ON ANOMALOUS ELECTROWEAK BARYON-NUMBER NON-CONSERVATION IN THE EARLY UNIVERSE

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We estimate the rate of the anomalous electroweak baryon-number non-conserving processes in the cosmic plasma and find that it exceeds the expansion rate of the universe at $T >$ (a few)$\times 10^2$ GeV We study whether these processes wash out the baryon asymmetry of the universe (BAU) generated at some earlier state (say, at GUT temperatures). We also discuss the possibility of BAU generation by the electroweak processes themselves and find that this does not take place if the electroweak phase transition is of second order No definite conclusion is made for the strongly first-order phase transition We point out that the BAU might be attributed to the anomalous decays of heavy ($M_F \gtrsim M_W/\alpha_W$) fermions if these decays are unsuppressed

1. Among various effects related to the θ-vacuum structure in gauge theories [1–3], the anomalous non-conservation of fermion quantum numbers is of particular interest. In the vacuum sector, this phenomenon is associated with instantons [4] describing tunneling transitions between vacua with different quantum numbers. In the weakly coupled theories, the probabilities of these transitions are exponentially suppressed; in particular, the corresponding suppression factor in the standard electroweak theory is $\exp(-4\pi/\alpha_W)$, where $\alpha_W = \alpha/\sin^2\theta_W$.

However, if the energy of the system is large enough, the system can, in principle, pass over the barrier between the different vacua instead of penetrating through the barrier [5–8]. In that case the rate of the anomalous non-conservation of the fermion number can be unsuppressed. This possibility was first discussed in the context of high-energy collisions [5]; however, it remains unclear whether the coherent gauge fields driving the non-conservation can be formed in the course of a collision On the other hand, it has been argued that the anomalous electroweak baryon-number non-conservation is indeed unsuppressed in decays of heavy fermions, or bound states of fermions like technibaryons [6,7] Yet another possibility, which has been mentioned in refs [8–10] is that the rates of the electroweak baryon-number non-conservation processes are large at sufficiently high temperatures It is worth noting that since the height of the barrier between the electroweak vacua with different baryon numbers is of the order M_W/α_W (~ 10 TeV) [8,6] the characteristic energy (or fermion masses, or temperatures) at which the rapid baryon-number non-conservation can take place, are of this order (or even less, see below).

The main purpose of this paper is to discuss some of the anomalous baryon-number non-conserving processes which could occur in the early universe. The obvious motivation is that the baryon-number non-

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conservation (together with the lack of CP non-conservation and thermodynamic equilibrium) can give rise to the baryon asymmetry of the universe (BAU) [11]. One could even hope to calculate both the magnitude and sign of the BAU in a reliable way, since the coupling constants and CP non-conserving phase in the standard electroweak theory are known experimentally. On the other hand, the electroweak baryon-number non-conserving processes can wash out BAU generated at some earlier stage (say, at GUT temperatures [12]). These possibilities are also discussed in this paper

We shall estimate the equilibrium rate of the anomalous fermion-number non-conserving processes in the cosmic plasma in the standard electroweak theory, and find that at sufficiently high temperature this rate exceeds the expansion rate of the universe Then we shall study whether the previously generated BAU is washed out by these processes and discuss the possibility of the BAU generation at the electroweak temperatures in the standard theory with light fermions. We shall find that if the Weinberg–Salam phase transition is of the second order, essentially no BAU is generated by the anomalous electroweak processes. In any case, the previously generated BAU survives only if the corresponding baryon and lepton numbers are not equal to each other In the context of GUTs this means that if $M_H \gtrsim M_W$ (M_H is the electroweak Higgs boson mass), the BAU can be explained only if $(B-L)$ is not conserved at the grand unified scale. There still remains a possibility of the BAU generation during the Weinberg–Salam phase transition, if the latter is of the first order

The situation might be different if there exist heavy elementary fermionic doublets ($M_F > M_W/\alpha_W$). Of course, the introduction of such fermions into the standard theory causes some difficulty, since their Yukawa coupling to the Higgs field is large ($h_\gamma^2 > \alpha_W^{-1} \gg 1$)[‡1]. Therefore, any scenario based on this assumption is extremely speculative. With this reservation, it seems plausible that such fermions rapidly decay due to the electroweak anomaly effects [6], we find it amusing to note in this paper that these decays may be responsible for the BAU generation at the electroweak temperature.

