

Status of electroweak phase transition and baryogenesis

J M CLINE

Department of Physics, McGill University, 3600 University Street, Montréal, Québec H3A 2T8,
Canada

Email: jcline@physics.mcgill.ca

Abstract. I review recent progress on the electroweak phase transition and baryogenesis, focusing on the minimal supersymmetric Standard Model as the source of new physics.

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1. Cosmological phase transitions – electroweak

It is possible that the universe has undergone a number of phase transitions, as illustrated in table 1. In most cases, it is difficult to find a signature of such a transition which survives to the present day. One important class of exceptions is when (meta)stable topological defects like cosmic strings are formed; I will not deal with this important topic in the present talk. Instead I will focus on the other main possibility of interest, the case of a first order transition.

A memorable example of the effect a first order phase transition could have was proposed by Witten [1] in 1984. He noted that when bubbles of the chiral symmetry broken phase form (where $\langle \bar{q}_L q_R \rangle \neq 0$), baryons tend to pile up on the bubble walls. Neutrons diffuse quickly into the bubble interiors, but protons diffuse more slowly, and the spatial separation of isospin was found to have an observable effect on helium production in primordial nucleosynthesis. Unfortunately, lattice studies have shown that the QCD chiral phase transition is a smooth crossover when quarks have masses, so that in fact there is no bubble formation.

If the QCD transition was not first order, what about the next lowest energy example on our list, the electroweak phase transition (EWPT)? In this case the bubbles would contain regions of nonvanishing Higgs field VEV, $\langle H \rangle \neq 0$. If CP is violated on the bubble wall, there can be a pile-up of chiral charge, which biases the anomalous sphaleron interactions of the Standard Model to produce baryons. We could thus explain the baryon asymmetry of the universe. Unfortunately, the lattice gauge theorists have again spoiled our fun by finding that for Higgs masses $m_H > 70\text{--}80$ GeV, the transition is a smooth crossover. (The latest limit from LEP is $m_H > 106$ GeV.) However, it is much easier to change

Table 1. Some phase transitions which may have taken place during cosmological history.

Energy scale	Transition	Order parameter
M_{Planck}	spacetime foam \rightarrow classical spacetime	$\langle g_{\mu\nu} \rangle$
M_{GUT}	GUT symmetry \rightarrow Standard Model	$\langle \text{GUT Higgs} \rangle$
Intermediate scale	Inflation	$\langle \text{Inflation} \rangle$
M_{SUSY}	SUSY breaking	$\langle (\text{gaugino})^2 \rangle?$
100 GeV	Electroweak: $\text{SU}(2) \times \text{U}(1) \rightarrow \text{U}(1)_{\text{em}}$	$\langle \text{SM Higgs} \rangle$
100 MeV	QCD chiral phase transition	$\langle \bar{q}_L q_R \rangle$

this negative conclusion by adding new physics (like supersymmetry) to the electroweak theory than it is for QCD. This will be the subject of the rest of this talk.

2. How strong is the EWPT?

One of the main tools for studying the strength of the EWPT analytically is the finite temperature effective potential, defined by

$$e^{-\beta \int d^3x V_{\text{eff}}[\Phi]} = \int \prod_i \mathcal{D}\phi_i e^{-\int_0^\beta d\tau \int d^3x S[\Phi, \phi_i]/\hbar}. \quad (1)$$

It is a path integral over fluctuating fields ϕ_i , around a constant background field Φ , in our case the Higgs field. The fields are in imaginary time with periodic boundary conditions (for bosons; antiperiodic for fermions) between $\tau = 0$ and $\tau = \beta = 1/T$. V_{eff} can be computed in perturbation theory, represented by Feynman diagrams like \bigcirc at one loop, and $\bigcirc\bigcirc$ or \ominus at two loops. The one loop term is the effect of a noninteracting boson or fermion gas, of which the particle masses depend on the background field Φ ,

