Nuclear Physics B287 (1987) 757–775
North-Holland, Amsterdam

BARYON ASYMMETRY OF THE UNIVERSE IN STANDARD ELECTROWEAK THEORY

M.E. SHAPOSHNIKOV

Aspen Center for Physics, Aspen, CO 81610, USA
*and Institute for Nuclear Research of the USSR Academy of Sciences, Moscow 117312, USSR**

Received 17 July 1986

We suggest a new scenario of the baryon asymmetry generation in the framework of standard electroweak theory which is based on the recent observation that anomalous baryon number violating processes are nonsuppressed at high temperatures. This scenario works if the ground state of gauge theories at high temperatures has nontrivial degeneracy with respect to the Chern-Simons number. We find that the correct baryon asymmetry is produced during the electroweak phase transition provided the microscopic asymmetry in B-violating processes $\delta_{ms} \geqslant 10^{-21}$ and the mass of the Higgs boson in near $M_H \simeq 10$ GeV or $M_H \simeq 56$–60 GeV. In that case the sign of baryon asymmetry is entirely determined by the sign of CP violation in K^0 decays. If $\delta_{ms} \leqslant 10^{-21}$ in the standard model then cosmology requires the additional generation of heavy fermions or second doublet of scalars. The constraints on the Higgs boson mass remain unchanged provided fermion and scalar masses are less than M_W.

1. Introduction

The appearance of the baryon asymmetry of the Universe (BAU) is usually related to the grand unified theories (GUT) [1]. The mechanism of the BAU generation in the decays of the heavy leptoquarks at the early stages of the Universe evolution is well elaborated and understood [2]. However, the experimental evidence for GUTs is lacking and therefore the search for other possible explanations of the BAU is desireable (see e.g., ref. [3]). There is also some important point which makes difficult the production of the baryonic excess in the framework of the minimal GUTs like SU(5) or SO(10). We mean the recent observation [4] that the anomalous electroweak baryon number nonconserving processes are nonsuppressed at high temperatures (at zero temperatures the suppression factor $\simeq \exp(-4\pi/\alpha_W)$ [5]). The latter fact implies the washing out of the primordial BAU with $B=L$ (B and L are the baryon and lepton numbers correspondingly) by the equilibrium anomalous processes with $\Delta B \neq 0$. Therefore, if the BAU is produced by the $B-L$ conserving interactions (as in the case of SU(5) or some versions of SO(10)) then it

* Permanent address.

0550-3213/87/$03.50 © Elsevier Science Publishers B.V.
(North-Holland Physics Publishing Division)

entirely disappears [4] to the moment of the SU(2) × U(1) phase transition. For that reason the GUTs with $B-L$ conservation fail in the explanation of the BAU. Although the construction of the theories with $B-L$ nonconservation and creation of nonzero $B-L$ excess may be carried out [6] these models seem to be rather artificial.

In the present paper we suggest a new mechanism of the BAU generation in the standard electroweak theory. It is based on the existence of rapid anomalous processes with B nonconservation [4, 6] and on the possible nontrivial degeneracy of the ground state of gauge theories at high temperatures with respect to the Chern-Simons number.

The main idea is as follows. At high temperatures the processes with baryon number nonconservation are almost (*not strictly*) in thermal equilibrium. Due to the B nonconservation, CP violation (which is described in this case by the CP violating phase, δ, in the Kobayashi-Maskawa matrix) and slight deviations from thermal equilibrium, a small amount of baryons is produced [4, 6]:

$$\frac{n_{\mathrm{B}}}{n_{\gamma}} \sim \frac{\mu_{\mathrm{B}}}{T} \sim \frac{\tau_0}{\tau_{\mathrm{u}}} \delta_{\mathrm{ms}}, \tag{1}$$

where δ_{ms} is the microscopic asymmetry in the processes with $B \neq 0$

$$\delta_{\mathrm{ms}} = \left[\sigma(\mathrm{in} \to \mathrm{out}) - \sigma(\overline{\mathrm{in}} \to \overline{\mathrm{out}})\right] / \left[\sigma(\mathrm{in} \to \mathrm{out}) + \sigma(\overline{\mathrm{in}} \to \overline{\mathrm{out}})\right], \tag{2}$$

n_{B} and n_{γ} are the densities of baryons and photons, μ_{B} is the baryonic chemical potential, τ_0 is the characteristic time of B-violating processes, τ_{u} is the age of the Universe. Note that the sign of the baryon asymmetry is connected with the sign of δ and therefore with the sign of CP violation in K^0 decays. In the presence of the nonzero fermionic density the effective action for the static gauge fields is modified [7, 8]

$$\Delta E \sim -\mu_{\mathrm{B}} N_{\mathrm{CS}}, \tag{3}$$

N_{CS} being the Chern-Simons term,

$$N_{\mathrm{CS}}(A) = -\frac{1}{16\pi^2} \mathrm{Tr}\left(F_{ij} A_k - \tfrac{2}{3} A_i A_j A_k\right) \varepsilon_{ijk}.$$

The visible gauge noninvariance of the effective action (3) have been explained in ref. [8]. Suppose that the gauge field slowly varies from $A=0$ at $t \to -\infty$ to some value A at $t \to \infty$. The difference between the free energies of the system at $t \to +\infty$ and $t \to -\infty$ coincides with the eq. (3). This ensures the gauge invariance of the action but makes it dependent on the path from $A=0$ to $A \neq 0$.

Consider now the effective potential for the Chern-Simons density

$$U(N_{CS}) = \Gamma(J) + JN_{CS},$$

$$\exp(\Gamma(J)) = \int DA \exp(-S - JN_{CS}),$$

$$\frac{d\Gamma}{dJ} = -N_{CS}. \tag{4}$$

Here S is the effective action for 3-dimensional pure gauge theory (which is a high temperature limit of 4-dimensional one), $\Gamma(J)$ is the generating functional for Green functions containing Chern-Simons density. The potential $U(N_{CS})$ cannot be calculated by the perturbative technique due to the infrared divergences at $T \neq 0$. The possible forms of the effective potential are pictured on fig. 1. In the first case (see fig. 1a) $U(N_{CS})$ has nontrivial degeneracy with respect to the Chern-Simons density. (Note that $U(N_{CS})$ is the convex function of N_{CS}.) This degeneracy definitely exists if the perturbative ground state at $T \neq 0$ is unstable against the spontaneous creation of the gauge magnetic field [9]. It is not excluded that the same situation takes place if at large distances $\gtrsim (g_W^2 T)^{-1}$ the electroweak plasma exhibits confining properties (see, e.g. [10]).