2 For the sake of simplicity, we disregard the interactions due to U(1) weak hypercharge, i e. we consider the limit $\sin^2\theta_W \to 0$, α_W = fixed, which is not far from reality. Then the electroweak lagrangian is

$$\mathcal{L} = (1/2g_W^2)\,\mathrm{Tr}\,F_{\mu\nu}^2 + (D_\mu\varphi)^\dagger\,(D_\mu\varphi) - \lambda(\varphi^\dagger\varphi - v^2/2)^2 + \mathcal{L}_F, \tag{1}$$

where $\mathcal{L}_F$ is the usual fermionic part To be specific, we first discuss the case when the Weinberg–Salam phase transition is of the second order, and consider the temperature of the universe T obeying $M_W < T < T_c$ where T_c is the critical temperature In that case the expectation value of the Higgs field $v(T)$ is non-vanishing (and different from its zero temperature value v).

We use the gauge $A_0 = 0$ in what follows. A subtle point about this gauge is that one has to perform the functional integration over the fields – being twisted boundary conditions at $t = 0$ and $t = T^{-1}$ [14], however, this will not be too essential for our purposes.

To discuss the anomalous fermion-number non-conservation it is convenient to introduce the notion of a number of real fermions in each fermionic doublet, N_f, i.e the number of occupied positive energy fermionic levels minus the number of unoccupied negative energy levels. In the vacuum case, this is just the fermion number, so we are primarily interested in its present value (which is conserved up to the electroweak instanton effects). As the gauge and Higgs fields fluctuate, N_f changes by plus (minus) one whenever a fermionic level crosses zero from below (above). Although this picture of the fermion-number non-conservation was first introduced in the context of the adiabatically varying fields [15], it turns out to be correct in the general case [5]. It is worth mentioning that this picture is valid also in the standard electroweak theory with massive fermions [16,6]. We note in passing that the level crossing occurs simultaneously for each fermionic doublet, so that the total changes in the numbers of real baryons and leptons are $\Delta B = \Delta L = \mathcal{N}$, where $\mathcal{N}$ is the number of fermion generations, i e $B-L$ is conserved (the latter point is clear already from the absence of the electroweak anomaly in $B-L$)

‡1 Another difficulty, which appears already at smaller h_γ is the instability of the Higgs vacuum [13], however, this can be avoided by the introduction of the additional heavy scalar doublet(s) or by the large Higgs self-coupling

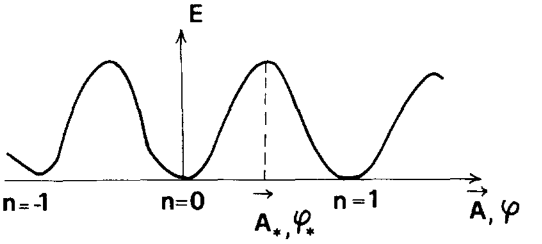

Fig. 1. Schematic dependence of the static energy of the gauge-Higgs system upon the configuration $(\boldsymbol{A}(x), \varphi(x))$ Minima of this curve correspond to the vacuum states with different topological (and baryon) numbers Configuration (A^*, φ^*) is the saddle point of the energy functional.

To estimate the rate of the fermion-number non-conservation (at thermodynamic equilibrium) we have to study the number of level crossings per unit time per unit volume; we do this to the lowest order in $\hbar$, in particular, we neglect the effects of fermions on the fluctuations of the bosonic fields.

