A NEW INFLATIONARY UNIVERSE SCENARIO: A POSSIBLE SOLUTION OF THE HORIZON, FLATNESS, HOMOGENEITY, ISOTROPY AND PRIMORDIAL MONOPOLE PROBLEMS

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A new inflationary universe scenario is suggested, which is free of the shortcomings of the previous one and provides a possible solution of the horizon, flatness, homogeneity and isotropy problems in cosmology, and also a solution of the primordial monopole problem in grand unified theories.

There is now considerable interest in the cosmological consequences of symmetry breaking phase transitions, which occur in grand unified theories (GUTs) with the decrease of temperature at the very early stages of the evolution of the universe [1–3]. These phase transitions typically are strongly first order [4,5]. The lifetime of the supercooled symmetric phase $\varphi = 0$ (φ is the Higgs scalar field which breaks the symmetry) in some theories may be extremely large [2,3,6–8]. In that case the energy–momentum tensor of particles $\sim T^4$ in the phase $\varphi = 0$ almost vanishes in the course of the expansion of the universe, and the total energy–momentum tensor reduces to the vacuum stress tensor (cosmological term) $T_{\mu\nu}^{\text{vac}} = g_{\mu\nu} V(0)$, where $V(\varphi)$ is the effective potential of the theory at vanishing temperature [2,3]. This leads to an exponentially fast expansion of the universe, $a \sim e^{Ht}$. Here a is the scale factor, and H is the Hubble constant at that time,

$$H = [(8\pi/3M_{\text{P}}^2)\,V(0)]^{1/2},$$

where $M_{\text{P}} \approx 10^{19}$ GeV is the Planck mass [9]. Then at some comparatively small temperature T_{c} the symmetry breaking phase transition takes place, all the vacuum energy $V_{(0)}$ transforms into thermal energy [2,3], the universe is reheated up to the temperature $T_1 \approx V_{(0)}^{1/4}$, and its further evolution proceeds in a standard way [10] [‡1].

A most detailed discussion of this scenario is contained in a very interesting paper of Guth [12], where it is shown that the existence of a sufficiently long period of exponential expansion (inflation) in the early universe would provide a natural solution of the horizon and flatness problems in cosmology and of the primordial monopole problem in grand unified theories [13].

Unfortunately, however, this scenario in the form suggested in ref. [12] leads to some unacceptable consequences, recognized by Guth himself and by other authors who have studied this problem later, see e.g. refs. [8,14–18].The phase transition from the symmetric vacuum state $\varphi = 0$ to the asymmetric state $\varphi = \varphi_0$ proceeds by creation and subsequent expansion of bubbles containing some nonvanishing fields φ. In ref. [12] it was implicitly assumed that inside these bubbles the scalar field φ rapidly grows to $\varphi = \varphi_0$, all energy of the bubbles becomes concentrated in their walls and thermalization occurs only after the collision of the walls. If this qualitative picture were correct, the exponential expansion would be finished at the temperature T_{c}, at which the phase transition occurs. For the flatness problem to be solved the universe (the scale factor a) should grow at least 10^{28}

‡1 For an alternative scenario of the exponential expansion at the very early stages of the evolution of the universe, which may occur due to quantum gravity effects, see ref. [11].

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times during the exponential expansion period [12]. Since at this period the value of aT is constant, the critical temperature T_c should be 10^{28} times smaller than the temperature T_0, at which the exponential expansion starts. In the simplest SU(5) model [19] $T_0 \sim 10^{14}$ GeV, so that

$$T_c \lesssim 10^{-14}\ \text{GeV} \sim 0.1\ \text{K}. \tag{1}$$

No GUTs with such a fantastically small value of the critical temperature have been suggested.

There is also another problem with the above mentioned scenario. If the bubble wall collisions are necessary for the reheating of the universe, then after such a phase transition the universe becomes greatly inhomogeneous and anisotropic, which would contradict cosmological data [8,12,17,18].

