ASTR 688: Term Paper

Baryogenesis

P. S. Bhupal Dev

Department of Physics, University of Maryland, College Park, MD 20742, USA

13 May 2008

Abstract

Baryogenesis is a typical example of interplay between particle physics and cosmology. This report provides a brief review of the present status of theories explaining the observed baryon asymmetry of the Universe. The main goal of all these theories is to understand how well-motivated particle physics theories in a cosmological setting might generate the observed value of baryon asymmetry while remaining consistent with all other particle physics and cosmological data. All models are closely tied to the nature of dark matter, especially in supersymmetric theories. In the near future, results from LHC and gamma-ray astronomy are expected to shed new light on the theory of baryogenesis.

1 Introduction

According to the CPT theorem in Quantum Field Theory, for any particle species X, there exists the antiparticle $\overline{X}$ with exactly the same mass, $m_X = m_{\overline{X}}$, decay width, $\Gamma_X = \Gamma_{\overline{X}}$, and opposite charges, $Q_X = -Q_{\overline{X}}$. This fundamental symmetry between particles and antiparticles, firmly established in particle accelerator experiments, would naturally lead us to conclude that the Universe should contain particles and antiparticles in equal number densities, $n_X = n_{\overline{X}}$. The observed Universe, however, is composed almost entirely of matter with little primordial antimatter. There is no observational evidence for primordial antimatter on scales certainly of our galaxy ($\sim$ 3 Kpc) and most probably on scales of clusters ($\sim$ 10 Mpc) [1]. If matter and antimatter were to coexist in clusters of galaxies, we would expect a detectable background of γ-rays from nucleon-antinucleon annihilations within the clusters. A much more severe bound on the observable scale ($\sim$ 300 Mpc) is potentially reachable by more powerful future γ-ray detectors (e.g. GLAST). Another signature for the presence of large domains of antimatter would be the distortion of the CMBR, which has never been observed. A numerical analysis of this problem [2] demonstrates that the Universe must consist entirely of matter on all scales up to the Hubble size.

2 The Matter-Antimatter Asymmetry Parameter

Although 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, n_b, does not remain constant during the evolution of the Universe; instead, it scales like a^{-3}, where a is the cosmological scale factor. It is therefore convenient to define the baryon asymmetry of the Universe in terms of the quantity

$$\eta \equiv \frac{n_B}{n_\gamma}, \tag{1}$$

where $n_B = n_b - n_{\bar{b}}$ is the difference between the number of baryons and antibaryons per unit volume, and $n_\gamma = \frac{2\zeta(3)}{\pi^2}T^3$ is the photon number density at temperature T. The parameter η is essential for determining the present cosmological abundances of all the light elements (viz. H, ^{3}He, ^{4}He, D, B and ^{7}Li) produced at the nucleosynthesis epoch [3]. The parameter η has been constant since nucleosynthesis.

Historically, η was determined using big bang nucleosynthesis. The theoretical predictions and experimental measurements are summarized in Figure 1 [4]. The boxes represent the regions which are consistent with experimental determinations, showing that different measurements obtain somewhat different values. The smallest error bars are for the Deuterium abundance, giving

$$\eta = 10^{-10} \times \begin{cases} 6.28 \pm 0.35 \\ 5.92 \pm 0.56 \end{cases} \tag{2}$$

These values are consistent with the ^{4}He abundance, though marginally inconsistent with ^{7}Li (the "Lithium problem").

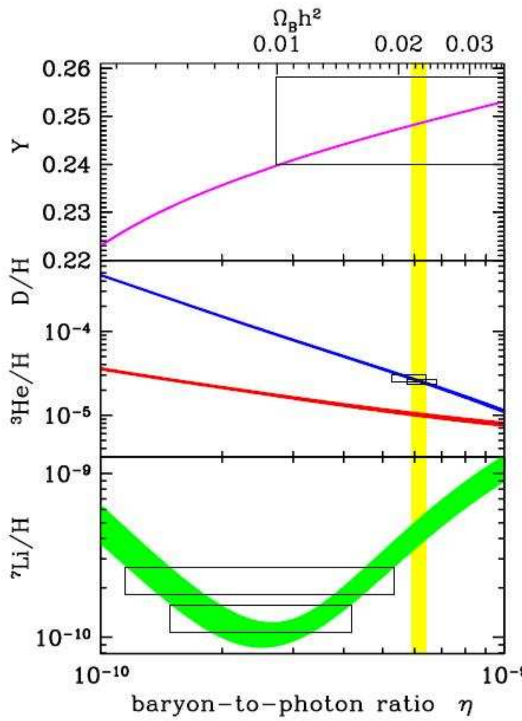

Figure 1: Primordial abundances of light elements versus the η parameter.