We recall the standard picture of the vacuum structure in gauge theories [17] (see fig. 1) The vacua with, say, $n = 1$ and $n = 0$ differ by a topologically non-trivial gauge transformation, and the level crossing occurs at $A = A_*$, $\varphi = \varphi_*$ [15]. Therefore, the rate of level crossings coincides with the rate of the transitions from the regions left to (A_*, φ_*) to that right to (A_*, φ_*). The latter can be evaluated with the use of the theory of the "vacuum" decay at finite temperature [9] which states that this rate is proportional to $\exp(-S)$ where S is the action for the appropriate periodic (up to twist) solution to the euclidean field equations with period T At sufficiently high temperatures (including those under discussion) the dominant contribution comes from static gauge and Higgs fields $\boldsymbol{A}_{\mathrm{cl}}(x)$, $\varphi_{\mathrm{cl}}(x)$, extremizing the free energy F [9]; this extremum should, in fact be a saddle point [9] (cf ref. [18]). In that case the transition rate is proportional to $\exp[-F(\boldsymbol{A}_{\mathrm{cl}}, \varphi_{\mathrm{cl}})/T]$.

At temperatures under discussion, the free energy functional is approximately equal to the static energy with $v(T)$, $g_{\mathrm{W}}(T)$, $\lambda(T)$ substituted for v, g_{W}, λ [19], i.e. $(i, j = 1, 2, 3)$

$$F = \int \mathrm{d}^3x \, \{-[1/2g_{\mathrm{W}}^2(T)] \operatorname{Tr} F_{ij}^2 + (\mathrm{D}_i\varphi)^\dagger (\mathrm{D}_i\varphi) + \lambda(T)[\varphi^\dagger\varphi - v^2(T)/2]^2\}. \tag{2}$$

The saddle point configuration extremizing the right-hand side of eq. (2) has been found in refs. [20,8] in a different context, it has the form (τ_a being Pauli matrices, $r^2 = \boldsymbol{x}^2$)

$$A^i_{\mathrm{cl}} = i(\epsilon_{ijk}\, x_j\, \tau_k/r^2) f(\xi), \tag{3a}$$

$$\varphi_{\mathrm{cl}} = [v(T)/\sqrt{2}]\, h(\xi)\, (i\tau_a\, x_a/r) \binom{0}{1}, \tag{3b}$$

where $\xi = g_{\mathrm{W}}(T)\, v(T) r$ and the functions $f(\xi)$ and $h(\xi)$ have the following asymptotics:

$$f(0) = h(0) = 0, \qquad f(\infty) = h(\infty) = 1. \tag{4}$$

The free energy of this configuration is

$$F = [2M_{\mathrm{W}}(T)/\alpha_{\mathrm{W}}(T)]\; B(\lambda(T)/\alpha_{\mathrm{W}}(T)), \tag{5}$$

where $M_{\mathrm{W}}(T) = \frac{1}{2} g_{\mathrm{W}} v(T)$ and the numerical values of $B(\lambda/\alpha_{\mathrm{W}})$ vary from 1 5 to 2.7 for $\lambda/\alpha_{\mathrm{W}}$ varying from 0 to ∞ [8]. We find that the rate of the fermion-number non-conservation is

$$\frac{\mathrm{d}N_{\mathrm{f}}}{N_{\mathrm{f}}\,\mathrm{d}t} = TC \exp\left(-\frac{2M_{\mathrm{W}}(T)}{T\alpha_{\mathrm{W}}(T)} B(\lambda(T)/\alpha_{\mathrm{W}}(T))\right), \tag{6}$$

where the pre-exponential factor T is introduced on dimensional grounds and C depends on $v(T)/T$ as well as on the coupling constants (the evaluation of C involves the calculation of various determinants in the background fields (3) and will not be attempted here).

