BARYOGENESIS

I. WHAT DO WE TRY TO EXPLAIN

A. How much antimatter?

The CPT theorem assures that for any particle species X there exists an antiparticle $\overline{X}$ with exactly the same mass, $m_X = m_{\overline{X}}$, and decay width, $\Gamma_X = \Gamma_{\overline{X}}$. The charges associated with particles and antiparticles are opposite, $Q_X = -Q_{\overline{X}}$. This striking symmetry would naturally lead us to conclude that the Universe contains particles and antiparticles in equal number densities, $n_X = n_{\overline{X}}$. The observed Universe is, however, drastically different. We do not observe any bodies of antimatter within the solar system. Antiprotons $\overline{p}$ are observed in the cosmic rays but are likely to be produced as secondaries in collisions $pp \to 3p + \overline{p}$ at a rate similar to the observed one, $n_{\overline{p}}/n_p \sim 3 \times 10^{-4}$. The observational limit on $\overline{n}_{^4\mathrm{He}}/n_{^4\mathrm{He}}$ is similarly of order 10^{-5}.

We cannot exclude that the dominance of matter over antimatter is only local and is only realized up to a certain length scale ℓ_B beyond which the picture is reveresed and islands of antimatter are found. The size of our matter domain must, however, be quite large, roughly speaking $\ell_B \gtrsim 10$ Mpc. Indeed, for smaller scales one would expect a significant amount of energetic γ-rays coming from the reactions of annihilation of $p\overline{p} \to \pi$'s followed by the subsequent decays $\pi^0 \to 2\gamma$, which would take place in the boundary area separating the matter and antimatter islands. Another signature for the presence of domains of antimatter would be the distortion of the spectrum of the CMBR. In such a case, the lower bound on ℓ_B might be weaker if voids separate matter and antimatter domains. The voids might be created because of an excessive pressure produced by the annihilations at early stages of the evolution of the Universe or because of low density matter and antimatter in the boundary regions, provided that the baryon asymmetry changes sign locally so that in the boundaries it is zero or very small.

All these considerations lead us to conclude that, if domains of matter and antimatter exist in the Universe, they are separated on scales certainly larger than the radius of our own galaxy (~ 3 Kpc) and most probably on scales larger than the Virgo cluster (~ 10 Mpc).

B. How much matter?

The baryon number density does not keep constant during the evolution of the Universe because 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

$$\eta \equiv \frac{n_B}{n_\gamma} \equiv \frac{n_b - n_{\overline{b}}}{n_\gamma}. \tag{1}$$

The parameter η is essential for determining the light element abundances produced at the nucleosynthesis epoch. The parameter η might have not changed since nucleosynthesis. At the relevant energy scales ($\sim$ 1 MeV) the baryon asymmetry is conserved if there are no processes which would have produced entropy to change the photon number density.

The baryon number density is

$$n_B = \frac{\rho_B}{m_B} = \frac{\Omega_B}{m_B}\rho_c = 1.1 \times 10^{-5}\ h^2\Omega_B\ \mathrm{cm}^{-3}, \tag{2}$$

while the photon number densiry is

$$n_\gamma = \frac{2\zeta(3)}{\pi^2}T^3 = 415\left(\frac{T_0}{2.735^o\ K}\right)^3\ \mathrm{cm}^{-3}. \tag{3}$$

Putting (2) and (3) together, we obtain

$$\eta = 2.65 \times 10^{-8}\Omega_B h^2\left(\frac{T_0}{2.735^o\ K}\right)^{-3}. \tag{4}$$

The range of η consistent with nucleosynthesis is

$$\eta = (6 \pm 1) \times 10^{-10}, \tag{5}$$

or, equivalently,

$$\Omega_B h^2 = 0.22 \pm 0.03. \tag{6}$$

It is often convenient to describe the baryon asymmetry normalized to the entropy density s rather than to n_γ:

$$B \equiv \frac{n_B}{s} = (8.4 \pm 1.4) \times 10^{-11}. \tag{7}$$

Thus, we would like to explain a small number density of baryons, $n_b/s \sim 10^{-10}$, and a vanishingly small number density of antibaryons, $n_{\overline{b}} = 0$.

C. Naive attempts

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 temperatues $T \lesssim 1$ GeV the equilibrium abundances of nucleons and antinucleons is

$$\frac{n_b}{n_\gamma} \simeq \frac{n_{\overline{b}}}{n_\gamma} \approx \left(\frac{m_p}{T}\right)^{3/2} e^{-m_p/T}. \tag{8}$$

When the Universe cools off, the number of nucleons and antinucleons decreases as long as the annihilation rate, $\Gamma_{\rm ann} = n_b\langle\sigma_A v\rangle$, is larger than the expansion rate of the Universe, $H = 1.66 g_*^{1/2}\frac{T^2}{m_{\rm Pl}}$. The thermally averaged cross section $\langle\sigma_A v\rangle$ is of the order of m_π^2. At $T \sim 20$ MeV, $\Gamma_{\rm ann} \simeq H$ and annihilations freeze out, nucleons and antinucleons being so rare that they cannot find each other and annihilate any longer. At that time,

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

which is much smaller than the value required by nucleosynthesis. In order to avoid the annihilation catastrophe, we may suppose that hypothetical new interactions separated matter from antimatter before $T \approx 38$ MeV, when $\eta \approx 10^{-10}$. At that time, $t \approx 10^{-3}$ sec. The causal region (horizon) was, however, small and contained only $\sim 10^{-7} M_\odot$. Hence, we cannot explain the asymmetry over galaxy scales.

This argument is not valid in cosmological models invoking inflation. There, the region of the universe that is causally connected today was connected even at times $\sim 10^{-3}$ sec. These scenarios pose other serious cosmological problems. If the processes responsible for the separation of matter from antimatter took place before inflation, then the baryon number density was diluted by an enormous factor $\sim \exp(200)$, because of the entropy production due to inflation. On the other side, if the separation took place after inflation, then it is not clear how to eliminate the boundaries separating matter from antimatter islands.

Another possibility is to try to explain the tiny value of η via statistical fluctuations in the baryon and antibaryon distributions. Our own galaxy contains at present approximately 10^{79} photons and 10^{69} baryons. But when the temperature was $T \gtrsim 1$ GeV, the comoving volume V that presently encompasses our galaxy contained about 10^{79} baryons and antibaryons. From pure statistical fluctuations, one might expect an asymmetry $(n_b - n_{\overline{b}})/n_b \simeq (n_b V)^{-1/2} \sim 10^{-40}$, which is, again, far too small to explain the observed η.