It is clear that the existence of the contribution [3] breaks the degeneracy. Therefore, if one starts from some point with, say, $N_{CS} = 0$ then after *sufficient* time the system finds itself in the state with the maximal possible N_{CS}. On dimensional grounds (see sect. 2) $N_{CS}^{max} \simeq (\alpha_W/\pi) g_W^4 T^3$. After the electroweak phase transition, gauge bosons acquire masses and the nontrivial ground state becomes unstable. Due to the electroweak anomaly the number of baryons produced coincides with the Chern-Simons density [11] (see the discussion of the level crossing phenomena in

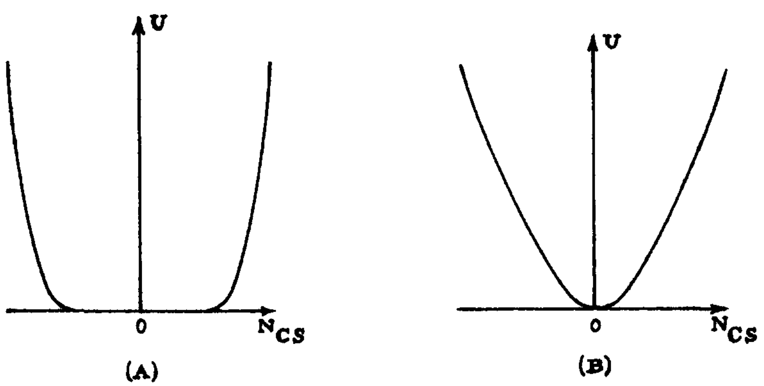

Fig. 1. Effective potential for the Chern-Simons density.

[12]). The fate of this asymmetry depends to a great extent on the details of the electroweak phase transition. In particular one should demand that after the phase transition the baryon number nonconserving processes move out of equilibrium. This gives the upper bound on the Higgs boson mass: $M_H \leqslant 60$ GeV. On the other hand the electroweak phase transition takes place at some (probably low) temperature. This gives the lower bound on the Higgs boson mass $M_H \geqslant 10$ GeV (if there are no very heavy fermions). Generally the baryon asymmetry produced by the ground state decay appears to be larger than the observed one. The sufficient dilution of the BAU takes place very near to the two quoted values of M_H: 10 GeV and 60 GeV. This can be considered as an almost exact prediction of the Higgs boson mass from cosmology.

In the second case (see fig. 1b) the ground state has zero Chern-Simons density, and the influence of the small fermionic density on it is negligible. No baryon asymmetry is produced during the electroweak phase transition. We believe that this possibility is not realized.

The paper is organized as follows: In sect. 2, we find the baryonic chemical potential arising due to CP nonconservation in processes with $\Delta B \neq 0$ and discuss the effective lagrangian for the gauge fields. Sect. 3 is devoted to the consideration of the Chern-Simons density evolution in the expanding Universe. In sect. 4 we discuss the electroweak phase transition and obtain the constraints on the Higgs boson mass. In sect. 5 we consider the magnitude of the relevant CP violation and estimate the BAU in the standard model. The possible modifications of the particle content of $SU(2) \times U(1)$ are also discussed. The last section is devoted to a discussion of the results.

2. Effective lagrangian for the gauge fields

The rate of the baryon nonconserving process at temperatures below the critical point has been estimated in ref. [4]:

$$V_B(t) = \beta_0 T \exp(-E_0/T), \qquad E_0 = \frac{2M_W(T)}{\alpha_W(T)} B\left(\frac{A}{\alpha_W}\right). \tag{5}$$

Here E_0 is the height of the barrier separating vacua with different baryon numbers, A is the Higgs self-coupling constant. Function B was numerically calculated in ref. [13]. It monotonically increased from 1.5 to 2.7 for A varying from 0 to ∞.

The evaluation of the pre-exponential factor β_0 involves the calculation of various determinants in the background of the solution of the static classical equations of motion which corresponds to the top of the barrier. At $T > T_c$ (T_c being critical temperature of the electroweak phase transition) the rate V_B cannot be calculated by the semiclassical technique but it is expected that $V_B(t)$ is large due to the absence of the exponential suppression [4]. In what follows we will assume that for $T > T_c$,

$V_B(t) = \beta T$ with β being some numerical coefficient depending on the gauge and scalar coupling constants.

The kinetic equation for the evolution of the baryonic chemical potential reads (see, e.g. [14])

$$\frac{\mathrm{d}}{\mathrm{d}t}\frac{\mu_B}{T} = -\frac{\mu_B}{T} V_B(t) + \delta_{ms} V_B(t)\big(f(t) - f_{eq}(t)\big), \tag{6}$$

where δ_{ms} is the microscopic asymmetry, defined in (2), $f(t)$ is the relative concentration of fermions, $f(t) = n_F/T^3$, f_{eq} is the equilibrium relative concentration.

The time dependence of the temperature is described by the usual formula $t = M_0/T^2 (M_0 = M_{Pl}/1.66\ N_{eff}^{1/2}$, N_{eff} being the effective number of massless degrees of freedom). To find the time evolution of μ_B one should consider the kinetic equation for $f(t)$:

$$\frac{\mathrm{d}f}{\mathrm{d}t} = -V_0(t)\big(f - f_{eq}\big), \tag{7}$$

where $V_0(t)$ is the rate of the processes establishing the kinetic equilibrium between the fermions. The main contributions to $V_0(t)$ give reactions like $qq \to qq, q\bar{q} \to WW$, etc. By the order of magnitude

$$V_0(t) \simeq \sigma n_\gamma \simeq \frac{\alpha_s^2 T}{\pi}, \tag{8}$$

where σ is the cross section of the above-mentioned processes, α_s being the strong gauge coupling constant. The solution of eq. (7),

$$f - f_{eq} \equiv \Delta f(t) = C \exp\left(-\int_{t_0}^{t} V_0(t')\,\mathrm{d}t'\right) + \int_{t_0}^{t} \frac{\mathrm{d}f_{eq}}{\mathrm{d}t'}\,\mathrm{d}t' \exp\left(-\int_{t'}^{t} V_0(t'')\,\mathrm{d}t''\right) \tag{9}$$

(C being the integration constant) depends on $\mathrm{d}f_{eq}/\mathrm{d}t$ which is of the order of

$$\frac{\mathrm{d}f_{eq}}{\mathrm{d}t} \simeq -\frac{\mathrm{d}}{\mathrm{d}t}\frac{M_f^2}{T^2}. \tag{10}$$

At high temperatures the effective fermion mass is nonzero [9], $M_f^2 \sim g_s^2 T^2$ (the contribution of the strong interactions is the largest one). This gives $\mathrm{d}f_{eq}/\mathrm{d}t \sim bT^2\alpha_s^2/M_0$, where b is the coefficient in the expression for the SU(3) β function,

$\mu^2\partial\alpha_s/\partial\mu^2 = b\alpha_s^2/\pi$, $b=\frac{1}{4}(11-\frac{2}{3}N_f)$, N_f is the number of quark flavours. For $t \geqslant V_0(t)^{-1}$ the first term in (9) gives no contribution to $\Delta f(t)$, while the second one is of the order

$$\Delta f(t) \simeq \pi b T/M_0 . \tag{11}$$

Let us emphasize that deviations from thermal equilibrium may arise only in scale noninvariant theories [2, 14, 15]. In our case scale invariance is broken by the varying of the gauge coupling with temperature.