$$V_{1\text{-loop}} = T \sum_i \mp \int \frac{d^3p}{(2\pi)^3} \ln \left(1 \pm e^{-\sqrt{p^2 + m_i^2(\Phi)}/T} \right), \quad (2)$$

where the upper (lower) sign is for fermions (bosons) in the loop. This can be approximated in a high-temperature expansion as

$$\begin{aligned} & \sum_i \frac{m_i^2(\Phi)T^2}{48} (\times 2 \text{ for bosons}) \sim \Phi^2 T^2 \\ & - \sum_i \frac{m_i^3(\Phi)T}{12\pi} (\text{bosons only}) \sim -\Phi^3 T \\ & \mp \frac{m_i^4(\Phi)}{64\pi^2} \left(\ln \frac{T^2}{\mu^2} + C_i \right) + O(m_i^6/T^2). \end{aligned} \quad (3)$$

The $\Phi^2 T^2$ term is responsible for symmetry restoration at high T , while the $-\Phi^3 T$ term gives a barrier or bump (\sim) in the potential, which is responsible for the first order transition, if it occurs.

Unfortunately, perturbation theory can be unreliable, especially near the phase transition [2]. If one starts with an arbitrarily complicated diagram contributing to $V_{n\text{-loop}}$ and adds an extra W boson propagator, the ‘cost’ of the new loop is a multiplicative factor which parametrically has the form

$$\epsilon = g^2 T \int \frac{d^3 p}{(2\pi)^3} (p^2 + m_W^2(\Phi))^{-2} \sim g^2 \frac{T}{m_W(\Phi)} \sim g \frac{T}{\Phi}. \quad (4)$$

The relevant value of Φ is Φ_c , the VEV in the broken phase at the critical temperature where $V_{\text{eff}}(0) = V_{\text{eff}}(\Phi_c)$. If we parametrize $V_{\text{eff}} \cong AT^2\Phi^2 - BT\Phi^3 + \lambda\Phi^4$ then $\Phi_c = 2BT/\lambda \sim g^3 T/\lambda \sim (m_w^2/m_H^2)gT$. Therefore the perturbative parameter ϵ goes like $\lambda/g^2 \sim m_w^2/m_H^2$, and perturbation theory breaks down for heavy Higgs bosons, $m_H > m_W$.

There are several ways to combat the breakdown of perturbation theory at finite temperature. The most brute force method is to use lattice gauge simulations for the full 4-D theory [3]. Somewhat easier is to use dimensional reduction—integrating out the heavy Matsubara modes (the Fourier modes of the compactified imaginary time direction, with masses $m_n \sim n\pi T$) to get an effective 3-D theory [4,5]. The 3-D theory can then be studied on the lattice [6], much more easily than the 4-D theory. A third method is to compute V_{eff} to higher order in perturbation theory [7–10]. This sounds unjustified, since perturbation theory was supposed to be breaking down, but experience shows that in fact it works rather well in the cases of interest—where the transition is strongly first order.

As mentioned above, the lattice studies have established that, although there is a line of first order phase transitions in the $T - m_H$ plane at small m_H , it comes to an end (at a point where the transition is 2nd order) around $m_H = 75$ GeV. For larger m_H there is no clear distinction between the unbroken and broken phases of the electroweak theory. For example, massive W bosons in the broken phase cannot be distinguished from massive composite objects ($H^\dagger \overleftrightarrow{D}_\mu H$) in the symmetric, confining phase.

To get a first order transition for realistic Higgs masses, we need to add new physics which couples significantly to the Higgs boson. Let’s see how supersymmetry can do this.