Finally, we could impose an initial asymmetry by hand. That would not explain anything and, furthermore, would be extremely unnatural. Moreover, such a scenario is inconsistent with inflation.

II. THE SAKHAROV CONDITIONS

Sakharov pointed out that a baryon asymmetry can be dynamically generated during the evolution of the Universe from an initial state with baryon number equal to zero is the following three conditions are met:

- Baryon number (B) violation;
- C and CP violation;
- Departure from thermal equilibrium.

1. B violation: Suppose that the baryon number B is conserved by the interactions. This means that the baryon number commutes with the Hamiltonian of the system H: $[B, H] = 0$. Therefore, supposing that $B(t_0) = 0$, we have $B(t) \propto \int_{t_0}^{t} [B, H] dt' = 0$ at all times and no baryon number production may take place.

2. C and CP violation: The baryon number operator,

$$\hat{B} = \frac{1}{3} \sum_q \int d^3x \ : q^\dagger(x,t) q(x,t) :, \tag{10}$$

is C-odd and CP-odd. This can be seen from the action of P, C and T on the quark fields:

$$\begin{aligned} Pq(x,t)P^{-1} &= \gamma^0 q(-x,t), \quad Pq^\dagger(x,t)P^{-1} = q^\dagger(-x,t)\gamma^0, \\ Cq(x,t)C^{-1} &= i\gamma^2 q^\dagger(x,t), \quad Cq^\dagger(x,t)C^{-1} = iq(x,t)\gamma^2, \\ Tq(x,t)T^{-1} &= -iq(x,-t)\gamma_5\gamma^0\gamma^2, \quad Tq^\dagger(x,t)T^{-1} = -i\gamma^2\gamma^0\gamma_5 q^\dagger(x,-t). \end{aligned} \tag{11}$$

Then

$$\begin{aligned} P : q^\dagger(x,t)q(x,t) : P^{-1} &= : q^\dagger(-x,t)q(-x,t) :, \\ C : q^\dagger(x,t)q(x,t) : C^{-1} &= - : q^\dagger(x,t)q(x,t) :, \\ T : q^\dagger(x,t)q(x,t) : T^{-1} &= : q^\dagger(x,-t)q(x,-t) :, \end{aligned} \tag{12}$$

and

$$P\hat{B}P^{-1} = \hat{B}, \quad C\hat{B}C^{-1} = -\hat{B}, \quad (CP)\hat{B}(CP)^{-1} = -\hat{B}. \tag{13}$$

Therefore, a non-zero expectation value $\langle\hat{B}\rangle$ requires that the Hamiltonian violates C and CP. More intuitively, C symmetry would guarantee that $\Gamma(i \to f) = \Gamma(i^\dagger \to f^\dagger)$, while CP symmetry would guarantee that $\Gamma(i \to f) = \Gamma(\bar{i} \to \overline{f})$ ($x^\dagger$ has opposite charges and same chiralities as x, $\overline{x}$ has opposite charges and opposite chiralities). Even though it may be possible (with CP) to create a baryon asymmetry in a certain region of phase space, integrating over all momenta and summing over all spins would produce a vanishing asymmetry.

3. Departure from thermal equilibrium: To give a formal proof, let us define $\Theta \equiv CPT$. Then, from (12) we obtain

$$\Theta\hat{B}(0)\Theta^{-1} = -\hat{B}(0). \tag{14}$$

A system in thermal equilibrium is described by a density operator $\rho = \exp(-H/T)$. Using $\hat{B}(t) = e^{iHt}\hat{B}(0)e^{-iHt}$, we have

$$\langle\hat{B}(t)\rangle_T = \mathrm{tr}[e^{-H/T}e^{iHt}\hat{B}(0)e^{-iHt}] = \mathrm{tr}[e^{-iHt}e^{-H/T}e^{iHt}\hat{B}(0)] = \langle\hat{B}(0)\rangle_T, \tag{15}$$

and

$$\langle\hat{B}(0)\rangle_T = \mathrm{tr}[e^{-H/T}\hat{B}(0)] = \mathrm{tr}[\Theta^{-1}\Theta e^{-H/T}\hat{B}(0)] = \mathrm{tr}[e^{-H/T}\Theta\hat{B}(0)\Theta^{-1}] = -\langle\hat{B}(0)\rangle_T. \tag{16}$$

How the average baryon number is kept equal to zero in thermal equilibrium is a bit tricky, as the following example shows. Consider an ensemble of a heavy particles species X that has a $\Delta B = +2/3$ decay mode, $X \to qq$, and a $\Delta B = -1/3$ decay mode, $X \to \ell\overline{q}$. Further assume that there is C and CP violation in these decays, such that an asymmetry in the partial decay rates of X and its antiparticle $\overline{X}$ is induced:

$$\Gamma(X \to qq) = (1+\epsilon)\Gamma_0, \quad \Gamma(\overline{X} \to \overline{qq}) = (1-\epsilon)\Gamma_0. \tag{17}$$

An asymmetry appears also in the other channel. CPT requires that $\Gamma_X = \Gamma_{\overline{X}}$. In the decays (17) a non-zero baryon number is generated. Yet, the ensemble is supposed to be in thermal equilibrium. One might be inclined to appeal to the principle of detailed balance, which would tell us that the inverse decay $qq \to X$ is more likely that $\overline{qq} \to \overline{X}$, and the temporary excess $\Delta B \neq 0$ is erased in this way. This principle is based, however, on T invariance. CPT implies that this symmetry is broken because of CP violation. In fact, applying a CPT transformation to the above decays, we learn that the inverse decay pushes ΔB into the same direction as (17):

$$\Gamma(qq \to X) = (1-\epsilon)\Gamma_0, \qquad \Gamma(\overline{qq} \to \overline{X}) = (1+\epsilon)\Gamma_0. \tag{18}$$

The elimination of the baryon number ΔB is achieved by the B-violating reactions $qq \to \ell\overline{q}$, $\overline{qq} \to \overline{\ell}q$ and the CPT-conjugate reactions (the $X, \overline{X}$ resonance contributions are to be taken out of the scattering amplitudes). It is the unitarity of the S matrix that does the job of keeping $\langle \hat{B} \rangle_T = 0$ in thermal equilibrium.

