RECENT PROGRESS IN BARYOGENESIS[*]

Antonio Riotto[1] and Mark Trodden[2]

[1]*CERN TH-Division, CH-1211 Geneva 23, Switzerland.*

[2]*Particle Astrophysics Theory Group, Department of Physics, Case Western Reserve University, 10900 Euclid Avenue, Cleveland, OH 44106-7079, USA.*

Abstract

We provide an up to date account of progress in understanding the origin of the observed baryon asymmetry of the universe. While our primary goal is to be current, we have attempted to give a pedagogical introduction to the primary areas of research in this field, giving a detailed description of the different scenarios. The very recent developments in GUT baryogenesis, leptogenesis, electroweak baryogenesis and the Affleck-Dine mechanism are presented. In particular, we focus on specific particle physics implementations, mostly in the context of supersymmetry, which lead to specific testable predictions.

I. INTRODUCTION

The symmetry between particles and antiparticles [1,2], firmly established in collider physics, naturally leads to the question of why the observed universe is composed almost entirely of matter with little or no primordial antimatter.

Outside of particle accelerators, antimatter can be seen in cosmic rays in the form of a few antiprotons, present at a level of around 10^{-4} in comparison with the number of protons (for example see [3]). However, this proportion is consistent with secondary antiproton production through accelerator-like processes, $p + p \rightarrow 3p + \bar{p}$, as the cosmic rays stream towards us. Thus there is no evidence for primordial antimatter in our galaxy. Also, if matter and antimatter galaxies were to coexist in clusters of galaxies, then we would expect there to be a detectable background of γ-radiation from nucleon-antinucleon annihilations within the clusters. This background is not observed and so we conclude that there is negligible antimatter on the scale of clusters (For a review of the evidence for a baryon asymmetry see [4]).

More generally, if large domains of matter and antimatter exist, then annihilations would take place at the interfaces between them. If the typical size of such a domain was small enough, then the energy released by these annihilations would result in a diffuse γ-ray background and a distortion of the cosmic microwave radiation, neither of which is observed. A careful numerical analysis of this problem [5] demonstrates that the universe must consist entirely of either matter or antimatter on all scales up to the Hubble size. It therefore seems that the universe is fundamentally matter-antimatter asymmetric.

While the above considerations put an experimental upper bound on the amount of antimatter in the universe, strict quantitative estimates of the relative abundances of baryonic matter and antimatter may also be obtained from the standard cosmology. The baryon number density does not remain constant during the evolution of the universe, instead scaling like a^{-3}, where a is the cosmological scale factor [6]. It is therefore convenient to define the baryon asymmetry of the universe in terms of the quantity

$$\eta \equiv \frac{n_B}{n_\gamma}, \tag{1.1}$$

where $n_B = n_b - n_{\bar{b}}$ is the difference between the number of baryons and antibaryons per unit volume and $n_\gamma = 2\frac{\zeta(3)}{\pi^2}T^3$ is the photon number density at temperature T. Primordial nucleosynthesis (for a review see [7]) is one of the most powerful predictions of the standard cosmological model. The theory allows accurate predictions of the cosmological abundances of all the light elements, H, ^{3}He, ^{4}He, D, B and ^{7}Li, while requiring only the single input parameter η which has been constant since nucleosynthesis. The range of η consistent with the deuterium and ^{3}He primordial abundances is [6]

$$4(3) \times 10^{-10} \lesssim \eta \lesssim 7(10) \times 10^{-10}, \tag{1.2}$$

where the most conservative bounds are in parentheses. Alternatively we may write the range as

$$0.015(0.011) \lesssim \Omega_B\, h^2 \lesssim 0.026(0.038)\ , \tag{1.3}$$

where Ω_B is the proportion of the critical energy density in baryons, and $0.5 \lesssim h \lesssim 0.9$ parametrizes the present value of the Hubble parameter via $h = H_0/(100\ \mathrm{Km\ Mpc^{-1}\ sec^{-1}})$.

To see that the standard cosmological model cannot explain the observed value of η, suppose that initially we start with $\eta = 0$. We can compute the final number density of nucleons b that are left over after annihilations have frozen out. At temperatures $T \lesssim 1$ GeV the equilibrium abundance of nucleons and antinucleons is [6]

$$\frac{n_b}{n_\gamma} \simeq \frac{n_{\bar{b}}}{n_\gamma} \simeq \left(\frac{m_p}{T}\right)^{3/2} e^{-\frac{m_p}{T}}. \tag{1.4}$$

When the universe cools off, the number of nucleons and antinucleons decreases as long as the annihilation rate $\Gamma_{\mathrm{ann}} \simeq n_b\langle\sigma_A v\rangle$ is larger than the expansion rate of the universe $H \simeq 1.66\, g_*^{1/2}\frac{T^2}{m_p}$. The thermally averaged annihilation cross section $\langle\sigma_A v\rangle$ is of the order of m_π^2, so at $T \simeq 20$ MeV, $\Gamma_{\mathrm{ann}} \simeq H$, and annihilations freeze out, nucleons and antinucleons being so rare that they cannot annihilate any longer. Therefore, from (1.4) we obtain

$$\frac{n_b}{n_\gamma} \simeq \frac{n_{\bar{b}}}{n_\gamma} \simeq 10^{-18}, \tag{1.5}$$

which is much smaller than the value required by nucleosynthesis.

In conclusion, in the standard cosmological model there is no explanation for the value of the ratio (1.2), if we start from $\eta = 0$. An initial asymmetry may be imposed by hand as an initial condition, but this would violate any naturalness principle. Rather, the guiding principle behind *modern* cosmology is to attempt to explain the initial conditions required by the standard cosmology on the basis of quantum field theories of elementary particles in the early universe. In this context, the generation of the observed value of η is referred to as *baryogenesis*.

The goal of this article is to describe recent progress in understanding scenarios for generating the baryon asymmetry of the universe (BAU), as quantified in equation (1.2), within the context of modern cosmology. Earlier reviews of this subject, with a different focus, can be found in Refs. [8–13]. We provide a brief account of the central ideas of baryogenesis, and then review the most recent views on the major implementations. We hope that this provides an up to date summary of the status of the field.

In the next section we shall describe the *Sakharov Criteria* which must be satisfied by any particle physics theory through which a baryon asymmetry is produced. In section (III) we describe how bayogenesis can proceed in certain Grand Unified theories and, in particular, we focus on how new ideas concerning preheating after inflation have modified this scenario. In section (IV) we review electroweak baryogenesis, concentrating on the implications of new results concerning the order of the electroweak phase transition (EWPT), the Higgs mass, and supersymmetric implementations. Section (V) contains an account of Affleck-Dine (AD) type baryogenesis scenarios. After describing the general ideas, we explain how this mechanism can find a natural home in the context of supersymmetric theories. We conclude by discussing the current status of baryogenesis and prospects for future progress.

A note about our conventions. Throughout we use a metric with signature $+2$ and, unless explicitly stated otherwise, we employ units such that $\hbar = c = k = 1$.

II. THE SAKHAROV CRITERIA

As pointed out by Sakharov [14], a small baryon asymmetry η may have been produced in the early universe if three necessary conditions are satisfied: *i)* baryon number (B) violation; *ii)* violation of C (charge conjugation symmetry) and CP (the composition of parity and C) and *iii)* departure from thermal equilibrium. The first condition should be clear since, starting from a baryon symmetric universe with $\eta = 0$, baryon number violation must take place in order to evolve into a universe in which η does not vanish. The second Sakharov criterion is required because, if C and CP are exact symmetries, then one can prove that the total rate for any process which produces an excess of baryons is equal to the rate of the complementary process which produces an excess of antibaryons and so no net baryon number can be created. That is to say that the thermal average of the baryon number operator B, which is odd under both C and CP, is zero unless those discrete symmetries are violated. CP violation is present either if there are complex phases in the lagrangian which cannot be reabsorbed by field redifinitions (explicit breaking) or if some Higgs scalar

3

field acquires a VEV which is not real (spontaneous breaking). We will discuss this in detail shortly.

Finally, to explain the third criterion, one can calculate the equilibrium average of B

$$\langle B \rangle_T = \text{Tr}\,(e^{-\beta H} B) = \text{Tr}\,[(CPT)(CPT)^{-1} e^{-\beta H} B)]$$
$$= \text{Tr}\,(e^{-\beta H}(CPT)^{-1} B(CPT)] = -\text{Tr}\,(e^{-\beta H} B)\,, \qquad (2.1)$$

where we have used that the Hamiltonian H commutes with CPT. Thus $\langle B \rangle_T = 0$ in equilibrium and there is no generation of net baryon number.

Of the three Sakharov conditions, baryon number violation and C and CP violation may be investigated only within a given particle physics model, while the third condition – the departure from thermal equilibrium – may be discussed in a more general way, as we shall see. Let us discuss the Sakharov criteria in more detail.

A. Baryon Number Violation

1. B-violation in Grand Unified Theories

Grand Unified Theories (GUTs) [15] describe the fundamental interactions by means of a unique gauge group G which contains the Standard Model (SM) gauge group $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$. The fundamental idea of GUTs is that at energies higher than a certain energy threshold M_{GUT} the group symmetry is G and that, at lower energies, the symmetry is broken down to the SM gauge symmetry, possibly through a chain of symmetry breakings. The main motivation for this scenario is that, at least in supersymmetric models, the (running) gauge couplings of the SM unify [16–18] at the scale $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV, hinting at the presence of a GUT involving a higher symmetry with a single gauge coupling.

Baryon number violation seems very natural in GUTs. Indeed, a general property of these theories is that the same representation of G may contain both quarks and leptons, and therefore it is possible for scalar and gauge bosons to mediate gauge interactions among fermions having different baryon number. However, this alone is not sufficient to conclude that baryon number is automatically violated in GUTs, since in some circumstances it is possible to assign a baryonic charge to the gauge bosons in such a way that at each boson-fermion-fermion vertex the baryon number is conserved. In the particular case of the gauge group $SU(5)$, it turns out that among all the scalar and gauge bosons which couple only to the fermions of the SM, five of them may give rise to interactions which violate the baryon number. The fermionic content of $SU(5)$ is the same as that of the SM. Fermions are assigned to the reducible representation $\overline{5}_f \oplus 10_f$ as $\overline{5}_f = \{d_L^c, \ell_L\}$ and $10_f = \{Q_L, u_L^c, e_L^c\}$. In addition, there are 24 gauge bosons which belong to the adjoint representation 24_V and which may couple to the fermions through the couplings

$$(g/\sqrt{2})24_V \left[(\overline{5}_f)^\dagger\,(\overline{5}_f) + (10_f)^\dagger\,10_f \right]\,. \qquad (2.2)$$

Because of this structure, the gauge bosons XY with SM gauge numbers $(3, 2, -5/6)$ may undergo baryon number violating decays. While in the gauge sector of $SU(5)$ the structure is uniquely determined by the gauge group, in the Higgs sector the results depend upon the choice of the representation. The Higgs fields which couple to the fermions may be in

the representation 5_H or in the representations 10_H, 15_H, 45_H and 50_H. If we consider the minimal choice 5_H, we obtain

$$h_U \, (10_f)^T \, (10_f) \, 5_H + h_D \, (\overline{5}_f)^T \, (10_f) \, \overline{5}_H \; , \tag{2.3}$$

where $h_{U,D}$ are complex matrices in the flavor space. The representation 5_H contains the Higgs doublet of the SM, $(1,2,1/2)$ and the triplet $H_3 = (3, 1, -1/3)$ which is B-violating [19]. Note that each of these bosons have the same value of the charge combination $B - L$, which means that this symmetry may not be violated in any vertex of $SU(5)$.

However, if we extend the fermionic content of the theory beyond that of minimal $SU(5)$, we allow the presence of more heavy bosons which may violate B and even $B - L$. In GUTs based on $SO(10)$, for instance, there exists a fermion which is a singlet under the SM gauge group, carries lepton number $L = -1$, and is identified with the right-handed antineutrino. However, this choice for the lepton number assignment leads to no new gauge boson which violates $B - L$. As we mentioned, the generation of a baryon asymmetry requires C violation. Since $SO(10)$ is C-symmetric, the C symmetry must be broken before a baryon asymmetry may be created. If $SO(10)$ breaks down to $SU(2)_L \otimes SU(2)_R \otimes SU(4)$ via the nonzero expectation value of a Higgs field in the 54_H representation [20–22], the C symmetry is not broken until $U(1)_{B-L}$ is broken at the scale M_{B-L}. In this case, the right-handed neutrino acquires a Majorana mass $M_N = \mathcal{O}(M_{B-L})$ and its out-of-equilibrium decays may generate a nonvanishing $B - L$ asymmetry [23].

B. B-violation in the Electroweak theory.

It is well-known that the most general Lagrangian invariant under the SM gauge group and containing only color singlet Higgs fields is automatically invariant under global abelian symmetries which may be identified with the baryonic and leptonic symmetries. These, therefore, are accidental symmetries and as a result it is not possible to violate B and L at tree-level or at any order of perturbation theory. Nevertheless, in many cases the perturbative expansion does not describe all the dynamics of the theory and, indeed, in 1976 't Hooft [24] realized that nonperturbative effects (instantons) may give rise to processes which violate the combination $B + L$, but not the orthogonal combination $B - L$. The probability of these processes occurring today is exponentially suppressed and probably irrelevant. However, in more extreme situations – like the primordial universe at very high temperatures [25–28] – baryon and lepton number violating processes may be fast enough to play a significant role in baryogenesis. Let us have a closer look.

