

Baryogenesis

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Abstract. Developments in understanding of baryogenesis are reviewed. We start with early motivations and the proposals in the context of GUTs. Next, the importance of the sphaleron solution and its implications are discussed. Studies of the Standard Model reveal that the latter has a Higgs structure incompatible with existence of observed B asymmetry. We then discuss a generic scenario for electroweak baryogenesis relying on bubble wall dynamics. We also summarise the status of the MSSM, and alternative scenarios utilising topological defects as the source of non-equilibrium behaviour and leptogenesis

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1. Why particularly, the sky? or cosmology as the laboratory for particle physics

Gauge symmetry of particle interactions and spontaneous breakdown of the same in unified theories imply interesting collective phenomena at high temperatures. Thus in the early Universe, we expect phase transitions, exotic states of matter, topological defects and so on. Some of these phenomena are expected to leave behind observable imprints. For example inflation results from the unusual equation of state obeyed by vacuum energy, or stable cosmic strings can bias the lumping of matter. This talk will focus on baryon asymmetry resulting from nonequilibrium conditions existing in the expanding Universe.

The baryon asymmetry can also be formulated as the baryon number to entropy density ratio of the Universe, $\Omega = n_B/s$. From direct observation of the number of galaxies, the average number of stars per galaxy and so on, Ω has the value 10^{-10} . This is corroborated by the relative abundance of light nuclei in Big Bang nucleosynthesis, $1.5 \times 10^{-10} < \Omega < 7 \times 10^{-10}$. This unnaturally small value demands a microscopic explanation. Intensive developments over the past decade have yielded important information about particle phenomenology, mostly in the form of vetoes. We present here a review for the non-expert.

Investigating the consistency of a given model of particle physics with the observed baryon asymmetry requires checking for the nature of high temperature phase transition in the theory and also for existence of requisite particle species content. We focus here on the general nature of these requirements, how they arise and then summarise the results for

some of the popular models. The interesting result is that most of the models accommodate the observed baryon asymmetry only with fine tuning of parameters.

1.1 Asymmetry observed

The need to understand baryon asymmetry arises in the first instant from the absence of any antimatter. How do we know the asymmetry observed in the immediate astral neighborhood is Universal? The three broad classes of data in this connection are

- The content of cosmic rays is baryonic. The observed ratio of \bar{p} to p in cosmic rays is 10^{-4} . This is consistent with secondary production $p + p \rightarrow 3p + \bar{p}$.
- If neighboring clusters of galaxies happened to contain matter and anti-matter, this would produce diffuse γ -ray background. This is not observed.
- Perhaps the regions of anti-matter are completely separated. This would reflect in inhomogeneities in the cosmic microwave background radiation (CMBR). These are also not observed.

We begin in §2 with a review of the general requirements for baryogenesis (B -genesis) followed by the more specific recipe of grand unified theory (GUT) baryogenesis. Section 3 contains the significance of the sphaleron solution and the beginning of the modern attack on the problem. It is shown that the Standard Model (SM) Higgs structure is incompatible with observed baryon asymmetry on fairly general grounds. Section 4 discusses the potential for B -genesis at the electroweak scale in extensions of the SM, with the example of the two Higgs doublet model (2HDM). The current status of the MSSM is also discussed. In §5 we present briefly the status of other mechanisms, viz., B -genesis induced by topological defects, and B -genesis through leptogenesis [5]. Section 6 contains the conclusion.

2. Baryogenesis in the beginning

Several peculiar features of the nuclear interactions such as parity and CP violation became known by the early 1960's. The discovery of the cosmic microwave background (CMBR) around the same time was confirming a cosmological arrow of time. Another puzzle that was noted around this time was the fact that baryon number, a symmetry of the strong and weak forces was only an algebraic symmetry and not a gauge symmetry which would have a fundamental justification for its exact conservation. Combined with CP violating interactions and the expanding Universe this presented the possibility of an *explanation* from physical laws of the baryon asymmetry [1]. A first explicit model by Sakharov [2] set forth the following salient issues, the so called Sakharov criteria: 1. B violating and 2. C violating interactions, 3. CP violation 4. Out of equilibrium conditions, i.e., the state of the system must be time asymmetric. The last ensures that CP violation becomes effective. We recapitulate here the model of Yoshimura [3] and Weinberg [4] proposed in the context of the GUTs.

Baryon number violation: Consider a species X that has two different modes of decay (even if the absolute baryon number X is not defined, it is the difference in the baryon number of the final states that matters).