This rate should be compared with the expansion rate of the universe, $\tau_{\mathrm{u}}^{-1} \sim N_{\mathrm{eff}}^{1/2}\, T^2/M_{\mathrm{Pl}}$, where N_{eff} is the effective number of degrees of freedom. One immediately obtains that the rate (6) exceeds τ_{u}^{-1} at $T > T^*$, where T^* is approximately given by the following relation.

$$T^* = [2M_{\mathrm{W}}(T^*)/\alpha_{\mathrm{W}} \ln(M_{\mathrm{Pl}}/T^*)]\, B(\lambda/\alpha_{\mathrm{W}}) \tag{7}$$

[we use the fact that α_{W} and λ are slowly varying functions of T, we also consider $\ln(M_{\mathrm{Pl}}/T)$ as a large number and neglect $\ln C$ in comparison with $\ln(M_{\mathrm{Pl}}/T)$]. So, at $T > T^*$ the fermion-number non-conserving processes are at equilibrium, while at $T < T^*$ the fermion-number non-conservation is switched off ‡2 The actual number of T^* depends on the value of λ. For $\lambda \sim g_{\mathrm{W}}^2$ it is typically well below T_{c} and still larger

‡2 At $T \ll T^*$ the dominant contribution to the fermion-number non-conservation rate comes from the time-dependent euclidean configurations of the gauge and Higgs fields (at $T = 0$ these are instantons) However, the rate turns out to be much less than τ_{u}^{-1}, so that in fact our conclusion remains.

than M_W. For instance, at $\lambda = g_W^2$ one finds $B = 2.1$, $T_c \approx 340$ GeV [19] and $T^* \approx 0.6\, T_c \approx 200$ GeV.

There is one point which has been missed in the above discussion. Namely, in the pure Yang–Mills theory the "magnetic" gauge bosons seem to acquire the magnetic mass M_{magn} of the order $\alpha_W T$ [19,14]. [The electric field of the configuration (3) is zero, so we need not discuss the electric mass.] For our results to be valid, the magnetic mass should be much less than $M_W(T)$. At $T = T^*$ this is indeed the case, $M_{\mathrm{magn}}/M_W(T^*) \approx 2B/\ln(M_{\mathrm{Pl}}/T^*) \ll 1$ At higher temperatures, in particular at $T > T_c$, the magnetic mass cannot be neglected However, the weight of the configurations of the form (3a) are believed to be unsuppressed at these temperatures [14], so that the fermion-number non-conserving rate is large, although it cannot be calculated within the semiclassical approach utilized here.

Turning to the possibility of the first order electroweak phase transition, we note that the estimate (6) remains valid for the stage *after* the phase transition. On the other hand, the above discussion implies that before the phase transition, when $\langle\varphi\rangle = 0$, the fermion-number non-conserving processes are rapid even at low temperature (which is possible because of the supercooling), although the corresponding rates cannot be reliably calculated by semiclassical technique.

3 Let us now consider the possible cosmological implications of the anomalous electroweak baryon-number non-conservation in the standard theory with light fermions. We begin with the case of the second-order electroweak phase transitions, in which the anomalous processes are at equilibrium up to $T = T^*$ If the BAU was generated at some earlier stage of the evolution of the universe (say, due to GUT interactions), some part of the BAU is washed out by the anomalous electroweak processes. Since $(B - L)$ is conserved in the latter processes, the evolution of the baryon (and lepton) asymmetry is governed by the following equations (see also ref. [21])

$$\mathrm{d}B/\mathrm{d}t = \mathrm{d}L/\mathrm{d}t = -\tau^{-1}(B + L), \tag{8}$$

where τ is the rate of the anomalous processes. As discussed above, $\tau^{-1} \sim \beta T$ at $T > T_c$, where β is an unknown factor depending on the coupling constants Therefore

$$B(T_c) = \tfrac{1}{2}(B_{\mathrm{in}} - L_{\mathrm{in}}) + \tfrac{1}{2}(B_{\mathrm{in}} + L_{\mathrm{in}})\,\mathrm{e}^{-A},$$

$$A \sim \beta M_{\mathrm{Pl}}/T_c\sqrt{N_{\mathrm{eff}}} \sim \beta \times 10^{15} \tag{9}$$

Clearly, $B(T_c) = \frac{1}{2}(B_{\mathrm{in}} - L_{\mathrm{in}})$ with great precision; this means that if the primordial baryon asymmetry is generated by the $(B - L)$ conserving processes (which is the case in the minimal SU(5) model [22]), it is completely washed out by the moment of the electroweak phase transition.