At the quantum level, the baryon and the lepton symmetries are anomalous [29,30]

$$\partial_\mu j_B^\mu = \partial_\mu j_L^\mu = n_f \left(\frac{g^2}{32\pi^2} W_{\mu\nu}^a \tilde{W}^{a\mu\nu} - \frac{g'^2}{32\pi^2} F_{\mu\nu} \tilde{F}^{\mu\nu} \right) \; , \tag{2.4}$$

where g and g' are the gauge couplings of $SU(2)_L$ and $U(1)_Y$, respectively, n_f is the number of families and $\tilde{W}^{\mu\nu} = (1/2)\epsilon^{\mu\nu\alpha\beta}W_{\alpha\beta}$ is the dual of the $SU(2)_L$ field strength tensor, with an analogous expression holding for $\tilde{F}$. To understand how the anomaly is closely related to the vacuum structure of the theory, we may compute the change in baryon number from time $t = 0$ to some arbitrary final time $t = t_f$. For transitions between vacua, the average

values of the field strengths are zero at the beginning and the end of the evolution. The change in baryon number may be written as

$$\Delta B = \Delta N_{CS} \equiv n_f[N_{CS}(t_f) - N_{CS}(0)] . \tag{2.5}$$

where the Cherns-Simon number is defined to be

$$N_{CS}(t) \equiv \frac{g^2}{32\pi^2} \int d^3x \, \epsilon^{ijk} \, \mathrm{Tr} \, \left(A_i\partial_j A_k + \frac{2}{3}igA_iA_jA_k \right) . \tag{2.6}$$

Although the Chern-Simons number is not gauge invariant, the change ΔN_{CS} is. Thus, changes in Chern-Simons number result in changes in baryon number which are integral multiples of the number of families n_f. Gauge transformations $U(x)$ which connects two degenerate vacua of the gauge theory may change the Chern-Simons number by an integer n, the winding number. If the system is able to perform a transition from the vacuum $\mathcal{G}_{\mathrm{vac}}^{(n)}$ to the closest one $\mathcal{G}_{\mathrm{vac}}^{(n\pm1)}$, the Chern-Simons number is changed by unity and $\Delta B = \Delta L = n_f$. Each transition creates 9 left-handed quarks (3 color states for each generation) and 3 left-handed leptons (one per generation). However, adjacent vacua of the electroweak theory are separated by a ridge of configurations with energies larger than that of the vacuum. The lowest energy point on this ridge is a saddle point solution to the equations of motion with a single negative eigenvalue, and is referred to as the *sphaleron* [26,27]. The probability of baryon number nonconserving processes at zero temperature has been computed by 't Hooft [24] and is highly suppressed by a factor $\exp(-4\pi/\alpha_W)$, where $\alpha_W = g^2/4\pi$. This factor may be interpreted as the probability of making a transition from one classical vacuum to the closest one by tunneling through an energy barrier of height ~ 10 TeV corresponding to the sphaleron. On the other hand, one might think that fast baryon number violating transitions may be obtained in physical situations which involve a large number of fields. Since the sphaleron may be produced by collective and coherent excitations containing $\sim 1/\alpha_W$ quanta with wavelength of the order of $1/M_W$, one expects that at temperatures $T \gg M_W$, these modes essentially obey statistical mechanics and the transition probability may be computed via classical considerations. The general framework for evaluating the thermal rate of anomalous processes in the electroweak theory was developed in [28]. Analogously to the case of zero temperature and since the transition which violates the baryon number is sustained by the sphaleron configuration, the thermal rate of baryon number violation in the *broken* phase is proportional to $\exp(-S_3/T)$, where S_3 is the three-diemensional action computed along the sphaleron configuration,

$$S_3 = E_{sp}(T) \equiv (M_W(T)/\alpha_W)\mathcal{E} , \tag{2.7}$$

with the dimensionless parameter $\mathcal{E}$ lying in the range $3.1 < \mathcal{E} < 5.4$ depending on the Higgs mass. The prefactor of the thermal rate was computed in [31] as

$$\Gamma_{sp}(T) = \mu \left(\frac{M_W}{\alpha_W T} \right)^3 M_W^4 \exp \left(-\frac{E_{sph}(T)}{T} \right) , \tag{2.8}$$

where μ is a dimensionless constant. Recent approaches to calculating the rate of baryon number violating events in the broken phase have been primarily numerical. The best calculation to date of the broken phase sphaleron rate is undoubtably that by Moore [32].

6

This work yields a fully nonperturbative evaluation of the broken phase rate by using a combination of multicanonical weighting and real time techniques.

Although the Boltzmann suppression in (2.8) appears large, it is to be expected that, when the electroweak symmetry becomes restored at a temperature of around $100\,\mathrm{GeV}$, there will no longer be an exponential suppression factor. Although calculation of the baryon number violating rate in the high temperature *unbroken* phase is extremely difficult, a simple estimate is possible. The only important scale in the symmetric phase is the magnetic screening length given by $\xi = (\alpha_W T)^{-1}$. Thus, on dimensional grounds, we expect the rate per unit volume of sphaleron events to be

$$\Gamma_{sp}(T) = \kappa(\alpha_W T)^4 \; , \qquad (2.9)$$

with κ another dimensionless constant. The rate of sphaleron processes can be related to the diffusion constant for Chern-Simons number by a fluctuation-dissipation theorem [33] (for a good description of this see [10]). In almost all numerical calculations of the sphaleron rate, this relationship is used and what is actually evaluated is the diffusion constant. The first attempts to numerically estimate κ in this way [34,35] yielded $\kappa \sim 0.1 - 1$, but the approach suffered from limited statistics and large volume systematic errors. Nevertheless, more recent numerical attempts [36–39] found approximately the same result. However, these approaches employ a poor definition of the Chern-Simons number which compromises their reliability.

This simple scaling argument leading to (2.9) has been recently criticized in refs. [40,41] where it has been argued that damping effects in the plasma suppress the rate by an extra power of α_W to give $\Gamma_{sp} \sim \alpha_W^5 T^4$. Indeed, since the transition rate involves physics at soft energies $g^2 T$ that are small compared to the typical hard energies $\sim T$ of the thermal excitations in the plasma, the simplest way of analyzing the problem is to consider an effective theory for the soft modes, where the hard modes have been integrated out, and to keep the dominant contributions, the so-called hard thermal loops. It is the resulting typical frequency ω_c of a gauge field configuration immersed in the plasma and with spatial extent $(g^2 T)^{-1}$ that determines the change of baryon number per unit time and unit volume. This frequency ω_c has been estimated to be $\sim g^4 T$ when taking into account the damping effects of the hard modes [40,41]. The effective dynamics of soft nonabelian gauge fields at finite temperature has also been recently addressed in [42,43], where it was found that $\Gamma_{sp} \sim \alpha_W^5 T^4 \ln(1/\alpha_W)$. Lattice simulations with hard-thermal loops included have been performed [44] and seem to indicate the $\Gamma_{sp} \sim 30\alpha_W^5 T^4$, which is not far from $\alpha_W^4 T^4$.

1. Other ways of achieving B-violation.

One of the features which distinguishes supersymmetric field theories [45,46] from ordinary ones is the existence of "flat directions" in field space on which the scalar potential vanishes. At the level of renormalizable terms, such flat directions are generic. Supersymmetry breaking lifts these directions and sets the scale for their potential. ¿From a cosmological perspective, these flat directions can have profound consequences. The parameters which describe the flat directions can be thought of as expectation values of massless chiral fields, the moduli. Many flat directions present in the minimal supersymmetric standard model

(MSSM), carry baryon or lepton number, since along them squarks and sleptons have non zero vacuum expectation values (VEVs). If baryon and lepton number are explicitly broken it is possible to excite a non-zero baryon or lepton number along such directions, as first suggested by Affleck and Dine [47] in a more general context. We shall explore these ideas more thoroughly in section V.

C. CP violation

1. CP-Violation in Grand Unified Theories.

CP violation in GUTs arises in loop-diagram corrections to baryon number violating bosonic decays. Since it is necessary that the particles in the loop also undergo B-violating decays, the relevant particles are the X, Y, and H_3 bosons in the case of $SU(5)$. In that case, CP violation is due to the complex phases of the Yukawa couplings h_U and h_D which cannot be reabsorbed by field redefinition. At tree-level these phases do not give any contribution to the baryon asymmetry and at the one-loop level the asymmetry is proportional to Im Tr $\left(h_U^\dagger h_U h_D^\dagger h_D \right) = 0$, where the trace is over generation indices [19]. The problem of too little CP violation in $SU(5)$ may be solved by further complicating the Higgs sector. For example, adding a Higgs in the 45 representation of $SU(5)$ leads to an adequate baryon asymmetry for a wide range of the parameters [48]. In the case of $SO(10)$, CP violation may be provided by the complex Yukawa couplings between the right-handed and the left-handed neutrinos and the scalar Higgs.

2. CP violation in the CKM matrix of the Standard Model.

Since only the left-handed electron is $SU(2)_L$ gauge coupled, C is maximally broken in the SM. Moreover, CP is known not to be an exact symmetry of the weak interactions. This is seen experimentally in the neutral Kaon system through K_0, $\bar{K}_0$ mixing. At present there is no accepted theoretical explanation of this. However, it is true that CP violation is a natural feature of the standard electroweak model. There exists a charged current which, in the weak interaction basis, may be written as $\mathcal{L}_W = (g/\sqrt{2})\bar{U}_L\gamma^\mu D_L W_\mu + $ h.c., where $U_L = (u, c, t, \ldots)_L$ and $D_L = (d, s, b, \ldots)_L$. Now, the quark mass matrices may be diagonalized by unitary matrices V_L^U, V_R^U, V_L^D, V_R^D via

$$\text{diag}(m_u, m_c, m_t, \ldots) = V_L^U M^U V_R^U \ , \tag{2.10}$$

$$\text{diag}(m_d, m_s, m_b, \ldots) = V_L^D M^D V_R^D \ . \tag{2.11}$$

Thus, in the basis of quark mass eigenstates, the charged current may be rewritten as $\mathcal{L}_W = (g/\sqrt{2})\bar{U}'_L K\gamma^\mu D'_L W_\mu + $ h.c., where $U'_L \equiv V_L^U U_L$ and $D'_L \equiv V_L^D D_L$. The matrix K, defined by $K \equiv V_L^U (V_L^D)^\dagger$, is referred to as the Kobayashi-Maskawa (KM) quark mass mixing matrix. For three generations, as in the SM, there is precisely one independent nonzero phase δ which signals CP-violation. While this is encouraging for baryogenesis, it turns out that this particular source of CP violation is not strong enough. The relevant effects are parameterized by a dimensionless constant which is no larger than 10^{-20}. This

appears to be much too small to account for the observed BAU and, thus far, attempts to utilize this source of CP violation for electroweak baryogenesis have been unsuccessful. In light of this, it is usual to extend the SM in some minimal fashion that increases the amount of CP violation in the theory while not leading to results that conflict with current experimental data.

3. CP violation in supersymmetric models

The most promising and well-motivated framework incorporating CP violation beyond the SM seems to be supersymmetry [45,46]. Let us consider the MSSM superpotential

$$W = \mu \hat{H}_1 \hat{H}_2 + h^u \hat{H}_2 \hat{Q} \hat{u}^c + h^d \hat{H}_1 \hat{Q} \hat{d}^c + h^e \hat{H}_1 \hat{L} \hat{e}^c, \tag{2.12}$$

where we have omitted the generation indices. Here $\hat{H}_1$ and $\hat{H}_2$ represent the two Higgs superfields, $\hat{Q}$, $\hat{u}$ and $\hat{d}$ are the quark doublet and the up-quark and down-quark singlet superfields respectively and $\hat{L}$ and $\hat{e}^c$ are the left-handed doublet and right-handed lepton singlet superfields.

The lepton Yukawa matrix h^e can be always taken real and diagonal while h^u and h^d contain the KM phase. There are four new sources for explicit CP violating phases, all arising from parameters that softly break supersymmetry. These are (i) the trilinear couplings

$$\Gamma^u H_2 \widetilde{Q} \widetilde{u}^c + \Gamma^d H_1 \widetilde{Q} \widetilde{d}^c + \Gamma^e H_1 \widetilde{L} \widetilde{e}^c + \text{h.c.} \ , \tag{2.13}$$

where we have defined

$$\Gamma^{(u,d,e)} \equiv A^{(u,d,e)} \cdot h^{(u,d,e)} \ , \tag{2.14}$$

and the tildes stand for scalar fields, (ii) the bilinear coupling in the Higgs sector $\mu B H_1 H_2$, (iii) the gaugino masses M (one for each gauge group), and (iv) the soft scalar masses $\widetilde{m}$. Two phases may be removed by redefining the phase of the superfield $\hat{H}_2$ in such a way that the phase of μ is opposite to that of B. The product μB is therefore real. It is also possible to remove the phase of the gaugino masses M by an R symmetry transformation. The latter leaves all the other supersymmetric couplings invariant, and only modifies the trilinear ones, which get multiplied by $\exp(-\phi_M)$ where ϕ_M is the phase of M. The phases which remain are therefore

$$\phi_A = \arg(AM) \ \text{ and } \ \phi_\mu = -\arg(B). \tag{2.15}$$

These new phases ϕ_A and ϕ_μ may be crucial for the generation of the baryon asymmetry.

In the MSSM, however, there are other possible sources of CP violation. In fact, when supersymmetry breaking occurs, as we know it must, the interactions of the Higgs fields H_1 and H_2 with charginos, neutralinos and stops (the superpartners of the charged, neutral gauge bosons, Higgs bosons and tops, respectively) at the one-loop level, lead to a CP violating contribution to the scalar potential of the form

$$V_{CP} = \lambda_1 (H_1 H_2)^2 + \lambda_2 |H_1|^2 H_1 H_2 + \lambda_3 |H_2|^2 H_1 H_2 + \text{h.c.}. \tag{2.16}$$

These corrections may be quite large at finite temperature [49,50]. One may write the Higgs fields in unitary gauge as $H_1 = (\varphi_1, 0)^T$ and $H_2 = (0, \varphi_2 e^{i\theta})^T$, where φ_1, φ_2, θ are real, and θ is the CP-odd phase. The coefficients $\lambda_{1,2,3}$ determine whether the CP-odd field θ is non zero, signalling the spontaneous violation of CP in a similar way to what happens in a generic two-Higgs model [51]. If, during the evolution of the early universe, a variation of the phase θ is induced, this CP breakdown may bias the production of baryons through sphaleron processes in electroweak baryogenesis models. The attractive feature of this possibility is that, when the universe cools down, the finite temperature loop corrections to the λ-couplings disappear, and the phase θ relaxes to zero. This is turn implies that, unlike for scenarios utilizing ϕ_A and ϕ_μ, one need not worry about present experimental constraints from the physics of CP violation. A recent study [52] has shown that when the constraints from experiment and the strength of the transition are taken into account, the relevant region of the parameter space for spontaneous CP violation in the MSSM at finite temperature is rather restricted. Nevertheless, in this small region, perturbative estimates need not be reliable, and non-perturbatively the region might be slightly larger (or smaller). More precise perturbative studies and/or lattice studies of the non-perturbative corrections are certainly needed.

Finally, in scenarios in which the BAU is generated from the coherent production of a scalar condensate along a flat direction of the supersymmetric extension of the SM, CP-violation is present in the nonrenormalizable superpotentials which lift the flat directions at large field values. We shall have more to say about this is section (V).