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$$\begin{aligned} X &\longrightarrow qq & \Delta B_1 &= 2/3 \\ &\longrightarrow \bar{q}\bar{l} & \Delta B_2 &= -1/3. \end{aligned}$$

Charge conjugation violation: The inequality of the amplitudes for the charge conjugated processes, $\mathcal{M}(X \rightarrow qq) \neq \mathcal{M}(\bar{X} \rightarrow \bar{q}\bar{q})$.

CP violation:

$$r_1 = \frac{\Gamma_1(X \rightarrow qq)}{\Gamma_1 + \Gamma_2} \neq \frac{\bar{\Gamma}_1(\bar{X} \rightarrow \bar{q}\bar{q})}{\bar{\Gamma}_1 + \bar{\Gamma}_2} = \bar{r}_1.$$

It is clear that phases appearing in the vertices of an effective Hamiltonian cannot enter in the rate formulae in the Born approximation. A crucial observation of [4] was that it must be the interference of a tree diagram with a higher order diagram with a CP violating phase which will result in $r \neq \bar{r}$.

Out of equilibrium conditions: The effect of the dominant forward reaction would be nullified by excess build-up of the species which would drive the reverse reaction and establish an equilibrium with the anti-species unless the state of the system prevented this from happening.

If these conditions are satisfied the rate of baryon number violation is

$$\begin{aligned} \mathcal{B} &= \Delta B_1 r_1 + \Delta B_2 (1 - r_1) + (-\Delta B_1) \bar{r}_1 + (-\Delta B_2) (1 - \bar{r}_1) \\ &= (\Delta B_1 - \Delta B_2) (r_1 - \bar{r}_1). \end{aligned} \quad (1)$$

The decay rate for the X particle in the early Universe has the temperature dependence $\propto 1/T$ while the Hubble parameter H varies as $\propto T^2$. The condition 4 above comes to be satisfied when $\Gamma \lesssim H$, which happens in GUTs very early in the Universe. The resulting Ω is estimated [4] to be $\sim \mathcal{B}/g_*$ where g_* is the effective number of degrees of freedom in equilibrium at that epoch. In order to obtain quantitative results the Boltzmann equations must be integrated taking account of each decay mode and the chemical potential of each of the conserved numbers.

This GUT based scenario has several problems if inflation has to also occur. However, more recently, the same scenario has been applied to the lepton number violation by heavy majorana neutrinos [5], as discussed in §5. In the meantime, an important discovery of the mid-1980's has completely revolutionized our understanding of the fate of the baryon number at a much lower energy scale, viz., the electroweak scale. This is what we turn to next.

3. The anomalous baryon number

Since the Standard Model is chiral it has the potential presence of anomalies. All the gauge currents and most of the global conserved currents are indeed anomaly free in the Standard Model. However, the number $B + L$ is in fact anomalous. Specifically, one finds that

$$\partial_\mu j_{B+L}^\mu = \frac{g^2}{32\pi^2} \epsilon^{\mu\nu\rho\sigma} F_{\mu\nu}^a F_{\rho\sigma}^a, \quad (2)$$

where g is the gauge coupling. This anomaly of the fermionic current is associated with another interesting fact of non-abelian gauge field theory discovered by Jackiw and Rebbi. There exist configurations of the gauge potentials A_μ^a each of which is pure gauge, i.e., with physical field strengths $F_{\mu\nu}^a = 0$, but which cannot be deformed into each other without turning on the physical fields. Such pure gauge vacuum configurations can be distinguished from each other by a topological charge called the Chern–Simons number

$$N_{C-S} = \frac{g^3}{32\pi^2} \epsilon^{ijk} \int d^3x \left\{ F_{ij}^a A_k^a - \frac{2}{3} \epsilon^{abc} A_i^a A_j^b A_k^c \right\} \quad (3)$$

which has integer values in vacuum configurations. On the other hand the RHS of the anomaly eq. (2) is equal to a total divergence $\partial_\mu K^\mu$ where

$$K^\mu = \epsilon^{\mu\nu\rho\sigma} \left\{ F_{\nu\rho}^a A_\sigma^a - \frac{2}{3} g \epsilon^{abc} A_\nu^a A_\rho^b A_\sigma^c \right\}. \quad (4)$$