In the last few years, the CMB anisotropy has given us an independent way of measuring the baryon asymmetry [5]. As illustrated in Figure 2, the relative sizes of the Doppler peaks of the temperature anisotropy are sensitive to η. The WMAP data fixed the combination

$$\Omega_B h^2 = 0.0224 \pm 0.0009, \tag{3}$$

corresponding to [4]

$$\eta = (6.14 \pm 0.25) \times 10^{-10} \tag{4}$$

which is more accurate than the nucleosynthesis determination. Apart from the Lithium problem, this value is seen to be in good agreement with the value given in Eq (2). The CMB-allowed range is shown as the vertical band in Figure 1.

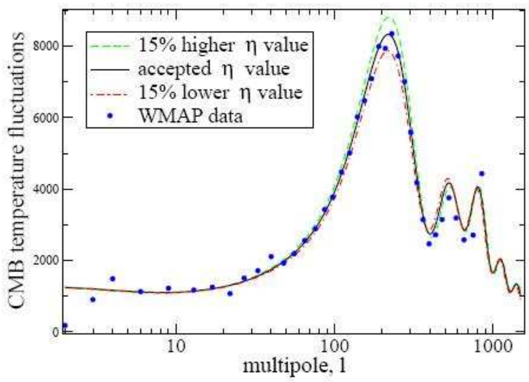

Figure 2: Dependence of the CMB Doppler peaks on η.

To see that the standard cosmological model cannot explain this observed value of η [6], suppose we start initially with $\eta = 0$. We can compute the final number density of nucleons b that are left over after annihilations have frozen out. At temperatures $T \lesssim m_p \sim 1$ GeV, the equilibrium abundance of nucleons and antinucleons is [3]

$$\frac{n_b}{n_\gamma} \simeq \frac{n_{\overline{b}}}{n_\gamma} \simeq \left(\frac{m_p}{T}\right)^{3/2} e^{-\frac{m_p}{T}} \tag{5}$$

When the Universe cools off, the number of nucleons and antinucleons decreases as long as the annihilation rate

$$\Gamma_{\rm ann} \simeq n_b \langle \sigma_A v \rangle \tag{6}$$

(where v is the velocity of the particles and σ_A is the nucleon-antinucleon annihilation cross section) is larger than the expansion rate of the Universe

$$H \simeq 1.66\ g_*^{1/2} \frac{T^2}{M_{\rm Pl}} \tag{7}$$

where g_* is the number of massless degrees of freedom. The thermally averaged annihilation cross section $\langle \sigma_A v \rangle$ is of the order of $1/m_\pi^2$ with $v \simeq c = 1$ (where m_π is the pion mass). Moreover, for nucleons (p, n) and antinucleons, $g_* = 9$ [7]. Hence the freeze out temperature is obtained by setting

$$\Gamma_{\rm ann} \simeq H, \quad \text{or} \quad T_* \simeq 20 \text{ MeV}, \tag{8}$$

below which the nucleons and antinucleons are so rare that they cannot annihilate any longer. Therefore, from Eq (5), we obtain

$$\frac{n_b}{n_\gamma} \simeq \frac{n_{\overline{b}}}{n_\gamma} \simeq 10^{-18} \tag{9}$$

which is much smaller than the value required by nucleosynthesis, cf. Eq (2). In order to avoid the annihilation catastrophe, we may suppose that some hypothetical new interactions separated matter from antimatter before $t \simeq 38$ MeV, when $\eta \simeq 10^{-10}$. At that time, $t \simeq 10^{-3}$ sec, however, the horizon was too small to explain the asymmetry over galactic scales today. This problem can be solved by invoking inflation; however, these scenarios pose other serious cosmological drawbacks. If the processes responsible for the separation of matter from antimatter took place before inflation, then the baryon number would be diluted by an enormous factor $\sim e^{200}$, because of inflation. On the other hand, if the separation took place after inflation, then it is not clear how to eliminate the boundaries separating matter and antimatter domains up to the observed scales today.

Another possible explanation for the value of η may be provided via statistical fluctuations in the baryon and antibaryon distributions. However, from pure statistical fluctuations one may expect an asymmetry $\eta \simeq 10^{-39.5}$ [8], which is again far too small to explain the observed asymmetry.

Thus, the standard cosmological model provides no explanation for the required value of η, 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 is to attempt to explain the initial conditions required by the standard cosmology on the basis of quantum field theories in the early Universe. The *baryogenesis* scenario provides a very attractive means by which the observed value of the crucial parameter η could arise dynamically in a Universe that is initially baryon symmetric. This report is a brief review of the present status of various baryogenesis models [9].

