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DECAY OF THE FALSE VACUUM IN THE VERY EARLY UNIVERSE

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ABSTRACT

The scenario whereby a false vacuum undergoes a first order phase transition to a true vacuum in the very early Universe is illustrated by considering the $SU(5)$ and $SO(10)$ unified gauge models. The initial number density of superheavy magnetic monopoles produced in such a first order phase transition can be strongly suppressed.

The very early Universe provides us with an important laboratory for testing some of the predictions that emerge from spontaneously broken unified gauge theories^{1),2)} (UGTs). In trying to deduce the implications of UGTs for cosmology, however, it is important to understand the nature of the phase transitions^{*)} that took place in the very early Universe. Studies^{4),5)} of spontaneously broken UGTs (such as $SU(5)$ ²⁾ and $SO(10)$ ⁶⁾) at finite temperatures have revealed that, as the Universe expands and cools after the hot big bang, the system may (and in some cases must) pass through intermediate stages that are not part of the usual gauge hierarchy at zero temperature. In particular, some of these phase transitions may be strongly first order and this could have, for instance, important implications for the expansion rate, the production density of magnetic monopoles^{4),7),8)} and the generation of baryon asymmetry⁹⁾ in the very early Universe.

Recently, Sato¹⁰⁾ has investigated the progress of a first order phase transition (which proceeds through bubble formation) from the false to the true vacuum in an expanding Universe. In particular, he has argued that if the nucleation rate of bubbles of the true vacuum in the false vacuum is sufficiently low (quantum tunnelling¹¹⁾ is expected to be a slow process) the Universe would remain in the metastable vacuum forever. This is because, after a period of time, the vacuum energy density of the false vacuum dominates the radiation energy density, subsequent to which the part of the Universe still in the false vacuum starts expanding exponentially. The progress of the phase transition is delayed even more by this rapid cosmic expansion (the bubbles can expand, at most, with the speed of light) and the phase transition is never completed. The Universe essentially remains in the false vacuum forever !

Finding the above scenario somewhat disturbing, to say the least, we discuss in this paper the decay of the false vacuum in the very early Universe in the context of $SU(5)$ and $SO(10)$ gauge models. We find, in both cases, that for a reasonable choice of Higgs parameters of the theory, the first order phase transition is completed, even if the initial nucleation rate is very slow. Indeed, in the scenario that we envisage, the initial nucleation of bubbles of the true vacuum can be neglected in the expanding Universe. Bubble formation only becomes important at temperatures very near some temperature T_{c2} at which the false vacuum ceases to exist even as a local minimum of the effective potential. The phase transition to the true vacuum takes place almost instantaneously at T_{c2} . The subsequent release of latent heat heats the Universe back up to some

*) Phase transitions in spontaneously broken gauge theories have first been studied in Ref. 3).

temperature T_R . The value of T_{C2} depends on the parameters of the Higgs potential and determines the amount of supercooling the system undergoes during the phase transition.

We also discuss the density of superheavy magnetic monopoles produced in such a first order phase transition⁴⁾. As examples, we consider the phase transition from $SU(5)$ to $SU(3)_C \times SU(2)_L \times U(1)$ via $SU(4) \times U(1)$ and from $SO(10)$ to $SO(6) \times SO(4)$ ($\equiv SU(4)_C \times SU(2)_L \times SU(2)_R$) via $SO(7) \times SO(3)$. We find, in both cases, that with a suitable choice of parameters, the initial monopole density is strongly suppressed.

PHASE TRANSITION FROM $SU(5)$ TO $SU(3)_C \times SU(2)_L \times U(1)$

The $SU(5)$ gauge group can be broken down to $SU(3)_C \times SU(2)_L \times U(1)$ by employing a 24-dimensional Higgs field ϕ . This field may be regarded as a Hermitean, traceless 5×5 matrix which transforms under $SU(5)$ as

$$\phi \longrightarrow U \phi U^\dagger, \quad U \in SU(5). \quad (1)$$

The most general $SU(5)$ invariant renormalizable Higgs potential for ϕ is

$$V(\phi) = \frac{a}{4} (\text{Tr } \phi^2)^2 + \frac{b}{2} \text{Tr } \phi^4 + \frac{c}{3} \text{Tr } \phi^3 - \frac{1}{2} \mu^2 \text{Tr } \phi^2. \quad (2)$$

We assume that $a \sim b \sim g^2$ and $c \sim g\mu$, where g is the $SU(5)$ gauge coupling. Radiative corrections to $V(\phi)$ will not be considered in this work.