Hence if we begin and end with configurations with $F_{\mu\nu}^a = 0$, then the change

$$\Delta Q_{B+L} \equiv \Delta \int j_{B+L}^0 d^3x = -\Delta \left(\frac{g^2}{32\pi^2} \int K^0 d^3x \right) = -\Delta N_{C-S} \quad (5)$$

using (2), (3) and (4). Thus the violation of the axial charge by unit occurs because of a quantum transition from one pure gauge configuration to another. At normal temperatures such transitions do not occur. As shown by 'tHooft, the tunneling probability is highly suppressed due to instanton effects. In case of gauge Higgs system with symmetry breaking, a barrier exists between the equivalent vacua. There is a configuration called the sphaleron [6,7] which is supposed to be a time independent configuration of gauge and Higgs fields which has maximum energy along a minimal path joining sectors differing by unit Chern–Simons number. In figure 1 the sphaleron corresponds to the peaks where N_{C-S} has half-integer values. It provides a lower bound on the amount of energy required for passage *over* the barrier. The sphaleron solution of the field equations has been discussed in detail [8] in many references. Here we shall focus on its consequences.

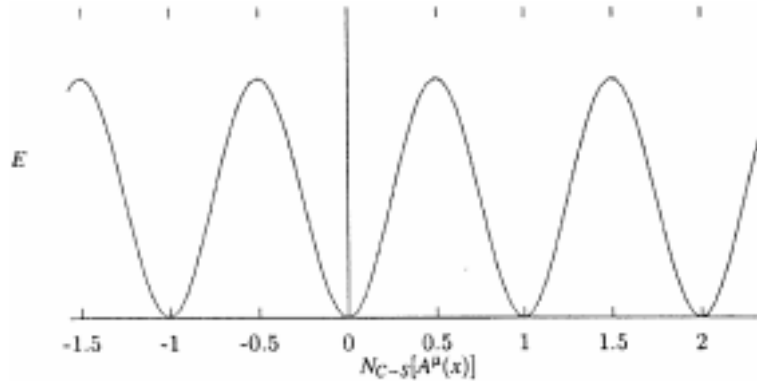


Figure 1. Energy profile of gauge field configurations.

3.1 The implications of the sphaleron

The energy of the sphaleron E_{sph} can be calculated from numerical solution of the gauge field equations satisfying the appropriate boundary conditions [7]. It is a weak function of λ/g^2 , i.e., of m_H/m_W in the SM. It varies over 7 to 14 TeV as this parameter varies from small values to infinity. At moderately high temperatures, by which is meant $m_W \ll T \ll m_W/\alpha_W \sim E_{\text{sph}}$, the barrier crossing rate is suppressed by a Boltzmann factor, and in saddle point approximation, the rate per unit four-volume as [9]

$$\Gamma = A\kappa(\mathcal{N}\mathcal{V})T^4 e^{-E_{\text{sph}}(T)/T}. \quad (6)$$

Here the \mathcal{N} and \mathcal{V} denote contributions from zero-mode integration, κ is the contribution of fluctuations and A is a dimensionless prefactor.

Suppose we have some mechanism for producing baryons above the electroweak scale. Once the symmetry breakdown occurs in the early Universe, sphalerons become possible. Then B number begins to deplete at a rate determined by the above rate (6). In order for the B -asymmetry to survive we need that the above rate is really too slow compared to the expansion rate of the Universe. We refer to this as the ‘wash-out’ constraint, and is the single most important implication of sphaleron physics, following entirely from the SM [10]. As a function of temperature, E_{sph} is given parametrically by

$$E_{\text{sph}}(T) \sim B \frac{m_W(T)}{\alpha_W(T)} \propto \frac{v(T_c)}{T_c}, \quad (7)$$

where B is a dimensionless quantity of order 1 and $v(T)$ is the temperature dependent vacuum value of the Higgs. E_{sph} should be large enough to prevent wash-out, which requirement is shown numerically to translate to

$$\frac{v(T_c)}{T_c} \gtrsim 1. \quad (8)$$

Since the order parameter has zero vacuum value before the phase transition, large $v(T_c)$ immediately after means that the phase transition should be strongly first order. From perturbative effective potential one learns that the SM phase transition is only mildly first order for $m_H \gtrsim 90$ GeV, the experimentally acceptable range. Thus the SM is contradicted by the presence of B -asymmetry in the Universe, unless $B - L$ is also not conserved, or that a net primordial value of this number has pre-existed the electroweak phase transition. In either case, the SM is demonstrated to be insufficient for a physical explanation of the B asymmetry.

