediated by sphaleron transitions in the unbroken phase, which are in equilibrium at temperatures $T \lesssim \left(\alpha_W^4 M_p T^2_{reheat}\right)^{1/3} \sim 10^4 \left(T_{reheat}/1 \text{ GeV}\right)^{2/3}$ GeV. We can now show that the requirement that the sphalerons are out-of-equilibrium in the broken phase is easier to fulfill if $T_{reheat} \ll T_c$ than in the standard cosmology. By once again comparing the rate Γ_s/T^3, given in eq. (1.1), with the new expansion rate of eq. (3.1) we find

$$\begin{aligned} \zeta(T_b) &\gtrsim 7\log\zeta(T_b) + 9\log 10 + \log\kappa \\ &+ 2\log(T_{reheat}/T_b) \\ &\gtrsim \zeta(T_b)^{S.C.} + 2\log(T_{reheat}/T_b)\,. \end{aligned} \quad (3.2)$$

This inequality is the standard one [1, 5], with one crucial difference: the presence of the last term. The presence of this term implies that, if the reheating temperature is much smaller than T_c (or equivalently the Universe is expanding very quickly) sphalerons are out-of-equilibrium in the broken phase easier than in the standard cosmology. Numerically, we can see that if we take $\zeta(T_b) \simeq 1.2\zeta(T_c)$ [12], then for $\kappa = 10^{-1}$ and $T_{reheat} \sim 1(10)$ GeV, we obtain that $\zeta(T_c) \gtrsim 28(33)$, which translates into

$$\frac{\langle\phi(T_c)\rangle}{T_c} \gtrsim 0.77(0.92). \quad (3.3)$$

This bound shows how the constraint on the order parameter changes compared to the standard result, $\langle\phi(T_c)\rangle/T_c \gtrsim 1$, obtained for the same value of κ. This finding clearly enlarges the available region in parameter space where the sphaleron bound is satisfied and relaxes the upper bound on the stop mass in the MSSM and on the Higgs mass in other extensions of the SM. The implication for the SM is that although current LEP bounds on the Higgs mass still rule out electroweak baryogenesis, for small values of the Higgs mass the phase transition is now strong enough for sphaleron transitions to be suppressed. From the lattice results of Ref. [12] we can determine that Eq. (3.3) implies that the electroweak phase transition would be strong enough for baryogenesis for $m_h \lesssim 50$ GeV. More interesting, for the MSSM in the region of allowed Higgs masses the new bound could increase the upper bound on the stop mass by about 10 GeV to $m_{\tilde{t}} \lesssim 180$ GeV for all other parameters fixed.

We have seen that preserving a baryon asymmetry is easier if $T_{reheat} \ll T_c$, however one must also check the effects on the production of the asymmetry in this context. The continous decays of the scalar field ϕ add entropy into the thermal bath when the temperature decreases from T_c to T_{reheat}. If we take B_c to be the baryon asymmetry to entropy density ratio n_B/s generated at the electroweak phase transition, one finds that the final baryon asymmetry is now

$$\frac{n_B}{s} \sim B_c \left(\frac{T_{reheat}}{T_c}\right)^5. \quad (3.4)$$

This means that, for $T_{reheat} \sim 10$ GeV, the mechanism of baryogenesis at the electroweak scale has to be more efficient by a factor $\sim 10^5$ than in the standard case. Although this is can be difficult it is not impossible to achieve.

4. Conclusions

We have presented an overview of the status of electroweak baryogenesis in the Standard Model and the (C)MSSM. We have shown that the available region in parameter space for the MSSM is highly constrained by experimental results. We have also presented alternative scenarios which can increase the allowed region of parameter space. In particular, we have discussed the scenario of non-standard cosmology with low reheat temperatures in which baryogenesis could still occur and the implications for electroweak baryogenesis.

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References

[1] For recent reviews, see W. Buchmuller, S. Fredenhagen, hep-ph/0001098; A. Riotto, hep-ph/9807454 ; A. Riotto and M. Trodden, Annual Reviews of Nuclear and Particle Science **49**, 35 (1999). V.A. Rubakov, M.E. Shaposhnikov, *Usp. Fiz.*

Nauk **166** 493 (1996).A. Cohen, D.B. Kaplan, and A.E. Nelson. *Annu. Rev. Nucl. Part. Sci.*, 43:27, 1993.
M. Quiros. *Helv. Phys. Acta*, 67:451, 1994.

