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LOCAL SUSY BREAKING INDUCED BY THE GUT PHASE TRANSITION
IN THE EARLY UNIVERSE

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A B S T R A C T

We show that the confining properties of the grand unification theory (GUT) may induce a transition to a ground state where both local SUSY and gauge symmetry are broken. The super-Higgs field belongs to the adjoint representation which breaks gauge symmetry taking a VEV of the order of the grand unification scale M_X . A suitable class of superpotentials is proposed. We show that the critical temperature of the phase transition is 10^{10} GeV. We get an approximate relation between M_X and the gravitino mass :

$$m_{3/2} M \simeq M_X^2 \exp(-4\pi/b_{\text{SU}(5)} \alpha_{\text{GUT}})$$

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Finite temperature effects and the GUT phase transition in the early universe have been recently analyzed in the context of $SU(5)$ global SUSY GUTs¹⁾⁻⁵⁾. When the universe cools down the vacuum remains in the symmetric phase until the confinement temperature (energy) is reached. Then the freezing of the degrees of freedom in this phase reduces the depth of the symmetric minimum which becomes a false vacuum. If the height of the barrier separating the two phases is smaller than the energy of thermal fluctuations the transition to the broken phase occurs.

The energy scales involved in this phenomenon are sufficiently high for us to be concerned about the possible effects induced by gravity. On the other hand the picture is not related to supersymmetry breaking. The simplest way to handle both problems is to study phase transitions at finite temperature in the context of $N = 1$ supergravity⁶⁾. It is well known that the super-Higgs phenomenon in $N = 1$ supergravity⁷⁾ induces a soft breaking of global SUSY in agreement with low energy phenomenology. In most semi-realistic models, recently proposed in the literature⁸⁾⁻¹⁵⁾, the super-Higgs mechanism is supposed to be driven by a singlet field which is completely decoupled in the flat limit (hidden sector) and takes a VEV of the order of the Planck mass M_P ⁷⁾. In this picture the hidden singlet is put in by hand just to break local SUSY but without any relation with the dynamics of the expanding early universe.

The main phenomena we know should occur in the very early universe are inflation and the GUT phase transition, hence the question arises: Is the breaking of local supersymmetry a consequence of any of these phenomena? Recently, primordial cosmological inflation has been analyzed in the context of $N = 1$ supergravity by Nanopoulos, Olive, Srednicki and Tamvakis¹⁶⁾ showing that cosmological inflation at the Planck time, with suitable scale independent density fluctuations, can be driven by a scalar field which does not break local supersymmetry. On the other hand, it has been recently proved^{17),18)} that a scalar field breaking local supersymmetry is unable to produce primordial cosmological inflation and create baryon number asymmetry. Therefore, we are led to investigate the relationship between local supersymmetry breaking and the GUT phase transition in the early universe. In this letter we want to present the scenario which emerges from such an analysis: It is the simultaneous breaking of local supersymmetry and gauge symmetry at M_X which actually occurs in the early universe. Stated in other words, when the universe cools down below the GUT confinement temperature, the symmetric phase confines and the broken phases become the true minima, but only the transition from the symmetric phase to the broken SUSY-broken gauge symmetry phase has a chance to occur. We derive the general conditions the superpotential must satisfy. Specific superpotentials are also given. Primordial cosmological inflation is

produced by the singlet field of Ref. 16)-18) which does not break local SUSY while creating baryon number asymmetry. The production of monopoles is suppressed at the critical temperature of the SU(5) phase transition, as in globally supersymmetric models.

Contrary to previously proposed scenarios we will not introduce any singlet field to break local SUSY. Instead, we will show that it is possible to break both local SUSY and¹⁹⁾ SU(5) gauge symmetry when the adjoint (24) representation Σ takes a non-vanishing VEV of the order of M_x , where $M_x \simeq 10^{16}$ GeV is the usual grand unification scale.