In the present paper we would like to suggest an improved inflationary universe scenario, which is free of the above mentioned difficulties. With this purpose we shall consider the phase transitions in GUTs with the Coleman–Weinberg mechanism of symmetry breaking [20]. Phase transitions in such theories have been studied recently by many authors [14,15,21–23]. In our opinion, however, several important features of these phase transitions have escaped their attention. A detailed discussion of the phase transitions in GUTs with the Coleman–Weinberg mechanism of symmetry breaking will be contained in a subsequent publication. Here we shall only outline the main idea, which is essential for the understanding of the new inflationary universe scenario.

For definiteness let us consider the SU(5) grand unified theory [19], though most of what will be discussed here will not depend on the details of the model under consideration. The one-loop effective potential for the symmetry breaking SU(5) $\to$ SU(3) $\times$ SU(2) $\times$ U(1) in the Coleman–Weinberg version of this model at finite temperature T is [1–3, 14,15]

$$V(\varphi, T) = (18T^4/\pi^2) \times \int_0^\infty dx\, x^2 \ln\{1 - \exp[-(x^2 + 25g^2\varphi^2/8T^2)^{1/2}]\} + (5625/512\pi^2) g^4 (\varphi^4 \ln(\varphi/\varphi_0) - \varphi^4/4 + \varphi_0^4/4], \tag{2}$$

where $\varphi_0 \sim 10^{14}$–10^{15} GeV, $g^2 \sim 1/3$ is the gauge coupling constant. At $T \gg \varphi_0$ the symmetry in this theory was restored, $\varphi = 0$ [1–3]. With a decrease of temperature the absolute minimum of $V(\varphi)$ appears at $\varphi \approx \varphi_0$. However at any $T \neq 0$ the point $\varphi = 0$ remains a local minimum of $V(\varphi, T)$, since near $\varphi = 0$

$$V(\varphi, T) = \tfrac{75}{16} g^2 T^2 \varphi^2 - (5625/512\pi^2) g^4 \varphi^4 \ln(M_X/T) + (9/32\pi^2) M_X^4, \tag{3}$$

where $M_X^2 = \frac{25}{8} g^2 \varphi_0^2$. The phase transition with symmetry breaking proceeds from a strongly supercooled state $\varphi = 0$ at temperature T_c, which is many orders of magnitude smaller than φ_0 [14,15,21–23]. The shape of the potential $V(\varphi, T)$ for $T \ll \varphi_0$ is shown in fig. 1. The phase transition begins with the formation of bubbles of the field φ, which is a tunneling process [6]. One may argue that for $T_c \ll \varphi_0$ this process does not depend on the properties of $V(\varphi, T_c)$ at $\varphi \sim \varphi_0$, and the maximal value of the field φ inside the bubble immediately after its formation should be of the order of φ_1, where

$$V(\varphi_1, T_c) = V(0, T_c), \tag{4}$$

$\varphi_1 \ll \varphi_0$, see fig. 1. Indeed, a detailed study of the bubble formation in this theory performed in refs. [24,25] by means of computer calculations shows that at the moment of the bubble formation the maximal value of the field φ inside the bubble equals approximately $3\varphi_1$. Therefore inside the bubble

$$\varphi \lesssim 3\varphi_1 = \frac{12\pi T_c}{5g} \left(\frac{2}{3\ln(M_X/T_c)}\right)^{1/2} \ll \varphi_0. \tag{5}$$

This means that the (negative) mass squared of the

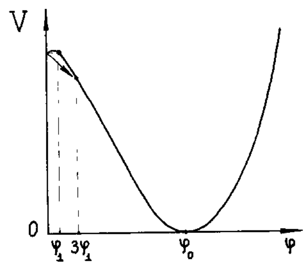

Fig. 1. Effective potential in the Coleman–Weinberg theory for $T \ll \varphi_0$. The arrow indicates the direction of the tunneling with bubble formation.