Above arguments relying on saddle point perturbation theory are physically transparent but only suggestive. More recently, a non-perturbative approach using lattice methods has been developed [11] and so far seems to not contradict the above conclusions.

3.2 The anomaly at higher temperatures

It is also natural to ask about the physics of the anomaly above the electroweak temperatures. With $m_H \sim 120$ to 150 GeV, the phase transition temperature T_c is 100 GeV. Above

this temperature, $v(T_c)$ vanishes. Then there is no sphaleron and it is reasonable to assume that there is no barrier either. Assuming that coherent fluctuations can exist on the scale of the magnetic screening length in the non-abelian plasma, the rate of N_{C-S} changing transitions is given in dimensional analysis by

$$\Gamma = C(\alpha_W T)^4. \quad (9)$$

In fact there are recent arguments to the effect that the prefactor C is $\sim \alpha_W$. In either case, there is no suppression and we assume that $B + L$ is freely violated. This conclusion on general grounds has been verified by real time simulations in lattice gauge theory [12]. It is shown that the N_{C-S} value oscillates around a given integer value corresponding to a specific potential well, but every once in a while, makes a sharp transition to a neighbouring well. The rate of transitions is then found to accord with above formula.

This raises an important issue in the explanation of baryon asymmetry of the Universe. $B + L$ asymmetry generated by any mechanism above the electroweak scale will be wiped out by the unsuppressed anomalous transitions at high temperatures. There are two solutions to this dilemma. One is to find physical mechanisms that produce a net $B - L$. A very attractive candidate in this class is the B -genesis through leptogenesis scenario to be discussed in §5. A tantalising possibility first suggested by Kuzmin, Rubakov and Shaposhnikov [13] is the possibility of processes operating at the electroweak scale itself that would produce the required baryon asymmetry. Since the interactions involved are or will soon be within the range of accelerator energies, this possibility is very exciting. Appropriately, it generated a decade long intense search for mechanisms that would work at electroweak phase transition temperatures, as discussed in the following section.

4. Electroweak baryogenesis

Any mechanism for B -genesis operating at the electroweak temperatures differs in an important way from the Yoshimura–Weinberg proposal in that we need a new source of time asymmetry. The point is that at the GUT scale the expansion of the Universe is fast enough to exceed the decay rates of relevant particle species, while at the electroweak epoch the expansion rate is orders of magnitude smaller. Any decays which are out of equilibrium at this epoch will in any case prove insufficient to the task.

A novel feature of the the electroweak physics is that the symmetry breaking transition could be first order, in the sense that the transition involves latent heat and the true vacuum is formed by tunneling from the false vacuum. Such formation would typically proceed in the shape of spontaneously forming ‘bubbles’ as per the mechanisms of Coleman [14] and Linde [15] (see figure 2). In the context of this picture we expect the expanding bubble walls to provide the out of equilibrium conditions needed for baryogenesis.

We now discuss this in greater detail. What needs to be verified is whether the given phase transition is indeed first order. One way to study this is to study the effective potential at finite temperature. The free energy of the Higgs field is given in the field theoretic formalism by the finite temperature effective potential. The one loop contribution to the effective potential of the Higgs from integrating out other particles is of the form

$$\Delta V_T^i \sim T^4 \int dx x^2 \ln(1 \pm e^{-\sqrt{x^2 + y_i^2}}), \quad (10)$$

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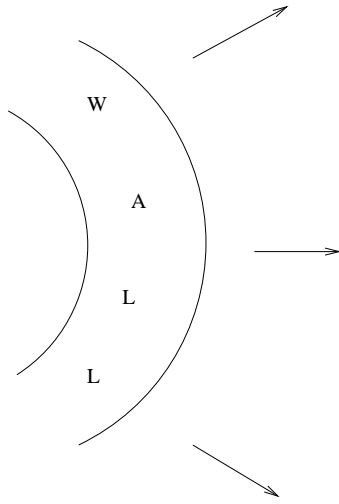


Figure 2. Cartoon of a bubble wall. The region to the left has $\langle\phi\rangle \neq 0$ while the region to the right into which the bubble is expanding has $\langle\phi\rangle = 0$. The region labelled WALL has interpolating values of $\langle\phi\rangle$.