D. Departure from Thermal Equlibrium

1. The out-of-equilibrium decay scenario

Scenarios in which the third Sakharov condition is satisfied due to the presence of superheavy decaying particles in a rapidly expanding universe, generically fall under the name of out-of-equilibrium decay mechanisms [6,19]. The underlying idea is fairly simple. If the decay rate Γ_X of the superheavy particles X at the time they become nonrelativistic (*i.e.* at the temperature $T \sim M_X$) is much smaller than the expansion rate of the universe, then the X particles cannot decay on the time scale of the expansion and so they remain as abundant as photons for $T \lesssim M_X$. In other words, at some temperature $T > M_X$, the superheavy particles X are so weakly interacting that they cannot catch up with the expansion of the universe and they decouple from the thermal bath while they are still relativistic, so that $n_X \sim n_\gamma \sim T^3$ at the time of decoupling. Therefore, at temperature $T \simeq M_X$, they populate the universe with an abundance which is much larger than the equilibrium one. This overbundance is precisely the departure from thermal equilibrium needed to produce a final nonvanishing baryon asymmetry when the heavy states X undergo B and CP violating decays. The out-of-equilibrium condition requires very heavy states: $M_X \gtrsim (10^{15} - 10^{16})\,\text{GeV}$ and $M_X \gtrsim (10^{10} - 10^{16})\,\text{GeV}$, for gauge and scalar bosons, respectively [19], if these heavy particles decay through renormalizable operators.

E. The Electroweak Phase Transition Scenario

At temperatures around the electroweak scale, the expansion rate of the universe in thermal units is small compared to the rate of baryon number violating processes. This means that the equilibrium description of particle phenomena is extremely accurate at electroweak temperatures. Thus, baryogenesis cannot occur at such low scales without the aid of phase transitions and the question of the order of the electroweak phase transition becomes central.

If the EWPT is second order or a continuous crossover, the associated departure from equilibrium is insufficient to lead to relevant baryon number production [28]. This means that for EWBG to succeed, we either need the EWPT to be strongly first order or other methods of destroying thermal equilibrium to be present at the phase transition.

Any thermodynamic quantity that undergoes such a discontinuous change at a phase transition is referred to as an *order parameter*, denoted by φ. For a first order transition there is an extremum at $\varphi = 0$ which becomes separated from a second local minimum by an energy barrier. At the critical temperature $T = T_c$ both phases are equally favored energetically and at later times the minimum at $\varphi \neq 0$ becomes the global minimum of the theory. The dynamics of the phase tansition in this situation is crucial to most scenarios of electroweak baryogenesis. The essential picture is that around T_c quantum tunneling occurs and nucleation of bubbles of the true vacuum in the sea of false begins. Initially these bubbles are not large enough for their volume energy to overcome the competing surface tension and they shrink and disappear. However, at a particular temperature below T_c, bubbles just large enough to grow nucleate. These are termed *critical* bubbles, and they expand, eventually filling all of space and completing the transition. As the bubble walls pass each point in space, the order parameter changes rapidly, as do the other fields, and this leads to a significant departure from thermal equilibrium. Thus, if the phase transition is strongly enough first order it is possible to satisfy the third Sakharov criterion in this way.

The precise evolution of critical bubbles in the electroweak phase transition is a crucial factor in determining which regimes of electroweak baryogenesis are both possible and efficient enough to produce the BAU. (See [53–60] for possible obstacles to the standard picture) We discuss this, and the issue of the order of the phase transition in detail in section IV.

For the bubble wall scenario to be successful, there is a further criterion to be satisfied. As the wall passes a point in space, the Higgs fields evolve rapidly and the Higgs VEV changes from $\langle \phi \rangle = 0$ in the unbroken phase to

$$\langle \phi \rangle = v(T_c) \tag{2.17}$$

in the broken phase. Here, $v(T)$ is the value of the order parameter at the symmetry breaking global minimum of the finite temperature effective potential. Now, CP violation and the departure from equilibrium occur while the Higgs field is changing. Afterwards, the point is in the true vacuum, baryogenesis has ended, and baryon number violation is exponentially supressed. Since baryogenesis is now over, it is imperative that baryon number violation be negligible at this temperature in the broken phase, otherwise any baryonic excess generated will be equilibrated to zero. Such an effect is known as *washout* of the asymmetry and the criterion for this not to happen may be written as

$$\frac{v(T_c)}{T_c} \gtrsim 1 \ . \tag{2.18}$$

11

Although there are a number of nontrivial steps that lead to this simple criterion, (2.18) is traditionally used to ensure that the baryon asymmetry survives after the wall has passed. It is necessary that this criterion be satisfied for any electroweak baryogenesis scenario to be successful.

1. Defect-Mediated Baryogenesis

A natural way to depart from equilibrium is provided by the dynamics of topological defects. Topological defects are regions of trapped energy density which can remain after a cosmological phase transition if the topology of the vacuum of the theory is nontrivial. Typically, cosmological phase transitions occur when a symmetry of a particle physics theory is spontaneously broken. In that case, the cores of the topological defects formed are regions in which the symmetry of the unbroken theory is restored. If, for example, cosmic strings are produced at the GUT phase transition, then the decays of loops of string act as an additional source of superheavy bosons which undergo baryon number violating decays. When defects are produced at the TeV scale, a detailed analysis of the dynamics of a network of these objects shows that a significant baryon to entropy ratio can be generated [61,64,62,63,65,66] if the electroweak symmetry is restored around such a higher scale ordinary defect [67–69]. Although a recent analysis has shown that B-violation is highly inefficient along nonsuperconducting strings [70], there remain viable scenarios involving other ordinary defects, superconducting strings or defects carrying baryon number [71,72].

III. RECENT DEVELOPMENTS IN GUT BARYOGENESIS

As we have seen, in GUTs baryon number violation is natural since quarks and leptons lie in the same irreducible representation of the gauge group. The implementation of the out-of-equilibrium scenario in this context goes under the name of GUT baryogenesis. The basic assumption is that the superheavy bosons were as abundant as photons at very high temperatures $T \gtrsim M_X$. This assumption is questionable if the heavy X particles are the gauge or Higgs bosons of Grand Unification, because they might have never been in thermal equilibrium in the early universe. Indeed, the temperature of the universe might have been always smaller than M_{GUT} and, correspondingly, the thermally produced X bosons might never have been as abundant as photons, making their role in baryogenesis negligible. All these considerations depend crucially upon the thermal history of the universe and deserve a closer look.

A. Inflation, reheating and GUT baryogenesis: the old wisdom

The flatness and the horizon problems of the standard big bang cosmology are elegantly solved if, during the evolution of the early universe, the energy density happened to be dominated by some form of vacuum energy, causing comoving scales grow quasi-exponentially [76–79]. The vacuum energy driving this *inflation* is generally assumed to be associated with the potential $V(\phi)$ of some scalar field ϕ, the *inflaton*, which is initially displaced from the minimum of its potential. As a by-product, quantum fluctuations of the inflaton field

may be the seeds for the generation of structure and the fluctuations observed in the cosmic microwave background radiation [79]. Inflation ended when the potential energy associated with the inflaton field became smaller than the kinetic energy of the field. By that time, any pre-inflationary entropy in the universe had been inflated away, and the energy of the universe was entirely in the form of coherent oscillations of the inflaton condensate around the minimum of its potential. Now, we know that somehow this low-entropy cold universe must be transformed into a high-entropy hot universe dominated by radiation. The process by which the energy of the inflaton field is transferred from the inflaton field to radiation has been dubbed *reheating*. The simplest way to envision this process is if the comoving energy density in the zero mode of the inflaton decays *perturbatively* into normal particles, which then scatter and thermalize to form a thermal background [80,81]. It is usually assumed that the decay width of this process is the same as the decay width of a free inflaton field. Of particular interest is a quantity known as the reheat temperature, denoted as T_r. This is calculated by assuming an instantaneous conversion of the energy density in the inflaton field into radiation when the decay width of the inflaton energy, Γ_ϕ, is equal to H, the expansion rate of the universe. This yields

$$T_r = \sqrt{\Gamma_\phi m_p}. \tag{3.1}$$

There are very good reasons to suspect that GUT baryogenesis does not occur if this is the way reheating happens. First of all, the density and temperature fluctuations observed in the present universe require the inflaton potential to be extremely flat – that is $\Gamma_\phi \ll M_\phi$. This means that the couplings of the inflaton field to the other degrees of freedom cannot be too large, since large couplings would induce large loop corrections to the inflaton potential, spoiling its flatness. As a result, T_r is expected to be smaller than 10^{14} GeV by several orders of magnitude. On the other hand, the unification scale is generally assumed to be around 10^{16} GeV, and B-violating bosons should have masses comparable to this scale. Furthermore, even the light B-violating Higgs bosons are expected to have masses larger than the inflaton mass, and thus it would be kinematically impossible to create them directly in ϕ decays, $\phi \to X\overline{X}$.

One might think that the heavy bosons could be created during the stage of thermalization. Indeed, it has recently been shown that particles of mass $\sim 10^4$ bigger than the reheating temperature T_r may be created by the thermalized decay products of the inflaton [82–85]. However, there is one more problem associated with GUT baryogenesis in the old theory of reheating, namely the problem of relic gravitinos [86]. If one has to invoke supersymmetry to preserve the flatness of the inflaton potential, it is mandatory to consider the cosmological implications of the gravitino – a spin-$(3/2)$ particle which appears in the extension of global supersymmetry to supergravity. The gravitino is the fermionic superpartner of the graviton and has interaction strength with the observable sector – that is the standard model particles and their superpartners – inversely proportional to the Planck mass. The slow gravitino decay rate leads to a cosmological problem because the decay products of the gravitino destroy ^{4}He and D nuclei by photodissociation, and thus ruin the successful predictions of nucleosynthesis. The requirement that not too many gravitinos are produced after inflation provides a stringent constraint on the reheating temperature, $T_r \lesssim (10^{10} - 10^{11})$ GeV [86]. Therefore, if $T_r \sim M_{\mathrm{GUT}}$, gravitinos would be abundant during nucleosynthesis and destroy the agreement of the theory with observations. However, if the

initial state after inflation was free from gravitinos, the reheating temperature is then too low to create superheavy X bosons that eventually decay and produce the baryon asymmetry.

B. Inflation, reheating and GUT baryogenesis: the new wisdom

The outlook for GUT baryogenesis has brightened recently with the realization that reheating may differ significantly from the simple picture described above [87–91]. In the first stage of reheating, called *preheating* [87], nonlinear quantum effects may lead to extremely effective dissipative dynamics and explosive particle production, even when single particle decay is kinematically forbidden. In this picture, particles can be produced in a regime of broad parametric resonance, and it is possible that a significant fraction of the energy stored in the form of coherent inflaton oscillations at the end of inflation is released after only a dozen or so oscillation periods of the inflaton. What is most relevant for us is that preheating may play an extremely important role in baryogenesis and, in particular, for the Grand Unified generation of a baryonic excess. Indeed, it was shown in [92,93] that the baryon asymmetry can be produced efficiently just after the preheating era, thus solving many of the problems that GUT baryogenesis had to face in the old picture of reheating.

The presence of a preheating stage is due to the fact that, for some parameter ranges, there is a new non-perturbative decay channel. Coherent oscillations of the inflaton field induce stimulated emissions of bosonic particles into energy bands with large occupancy numbers [87]. The modes in these bands can be understood as Bose condensates, and they behave like classical waves.

The idea of preheating is relatively simple, the oscillations of the inflaton field induce a time-dependent mass for fields coupled to the inflaton. This leads to a mixing of positive and negative frequencies in the quantum state of these fields. For the sake of simplicity, let us focus on the case of chaotic inflation, with a massive inflaton ϕ, with quadratic potential $V(\phi) = \frac{1}{2}M_\phi^2\phi^2$, $M_\phi \sim 10^{13}$ GeV, coupled to a massive scalar field X via the quartic coupling $g^2\phi^2|X|^2$. Let us also neglect the expansion of the universe.

The evolution equation for the Fourier modes of the X field with momentum k is

$$\ddot{X}_k + \omega_k^2 X_k = 0, \tag{3.2}$$

with $\omega_k^2 = k^2 + M_X^2 + g^2\phi^2(t)$. This Klein-Gordon equation may be cast in the form of a Mathieu equation

$$X_k'' + [A(k) - 2q\cos 2z]X_k = 0, \tag{3.3}$$

where $z = M_\phi t$ and

$$A(k) = \frac{k^2 + M_X^2}{M_\phi^2} + 2q,$$

$$q = g^2\frac{\Phi^2}{4M_\phi^2}, \tag{3.4}$$

where Φ is the amplitude and M_ϕ is the frequency of inflaton oscillations, defined through $\phi(t) = \Phi(t)\sin(M_\phi t)$. Notice that, at least initially when $\Phi = cm_p \lesssim m_p$

14

$$g^2 \frac{\Phi^2}{4M_\phi^2} \sim g^2 c^2 \frac{m_p^2}{M_\phi^2} \sim g^2 c^2 \times 10^{12} \gg 1 \qquad (3.5)$$

and the resonance is broad. A crucial observation for baryogenesis is that particles with mass larger than that of the inflaton may be produced during preheating [92]. Particle production occurs above the line $A = 2q$, and the width of the instability strip scales as $q^{1/2}$ for large q, independent of M_X. The condition for broad resonance [87,92] $A - 2q \lesssim q^{1/2}$ becomes $(k^2 + M_X^2)/M_\phi^2 \lesssim g(\Phi/M_\phi)$, which yields $E_X^2 = k^2 + M_X^2 \lesssim g\Phi M_\phi$ for the typical energy of X bosons produced in preheating. By the time the resonance develops to the full strength, $\Phi^2 \sim 10^{-5} m_p^2$, and so the resulting estimate for the typical energy of particles at the end of the broad resonance regime for $M_\phi \sim 10^{13}$ GeV is

$$E_X \sim 10^{-1} g^{1/2} \sqrt{M_\phi m_p} \sim g^{1/2} 10^{15} \text{ GeV}. \qquad (3.6)$$

Supermassive X bosons can be produced by the broad parametric resonance for $E_X > M_X$, which leads to the condition $M_X < g^{1/2} 10^{15}$ GeV [87,92]. Thus, for $g^2 \sim 1$ one would have copious production of X particles as heavy as 10^{15} GeV, *i.e.*, 100 times greater than the inflaton mass.