Using (11) we may find the solution of eq. (6). For the time $t \geqslant (\beta T)^{-1}$ is has the form

$$\mu_B(t) = \pi b T^2 \delta_{ms}/M_0 . \tag{12}$$

The nontrivial contribution of the fermions to the effective action for the static magnetic fields is well known [7, 8]:

$$\Delta E = -2\mu_B N_f N_{CS}(A) - \mu_B\Big(\sum_{\text{doublets}} Y_L^2 - \tfrac{1}{2}\sum_{\text{singlets}} Y_R^2\Big) N_{CS}(B) . \tag{13}$$

Here Y_L, Y_R are hypercharges of quarks and leptons, $N_{CS}(B)$ is the Chern-Simons density of the abelian U(1) gauge field.

Let us also outline another way for obtaining contributions proportional to N_{CS}. Consider the two-point Green function of the gauge fields

$$G_{ij}(k) = \int e^{ikx} d^3x \, \mathrm{Tr}\{A_i(x)A_j(0)\rho(t)\} , \tag{14}$$

$\rho(t)$ being the statistical operator of the system. In thermodynamical equilibrium $\rho(t) = (1/Z)\exp(-H/T)$. *CPT* invariance implies $G_{ij}(k) = G_{ji}(k)$, that is, the Chern-Simons term (which is asymmetric on ij) is forbidden. In the *CP*-invariant theory N_{CS} is absent even in the nonstationary case (if $\rho(t)$ commutes with *CP* transformation) because *CP* invariance requires $G_{ij}(k) = G_{ij}(-k)$. If both *CP* and *CPT* symmetries are absent (Universe evolution breaks *CPT*) the N_{CS} term arises naturally. The contribution of some specific diagram has the form (see fig. 2)

$$\left(\left(\delta_{ij} - \frac{k_i k_j}{k^2}\right)\Pi_1(k^2) + \varepsilon_{ijk} k_k \Pi_2(k^2)\right)\prod h_k \tag{15}$$

owing to the transversality of G_{ij} on k and the isotropy of the space. Πh_k represents the product of the relevant Yukawa coupling constants. $\Pi_2(k^2)$ is a nonzero and is expected to be proportional to the Hubble constant $H = \dot{a}/a = T^2/2M_0$ which is the natural measure of the *CPT* breaking due to the Universe expansion.

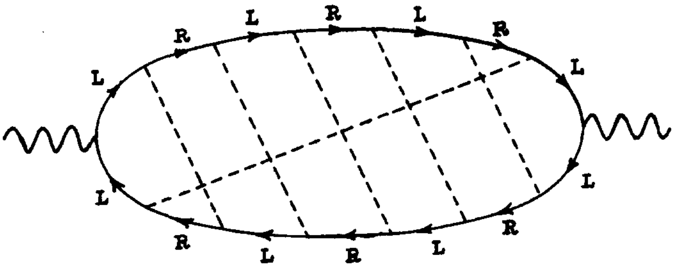

Fig. 2. Lowest order CP violating diagram contributing to the Chern-Simons density.

The contribution of the CP conjugated diagram can be read off from (15) with the substitution $k \to -k$, $h_k \to h_k^*$. The sum of these two diagrams has nontrivial contribution proportional to the Chern-Simons density

$$G_{ij}^{(\mathrm{CS})} \simeq \varepsilon_{ijk} k_k \Pi_2(k^2) \mathrm{Im} \prod h_k . \tag{16}$$

Note that the amplitude of CP violation in the B nonconserving processes is also proportional to the imaginary part of the product of the Yukawa couplings. That means that the contribution (16) has essentially the same order of magnitude as the "one-loop" term proportional to μ_{B}.

We will discuss the magnitude of δ_{ms} in the electroweak theory in sect. 4.

3. Evolution of the Chern-Simons density in the expanding Universe

In this section we will consider the influence of the nonzero fermionic density on the structure of the ground state of gauge theories at high temperatures.

It is well known that the perturbative calculations at $T \neq 0$ break down at small momentum $p \leqslant g_{\mathrm{W}}^2 T$ [9,16]. This fact, for example, makes impossible the calculation of $\mathrm{O}(g_{\mathrm{W}}^6)$ contributions to the thermodynamical potential. One of the possible solutions of the infrared problem in the thermodynamics of Yang-Mills fields consists of the dynamical mass generation for the static magnetic fields, $m_{\mathrm{magn}} \sim g_{\mathrm{W}}^2 T$ [9,16]. On the other hand the power infrared divergences may signal the incorrectness of our choice of the ground state [9], which is in fact unstable against the creation of the inhomogeneous magnetic field. The latter possibility seems to be quite natural from the point of view of the perturbative calculations of the infrared behavior of the nonabelian gauge self-energy at $T \neq 0$ [9]:

$$G_{ij}(\omega = 0, \boldsymbol{p} \to 0) = \left(\delta_{ij} - \frac{p_i p_j}{p^2}\right)\left(p^2 - \tfrac{1}{64} g^2 N |p| T(11 + 2\alpha + \alpha^2)\right), \tag{17}$$

where α is a gauge fixing parameter, $N = 2$ in our case of the SU(2) group. The physical interpretation of the result (17) is not transparent due to its gauge dependence. However, the minus sign in front of the term linear in p does not depend on the gauge and may give rise to the unstability of the state with $A=0$.

The magnitude of the gauge field in the ground state can be easily estimated by power counting (see the analogous consideration in [9, 16]). The high temperature limit of the four-dimensional theory coincides with three-dimensional theory and with the unique dimensionful gauge coupling constant $\tilde{g}=g_{\mathrm{W}}\sqrt{T}$, $\tilde{A}_i=A_iT^{-1/2}$. Thus $\tilde{A}_\mu \sim g_{\mathrm{W}}\sqrt{T}$, $p \sim g_{\mathrm{W}}^2 T$, i.e. $A_i \sim g_{\mathrm{W}} T$, $H_i \sim g_{\mathrm{W}}^3 T$ (H is 4-dimensional magnetic field). The contribution of these fields to the free energy density is of the order of $g_{\mathrm{W}}^6 T^4$, coinciding with the first noncalculable by the perturbation theory term.