The stable phase of the $SU(5)$ gauge theory with one 24-plet of Higgs fields is determined by the absolute minimum of $V(\phi)$. There are three possible phases^{4),12)} which correspond to the unbroken groups $SU(5)$, $SU(4) \times U(1)$ and $SU(3)_C \times SU(2)_L \times U(1)$. The range of parameters of $V(\phi)$ for which the various stable phases are realized is discussed in Ref. 4).

In the following, we shall be interested not only in the absolute but also in the relative minima of the potential $V(\phi)$ which correspond to metastable phases of the $SU(5)$ system. It has recently been shown¹³⁾ that even at the local minima of $V(\phi)$, the possible subgroups of $SU(5)$ that leave $\langle \phi \rangle$ invariant are $SU(5)$, $SU(4) \times U(1)$ and $SU(3)_C \times SU(2)_L \times U(1)$. We wish to investigate the range of parameters of $V(\phi)$ for which each of the above three phases can exist either as a stable or a metastable phase. To this end, it is convenient to define

$$\begin{aligned}\hat{V}(\hat{\phi}) &\equiv \frac{2b^3}{c^4} V(\phi) \\ &= \text{Tr} \hat{\phi}^4 + \frac{2}{3} \text{Tr} \hat{\phi}^3 + \frac{1}{2} \eta (\text{Tr} \hat{\phi}^2)^2 + \xi \text{Tr} \hat{\phi}^2,\end{aligned}\quad (3)$$

where

$$\begin{aligned}\hat{\phi} &= \frac{b\phi}{c}, \quad \eta = \frac{a}{b}, \quad \xi = -\frac{\mu^2 b}{c^2}, \\ b > 0, \quad \eta > -\frac{7}{15}, \quad c > 0.\end{aligned}\quad (4)$$

i) SU(5) phase

This phase can exist, at least as a local minimum, for $\xi > 0$, i.e., to the right of curve C in Fig. 1.

ii) SU(4) $\times$ U(1) phase

The expectation value $\langle \hat{\phi} \rangle$ takes the form

$$\langle \hat{\phi} \rangle = \text{diag} (x, x, x, x, -4x),$$

where^{*)}

$$x = \frac{3}{4} \cdot \frac{1 + \sqrt{1 - \frac{8}{9} \xi (10\eta + 13)}}{10\eta + 13}.\quad (5)$$

Let us now discuss the fluctuations $\delta \hat{\phi}$ about $\langle \hat{\phi} \rangle$. The most general $\delta \hat{\phi}$ that is orthogonal to the gauge orbit of $\langle \hat{\phi} \rangle$ can be written as

$$\delta \hat{\phi} = \sum_{i=1}^{15} \delta x_i L_i + \delta x_{16} U,$$

where

$$L_i = \left(\begin{array}{c|c} \lambda_i & 0 \\ \hline 0 & 0 \end{array} \right), \quad U = \left(\begin{array}{c|c} 1_4 & 0 \\ \hline 0 & -4 \end{array} \right), \quad (6)$$

^{*)} Note that $x = (3/4) \cdot [1 - \sqrt{1 - (8/9)\xi(10\eta+13)}] / (10\eta+13)$ never corresponds to a local or an absolute minimum.