The corresponding number density n_k grows exponentially with time because these particles belong to an instability band of the Mathieu equation (for a recent comprehensive review on preheating after chaotic inflation [94] and references therein)

$$X_k \propto e^{\mu_k M_\phi t} \Rightarrow n_k \propto e^{2\mu_k M_\phi t}, \qquad (3.7)$$

where the parameter μ_k depends upon the instability band and, in the broad resonance case, $q \gg 1$, it is ~ 0.2.

These instabilities can be interpreted as coherent "particle" production with large occupancy numbers. One way of understanding this phenomenon is to consider the energy of these modes as that of a harmonic oscillator,

$$E_k = |\dot{X}_k|^2/2 + \omega_k^2 |X_k|^2/2 = \omega_k(n_k + 1/2) . \qquad (3.8)$$

The occupancy number of level k can grow exponentially fast and these modes soon behave like classical waves. The parameter q during preheating determines the strength of the resonance.

This is a significant departure from the old constraints of reheating. Production of X bosons in the old reheating picture was kinematically forbidden if $M_\phi < M_X$, while in the new scenario it is possible because of coherent effects. Not only are particles produced out-of-equilibrium, thus satisfying one of the basic requirements to produce the baryon asymmetry, but also gravitinos may be shown to be produced in sufficiently small amounts [92] to avoid the gravitino problem we discussed earlier.

Back-reaction, or scattering of X fluctuations off the zero mode of the inflaton field, limits the maximum magnitude of X fluctuations $\langle X^2 \rangle$ [90]. This in turn restricts the corresponding number density of created X-particles and therefore weakens the whole mechanism. Parametric resonance is also rendered less efficient when the X particles have a decay width Γ_X, a crucial parameter in GUT baryogenesis. Using the methods developed in Refs. [88–90], the problem of the production of massive, unstable X particles in the process of the inflaton

decay has been recently addressed [93] for the case of a quadratic inflaton potential. It was shown that the observed baryon asymmetry may be explained by GUT baryogenesis after preheating if

$$\frac{\langle X^2 \rangle}{m_p^2} \simeq 5 \times 10^{-13} \left(\frac{10^{-2}}{\epsilon} \right) \left(\frac{\Gamma}{5 \times 10^{-5}} \right), \tag{3.9}$$

where $\Gamma \equiv \Gamma_X / M_\phi$, and ϵ is a dimensionless measure of the CP violation in the system. This may happen if [93] $\Gamma_X \lesssim 10^{-3} M_X$, which is a relatively mild constraint.

Grand Unified baryogenesis after preheating solves many of the serious drawbacks of GUT baryogenesis in the old theory of reheating where the production of superheavy states after inflation was kinematically forbidden. Moreover, the out-of-equilibrium condition is naturally achieved in this scenario, since the distribution function of the X-quanta generated at the resonance is far from a thermal distribution. This situation is considerably different from the one present in the thermal GUT scenario where superheavy particles usually decouple from the thermal bath when still relativistic and then decay producing the baryon asymmetry. It is quite intriguing that out of all possible ways the parametric resonance may develop, Nature might have chosen only those ways without instantaneous thermalization and also with a successful baryogenesis scenario.

C. Baryogenesis via leptogenesis

Since the linear combination $B - L$ is left unchanged by sphaleron transitions, the baryon asymmetry may be generated from a lepton asymmetry [23]. Indeed, sphaleron transition will reprocess any lepton asymmetry and convert (a fraction of) it into baryon number. This is because $B + L$ must be vanishing and the final baryon asymmetry results to be $B \simeq -L$. Once the lepton number is produced, thermal scatterings redistribute the charges. In the high temperature phase of the SM, the asymmetries of baryon number B and of $B - L$ are therefore proportional: [33]

$$B = \left(\frac{8n_f + 4n_H}{22n_f + 13n_H} \right) (B - L), \tag{3.10}$$

where n_H is the number of Higgs doublets. In the SM as well as in its unified extension based on the group $SU(5)$, $B - L$ is conserved and no asymmetry in $B - L$ can be generated. However, adding right-handed Majorana neutrinos to the SM breaks $B - L$ and the primordial lepton asymmetry may be generated by the out-of-equilibrium decay of heavy right-handed Majorana neutrinos N_L^c (in the supersymmetric version, heavy scalar neutrino decays are also relevant for leptogenesis). This simple extension of the SM can be embedded into GUTs with gauge groups containing $SO(10)$. Heavy right-handed Majorana neutrinos can also explain the smallness of the light neutrino masses via the see-saw mechanism [95].

The relevant couplings are those between the right-handed neutrino, the Higgs doublet Φ and the lepton doublet ℓ_L

$$\mathcal{L} = \overline{\ell_L} \, \Phi \, h_\nu \, N_L^c + \frac{1}{2} \overline{N_L^c} \, M \, N_L^c + \text{h.c.} \tag{3.11}$$

16

The vacuum expectation value of the Higgs field $\langle\Phi\rangle$ generates neutrino Dirac masses $m_D = h_\nu\langle\Phi\rangle$, which are assumed to be much smaller than the Majorana masses M. When the Majorana right-handed neutrinos decay into leptons and Higgs scalars, they violate the lepton number since right-handed neutrino fermionic lines do not have any preferred arrow

$$N_L^c \to \overline{\Phi} + \ell,$$
$$N_L^c \to \Phi + \overline{\ell}. \tag{3.12}$$

The interference between the tree-level decay amplitude and the absorptive part of the one-loop vertex leads to a lepton asymmetry of the right order of magnitude to explain the observed baryon asymmetry and has been discussed extensively in the literature [23,96–98]. Recently, much attention has been paid to the effects of finite temperature on the CP violation [99], and the contribution to the lepton asymmetry generated by the tree-level graph and the absorptive part of the one-loop self-energy. In particular, it has been observed that CP violation may be considerably enhanced if two heavy right-handed neutrinos are nearly degenerate in mass [100]. It is also important to remember that large lepton number violation at intermediate temperatures may potentially dissipate away the baryon number in combination with the sphaleron transitions. Indeed, $\Delta L = 2$ interactions of the form

$$\frac{m_\nu}{\langle\Phi\rangle^2}\ell_L\ell_L\Phi\Phi + \text{h.c.}, \tag{3.13}$$

where m_ν is the mass of the light left-handed neutrino, are generated through the exchange of heavy right-handed neutrinos. The rate of lepton number violation induced by this interaction is therefore $\Gamma_L \sim (m_\nu^2/\langle\Phi\rangle^4)T^3$. The requirement of harmless letpon number violation, $\Gamma_L \lesssim H$ imposes an interesting bound on the neutrino mass

$$m_\nu \lesssim 4\,\text{eV}\left(\frac{T_X}{10^{10}\,\text{GeV}}\right)^{-1/2}, \tag{3.14}$$

where $T_X \equiv \text{Min}\{T_{B-L}, 10^{12}\,\text{GeV}\}$, T_{B-L} is the temperature at which the $B-L$ number production takes place, and $\sim 10^{12}$ GeV is the temperature at which sphaleron transitions enter in equilibrium. One can also reverse the argument and study leptogenesis assuming a similar pattern of mixings and masses for leptons and quarks, as suggested by $SO(10)$ unification [101]. This implies that $B-L$ is broken at the unification scale $\sim 10^{16}$ GeV, if $m_{\nu_\mu} \sim 3 \times 10^{-3}$ eV as preferred by the MSW explanation of the solar neutrino deficit [102,103].

IV. ELECTROWEAK BARYOGENESIS

Scenarios in which anomalous baryon number violation in the electroweak theory implements the first Sakharov criterion are refered to as *electroweak baryogenesis*. Typically these scenarios require a strongly first order phase transition, either in the minimal standard model, or in a modest extension. Another possibility is that the out of equilibrium requirement is achieved through TeV scale topological defects, also present in many extensions of the electroweak theory. Finally, a source of CP-violation beyond that in the CKM matrix is usually involed.

17

Historically, the ways in which baryons may be produced as a bubble wall, or phase boundary, sweeps through space, have been separated into two categories.

1. *local baryogenesis*: baryons are produced when the baryon number violating processes and CP violating processes occur together near the bubble walls.

2. *nonlocal baryogenesis*: particles undergo CP violating interactions with the bubble wall and carry an asymmetry in a quantum number other than baryon number into the unbroken phase region away from the wall. Baryons are then produced as baryon number violating processes convert the existing asymmetry into one in baryon number.

In general, both local and nonlocal baryogenesis will occur and the BAU will be the sum of that generated by the two processes. However, if the speed of the wall is greater than the sound speed in the plasma, then local baryogenesis dominates [104]. In other cases, nonlocal baryogenesis is usually more efficient and we shall focus on that here.

Nonlocal baryogenesis typically involves the interaction of the bubble wall with the various fermionic species in the unbroken phase. The main picture is that as a result of CP violation in the bubble wall, particles with opposite chirality interact differently with the wall, resulting in a net injected chiral flux. This flux thermalizes and diffuses into the unbroken phase where it is converted to baryons. In this section, for definiteness when describing these effects, we assume that the CP violation arises because of a two-Higgs doublet structure.

The chiral asymmetry which is converted to an asymmetry in baryon number is carried by both quarks and leptons. However, the Yukawa couplings of the top quark and the τ-lepton are larger than those of the other quarks and leptons respectively. Therefore, it is reasonable to expect that the main contribution to the injected asymmetry comes from these particles and to neglect the effects of the other particles.

When considering nonlocal baryogenesis it is convenient to write the equation for the rate of production of baryons in the form [105]

$$\frac{dn_B}{dt} = -\frac{n_f \Gamma_{sp}(T)}{2T} \sum_i \mu_i \ , \tag{4.1}$$

where the rate per unit volume for electroweak sphaleron transitions is given by (2.9). Here, n_f is again the number of families and μ_i is the chemical potential for left handed particles of species i. The crucial question in applying this equation is an accurate evaluation of the chemical potentials that bias baryon number production.

There are typically two distinct calculational regimes that are appropriate for the treatment of nonlocal effects in electroweak baryogenesis. Which regime is appropriate depends on the particular fermionic species under consideration.

1. The thin wall regime

The thin wall regime [106,107,105] applies if the mean free path ℓ of the fermions being considered is much greater than the thickness L_w of the wall, i.e. if

$$\frac{L_w}{\ell} \lesssim 1 \; . \tag{4.2}$$

In this case we may neglect scattering effects and treat the fermions as free particles in their interactions with the wall.

In this regime, in the rest frame of the bubble wall, particles see a sharp potential barrier and undergo CP violating interactions with the wall due to the gradient in the CP odd Higgs phase. As a consequence of CP violation, there will be asymmetric reflection and transmission of particles, thus generating an injected current into the unbroken phase in front of the bubble wall. As a consequence of this injected current, asymmetries in certain quantum numbers will diffuse both behind and in front of the wall due to particle interactions and decays [106,107,105]. In particular, the asymmetric reflection and transmission of left and right handed particles will lead to a net injected chiral flux from the wall. However, there is a qualitative difference between the diffusion occurring in the interior and exterior of the bubble.

Exterior to the bubble the electroweak symmetry is restored and weak sphaleron transitions are unsuppressed. This means that the chiral asymmetry carried into this region by transport of the injected particles may be converted to an asymmetry in baryon number by sphaleron effects. In contrast, particles injected into the phase of broken symmetry interior to the bubble may diffuse only by baryon number conserving decays since the electroweak sphaleron rate is exponentially suppressed in this region. Hence, we concentrate only on those particles injected into the unbroken phase.

The net baryon to entropy ratio which results via nonlocal baryogenesis in the case of thin walls has been calculated in several different analyses [106,107] and [105]. In the following we give a brief outline of the logic of the calculation, following [105]. The baryon density produced is given by (4.1) in terms of the chemical potentials μ_i for left handed particles. These chemical potentials are a consequence of the asymmetric reflection and transmission off the walls and the resulting chiral particle asymmetry. Baryon number violation is driven by the chemical potentials for left handed leptons or quarks. To be concrete, we focus on leptons [105] (for quarks see e.g. [106]). If there is local thermal equilibrium in front of the bubble walls then the chemical potentials μ_i of particle species i are related to their number densities n_i by

$$n_i = \frac{T^2}{12} k_i \mu_i \; , \tag{4.3}$$

where k_i is a statistical factor which equals 1 for fermions and 2 for bosons. In deriving this expression, it is important [108] to correctly impose the constraints on quantities which are conserved in the region in front of and on the wall.

Using the above considerations, the chemical potential μ_L for left handed leptons can be related to the left handed lepton number densities L_L. These are in turn determined by particle transport. The source term in the diffusion equation is the flux J_0 resulting from the asymmetric reflection and transmission of left and right handed leptons off the bubble wall.

For simplicity assume a planar wall. If $|p_z|$ is the momentum of the lepton perpendicular to the wall (in the wall frame), the analytic approximation used in [105] allows the asymmetric reflection coefficients for lepton scattering to be calculated in the range

$$m_l < |p_z| < m_h \sim \frac{1}{L_w} \ , \tag{4.4}$$

where m_l and m_h are the lepton and Higgs masses, respectively, and results in

$$\mathcal{R}_{L \to R} - \mathcal{R}_{R \to L} \simeq 2\Delta\theta_{CP} \frac{m_l^2}{m_h |p_z|} \ . \tag{4.5}$$

The corresponding flux of left handed leptons is

$$J_0 \simeq \frac{v_w m_l^2 m_h \Delta\theta_{CP}}{4\pi^2} \ , \tag{4.6}$$

where v_w is the velocity of the bubble wall. Note that in order for the momentum interval in (4.4) to be non-vanishing, the condition $m_l L_w < 1$ needs to be satisfied.

The injected current from a bubble wall will lead to a "diffusion tail" of particles in front of the moving wall. In the approximation in which the persistence length of the injected current is much larger than the wall thickness we may to a good approximation model it as a delta function source and search for a steady state solution. In addition, assume that the decay time of leptons is much longer than the time it takes for a wall to pass so that we may neglect decays. Then the diffusion equation for a single particle species is easily solved by

$$L_L(z) = \begin{cases} J_0 \frac{\xi_L}{D_L} e^{-\lambda_D z} & z > 0 \\ 0 & z < 0 \end{cases} \ , \tag{4.7}$$

with the diffusion root

$$\lambda_D = \frac{v_w}{D_L} \ , \tag{4.8}$$

where D_L is the diffusion constant for leptons, and ξ^L that is called the *persistence length* of the current in front of the bubble wall.