3. Adding supersymmetry

Recall that at one loop, it is the cubic term, $-Tm^3(H)/12\pi$ that gives bump in the potential (\sim) hence a first order transition. In the minimal supersymmetric Standard Model (MSSM), we have two Higgs doublets, and the top squarks (\tilde{t}_L and \tilde{t}_R) couple strongly to the second one, H_2 . Ignoring small terms involving the weak gauge coupling g , the $\tilde{t}_L - \tilde{t}_R$ mass matrix has the form

$$M_{\tilde{t}_L, \tilde{t}_R}^2 = \begin{pmatrix} m_Q^2 + y^2 H_2^2 & y(A_t H_2 - \mu H_1) \\ y(A_t H_2 - \mu H_1) & m_U^2 + y^2 H_2^2 \end{pmatrix}. \quad (5)$$

We need at least one of the SUSY breaking soft masses, m_U^2 or m_Q^2 , to be small so that there will be an eigenvalue $m(H)$ which is really cubic, $\sim H^3$, not $(m_0^2 + y^2 H^2)^{3/2}$, since

the latter form does not give a true bump in the potential [8]. On the other hand, one of m_U^2 or m_Q^2 should be large so that the stop radiative correction to the Higgs mass can be big enough to satisfy the experimental constraint:

$$m_H^2 \sim m_Z^2 + O \left[\frac{m_t^4}{v^2} \ln \left(\frac{m_{\tilde{t}_L} m_{\tilde{t}_R}}{m_t^2} \right) \right] > (96 \text{ GeV})^2. \quad (6)$$

(The limit on m_H in the MSSM is about 10 GeV weaker than in the SM.) The precision electroweak ρ parameter (a.k.a. ϵ_1) dictates that the \tilde{t}_L -like squark should be the heavy one, hence the \tilde{t}_R -like squark is light.

It turns out that two-loop effects are crucial for getting a strong enough phase transition [7–10]. Diagrams like \ominus , with a gluon or Higgs boson as one of the internal lines, contribute a term

$$\begin{aligned} \Delta V_{\text{eff}} &= -(8g_s^2 - 3y^2 \sin^2 \beta) T^2 m_{\tilde{t}_R}^2 (H) \ln \left(\frac{m_{\tilde{t}_R}(H)}{T} \right) \\ &\sim -CT^2 H^2 \ln \left(\frac{H}{T} \right) \end{aligned} \quad (7)$$

to V_{eff} , whose form is unlike any generated at one loop. One can show that such a term shifts the critical value of the Higgs VEV according to

$$\frac{\langle H \rangle}{T} \sim \frac{B}{\lambda} + \sqrt{\left(\frac{B}{\lambda} \right)^2 + \frac{2C}{\lambda}}, \quad (8)$$

where B is the part coming from the cubic term. $\langle H \rangle/T$ is the relevant measure of the strength of the transition, as we will discuss below.

In ref. [10] we have performed a Monte Carlo search of the MSSM parameter space to find those which give a strong enough phase transition for electroweak baryogenesis. We can summarize the resulting constraints on the squark masses as follows:

$$120 \text{ GeV} \lesssim m_{\tilde{t}_R} \lesssim m_{\text{top}} \quad (9)$$

$$m_{\tilde{t}_L} > 265 \text{ GeV} \times e^{(m_H - 95 \text{ GeV})/9.2}. \quad (10)$$

It is often said that electroweak baryogenesis puts an upper limit on the Higgs boson mass, but this is not correct in the MSSM, where the light right stop is doing most of the job of making the transition first order. However, as we can see from eq. (10), it is true that the left stop quickly becomes unnaturally heavy as m_H is increased. Thus we are pushed to a rather strange corner of parameter space, where \tilde{t}_R is extremely light, and \tilde{t}_L extremely heavy. If m_H turns out to be much heavier than its current experimental limit, there must be some other new physics accounting for its mass.

More promising is the prediction (9) for the stop mass. Although Tevatron searches for the decay $\tilde{t} \rightarrow \tilde{\chi}_1^0 c$ are only beginning to probe the region of interest (figure 1a) [11], in Run II this decay will probe up to the top mass if $m_{\tilde{\chi}_1^0}$ is in the right range (dark region of figure 1b), and the decays $\tilde{t} \rightarrow b \tilde{\chi}_1^+$ or $\tilde{t} \rightarrow bW \tilde{\chi}_1^0$ may be more revealing, if the chargino is light (lighter regions) [12].