It is clear that the free energy of this configuration does not depend on the Chern-Simons density. This means that in thermodynamical equilibrium the ground state has nontrivial degeneracy with respect to the changing of N_{CS}:

$$N_{\mathrm{CS}}=N_{\mathrm{CS}}(0)\cos\varphi, \qquad N_{\mathrm{CS}}(0) \simeq \frac{\alpha_{\mathrm{W}}}{\pi} g_{\mathrm{W}}^4 T^3, \qquad 0<\varphi \leqslant 2\pi. \tag{18}$$

Consider for example the abelian configuration of the form [8]

$$\begin{aligned} A_1^{(3)} &= a\sin(kx_3+\varphi), \\ A_2^{(3)} &= a\sin kx_3, \\ A_i^{(1)} &= A_i^{(2)}=0. \end{aligned} \tag{19}$$

It has $\langle H_1^2+H_2^2\rangle = k^2a^2$ and $N_{\mathrm{CS}}=a^2k\cos\varphi$ i.e. the energy does not depend on the phase shift φ but N_{CS} does.

There is one point where the discussion of the spontaneous magnetization of the ground state could be incorrect. The effective 3-dimensional theory which governs the dynamics of the magnetic fields is believed to be confining [16, 10]. This probably forbids the existence of long-ranged magnetic fields with size $\gtrsim 1/g_{\mathrm{W}}^2 T$. However, it is not excluded that the ground state in this case has the same degeneracy.

In what follows we shall consider only the ground state containing nonzero magnetic field. Because all of our estimates are based on power counting we expect that the results do not depend on the exact structure of the ground state at high temperatures. The only thing that is really needed is the degeneracy of the ground state with respect to the Chern-Simons number.

The existence of thermodynamical nonequilibrium breaks this degeneracy and the configuration with the maximal possible N_{CS} becomes preferable (see eq. (13)). If the system has enough time then it finds itself finally in the state with $N_{\mathrm{CS}}=N_{\mathrm{CS}}(0)$. After the electroweak phase transition W and Z bosons become massive and the ground state with nonzero weak magnetic field becomes unstable and decays, producing $N_B = \frac{1}{2}N_f N_{CS}$ baryons due to the electroweak anomaly in the fermionic current*. The sign of BAU produced by this process is defined by the CP violation in K^0 decays through the CP phase δ in the Kobayashi-Maskawa matrix. The magnitude of the BAU depends on the details of the electroweak transition and will be considered in sects. 4, 5.

Consider now the evolution of the Chern-Simons density of the ground state in the expanding Universe. The change of the phase φ leads to the appearance of a weak electric field proportional to $\dot{\varphi}$. This gives rise to the electric energy density:

$$E_e = \tfrac{1}{2}\sqrt{-g}\, g_{ik} A_a^i(T) A_a^k(T) \dot{\varphi}^2, \tag{20}$$

$A_a^i(T)$ being the amplitude of the gauge field and g_{ik} are the spatial components of the Freedmann-Robertson-Walker metric, $ds^2 = (dt^2 - a^2(t)dx^2)$. Taking into account the Chern-Simons contribution to the free energy (12), (13), (18), the following equation for the phase φ can be obtained:

$$\ddot{\varphi} + 3H\dot{\varphi} + \alpha_W b g_W^2 \frac{T^3}{M_0} \delta_{MS} \sin\varphi = 0, \tag{21}$$

where H is the Hubble constant $H = \dot{a}/a = 1/2t$ in the radiation dominated Universe. The solution of the linearized version of (21) with the initial conditions $\varphi(0) = \varphi_0$, $\dot{\varphi}(0) = 0$ has the form

$$\varphi = 2\varphi_0 \frac{1}{z} J_1(z), \qquad z = 4\left(\frac{M_0}{T} b g_W^2 \alpha_W \delta_{ms}\right)^{1/2}. \tag{22}$$

Here $J_1(z)$ is the Bessel function of the first kind. For $z \gg 1$, φ oscillates near the point $\varphi = 0$ with the decreasing amplitude:

$$\varphi(z) = 2\varphi_0 \left(\frac{2}{\pi z}\right)^{1/2} \cdot \frac{1}{z} \cos\left(z - \tfrac{3}{4}\pi\right). \tag{23}$$

So, if the temperature of electroweak phase transition is less than $M_0 b g_W^2 \alpha_W \delta_{ms}$ then the ground state has the maximal Chern-Simons density. For $z \ll 1$

$$\varphi(z) = \varphi_0\left(1 - \tfrac{1}{8}z^2\right), \tag{24}$$

* In the absence of CP violation the number of baryons produced equals zero due to the equipartition distribution of the N_{CS} density

$$n_B = \frac{1}{2\pi}\int_0^{2\pi} N_{CS}(\varphi)\, d\varphi = 0.$$

that is the average N_{CS} is of the order of $z^2/8N_{CS}(0)$*. The number density of baryons generated during the decay of the ground state at the moment T_c equals

$$n_B \simeq \tfrac{1}{2} N_f N_{CS}(\varphi(T_c)). \tag{25}$$

Up to now we have discussed the contribution of the nonabelian fields to the baryon asymmetry. What about the abelian field? The main difference between abelian and nonabelian fields is that the state with zero abelian magnetic field is definitely stable at high temperatures and $\mu_B = 0$ [17]. Of course, the nonzero contribution of the fermionic density gives rise to instability, but only modes with very small momentum $k \simeq \delta_{ms} T^2/M_0$ grow. Because the rate of N_{CS} changing is in turn proportional to k the number of baryons generated by the decay of the abelian configuration is of the order of δ_{ms}^2. This result can be easily checked by the analysis of the corresponding equation of motion.

4. Electroweak phase transition and the bounds on the Higgs boson mass

The baryonic excess generated during the electroweak phase transition must survive till the present time. The necessary condition for that is the freezing out of the baryon number nonconserving processes in the phase with broken SU(2) × U(1) at the critical point T_c. Thus, the phase transition should be of the first order: in the opposite case the vacuum expectation value of the scalar field gradually changed from zero to $\sigma \simeq 250$ GeV and one expects the washing out of any baryonic number [4, 6]. The order of the electroweak phase transition depends on the Higgs boson mass M_H [18]. To find the bounds on the M_H we have to consider only the case $g_W^4/64\pi^2 \ll A \ll g_W^2$ where A is the Higgs self-coupling constant. If $A \simeq g_W^2$ then the electroweak phase transition is of the second order [18] and final BAU is exponentially close to zero. For the small $A \sim g_W^4/64\pi^2$ phase transition (PT) proceeds with the strong supercooling and W-boson mass after PT is near their zero temperature value. This ensures the exponential suppression of the baryon nonconserving processes, so the number of baryons does not change.

The one-loop effective potential for the Higgs field φ at zero temperatures has the form [19]

$$V(\varphi)_{T=0} = -(2A + B)\sigma^2\varphi^2 + A\varphi^4 + B\varphi^4 \ln\frac{\varphi^2}{\sigma^2}, \tag{26}$$

where $\varphi = \sigma \simeq 248$ GeV corresponds to minimum $V(\varphi)$,

$$B = \tfrac{3}{64}\alpha_W^2\left(2 + \frac{1}{\cos^4\theta_W}\right). \tag{27}$$

* Note that the rate of the changing of $N_{CS}\dot{\varphi} = \delta_{ms} T$ is small compared with the rate of baryon number nonconserving processes $V_B \simeq \beta T$. This fact justifies the neglection of influence of the ground state on the value of fermionic chemical potential μ_B (see eq. (6)).