III. CP AND B VIOLATION IN THE STANDARD MODEL

CP (and C) are violated, within the Standard Model (SM), by the quark charged current interactions:

$$\mathcal{L}_{\rm cc} = -\frac{g}{\sqrt{2}}\overline{U_L}\gamma^\mu V_{\rm CKM} D_L W_\mu^+ + {\rm h.c.}. \tag{19}$$

Here $U_L = (u_L, c_L, t_L)^T$ and $D_L = (d_L, s_L, b_L)$. The quark mixing matrix $V_{\rm CKM}$ is a 3×3 unitary matrix. Taking into account the freedom in redefining quark field phases, one finds that $V_{\rm CKM}$ is still, in general, complex, depending on a single phase (the Kobayashi-Maskawa phase) and, consequently, $\mathcal{L}_{\rm cc}$ violates CP.

There is also baryon number violation in the Standard Model, but this is a subtle, nonperturbative effect which is completely negligible for particle interactions in the laboratories at present-day collision energies. It is, however, very significant for the physics of the early Universe. Let us outline how this arises. Baryon number, under which all quarks (antiquarks) carry charge $+1/3$ $(-1/3)$ and all other particles are neutral, has never been experimentally observed to be violated. Within the SM, this is explained by the Lagrangian, $\mathcal{L}_{\rm SM}$, when subject to the SM gauge group, $G_{\rm SM} = SU(3)_{\rm C} \times SU(2)_{\rm W} \times U(1)_{\rm Y}$, and containing the three generations of $(3,2)_{1/6} + (\bar{3},1)_{-2/3} + (\bar{3},1)_{+1/3} + (1,2)_{-1/2} + (1,1)_{+1}$ $G_{\rm SM}$-representations, acquires an accidental $U(1)_B$ (and $U(1)_e \times U(1)_\mu \times U(1)_\tau$ and, in particular, $U(1)_L$) accidental symmetry. The $U(1)_B$ and $U(1)_L$ transformation laws are

$$U(1)_B: \quad q \to e^{i\omega/3}q, \quad \ell \to \ell; \qquad U(1)_L: \quad \ell \to e^{i\theta}\ell, \quad q \to q. \tag{20}$$

Applying Noether's theorem, we obtain the associated symmetry currents, J_μ^B and J_μ^L, which are conserved at the Born level:

$$\partial^\mu J_\mu^B = \partial^\mu \sum_q \frac{1}{3}\overline{q}\gamma_\mu q = 0; \qquad \partial^\mu J_\mu^L = \partial^\mu \sum_\ell \overline{\ell}\gamma_\mu \ell = 0. \tag{21}$$

The associated charge operators,

$$\hat{B} = \int d^3x \; J_0^B(x); \qquad \hat{L} = \int d^3x \; J_0^L(x), \tag{22}$$

are time-independent. At the level of quantum fluctuations, however, these symmetries are broken because eqs. (21) no longer hold. This is seen as follows. Decompose the vector current into left- and right-handed pieces:

$$\overline{f}\gamma_\mu f = \overline{f_L}\gamma_\mu f_L + \overline{f_R}\gamma_\mu f_R, \tag{23}$$

where $f = q, \ell$. Because of the clash between gauge and chiral symmetries at the quantum level, the gauge-invariant chiral currents are not conserved: the current-divergences suffer from *anomaly.* For a gauge theory based on gauge group G, which is a simple Lie group of dimension d_G, the anomaly equations read

$$\partial^\mu \overline{f_L}\gamma_\mu f_L = -c_L \frac{g^2}{32\pi^2} F^a_{\mu\nu}\tilde{F}^{a\mu\nu}, \quad \partial^\mu \overline{f_R}\gamma_\mu f_R = +c_R \frac{g^2}{32\pi^2} F^a_{\mu\nu}\tilde{F}^{a\mu\nu}, \tag{24}$$

where $F^{a\mu\nu}$ is the (non)Abelian field strength tensor ($a = 1, 2 \ldots, d_G$) and $\tilde{F}^{a\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\alpha\beta}F^a_{\alpha\beta}$ is the dual tensor, and $c_{L,R}$ depend on the G-representation of $f_{L,R}$. Let us apply eqs. (23) and (24) to the baryon and lepton number currents of the SM. Gluons couple to left- and right-handed quark currents in the same way, so that $c_L^{\mathrm{C}} = c_R^{\mathrm{C}}$. Consequently, J_μ^B and J_μ^L have no QCD anomaly. But, since W_μ^a ($a = 1, 2, 3$) couple only to left-handed quarks and leptons, we have $c_R^W = 0$, and, since B couples differently to f_L and f_R, we have $c_L^{\mathrm{Y}} \neq c_R^{\mathrm{Y}}$. One obtains:

$$\partial^\mu J_\mu^B = \partial^\mu J_\mu^L = \frac{n_F}{32\pi^2}(-g^2 W^a_{\mu\nu}\tilde{W}^{a\mu\nu} + g'^2 B_{\mu\nu}\tilde{B}^{\mu\nu}), \tag{25}$$

where $n_F = 3$ is the number of generations.

Eq. (25) implies that

$$\partial^\mu (J_\mu^B - J_\mu^L) = 0. \tag{26}$$

Thus, *the quantum number $B - L$ is conserved in the SM.*

How is $B+L$ violated? The RHS of eq. (25) can be written as the divergence of a current K^μ, that is $\partial^\mu J^B_\mu = \partial^\mu J^L_\mu = n_F \partial^\mu K_\mu$, where

$$K^\mu = -\frac{g^2}{32\pi^2}\epsilon^{\mu\nu\alpha\beta} W^a_\nu \left(2\partial_\alpha W^a_\beta + \frac{2}{3} g \epsilon^{abc} W^b_\alpha W^c_\beta \right) + \frac{g'^2}{32\pi^2}\epsilon^{\mu\nu\alpha\beta} B_\nu B_{\alpha\beta}. \tag{27}$$

Let us integrate the RHS of (25) over d^4x, using the definition of K_μ. Using Gauss law we convert these integrals into integrals over a surface at infinity. The Abelian part does not contribute, and we obtain

$$\int d^4x \partial^\mu K_\mu = \frac{g^3}{96\pi^2} \int_{\partial V_4} dn^\mu \epsilon_{\mu\nu\alpha\beta} \epsilon^{abc} W^{a\nu} W^{b\alpha} W^{c\beta}. \tag{28}$$