The relevant scalar potential, in $N = 1$ supergravity with minimal coupling, corresponding to the Σ field is given by²⁰⁾

$$V(\Sigma) = e^{\text{tr} \Sigma^2} \left\{ |F_\Sigma|^2 - 3 |f|^2 \right\} \quad (1)$$

where

$$F_\Sigma = \frac{\partial f}{\partial \Sigma} - \frac{1}{5} \text{tr} \frac{\partial f}{\partial \Sigma} + f \Sigma^* \quad (2)$$

and $f(\Sigma)$ is a general non-renormalizable superpotential (we are working in units of $M = M_P / \sqrt{8\pi} = 2.4 \cdot 10^{18}$ GeV).

Let us classify the minima which may appear in the potential (1) :

- i) The minimum v_0 conserving local SUSY and gauge symmetry : $\Sigma_0 = 0$, $v_0 \leq 0$.
- ii) The minima v_1 breaking both local SUSY and gauge symmetry : $\Sigma_1 \neq 0$, $F_{\Sigma_1} \neq 0$.
- iii) The minima v_2 conserving local SUSY and breaking gauge symmetry : $\Sigma_2 \neq 0$, $F_{\Sigma_2} \neq 0$, $v_2 \leq 0$.

The VEVs Σ_1 and Σ_2 can produce a variety of breakings of SU(5). In what follows we shall restrict to the breaking of SU(5) into SU(3) $\times$ SU(2) $\times$ U(1). We shall show later on that the phase SU(4) $\times$ U(1) is harmless for our purpose.

Next, we shall give an overall picture of the scenario we are proposing in this paper. First of all we shall assume two different scales for the minima breaking and conserving local SUSY. The minimum Σ_1 breaking local SUSY should

occur at a value of the order of the grand unification scale, M_X . Some locally supersymmetric minimum may (or not) occur at a scale of the order of the Planck mass M . In general we are not allowed to choose Σ_1 as the ground state of the theory since it is between two other minima Σ_0 and Σ_2 with lower energies. The physical situation corresponds to the case where the three types of minima are degenerate in energy, $v_0 = v_1 = v_2 = 0$, (Fig. 1) as we shall see later. This corresponds to a particular class of superpotentials to be specified below. It turns out that the height of the barrier separating Σ_0 and Σ_1 is $a \approx (m_{3/2} M)^2$ while the second barrier, separating Σ_1 and Σ_2 is many orders of magnitude larger (see later). This guarantees that the tunnelling from Σ_1 into Σ_2 is enormously suppressed at the critical temperature $T_c \approx a^{1/4}$ for the phase transition from Σ_0 to Σ_1 . The simplest situation, where only Σ_0 and Σ_1 exist (but not Σ_2) is depicted in Fig. 2.

Let us now consider the evolution of the ground state as the universe cools down. The effective potential at high temperature is given by^{1),2),21)}

$$V(\Sigma, T) \simeq V(\Sigma, 0) + \frac{T^2}{24} \sum_f t_f m_f^2 - \frac{\pi^2}{90} N(T) T^4 \quad (3)$$

where $N(T) = N_B + 7/8 N_F$ is the number of light degrees of freedom at the temperature T . At high temperatures the T^4 term dominates (3) and we need to estimate the value of $N(T)$ for the different phases involved. A simple calculation gives $N_B = 188$ for the $SU(5)$ phase, $N_B = 156$ for the $SU(4) \times U(1)$ phase and $N_B = 140 + 8N_H + 4N_{10}$ for the $SU(3) \times SU(2) \times U(1)$ phase, where N_H is the number of pairs of light complex doublets and N_{10} is the number of pairs of complex Higgs multiplets in the $\underline{10} + \overline{\underline{10}}$ representation of $SU(5)$. To generate a non-vanishing baryon asymmetry we must consider $N_H = 2$ with "light" triplets at an intermediate scale (10^{10} GeV)¹⁾ while $N_{10} = 3$ in order to get reasonable values of $\sin^2 \theta_w$ and M_X . In this way $N(SU(5)) > N(SU(3) \times SU(2) \times U(1)) > N(SU(4) \times U(1))$ at temperatures of the order of 10^{10} GeV, therefore the ground state is confined in the $SU(5)$ symmetric phase. However, at the critical temperature $T_c \approx O(\Lambda_{SU(5)})$ the $SU(5)$ theory becomes confining. The main effect of confinement is to reduce the number of light degrees of freedom and shift the symmetric minimum upwards^{1),2)} (Fig. 3). The transition to the broken vacuum becomes possible and proceeds at a reasonable rate provided that the height of the barrier is not much larger than the thermal fluctuations.