where $-$ is for bosons and $+$ for fermions, and for each species i ,

$$y_i = M_i \langle\phi\rangle / v_0 T, \quad (11)$$

M_i being the respective mass and v_0 is the zero temperature expectation value of the Higgs field, $v_0 = 246$ GeV. If y_i can be treated as small parameters, the effective potential for the Higgs becomes

$$V_{\text{eff}}^T[\phi] = D(T^2 - T_0^2)\phi^2 - ET\phi^3 + \frac{\lambda_T}{4}\phi^4, \quad (12)$$

where D , T_0 and E (all positive) are parameters determined in terms of the zero temperature masses of the gauge bosons and the top quark. λ_T is only mildly temperature dependent. At $T \gg T_0$, only the $\phi = 0$ minimum exists. For $T < T_0$, this minimum is destabilized. But there exists a $T_c > T_0$, when another minimum with nontrivial value of ϕ becomes possible, and

$$\frac{\phi(T_c)}{T_c} = \frac{2E}{\lambda} \simeq \frac{m_W^2}{m_H^2}. \quad (13)$$

A barrier separates the two vacua and tunneling across the barrier via thermal and quantum fluctuations becomes possible. Whenever tunneling to the true vacuum occurs in any region of space, it results in a bubble. According to a well developed formalism [15], the tunneling probability per unit volume per unit time is given by

$$\gamma = CT^4 e^{-S_{\text{bubble}}}, \quad (14)$$

where S_{bubble} is value of

$$S = 4\pi \int r^2 dr \left\{ \frac{1}{2} \phi'^2 + V_{\text{eff}}^T[\phi] \right\} \quad (15)$$

extremised over ϕ configurations which satisfy the bubble boundary conditions $\phi(r = 0) = \phi(T_c)$, $\phi \rightarrow 0$ as $r \rightarrow \infty$. Once a bubble forms, energetics dictates that it keeps expanding, converting more of the medium to the true vacuum. The expansion is irreversible and provides one of the requisite conditions for producing baryon asymmetry.

4.1 Two obstacles to SM B-genesis

Thus the SM seems to possess all the ingredients necessary for producing B -asymmetry. But there are two important issues to be faced. First we note that the CP violation available in SM is far too small. A model independent dimensionless parameter characterising the scale of this effect has the value [16] $\delta_{CP} \sim 10^{-20}$. Thus the explanation of B -asymmetry is not possible purely within the SM.

But even assuming that there may be non-SM sources of CP violation, an important question is whether the phase transition at the electroweak scale is first order or second order. In the formalism discussed above we used V_{eff}^T extensively. Unfortunately, perturbative V_{eff}^T is not always a reliable tool for studying the nature of the phase transition. One way to understand this problem [17] is to observe that the expansion parameters in (11) are not small. Specifically, the gauge boson contributions are expanded in powers of (number) $\times m_H^2/m_W^2$. Since the ratio of the masses is now known to be at least unity, the validity of the perturbation series hinges crucially on numerical factors. The only way to test the validity would be to carry the expansion to higher orders. In particular, it is found that the two-loop corrections are not small compared to the one-loop contributions [18,19].

Several other approaches have been tried. One is to improve the perturbation theory. This is done by the technique of dimensional reduction possible in high temperature expansion. This method exploits the effective three dimensional nature of the Euclidean thermal theory if only the zero Matsubara frequency modes are retained [20,21]. The higher frequency modes can be accounted for in a modification of this approach by including additional effective terms [22]. The most direct approach involves lattice simulations, both in $4d$ theory and the effective $3d$ theories. The upshot of these is that the SM electroweak phase transition is mildly first order for $m_H \sim 100$ GeV, but becomes second order for higher m_H . Since the accelerator limits on m_H have exceeded 90 GeV, we clearly need physics beyond the SM, both for the CP violation as also for obtaining out of equilibrium conditions.

4.2 Electroweak baryogenesis in other models

We seek an extension of the SM that satisfies the requirements spelt out above. In any model which suits this purpose a detailed mechanism is necessary to bring the non-equilibrium behaviour of the bubble walls into play. In particular, completely different mechanisms would be successful depending upon the nature of the bubble walls. There are two broad classes:

Fast moving, thin walls: If the wall thickness is less than the inverse thermal mass of some particle species, the interaction of the latter with the wall can be treated in the sudden approximation. The scenario considered by Cohen Kaplan and Nelson [24] involves scattering of neutrinos from the expanding walls in this approximation. Lepton number violation occurs if the neutrino obtains a majorana mass from the scalar forming the wall. CP phase enters from the majorana coupling, and the resulting rate of reflection for ν_L can be different from that for the CP conjugate state $\bar{\nu}_R$. Finally the excess L number generated in front of the wall is converted to B -number by the unsuppressed anomalous transitions which continue to set $B + L = 0$. Once the B excess is engulfed by the interior of the bubble, it is protected if sphalerons are sufficiently heavy.