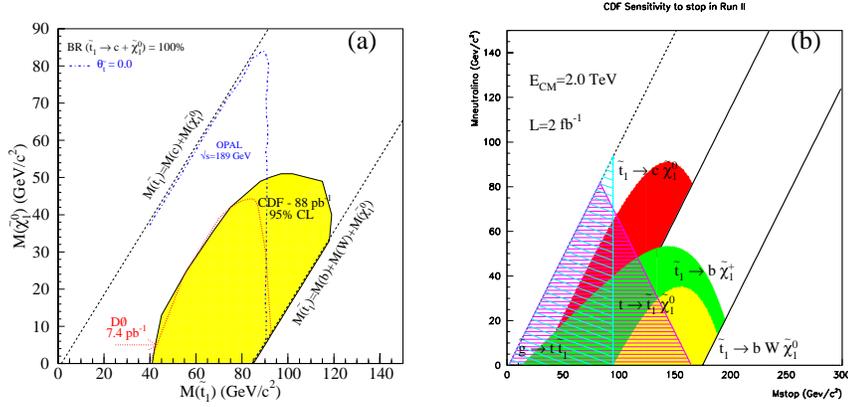


Figure 1. (a) Present Tevatron limits [11] and (b) sensitivity in Run II for the stop mass [12] in the (neutralino mass)–(stop mass) plane.

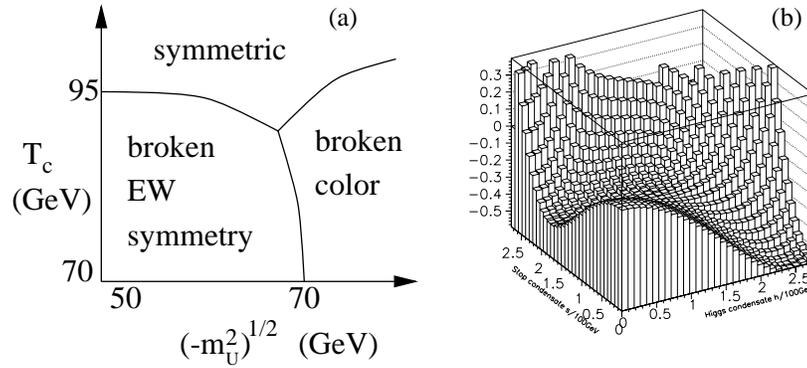


Figure 2. (a) Phase diagram and (b) effective potential, $V_{\text{eff}}(\tilde{t}_R, H)$, for color breaking due to stop condensation.

An interesting consequence of such a light stop is that its bare (mass)² must be negative. In the absence of left–right mixing, $m_{\tilde{t}_R}^2 = m_U^2 + m_t^2$, so $m_{\tilde{t}_R} < m_t$ implies $m_U^2 < 0$. This can cause an instability toward condensation of the \tilde{t}_R field in the early universe, when $\langle H \rangle$ is still zero, which would break $SU(3)_{\text{color}}$ [9]. Indeed, the lattice study of ref. [13] finds a phase diagram similar to figure 2a, which shows that color-breaking occurs when $(-m_U^2)^{1/2}$ exceeds 60–70 GeV, depending on the critical temperature T_c . In ref. [14] we have constructed the two loop effective potential $V_{\text{eff}}(\tilde{t}_R, H)$ for stop and Higgs fields, and studied the possibility that the universe might temporarily enter the color-breaking phase, before finally tunneling to the EW-breaking vacuum which we inhabit now. We find that because of the potential barrier separating the two minima, shown in figure 2b, the rate of tunneling is so small that if the universe ever enters the color-breaking minimum, it stays

there forever. This conclusion could however change in the presence of R -parity violating interactions like $y_{332} A_R^a \bar{l}_R^b \bar{s}_R^c \epsilon_{abc}$ which would lower the barrier.