Can the additional BAU be generated *after* this phase transition? In spite of the fact that the necessary conditions for the BAU generation are satisfied at $T = T^*$, the answer is negative for the following reason. As shown in ref. [23], the most effective BAU generation takes place at the time when the kinetic equilibrium between the relevant particles is violated (and not just at the time when the processes with $\Delta B \neq 0$ come out of the equilibrium) In our case the kinetic equilibrium persists up to $T \sim M_W/\ln(M_{\mathrm{Pl}}/M_W)$, but at this temperature the anomalous electroweak processes are inoperative. An estimate for the BAU generated at $T \sim T^*$ is ($\Delta \equiv n_B/n_\gamma$, n_B and n_γ are baryon and photon number densities respectively)

$$\Delta \sim (\tau_0/\tau_u)\,\delta, \tag{10}$$

where τ_0 is the rate of the reactions establishing the kinetic equilibrium at $T \simeq T^*$, δ is the microscopic asymmetry generated in the reactions with $\Delta B \neq 0$ Since $\tau_0 \sim (\alpha^2 T^*)^{-1}$, one has

$$\Delta \sim (\sqrt{N_{\mathrm{eff}}}\,T^*/\alpha^2 M_{\mathrm{Pl}})\,\delta \sim 10^{-13}\,\delta, \tag{11}$$

which is small compared to the observational value $\Delta_{\mathrm{obs}} \sim 10^{-8}$–$10^{-10}$.

Thus, if the Weinberg–Salam phase transition is of the second order (i e $M_H \gtrsim M_W$ [19]), the anomalous electroweak processes cannot generate the BAU. This conclusion may in principle be verified if it would turn out that $M_H \gtrsim M_W$, then the BAU could be explained only by the $(B - L)$ non-conservation at some large (say, GUT) scale, so that one would predict the non-standard proton decay modes.

Let us briefly discuss the case of the first order electroweak phase transition. As before, the primordial $(B - L)$ is washed out at the $\mathrm{SU}(2)_L$ symmetric stage. However, the new BAU might be generated by the anomalous electroweak processes in the course of the phase transition due to the highly non-equilibrium nature of the latter (formations and collisions of bubbles of the new phase, etc.). If this is indeed the case, then the fate of the new BAU depends on the temperature after reheating, T_{reh} if $T_{reh} > T^*$, the new BAU disappears, while for $T_{reh} < T^*$ (strong supercooling) the BAU survives up to now. The latter possibility, which is the most interesting one, requires further examination; it is worth pointing out that the corresponding values of the Higgs boson mass are close to the Coleman–Weinberg one, $M_H \sim 10.4$ GeV [19].

4. Let us now consider the possibility of the BAU production in the framework of the Weinberg–Salam model with additional doublets of very heavy leptons (although the same effects might take place in processes involving heavy quarks). By the latter we mean the particles with masses of order $M_L \sim M_W/\alpha_W \sim 10$ TeV. Although we cannot evaluate the properties of these particles in a reliable way, it seems plausible that the lightest member of the most heavy doublet can more or less rapidly decay due to the anomalous electroweak effects [6,7] (the heaviest lepton is too short living due to the usual weak decay and so it is not of interest for us), producing $3\mathcal{N}$ quarks and $(\mathcal{N} - 1)$ leptons (the factor 3 accounts for color, as before, $\mathcal{N}$ is the total number of the fermion generations). In fact, these decays can be energetically forbidden, but we assume the masses to be chosen in such a way that this does not happen. One can parametrize the anomalous decay width as follows, $\Gamma = f^2 M_L$. We assume, following ref. [6], that $f \sim \alpha_W$ for M_L exceeding some critical mass $M_c \sim M_W/\alpha_W$, while $f(M_L)$ rapidly falls down to $\exp(-4\pi/\alpha_W)$ for M_L varying from M_c to zero.