Slow moving, thick walls: If the wall thickness is large compared to thermal correlation lengths, coherent gauge and scalar fluctuations are possible within the walls. Thus Chern–Simons number changing transitions are possible within the thickness of the wall. Further, as recognised by McLerran, Shaposhnikov, Turok and Voloshin [25], the nontrivial background field of the wall induces a non-trivial chemical potential for the N_{C-S} as discussed later. If the scalar sector of the theory also contains CP violation then we may expect a non-trivial asymmetry to result. The B number thus generated falls in the interior of the bubble and is again assumed to not be depleted significantly by the sphalerons of the theory.

4.3 Baryogenesis in 2HDM

We shall now discuss the second scenario in greater detail. A mechanism for biasing N_{C-S} is available in scalar couplings but to achieve the desired CP violation we need the two Higgs doublet model (2HDM). Consider the diagram of figure 3. The heavy quark in the loop is assumed to couple only to ϕ_1 to avoid flavour changing neutral currents. This diagram induces in the effective action a term which has a nontrivial expectation value only where the fields are space-time dependent.

$$\begin{aligned} \Delta S = & \frac{-7}{4} \zeta(3) \left(\frac{m_t}{\pi T} \right)^2 \frac{g}{16\pi^2} \frac{1}{v_1^2} \\ & \times \int (\mathcal{D}_i \phi_1^\dagger \sigma^a \mathcal{D}_0 \phi_1 - \mathcal{D}_0 \phi_1^\dagger \sigma^a \mathcal{D}_i \phi_1) \epsilon^{ijk} F_{jk}^a d^4x, \end{aligned} \quad (16)$$

where m_t is the top quark mass, ζ is the Riemann zeta function, and the σ^a are the Pauli matrices. Assuming the time derivatives dominate compared to space derivatives, in the gauge $A_0^a = 0$, we can rewrite this in the form

$$\begin{aligned} \Delta S = & \frac{-i7}{4} \zeta(3) \left(\frac{m_t}{\pi T} \right)^2 \frac{2}{v_1^2} \\ & \times \int dt [\phi_1^\dagger \mathcal{D}_0 \phi_1 - (\mathcal{D}_0 \phi_1)^\dagger \phi_1] N_{CS} \\ \equiv & \mathcal{O} N_{CS}. \end{aligned} \quad (17)$$

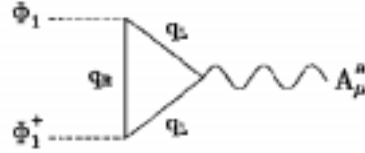


Figure 3. A nontrivial contribution to the S_{eff} .

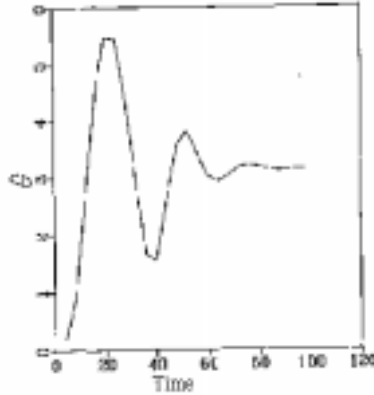


Figure 4. Oscillations of the CP phase at a point swept by the advancing wall.

This leads to [25] an estimate of $n_B/n_\gamma \sim 10^{-3} \alpha_W^4 \sin 2\xi(T_c)$ where ξ is the value of CP violating angle at the relevant temperature. If $\sin 2\xi(T_c)$ is $O(1)$, this leads to an answer in the correct range of values. It may be noted that the mechanism in spirit is very similar to the generic ‘spontaneous baryogenesis’ proposal of Cohen, Kaplan and Nelson [27], but which did not consider the specifics of a phase transition bubble wall.