4. Baryon asymmetry of the universe and baryogenesis

I will now review some basics about baryogenesis and describe some recent developments, before giving the latest details on electroweak baryogenesis in the MSSM. It is very unlikely that we live in a baryon–antibaryon symmetric universe since there is no evidence of antigalaxies colliding with galaxies, which would give an intense source of gamma rays. If there are regions consisting of antimatter outside of our Hubble volume, it is difficult to imagine how separation from ordinary matter could occur on such a large distance scale. Big bang nucleosynthesis tells us that

$$2 \lesssim \eta_{10} \equiv \frac{n_B - n_{\bar{B}}}{n_\gamma} \times 10^{-10} \lesssim 3. \quad (11)$$

Yet in the early universe the most natural initial condition is equal numbers of baryons and antibaryons, since there are many possible B -violating interactions that could be in thermal equilibrium at high temperatures, for example a dimension 9 operator like $(udd)^2 / \Lambda^5$, which would cause neutron–antineutron oscillations, or the interactions with heavy X gauge bosons in grand unified theories. In fact we need not look so far afield, since sphalerons violate baryon number within the Standard Model itself!

It is well known that Sakharov’s three conditions must be fulfilled to generate a baryon asymmetry: (1) baryon number violation; (2) C and CP violation; (3) loss of thermal equilibrium for the B -violating interactions. The first and second conditions are present within the SM, but (2) is too weak for baryogenesis, and (3) is not fulfilled at all. We have already seen how the MSSM can increase $\langle H \rangle / T$; this is what is needed to make the sphalerons go out of equilibrium inside the bubbles that form during the EWPT. The MSSM can also cure the problem of getting strong enough CP violation, as we will discuss in the next section. Here, I will only mention some of the other proposals for baryogenesis which are new or currently popular.

Baryogenesis via leptogenesis [15] is one of the most plausible alternatives to electroweak baryogenesis. In analogy to GUT baryogenesis, heavy sterile neutrinos decay out of equilibrium, producing a lepton asymmetry, which is converted by sphalerons into the baryon asymmetry. The predictions can be related to neutrino masses in a GUT framework like $SO(10)$ [16].

The Affleck–Dine mechanism [17] has long been one of the most efficient baryogenesis mechanisms. It uses the fact that SUSY scalar potentials from D -terms can often have flat directions which generate a huge baryon number when the flatness is lifted and the field evolves by spiraling in the complex plane. It was recently pointed out that this can be combined with leptogenesis to give a minimal supersymmetric baryogenesis model in which the flat direction is a linear combination of H_2 (the second Higgs doublet) and L_e (the selectron), using physics which is already needed for generating neutrino masses [18].

Large extra dimensions can present a serious challenge to baryogenesis by constraining the reheat temperature after inflation to be very low [19]. The Randall–Sundrum alternative of warped compactification [20] evades this problem; so perhaps do intrinsically ‘braney’ approaches to baryogenesis [21].

Other novel ideas make use of the phase transition in left–right symmetric models [22] and decaying primordial black holes [23].

5. Electroweak baryogenesis in the MSSM

Electroweak baryogenesis in the MSSM is one of the most carefully studied ideas for baryogenesis, owing to its close ties to present-day phenomenology and accelerator searches [24–29]. Its basic mechanism [30] is intuitively easy to understand: particles interact in a CP-violating manner with bubble walls, which form during the first order electroweak phase transition, when the temperature of the universe was near $T = 100$ GeV. This causes a buildup of a left-handed quark density in excess of that of the corresponding antiquarks, and an equal and opposite right-handed asymmetry, so that there is initially no net baryon number. The left-handed quark asymmetry biases anomalous sphaleron interactions, present within the Standard Model, to violate baryon number preferentially to create a net quark density. The resulting baryon asymmetry of the universe (BAU) soon falls inside the interiors of the expanding bubbles, where the sphaleron interactions are shut off (provided that $\langle H \rangle > T$), and thus baryon number is safe from subsequent sphaleron-induced relaxation to zero.