Let there exist N_h new heavy fermion generations. If one takes $N_h = 3$ then at temperatures $T \sim M_W$ there would in any case be preserved *CP*-violating phases in the lepton sector. Thus we shall now treat the charge asymmetry in the decays of the heavy leptons in thermodynamic non-equilibrium plasma as a possible origin of the BAU. The value Δ of the BAU is of the order [11,12]

$$\Delta \sim (n_L/N_{eff} n_\gamma)\, \delta_{micro}, \tag{12}$$

where δ_{micro} is the average baryon asymmetry in decays of heavy leptons in vacuum,

$$\delta_{micro} \equiv \frac{\Gamma(L \to \bar{q}\text{'s}) - \Gamma(\bar{L} \to q\text{'s})}{\Gamma(L \to \bar{q}\text{'s}) + \Gamma(\bar{L} \to q\text{'s})}. \tag{13}$$

Let us first consider that the W–S phase transition is the second order one. In this case n_L is to be taken equal to the equilibrium number density of L's at the temperature T_f when the heavy leptons decouple from the cosmological plasma. Let us now estimate δ_{micro} in the model. The diagrams contributing to δ_{micro} are shown schematically in fig. 2, where $\Gamma_L \sim g^2 M_L$ gives the width of the weak decay of the heavy lepton into lighter leptons from other generations due to the generation mixing, g being of the generic form $g = g_W \sin\theta_i\, e^{i\varphi_i}$, where θ_i are the (Cabibbo-like) lepton generation mixing angles, φ_i are the *CP* non-conserving phases

The BAU production in decays of heavy particles is possible only if their masses and decay widths obey the following inequalities [10]:

$$M_L > T_f, \quad M_L > (\Gamma_{tot} M_{Pl}/N_{eff})^{1/2}, \tag{14}$$

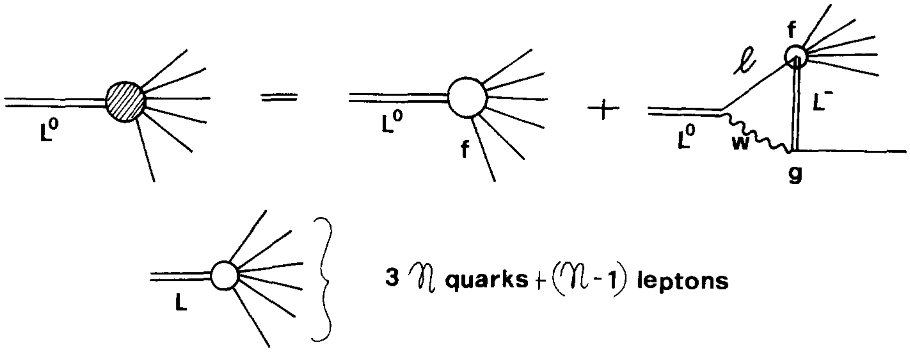

Fig. 2 Diagrams for the anomalous *CP* non-conserving decays of heavy lepton. f is the effective coupling constant for baryon-number non-conserving processes and $g \sim g_W \sin\theta\, e^{i\varphi}$ is the constant for weak decays of heavy lepton into lighter leptons from other generations.

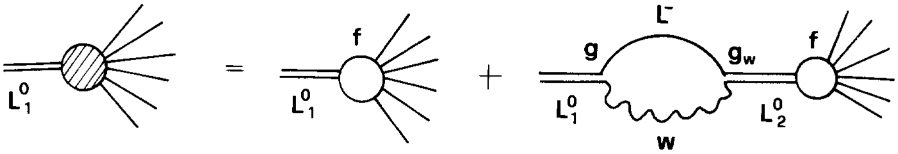

Fig 3. The enhancement of the CP-violating amplitude due to mixing between degenerate heavy leptons from different generations

$$\Gamma_{tot} \sim (f^2 + g^2) M_L, \qquad (14 \text{ cont'd})$$

which at $M_L \sim M_W/\alpha_W$ results in

$$f^2 + g^2 \lesssim 10^{-14}, \quad |f| \lesssim 10^{-7}, \quad \theta \lesssim 10^{-6}. \qquad (15)$$

It is seen from fig. 2 and eq. (15) that

$$\delta_{micro} \sim f^2 g^2/(f^2 + g^2) \lesssim 10^{-14} \qquad (16)$$

One has therefore to conclude that, even without taking into account the macroscopic suppression factor in eq (12), the value of Δ is again too small in comparison with the observational value unless there exists some mechanism of the enhancement of δ_{micro}.