and λ_i are the $SU(4)$ generators normalized according to $\text{Tr}(\lambda_i \lambda_j) = \delta_{ij}$. Substituting $\langle \hat{\phi} \rangle + \delta \hat{\phi}$ for $\hat{\phi}$ in (3), we obtain

$$\begin{aligned} \hat{V}(\langle \hat{\phi} \rangle + \delta \hat{\phi}) &= \hat{V}(\langle \hat{\phi} \rangle) + 2 \left[(10\eta + 3)x^2 + x + \frac{1}{2} \xi \right] \sum_{i=1}^{15} \delta x_i^2 \\ &\quad + 120 \left[(10\eta + 13)x^2 - x + \frac{1}{6} \xi \right] \delta x_{16}^2 + O(\delta x^3). \end{aligned} \quad (7)$$

The local stability of the $SU(4) \times U(1)$ phase requires that

$$\alpha = (10\eta + 3)x^2 + x + \frac{1}{2} \xi > 0$$

and

$$\beta = (10\eta + 13)x^2 - x + \frac{1}{6} \xi > 0. \quad (8)$$

It then follows that the $SU(4) \times U(1)$ phase can exist either as a local or an absolute minimum for

$$-\frac{1}{8}(10\eta + 7) < \xi < \frac{9}{8} \cdot \frac{1}{10\eta + 13}, \quad \eta > -\frac{7}{15}, \quad (9)$$

i.e., between curves A and F in Fig. 1. Note that $\alpha > 0$ for

$$-\frac{1}{8}(10\eta + 7) < \xi \leq \frac{9}{8} \cdot \frac{1}{10\eta + 13},$$

$\alpha = 0$ for $\xi = -(10\eta + 7)/8$ and $\alpha < 0$ for

$$\xi < -\frac{1}{8}(10\eta + 7).$$

On the other hand, $\beta = 0$ for

$$\xi = \frac{9}{8} \cdot \frac{1}{10\eta + 13}$$

and $\beta > 0$ for

$$\xi < \frac{9}{8} \cdot \frac{1}{10\eta + 13}.$$

iii) $SU(3)_C \times SU(2)_L \times U(1)$ phase

The expectation value $\langle \hat{\phi} \rangle$ now takes the form

$$\langle \hat{\phi} \rangle = \text{diag} \left(y, y, y, -\frac{3}{2}y, -\frac{3}{2}y \right),$$

where^{*)}

$$y = \frac{1}{2} \cdot \frac{1 + \sqrt{1 - 8\xi(15\eta + 7)}}{15\eta + 7}. \quad (10)$$

The most general fluctuation $\delta \hat{\phi}$ about $\langle \hat{\phi} \rangle$, which is orthogonal to its gauge orbit, can now be written as

$$\delta \hat{\phi} = \sum_{i=1}^8 \delta y_i S_i + \sum_{i=9}^{11} \delta y_i T_i + \delta y_{12} V, \quad (11)$$

where

$$S_i = \left(\begin{array}{c|c} \lambda_i & 0 \\ \hline 0 & 0 \end{array} \right), \quad T_i = \left(\begin{array}{c|c} 0 & 0 \\ \hline 0 & \sigma_i \end{array} \right), \quad V = \left(\begin{array}{c|c} 1_3 & 0 \\ \hline 0 & -\frac{3}{2}1_2 \end{array} \right),$$

with λ_i, σ_i the $SU(3)_C$ and $SU(2)_L$ generators respectively, normalized according to $\text{Tr}(\lambda_i \lambda_j) = \delta_{ij}$, $\text{Tr}(\sigma_i \sigma_j) = \delta_{ij}$. Substituting $\langle \hat{\phi} \rangle + \delta \hat{\phi}$ in (3), we obtain

$$\begin{aligned} \hat{V}(\langle \hat{\phi} \rangle + \delta \hat{\phi}) &= \hat{V}(\langle \hat{\phi} \rangle) + \frac{1}{2} \left[(15\eta + 12)y^2 + 4y + 2\xi \right] \sum_{i=1}^8 \delta y_i^2 \\ &\quad + \frac{1}{2} \left[(15\eta + 27)y^2 - 6y + 2\xi \right] \sum_{i=9}^{11} \delta y_i^2 \\ &\quad + \frac{45}{4} \left[(15\eta + 7)y^2 - \frac{2}{3}y + \frac{2}{3}\xi \right] \delta y_{12}^2 + O(\delta y^3). \end{aligned} \quad (12)$$