The Higgs boson mass is given by

$$M_{\mathrm{H}}^2 = 4\sigma^2(3B + 2A). \tag{28}$$

We will suppose that there are no heavy fermions and $m_t \lesssim M_{\mathrm{W}}$. For $A \ll g^2$ the main contribution to the finite temperature effective potential comes from the gauge loops, the result of the standard calculation is [20, 18]

$$\Delta V_T = \tfrac{1}{32} T^2 g_{\mathrm{W}}^2 \varphi^2 \left(2 + \frac{1}{\cos^2\theta_{\mathrm{W}}}\right) - \frac{T}{32\pi} g_{\mathrm{W}}^3 \varphi^3 \left(2 + \frac{1}{\cos^3\theta_{\mathrm{W}}}\right)$$

$$+ \tfrac{3}{64}\alpha_{\mathrm{W}}^2 \varphi^4 \left\{ 2\ln\frac{T^2\xi}{\frac{1}{4}g_{\mathrm{W}}^2\varphi^2} + \frac{1}{\cos^4\theta_{\mathrm{W}}} \ln\frac{T^2\xi\cos^2\theta_{\mathrm{W}}}{\frac{1}{4}g_{\mathrm{W}}^2\varphi^2} \right\},$$

$$\ln\xi = 2\ln 4\pi - 2\gamma = 3.91,$$

$$V_T = V_{T=0} + \Delta V_T. \tag{29}$$

The moment of the Weinberg-Salam phase transition practically coincides with the moment of the absolute instability of the phase with $\varphi = 0$ [21] because the tunneling transitions are strongly suppressed. The critical temperature is

$$T_{\mathrm{c}}^2 = \frac{(B + 2A)\sigma^2}{\frac{1}{32}g_{\mathrm{W}}^2\left(2 + 1/\cos^2\theta_{\mathrm{W}}\right)}. \tag{30}$$

After the rapid phase transition condensate of the field, φ is nonzero and equals

$$\varphi = \frac{3T_{\mathrm{c}}}{128\pi A_{\mathrm{eff}}} g_{\mathrm{W}}^3 \left(2 + \frac{1}{\cos^3\theta_{\mathrm{W}}}\right),$$

$$A_{\mathrm{eff}} = A + \tfrac{3}{64}\alpha_{\mathrm{W}}^2 \left(2\ln\frac{T^2\xi}{M_{\mathrm{W}}^2} + \frac{1}{\cos^4\theta_{\mathrm{W}}}\ln\frac{T^2\xi}{M_{\mathrm{z}}^2}\right). \tag{31}$$

The kinetic equation describing the washing out of the baryon asymmetry reads (see eq. (5)):

$$\frac{\mathrm{d}}{\mathrm{d}t}\left(\frac{n_{\mathrm{B}}}{n_\gamma}\right) = -\left(\frac{n_{\mathrm{B}}}{n_\gamma}\right)\beta_0 T \exp\left\{-\zeta\frac{T_{\mathrm{c}}}{T}\right\}, \qquad \zeta = \frac{3\pi\alpha_{\mathrm{W}}}{8A_{\mathrm{eff}}}\left(2 + \frac{1}{\cos^3\theta_{\mathrm{W}}}\right) B\left(\frac{A}{\alpha_{\mathrm{W}}}\right). \tag{32}$$

We substitute $M_{\mathrm{W}}(T)$ by $M_{\mathrm{W}}(T_{\mathrm{c}})$ because the most effective dilution of the baryon number is expected to take place near the critical temperature. The function $B(A/\alpha_{\mathrm{W}})$ equals 1.5 for the values of A considered.

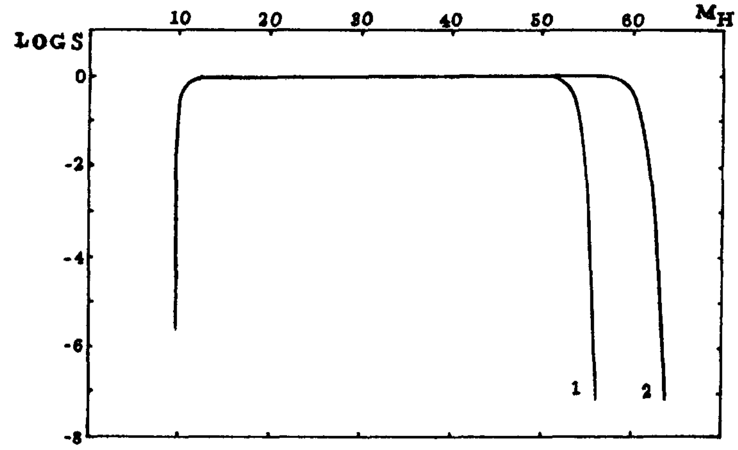

Fig. 3. The dependence of the baryon asymmetry suppression factor $S(M_{\rm H})$ on $M_{\rm H}$. Curve 1 corresponds to $\beta_0 = 1$ and curve 2 to $\beta_0 = 10^{-3}$.

The solution of (31) is given by

$$\left(\frac{n_{\rm B}}{n_\gamma}\right)\bigg|_{t\to\infty} = \left(\frac{n_{\rm B}}{n_\gamma}\right)_{T_{\rm c}} S(M_{\rm H}), \qquad S(M_{\rm H}) = \exp\left(-\beta_0 \frac{2M_0}{T_{\rm c}}\frac{1}{\zeta}\exp(-\zeta)\right). \quad (33)$$

The dependence of the suppression factor S on $M_{\rm H}$ is presented on fig. 3 for $\beta_0 = 1$ and $\beta_0 = 10^{-3}$. Note the very sharp changes of S near $M_{\rm H} = M_{\rm crit} \simeq 56$–$60$ GeV which come from the "twice" exponential behaviour of S. From that we obtain the upper bound on the Higgs boson mass $M_{\rm H} \simeq 60$ GeV slightly depending on the magnitude of β. The maximal admissible temperature of the PT is near 120 GeV.

Now let us turn to the light Higgs boson. If $M_{\rm H}$ is equal to the Coleman-Weinberg value $M_{\rm CW} \simeq 10$ GeV [22] then the electroweak PT occurs at rather low temperatures. In this case PT begins at the moment of chiral symmetry breaking by strong interactions [23]. The baryon asymmetry decreases by the factor $S_{\rm in}/S_{\rm out} \sim 10^{-5} \div 10^{-6}$ [23, 24] where $S_{\rm in}$ and $S_{\rm out}$ are the entropy densities before and after PT.