We choose the surface ∂V_4 to be a large cylinder with top and bottom surfaces at t_f and t_i, and let the volume of the cylinder tend to infinity. We choose to work in the temporal gauge, $W^a_0 = 0$. Then there is no contribution from the integral over the coat of the cylinder and we obtain

$$\int d^3x dt \; \partial^\mu K_\mu = N_{\rm CS}(t_f) - N_{\rm CS}(t_i) \equiv \Delta N_{\rm CS}, \tag{29}$$

where

$$N_{\rm CS}(t) = \frac{g^3}{96\pi^2} \int d^3x \; \epsilon_{ijk} \epsilon^{abc} W^{ai} W^{bj} W^{ck} \tag{30}$$

is the Chern-Simons number. This integral assigns a topological 'charge' to a classical gauge field. Note that $N_{\rm CS}$ is not gauge invariant, but $\Delta N_{\rm CS}$ is. A nonAbelian gauge theory, like the weak-interaction $SU(2)_{\rm W}$, is topologically non-trivial, which is reflected by the fact that it has an infinite number of ground states whose vacuum-field configurations have different topological charges $\Delta N_{\rm CS} = 0, \pm 1, \pm 2, \ldots$. Imagine the set of gauge and Higgs fields and consider the energy functional $E[W^a_\mu, \phi]$ that forms a hypersurface over the infinite-dimensional space. The ground states with different topological charge are separated by a potential barrier.

Integrating also the LHS of (25), we get $(\Delta \hat{Q} \equiv \hat{Q}(t_f) - \hat{Q}(t_i))$

$$\Delta \hat{B} = \Delta \hat{L} = n_F \Delta N_{\rm CS}. \tag{31}$$

The interpretation of eq. (31) is as follows. As long as we consider small gauge field quantum fluctuations around the perturbative vacuum configuration $W^a_\mu = 0$, the RHS of (31) is zero, and B and L numbers remain conserved. This is the case in perturbation theory to arbitrary order, where B- and L-violating processes have zero amplitudes. However, large gauge fields,

$W^a_\mu \sim 1/g$, with non-zero topological charge $\Delta N_{\rm CS} = \pm 1, \pm 2, \ldots$ exist. They can induce transitions between fermionic states $|i, t_i\rangle$ and $|f, t_f\rangle$ with baryon and lepton numbers that differ according to the rule (31). This selection rule tells us that B and L must change by at least three units. (This means that, even when these non-perturbative effects are taken into account, the SM still predicts a stable proton.) Actually, the selection rule can be refined:

$$\Delta L_e = \Delta L_\mu = \Delta L_\tau = \Delta B/3 = \Delta N_{\rm CS}. \tag{32}$$

The dominant B- and L-violating transitions are between states $|i, t_i\rangle$ and $|f, t_f\rangle$ where $|\Delta B| = |\Delta L| = 3$. At $T = 0$, such transitions are induced by the (anti)instanton, a field configuration whose topological charge differ by ± 1. The corresponding amplitudes $\langle f, t_f|i, t_i\rangle$ involve 9 left-handed quarks (right-handed antiquarks), with each generation participating with three different color states, and 3 left-handed leptons (right-handed antileptons), one of each generation. Hence we have, for example, the anti-instanton induced reaction with $\Delta B = \Delta L = -3$:

$$u + d \to \bar{d} + 2\bar{s} + \bar{c} + 2\bar{b} + \bar{t} + \bar{\nu}_e + \bar{\nu}_\mu + \bar{\nu}_\tau. \tag{33}$$

What is the probability for such a transition to occur? The transition corresponds to a tunneling process and is thus exponentially suppressed. For $E_{\rm c.m.} \lesssim 1$ TeV, the cross section is

$$\sigma_{|\Delta B|=|\Delta L|=3} \propto e^{-4\pi/\alpha_2} \sim 10^{-150}. \tag{34}$$

When the SM is coupled t a heat bath of temperature T, the situation changes. At high temperatures, $T \gtrsim T_{\rm EW} \sim 100$ GeV, the B- and L-violating processes in the SM are fast enough to play a significant role in baryogenesis. The ground states with different $N_{\rm CS}$ are separated by a potential barrier of minimal height,

$$E_{\rm sph}(T) = \frac{4\pi}{g} v_T f(\lambda/g), \tag{35}$$

where $v_T \equiv \langle\phi\rangle_T$ is the VEV of the SM Higgs field at temperature T (e.g. $v_{T=0} = 246$ GeV). The parameter f varies between $1.6 < f < 2.7$ depending on the ratio between the Higgs quartic coupling λ and the gauge coupling g. This yields $E_{\rm sph}(T = 0) \simeq 8 - 13$ TeV. The subscript "sph" refers to the sphaleron, a gauge and Higgs field configuration of Chern-Simons number 1/2(+integer) which is an (unstable) solution of the classical equations of the SM gauge-Higgs sector. These kinds of field configurations lie on the respective minimum energy path from one ground state to another with a different Chern-Simons number. The rate of fermion number non-conserving transitions will be proportional to the Boltzmann factor $\exp[-E_{\rm sph}(T)/T]$ as long as the energy of the thermal excitations is smaller than that of the barrier, while unsuppressed transitions will occur above the barrier.

At this point we recall that the EW gauge symmetry is unbroken at high temperatures, that is, in the early Universe. The critical temperature $T_{\rm EW}$ where the transition from a symmetric phase with $v_T = 0$ to a broken phase with $v_T \neq 0$ occurs is, in the SM, about 100 GeV. Hence, B- and L-violating transition rates are not exponentially suppressed above this temperature. Detailed investigations have led to the following results:

• In the broken phase, i.e. $T < T_{\rm EW} \sim 100$ GeV, the sphaleron induced $(B+L)$-breaking transition rate per volume V is given by

$$\frac{\Gamma^{\rm sph}_{\not B+\not L}}{V} = \kappa_1 \left(\frac{m_W}{\alpha_W T}\right)^3 m_W^4 \exp[-E_{\rm sph}(T)/T], \tag{36}$$

where $m_W(T) = g v_T/2$ is the temperature-dependent mass of the W-boson and κ_1 is a dimensionless constant.