In [25] it was assumed that the relative phase between the vacuum expectation values of the Higgs remains static. This however is too restrictive and leads to a suppression of this effect in the lowest adiabatic order as noted by the authors. An appropriate generalisation is to let this phase θ be time dependent and make the ansatz [26]

$$\phi_1^0 = \rho(r, t) \cos \gamma e^{-i\theta(t)} \tag{18}$$

$$\phi_2^0 = \rho(r, t) \sin \gamma e^{i\omega(t)}. \tag{19}$$

Here γ is the particular direction in the $\phi_1^0 - \phi_2^0$ plane along which the scalar fields tunnel to form a bubble. The angles ω and θ are related by the unitary gauge requirement that the Goldstone boson eaten by the Z^0 should be orthogonal to the remaining physical pseudoscalar: $\partial_\mu \omega = (\rho_1/\rho_2)^2 \partial_\mu \theta$ [27]. The variation of θ as the wall sweeps past a particular point is shown in figure 4 [26]. The first peak of this graph occurs over the time

scale it takes for the wall to sweep past a particular point. The remaining oscillations occur in the wake of the wall. Thus, $\text{Im } \phi$ which was static and gave no leading order contribution is now replaced by a phase which oscillates over its full range of values. This $\Delta\theta$ acts as the transient CP violation parameter.

To estimate the generated baryon asymmetry we proceed as in the general case outlined earlier. The number of fermions created per unit time in the bubble wall is given by

$$B = \kappa(\alpha_w T)^4 l S \times \frac{1}{T} \frac{\mathcal{O}}{l}, \quad (20)$$

where [25,27], l and S are the thickness and the surface area of the bubble wall respectively. From which we get the baryon to photon ratio to be

$$\Delta \equiv \frac{n_B}{s} \simeq \frac{1}{g_*} (\alpha_W)^4 \mathcal{O} \simeq 10^{-8} \times \left(\frac{E}{K_1} \right)^2 \Delta\theta, \quad (21)$$

where we have used $\alpha_W \sim 10^{-3/2}$ and $g_* \sim 100$. The parameters E and K_1 are given in terms, respectively, of the cubic and quartic couplings in the 2HDM. This answer easily accommodates the observed value of this number.

It is worth emphasizing the physics of this answer which is robust against changes in the specific particle physics model [28]. The thermal rate contributes 10^{-6} through α_W^4 , and another 10^{-2} is contributed by N_{eff} . Further, the temperature induced cubic coupling E is generically small compared to K_1 so that in ‘bubble wall’ scenarios we would never need CP violation larger than 10^{-3} , and more likely much less.

4.4 Baryogenesis in the MSSM

The minimal supersymmetric standard model is an attractive extension of the SM. Extensive investigations [37–39] have been made of this model leading to reasonably definitive conclusions about the viability of B -genesis in it. The sphaleron does occur [40] in the model and high temperature anomalous processes are also expected to occur. As was seen in the SM case, if the Higgs is too heavy it tends to make the phase transition second order. In MSSM, the phase transition is first order if the lightest Higgs and the scalar superpartner of the top quark, the stop are sufficiently light. The CP violation cannot arise in the Higgs sector but can occur in the soft SUSY breaking parameters associated with mixing of left and right stops. However, too large a mixing between L and R stops reduces the strength of the first order phase transition. The baryon creation mechanism relies on wall motion and the interaction of CP odd currents with the wall. It is shown that the generated B asymmetry is proportional to change in the $\tan\beta$ parameter (the ratio of the vacuum values of the two Higgs) across the bubble walls. If this variation has to be significant, say $\sim 10^{-2}$ then m_A , the mass of the pseudoscalar Higgs $\lesssim 150$ –200 GeV. On the other hand, a value of m_A less than above range tends to weaken the strength of the first order phase transition. Thus baryogenesis is viable in the model in a fairly restricted window in the parameter space. Numerical calculations then show that the lightest Higgs should be in the range 85–110 GeV, accessible to LEP2 and the right stop mass in the range 120–172 GeV. Further, $\tan\beta$ should be $\lesssim 4$. Since the perturbative consistency of the model up to GUT scales requires $\tan\beta \gtrsim 1.2$, this is a significant constraint.

5. Other approaches

Topological defects: It was pointed out in [29] that the presence of cosmic strings at the electroweak phase transition would simplify some of the nettlesome issues related to bubble wall thickness and velocity. The bubbles described by the Coleman–Linde formalism are distributed in their sizes and the times of their occurrence. The tunneling probability is a very sensitive function of the parameters in the potential and it is not possible to work with a generic B -genesis scenario but each model must be examined in detail. By contrast, if the phase transition is induced by the presence of cosmic strings then the parameters of the bubbles are *decided entirely by electroweak physics* regardless of the detail of the unified model giving rise to the strings.

