INFLATION CAN BREAK SYMMETRY IN SUSY

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It is shown that the exponential expansion of the universe $a(t) \sim \exp(Ht)$ leads to a simultaneous generation of many different types of classical scalar fields φ with the amplitude $\varphi \gg H$. These effects may help us to solve the problem of symmetry breaking in SUSY GUTs and the problem of primordial monopoles which appear after primordial inflation.

Attempts to construct a unified supersymmetric theory of all fundamental interactions are of great interest now. Supersymmetric theories have many advantages over the standard theories. However there exist also some problems specific to supersymmetric theories. One of such problems is the problem of symmetry breaking in SUSY GUTs.

It is well-known that the effective potentials $V(\Phi)$ in supersymmetric theories typically have many different minima of the same depth, see e.g. a discussion of this question in refs. [1–3]. For example, in the minimal SU(5) supersymmetric theory [4,5] with the superpotential

$$W = \lambda\left(\tfrac{1}{2}m \operatorname{tr}\Phi^2 + \tfrac{1}{3}\operatorname{tr}\Phi^3\right), \tag{1}$$

where Φ is the adjoint Higgs field, the effective potential $V(\Phi)$ has four degenerate minima: SU(5) invariant minimum, SU(4) × U(1) minimum, SU(3) × SU(2) × U(1) minimum, and the Dragon minimum [2]. One can increase the energy of the SU(5) minimum e.g. by adding the term $W' = \alpha z(\operatorname{tr}\Phi^2 - \mu^2)$ to the superpotential (1) (here z is some chiral singlet superfield); however, the degenerate minima SU(4) × U(1) and SU(3) × SU(2) × U(1) are still present in such a non-minimal theory [5]. The number of different degenerate minima of the effective potential increases considerably when one adds to W(1) the superpotential of other Higgs fields φ[3].

The degeneracy of the energy of different minima of $V(\Phi)$ is removed by gravitational effects, but in a rather unfortunate way: the lowest energy state in the minimal theory (which will be studied in the present paper) becomes the SU(5) invariant state [1]. Moreover, even in the cases in which gravitational effects are small, cosmological considerations pick up not the SU(3) × SU(2) × U(1) phase, but just the SU(5) symmetric phase.

Indeed, in the very early universe at the temperature $T \gg m$ the only minimum of $V(\Phi)$ was the SU(5) minimum, corresponding to $\Phi = 0$, and the SU(5) symmetry was restored [6]. All other minima appear only at a sufficiently small temperature T. The leading contribution to $V(\Phi, T)$ due to high-temperature effects is given by $-\frac{1}{90}\pi^2 T^4(N_B + \frac{7}{8}N_F)$, where N_B and N_F are the (effective) numbers of bosons and fermions with masses $M \lesssim T$ [6]. The value of $N_B + \frac{7}{8}N_F$ is maximal in the SU(5) symmetric phase, in which $\Phi = 0$ and all vector bosons and fermions are massless [7]. This means, that the SU(5) minimum at all temperatures is a global minimum of the effective potential, which makes it absolutely unclear how the phase transition from this phase to the phase SU(3) × SU(2) × U(1) could occur.

An interesting possibility to overcome this difficulty was suggested in refs. [7,8]. Let us consider e.g. the phase transition SU(5) → SU(3) × SU(2) × U(1) in the minimal supersymmetric SU(5) theory, which occurs due to the ap-

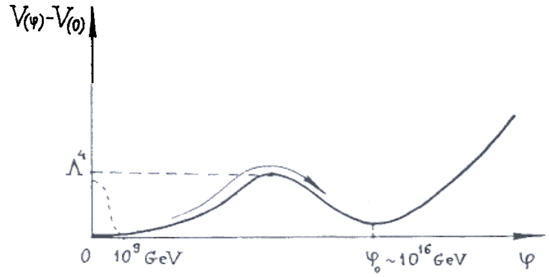

Fig. 1. Effective potential $V(\Phi, T)$ in the minimal supersymmetric SU(5) theory at $T \leqslant 10^9$ GeV for $\Phi = \varphi(\frac{2}{15})^{1/2}\text{diag}(1, 1, 1, -\frac{3}{2}, -\frac{3}{2})$. Broken line corresponds to $V(\varphi, T)$ in the SU(5) confinement phase at $\varphi \leqslant 10^9$ GeV. The arrow shows the behaviour of the field $\varphi = \langle\varphi^2\rangle^{1/2}$ in the inflationary universe.

pearance of the classical scalar field

$$\Phi = \varphi(\tfrac{2}{15})^{1/2}\text{diag.}(1, 1, 1, -\tfrac{3}{2}, -\tfrac{3}{2}) . \tag{2}$$