We would like to note that this enhancement could take place if there existed degeneracy in masses between heavy leptons from different generations resulting in the strong mixing and the enhancement of the CP-violating amplitude due to the resonant intermediate state (similar to the K^0–$\bar{K}^0$ system). The corresponding diagrams are shown in fig 3. The value of δ_{micro} is then of order $\delta_{micro} \sim \sin\theta \sin\varphi$, where φ is a CP-violating phase with the maximum CP-violation, $\sin\varphi \sim 1$, the value of δ_{micro} may be as large as $\delta_{micro} \sim 10^{-6}$ The macroscopic factor in eq (12) is in this case of order $N_{eff}^{-1} \sim 10^{-2}$ provided that the conditions (15) are fulfilled, so $\Delta \sim 10^{-8}$.

A more natural possibility of the BAU generation in the anomalous decays of heavy leptons is provided by the strongly first order electroweak phase transition and supercooling. Since there appear large fluctuations of the Higgs field (bubble walls, etc.) during the transition process, it seems reasonable that an appreciable part of the energy of the false vacuum is transformed into the heavy fermions which strongly interact with the Higgs field. So, we estimate the number density of the heavy fermions (leptons) immediately after the phase transition from the energetic reasons,

$$4N_L n_L M_L \sim \epsilon_{vac}, \qquad (17)$$

where M_L is the typical mass of a heavy lepton, N_L is the number of heavy doublets and $\epsilon_{vac} = \kappa M_W^4$ is the energy density of the false vacuum (κ is some function of the coupling constants, in the Coleman–Weinberg theory without heavy fermions $\kappa \sim \alpha_W^2$). The resulting number density much exceeds the equilibrium one, which is of order $e^{-M_L/T}$; the anomalous decays of heavy leptons are automatically out of equilibrium. The microscopic baryon asymmetry is still of the order

$$\delta_{micro} \sim f^2 g_W^2 \sin^2\theta \sin\varphi/(f^2 + g_W^2 \sin^2\theta), \qquad (18)$$

but now the constants need not obey the constraints (14), (15). Therefore for $M_L > M_c$ one has $\delta_{micro} \sim \alpha_W \sin^2\theta \sin\varphi$ and for "natural" choice of the parameters ($\sin\theta \sim 1$, $\sin\varphi \sim 1$) one obtains $\delta_{micro} \sim 10^{-2}$ The macroscopic BAU is

$$\Delta \sim (n_{decay}/T_{reh}^3 N_{eff})\, \delta_{micro}, \quad N_{eff}\, T_{reh}^4 \sim \epsilon_{vac}, \qquad (19)$$

where n_{decay} is the number density of the heavy leptons which decays. In general $n_{decay} \neq n_L$ because the leptons can annihilate (mostly into Higgs scalars which are strongly coupled to the lepton)

However, for a large range of the parameters of the model, the annihilation of the heavy leptons is presumably negligible. Indeed, the annihilation rate is $t_{ann}^{-1} \sim \sigma n_L v$, where σ is the annihilation cross section. A "natural" estimate of σ is provided by the "geometric" cross section, $\sigma \sim M_H^{-2} V^{-1}$. Comparing t_{ann}^{-1} with the decay rate (which was assumed to be of order M_W at $M_L > M_c$) one finds that the annihilation is negligible if

$$M_H^2 \gg \epsilon_{vac}/4N_L M_L M_W \qquad (20)$$

which is not too restrictive. If we neglect the annihilation, i e. set $n_{decay} = n_L$ in eq. (19), we obtain

$$\Delta \sim (\epsilon_{vac}/N_{eff})^{1/4}\, (1/\varphi N_L M_L)\, \delta_{micro} \qquad (21)$$

Taking as an example, $\epsilon_{vac} \sim \alpha_W^2 M_W^4$, $M_4 \sim 10$ TeV, $N_4 = 3$ we get

$$\Delta \sim 10^{-5}\, \delta_{micro}. \qquad (22)$$

Together with eq. (18) this implies that the observed

BAU might indeed be attributed to the anomalous electroweak decays of heavy particles, provided that the electroweak phase transition is the strongly first order one.