3 The Sakharov Criteria

It has been suggested by Sakharov long ago [10] that a small baryon asymmetry η may have been produced in the early Universe if the following three necessary conditions are satisfied:

1. Baryon number (B) violation,
2. C and CP violation,
3. Departure from thermal equilibrium

The first condition is obvious because, starting from a baryon-symmetric Universe with $\eta = 0$, baryon-number violating processes (with $\Delta B \neq 0$) must take place in order to evolve into a Universe where $\eta \neq 0$. Grand Unified Theories (GUTs) predict the existence of B-violating interactions which also mediate proton decay; in such a case phenomenological constraints are provided by the lower bound on the proton lifetime, $\tau_p \gtrsim 5 \times 10^{32}$ years.

The second condition is required because, even if B is violated, one can never establish baryon-antibaryon asymmetry unless C and CP are violated. This can be seen as follows: Were C an exact symmetry, the probability of the process $i \rightarrow f$ would be exactly equal to that of $\overline{i} \rightarrow \overline{f}$. Since $B(f) = -B(\overline{f})$, the net baryon number B would vanish. Furthermore, because of the CPT theorem, CP invariance implies time-reversal invariance which, in turn, assures that the rate of the process $i(\mathbf{r}_i, \mathbf{p}_i, \mathbf{s}_i) \rightarrow f(\mathbf{r}_j, \mathbf{p}_j, \mathbf{s}_j)$ and that of its time-reversed process $f(\mathbf{r}_j, -\mathbf{p}_j, -\mathbf{s}_j) \rightarrow i(\mathbf{r}_i, -\mathbf{p}_i, -\mathbf{s}_i)$ are equal. Thus, even though it is possible to create a baryon asymmetry in a certain region of phase space, integrating over all momenta $\mathbf{p}$ and summing over all spins $\mathbf{s}$ would produce a vanishing baryon asymmetry unless CP is violated. C is maximally violated by weak interactions. CP violation occurs either if there are complex phases in the Lagrangian that cannot be reabsorbed by field redefinitions (explicit breaking) or if some Higgs scalar field acquires a complex vacuum expectation value (spontaneous breaking). CP violation has been observed in K^0-$\overline{K}^0$ and B^0-$\overline{B}^0$ systems. However, a fundamental understanding of the origin of CP violation is still lacking. Hopefully, studies of baryogenesis may shed some light on it.

The third condition is essential for a net nonzero baryon asymmetry because, the equilibrium average of B vanishes:

$$\begin{aligned} \langle B \rangle_T &= \mathrm{Tr}(e^{-\beta H} B) = \mathrm{Tr}[(CPT)(CPT)^{-1} e^{-\beta H} B] \\ &= \mathrm{Tr}[e^{-\beta H} (CPT)^{-1} B (CPT)] = -\mathrm{Tr}(e^{-\beta H} B) = -\langle B \rangle_T \end{aligned} \tag{10}$$

where $\beta = 1/T$. We have used the facts that the Hamiltonian H commutes with CPT, while B is odd under CPT (odd under C, even under P and T). Thus $\langle B \rangle_T = 0$ in thermal equilibrium. Hence to establish baryon asymmetry dynamically, B violating processes must be out-of-equilibrium in the Universe.

Of the three Sakharov conditions, B violation and C and CP violation may be investigated thoroughly only within a given particle physics model [9, 11], whereas the third condition – the departure from thermal equilibrium – may be discussed in a more general way.

4 Departure from Thermal Equilibrium

In this context, the various models of baryogenesis proposed so far roughly fall into two categories:

- The *standard out-of-equilibrium decay scenario* [3, 12] in which the third Sakharov condition is satisfied owing to the presence of heavy decaying particles in a rapidly expanding Universe.
- The *phase transition scenario* [11, 13] in which the departure from thermal equilibrium is attained during the phase transitions leading to the spontaneous breaking of some global and/or gauge symmetry.

Let us first analyze the **standard out-of-equilibrium decay scenario** in which the underlying idea is fairly simpler than the phase transition approach.

It is obvious that in a static Universe any particle, even very weakly interacting, will eventually attain thermal equilibrium with the surrounding plasma. The expansion of the Universe, however, introduces a finite time-scale, $\tau_H \sim H^{-1}$. Suppose X is a *baryon number violating* superheavy boson (vector or scalar) field which is coupled to light fermionic fields with a strength $\alpha_X^{1/2}$ (either a gauge coupling or a Yukawa coupling). For renormalizable couplings (no negative mass dimension for α_X), the decay rate Γ_X of the superheavy boson is given by

$$\Gamma_X \sim \alpha_X M_X \tag{11}$$

On the other hand, if the X boson is a gauge singlet field coupling to light matter only through gravitational interactions (e.g. singlets in the hidden sector of supergravity models [14]), then from dimensional arguments, the decay rate is

$$\Gamma_X \sim \frac{M_X^3}{M_{\rm Pl}^2} \tag{12}$$

At very large temperatures $T \gg M_X$, it is assumed that all the particle species are in thermal equilibrium, i.e. $n_X \simeq n_{\overline{X}} \simeq n_\gamma$ and the net baryon number $B = 0$. At $T \lesssim M_X$ the equilibrium abundance of X and $\overline{X}$ relative to photons is given by

$$\frac{n_X^{\rm EQ}}{n_\gamma} \simeq \frac{n_{\overline{X}}^{\rm EQ}}{n_\gamma} \simeq \left(\frac{M_X}{T}\right)^{3/2} e^{-\frac{M_X}{T}} \tag{13}$$