Note that in this approximation the injected current does not generate any perturbation behind the wall. This is true provided $\xi^L \gg L_w$ is satisfied. If this inequality is not true, the problem becomes significantly more complex [105].

In the massless approximation the chemical potential μ_L can be related to L_L by

$$\mu_L = \frac{6}{T^2} L_L \tag{4.9}$$

(for details see [105]). Inserting the sphaleron rate and the above results for the chemical potential μ into (4.1), and using $1/D_L \simeq 8\alpha_W^2 T$, the final baryon to entropy ratio becomes

$$\frac{n_B}{s} \sim 0.2 \, \alpha_W^2 (g^*)^{-1} \kappa \Delta\theta_{CP} \frac{1}{v_w} \left(\frac{m_l}{T}\right)^2 \frac{m_h}{T} \frac{\xi^L}{D_L} \ . \tag{4.10}$$

where we have assumed that sphalerons do not equilibrate in the diffusion tail.

Now consider the effects of top quarks scattering off the advancing wall [106,107]. Several effects tend to decrease the contribution of the top quarks relative to that of tau leptons. Firstly, for typical wall thicknesses the thin wall approximation does not hold for top quarks.

This is because top quarks are much more strongly interacting than leptons and so have a much shorter mean free path. An important effect is that the diffusion tail is cut off in front of the wall by *strong sphalerons* [109,110]. There is an anomaly in the quark axial vector current in QCD. This leads to chirality non-conserving processes at high temperatures. These processes are relevant for nonlocal baryogenesis since it is the chirality of the injected current that is important in that scenario. In an analogous expression to that for weak sphalerons, we may write the rate per unit volume of chirality violating processes due to strong sphalerons in the unbroken phase as

$$\Gamma_s = \kappa_s (\alpha_s T)^4 \ , \tag{4.11}$$

where κ_s is a dimensionless constant [111]. Note that the uncertainties in κ_s are precisely the same as those in κ defined in (2.9). As such, κ_s could easily be proportional to α_s, in analogy with (2.9), perhaps with a logarithmic correction. These chirality-changing processes damp the effect of the injected chiral flux and effectively cut off the diffusion tail in front of the advancing bubble wall. Second, the diffusion length for top quarks is intrinsically smaller than that for tau leptons, thus reducing the volume in which baryogenesis takes place. Although there are also enhancement factors, e.g. the ratio of the squares of the masses m_t^2/m_τ^2, it seems that leptons provide the dominant contribution to nonlocal baryogenesis.

2. The thick wall regime

Now we turn to the thick wall, or adiabatic, regime. This is relevant if the mean free path of the particles being considered is smaller than the width of the wall, $\ell \lesssim L_w$. When the walls are thick, most interactions within the wall will be almost in thermal equilibrium. The equilibrium is not exact because some interactions, in particular baryon number violation, take place on a time scale much slower than the rate of passage of the bubble wall. These slowly-varying quantities are best treated by the method of chemical potentials. The basic idea is to produce asymmetries in some local charges which are (approximately) conserved by the interactions inside the bubble walls, where local departure from thermal equilibrium is attained. These local charges will then diffuse into the unbroken phase where baryon number violation is active thanks to the unsuppressed sphaleron transitions. The latter convert the asymmetries into a baryon asymmetry. Therefore, one has to *i)* identify those charges which are approximately conserved in the symmetric phase, so that they can efficiently diffuse in front of the bubble where baryon number violation is fast, *ii)* compute the CP violating currents of the plasma induced inside the bubble wall and *iii)* solve a set of coupled differential diffusion equations for the local particle densities, including the CP violating sources.

There are a number of possible CP violating sources. To be definite, consider the example where CP violation is due to a CP odd phase θ in the two-Higgs doublet model [112]. This goes generically under the name of spontaneous baryogenesis. In order to explicitly see how θ couples to the fermionic sector of the theory (to produce baryons) we may remove the θ-dependence of the Yukawa couplings arising from the Higgs terms. We do this by performing an anomaly-free hypercharge rotation on the fermions [112], inducing a term in the Lagrangian density of the form $\mathcal{L}_{CP} \propto \partial_\mu \theta J_Y^\mu$, where J_Y^μ is the fermionic part of

the hypercharge current. Therefore, a nonvanishing $\dot{\theta}$ provides a preferential direction for the production of quarks and leptons; in essence a chemical potential $\mu_B \propto \dot{\theta}$ (in the rest frame of the thermal bath) for baryon number. Of course, strictly speaking, this is not a chemical potential since it arises dynamically rather than to impose a constraint. For this reason, the quantity μ_B is sometimes referred to as a "charge potential". A more complete field theoretic approach to the computation of particle currents on a space-time dependent and CP violating Higgs background was provided in [113,114] where it was shown that fermionic currents arise at one loop, and are proportional to $(hv(T_c)/\pi T)^2 \dot{\theta}$, where h is the appropriate Yukawa coupling. The relevant baryon number produced by spontaneous baryogenesis is then calculated by integrating

$$\frac{dn_B}{dt} = -9 \frac{\Gamma_{sp}(T)}{T} \mu_B \; . \tag{4.12}$$

Initially, spontaneous baryogenesis was considered as an example of local baryogenesis. However, it has become clear that diffusion effects can lead to an appreciable enhancement of the baryon asymmetry produced by this mechanism [115–117].

An alternative way of generating a source for the diffusion equation was suggested by Joyce et al. [118]. The essential idea is that there exists a purely classical chiral force that acts on particles when there is a CP violating field on the bubble wall. The Lagrangian for a fermion ψ in the background of a bubble wall in the presence of the CP-odd phase θ is

$$\tilde{\mathcal{L}} = i\bar{\psi}\gamma^\mu \left(\partial_\mu + \frac{i}{2} \frac{v_2^2}{v_1^2 + v_2^2} \gamma^5 \partial_\mu \theta \right) \psi - m\bar{\psi}\psi \; , \tag{4.13}$$

where v_1 and v_2 are the VEVs of the Higgs fields and m is the fermion mass. As usual we assume a planar wall propagating at speed v in the z-direction. Further, if interactions of the Higgs and gauge fields with the wall have reached a stationary state, then these fields are functions only of $z - vt$. Then, in the wall rest frame plane-wave fermion solutions $\psi \propto e^{-ip.x}$ have the dispersion relation

$$E = \left[p_\perp^2 + \left(\sqrt{p_z^2 + m^2} \pm \frac{1}{2} \frac{v_2^2}{v_1^2 + v_2^2} \partial_z \theta \right)^2 \right]^{1/2} \; , \tag{4.14}$$

where $\pm$ corresponds to $s_z = \pm 1/2$, with s_z the z-component of the spin.

In the WKB approximation, the Hamilton equations derived from this relation allow calculation of the acceleration of a WKB wave-packet (i.e. the chiral force). This force acts as a potential well that induces an excess of chiral charge on the wall. In the diffusion tail in front of the wall, there exists a corresponding deficit of chiral charge. This acts as a chemical potential which sources the production of baryon number.

This calculation is performed by following a fluid approach to solving the associated Boltzmann equation, By concentrating on particles with $|p_z| \sim T \gg m$, the particle and antiparticle excitations (4.14) may be treated as separate fluids in the WKB approximation. This permits an analytic solution to the diffusion equation which yields a relation between the chemical potential on the wall for the fermion species considered and the chiral force. Integrating this in front of the wall yields a chemical potential that may then be fed into (4.1) to give an approximate answer for the BAU.

22

To summarize, the effects of transport in the plasma external to the expanding bubble allow baryon violating transitions in the unbroken phase to transform charge asymmetries produced on the wall into a baryon asymmetry. This means that in the case of nonlocal baryogenesis we do not need to rely on baryon number violating processes occurring in the region where the Higgs fields are changing. The diffusion equation approach to the problem of nonlocal baryogenesis has been very successful. In the thin wall case, it is a valid approximation to assume that the source for the diffusion equation is essentially a δ-function on the wall. This is because one may ignore the effects of particle scattering in this picture. However, in the case of thick walls, significant particle scattering occurs and as a result it is necessary to consider sources for the diffusion equations that extend over the wall.

A. Electroweak baryogenesis in the MSSM

As we mentioned, the most promising and well-motivated framework incorporating CP violation beyond the SM and an enhanced electroweak phase transition seems to be supersymmetry. Electroweak baryogenesis in the framework of the MSSM has attracted much attention in the past years, with particular emphasis on the strength of the phase transition [119–122] and the mechanism of baryon number generation [123,124,50,125–128].

1. The strength of the electroweak phase transition in the MSSM

The behavior of the electroweak phase transition in the minimal supersymmetric standard model is dependent on the mass of the lightest Higgs particle, and the mass of the top squark. Recent analytical [129–137] and lattice computations [138,140,139,141] have revealed that the phase transition can be sufficiently strongly first order in the presence of a top squark lighter than the top quark. In order to naturally suppress contributions to the ρ-parameter, and hence preserve a good agreement with precision electroweak measurements at LEP, the top squark should be mainly right handed. This can be achieved if the left handed stop soft supersymmetry breaking mass m_Q is much larger than M_Z. For moderate mixing, the lightest stop mass is then approximately given by

$$m_{\tilde{t}}^2 \simeq m_U^2 + m_t^2(\phi) \left(1 - \frac{\left| \tilde{A}_t \right|^2}{m_Q^2} \right) \ , \tag{4.15}$$

where $\tilde{A}_t = A_t - \mu^*/\tan\beta$ is the particular combination appearing in the off-diagonal terms of the left-right stop squared mass matrix and m_U is the soft supersymmetry breaking mass parameter of the right handed stop.

The preservation of the baryon number asymmetry requires the order parameter $\langle\phi(T_c)\rangle/T_c$ to be larger than one, see Eq. (2.18). This quantity is bounded from above

$$\frac{\langle\phi(T_c)\rangle}{T_c} < \left(\frac{\langle\phi(T_c)\rangle}{T_c} \right)_{\text{SM}} + \frac{2\,m_t^3\left(1 - \tilde{A}_t^2/m_Q^2\right)^{3/2}}{\pi\,v\,m_h^2} \ , \tag{4.16}$$

where $m_t = \overline{m}_t(m_t)$ is the on-shell running top quark mass in the $\overline{\text{MS}}$ scheme, m_h is the lightest Higgs boson mass and $v = \sqrt{v_1^2 + v_2^2}$ is the vev of the Higgs fields. The first term on the right hand side of expression (4.16) is the Standard Model contribution

$$\left(\frac{\langle \phi(T_c) \rangle}{T_c} \right)_{\text{SM}} \simeq \left(\frac{40}{m_h[\text{GeV}]} \right)^2 , \tag{4.17}$$

and the second term is the contribution that would be obtained if the right handed stop plasma mass vanished at the critical temperature (see Eq. (4.18)). The difference between the SM and the MSSM is that light stops may give a large contributions to the effective potential in the MSSM and therefore overcome the Standard Model constraints. The stop contribution strongly depends on the value of m_U^2, which must be small in magnitude, and negative, in order to induce a sufficiently strong first order phase transition. Indeed, large stop contributions are always associated with small values of the right handed stop plasma mass

$$m_{\tilde{t}}^{\text{eff}} = -\widetilde{m}_U^2 + \Pi_R(T) , \tag{4.18}$$

where $\widetilde{m}_U^2 = -m_U^2$, and

$$\Pi_R(T) \simeq 4g_3^2 T^2/9 + h_t^2/6[2 - \widetilde{A}_t^2/m_Q^2]T^2 \tag{4.19}$$

is the finite temperature self-energy contribution to the right-handed squarks. Moreover, the trilinear mass term, $\widetilde{A}_t$, must satisfy $\widetilde{A}_t^2 \ll m_Q^2$ in order to avoid the suppression of the stop contribution to $\langle \phi(T_c) \rangle / T_c$. Note that, although large values of $\widetilde{m}_U$, of order of the critical temperature, are useful to get a strongly first order phase transition, $\widetilde{m}_U$ is bounded from above to avoid the appearance of dangerous charge and color breaking minima [136].

As is clear from (4.16), in order to obtain values of $\langle \phi(T_c) \rangle / T_c$ larger than one, the Higgs mass must take small values, close to the present experimental bound. For small mixing and large m_A, the one-loop Higgs mass has a very simple form

$$m_h^2 = M_Z^2 \cos^2 2\beta + \frac{3}{4\pi^2} \frac{\overline{m}_t^4}{v^2} \log \left(\frac{m_{\tilde{t}}^2 m_{\tilde{T}}^2}{\overline{m}_t^4} \right) \left[1 + \mathcal{O} \left(\frac{\widetilde{A}_t^2}{m_Q^2} \right) \right], \tag{4.20}$$

where $m_{\tilde{T}}^2 \simeq m_Q^2 + m_t^2$. Hence, small values of $\tan \beta$ are preferred. The larger the left handed stop mass, the closer to one $\tan \beta$ must be. This implies that the left handed stop effects are likely to decouple at the critical temperature, and hence that m_Q mainly affects the baryon asymmetry through the resulting Higgs mass. A detailed analysis, including all dominant two-loop finite temperature corrections to the Higgs effective potential and the non-trivial effects arising from mixing in the stop sector, has been recently performed [136]. This analysis has identified the region of parameter space for which MSSM electroweak baryogenesis can happen. Taking into account the experimental bounds as well as the requirement of avoiding dangerous color breaking minima, it was found that the lightest Higgs should be lighter than about 105 GeV, while the stop mass may be close to the present experimental bound and must be smaller than, or of order of, the top quark mass. This lower bound has been essentially confirmed by lattice simulations [141], providing a motivation for the search for Higgs and stop particles at the LEP and Tevatron colliders.

As we have seen, the MSSM contains additional sources of CP-violation besides the CKM matrix phase. These new phases are essential for the generation of the baryon number since large CP violating sources may be locally induced by the passage of the bubble wall separating the broken from the unbroken phase during the electroweak phase transition. The new phases appear in the soft supersymmetry breaking parameters associated with the stop mixing angle and the gaugino and neutralino mass matrices. However, large values of the stop mixing angle are strongly restricted in order to preserve a sufficiently strong first order electroweak phase transition. Therefore, an acceptable baryon asymmetry may only be generated from the stop sector 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 [124]. As a result, the contribution to the final baryon asymmetry from the stop sector turns out to be negligible. On the other hand, charginos and neutralinos may be responsible for the observed baryon asymmetry if a large CP violating source for the higgsino number $\gamma_{\widetilde{H}}$ is induced in the bubble wall. This can happen for sufficiently large values of the phase of the parameter μ [124,128].