• The calculation of the transition rate in the symmetric phase is very difficult. On dimensional grounds, we expect this rate per volume to be $\propto T^4$. Recent investigations yield, for $T > T_{\rm EW} \sim 100$ GeV,

$$\frac{\Gamma^{\rm sph}_{\not B+\not L}}{V} = \kappa_2 \alpha_W^5 T^4, \tag{37}$$

with $\kappa_2 \sim 20$.

By comparing $\Gamma^{\rm sph}_{\not B+\not L}$ above $T_{\rm EW}$ with the expansion rate H, we can assess whether the $(B+L)$-violating [$(B-L)$-conserving] SM interactions are fast enough to keep up with the expansion of the early, radiation dominated Universe. We find that $\Gamma^{\rm sph}_{\not B+\not L} \gg H$, so that these interactions are in thermal equilibrium, for temperatures T in the range

$$T_{\rm EW} \lesssim T \lesssim 10^{12}\ {\rm GeV}. \tag{38}$$

This result provides an important constraint on any baryogenesis mechanism that operates above $T_{\rm EW}$. If the B- and L-violating interactions involved in this mechanism conserve $B-L$, then any excess of baryon and lepton number generated above $T_{\rm EW}$ will be washed out by the B- and L-violating SM sphaleron-induced interactions. Hence, baryogenesis scenarios above $T_{\rm EW}$ must be based on particle physics models that violate also $B-L$.

IV. THE ELECTROWEAK PHASE TRANSITION (EWPT)

A baryogenesis scenario based on the SM requires that the departure from thermal equilibrium must come from the electroweak phase transition (EWPT). The reason for that is that the expansion rate of the Universe at temperatures $\lesssim 10^{12}$ GeV is too slow for causing a departure from local thermal equilibrium: the reaction rates of most of the SM particles are typically $\Gamma \gtrsim \alpha_W^2 T \gg H$. If $\Delta B \neq 0$ was created at the EWPT, and this phase transition was strongly first order, B would be frozen in during the later evolution of the Universe, as the B- and L-violating reactions would be strongly suppressed below $T_{\rm EW}$. Investigations have shown, however, that the EWPT in the SM fails to provide the thermal instability. We now explain why the SM fails.

At $T = 0$, the tree-level effective potential is just the classical Higgs potential,

$$V_{\rm tree} = -\mu^2 \Phi^\dagger \Phi + \lambda(\Phi^\dagger \Phi)^4, \tag{39}$$

leading to $v_{T=0} = \sqrt{2\mu^2/\lambda}$, and

$$m_H = \sqrt{2\lambda} v_{T=0} + \text{quantum corrections}. \tag{40}$$

LEP established $m_H > 114$ GeV and, hence, $\lambda > 0.33$.

At $T \neq 0$, the SM effective potential gets finite temperature corrections. Neglecting g' and small Yukawa couplings, one obtains

$$V_{\rm eff}(\phi, T) = \frac{1}{2} a (T^2 - T_1^2)\phi^2 - \frac{1}{3} b T \phi^3 + \frac{1}{4}\lambda \phi^4, \tag{41}$$

where

$$a = \frac{3g^2}{16} + \left(\frac{1}{2} + \frac{m_t^2}{m_H^2}\right)\lambda, \quad b = \frac{9g^3}{32\pi}, \quad T_1 = \frac{m_H}{2\sqrt{a}}. \tag{42}$$

Eq. (41) confirms that at high temperature $v_T = 0$. At $T_c = T_1/\sqrt{1 - 2b^2/(9a\lambda)} > T_1$ a first order phase transition occurs: $V_{\rm eff}$ has two degenerate minima, separated by an energy barrier, one at $\phi = 0$ and the other at

$$v_{T_c} \equiv \phi_{\rm crit} = \frac{2b}{3\lambda} T_c. \tag{43}$$

At T_c the free energy of the symmetric and the broken phase are equal. The Universe remains, however, for a while in the symmetric phase because of the energy barrier. As the Universe expands and cools further down, bubbles filled with the Higgs condensate start to nucleate at some temperature below T_c. These bubbles become larger by releasing latent heat, percolate, and eventually fill the whole volume at $T = T_1$.

In the case of a second order phase transition, there are no energetically degenerate minima separated by a barrier, i.e. no bubble nucleation and expansion. The Higgs field gradually condenses uniformly at $T \lesssim T_c$ and grows to its present value as the system cools off.

The value of the critical temperature depends on the parameters of the respective model and is obtained by detailed computations. Nevertheless, we may use the above formula for a crude estimate, $T_c \sim 70$ GeV for $m_H = 100$ GeV. (This implies that the EWPT took place at $t \sim 5 \times 10^{-11}$ s, with a causal domain, $d_H(t) = 2t$, of order of a few centimeters.)

It should be clear by now why a first order EWPT is required for baryogenesis. In this case, the time scale associated with the nucleation and expansion of the Higgs bubbles is comparable with the time scales of the particle reactions. This causes a departure from thermal equilibrium. Let us quantify this situation. Consider one of the bubbles with $v_T \neq 0$ which, after expansion and percolation, becomes our Universe. The bubble must get filled with more quarks than antiquarks such that $n_B/s \sim 10^{-10}$ and this ratio remains conserved. This means that baryogenesis has to take place outside the bubble, while the sphaleron-induced $(B + L)$-violating reactions within the bubble must be strongly suppressed. In order that the sphaleron rate, which in the broken phase is given by eq. (36), $\Gamma^{\rm sph}_{\not B + \not L} \propto \exp(-4\pi f v_T/gT)$, is practically switched off, the order parameter must jump at T_c from $\phi = 0$ in the symmetric phase to a value v_{T_c} in the broken phase such that

$$\frac{v_{T_c}}{T_c} \gtrsim 1. \tag{44}$$

This is the condition for a first order transition to be strong.

In view of the experimental bound, $m_H^{\rm SM} \geq 114$ GeV, eqs. (41) and (42), which are valid only for a very light Higgs, no longer apply. Nevertheless, eq. (43) shows that the discontinuity gets weaker for a larger Higgs mass. Various studies, analytical, numerical and lattice calculations, confirm this generic picture: a strongly first order phase transition occurs for $m_H \lesssim 40$ GeV; a weakly first order phase transition occurs for 40 GeV $\lesssim m_H \lesssim m_W$; for a higher Higgs mass, the phase transition is second order.