In order to maintain their equilibrium abundances, the X and $\overline{X}$ particles must be able to diminish their number rapidly with respect to the Hubble rate $H(T)$. It is therefore useful to define the quantity

$$K \equiv \left.\frac{\Gamma_X}{H}\right|_{T=M_X} \tag{14}$$

which measures the effectiveness of decays at the crucial epoch ($T \sim M_X$) when the X and $\overline{X}$ particles must decrease in number if they are to stay in equilibrium. If $K \gg 1$, and therefore,

$$\Gamma_X \gg H|_{T=M_X}, \tag{15}$$

then the X and $\overline{X}$ particles will adjust their abundances by decaying to their equilibrium abundances and no baryogenesis can be introduced simply because the out-of-equilibrium condition is not satisfied. Using Eq (7) for the expansion rate of the Universe, the condition (15) is equivalent to

$$M_X \ll g_*^{-1/2} \alpha_X M_{\rm Pl} \tag{16}$$

for strongly coupled scalar bosons, and to

$$M_X \gg g_*^{1/2} M_{\rm Pl} \tag{17}$$

for gravitationally coupled X particles. Obviously, this last condition is never satisfied for $M_X \lesssim M_{\rm Pl}$.

However, if the decay rate is such that $K \lesssim 1$, and therefore

$$\Gamma_X \lesssim H|_{T=M_X}, \tag{18}$$

then the X and $\overline{X}$ particles cannot decay on the expansion time scale τ_H 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 much larger than the equilibrium one. This overabundance is precisely the departure from thermal equilibrium needed to produce a final nonvanishing baryon asymmetry when the heavy bosons X undergo CP violating weak decays. The condition (18) is equivalent to

$$M_X \gtrsim g_*^{-1/2} \alpha_X M_{\rm Pl} \tag{19}$$

for strongly coupled scalars, and to

$$M_X \lesssim g_*^{1/2} M_{\rm Pl} \tag{20}$$

for gravitationally coupled X particles. This last condition is always satisfied, whereas the condition (19) is based on the smallness of the quantity $g_*^{-1/2}\alpha_X$. In particular, if X is a gauge boson, $\alpha_X = \alpha_{\rm gauge}$ can span the range $(2.5 \times 10^{-2} - 10^{-1})$, while g_* is about 10^2. Hence the out-of-equilibrium condition (19) can be satisfied for

$$M_X \gtrsim (10^{-4} - 10^{-3}) M_{\rm Pl} \simeq (10^{15} - 10^{16}) \text{ GeV} \tag{21}$$

Instead, if X is a scalar boson, its Yukawa coupling, α_Y, to light fermions, f, is proportional to m_f^2:

$$\alpha_Y \sim \left(\frac{m_f}{m_W}\right)^2 \alpha_{\rm gauge} \tag{22}$$

and is typically in the range $(10^{-2} - 10^{-7})$, so that the condition (19) is satisfied for

$$M_X \gtrsim (10^{-8} - 10^{-3}) M_{\rm Pl} \simeq (10^{10} - 10^{16}) \text{ GeV} \tag{23}$$

Thus we conclude that baryogenesis is more easily produced through the decay of superheavy scalar bosons whereas the out-of-equilibrium condition is automatically satisfied for gravitationally interacting particles.

When the Universe becomes as old as the lifetime of the X and $\overline{X}$ particles, $t \sim H^{-1} \sim \Gamma_X^{-1}$, they start decaying. This takes place at a temperature T_D defined by the condition

$$\Gamma_X \simeq H|_{T=T_D} \tag{24}$$

Using Eqs (7) and (11) we get for particles with unsuppressed couplings

$$T_D \simeq g_*^{-1/4} \alpha_X^{1/2} (M_X M_{\rm Pl})^{1/2} < M_X \tag{25}$$

whereas for particles with only gravitational interactions, cf. Eq (12),

$$T_D \simeq g_*^{-1/4} M_X \left(\frac{M_X}{M_{\rm Pl}}\right)^{1/2} < M_X \tag{26}$$

Note that the expressions (19) and (20) imply that $T_D < M_X$ in both the cases.