It is not difficult to find this suppression factor for the $M_{\rm H}$ close to $M_{\rm CW}$. One must just equate the energy densities before and after the PT defining in that way the reheating temperature $T_{\rm r}$. Then the calculation of $S_{\rm in}$ and $S_{\rm out}$ is straightforward. The result is

$$\frac{S_{\rm in}}{S_{\rm out}} = \frac{N_{\rm in} T_{\rm c}^3}{N_{\rm out} T_{\rm r}^3} = \left(\frac{N_{\rm in}}{N_{\rm out}}\right)^{1/4} \left(1 + \frac{\left(2 + 1/\cos^4\theta_{\rm W}\right)^2}{2 + 1/\cos^2\theta_{\rm W}} \frac{5}{2N_{\rm in}} \frac{x - \frac{1}{2}}{(x-1)^2}\right)^{-3/4}, \quad (34)$$

$$x = M_{\rm H}^2/M_{\rm CW}^2.$$

Here $N_{\text{in}} = 90.75$ and $N_{\text{out}} = 71.25$ are the effective number of massless helicity states of the particles before and after the PT, respectively. Expression (34) is correct for $T_c > T_{\text{chiral}} \sim 100$ MeV. It is seen that $S_{\text{in}}/S_{\text{out}}$ is very close to 1, starting practically from the Coleman-Weinberg value of the Higgs boson mass (see fig. 3). The existence of the plateaux from $M_H \simeq M_{CW}$ to $M_H \simeq M_{\text{crit}}$ in the total suppression factor $S(M_H)$ taking into account both the entropy generation and partial equilibrium of $\Delta B \neq 0$ processes makes it possible to give the almost unambiguous prediction of the Higgs boson mass. We will discuss this point in the next section.

5. Baryon asymmetry of the Universe in the electroweak theory

Now we are ready to write down the final expression for the baryon asymmetry of the Universe:

$$\Delta = \frac{n_B}{n_\gamma} \simeq \frac{45}{4\pi^2 N_{\text{in}}} N_f N_{CS}(\varphi(T_c)) S(M_H). \tag{35}$$

The last quantity which should be found is the magnitude of the CP violation in $\Delta B \neq 0$ processes. CP breaking in the electroweak theory at high temperatures arises from the interaction of quarks with Higgs field through the Yukawa couplings. The relevant part of the lagrangian is

$$L_Y = \frac{g_W}{\sqrt{2} M_W} \{ \bar{Q}_L K M_d D_R \varphi + \bar{Q}_L M_u U_R \tilde{\varphi} + \text{h.c.} \}, \tag{36}$$

where Q_L^α are the left-handed quark doublets (α is the index of generation), $U_R^{(\alpha)}$ and $D_R^{(\alpha)}$ are the right-handed quarks with electric charges $\frac{2}{3}$ and $-\frac{1}{3}$ correspondingly, K is the Kobayashi-Maskawa matrix, M_u and M_d are diagonal matrices of the fermions.

It is easy to see that the imaginary part of the product of the Yukawa coupling constants arises only in the 12th order of perturbation theory. The lowest contribution is:

$$D \equiv \operatorname{Im} \prod h_K = \operatorname{Im} \operatorname{Tr} \mathcal{M}_u^2 \mathcal{M}_d^2 \mathcal{M}_u \mathcal{M}_d = \left(\frac{g_W^2}{2M_W^2} \right)^6 s_1^2 s_2 s_3 \sin\delta \, m_t^4 m_b^4 m_c^2 m_s^2,$$

$$\mathcal{M}_d \equiv K M_d^2 K^+, \qquad \mathcal{M}_u = M_u^2. \tag{37}$$

Here $s_i = \sin\theta_i$, θ_1 is the Cabibbo angle, δ is the CP violating phase. The most difficult task is the calculation of the 7-loop diagrams contributing to the Chern-Simons term (see fig. 2). The number of different diagrams is $N_d = 6! \cdot 5 \cdot 6/2 =$

10 800 (the factor 6! is the number of the different scalar field pairing, $5 \cdot 6/2$ is the number of positions of the gauge vertices). Because of this tremendous number and very high order of the diagrams we do not try to calculate them here. For the same reasons we do not attempt here to find the sign of $\delta_{\rm ms}$ and therefore we do not know the sign of BAU, though this problem in principle can be solved. The best that we can do is to get more or less reliable estimates of $\delta_{\rm ms}$. At zero temperatures the usual "$1/4\pi^2$ for a loop" counting gives

$$\delta_{\rm ms} \simeq DN_{\rm d}\left(1/4\pi^2\right)^6 \simeq 2\times 10^{-27}. \tag{38}$$

We use the upper limit on the product $s_1^2 s_2 s_3 \leqslant 3\times 10^{-4}$ [25] and $\sin\delta \simeq 1$, $m_{\rm t} \simeq M_{\rm W}$. Note that the seventh $(4\pi^2)^{-1}$ factor has been taken into account in the definition of Chern-Simons terms (see eq. (3)). Actually, the presence of nonzero temperature strongly enhances the magnitude of the $\delta_{\rm ms}$.

If at high temperatures there is no nonzero classical scalar field (see, however, the discussion of this point in the appendix) then the estimate for $\delta_{\rm ms}$ should be:

$$\delta_{\rm ms} \simeq DN_{\rm d}\left(\langle\varphi^\dagger\varphi\rangle_T^{1/2}\langle p^{-1}\rangle\right)^{12} \simeq 2\times 10^{-22}, \tag{39}$$

for $\langle\varphi^\dagger\varphi\rangle_T = \frac{1}{3}T^2$, $\langle p^{-1}\rangle = \pi^2/12\, T\zeta(3)$.

Here quantity $\langle\varphi^\dagger\varphi\rangle_T$ represents the fluctuations of the scalar field, $\langle p^{-1}\rangle$ takes into account the average virutality in the quark propagator due to the absorption and emission of the scalar particles.

If at high temperatures there is a nonzero scalar condensate with $\langle\varphi\rangle \simeq g_{\rm W}T$ and typical inhomogeneities of order $(g_{\rm W}^2 T)^{-1}$ then the diagram in fig. 2 becomes effectively a two-loop one (the φ field goes to vacuum with virtually less than the effective fermion mass at high temperatures, $m_{\rm q}^2 = \frac{1}{6}g_{\rm s}^2T^2$ [9] and therefore the φ-lines should be considered as mass insertions). A simple estimate gives

$$\delta_{\rm ms} \simeq \frac{{\rm O}(10)}{4\pi^2} D\left(\frac{\langle\varphi\rangle}{m_{\rm q}}\right)^{12} \simeq 10^{-16}\text{–}10^{-20} \tag{40}$$

for, e.g. $\langle\varphi\rangle \simeq 2g_{\rm W}T \div g_{\rm W}T$. (One does not take into account the factor 6! due to the absence of scalar field propagators.)