One-loop effective potential $V(\varphi, T)$ in this theory is shown in fig. 1. Now let us take into account, following refs. [7,8], that at $\varphi = 0$, $T \leqslant 10^9$ GeV all matter in the SU(5) theory should be in the confinement phase. In this phase the one-loop results for $V(\varphi, T)$ are unreliable. The number of degrees of freedom in the confinement phase is considerably reduced as compared with the normal phase, which leads to an increase of $V(0, T)$ by $O(T^4)$. It was concluded therefore that at $T \leqslant 10^9$ GeV the phase transition from $\varphi = 0$ to $\varphi = \varphi_0 \sim 10^{16}$ GeV (φ_0 corresponds to the minimum of $V(\varphi)$ in the SU(3) × SU(2) × U(1) phase) becomes possible [7]. Later it was understood that such a phase transition may occur only if the effective potential $V(\varphi)$ is extremely flat. This is necessary in order to reduce the gravitational corrections to $V(\varphi_0, T)$ [1], and to make possible the tunneling through the barrier ΔV of the height Λ^4 between $\varphi = 0$ and $\varphi = \varphi_0$, see fig. 1. In refs. [7,8] it was suggested to consider the theories in which the effective coupling constant $\lambda \sim 10^{-12}$ (1), and $\Lambda \sim 10^{10}$ GeV. In such theories gravitational corrections to $V(\varphi)$ actually are very small [9]. However, by the use of the methods developed in ref. [10] it can be shown that the rate of tunneling through the barrier of the height $\sim\Lambda^4$ at temperature $T \leqslant \Lambda$ is suppressed by a factor of the order $\exp(-\lambda^{3/2}) \sim \exp(-\varphi_0^3/\Lambda^3) \sim \exp(-10^{18})$. Therefore the phase transition could actually occur only if there would be no barrier between the phases $\varphi = 0$ and $\varphi = \varphi_0$. The situation becomes even more complicated if one takes into account, that the confinement phase at $T \leqslant 10^9$ GeV may exist only at $\varphi \leqslant 10^9$ GeV, where all SU(5) vector fields have small masses $M \leqslant 10^9$ GeV. At $\varphi \gg 10^9$ GeV only SU(3) × SU(2) × U(1) vector particles are massless, whereas all other particles acquire masses $M \gg 10^9$ GeV, and the SU(5) confinement disappears. Therefore the SU(5) confinement cannot lift up the whole SU(5) minimum; it can create only a local maximum of $V(\varphi, T)$ at $\varphi \leqslant 10^9$ GeV, see fig. 1, but the value of $V(\varphi, T)$ at $\varphi \sim 10^{10}$ GeV remains smaller than $V(\varphi_0, T)$, and the phase transition to the SU(3) × SU(2) × U(1) phase cannot occur.

The main aim of the present paper is to suggest a possible solution of the problem of symmetry breaking in SUSY GUTs in the context of the new inflationary universe scenario [11–15] (for a review of the present status of this scenario see ref. [16]).

The possibility to be discussed in this paper is based on the following observation. As it is shown in ref. [17], fluctuations of a scalar field φ with a small mass $m^2 \ll H^2$ in the de Sitter universe expanding as $a(t) \sim e^{Ht}$ are extremely large:

$$\langle\varphi^2\rangle = 3H^4/8\pi^2m^2 . \tag{3}$$

To be more precise, if the field φ interacts with the scalar curvature R as $\frac{1}{2}\xi R\varphi^2$, one should write $m^2 + \xi R$ instead of m^2 in (3). However we will consider here the theories without such interactions. Note, that such an interaction does not appear in the theory of scalar fields coupled to $N = 1$ supergravity [18].

Eq. (3) is valid for the eternally existing de Sitter universe. In the inflationary universe scenario, in which the hot Friedmann universe becomes exponentially expanding at some moment t_0, the corresponding expression for $\langle\varphi^2\rangle$ looks as follows [11,15,19]:

$$\langle\varphi^2\rangle = (3H^4/8\pi^2m^2) \times \{1 - \exp[-(2m^2/3H)(t - t_0)]\} . \tag{4}$$

Thus, at $t - t_0 \ll 3H/m^2$ the value of $\langle\varphi^2\rangle$ grows linearly,

$$\langle\varphi^2\rangle = (H^3/4\pi^2)(t - t_0)\,, \tag{5}$$

and then it approaches the limiting value $3H^4/8\pi^2m^2$ (3). It is very important (and rather unusual) that the leading contribution to $\langle\varphi^2\rangle$ goes from the fluctuations of the field φ with the exponentially large wavelength, $k^{-1} \sim H^{-1}\,e^{H(t-t_0)}$. Therefore at a scale $l \lesssim H^{-1}\,e^{H(t-t_0)}$ the fluctuations of the field φ practically cannot be distinguished from the homogeneous classical field $\varphi = \langle\varphi^2\rangle^{1/2}$. This effect initially served as a basis for the version of the new inflationary universe scenario suggested in the last paper of ref. [11] and also in refs. [15,20]. However eqs. (4), (5) actually are valid not only for the field φ which is responsible for the inflationary phase transition, but for *any* scalar field with $m^2 \ll H^2$.