[2] A.D. Sakharov, *JETP Lett.* **5** (1967) 24.

[3] G.Farrar and M. Shaposhnikov, *Phys. Rev. Lett.* 2833 (1993); (E) **71** 210 1993.

[4] M.B. Gavela, P. Hernandez, J. Orloff, O. Pène and C. Quimbay, *Mod. Phys. Lett.* **9** 795 1994.

[5] A.I. Bochkarev, S.V. Kuzmin and M.E. Shaposhnikov, Phys. Rev. **D43**, 369 (1991).

[6] L. Carson, Xu Li, L. McLerran and R.-T. Wang, Phys. Rev. **D42**, 2127 (1990).

[7] V.A. Kuzmin, V.A. Rubakov and M.E. Shaposhnikov, Phys. Lett. **B155**, 36 (1985).

[8] M. Losada. *Phys. Rev.*, **D6** 2893, 1997.
M. Laine. *Nucl. Phys.*, **B481** 43, 1996.

[9] see *e.g.*, P. Arnold, O. Espinosa, *Phys. Rev.* **D47** (1993) 3546.

[10] M. Dine, R. G. Leigh, P. Huet, A. Linde, D. Linde, *Phys.Rev.* **D46** (1992) 550.

[11] K. Kajantie, M. Laine, K. Rummukainen, M. Shaposhnikov, *Nucl.Phys.* **B458** (1996) 90; hep-ph/9508379.

[12] K. Kajantie, M. Laine, K. Rummukainen and M. Shaposhnikov, Nucl. Phys. **B466**, 189 (1996).

[13] M. Carena, M. Quiros and C.E.M. Wagner, Nucl. Phys. **B524**, 3 (1998);

[14] M. Laine and M. Losada, Nucl. Phys. **B582**, 277 (2000).

[15] M. Losada, Nucl. Phys. **B537**, 3 (1999); M. Losada, Nucl Phys **B569**, 125 (2000).

[16] M. Carena, M. Quiros and C.E.M. Wagner, Phys. Lett. **B380** 81 (1996); D. Delepine, J.M. Gerard, R. Gonzalez Felipe and J. Weyers, Phys. Lett. **B386**, 183 (1996); J.R. Espinosa, Nucl. Phys. **B475**, 273 (1996); B. de Carlos and J.R. Espinosa, Nucl. Phys. **B503**, 24 (1997);

[17] M. Laine and K. Rummukainen, Nucl. Phys. **B535**, 423 (1998).

[18] D. Bodeker, P. John, M. Laine, and M.G. Schmidt. *Nucl. Phys.*, B497:387, 1997.

[19] J. Cline, G.D. Moore, and G. Servant. Technical Report (hep-ph/9902220), 1999.

[20] S. Davidson, T. Falk and M. Losada, Phys. Lett. **B463**, 214 (1999).

[21] J. Ellis, T. Falk, K.A. Olive and M. Schmitt, Phys. Lett. **B388** (1996) 97; J. Ellis, T. Falk, K.A. Olive and M. Schmitt, Phys. Lett. **B413** (1997) 355; J. Ellis, T. Falk, G. Ganis, K.A. Olive and M. Schmitt, Phys. Rev. **D58** (1998) 095002.

[22] M. Carena, J. Ellis, A. Pilaftsis and C. Wagner, hep-ph/0003180.

[23] G. Kane and L.T. Wang, hep-ph/0003198.

[24] Kallosh, L. Kofman, A. Linde and A. Van Proeyen, hep-th/9907124; G. F. Giudice, I. Tkachev and A. Riotto, JHEP **9908**, 009 (1999); G. F. Giudice, A. Riotto and I. Tkachev, JHEP **9911**, 036 (1999); G. Felder, L. Kofman and A. Linde, hep-ph/9909508.

[25] N. Arkani-Hamed, S. Dimopoulos and G. Dvali, *Phys.Rev.* **D59** 086004,1999,

[26] D.J.H. Chung, E.W. Kolb and A. Riotto, Phys. Rev. **D60** (1999) 063504, hep-ph/9809453.

[27] S. Davidson, M. Losada and A. Riotto, Phys. Rev. Lett **84** 4284 (2000).

[28] G. Giudice, E. Kolb and A. Riotto, hep-ph/0005123.