For $\delta_{\rm ms}$ as in (39), $z \ll 1$ and baryon asymmetry equals

$$\Delta \simeq 10^{-11} - 10^{-12},$$

which is not very far from reality $\Delta \sim 10^{-8} - 10^{-10}$. (We assumed that there is no dilution of BAU i.e., $M_{\rm CW} \leqslant M_{\rm H} \leqslant M_{\rm crit}$.) For $\delta_{\rm ms} \sim 10^{-16}$ as in (40), $z \simeq 2-6 > 1$ and

$$\Delta \simeq (10^{-4} - 10^{-5})S(M_{\rm H}). \tag{41}$$

Of course, there is a large uncertainty in these estimates because we do not know exactly the magnitude of the maximal possible Chern-Simons density of the ground state.

Let us discuss now the value of the Higgs boson mass. There are three possibilities:

(i) the maximal BAU which may be generated by the decay of the ground state is in fact less than the observed value. This can happen if we have overestimated the maximal possible Chern-Simons density by a factor of the order of 10^6. That possibility seems to be rather improbable though not excluded. In this case the electroweak theory probably fails in explanation of BAU. We shall not discuss this possibility further.

(ii) The predicted BAU appears to be *strictly* equal to the observed value without any suppression. This seems to be rather unnatural due to the necessity of the fine tuning of the magnitude of the microscopic asymmetry. In this case cosmology implies the following bounds on the Higgs boson mass:

$$M_{\mathrm{CW}} \leqslant M_{\mathrm{H}} \leqslant M_{\mathrm{crit}}.$$

(iii) The number of baryons produced in the electroweak phase transition is larger than is needed. In this case baryon asymmetry should be diluted by one of the two possible mechanisms. Due to the very special dependence of the suppression factor on M_{H} this can happen only if $M_{\mathrm{H}} \simeq M_{\mathrm{CW}}$ or $M_{\mathrm{H}} \simeq M_{\mathrm{crit}}$. These values of the Higgs boson mass seem to be the most natural from the cosmological point of view. The discovery of the Higgs boson with these properties will strongly support the discussed scenario of the BAU generation.

If the magnitude of the *CP* violation in the standard model with three fermionic generations appears to be small, $\delta_{\mathrm{ms}} \leqslant 10^{-21}$ then one should find ways of increasing δ_{ms}. Consider possible extensions of the standard theory predicting the larger values of δ_{ms}.

(i) In the electroweak theory with the additional fourth generation of heavy fermions the magnitude of *CP* violation is larger than the quoted values (40), (39) by a factor of

$$\left(\frac{m'_{\mathrm{t}}}{m_{\mathrm{t}}}\right)^4\left(\frac{m_{\mathrm{t}}}{m_{\mathrm{c}}}\right)^2\left(\frac{m'_{\mathrm{b}}}{m_{\mathrm{b}}}\right)^4\left(\frac{m_{\mathrm{b}}}{m_{\mathrm{s}}}\right)^2 \simeq 10^{11}$$

for $m'_{\mathrm{t}} \simeq m'_{\mathrm{b}} \simeq M_{\mathrm{W}}$. In this case the variable z is definitely larger than one and the maximal possible BAU is generated. The Higgs boson mass should be M_{CW} or M_{crit} if the fourth generation is not too heavy ($m'_{\mathrm{t}} \simeq m'_{\mathrm{b}} \lesssim M_{\mathrm{W}}$). The sign of BAU is not directly related to the sign of *CP* violation in K^0 physics due to the appearance of the additional *CP* phases in the interactions of heavy quarks with W bosons.

(ii) The introduction of the additional Higgs doublet with all possible Yukawa couplings h gives $\delta_{ms} \sim O(10^{-4})$ if $h \simeq O(g_W)$. This is because CP violation arises already on the two-loop level. To find the sign of BAU one should know not only the CP phase in interactions of quarks with W bosons but CP phases in Yukawa couplings too. If the second doublet of scalars is not very heavy then only two values of the standard neutral Higgs boson mass are admitted.

(iii) There are two Higgs doublets in the supersymmetric extensions of the electroweak theory. CP violation arises only in twelfth order in Yukawa coupling constants because the first Higgs interacts only with u-type right-handed quarks and the second one only with d-type right-handed quarks. The amplification of δ_{ms} is given by the factor

$$P_{susy} \simeq \left(\frac{V_1}{V_2} + \frac{V_2}{V_1} \right)^6 ,$$

where V_1 and V_2 are vacuum expectation values of the first and the second Higgs fields, respectively. Say, for $V_2/V_1 \simeq 10$ we have $P_{susy} \simeq 10^6$ which ensures $z \gg 1$. The sign of BAU is determined by the K^0 decays. The constraints of the neutral Higgs boson mass remain the same as before if the Yukawa couplings are less than g_W*.

6. Conclusion

The main assumption which was made in the present paper is that the ground state of the Yang-Mills theories at high temperatures has nontrivial degeneracy with respect to the Chern-Simons density. Although this seems to be quite natural from the perturbative analysis of the infrared problem in thermodynamics of the gauge theories it is clear that this point deserves further investigation. To do this one should elaborate on the nonperturbative methods efficient in the strong coupling (infrared) limit of three-dimensional gauge theory. Probably, lattice calculations can clarify this matter.

Due to large uncertainties in our knowledge of the ground state of high temperature gauge-Higgs system and very high order of the CP violation in the standard electroweak theory with three fermionic generations we cannot definitely decide whether it is possible to generate the observed BAU in the minimal scheme. It is not excluded that cosmology requires additional structures like heavy fermionic generations, supersymmetry or two Higgs doublets. In any case the most natural values of

* Note that this prediction does not change if the masses of superpartners of quarks and leptons are larger than M_W because their masses do not originate from the standard scalar condensate.

the Higgs boson mass lie near $M_H \simeq M_{crit}$ or $M_H \simeq M_{CW}$ (in the absence of large Higgs or Yukawa coupling constants). It seems that the experimental search for Higgs bosons in this mass interval is the most important.

The author is deeply indebted to A. Yu. Ignatiev, V.A. Kuzmin, V.A. Matveev, V.A. Rubakov, A.N. Tavkhelidze, I.I. Tkachev, A.I. Bochkarev and S. Yu. Khlebnikov for numerous helpful discussions related to this work. I am grateful to S. Coleman, M. Dine, C. DeTar and L. Yaffe for useful conversations, criticism and suggestions and to M. Dine for careful reading of the manuscript. It is a pleasure to thank the Aspen Center for Physics, where this work was completed, for kind hospitality.

Note added in proof

The preexponential factor for the rate of the baryon number non-conservation in electroweak theory was estimated by P. Arnold, F. Accetta, E. Kolb, L. McLerran and M. Turner and appears to be large for the temperature of the first-order phase transition T_c. It corresponds to $\beta_0(T_c) \sim 10^6$ in our eq. (5). This results in a decrease of the upper bound on the Higgs boson mass up to $M_{crit} \sim 45$ GeV. I am grateful to Peter Arnold and Larry McLerran for explaining to me their results prior to publication and many interesting discussions.