field φ inside the bubble is

$$-m^2 = -\tfrac{2}{15}\,\mathrm{d}^2V/\mathrm{d}\varphi^2 \lesssim 75g^2T_c^2 \sim 25T_c^2 . \qquad (6)$$

After the bubble formation the field φ inside the bubble gradually grows up to its equilibrium value $\varphi(T_1) \sim \varphi_0$. At the first stages of this process the field φ grows approximately as e^{mt}. Therefore it approaches its equilibrium value $\varphi(T_1)$ only after some period of time $\tau \gtrsim m^{-1} \sim 0.2T_c^{-1}$. A more complete investigation shows that τ is several times greater than m^{-1}; here for simplicity we shall take as an estimate

$$\tau \sim T_c^{-1} . \qquad (7)$$

It can also be easily shown that during most of this period the field φ inside the bubble remains much less than φ_0. Therefore during some time of the order of $\tau \sim T_c^{-1}$ the vacuum energy density $V(\varphi)$ remains almost equal to $V(0)$, and the part of the universe inside the bubble expands exponentially just as it expanded before the bubble creation. This simple observation has very important consequences for the theory of the phase transitions in the Coleman–Weinberg model.

Let us suppose that the phase transition in the Coleman–Weinberg SU(5) theory occurs at $T_c \sim 2 \times 10^6$ GeV, as it was claimed in ref. [15]. From eq. (3) it follows that with the parameters of the theory used in ref. [15] ($M_X \sim 6 \times 10^{14}$ GeV) the Hubble constant

$$H = [(8\pi/3M_P^2)\,V(0)]^{1/2}$$

is equal to 1.5×10^{10} GeV. Therefore during the exponential expansion period $\tau \sim T_c^{-1}$ the universe should grow $e^{H\tau}$ times, where

$$e^{H\tau} \sim e^{H/T_c} \sim e^{7500} \sim 10^{3260} . \qquad (8)$$

A typical size of the bubble at the moment of its creation is $O(T_c^{-1}) \sim 10^{-20}$ cm [25]. After the period of the exponential expansion this bubble will have a size of

$$10^{-20} \cdot e^{H\tau}\ \mathrm{cm} \sim 10^{3240}\ \mathrm{cm},$$

which is much greater than the size of the observable part of the universe $l \sim 10^{28}$ cm. Therefore the whole observable part of the universe is contained *inside one bubble,* so we see no inhomogeneities caused by the wall collisions. After some time of the order of τ after the bubble creation all the vacuum energy density $V(0)$ transforms into thermal energy $\sim T_1^4$, where in our model $T_1 \approx 0.15M_X \sim 10^{14}$ GeV. However the thermalization occurs now not due to the wall collisions, but due to the interactions of particles created by the classical homogeneous field φ, convergently oscillating near its equilibrium value $\varphi(T_1) \approx \varphi_0$ with a frequency of about 10^{14} GeV.

One can easily verify that the size of the particle horizon at the time of the phase transition was much greater than the size of the bubble $\sim T_c^{-1}$, i.e. all points inside the bubble were causally connected. After the exponential expansion period this causally connected domain covers the whole observable part of the universe, which solves the horizon problem [12].

Now let us remember that particle creation in the very early universe in general cannot make it completely isotropic, but makes it quasi-isotropic [10], i.e. locally isotropic in small domains of space of the size of the same order as or greater than the Planck length $l_P \sim 10^{-33}$ cm $\sim M_P^{-1}$ at the Planck time $t_P \sim 10^{-43}$ s, when the temperature was $T_P \sim M_P \sim 10^{19}$ GeV. Since before the phase transition the quantity aT was constant inside each isotropic domain of the universe, at the moment of the phase transition the typical size of the isotropic domain exceeds the bubble size $\sim T_c^{-1}$. Therefore the space–time inside the bubble was isotropic and the exponential expansion extends this isotropy to the whole observable part of the universe. (Moreover, the remaining small anisotropy inside the bubble decreases rapidly during the exponential expansion period [18].) This may solve the long-standing problem of the space–time isotropy in our universe.