To illustrate the mechanism of baryogenesis, let us assume that the X particle can decay into two channels, say a and b, with different baryon numbers B_a and B_b, respectively (e.g. $X \to qq$ and $X \to \overline{q}\overline{l}$ final states [3]). CPT invariance plus unitarity require the equality of the decay rates of the X and $\overline{X}$ particles; C and CP are violated if the branching ratio of $X \to a$ is unequal to the branching ratio

Decay Channel	Branching ratio	Baryon Number
$X \to a$	r	B_a
$X \to b$	$1-r$	B_b
$\overline{X} \to \overline{a}$	$\overline{r}$	$-B_a$
$\overline{X} \to \overline{b}$	$1-\overline{r}$	$-B_b$

Table 1: Final states and branching ratios for X, $\overline{X}$ decays.

of $\overline{X} \to \overline{a}$, i.e. $r \neq \overline{r}$ (see Table 1). The average net baryon number produced in the X decays is $rB_a + (1-r)B_b$, and that produced by $\overline{X}$ decays is $-\overline{r}B_a - (1-\overline{r})B_b$. Hence the mean net baryon number produced in X and $\overline{X}$ decays is

$$\Delta B = (r - \overline{r})B_a + [(1-r) - (1-\overline{r})]B_b = (r-\overline{r})(B_a - B_b) \tag{27}$$

Eq (27) can be easily generalized to the case in which $X(\overline{X})$ decays into a set of final states $f_n(\overline{f}_n)$ with baryon numbers $B_n(-B_n)$:

$$\Delta B = \frac{1}{\Gamma_X}\sum_n B_n \left[\Gamma(X \to f_n) - \Gamma(\overline{X} \to \overline{f}_n)\right] \tag{28}$$

At $T \simeq T_D$, the net baryon number density produced by the out-of-equilibrium decay is

$$n_B = \Delta B n_X, \tag{29}$$

and hence ΔB coincides with the parameter η defined in Eq (1) if $n_X \simeq n_{\overline{X}} \simeq n_\gamma$. From Eq (27) we see that only tiny C and CP violations are required to generate $\eta \sim 10^{-10}$.

We have considered only the cases $K \gg 1$ and $K \ll 1$; the intermediate regime, $K \sim 1$, is more interesting but one needs to invoke numerical analysis involving Boltzmann equations for the evolution of B [12].

5 GUT Baryogenesis

Baryon number violation is natural in GUTs because quarks and leptons lie in the same irreducible representation of the gauge group and superheavy gauge bosons mediate B-changing processes. However, the out-of-equilibrium scenario depicted in the previous section is operative only if 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 the Grand Unification, because they might never have been in thermal equilibrium in the early Universe. Secondly, the temperature of the Universe might always have been smaller than $M_{\rm GUT}$ and, correspondingly, the thermally produced X bosons might never have been as abundant as photons, making their role in baryogenesis negligible. This depends crucially upon the thermal history of the Universe. Of particular interest is a quantity known as the *reheating*[1] temperature, $T_{\rm RH}$. 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 the Hubble expansion rate, H. The result is

$$T_{\rm RH} \simeq \sqrt{\Gamma_\phi M_{\rm Pl}} \tag{30}$$

In the simplest chaotic inflation model, the inflaton potential is

$$V(\phi) = \frac{1}{2}M_\phi^2\phi^2, \quad \text{with} \quad M_\phi \sim 10^{13}\ \text{GeV} \tag{31}$$

[1]The process by which the energy of the inflaton field is transferred from the inflaton field to radiation is known as reheating.

in order to reproduce the observed CMB anisotropies [15]. Writing $\Gamma_\phi = \alpha_\phi M_\phi$, one finds

$$T_{\rm RH} \simeq 10^{16}\sqrt{\alpha_\phi}\ {\rm GeV} \tag{32}$$

There are several problems with GUT baryogenesis in this old theory of reheating. First, the density and temperature fluctuations observed in the present Universe, $\delta T/T \sim 10^{-5}$, require the inflaton potential to be extremely flat, i.e. $\alpha_\phi \ll 1$. 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, $T_{\rm RH}$ is expected to be smaller than 10^{16} 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 X 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, $\phi \to X\overline{X}$.

There is another problem associated with GUT baryogenesis, namely the problem of relic gravitinos [16]. The gravitino is the fermionic superpartner of graviton and has interaction strength with the observed standard model sector inversely proportional to $M_{\rm Pl}$. 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 be produced after inflation provides an upper bound on the reheating temperature,

$$T_{\rm RH} \lesssim (10^{10} - 10^{11})\ {\rm GeV} \tag{33}$$

which is too low to create superheavy X bosons that eventually decay and produce the baryon asymmetry.

These problems of GUT baryogenesis can be resolved in the new theory of reheating which differs significantly from the simple picture described above. In the first stage of reheating, called *preheating* [17], 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 parametric resonance. 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. Indeed, it has been shown [18] that the baryon asymmetry can be produced efficiently just after the preheating era, thus solving many of the drawbacks of GUT baryogenesis in the old theory of reheating.