However, the correct way to treat the generation of the chiral quark asymmetry in front of the bubble wall is controversial. In the simplest model, the top quark has a spatially varying complex mass, $m(x) = |m(x)|e^{i\theta(x)}$, which gives rise to CP-violating quantum mechanical reflection of quarks as they pass through the bubble wall. It also induces a CP-violating *classical* force on the quarks. In ref. [31] it was shown that the latter is the more appropriate treatment when the bubble wall thickness l_w is large compared to the inverse temperature, which is the case in the MSSM: $l_w = (10 \pm 4)/T$ [29]. Furthermore it is possible to rigorously derive the way in which the CP violating force influences the particle transport, using the Boltzmann equation; no such complete derivation yet exists in the quantum reflection formalism.

In the MSSM, the top quark mass does not have a CP-violating phase; however the charginos do; their mass matrix has the form

$$\overline{\psi}_R M_\chi \psi_L = (\overline{\tilde{w}}^+, \overline{\tilde{h}}_2^+)_R \begin{pmatrix} m_2 & gH_1(x) \\ gH_2(x) & \mu \end{pmatrix} \begin{pmatrix} \tilde{w}_1^+ \\ \tilde{h}_1^+ \end{pmatrix}_L + \text{h.c.} \quad (12)$$

Since the SUSY parameters μ and m_2 can be complex, the mass eigenstates can have spatially varying complex phases in the wall. There thus arises a classical force which separates the two kinds of Higgsinos, $\tilde{h}_{2,L}^+$ and $\tilde{h}_{1,L}^+$, in front of the wall. This kind of asymmetry is not enough to bias the sphaleron interactions, but scatterings, such as $\tilde{h}_2 \tilde{g} \rightarrow t_L \bar{t}_R$, will partially convert the Higgsino asymmetry into a chiral quark asymmetry, n_{q_L} . Once the latter is determined, it is straightforward to integrate the rate of baryon violation by sphalerons, governed by the equation

$$\frac{dn_B}{dt} = 27 \frac{\Gamma_{\text{sph}}}{T^3} n_{q_L}. \quad (13)$$

The rate of sphaleron interactions per unit volume has been measured by lattice simulations to be $\Gamma_{\text{sph}} = (20 \pm 2)\alpha_w^5 T^4$ [32].

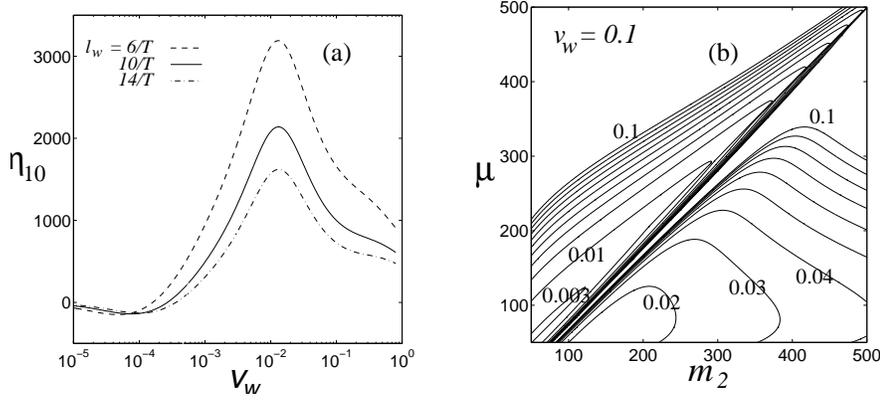


Figure 3. (a) Baryon asymmetry versus wall velocity for maximal CP violation and three different wall thicknesses; and (b) contours of constant CP-violating phase giving the desired baryon asymmetry, in the chargino mass parameter plane.