Density fluctuations inside the bubble immediately after its formation are negligibly small as compared with $V(0)$, i.e. the space inside the bubble is almost homogeneous. Then the exponential expansion extends this homogeneity to the whole observable part of the universe, which explains the large-scale homogeneity of the universe.

One may argue that it is not very good to obtain an absolutely homogeneous universe, since in that case it would be difficult to understand the origin of galaxies. However, as will be explained in a separate publication (see also refs. [7,26]), the necessary inhomogeneities may be generated after a subsequent phase transition with a smaller degree of supercooling. Moreover, as is shown in ref. [27], the perturbations necessary for galaxy formation arise due to quantum gravity effects just after phase transitions of the type considered above in GUTs with the unification scale $\Lambda \sim 10^{17}-10^{18}$ GeV, which is not unrealistic [28].

From our results it follows that the size of the universe l_1 after the phase transition should exceed 10^{3240} cm, and the temperature T_1 is of the order of 10^{14} GeV. Therefore the total entropy of the universe should exceed $(l_1 T_1)^3 \sim 10^{10000}$, which explains why the total entropy of the universe exceeds 10^{85} and simultaneously solves the flatness problem [12,21].

It is known that the primordial monopoles in GUTs are created only in the points, in which bubbles with different types of Higgs field φ collide [13]. Therefore in our scenario no monopoles are created in the observable part of the universe, which solves the primordial monopole problem in GUTs [13]. For the same reason there will be no domain walls in the observable part of the universe in the theories with broken discrete symmetries [29], and in particular in the theories with spontaneously broken *CP* invariance. This helps to solve the problem of the baryon asymmetry of the universe [30]. Moreover, in the scenario under consideration there appears an additional source of baryon asymmetry. The standard mechanism is connected with the decay of the X, Y bosons and Higgs mesons, which appear in the course of the reheating of the universe during the phase transition. It is clear, however, that the baryon asymmetry generated by the decay of the Higgs mesons may be generated by the decay of the classical Higgs field vacuum $\varphi = 0$ as well. Note also that, whereas in a standard scenario the particles created by the decay of the X, Y and Higgs mesons are only a few percent of the total amount of particles [30], in our case all particles which appear after the phase transition are created by the oscillating classical Higgs field φ during the phase transition.

The new inflationary universe scenario discussed above is, of course, oversimplified. To get a complete scenario, one should analyse the phase transitions in the Coleman–Weinberg model more accurately taking into account the renormalization group equation [21] and the nonperturbative effects [15]. This analysis should be performed simultaneously with the investigation of the effects connected with the nonvanishing curvature and rapid expansion of the universe, which become important for $H > T_c$. One may ask e.g. whether it is possible for the universe to be in a state with temperature T_c smaller than the Hawking temperature $T_H = H/2\pi$, to which consequences the terms $\sim R\varphi^2$ in the effective potential may lead etc. [23]. An investigation of these problems is rather involved and will be contained in a separate publication. Our preliminary result is that there exists an *improved Coleman–Weinberg theory*, in which $d^2V/d\varphi^2 = 0$ at $\varphi = 0$ not in Minkowski space, but rather in de Sitter space with the curvature R determined by the vacuum energy density $V(\varphi)$ at the symmetric point $\varphi = 0$. In some versions of this theory the phase transition with symmetry breaking occurs due to nonperturbative effects [15]. The kinetics of this phase transition is somewhat more complicated than that described above, but the main feature of the new inflationary universe scenario remains intact: The field φ approaches its equilibrium value $\varphi(T_1) \sim \varphi_0$ during the period $\tau \gg H^{-1}$, which just leads to the desirable inflation of the universe discussed in the present paper

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