A more direct role for cosmic strings was proposed in [30] where unstable cosmic strings themselves act as locations of non-equilibrium processes. A set of model independent defect induced scenarios was also proposed [31,32]. Recent numerical work [33] shows that of these the string mediated mechanism is not efficient enough at B -asymmetry production. However, the possibility that unstable domain walls could lead to B -genesis is still open [34,35]. In particular, many interesting topological objects including unstable domain walls have been shown to occur in the left–right symmetric model [36] where such scenarios can work.

Leptogenesis: The early proposal of Fukugita and Yanagida of B -genesis from leptogenesis gains a tantalising prospect after the recent strong indications that the neutrino has mass. Models that accommodate the indicated minuscule masses should incorporate the see-saw mechanism. If this is true there must be heavy neutrinos possessing majorana masses. Further the Yukawa couplings giving rise to such masses can themselves be CP violating. We then have all the ingredients necessary to play out the scenario of [5], further developed in [41]. We report here the results of Buchmüller and Plümacher [42]. Consider the generic leptonic Yukawa couplings

$$\mathcal{L}_{\text{Yuk}} = -\bar{l}_L \tilde{\phi} g_l e_R - \bar{l}_L \phi g_\nu \nu_R - \frac{1}{2} \nu_R^c M \nu_R + \text{h.c.} \tag{22}$$

This results in mass $m_l = g_l v$ for the charged lepton and Dirac mass $m_D = g_\nu v$ for the neutrino if $\langle \phi \rangle = v$. M results from unknown physics but which is a CP violating coupling. By see-saw mechanism we get $m_\nu \sim m_D^2/M$ and for the heavy species, $m_N \sim M$.

Consider the decays of the N as in figure 5, taken from [42]. From the interference of these diagrams (see also [43]) we get for the quantity \mathcal{B} of eq. (1),

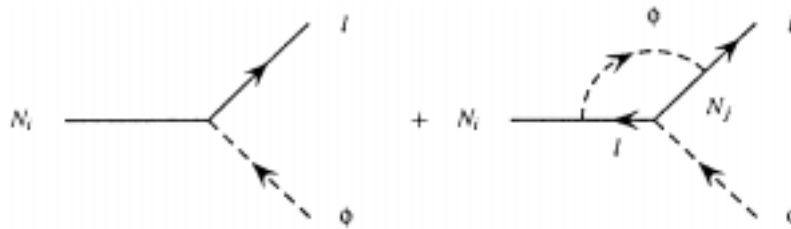


Figure 5. The contributions to ν decay whose interference results in CP violating rates.

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$$\mathcal{B} = \frac{r - \bar{r}}{r + \bar{r}} \sim \frac{1}{v^2 m_D^2} \text{Im}(m_D^\dagger m_D)^2. \quad (23)$$

In a model with $M \sim 10^{14}$ GeV and with the assumption that the neutrinos obey the hierarchy $m_{\nu_\mu} \sim 10^{-3}$ eV, $m_{\nu_e} \sim 10^{-5}$ eV and $m_{\nu_\tau} \sim 0.1$ eV, it has been shown that $\Omega \sim 10^{-10}$ is possible to achieve.

6. Conclusion

All the physics needed for arriving at the important phenomenon of sphalerons and of the high temperature anomaly was known since mid-70's. However a systematic investigation did not begin until the mid-80's. The intensive development of many conceptual issues and of calculation techniques now puts us in a position to veto models of particle physics based on their potential for dynamical explanation of the baryon asymmetry. Some of the important techniques include improved calculations of the effective potential for the study of the nature of phase transition and the evolution of true vacuum bubbles as sites of B -symmetry generation. These methods however are perturbative in nature and if the problems are really hard lattice methods may be inevitable.

The generic but non-supersymmetric 2HDM is adequate for B -genesis at the electroweak scale. This however is not natural in any unification scheme. The MSSM has been explored extensively and this has resulted in constraints on $\tan \beta$, and masses of the lightest Higgs and the stop. The other very promising alternative is leptogenesis, since it may help to constrain neutrino physics, for which interesting experimental evidence is emerging.

A generic mechanism applicable to SUSY GUTs and supergravity models is the Affleck–Dine mechanism [44]. It bears investigation in specific models, where it may place specific constraints on the model. It was not possible to discuss this here due to lack of space.

Finally, the reader is referred to some of the more extensive recent reviews such as [45–47].

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