The local stability of the $SU(3)_C \times SU(2)_L \times U(1)$ phase requires that

^{*)} Note that $y = \frac{1}{2} \cdot [1 - \sqrt{1 - 8\xi(15\eta + 7)}] / (15\eta + 7)$ corresponds to a local minimum for $15\eta + 2\xi + 8 < 0$ (i.e., to the left of curve G in Fig. 1) which lies outside the region of interest here.

$$\gamma = (15\eta + 12)y^2 + 4y + 2\xi > 0,$$

$$\delta = (15\eta + 27)y^2 - 6y + 2\xi > 0,$$

(13)

$$\epsilon = (15\eta + 7)y^2 - \frac{2}{3}y + \frac{2}{3}\xi > 0.$$

It follows that the $SU(3)_C \times SU(2)_L \times U(1)$ phase can exist either as a local or an absolute minimum for the following range of parameters:

$$\eta \geq -\frac{1}{3}, \quad \xi < -\frac{15\eta + 3}{32}$$

and

$$-\frac{7}{15} < \eta \leq -\frac{1}{3}, \quad \xi < \frac{1}{8(15\eta + 7)} \quad (14)$$

i.e., to the left of curve D in Fig. 1.

At zero temperature, the parameters ξ and η must be chosen so that the $SU(3)_C \times SU(2)_L \times U(1)$ phase corresponds to the absolute minimum of the Higgs potential. This is arranged by choosing the point $(\xi, \eta)_{T=0}$ to lie to the left of curve E (Fig. 1). If $(\xi, \eta)_{T=0}$ lies between the curves E and F, the Higgs potential possesses, in addition, a local minimum corresponding to a metastable $SU(4) \times U(1)$ vacuum. However, if $(\xi, \eta)_{T=0}$ lies between the curves F and G, the $SU(3)_C \times SU(2)_L \times U(1)$ vacuum is the sole minimum of $V(\phi)$ *).

In order to discuss the evolution of the $SU(5)$ gauge system in the very early Universe, we need to study the temperature dependent effective potential $V(\phi, T)$. For $T \sim \mu/g$, $V(\phi, T)$ can be obtained^{3), 14)} from $V(\phi)$ by replacing μ^2 with $\mu^2 - \sigma T^2$ **, where⁴⁾

*) To the left of G there appears a local minimum that also corresponds to $SU(3)_C \times SU(2)_L \times U(1)$ but this region will not concern us here.

**) The finite temperature correction to the ϕ^3 term of $V(\phi)$ is $\sim g^3 T \phi^3$. For $T \sim \mu/g$, this correction is $\sim g^2 \mu \phi^3$ and thus has been neglected.

$$\sigma = \frac{1}{60} (130a + 94b + 75g^2) \quad .$$

From Fig. 1 we see that shortly after the hot big bang, the Universe is in the SU(5) symmetric phase. As the Universe expands and cools, $V(\phi, T)$ develops a local minimum corresponding to the $SU(4) \times U(1)$ phase which, to start with, is thermodynamically unfavourable compared to the SU(5) phase^{*}). However, the $SU(4) \times U(1)$ phase quickly becomes the absolute minimum of $V(\phi, T)$. After further cooling, the SU(5) phase is not even a local minimum of the potential. The system undergoes a rapid transition at $T_{c1} \approx \mu/\sqrt{\sigma}$ to the $SU(4) \times U(1)$ phase.