6 Electroweak Baryogenesis

The departure from thermal equilibrium may also be attained by a different mechanism, namely the phase transition mechanism. At temperatures around the electroweak scale, the expansion rate of the Universe in thermal units is small compared to the rate of B-violating processes. This means that the equilibrium description of particle phenomena is extremely accurate at electroweak temperatures. Thus, baryogenesis cannot occur at such low energy scales without the aid of phase transitions.

It has been shown [19] that if the electroweak phase transition is second-order or a continuous crossover, the associated departure from equilibrium is insufficient for the required baryon asymmetry. This means that for electroweak baryogenesis to succeed, either the electroweak phase transition must be strongly first-order or other methods of destroying thermal equilibrium (e.g. topological defects) must be present at the phase transition [11, 13, 20].

For a first-order phase transition, there is an extremum at $\phi = 0$ (where ϕ is the order parameter), 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 $\phi \neq 0$ becomes the global minimum of the theory. A first-order electroweak phase transition proceeds via nucleation and growth of bubbles [21, 11]. The essential picture is that at around T_c, quantum tunneling occurs and nucleation of bubbles of the true vacuum in the sea of false vacua 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 nucleate that are just large enough to grow. These are called the *critical bubbles*, and they expand, eventually filling all of space and completing the phase transition. As the bubble walls pass each point in space, the order parameter changes rapidly, as do other fields, and this leads to a significant departure from thermal equilibrium. CP violating reflections and transmissions at the bubble surface then generate an asymmetry in baryon number, and for a sufficiently strong phase transition this asymmetry is frozen in the true vacuum inside the bubble.

In the standard model (SM) the electroweak transition is just a smooth crossover, because the experimental lower bound on the Higgs mass, $m_H > 114$ GeV, is larger than the critical Higgs mass, $m_H^c \simeq 72$ GeV, required for the phase transition to be first-order [22]. Hence, there is no departure from thermal equilibrium and baryogenesis cannot take place. The situation changes in many extensions of the SM such as two-Higgs doublet models and supersymmetric extensions of the SM where one can have a sufficiently strong first-order phase transition [11]. This requires, however, a rather exceptional mass spectrum of superparticles. Even more stringent constraints are obtained if the lightest neutralino is required to be the dominant component of cold dark matter. This suggests that, should supersymmetry be discovered at the LHC, the consistency of WIMP dark matter and electroweak baryogenesis will be a highly non-trivial test of supersymmetric extensions of the SM.

7 Affleck-Dine Baryogenesis

The Affleck-Dine (AD) mechanism [23] involves the cosmological evolution of scalar fields carrying baryonic charge. These scenarios are naturally implemented in the context of supersymmetric models. An important feature of supersymmetric field theories is the existence of *flat directions* in field space, on which the scalar potential vanishes. Let us 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 and slepton fields. Consider the case in which some component of the field χ lies along a flat direction. This means 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 SUSY breaking and nonrenormalizable operators lift the flat directions and set the scale for their potential.

During inflation, the χ field generically develops a large expectation value, establishing the initial conditions for its subsequent evolution. After inflation, when the Hubble rate becomes of the order of the curvature of the potential $V(\chi)$, the condensate leads to coherent oscillations around its present minimum. At this time, the B-violating terms in the potential are of comparable importance to the mass term, thereby imparting a large baryon number to the condensate. The decay of this condensate eventually converts the scalar charge densities to ordinary fermionic baryon number. After this time, the B-violating operators become negligible, so that the net baryon number of the Universe is preserved by the subsequent cosmological evolution.

This AD mechanism is a prominent example of non-thermal baryogenesis [13]. So far no 'standard model' of AD baryogenesis has emerged, and it also appears difficult to falsify this scenario. The most recent implementations of the AD scenario have been in the context of the MSSM[2] [24, 25]. In the MSSM models with large numbers of fields, flat directions occur because of accidental degeneracies in the field space. In general, flat directions carry a global $U(1)$ quantum number, and in particular, the MSSM contains many combinations of fields for constructing flat directions carrying a nonvanishing $B-L$ quantum number [25]. A successful AD baryogenesis mechanism is achieved if the effective potential during inflation contains a negative effective mass-term, so that a large expectation value for a flat direction may develop, and nonrenormalizable terms that lift the flat directions.