To determine n_{qL} , we must solve a set of coupled diffusion equations for the various species i of particles in the plasma. They have the form

$$-D_i n_i'' - v_w n_i + \Gamma_{ijk}(n_i - n_j - n_k) = S_i, \quad (14)$$

where D_i is a diffusion coefficient (of order the inverse mean free path), v_w is the bubble wall velocity, Γ_{ijk} is the rate of interactions of the type $i \rightarrow j + k$ (given as an example), and S_i is a CP-violating source term arising from the force on the particles or from quantum reflections. It is possible to show that the source term in the Higgsino diffusion equation is related to the force F by the thermal averages

$$S(x) = -\frac{v_w D}{\langle \vec{v}^2 \rangle} \langle v_x F(x) \rangle', \quad (15)$$

and, by solving the Dirac equation in the WKB approximation, the CP-violating part of the force is related to the complex Higgsino mass $me^{i\theta}$ by $F = (s/2E^2)(m^2\theta)'$, where $s = \pm 1$ is the spin. Moreover the combination $m^2\theta'$ is given by

$$m^2\theta' = \frac{g^2 \text{Im}(m_2\mu)}{2(m_+^2 - m_-^2)} (H_1 H_2' + H_1' H_2), \quad (16)$$

where m_{\pm}^2 are the two eigenstates of the mass matrix in (12). A remarkable feature of this expression is the relative + sign between $H_1 H_2'$ and $H_1' H_2$. Other authors [24,26] had previously obtained a source proportional to $H_1 H_2'$ minus $H_1' H_2$, which is highly suppressed, because H_1/H_2 tends to be constant within the wall, to within a part in 10^2 or 10^3 [33,10]. The origin of the discrepancy is that (for technical reasons) the previous authors considered only the linear combination of Higgsino densities $n_{\bar{h}_1} - n_{\bar{h}_2}$, whereas our classical force is providing a source for $n_{\bar{h}_1} + n_{\bar{h}_2}$ [29].

By solving the diffusion equations and numerically integrating the baryon violation rate equation, we obtained the baryon asymmetry as a function of the model-dependent

parameters of the MSSM [29]. Figure 3a shows how η_{10} varies with the bubble wall velocity v_w in a typical case, for the allowed range of values of the bubble wall thickness, where we assumed maximal CP violation, $\text{Im}(m_2\mu) = |m_2\mu|$. Since typically η_{10} is ~ 1000 , but we only need (2–3), we see that the CP-violating phase need not be maximal, but could be $(2-3)\times 10^{-3}$. This is good news, since the neutron and electric dipole moment searches give constraints which are of this order, unless some kind of fine tuning is invoked. Baryon production peaks at small wall velocities near $v_w = 10^{-2}$. Interestingly, recent estimates of v_w in the MSSM give values which are this small [34]. Figure 3b shows the contours of constant CP phase which yield $\eta_{10} = 3$ in the chargino mass parameter plane.

6. Conclusions

In this talk I have discussed one of the main ‘applications’ of cosmological phase transitions, baryogenesis. Although there are many imaginative ideas for getting the baryon asymmetry, the most mainstream ones are those that are most closely related to phenomenology. Baryogenesis from leptogenesis is appealing because it can potentially make contact with neutrino masses. Electroweak baryogenesis is indirectly testable by searches for Higgs bosons, the top squark, charginos, neutralinos and electric dipole moments. We have seen that it is relatively easy to generate a large enough baryon asymmetry in the electroweak model with the MSSM, but it is not so easy to get a strong enough phase transition to safeguard it against washout by sphalerons: a large hierarchy between the left and right stop masses is needed. It would be nice to have some more robust way of strengthening the transition. One possibility is adding a singlet Higgs field [35]. An interesting proposal is to reheat after inflation to temperatures below the electroweak transition, but rely upon nonequilibrium effects, similar to parametric resonance, to produce sphalerons [36]. Another idea is to modify the expansion rate of the universe instead. It has been noted that extra dimensions could have this effect at sufficiently early times [37].

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