As the Universe cools further in the $SU(4) \times U(1)$ phase, the effective potential develops a local minimum corresponding to the $SU(3)_C \times SU(2)_L \times U(1)$ phase which becomes degenerate with the former phase at some critical temperature T_C . After more cooling, the $SU(3)_C \times SU(2)_L \times U(1)$ phase has a lower free energy than the $SU(4) \times U(1)$ phase. For reasonable choice of Higgs parameters, a large barrier separates the two minima. The transition from the false ($SU(4) \times U(1)$) phase to the true ($SU(3)_C \times SU(2)_L \times U(1)$) phase proceeds through barrier tunnelling (bubble formation) which, at least to start with, is a very slow process^{**}). The Universe will therefore supercool in the false ($SU(4) \times U(1)$) phase until a temperature T_{c2} is approached where the tunnelling process becomes large. Indeed, at temperature T_{c2} , the false vacuum is not even a local minimum of the potential and the Universe falls almost instantaneously into the $SU(3)_C \times SU(2)_L \times U(1)$ phase. In order to make sure that this scenario can be realized we only need to ensure that the zero temperature point $(\xi, \eta)_{T=0}$ is chosen to lie to the left of curve F. If, however, $(\xi, \eta)_{T=0}$ lies between the curves E and F, the false ($SU(4) \times U(1)$) vacuum remains a local minimum even at $T = 0$. In that case the difficulty discussed by Sato¹⁰⁾ that the phase transition from the false vacuum (in our case $SU(4) \times U(1)$) to the true ($SU(3)_C \times SU(2)_L \times U(1)$) vacuum may never be completed is likely to be present.

To summarize this part, we propose that the fate of the false ($SU(4) \times U(1)$) vacuum in the SU(5) gauge system is that it is not even a local minimum below some temperature T_{c2} . It exists as an absolute or a local minimum only between T'_{c1} ($> T_{c1}$) and T_{c2} . At T_{c2} the Universe undergoes a sudden transition to the true ($SU(3)_C \times SU(2)_L \times U(1)$) phase. The latent heat released in this transition reheats the Universe to a temperature $T_R > T_{c2}$.

^{*}) We assume from now on that $\eta > -1/5$.

^{**}) The probability of bubble formation is suppressed due to the factor $\exp(-16\pi a^3/3\epsilon^2 T)$, where a is the surface energy density of the bubble wall and ϵ the free energy density difference of the false and the true vacuum.

In order to make the above discussion more quantitative, let us be more specific and consider a typical case $a = b = g^2$ ($\eta = 1$) and $\xi(T=0) = -2.13$. The zero temperature point $(\xi, \eta)_{T=0}$ then lies to the left of curve F (Fig. 1). Taking $\alpha_G = g^2/4\pi = 1/45$ and $M_X(T=0) = 6 \times 10^{14}$ GeV, one finds

$$\begin{aligned} T_{c1} &= \frac{\mu}{\sqrt{\sigma}} \simeq 10^{15} \text{ GeV}, \\ T_c &\simeq 8 \times 10^{14} \text{ GeV}, \\ T_{c2} &\simeq 5 \times 10^{13} \text{ GeV}. \end{aligned} \tag{15}$$

The temperature T_R that is reached in the true $(SU(3)_C \times SU(2)_L \times U(1))$ vacuum after the release of latent heat and subsequent thermalization can be obtained from energy conservation considerations. One finds

$$\frac{8}{3} T_c^4 + (f+8) T_{c2}^4 = f T_R^4, \tag{16}$$

where $f = 106.75$. This gives

$$T_R \simeq 3 \times 10^{14} \text{ GeV}. \tag{17}$$

We see in the above example that $T_{c2} = (1/16)T_c$. Thus, the Universe does not supercool in the false $(SU(4) \times U(1))$ phase far below the equilibrium transition temperature T_c . However, a much larger amount of supercooling can be achieved if the point $(\xi, \eta)_{T=0}$ is chosen to lie very close to curve F. Even in this case, the temperature T_R is practically the same as in (17).