The growth of the "classical field" $\varphi = \langle\varphi^2\rangle^{1/2}$ is a rather surprising effect, since the field φ grows practically independently of the sign of m^2 (at $|m^2| \ll H^2$). In particular, the field φ can grow from the minimum to a maximum (or *over* the maximum) of the effective potential, though it could seem energetically unfavourable. The reason of such a strange effect is analogous to the reason of particle creation in expanding universe (which at the first glance also could seem energetically unfavourable). Long-range fluctuations in the inflationary universe grow due to the vacuum rearrangement during the transition from the hot Friedmann universe, in which long-range fluctuations are suppressed by the high-temperature effects, to the de Sitter universe in which the density of long-range fluctuations is extremely large [16].

In the inflationary universe scenario exponential expansion should occur during some time $\Delta t = t - t_0 \gtrsim 10^2H^{-1}$ [21,16]. If the mass of the field φ is sufficiently small (one or two orders smaller than H), then during the whole interval Δt the field grows linearly (5). In this case

$$\varphi = \langle\varphi^2\rangle^{1/2} = (H/2\pi)(H\Delta t)^{1/2}\,, \tag{6}$$

where $H\Delta t \gtrsim 10^2$. From (6) it follows therefore that during the inflation the amplitude of the field φ approaches some value $\varphi \gtrsim H$. (In the theories, in which density perturbations $\delta\rho/\rho$ are of the order 10^{-4} after inflation, the typical value of $H\Delta t$ is 10^3–10^5, which yields $\varphi \gtrsim 10H$.) Now let us consider the new inflationary universe scenario with $H \sim 10^{16}$–10^{17} GeV (such values of H may appear e.g. due to symmetry breaking in supergravity in the primordial inflation scenario [22,23]). Typical masses of the Higgs fields Φ in the minimal supersymmetric SU(5) model [4,5] are $\sim 10^{15}$–10^{16} GeV, but they may be many orders smaller [8,9]. The energy density of the SU(5) fields Φ is negligibly small compared with the energy density of some other scalar fields φ which are responsible for the (primordial) inflation with large $H \sim 10^{16}$–10^{17} GeV. Therefore the SU(5) fields Φ have almost no effect on the process of inflation and reheating in our scenario. On the other hand, fluctuations $\langle\Phi^2\rangle$ of the field Φ, just as of all other scalar fields with $m^2 \ll H^2$, grow in time according to (4)–(6). Therefore during the inflation the classical fields Φ of all types with the amplitude $\Phi \gtrsim H \sim 10^{16}$–$10^{17}$ GeV are generated. These fields are practically homogeneous at a scale $l \lesssim H^{-1}\,e^{H\Delta t}$.

After the end of inflation the field Φ stops its growth and becomes convergently oscillating in the vicinity of a nearest minimum of $V(\Phi)$. Therefore after the end of inflation the universe becomes divided into many domains with all possible types of symmetry breaking, the typical size of each domain being many orders greater than the size of the observable part of the universe $l \sim 10^{28}$ cm. In particular, there will be many (in open universe – infinitely many) domains of the phase SU(3) × SU(2) × U(1) in one of which we live now [16].

We would like to emphasize the difference between the above-mentioned mechanism of symmetry breaking and the standard one. Usually it is assumed that it is impossible for the phase transition to occur from the global minimum of the effective potential $V(\varphi)$ to any local minimum of $V(\varphi)$. From our results it follows, however, that such a phase transition becomes effectively possible in the new inflationary universe scenario due to the anomalous growth of the long-range scalar field fluctuations in the exponentially expanding universe. Usually the phase transition proceeds to one preferred phase in the whole universe. In our case the phase transition proceeds with a comparable probability to many different phases. Each of these phases fills a domain (mini-universe) of a size exceeding the size of the observable part of our universe. All these phases with a non-negative vacuum energy density are practically stable. The only phase in which life of our type may exist is the phase $SU(3)\times U(1)$ with vanishing vacuum energy [16]. This phase appears after symmetry breaking in the $SU(3)\times SU(2)\times U(1)$ domains, and proves to be absolutely stable in the theories under consideration [1].

One could argue, however, that after reheating of the universe in the end of inflation the phase transition with the SU(5) symmetry restoration may occur, and the problem of symmetry breaking in SUSY GUTs may appear again. Fortunately, as it was shown in refs. [24,23], reheating in the theories with large H is rather ineffective, and the temperature T_R of the universe after reheating typically is many orders smaller than the critical temperature of symmetry restoration in SUSY GUTs $T_c \sim 10^{16}$ GeV [23].

We would like to note also, that the large size of domains, just as the large size of bubbles in the first version of the new inflationary universe scenario [11], implies that no monopoles appear in the observable part of the universe after the phase transition. This solves the primordial monopole problem, which was claimed to exist in the primordial inflation scenario [22]. In the non-minimal SU(5) theory containing e.g. the term $W' = \alpha z(\mathrm{tr}\,\Phi^2 - \mu^2)$ in the superpotential the primordial monopole problem can be solved even in a more simple way [23].

At present it seems that the new inflationary universe scenario can be completely realized in the context of $N = 1$ supergravity coupled to matter [22–25]. It is pleasant to note that the inflationary universe is not ungrateful to supersymmetric theories. It provides a possible solution to the gravitino problem [26,23] and to the problem of symmetry breaking in SUSY GUTs considered in the present paper.

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