The first step in the computation of the baryon asymmetry is to solve a set of coupled differential equations describing the effects of diffusion, particle number changing reactions, and CP violating source terms [123,124,128]. Major simplifications of the diffusion equations take place when neglecting all the couplings except for gauge interactions and the top Yukawa coupling. Neglecting the weak sphalerons (in the first step) allows us to ignore leptons in the diffusion equations, and turns out to be a good approximation when computing Higgs and quark densities. A further major simplification is to assume that the top Yukawa and the strong sphaleron interactions are fast, thus reducing the diffusion equations to an equation governing the baryon asymmetry n_B. We write the integral of the CP violating source as

$$\mathcal{A} \propto \frac{1}{\overline{D} \, \lambda_+} \int_0^\infty du \, \widetilde{\gamma}_{\widetilde{H}}(u) e^{-\lambda_+ u} \ , \tag{4.21}$$

where

$$\lambda_+ = \frac{v_w + \sqrt{v_w^2 + 4\widetilde{\Gamma}\overline{D}}}{2\overline{D}}, \tag{4.22}$$

and $\widetilde{\Gamma}$ is the effective decay constant governing the particle number changing reaction [123]. Then, the solution for the baryon to entropy ratio is [123,124]

$$\frac{n_B}{s} = -g(k_i) \frac{\mathcal{A}\overline{D}\Gamma_{sp}}{v_w^2 s} \ , \tag{4.23}$$

where $g(k_i)$ is a numerical coefficient depending upon the light degrees of freedom present in the thermal bath, and $\overline{D}$ is an effective diffusion constant. CP violating densities are therefore non zero for a time $t \sim \overline{D}/v_w^2$ and the assumptions leading to the above expressions are valid as long as top Yukawa and the strong sphaleron interactions are faster than $\sim \overline{D}/v_w^2$ [123,124]. On larger time scales, diffusion takes over, washing out the sources.

The next step is the evaluation of the CP violating source for higgsino number. In this case, non-equilibrium quantum field theory provides the necessary tools to write down a set

of quantum Boltzmann equations describing the local particle densities and automatically incorporating the CP violating sources [125–127]. The closed time-path formalism is a powerful Green's function formulation for describing non-equilibrium phenomena in field theory. It leads to a complete non-equilibrium quantum kinetic theory approach, and a rigorous computation of the CP violating sources for the stop and Higgsino numbers [125–127]. CP violating sources are built up when Higgsinos scatter off the advancing bubble wall, and CP is violated at the vertices of interactions. In the classical kinetic theory the "scattering term" does not include any integral over the past history of the system. This is equivalent to assuming that any collision in the plasma does not depend upon the previous ones. In contrast, the quantum approach reveals that the CP violating source is manifestly non-Markovian.

Let us just comment on some crucial findings. First of all, since the higgsino CP violating source builds up when the charginos and neutralinos scatter off the advancing bubble wall, there appears to be a resonant behaviour associated with the fact that the Higgs background is carrying a very low momentum (of order of the inverse of the bubble wall width L_w) and with the possibility of absorption or emission of Higgs quanta by the propagating supersymmetric particles [124]. The resonance can only take place when the higgsino and the wino 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 -inos. A detailed computation of the thermalization rate of the supersymmetric particles from the imaginary part of the two-point Green function has recently been performed in [142,143] by making use of improved propagators and including resummation of hard thermal loops, showing that the damping rate of -inos in the plasma is about $0.065\,T$. Therefore, the CP violating source for the Higgs fermion number is enhanced if $M_2, M_1 \sim \mu$, and low momentum particles are transmitted over the distance L_w. This means that the classical approximation is not entirely adequate to describe the quantum interference nature of CP violation, and only a quantum approach is suitable for the computation of how the CP violating sources build up. Secondly, the source is built up by integrating over all the history of the system. This leads to "memory" effects that are responsible for some enhancement of the final baryon asymmetry. These memory effects lead to "relaxation" times for the CP violating sources which are typically longer than the ones dictated by the thermalization rates of the particles in the thermal bath.

Taking into account the detailed analysis of the bubble wall profile [144], the final baryon asymmetry (4.23) has been computed to be [126]

$$\frac{n_B}{s} \simeq \left(\frac{|\sin(\phi_\mu)|}{10^{-3}} \right) 4 \times 10^{-11}, \tag{4.24}$$

for $v_w \simeq 1$. It is intriguing that these small values of the phases are perfectly consistent with the constraints from the electric dipole moment of the neutron and squarks of the first and second generation as light as ~ 300 GeV may be tolerated.

V. AFFLECK-DINE BARYOGENESIS

We now turn to a third baryogenesis scenario, that has received a lot of attention. This mechanism was introduced by Affleck and Dine (AD) [145] and involves the cosmological

evolution of scalar fields carrying baryonic charge. As we shall see, these scenarios are naturally implemented in the context of supersymmetric models. Let us quickly describe this before presenting a detailed analysis.

Consider a colorless, electrically neutral combination of quark and lepton fields. In a supersymmetric theory this object has a scalar superpartner, χ, composed of the corresponding squark $\tilde{q}$ and slepton $\tilde{l}$ fields.

Now, as we mentioned earlier, an important feature of supersymmetric field theories is the existence of "flat directions" in field space, on which the scalar potential vanishes. Consider the case where some component of the field χ lies along a flat direction. By this we mean that there exist directions in the superpotential along which the relevant components of χ can be considered as a free massless field. At the level of renormalizable terms, flat directions are generic, but supersymmetry breaking and nonrenormalizable operators lift the flat directions and sets the scale for their potential.

During inflation it is likely that the χ field is displaced from the position $\langle \chi \rangle = 0$, establishing the initial conditions for the subsequent evolution of the field. An important role is played at this stage by baryon number violating operators in the potential $V(\chi)$, which determine the initial phase of the field. When the Hubble rate becomes of the order of the curvature of the potential $\sim m_{3/2}$, the condensate starts oscillating around its present minimum. At this time, B-violating terms in the potential are of comparable importance to the mass term, thereby imparting a substantial baryon number to the condensate. After this time, the baryon number violating operators are negligible so that, when the baryonic charge of χ is transferred to fermions through decays, the net baryon number of the universe is preserved by the subsequent cosmological evolution.

In this section we will focus on the most recent developments about AD baryogenesis, referring the reader to [13] for more details about the scenario as envisaged originally in [47].

A. Supersymmetric Implementations

The most recent implementations of the Affleck-Dine scenario have been in the context of the minimal supersymmetric standard model [146,147]. In models like the MSSM, with large numbers of fields, flat directions occur because of accidental degeneracies in field space. Although a flat direction is parameterized by a chiral superfield, we shall here focus on just the scalar component, since it is the dynamics of this that we shall be interested in. In general, flat directions carry a global $U(1)$ quantum number. We shall be interested in those directions in the MSSM that carry $B - L$, since the Affleck-Dine condensate will decay above the weak scale and we must avoid erasure of any asymmetry through rapidly occurring sphaleron processes. Although there are many such directions in the MSSM [147], it is sufficient to treat a single one as typical of the effect. As mentioned above, flat directions can be lifted by supersymmetry breaking and nonrenormalizable effects in the superpotential. An important obeservation is that inflation, by definition, breaks global supersymmetry because of a nonvanishing cosmological constant (the false vacuum energy density of the inflaton). In supergravity theories, supersymmetry breaking is transmitted by gravitational interactions and the supersymmetry breaking mass squared is naturally CH^2, where $C = \mathcal{O}(1)$ [148,149]. To illustrate this effect, consider a term in the Kähler potential of the form

$$\delta K = C \int d^4\,\theta \frac{1}{m_p^2} \chi^\dagger \chi \phi^\dagger \phi, \tag{5.1}$$

where ϕ is the field which dominates the energy density ρ of the universe, that is $\rho \simeq \langle \int d^4\theta \phi^\dagger \phi \rangle$. During inflation, ϕ is identified with the inflaton field and $\rho = V(\phi) = 3H^2 m_p^2$. The term (5.1) therefore provides an effective χ mass $m_\chi = -3CH^2$. This example can easily be generalized, and the resulting potential of the AD flat direction during inflation is of the form

$$V(\chi) = H^2 m_p^2 f(\chi/m_p) + H m_p^3 g(\chi^n/m_p^n), \tag{5.2}$$

where f and g are some functions, and the second term is the generalized A-term coming from nonrenormalizable operators in the superpotential.

Crucial in assessing the possibility of baryogenesis is the sign of the induced mass squared at $\chi = 0$. If the sign of the induced mass squared is negative, a large expecation value for a flat direction may develop, which is set by the balance with the nonrenormalizable terms in the superpotential. In such a case the salient features of the post-inflationary evolution are as follows [147]:

$i)$ Suppose that during inflation the Hubble rate is $H_I \gg m_{3/2}$ and the potential is of the form

$$V(\chi) = \left(m_{3/2} - |C|H_I^2 \right) |\chi|^2 + \left[\frac{\lambda(aH_I + Am_{3/2})\chi^n}{nM^n} + \text{h.c.} \right] + |\lambda|^2 \frac{|\chi|^{2n-2}}{M^{2n-6}}, \tag{5.3}$$

where a is a constant of order unity, and M is some large mass scale such as the GUT or the Planck scale. The AD field evolves exponentially to the minimum of the potential, $\langle |\chi| \rangle \sim (H_I M^{n-3}/\lambda)^{1/n-2}$. The A-term in (5.3) violates the $U(1)$ carried by χ, and the potential in the angular direction is proportional to $\cos(\theta_a + \theta_\lambda + n\theta)$ where $\chi = |\chi|e^{i\theta}$, and θ_a and θ_λ are the respective phases of a and λ. This creates n discrete minima for the phase of χ. During inflation the field quickly settles into one of these minima, establishing the initial phase of the field.

$ii)$ Subsequent to inflation, the minimum of the potential is time dependent since H changes with time and the AD field oscillates near the time dependent minimum $\langle |\chi| \rangle(t)$ with decreasing amplitiude.

$iii)$ When $H \sim m_{3/2}$ the sign of the mass squared becomes positive and the field begins to oscillate about $\chi = 0$ with frequency $m_{3/2}$ and amplitude $\langle |\chi| \rangle(t \sim m_{3/2}^{-1})$. At this time the B-violating A-terms are as important as the mass term, and there is no sense in which the baryon number is conserved. When the field begins to oscillate freely, a large fractional baryon number is generated during the initial spiral motion of the field. The important role of CP violation is also dictated by the A-terms, since the angular term $\cos(\theta_A + \theta_\lambda + n\theta)$ becomes important at this stage, and a nonzero $\dot{\theta}$ is generated if $\theta_a \neq \theta_A$. This is of course necessary since the baryon number is given by $n_b = 2|\chi|^2\dot{\theta}$. When the condensate decays, the baryon asymmetry is transferred to fermions.

$iv)$ In the case in which the inflaton decays when $H < m_{3/2}$ (consistently with the requirement that not too many gravitinos are produced at reheating), the resulting baryon to entropy ratio is

$$B \simeq \frac{n_b}{n_\chi} \frac{T_r}{m_{3/2}} \frac{\rho_\chi}{\rho_\phi}, \tag{5.4}$$

where $n_\chi = \rho_\chi/m_{3/2}$, $\rho_\phi \sim m_{3/2}^2 m_p^2$ is the energy density of the inflaton field when $H \sim m_{3/2}$, and T_r is the reheating temperature. For example $n = 4$ gives [147]

$$B \simeq 10^{-10} \left(\frac{T_r}{10^6 \, \text{GeV}} \right) \left(\frac{10^{-3} \, M}{\lambda m_p} \right). \tag{5.5}$$

The MSSM contains many combinations of fields for constructing flat directions carrying a nonvanshing $B - L$ [147]. Particularly appealing directions are the ones which carry $B - L$ and can be lifted at $n = 4$ level. They are the $\chi^2 = LH_2$ directions. The nonrenormalizable operator is then

$$W = \frac{\lambda}{M} (LH_2)^2, \tag{5.6}$$

which may be present directly at the Planck scale, or could be generated, as in $SO(10)$ GUTs, by integrating out right-handed neutrino superfields which are heavy standard model singlets.

B. Initial conditions for viable AD baryogenesis

¿From the above picture it is clear that a successful AD baryogenesis mechanism is achieved if the effective potential during inflation contains a negative effective mass term and nonrenormalizable terms that lift the flat directions of the potential.

With a minimal Kähler potential only, the effective mass squared $m^2 \sim H_I^2$ during inflation is positive, $\chi = 0$ is stable, and the large expectation values required for baryogenesis do not result. Quantum de Sitter fluctuations do excite the field with $\langle \chi^2 \rangle \sim H_I^2$, but with a correlation length of the order of H_I^{-1}. Any resulting baryon asymmetry then averages to zero over the present universe. In addition, the relative magnitude of the B violating A-terms in the potential is small for $H_I \ll M$.

A negative sign for the effective mass squared is possible if one either considers general Kähler potentials, or in the case in which inflation is driven by a Fayet-Iliopoulos D term, which preserves the flat directions of global supersymmetry, and in particular keeps the inflaton potential flat, making D-term inflation particularly attractive [79]. If the AD field carries a charge of the appropriate sign under the $U(1)$ whose Fayet-Iliopoulos term is responsible for inflation, a nonvanishing vacuum expectation value may develop [150]. Notice that, if the AD field is neutral under the $U(1)$, it may nonetheless acquire a negative mass squared after the end of D-term inflation through nonminimal Kähler potentials relating the AD field to the fields involved in the inflationary scenario [151].

Another (marginal) possibility is that the effective mass squared during inflation is positive, but much smaller than H_I^2, $0 < C \ll 1$. This happens either in supergravity models that possess a Heisenberg symmetry in which supersymmetry breaking by the inflationary vacuum energy does not lift flat directions at tree level [152], or in models of hybrid inflation based on orbifold constructions, in which a modulus field T is responsible for the large value

of the potential during inflation, and a second field ϕ with appropriate modular weight is responsible for the roll-over [150]. The correlation length for de Sitter fluctuations in this case is $\ell_c = H_I \exp(3H_I^2/2m^2)$. This is large compared to the horizon size. In fact, the present length corresponding to ℓ_c should be larger than the horizon size today. Using $H_I = 10^{13}$ GeV, this gives $(m/H_I)^2 \lesssim 40$, which is easily satisfied. The flat direction is truly flat and the AD baryogenesis may be implemented as originally envisaged [47]. After inflation, these truly flat directions generate a large baryon asymmetry, typically $B = \mathcal{O}(1)$. Mechanisms for suppressing this asymmetry to the observed level have been considered in [152]. These include dilution from inflaton or moduli decay, and higher dimensional operators in both GUT models and the MSSM. The observed BAU can easily be generated when one or more of these effects is present.