It is interesting to discuss the expansion of the Universe during the phase transition from $SU(4) \times U(1)$ to $SU(3)_C \times SU(2)_L \times U(1)$. For temperatures $T \lesssim T_e \simeq 0.4 T_c$, the vacuum energy density of the $SU(4) \times U(1)$ phase dominates the radiation energy density. Between T_e and T_{c2} , the Universe therefore undergoes an exponential expansion¹⁰⁾. In the example discussed above, one finds that the transition from the $SU(5)$ to the $SU(4) \times U(1)$ phase takes place at a cosmic time $t_{c1} \simeq 2.5 \times 10^{-37}$ sec. The transition to the $SU(3)_C \times SU(2)_L \times U(1)$ phase occurs at $t_{c2} \simeq 10^{-35}$ sec. Between $t_e \simeq 2.3 \times 10^{-36}$ sec and t_{c2} , the Universe expands exponentially with time. The cosmic scale factor $R(t_{c2})$ is found to be three times larger than it otherwise would be if there was no exponential expansion. Clearly, if $t_{c2} \gg t_e$ (and this can be arranged in the $SU(5)$ model), then $R(t_{c2})$ can be even larger.

SU(5) MAGNETIC MONOPOLES IN THE VERY EARLY UNIVERSE

We have seen above that at a temperature T_{c2} , the $SU(4) \times U(1)$ vacuum becomes locally unstable and decays almost instantaneously into the $SU(3)_C \times SU(2)_L \times U(1)$ vacuum. Magnetic monopoles carrying $SU(3)_C$, $SU(2)_L$ and $U(1)$ magnetic fields can then be produced. To estimate the number density of these monopoles, we note that the rate at which the Higgs field tends to its $SU(3)_C \times SU(2)_L \times U(1)$ symmetric value is much bigger than the expansion rate of the Universe [cf. Ref. 15)]. Thermal equilibrium is therefore established soon after the phase transition. The relative initial monopole density is then given by^{*)}

$$r(T_R) \equiv \frac{n(T_R)}{T_R^3} = \left(\frac{m(T_R)}{2\pi T_R} \right)^{\frac{3}{2}} \exp\left(-\frac{m(T_R)}{T_R}\right), \quad (18)$$

where $n(T_R)$ is the number density of monopoles at T_R and $m(T_R)$ is the monopole mass at T_R .

Estimating $r(T_R)$ in the numerical example above we obtain

$$r(T_R) \simeq 3 \times 10^{-36}. \quad (19)$$

We see therefore that the estimated relative monopole density is much smaller than 10^{-19} , the allowed upper bound^{**) 7)}. A first order phase transition can indeed suppress monopole production, as was first suggested in the first paper in Ref. 7) and in Ref. 4). There is no conflict between $SU(5)$ gauge theory and the hot big bang cosmology.

Finally, we note that since the temperature T_R is of order M_X , the lepto-quark Higgs and gauge bosons, are copiously produced at T_R . Baryon asymmetry in the Universe can then be generated following the usual scenario⁹⁾.

PHASE TRANSITION FROM $SO(10)$ TO $SU(4)_C \times SU(2)_L \times SU(2)_R$

It was recently shown⁵⁾ that the breaking of $SO(10)$ down to $SU(3)_C \times SU(2)_L \times U(1)$ via $SU(4)_C \times SU(2)_L \times SU(2)_R$ necessarily involves a first order phase transition. The fate of the false vacuum in the very early Universe can be discussed for $SO(10)$ similarly to the $SU(5)$ case.

^{*)} We assume that the probability of converting an $SU(4) \times U(1)$ monopole into an $SU(3)_C \times SU(2)_L \times U(1)$ monopole is negligible. See Ref. 4).

^{**) Note that, for $a = b = g^2$ ($\eta = 1$), we must have $-2.8 \leq \xi < -2.125$, in order that $r(T_R) \leq 10^{-19}$. The allowed range for ξ is not unnaturally small.}

By employing a 54-dimensional Higgs field ϕ , $SO(10)$ can break down to $SO(6) \times SO(4)$ ^{5) *}. At finite temperatures the breaking of $SO(10)$ proceeds as follows:

$$SO(10) \rightarrow SO(9) \rightarrow SO(8) \times U(1) \rightarrow SO(7) \times SO(3) \rightarrow SO(6) \times SO(4).$$