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References

[1] G 't Hooft, Phys Rev Lett 37 (1976) 8, Phys Rev D14 (1976) 3432.
[2] C.G. Callan, R.F. Dashen and D.J. Gross, Phys. Lett 63B (1976) 334,
R Jackiw and C. Rebbi, Phys Rev Lett. 37 (1976) 172
[3] N.V. Krasnikov, V.A Matveev, V A Rubakov, A.N. Tavkhelidze and V F Tokarev, Teor. Mat Fiz 45 (1980) 313.
[4] A.A Belavin, A M. Polyakov, A S Schwarz and Yu.S. Tyupkin, Phys Lett. 58B (1975) 85.
[5] N.H Christ, Phys Rev. D21 (1980) 1591
[6] V.A Rubakov, Unsuppressed electroweak baryon-number non-conservation in decays of heavy particles, NORDITA preprint (1984)
[7] J Ambjørn and V.A. Rubakov, Classical versus semi-classical electroweak decay of a techni-skyrmion, NORDITA preprint (1984).
[8] R.F. Klinkhamer and N.S. Manton, Phys. Rev D30 (1984) 2212.
[9] A D Linde, Phys. Lett. 70B (1977) 306; 100B (1981) 37
[10] S Dimopoulos and L Susskind, Phys. Rev D18 (1978) 4500
[11] A D Sakharov, Pis'ma Zh. Eksp Teor. Fiz. 5 (1967) 32,
V A Kuzmin, Pis'ma Zh Eksp Teor Fiz 13 (1970) 335
[12] A Yu Ignatiev, N V. Krasnikov, V.A Kuzmin and A N Tavkhelidze, Proc Intern Conf. Neutrino-77, Vol 2 (Nauka, Moscow, 1978) p. 293; Phys Lett 76B (1978) 436,
M Yoshimura, Phys. Rev. Lett 41 (1978) 281, 42 (1979) 476 (E),
S. Weinberg, Phys Rev Lett. 42 (1979) 850;
A Yu Ignatiev, V A Kuzmin and M E Shaposhnikov, Phys Lett 87B (1979) 114.
[13] N.V Krasnikov, Ya Fiz. 28 (1978) 549
[14] D J. Gross, R D Pisarski and L G Jaffe, Rev. Mod Phys. 53 (1981) 43
[15] C G Callan, R F Dashen and D J Gross, Phys Rev. D17 (1978) 2717,
J Kiskis, Phys. Rev. D18 (1978) 3690,
J. Ambjørn, J Greensite and C Peterson, Nucl. Phys B221 (1983) 381 and references therein
[16] N V Krasnikov, V.A Rubakov and V.F. Tokarev, J. Phys. (NY) 12 (1979) L343
[17] R Jackiw, Rev. Mod Phys. 49 (1977) 681
[18] C G Callan and S. Coleman, Phys. Rev D16 (1977) 1762.
[19] A D. Linde, Rep Prog. Phys. 42 (1979) 389 and references therein; Phys Lett 96B (1980) 289
[20] R F. Dashen, B Hasslacher and A Neveu, Phys Rev D10 (1974) 4138,
J Boguta, Phys Rev Lett 50 (1983) 148,
P. Forgacs and Z Horvath, Phys. Lett 138B (1984) 397.
[21] S.B Treiman and F Wilczek, Phys. Lett 95B (1980) 222,
V A. Kuzmin and M.E Shaposhnikov, Phys. Lett 105B (1981) 163
[22] H Georgi and S L Glashow, Phys Rev Lett 32 (1974) 438
[23] A D Dolgov, Pis'ma Zh Eksp Teor. Fiz 29 (1979) 254;
A.D Dolgov and Ya B. Zeldovich, Usp. Fiz Nauk 130 (1980) 559