C. Preheating and AD baryogenesis

As we have described in subsection (III B), in the first stage of reheating after inflation, called preheating, resonance effects may lead to explosive particle production. Fluctuations of scalar fields produced at the stage of preheating are so large that they may break supersymmetry much more strongly than inflation itself. These fluctuations may lead to symmetry restoration along flat directions of the effective potential even in the theories where the usual high temperature corrections are exponentially suppressed. These nonthermal phase transitions after preheating can play a crucial role in the generation of the primordial baryon asymmetry by the AD mechanism. In particular, the baryon asymmetry may be generated at the very early stage of the evolution of the universe, at the preheating era, and not when the Hubble parameter becomes of order the gravitino mass $m_{3/2}$ [153].

Let us consider the chaotic inflation scenario where the inflaton field ϕ couples to a complex X-field with a potential of the form

$$V = \frac{M_\phi^2}{2}\,\phi^2 + g^2\,\phi^2\,|X|^2 \; . \tag{5.7}$$

After inflation, the initial energy of the inflaton field is transferred to the X-quanta in the regime of parametric resonance [87]. It is especially important that the maximum amplitude of field fluctuations produced at that stage may be sizeable, $\langle X^2 \rangle \sim M_\phi^2/g^2$ [89], or even larger. These large fluctuations occur only in the bosonic sector of the theory, thus breaking supersymmetry by a large amount at the quantum level. Suppose the field χ labelling the AD flat direction couples to the generic supermultiplet X by a renormalizable interaction of the form $\mathcal{L}_{\mathrm{int}} = h^2|\chi|^2|X|^2$. Preheating gives an additional contribution to the effective squared mass along the χ-direction beyond the contribution of the $-|C|H^2$ term discussed in the previous subsection, which is responsible for a nonvanishing vacuum expectation value $\langle \chi_0 \rangle$ of the AD field during inflation. The curvature of the effective potential becomes large and positive, and the symmetry rapidly restores for $h^2\langle X^2 \rangle > |C|\,H^2$. Thus, for $C = \mathcal{O}(1)$, symmetry becomes restored for $h^2 \gtrsim g^2$. In other words, supersymmetry breaking due to preheating is much greater than the supersymmetry breaking proportional to H unless the coupling constant h^2 is anomalously small. This leads to symmetry restoration in the AD model soon after preheating. The AD field rapidly rolls towards the origin $\chi = 0$ and makes fast oscillations about it with an initial amplitude given by $\chi \sim |\chi_0|$. The frequency of

the oscillations is of order of $h\langle X^2\rangle^{1/2}$ and is much higher than the Hubble parameter at preheating. The field is underdamped and relaxes to the origin after a few oscillations. Baryon number can be efficiently produced during this fast relaxation of the field to the origin. This picture is similar to that of ref. [47], but there are considerable differences. According to [47], the baryon asymmetry is produced at very late times when the Hubble parameter H becomes of order of the gravitino mass $m_{3/2}$. At preheating, baryon number production occurs right after inflation. This is possible because the same nonthermal effects generating a large curvature at the origin give rise also to sizable CP-violating terms along the AD flat direction [153].

D. Q-balls and AD baryogenesis

Q-balls are stable, charged solitons appearing in a scalar field theory with a spontaneously broken global $U(1)$ symmetry [154]. The Q-ball solution arises if the scalar potential $V(\phi)$ is such that $V(\phi)/|\phi|^2$ has a minimum at non-zero ϕ. In the MSSM such solitons are part of the spectrum of the theory [155]. This is due to the form of the MSSM scalar potential and the fact that squarks and sleptons carry a global $B-L$ charge. An interesting possibility for Q-ball formation is provided by the fact that there are many flat directions in the MSSM [147]. During inflation the MSSM scalar fields are free to fluctuate along these flat directions and to form condensates. This is closely related to Affleck-Dine (AD) baryogenesis [47] with the difference that the condensate is naturally unstable with respect to the formation of Q-balls [154,156,155] with a very large charge [157]. The properties of the MSSM Q-balls depend upon the scalar potential associated with the condensate, which in turn depends upon the supersymmetry breaking mechanism and on the order at which the non-renormalizable terms lift the degeneracy of the potential. If supersymmetry breaking occurs via the supergravity hidden sector [45], B-balls are allowed to form if the charge density inside the initial lump is small enough. Since this charge density scales as the inverse of the reheating temperature T_r, this can be translated into the condition [158,159]

$$T_r \gtrsim \left(\frac{m_{3/2}}{100\,\text{GeV}}\right)\,\text{GeV}. \tag{5.8}$$

On the other hand, after formation, B-balls could be dissociated by the bombardment of thermal particles, or dissolve by charge escaping from the outer layers. Both phenomena may be prevented if $T_r \lesssim (10^3 - 10^5)$ GeV. Therefore, surviving B-balls must have very large charges, $B \gtrsim 10^{14}$ [158]. Since the decay rate of the B-ball also depends on its charge, for $B \gtrsim 10^{14}$, B-balls will indeed decay well below the electroweak phase transition temperature. The only requirement is relatively low reheating temperature after inflation, typically of the order of 1 GeV, which is fixed by the observed baryon asymmetry when the CP violating phase responsible for the baryon asymmetry is of the order of unity [159]. In particular, this is expected to be true for D-term inflation models [151,160]. Such B-balls can therefore decay at temperatures less than that of the electroweak phase transition, and consequently lead to a novel version of Affleck-Dine baryogenesis, in which the baryon asymmtry asymmetry comes from B-ball decay rather than condensate decay. Since the baryon asymmetry production takes place when sphaleron processes are no longer active, this mechanism can work even in

the presence of additional L violating interactions or $B - L$ conservation, which would rule out conventional AD baryogenesis along B violating directions.

VI. CONCLUSIONS

The origin of the baryon to entropy ratio of the universe is one of the fundamental initial condition challenges of the standard model of cosmology. Over the past several decades, many mechanisms utilizing quantum field theories in the background of the expanding universe have been proposed as explanations for the observed value of the asymmetry. In this review, with space constraints, we have focused on those that we feel hold the most promise. We have given what we hope is an up to date review of three popular theories of baryogenesis; GUT scenarios, electroweak baryogenesis, and the Affleck-Dine mechanism. Each of these mechanisms has both attractive and problematic aspects, as we'll attempt to summarize here.

Grand Unified baryogenesis was the first implementation of Sakharov's baryon number generation ideas. There are at least three reasons why this is an attractive mechanism. First, baryon number violation is an unavoidable consequence of Grand Unification. If one wishes to have quarks and leptons transforming in the same representation of a single group, then there necessarily exist processes in which quarks convert to leptons and vice-versa. Hence, baryon number violation is manifest in these theories. Second, it is a simple and natural task to incorporate sufficient amounts of CP-violation into GUTs. Among other possibilities, there exist many possible mixing phases that make this possible. Finally, at such early epochs of cosmic evolution, the relevant processes are naturally out of equilibrium since the relevant timescales are slow compared to the expansion rate of the universe. The simplicity with which the Sakharov criteria are achieved in GUTs makes this mechanism very attractive.

While GUT baryogenesis is attractive, it is not likely that the physics involved will be directly testable in the foreseeable future. While we may gain indirect evidence of grand unification with a particular gauge group, direct confirmation in colliders seems unrealistic. A second problem with GUT scenarios is the issue of erasure of the asymmetry. As we have seen, the electroweak theory ensures that there are copious baryon number violating processes between the GUT and electroweak scales. These sphaleron events violate $B + L$ but conserve $B - L$. Thus, unless a GUT mechanism generates an excess $B - L$, any baryonic asymmetry produced will be equilibrated to zero by anomalous electroweak interactions. While this does not invalidate GUT scenarios, it is a constraint. For example, $SU(5)$ will not be suitable for baryogenesis for this reason, while $SO(10)$ may be. This makes the idea of baryogenesis through leptogenesis particularly attractive.

In recent years, perhaps the most widely studied scenario for generating the baryon number of the universe has been electroweak baryogenesis. The popularity of this idea is tightly bound to its testability. The physics involved is all testable in principle at realistic colliders. Furthermore, the small extensions of the model involved to make baryogenesis successful can be found in supersymmetry, which is an independently attractive idea, although electroweak baryogenesis does not depend on supersymmetry.

However, the testability of electroweak scenarios also leads to tight constraints. At present there exists only a small window of parameter space in extensions of the electroweak

theory in which baryogenesis is viable. This is because electroweak baryogenesis to be effective requires a strong enough first order phase transition. This translates into a severe upper bound on the lightest Higgs boson mass, $m_h \lesssim 105$ GeV, in the case in which the mechanism is implemented in the MSSM. Furthermore, the stop mass may be close to the present experimental bound and must be smaller than, or of order of, the top quark mass. The most direct experimental way of testing this scenario is through the search for the lightest Higgs at LEP2. In this sense, we are close to knowing whether electroweak processes were responsible for the BAU.

If the Higgs is found, the second test will come from the search for the lightest stop at the Tevatron collider (the stop mass is typically too large for the stop to be seen at LEP). If both particles are found, the last crucial test will come from B physics, more specifically, in relation to the CP-violating effects. Depending on their gaugino-Higgsino composition, light neutralinos or charginos, at the reach of LEP2, although not essential, would further restrict the allowed parameter space consistent with electroweak baryogenesis.

Moreover, the selected parameter space leads to values of the branching ratio $\mathrm{BR}(b \to s\gamma)$ different from the Standard Model case. Although the exact value of this branching ratio depends strongly on the value of the μ and A_t parameters, the typical difference with respect to the Standard Model prediction is of the order of the present experimental sensitivity and hence in principle testable in the near future. Indeed, for the typical spectrum considered here, due to the light charged Higgs, the branching ratio $\mathrm{BR}(b \to s\gamma)$ is somewhat higher than in the SM case, unless negative values of $A_t\mu$ are present.

Affleck-Dine baryogenesis is a particularly attractive scenario, and much progress has been made in understanding how this mechanism works. As was the case for electroweak baryogenesis, this scenario has found its most promising implementations in supersymmetric models, in which the necessary flat directions are abundant. Particularly attractive is the fact that these moduli, carrying the correct quantum numbers, are present even in the MSSM.

The challenges faced by Affleck-Dine models are combinations of those faced by the GUT and electroweak ideas. In particular, it is necessary that $B - L$ be violated along the relevant directions (except perhaps in the Q-ball implementations) and that there exist new physics at scales above the electroweak. If supersymmetry is not found, then it is hard to imagine how the appropriate flat directions can exist in the low energy models.

Understanding how the baryon asymmetry of the universe originated is one of the fundamental goals of modern cosmology. Particle physics has provided us with a number of possibilities, all of which involve fascinating physics, but which have varying degrees of testability. Our primary goal in this review has been to provide an as up to date as possible account of current efforts in three primary avenues of baryogenesis research. We have also tried to be as pedagogical as space allows. We hope that the reader has found this helpful, and that friends and colleagues will forgive us for any inadvertent omissions.

ACKNOWLEDGEMENTS

REFERENCES

[1] P.A.M. Dirac, Proc. Roy. Soc **A126**, 360 (1930).

[2] P.A.M. Dirac, Proc. Roy. Soc **133**, 60 (1931).

[3] S.P. Ahlen, Phys. Rev. Lett. **61**, 145 (1988).

[4] G. Steigman, Ann. Rev. Astron. Astrophys. **14**, 336 (1976).

[5] A.G. Cohen, A. De Rujula and S.L. Glashow, astro-ph/9707087 (1997).

[6] See, for instance, E.W. Kolb and M.S. Turner, *The Early Universe*, Addison-Wesley (1990).

[7] C.J. Copi, D.N. Schramm and M.S. Turner, Science **267**, 192 (1995).

[8] A. Riotto, Lectures given at the *ICTP Summer School in High-Energy Physics and Cosmology*, Miramare, Trieste, Italy, 29 Jun - 17 Jul 1998, hep-ph/9807454, (1998).

[9] M. Trodden, "Electroweak Baryogenesis", to appear in Rev. Mod. Phys. October 1999, hep-ph/9803479, (1998).

[10] V.A. Rubakov and M. E. Shaposhnikov, Phys. Usp. **39**, 461 (1996).

[11] N. Turok, "Electroweak Baryogenesis", in *Perspectives on Higgs Physics*, edited by G.L. Kane, p. 300 (1992).

[12] A.G. Cohen, D.B. Kaplan and A.E. Nelson, Ann. Rev. Nucl. Part. Sci. **43**, 27 (1993).

[13] A. Dolgov, Physics Reports, **222**, 309 (1992).

[14] A.D. Sakharov, Zh. Eksp. Teor. Fiz. Pis'ma **5**, 32 (1967); JETP Lett. **91B**, 24 (1967).

[15] For a review, see P. Langacker, Phys. Rept. **72**, 185 (1981).

[16] U. Amaldi, W. de Boer and H. Furstenau, Phys. Lett. **B260**, 447 (1991).

[17] J. Ellis, S. Kelley and D.V. Nanopoulos, Phys. Lett. **B249**, 441 (1990); **B260**, 131 (1991).

[18] P. Langacker and M.-X. Luo, Phys. Rev. **D44**, 817 (1991).

[19] For a comprehensive review, see E.W. Kolb and M.S. Turner, Ann. Rev. Nucl. Part. Sci. **33**, 645 (1983).

[20] J.C. Pati and A. Salam, Phys. Rev. **D10**, 275 (1974).

[21] R.N. Mohapatra, J.C. Pati, Phys. Rev. **D11**, 566 (1975).

[22] R.N. Mohapatra and G. Senjanovic, Phys. Rev. **D12**, 1502 (1975).

[23] M. Fukugita and T. Yanagida, Phys. Lett. **174**, 45 (1986).

[24] G. 't Hooft, Phys. Rev. Lett. **37**, 37 (1976); Phys. Rev. **D14**, 3432 (1976).

[25] S. Dimopoulos and L. Susskind, Phys. Rev. **D18**, 4500 (1978).

[26] N.S. Manton, Phys. Rev. **D28**, 2019 (1983).