To simplify the discussion, let us concentrate on the transition from the $SO(7) \times SO(3)$ phase to the $SO(6) \times SO(4)$ phase. An analysis similar to the previous discussion shows that the $SO(7) \times SO(3)$ phase can exist either as a local or an absolute minimum for the following range of parameters ^{**)}:

$$\eta \geq -\frac{1}{7}, \quad -\frac{105\eta + 33}{2} < \xi < -\frac{105\eta - 7}{242},$$

and

$$-\frac{1}{7} \geq \eta > -\frac{1}{5}, \quad -\frac{105\eta + 33}{2} < \xi < \frac{2}{105\eta + 37}, \quad (20)$$

i.e., between curves A and C in Fig. 2. If $(\xi, \eta)_{T=0}$ is chosen to lie to the left of curve C (Fig. 2), there exists a temperature $T_{c2} < T_c$ (the temperature at which the $SO(7) \times SO(3)$ and $SO(6) \times SO(4)$ phases are degenerate) at which the $SO(7) \times SO(3)$ phase decays almost instantaneously into the $SO(6) \times SO(4)$ phase.

Consider a typical example with $a = b = g^2$ ($\eta = 1$) and $\xi(T=0) = -70$. One finds that

$$T_{c2} \simeq 0.12 T_c \quad \text{and} \quad T_R \simeq 0.35 T_c. \quad (21)$$

The relative monopole density is estimated to be

$$r(T_R) \simeq 7 \times 10^{-38} \ll 10^{-19}. \quad (22)$$

Clearly, there is no conflict between $SO(10)$ and the hot big bang cosmology.

^{*}) The most general renormalizable Higgs potential for ϕ has the same form as (2), with ϕ now regarded as a symmetric, traceless 10×10 matrix.

^{**}) The parameters ξ and η are defined as in the $SU(5)$ case.

Note once again that if $(\xi, n)_{T=0}$ is chosen to lie very close to curve C (Fig. 2), then $T_{C2} \ll T_C$. However, T_R remains essentially unchanged.

To summarize, we have shown that in an expanding Universe, a first order phase transition from a false to a true vacuum can take place in a finite amount of time. In particular, for SU(5) and SO(10) gauge models, there exists a reasonable choice of parameters of the Higgs potential for which this transition is completed. The Universe supercools in the false vacuum until some temperature T_{C2} at which the false vacuum becomes locally unstable and decays almost instantaneously into the true vacuum^{*)}. We also find that the density of monopoles produced in such a phase transition is strongly suppressed. There is no conflict between unified gauge theories and the hot big bang model of cosmology.

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^{*)} Symmetry breaking of SU(5) (or SO(10)) down to $SU(3)_C \times SU(2)_L \times U(1)$ can also be achieved by the Coleman-Weinberg mechanism. In this case, however, the Universe is likely to remain in the false (SU(5) or SO(10) symmetric) phase till $T = 0$.

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FIGURE CAPTIONS

Fig. 1 : Phase diagram for the spontaneous breaking of $SU(5)$ gauge symmetry induced by the vacuum expectation value of a Higgs 24-plet. Between curves B and E, the $SU(4) \times U(1)$ phase is the absolute minimum of the Higgs potential. To the left of D there appears a local minimum corresponding to the $SU(3)_C \times SU(2)_L \times U(1)$ phase which becomes the absolute minimum to the left of curve E. To the left of F, the $SU(4) \times U(1)$ phase ceases to exist even as a local minimum of the Higgs potential. The cross hatched region ($\eta \leq -7/15$) is forbidden.

Fig. 2 : Partial phase diagram for the spontaneous breaking of $SO(10)$ gauge symmetry induced by the vacuum expectation value of a Higgs 54-plet. Between curves A and C, the $SO(7) \times SO(3)$ exists either as an absolute or a local minimum of the Higgs potential. To the left of curve B, the $SO(6) \times SO(4)$ phase is the absolute minimum of the potential. To the left of C, the $SO(7) \times SO(3)$ phase ceases to exist even as a local minimum of the potential. The cross hatched region ($\eta \leq -1/5$) is forbidden.

fig.1

fig.2