The fact that there are many flat directions in the MSSM provides an interesting possibility for the

[2]MSSM stands for the 'Minimal Supersymmetric Standard Model'.

formation of Q-balls[3] [26]. These macroscopic objects with large baryon number and mass,

$$B_Q \sim 10^{26}, \qquad M_Q \sim 10^{24}\ \text{GeV}, \tag{34}$$

can lead to striking signatures at Super-Kamiokande and ICECUBE. Alternatively, the decay of Q-balls at temperatures less than that of the electroweak phase transition can nonthermally produce higgsinos which yield the observed cold dark matter density. A discovery of higgsino dark matter in direct search experiments could be a hint for Q-balls as a possible nonthermal production mechanism. In this way, as in the case of electroweak baryogenesis, the nature of dark matter would provide a clue also for the origin of ordinary matter.

8 Baryogenesis via Leptogenesis

It is well-known that the SM Lagrangian is invariant under global Abelian symmetries which may be identified as the baryonic and leptonic symmetries. Therefore, B and L are *accidental* symmetries. As a result, it is not possible to violate B and L at the tree-level and at any order of perturbation theory. As a consequence, the proton is stable in the SM and any perturbative process which violates B and/or L in the GUTs is necessarily suppressed by powers of $M_{\rm GUT}/M_W$. Nevertheless, in many cases the perturbative expansion does not describe all the dynamics of the theory. It was realized by 't Hooft [27] that nonperturbative effects (instantons) may give rise to processes which violate the combination $B+L$ due to *chiral anomaly*, but not the orthogonal combination $B-L$.

There exists a deep connection between the chiral anomaly, the topological structure of $SU(2)$ and the baryon number violation [28]. Due to the topological properties of $SU(2)$, their gauge transformations may be divided in two categories based on whether or not they change the *Chern-Simons number*. Chern-Simons number has a topological nature which can be seen when considering configurations which are pure gauge at some initial and final times, t_0 and t_1, respectively. It can be shown that

$$N_{\rm CS}(t_1) - N_{\rm CS}(t_0) = n, \tag{35}$$

which is an integer, called the *winding number*. As a result, in the limit of zero Weinberg mixing angle (i.e. pure $SU(2)_L$ gauge) we may go from the classical vacuum to an infinite number of other vacua which are classically degenerate and have different Chern-Simons number. If the system is able to perform a transition from the vacuum $\mathcal{G}^{(n)}_{\rm vac}$ to the closest one $\mathcal{G}^{(n\pm1)}_{\rm vac}$, the Chern-Simons number is changed by one unity and

$$\Delta B = \Delta L = N_F, \tag{36}$$

where N_F is the number of fermionic families. In the SM, each transition creates 9 left-handed ($SU(2)_L$ doublet) quarks (3 color states for each generation) and 3 left-handed leptons (one per generation). It violates B and L by 3 units each, $\Delta B = \Delta L = 3$.

To quantify the probability of transition between two different vacua, it is important to understand the properties of the field configurations which interpolate the two vacua. It was found [29] that there exist static configurations corresponding to unstable solutions of the equation of motion. These solutions are called *sphalerons* and possess Chern-Simons number equal to 1/2. They correspond to saddle points of the energy functional and are localized in space (though unstable, contrary to the case of solitons). The energy of the sphaleron configuration is basically the result of the competition between the energy of the gauge configuration and the energy of the Higgs field. The typical value is [29]

$$E_{\rm sph} \sim 10\ \text{TeV} \tag{37}$$

The probability of baryon number nonconserving processes at zero temperature is highly suppressed by a factor of $\sim 10^{-150}$ [27]. This factor may be interpreted as the probability of making a transition

[3] Q-balls are stable, charged solitons appearing in a scalar field theory with a spontaneously broken global $U(1)$ symmetry. Such solitons are part of the MSSM spectrum.

from one classical vacuum to the closest one by tunneling, by going through the barrier of ~ 10 TeV corresponding to the sphaleron. However, in more extreme situations – like the primordial Universe at very high temperatures – these processes may be fast enough to play a significant role in baryogenesis.

The fact that the combination $B-L$ is left unchanged by sphaleron transitions opens up the possibility of generating baryon asymmetry from a lepton asymmetry [30]. The basic idea is that sphaleron transition will reprocess any lepton asymmetry produced and will convert some fraction of it into baryon number. This is because $B+L$ must be vanishing and therefore the final baryon asymmetry that results is $B \simeq -L$. The primordial lepton asymmetry is generated by the out-of-equilibrium decay of heavy right-handed Majorana neutrinos. Once the lepton number is produced, thermal scatterings redistribute the charges in such a way that in the high temperature phase of the SM the asymmetries of B and $B-L$ are proportional [31]:

$$B = \left(\frac{8N_F + 4N_H}{22N_F + 13N_H} \right) (B-L), \tag{38}$$

where N_H is the number of Higgs doublets.

In the SM as well as its unified extensions based on the group $SU(5)$, $B-L$ is conserved, and hence, no asymmetry in $B-L$ can be generated. However, a nonvanishing $B-L$ asymmetry may be obtained by adding right-handed Majorana neutrinos to the SM. This 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.