[27] R.F. Klinkhamer and N.S. Manton, Phys. Rev. **D30**, 2212 (1984).

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

[29] S. Adler, Phys. Rev. **177**, 2426 (1969).

[30] J.S. Bell and R. Jackiw, 1969, Nuovo Cimento **51**, 47 (1969).

[31] L. Carson et *al.*, Phys. Rev. **D42**, 2127 (1990).

[32] G.D. Moore, hep-ph/9801204.

[33] S. Yu Khlebnikov and M. Shaposhnikov, Nucl. Phys. **B308**, 885 (1988).

[34] J. Ambjørn, T. Askgaard, H. Porter and M. Shaposhnikov, Phys. Lett. **B244**, 479 (1990).

[35] J. Ambjørn, T. Askgaard, H. Porter and M. Shaposhnikov, Nucl. Phys. **B353**, 346 (1991).

[36] J. Ambjørn and A. Krasnitz, 1995, Phys. Lett. **B362**, 97 (1995).

[37] G.D. Moore, Nucl. Phys. **B480**, 657 (1996).

[38] W.-H. Tang and J. Smit, 1996, Nucl. Phys. **B482**, 265 (1996).

[39] G.D. Moore and N. Turok, 1997, Phys. Rev. **D55**, 6538 (1997).

[40] P. Arnold, D. Son and L.G. Yaffe, Phys. Rev. **D55**, 6264 (1997).

[41] P. Arnold, Phys. Rev. **D55**, 7781 (1997); *ibidem* hep-ph/9706305.

[42] D. Bodeker, Phys. Lett. **B426**, 351 (1998).

[43] P. Arnold, D.T. Son and L.G. Yaffe, hep-ph/9810216.

[44] G.D. Moore, C. Hu and B. Muller, hep-ph/9710436.

[45] H. P. Nilles, Phys. Rep. **110**, 1 (1984).

[46] H.E. Haber and G.L. Kane, Phys. Rept. **117**, 75 (1985).

[47] I. Affleck and M. Dine, Nucl. Phys. B **249** (1985) 361.

[48] J.A. Harvey, E.W. Kolb, D.B. Reiss and S. Wolfram, Nucl. Phys. **B201**, 16 (1982).

[49] D. Comelli, M. Pietroni and A. Riotto, Nucl. Phys. **B412**, 441 (1994); Phys. Rev. **D50**, 7703 (1994).

[50] A. Riotto, hep-ph/9709286.

[51] J.F. Gunion, H.E. Haber, G.L. Kane and S. Dawson, 1989, *The Higgs Hunter's Guide*, (Addison-Wesley).

[52] M. Laine and K. Rummukainen, hep-ph/9811369.

[53] M. Gleiser, Phys. Rev. **D49**, 2978 (1993).

[54] M. Gleiser, Phys. Rev. Lett. **73**, 3495 (1994).

[55] M. Gleiser and R. Ramos, Phys. Rev. **D50**, 2441 (1994).

[56] M. Gleiser, Phys. Rev. Lett. **75**, 589 (1995).

[57] J. Borrill and M. Gleiser, Phys. Rev. **D51**, 4111 (1995).

[58] M. Gleiser and A.F. Heckler, Phys. Rev. Lett. **76**, 180 (1996).

[59] J. Borrill and M. Gleiser, Nucl. Phys. **B483**, 416 (1997).

[60] M. Gleiser, A.F. Heckler and E.W. Kolb, Phys. Lett. **B405**, 121 (1997).

[61] R. H. Brandenberger, A.-C. Davis, and M. Trodden, *Phys. Lett.* **B335**, 123 (1994).

[62] R. H. Brandenberger, A.-C. Davis, T. Prokopec and M. Trodden, *Phys. Rev.* **D53**, 4257 (1996).

[63] T. Prokopec, R. H. Brandenberger, A.-C. Davis and M. Trodden, *Phys. Lett.* **B384**, 175 (1996).

[64] M. Trodden, A.-C. Davis and R. H. Brandenberger, *Phys. Lett.* **B349**, 131 (1995).

[65] R.H. Brandenberger and A.-C. Davis, Phys. Lett. **B308**, 79 (1993).

[66] R.H. Brandenberger, A.-C. Davis and M. Hindmarsh, Phys. Lett. **B263**, 239 (1991).

[67] W. Perkins and A.-C. Davis, *Nucl. Phys.* **B406**, 377 (1993).

[68] J. Ambjørn, N. Nielsen and P. Olesen, Nucl. Phys. **B310**, 625 (1988).

[69] M. Trodden, *Mod. Phys. Lett.* **A9**, 2649 (1994).

[70] J. Cline , J. Espinosa , G.D. Moore and A. Riotto, hep-ph/9810261.

[71] A. Riotto, Phys. Lett. **B413**, 22 (1997).

[72] R. Brandenberger and A. Riotto, hep-ph/9801448.

[73] E.W. Kolb and S. Wolfram, Nucl. Phys. **B172**, 224 (1980); Phys. Lett. **B91**, 217 (1980).

[74] J.N. Fry, K.A. Olive and M.S. Turner, Phys. Rev. **D22**, 2953 (1980); Phys. Rev. **D22**, 2977 (1980); Phys. Rev. Lett. **45**, 2074 (1980).

[75] J.A. Harvey, E.W. Kolb, D.B. Reiss and S. Wolfram, Nucl. Phys. **B201**, 16 (1982).

[76] Guth, A. H., Phys. Rev. **D23**, 347 (1981).

[77] A. D. Linde, Phys. Lett. **108B** (1982).

[78] A. Albrecht and P. J. Steinhardt, Phys. Rev. Lett. **48**, 1220 (1982) .

[79] For a recent review, see D.H. Lyth and A. Riotto, hep-ph/9807278, to appear in Phys. Rept.

[80] A. D. Dolgov and A. D. Linde, Phys. Lett. **B116**, 329 (1982).

[81] L. F. Abbott, E. Farhi and M. Wise, Phys. Lett. **B117**, 29 (1982).

[82] D.J.H. Chung, E.W. Kolb and A. Riotto, Phys. Rev. **D59**, 23501 (1999), hep-ph/9802238.

[83] D.J.H. Chung, E.W. Kolb and A. Riotto, Phys. Rev. Lett. **81**, 4048 (1998), hep-ph/9805473.

[84] D.J.H. Chung, E.W. Kolb and A. Riotto, hep-ph/9809453.

[85] D.J.H. Chung, E.W. Kolb and A. Riotto, hep-ph/9810361.

[86] J. Ellis, J. Kim and D.V. Nanopoulos, Phys. Lett. **B145**, 181 (1984).

[87] L. A. Kofman, A. D. Linde and A. A. Starobinsky, Phys. Rev. Lett. **73**, 3195 (1994);

[88] S.Yu. Khlebnikov and I.I. Tkachev, Phys. Rev. Lett. **77**, 219 (1996).

[89] S.Yu. Khlebnikov and I. Tkachev, Phys. Lett. **B390**, 80 (1997).

[90] S.Yu. Khlebnikov and I. Tkachev, Phys. Rev. Lett. **79**, 1607 (1997).

[91] S.Yu. Khlebnikov and I. Tkachev, Phys. Rev. **D56**, 653 (1997).

[92] E.W. Kolb, A.D. Linde and A. Riotto, Phys. Rev. Lett. **77**, 4290 (1996), hep-ph/9606260.

[93] E.W. Kolb, A. Riotto and I.I. Tkachev, Phys. Lett. **B423**, 348 (1998), hep-ph/9801306.

[94] L. Kofman, astro-ph/9802285.

[95] M. Gell-Mann, P. Ramond and R. Slanski, in *Supergravity*, ed. P. Van Nieuwenhuizen and D.Z. Freedman, North Holland (1979).

[96] P. Langacker, R.D. Peccei and T. Yanagida, Mod. Phys. Lett. **A1**, 541 (1986).

[97] M. A. Luty, Phys. Rev. **D45**, 455 (1992).

[98] C.E. Vayonakis, Phys. Lett. **B286**, 92 (1992).

[99] L. Covi, N. Rius, E. Roulet and F. Vissani, Phys. Rev. **D57**, 93 (1998).

[100] For a comprehensive, review on baryogenesis through leptogenesis, see A. Pilaftsis, hep-ph/9812256.

[101] W. Buchmuller and M. Plumacher, Phys. Lett. **B389**, 73 (1996).

[102] L. Wolfenstein, Phys. Rev. **D17**, 2369 (1978).

[103] S.P. Mikheyev and A.Y. Smirnov, Nuovo Cim. **9C**, 17 (1986).

[104] A. Lue, K. Rajagopal and M. Trodden, Phys. Rev. **D56**, 1250 (1997)

[105] M. Joyce, T. Prokopec and N. Turok, Phys. Lett. **B339**, 312 (1994).

[106] A.G. Cohen, D. B. Kaplan and A. E. Nelson, Phys. Lett. **B245**, 561 (1990).

[107] A.G. Cohen, D. B. Kaplan and A. E. Nelson, Nucl. Phys. **B349**, 727 (1991).

[108] M. Joyce, T. Prokopec and N. Turok, Phys. Lett. **B338**, 269 (1994).

[109] R.N. Mohapatra and X. Zhang, Phys. Rev. **D45**, 2699 (1992).

[110] G.F. Giudice and M. E. Shaposhnikov, Phys. Lett. **B326**, 118 (1994).

[111] G.D. Moore, Phys. Lett. **B412**, 359 (1997).

[112] A.E. Nelson, D. B. Kaplan and A. G. Cohen, Nucl. Phys. **B373**, 453 (1992).

[113] D. Comelli, M. Pietroni and A. Riotto, Phys. Lett. **B354**, 91 (1995).

[114] D. Comelli, M. Pietroni and A. Riotto, Phys. Rev. **D53**, 4668 (1996).

[115] A.G. Cohen, D. B. Kaplan and A. E. Nelson, Phys. Lett. **B336**, 41 (1994).

[116] D. Comelli, M. Pietroni and A. Riotto, hep-ph/9406369.

[117] D. Comelli, M. Pietroni and A. Riotto, Astropart. Phys. **4**, 71 (1995).

[118] M. Joyce, T. Prokopec and N. Turok, Phys. Rev. Lett. **75**, 1695; Erratum-ibid **75**, 3375 (1994).

[119] G.F. Giudice, Phys. Rev. **D45**, 3177 (1992).

[120] S. Myint, Phys. Lett. **B287**, 325 (1992).

[121] J.R. Espinosa, M. Quirós and F. Zwirner, Phys. Lett. **B307**, 106 (1993).

[122] A. Brignole, J.R. Espinosa, M. Quirós and F. Zwirner, Phys. Lett. **B324**, 181 (1994).

[123] P. Huet and A.E. Nelson, Phys. Rev. **D53**, 4578 (1996).

[124] M. Carena, M. Quiros, A. Riotto, I. Vilja and C.E.M. Wagner, Nucl. Phys. **B503**, 387 (1997), [hep-ph/9702409].

[125] A. Riotto, Nucl. Phys. **B518**, 339 (1998), hep-ph/9712221.

[126] A. Riotto, hep-ph/9803357, to appear in Phys. Rev. **D**.

[127] A. Riotto, hep-ph/9802240.

[128] J.M. Cline, M. Joyce and K. Kainulainen, Phys.Lett. **B417**, 79 (1998).

[129] M. Carena, M. Quiros and C.E.M. Wagner, Phys. Lett. **B380**, 81 (1996).

[130] D. Delepine, J.M. Gerard, R. Gonzalez Felipe and J. Weyers, Phys. Lett. **B386**, 183 (1996).

[131] J.R. Espinosa, Nucl. Phys. **B475**, 273 (1996).

[132] D. Bodeker, P. John, M. Laine, M.G. Schmidt, Nucl. Phys. **B497**, 387 (1997).

[133] M. Losada, Phys. Rev. **D56**, 2893 (1997).

[134] G.R. Farrar and M. Losada, Phys. Lett. **B406**, 60 (1997).

[135] B. de Carlos and J.R. Espinosa, Nucl. Phys. **B503**, 24 (1997).

[136] M. Carena, M. Quiros and C.E.M. Wagner, CERN-TH/97-190, Nucl. Phys. **B524**, 3 (1998).

[137] M. Losada, hep-ph/9806519..

[138] M. Laine, Nucl. Phys. **481**, 43 (1996).

[139] M. Laine and K. Rummukainen, Phys. Rev. Lett. **80**, 5259 (1998).

[140] J.M. Cline and K. Kainulainen, Nucl. Phys. **B482**, 73 (1996).

[141] M. Laine and K. Rummukainen, hep-lat/9804019.

[142] K. Enqvist, A. Riotto and I. Vilja, hep-ph/9710373.

[143] P. Elmfors, K. Enqvist, A. Riotto and I. Vilja, hep-ph/9809529.

[144] J.M. Moreno, M. Quiros and M. Seco, hep-ph/9801272.

[145] I. Affleck and M. Dine, Nucl. Phys. **B249**, 361 (1985).

[146] M. Dine, L. Randall and S. Thomas, Phys. Rev. Lett. **75**, 398 (1995).

[147] M. Dine, L. Randall and S. Thomas, Nucl. Phys. **B458**, 291 (1996).

[148] G. D. Coughlan, R. Holman, P. Ramond and G. G. Ross, Phys. Lett. **140B**, 44 (1984).

[149] M. Dine, W. Fischler and D. Nemeschansky, Phys. Lett. **136B**, 169 (1984).

[150] A. Casas and G.B. Gelmini, Phys. Lett. **B410**, 36 (1997).

[151] C. Kolda and J. March-Russell, hep-ph/9802358.

[152] B.A. Campbell, M.K. Gaillard, H. Murayama and K.A. Olive, hep-ph/9805300.

[153] G.W. Anderson, A. Linde and A. Riotto, Phys. Rev. Lett. **77**, 3716 (1996).

[154] S. Coleman, Nucl. Phys. **B262**, 263 (1985).

[155] A. Kusenko, Phys. Lett. **B405**, 108 (1997).

[156] A. Cohen, S. Coleman, H. Georgi and A. Manohar, Nucl. Phys **B272**, 301 (1986).

[157] A. Kusenko and M. Shaposhnikov, Phys. Lett. **B418**, 104 (1998).

[158] K. Enqvist and J. McDonald, Phys. Lett. **B425**, 309 (1998), hep-ph/9711514.

[159] K. Enqvist and J. McDonald, hep-ph/9803380.

[160] K. Enqvist and J. McDonald, hep-ph/9806213.