Thus, the LEP result, $m_H^{\rm SM} \geq 114$ GeV, leads to the following conclusion: if the SM Higgs mechanism provides the correct description of the EWSB, then the EWPT in the early Universe does not provide the thermal instability required for baryogenesis. The B-violating sphaleron processes are only adiabatically switched off during the transition from $T > T_c$ to $T < T_c$; they are still thermal for $T \lesssim T_c$. Thus, the SM cannot explain η.

Simple extensions of the SM Higgs sector do have room for a first order EWPT. Thus, for example, the minimal supersymmetric SM (MSSM) predicts a sufficiently strong phase transition if

$$m_h \lesssim 105 - 115 \text{ GeV}, \quad m_{\tilde{t}_R} \lesssim 170 \text{ GeV}. \tag{45}$$

In the next-to-minimal SUSY model (NMSSM), a strong first order transition can be easily arranged. Similarly, various two Higgs doublet models (2HDM) can account for a viable thermal instability.

V. ELECTROWEAK BARYOGENESIS

Electroweak baryogenesis works only in extensions of the SM. The required departure from thermal equilibrium is provided by the expansion of the Higgs bubbles, the true vacuum. When the bubble walls pass through a point in space, the classical Higgs fields change rapidly in the vicinity of such a point, as do other fields that couple to the Higgs.

Let us see how baryogenesis occurs. We consider for simplicity only the so-called thin-wall regime, which applies if the mean free path of a fermion, l_ψ, is larger than the wall thickness, l_w. Then the quarks and leptons can be treated as free particles, interacting only in a small region with a non-trivial Higgs background field. The expansion of the wall is supposed to be spherically symmetric and a one-dimensional description applies. Consider left-handed and right-handed quarks, q_L and q_R, incident from the symmetric phase, which hit the moving wall and are reflected by the Higgs bubble into right-handed and left-handed quarks, respectively.

In the frame where the wall is at rest, the fermion interactions with the bubble wall are described by the Dirac equation,

$$(i\gamma^\mu \partial_\mu - m(z)P_R - m^*(z)P_L)\psi(z,t) = 0, \tag{46}$$

where $P_{R,L} = (1 \pm \gamma_5)/2$ and z is the distance from the wall. Instead of solving the equation, let us make a few general considerations. Let us consider the scattering process where, in the symmetric phase ($z > 0$), a left-handed quark q_L (having momentum $k_z < 0$) is reflected at the wall into a right-handed q_R. Conservation of electric charge guarantees that a quark is reflected into a quark and not into an antiquark. Angular momentim conservation tells us that q_L is reflected as q_R and vice versa. If CP were conserved, then

$$\mathcal{R}_{L\to R} = \mathcal{R}_{\overline{R}\to\overline{L}}, \quad \mathcal{R}_{R\to L} = \mathcal{R}_{\overline{L}\to\overline{R}}, \tag{47}$$

would hold. CPT, respected by all particle physics models that we consider, implies

$$\mathcal{R}_{L\to R} = \mathcal{R}_{\overline{L}\to\overline{R}}, \quad \mathcal{R}_{R\to L} = \mathcal{R}_{\overline{R}\to\overline{L}}, \tag{48}$$

The mechanism works as follows. At some initial time we have equal numbers of quarks and antiquarks in the symmetric phase, in particular, $n_{q_L} = n_{\overline{q}_R}$ and $n_{q_R} = n_{\overline{q}_L}$, which hit the expanding bubble wall. Reflection converts $q_L \to q_R$, $\overline{q}_R \to \overline{q}_L$, $q_R \to q_L$, and $\overline{q}_L \to \overline{q}_R$, and the particles move back to the region where the Higgs fields are zero. Because the interaction with the bubble wall is assumed to be CP violating, the relations (47) for the reflection probabilities no longer hold,

$$\Delta\mathcal{R}_{CP} \equiv \mathcal{R}_{L\to R} - \mathcal{R}_{\overline{R}\to\overline{L}} = \mathcal{R}_{\overline{L}\to\overline{R}} - \mathcal{R}_{R\to L} \neq 0. \tag{49}$$

An explicit computation yields

$$\Delta\mathcal{R}_{CP}(k_z) \propto \int_{-\infty}^{\infty} dz \cos(2k_z z)\mathcal{I}m[m_q(z)M_q^*], \tag{50}$$

where $M_q = m_q(z = -\infty)$ is the quark mass in the broken phase. Note that at this stage the baryon number is still zero. This is because the difference J_q^L of the fluxes if $\overline{q}_R$ and q_L, injected from the wall back into the symmetric phase, is equal to J_q^R which we define as the difference of the fluxes of q_R and $\overline{q}_L$. However, the $(B + L)$-violating weak sphaleron interactions, which are unsuppressed in the symmetric phase away from the wall, act only on the (massless) left-handed quarks and right-handed antiquarks. For example, the reaction (33) decreases baryon number by 3 units, while the corresponding reaction with right-handed antiquarks in the initial state increases B by the same amount. Thus, if the functional form of the CP violating part $\mathcal{I}m[m_q(z)M_q^*]$ is such that $J_q^L > 0$ then, after the anomalous weak interactions took place, there are more left-handed quarks than right-handed antiquarks. The fluxes of the reflected $\overline{q}_L$ and q_R are not affected by the anomalous weak sphaleron interactions. Adding it all up, we see that some place away from the wall a net baryon number, $\Delta B > 0$, is produced. Some instant later, the expanding bubble wall sweeps over that region and the associated non-zero Higgs fields strongly suppress the $(B+L)$-violating backreactions that would wash out ΔB. Thus the non-zero B number produced before is frozen in.