A.I. Bochkarev and the present author have found that the cosmological bound on the Higgs boson mass practically does not depend on the t-quark mass. Taking into account the bound on m_t, $4m_t^4 < 2M_W^4 + M_Z^4 + \frac{3}{4}M_H^4$ (see, e.g. ref. [26]) we find $m_t < 80$ GeV. The details of this consideration will be published elsewhere.

Appendix

In this appendix we shall give arguments showing that the ground state of gauge theories at high temperatures may contain also a nonzero scalar field.

Consider the high temperature limit of the gauge-Higgs system. The lagrangian is

$$L = \tfrac{1}{4}\tilde{F}^a_{ij}\tilde{F}^a_{ij} + (D_i\tilde{\varphi})^\dagger(D_i\tilde{\varphi}) + M^2(\tilde{\varphi}^\dagger\tilde{\varphi}) + \lambda(\tilde{\varphi}^\dagger\tilde{\varphi})^2, \tag{A.1}$$

where $M^2 \simeq g_W^2 T$ is the effective scalar mass at high temperatures. The connection between 4-dimensional and 3-dimensional quantities reads $A_i = \sqrt{T}\tilde{A}$, $\varphi = \sqrt{T}\tilde{\varphi}$, $\tilde{g} = g\sqrt{T}$. Let us estimate the contribution to the 3-dimensional effective potential coming from gauge loops (scalar loops may be neglected due to the nonzero effective mass of the scalar). By power counting we get:

$$V^{(3)}(\tilde{\varphi}) = \tilde{\varphi}^2 M_W^2 \sum_{n=1}^{\infty} C_N \left(\frac{\tilde{g}^2}{M_W}\right)^N, \tag{A.2}$$

where M_W is the gauge boson mass which originates from the gauge-Higgs interaction, $M_W \simeq \tilde{g}\tilde{\varphi}$, C_N are some numerical factors, N is the number of loops. Coming back to the 4-dimensional quantities we obtain

$$V(\varphi)_{\text{magn}} = g^3\varphi^3 T \sum_{N=0}^{\infty} C_N \left(\frac{gT}{\varphi}\right)^N . \tag{A.3}$$

This means that nothing is known about the form of the effective potential for the small fields $\varphi \leqslant gT$. It is not excluded that $V(\varphi)$ has a minimum at some nonzero point $\varphi \sim gT$. In the latter case one expects the appearance of the slowly varying (the typical size of inhomogeneities is fixed by the Wφ interaction and is of the order of $(g^2T)^{-1}$) scalar field with the magnitude $\varphi \simeq gT$.

References

[1] J.C. Pati and A. Salam, Phys. Rev. Lett. 31 (1973) 661;
H. Georgi and S.L. Glashow, Phys. Rev. Lett. 32 (1974) 438
[2] A.D. Sakharov, Pisma ZhETF 5 (1967) 32;
V.A. Kuzmin, Pisma ZhETF 1 (1970) 335;
A. Yu. Ignatiev, N.V. Krasnikov, V.A. Kuzmin and A.N. Tavkhelidze, Proc. Int. Conf. Neutrino-77, vol. 2 (Nauka, Moscow, 1978) p. 283; Phys. Lett. 76B (1978) 436;
S. Weinberg, Phys. Rev. Lett. 4 (1979) 850;
A.Yu. Ignatiev, V.A. Kuzmin and M.E. Shaposhnikov, Phys. Lett. 87B (1979) 114
[3] I. Affleck and M. Dine, Nucl. Phys. B249 (1985) 361
[4] V.A. Kuzmin, V.A. Rubakov and M.E. Shaposhnikov, Phys. Lett. 155B (1985) 36
[5] G. 't Hooft, Phys. Rev. Lett. 37 (1976) 8; Phys. Rev. D14 (1976) 3432
[6] V.A. Kuzmin, V.A. Rubakov and M.E. Shaposhnikov, Proc. Int. Seminar Quarks-86, Tbilisi (15–17 April 1986)
[7] A.N. Redlich and L.C.R. Wijewardhana, Phys. Rev. Lett. 54 (1985) 970;
A.F. Niemi and G.W. Semenoff, Phys. Rev. Lett. 54 (1985) 2166;
K. Tsokos, Phys. Lett. 157B (1985) 413
[8] V.A. Rubakov and A.N. Tavkhelidze, Phys. Lett. 165B (1985) 109;
V.A. Rubakov, Prog. Theor. Phys. 75 (1986) 66;
V.A. Matveev, V.A. Rubakov, A.N. Tavkhelidze and V.F. Tokarev, Proc. Int. Seminar Quarks-86, Tbilisi (15–17 April 1986)
[9] A.D. Linde, Phys. Lett. 96B (1980) 289;
O.K. Kalashnikov, V.V. Klimov and E. Casado, Phys. Lett. 114B (1982) 49;
O.K. Kalashnikov, Fortschr. Phys. 32 (1984) 525
[10] C. DeTar, Phys. Rev. D32 (1985) 276
[11] N.H. Christ, Phys. Dev. D121 (1980) 1591
[12] 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
[13] R.F. Klinkhamer and N.S. Manton, Phys. Rev. D30 (1984) 2212
[14] V.A. Kuzmin and M.E. Shaposhnikov, Phys. Lett. 105B (1981) 163; Inst. for Nucl. Res. preprint (Moscow) P-0213 (1981);
A.D. Dolgov, Pisma ZhETF 29 (1979) 254;
A.D. Dolgov and Ya.B. Zeldovich, Usp. Fiz. Nauk 130 (1980) 559
[15] S. Dimopoulos and L. Susskind, Phys. Rev. D18 (1978) 4500
[16] D.J. Gross, R.D. Pisarski and L.G. Yaffe, Rev. Mod. Phys. 53 (1981) 43

[17] E.S. Fradkin, ZhETF 38 (1960) 157
[18] A.D. Linde, Rep. Prog. Phys. 42 (1979) 389
[19] A.D. Linde, Pisma ZhETF 23 (1976) 64
S. Weinberg, Phys. Rev. Lett. 36 (1976) 294
[20] L. Dolan and R. Jackiw, Phys. Rev. D9 (1979) 3320
[21] A.D. Linde, Phys. Lett. 70B (1977) 306; Nucl. Phys. B216 (1983) 421
[22] S. Coleman and E. Weinberg, Phys. Rev. D7 (1973) 1888
[23] E. Witten, Nucl. Phys. B177 (1981) 477
[24] A. Guth and E. Weinberg, Phys. Rev. Lett. 45 (1980) 1131
[25] Review of Particle Properties, Rev. Mod. Phys. 56 no. 2, part II (1984)
[26] N.V. Krasnikov, Yad. Fiz. 28 (1978) 549