However, one has to avoid a large lepton number violation at intermediate temperatures which may potentially dissipate away the baryon number in combination with the sphaleron transitions. The requirement of harmless lepton number violation imposes an interesting bound on the neutrino mass [32]:

$$m_\nu \lesssim 4 \text{ eV} \left(\frac{T_X}{10^{10} \text{ GeV}} \right)^{-1/2}, \tag{39}$$

where $T_X \equiv \text{Min}\{T_{B-L}, 10^{12} \text{ GeV}\}$ and 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 $SO(10)$ models [33]. The consistency with the experimental evidence for neutrino masses has dramatically increased the popularity of the leptogenesis mechanism.

9 Conclusion

Modern cosmology provides really strong arguments in favor of the nonconservation of baryon number. Understanding how the baryon asymmetry of the Universe originated is one of its fundamental goals. Particle physics has provided us with a number of possibilities, all of which involve fascinating physics, but they have varying degrees of testability. The Standard Model, which is so successful in explaining the experimental data obtained at the current generation accelerators up to energy scales of about 100 GeV, seems unable to explain the observed baryon asymmetry of the Universe. This is a strong indication that there exists some New Physics beyond the SM. We do not know whether this is just the low energy supersymmetric extension of the SM. If so, the next generation of accelerators, such as LHC, will tell us if this a tenable option. A discovery of the standard supergravity scenario at LHC could be consistent with electroweak baryogenesis. On the other hand, evidence for gravitino dark matter can be consistent with leptogenesis. Finally, the discovery of macroscopic dark matter like Q-balls would point toward nonperturbative dynamics of scalar fields in the early Universe and therefore favor Affleck-Dine baryogenesis.

It might be that this cosmological puzzle has been taken care of by some new physics at energy scales much higher than the weak scale, the GUT scale $\sim 10^{16}$ GeV, as suggested by gauge coupling unification. Even though this option is not testable at particle colliders, the most striking evidence of baryon number violation might come from the detection of proton decay. Hopefully in the next few years we will be able to confirm (or disprove) some of the theories of baryogenesis.

References

[1] G. Steigman, *Ann. Rev. Astron. Astrophys.* **14**, 339 (1976).

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

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

[4] R. H. Cyburt, *Phys. Rev.* **D70**, 023505 (2004).

[5] W. Hu and S. Dodelson, *Ann. Rev. Astron. Astrophys.* **40**, 171 (2002).

[6] E. R. Harrison, *Ann. Rev. Astron. Astrophys.* **11**, 155 (1973).

[7] Riazuddin, `hep-ph/0302020` (2003).

[8] R. A. Alpher, J. W. Follin, Jr., and R. C. Herman, *Phys. Rev.* **92**, 1347 (1953); R.A. Alpher and R.C. Herman, *Science* **128**, 904 (1958).

[9] A. Riotto and M. Trodden, *Ann. Rev. Astron. Astrophys.* **49**, 35 (1999); A. Riotto, `hep-ph/9807454` (1998).

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

[11] J. M. Cline, `hep-ph/0609145` (2006).

[12] E. W. Kolb and M. S. Turner, *Ann. Rev. Nucl. Part. Sci.* **33**, 645 (1983).

[13] A. D. Dolgov, *Phys. Rep.* **222**, 309 (1992).

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

[15] D. H. Lyth and A. Riotto, *Phys. Rep.* **314**, 1 (1999).

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

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

[18] E. W. Kolb, A. D. Linde and A. Riotto, *Phys. Rev. Lett.* **77**, 4290 (1996); E.W. Kolb, A. Riotto and I. Tkachev, *Phys. Lett.* **B423**, 348 (1998).

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

[20] M. Trodden, *Rev. Mod. Phys.* **71**, 1463 (1999); A. G. Cohen, D. B. Kaplan and A. E. Nelson, *Ann. Rev. Nucl. Part. Sci.* **43**, 27 (1993).

[21] W. Bernreuther, `hep-ph/0205279` (2002).

[22] W. Buchmüller and O. Philipsen, *Nucl. Phys.* **B443**, 47 (1995); K. Kajantie et al., *Phys. Rev. Lett.* **77**, 2887 (1996).

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

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

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

[26] M. Dine and A. Kusenko, *Rev. Mod. Phys.* **76**, 1 (2004).

[27] G. 't Hooft, *Phys. Rev. Lett.* **37**, 8 (1976).

[28] V. A. Rubakov and M. E. Shaposhnikov, `hep-ph/9603208` (1996).

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

[30] M. Fukugita and T. Yanagida, *Phys. Lett.* **B174**, 45 (1986).

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

[32] A. Pilaftsis, *Int. J. Mod. Phys.* **A14**, 1811 (1999).

[33] W. Buchmüller and M. Plümacher, *Phys. Lett.* **B389**, 73 (1996).