We must also take into account that (anti)particles in the broken phase can be tranmitted into the symmetric phase and contribute to the (anti)particle fluxes discussed above. Using CPT invariance and unitarity, the transmission and reflection probabilities are related:

$$\begin{aligned}\mathcal{T}_{L\to L} &= 1-\mathcal{R}_{R\to L} = 1-\mathcal{R}_{\overline{R}\to\overline{L}} = \mathcal{T}_{\overline{L}\to\overline{L}},\\ \mathcal{T}_{R\to R} &= 1-\mathcal{R}_{L\to R} = 1-\mathcal{R}_{\overline{L}\to\overline{R}} = \mathcal{T}_{\overline{R}\to\overline{R}}.\end{aligned} \tag{51}$$

We can now write down a formula for the current J_q^L, that is the difference of $\overline{q}_R$ and q_L fluxes injected from the wall into the symmetric phase. The contribution from the reflected particles involves $\Delta\mathcal{R}_{CP} f_s$, where f_s is the Fermi-Dirac phase-space distribution of the (anti)quarks in the region $z > 0$ that move to the left, i.e. towards the wall. The contribution from the (anti)quarks that return from the broken phase involves $(\mathcal{T}_{\overline{R}\to\overline{R}} - \mathcal{T}_{L\to L}) f_b = -\Delta\mathcal{R}_{CP} f_b$, where f_b is the phase-space distribution of the transmitted (anti)quarks that move to the right. The reference frame is the wall frame. Note that f_s and f_b differ because the wall moves (in our convention - from left to right) with a velocity $v_w \neq 0$. Then

$$J_q^L = \int_{k_z<0} \frac{d^3k}{(2\pi)^3}\frac{|k_z|}{E}(f_s - f_b)\Delta\mathcal{R}_{CP}, \tag{52}$$

where $|k_z|/E$ is the group velocity. The current is non-zero because two of the three Sakharov conditions, CP violation and departure from thermal equilibrium, are met. The current would vanish if the wall were at rest in the plasma frame, which leads to thermal equilibrium, because then $(f_s - f_b) = 0$.

A. The Kobayashi Maskawa Mechanism

The SM fails to account for η in two aspects. First, as discussed above, the SM Higgs sector cannot generate a strongly first order phase transition. Second, as we explain in this section, CP violation from the Kobayashi Maskawa phase contributes to η with a much too strong suppression.

CP is violated within the Standard Model if, and only if, the following quantity is different from zero:

$$J_{CP} = (m_t^2 - m_c^2)(m_t^2 - m_u^2)(m_c^2 - m_u^2)(m_b^2 - m_s^2)(m_b^2 - m_d^2)(m_s^2 - m_d^2)\mathcal{I}mQ, \tag{53}$$

where $\mathcal{I}mQ$ can be defined as $\mathcal{I}m(V_{ud}V_{cb}V_{ub}^*V_{cd}^*)$. In the standard parametrization,

$$\mathcal{I}mQ = s_{12}s_{23}s_{13}c_{12}c_{13}^2c_{23}s_\delta \simeq 2 \times 10^{-5}. \tag{54}$$

Baryogenesis driven by the KM mechanism of CP violation generates a baryon asymmetry that is proportional to J_{CP}. Actually, it is suppressed by the following dimensionless measure of CP violation:

$$\eta^{\mathrm{KM}} \propto d_{CP} \equiv \frac{J_{CP}}{T_c^{12}} \sim 10^{-19}. \tag{55}$$

The suppression factor from the KM mechanism is thus much too strong to account for a baryon asymmetry as large as $\eta \sim 6 \times 10^{-10}$. The conclusion is that there must exist sources of CP violation beyond the Kobayashi Maskawa phase.

B. Intermediate summary

Baryogenesis implies that the SM cannot be the complete picture of particle interactions. There must exist additional sources of CP violation, beyond the single phase in the quark mixing matrix. The scalar sector must be extended beyond the single Higgs field of the SM. The relevant extension might take place at an energy scale not far above the electroweak breaking scale, inducing the baryon asymmetry via the EWPT. Or, they could take place at yet higher energy scales, via other mechanisms of departure from thermal equilibrium.

Historically, the first type of baryogenesis scenario which was developed in detail is the out-of-equilibrium decay of super-heavy particles. The basic idea is that at a very early stage of the expanding Universe, a super-heavy particle species X existed, the raction rate of which became smaller than the expansion rate H of the Universe at temperatures $T \gg T_{\mathrm{EW}}$. Therefore, these paricles decoupled from the thermal bath and became over-abundant. The decays of the X and of the antiparticle $\overline{X}$ are supposed to be CP- and B-violating, such that a net baryon number $\Delta B \neq 0$ is produced when $X, \overline{X}$ have decayed. This scenario has its natural setting in the framework of grand unified theories. Another viable variant is baryogenesis via leptogenesis through the lepton-number violating decays of a heavy Majorana neutrino.

VI. GUT BARYOGENESIS

The out-of-equilibrium decay scenario is natural in the context of grand unified theories. Grand unification aims at unifying the strong and electroweak interactions at some high energy scale. The hope is to get a consistent coupling unification, that is, a situation where all gauge couplings assume the same value at some high energy scale, and representation unification, that is, a situation where the SM matter representations combine into a smaller number of representations of the unifying gauge group. Indeed, in the context of supersymmetric GUT, the three SM gauge couplings unify at $M_{\rm GUT} \simeq 10^{16}$ GeV.

The matter representations of the simplest GUT theory, $G_{\rm GUT} = SU(5)$, and its extension with $SO(10)$ or E_6 groups, contain both quarks and leptons. Gauge bosons mediate transitions between members of such multiplets and, consequently, induce B-violating processes quite generically. Additional scalar representations are necessary to spontaneously break $G_{\rm GUT} \to G_{\rm SM}$. These super-heavy scalars often have baryon number violating and CP violating couplings to quarks and leptons.

The simplest GUT model is based, as mentioned above, on $G_{\rm GUT} = SU(5)$. The simplest verions of these theories are excluded because of a conflict between the constraints from coupling unification and those from proton decay. More complicated versions are still viable. Irrespective of this aspect, the minimal versions of $SU(5)$-GUTs cannot account for the observed baryon asymmetry because they conserve $B - L$. Thus, any baryon asymmetry induced by the GUT-particle decays is washed out by sphaleron interactions. In $SO(10)$ models, it is possible to overcome the proton decay problems and to generate a successful (i.e. $(B - L)$-violating) baryogenesis.

Rather than going into the details of a specific GUT, let us illustrate the baryogenesis mechanism with a toy model. Consider a super-heavy leptoquark gauge boson X with decay channels $X \to qq, \ell\overline{q}$:

$$\begin{aligned}
\mathcal{B}r(X \to qq) &= r, && B = +2/3, \\
\mathcal{B}r(X \to \ell\overline{q}) &= 1 - r, && B = -1/3, \\
\mathcal{B}r(\overline{X} \to \overline{qq}) &= \overline{r}, && B = -2/3, \\
\mathcal{B}r(\overline{X} \to \overline{\ell}q) &= 1 - \overline{r}, && B = +1/3. \qquad (56)
\end{aligned}$$

The baryon number produced per X or $\overline{X}$ decay is

$$\begin{aligned} B_X &= \frac{2}{3}r - \frac{1}{3}(1-r), \\ B_{\overline{X}} &= -\frac{2}{3}\overline{r} + \frac{1}{3}(1-\overline{r}), \end{aligned} \tag{57}$$

and the net baryon number produced is

$$\begin{aligned} \Delta B_X &\equiv B_X + B_{\overline{X}} = r - \overline{r} \\ &= \frac{\Gamma(X \to qq) - \Gamma(\overline{X} \to \overline{qq})}{\Gamma}. \end{aligned} \tag{58}$$

Obviously, if C or CP were conserved, then $\Delta B_X = 0$. It is natural to assume that C is already violated in the tree level interactions of X with fermions. CP violating couplings, on the other hand, can affect the $X, \overline{X}$ decay rates only beyond tree level. One can write generically

$$A(X \to qq) = A_0 \left(1 + |A_1/A_0| e^{i(\delta+\phi)}\right), \tag{59}$$

where A_0 is a tree amplitude, chosen here by convention to be real, A_1 is a one-loop amplitude, ϕ is a CP violating phase ('weak phase'), and δ is a CP conserving phase ('strong phase'), coming from the contribution of intermediate on-shell particles to the loop diagram. Then,

$$A(\overline{X} \to \overline{qq}) = A_0 \left(1 + |A_1/A_0| e^{i(\delta-\phi)}\right). \tag{60}$$

Putting (59) and (60) into (58), we obtain

$$\Delta B_X \propto \frac{2|A_0 A_1| \sin\delta \sin\phi}{\Gamma}. \tag{61}$$

As long as the interactions of these bosons, which include decays, inverse decays, annihilation, the B-violating reactions $qq \to \ell\overline{q}$, $\overline{qq} \to \overline{\ell}q$, etc., are fast compared to the expansion rate H, the X, $\overline{X}$ have thermal distributions and the average baryon number of the plasma remains zero. Therefore, the interactions of these bosons must be weak enough that they can fall out of equilibrium. This is a delicate issue, since these particles carry gauge quantum numbers and can couple quite strongly to the plasma of the early Universe.

The scenario works as follows. At temperatures $T \gg m_X$, the X, $\overline{X}$ particles have relativistic velocities and are assumed to be in thermal equilibrium. Then $n_X = n_{\overline{X}} \sim n_\gamma$. At lower temperatures, the X, $\overline{X}$ bosons become non-relativistic and, as long as they remain in thermal equilibrium, their densities get Boltzmann-suppressed with decreasing temperatue,

$n_X = n_{\overline{X}} \sim (m_X T)^{3/2} \exp(-m_X/T)$. Because $\Gamma_{\rm ann} \propto n_X$, the total rate $\Gamma_X \sim \alpha m_X$ of X and $\overline{X}$ decays is the relevant rate to compare with H. If $\Gamma_X < H$, an excess of X, $\overline{X}$ with respect to their equilibrium numbers will develop. The X, $\overline{X}$ drift along in the expanding Universe for a little while and decay. Notice that the inverse decays, $f \to X$, $\overline{f} \to \overline{X}$, by which bosons are created again by quark-quark annihilations, are blocked, because the fraction of these fermions with sufficient energy to produce a super-heavy boson is Boltzmann-suppressed for $T < m_X$. At the time of their decay, $t \sim \Gamma_X^{-1}$, there is quite an over-abundance: $n_X = n_{\overline{X}} \sim n_\gamma(T_{\rm decay})$. Using that the entropy density $s \sim g_* n_\gamma$, one gets the produced asymmetry:

$$\frac{n_B}{s} \sim \frac{\Delta B n_\gamma}{g_* n_\gamma} \sim \frac{\Delta B}{g_*}, \tag{62}$$

where ΔB is the baryon number produced per boson decay. For a (GUT) extension of the SM, one may expect that g_* is in the $10^2 - 10^3$ range. Thus, only a tiny CP asymmetry, $\Delta B \sim 10^{-8} - 10^{-7}$, is required to obtain $n_B/s \sim 10^{-10}$. Of course, these crude estimates must be made quantitative by computing the relevant reaction rates using a specific particle physics model, and tracking the time evolution of the particle densities by solving the Boltzmann equations.

The above condition for the decoupling of the X particles from the thermal bath, $\Gamma_X \sim \alpha m_X < H$, translates into a condition on the mass of the spin-one gauge boson: $m_X > \alpha g_*^{-1/2} m_{\rm Pl} \sim 10^{16}$ GeV. The lower bound on super-heavy scalars is typically weaker.

There are several pitfalls that constrain this type of baryogenesis. First, a scenario that makes an attempt to explain η by a mechanism that operates above $T_{EW} \sim 100$ GeV, must involve interactions that violate $B - L$. In the context of GUTs, minimal SU(5) models conserve $B - L$ and cannot explain η, but SO(10) models can. These models have several attractive features, also in the context of leptogenesis.

GUT baryogenesis may be in conflict with *inflation*. This is a serious problem. An essential assumption in the above scenario was that at very high temperatures, $\gg m_X$, the X, $\overline{X}$ particles were in thermal equilibrium, and were as abundant as photons. If these particles are the super-heavy gauge bosons or scalars of a GUT, this assumption may be wrong. It might be that the temperature of the quasi-adiabatically expanding plasma of particles in the early Universe was always smaller than $M_{\rm GUT}$. There are a number of reasons to believe that the energy density of the very early Universe was dominated by vacuum energy, which led to exponential expansion of the cosmos. This is the basic assumption of inflationary models. These models solve a number of fundamental cosmological problems, including the monopole problem. A number of GUTs predict super-heavy, stable magnetic monopoles. Their contribution to the energy density of the Universe would over-close it, but that is not observed. Inflation would strongly (exponentially) dilute the monopoles, together with other particles, leaving empty space. At the end of inflation, the vacuum energy is converted through quantum fluctuations into pairs of relativistic particles and antiparticles which then thermalize. This process is called *reheating*, and it can be characterized by an energy scale called the reheat temperature $T_{\rm RH}$. If $T_{\rm RH} > M_{\rm GUT}$, the monopoles are re-created. If $T_{\rm RH} \ll M_{\rm GUT}$, then the reproduction of the X, $\overline{X}$ super-heavy gauge bosons and of GUT-scale scalars, which were to initiate baryogenesis, is inhibited or should, at least, be suppressed.