The transition temperature T_c is of the order of the confining scale of SU(5)

$$T_c \simeq \Lambda_{su(5)} = M_x \exp \left\{ -\frac{2\pi}{b_5 \alpha_{GUT}} \right\} \quad (4)$$

which implies a non-trivial relation between the scales $m_{3/2}$ and M_x

$$m_{3/2} M \simeq M_x^2 \exp \left\{ -\frac{4\pi}{b_5 \alpha_{GUT}} \right\} \quad (5)$$

For the particle content considered above⁵⁾ the values of α_{GUT} and b_5 are given by $\alpha_{GUT}^{-1} = 5.8$, $b_5 = 3$, so one has

$$\log_{10} M_x \simeq \frac{1}{2} \log_{10} m_{3/2} + 15 \quad (6)$$

Notice that Eq.(5) is equivalent to a relation between $m_{3/2}$ and the QCD scale, Λ_{QCD} , defined at m_w

$$m_{3/2} M \simeq M_x^{2(1 - \frac{b_3}{b_5})} \Lambda_{QCD}^{\frac{2b_3}{b_5}} \quad (7)$$

For a GUT scale $M_x \simeq 10^{16}$ GeV, one derives from Eq.(6) $m_{3/2} \simeq 10^2$ GeV.

Let us now discuss the conditions the superpotential should satisfy for the validity of our general picture. We take a general non-renormalizable superpotential as

$$f(\Sigma) = \sum_{n=2}^{\infty} a_n \text{tr} \Sigma^n \quad (8)$$

and choose the direction which breaks SU(5) into SU(3) $\times$ SU(2) $\times$ U(1)

$$\Sigma = \omega \sqrt{\frac{2}{15}} \text{diag} \left\{ 1, 1, 1, -\frac{3}{2}, -\frac{3}{2} \right\} \quad (9)$$

In this case the scalar potential (1) can be written as a function of a simple complex field ω

$$V = e^{|\omega|^2} \left\{ |F|^2 - 3|f|^2 \right\} \quad (10)$$

where

$$F = \sum_{n=2}^{\infty} c_n \left\{ n \omega^{n-1} + \omega^n \omega^* \right\} \quad (11)$$

and
$$f = \sum_{n=2}^{\infty} c_n \omega^n \quad (12)$$

The coefficients c_n and a_n are simply related by

$$c_n = a_n 3 \left(\frac{2}{15} \right)^{\frac{n}{2}} \left\{ 1 - \left(-\frac{3}{2} \right)^{n-1} \right\}, \quad (c_1 = 0) \quad (13)$$

The condition to have a minimum in the complex plane at $\omega = 0$ (Σ_0 in Fig. 1) is $c_0 = 0$, $c_2 \neq 0$, and the singlet ω takes a mass $m_0 = 2|c_2|$.

As we mentioned before, supersymmetric minima may (or may not) appear at large scales $\omega \approx 1$. In order to describe the most general situation, we shall assume in the following that the potential has a minimum at $\omega = 1$.

The condition to have a vanishing cosmological constant at $\omega = 1$ (Σ_2 in Fig. 1) is

$$\sum_{n=2}^{\infty} c_n = \sum_{n=2}^{\infty} n c_n = 0 \quad (14)$$

which implies that the potential is stationary at that point. On the other hand, $\omega = 1$ is a minimum in the complex plane provided that

$$m_2 = e^{1/2} \left| \sum_{n=2}^{\infty} n^2 c_n \right| \neq 0 \quad (15)$$

Finally, the conditions for the physical minimum, at $\omega = x$, which breaks local SUSY, are :

$$V(x) = 0 \implies \sum_{n=2}^{\infty} n c_n x^{n-1} = (\sqrt{3} - x) e^{-\frac{x^2}{2}} m_{3/2} \quad (16)$$

$$V'(x) = 0 \implies \sum_{n=2}^{\infty} n(n-1) c_n x^{n-2} = \{(\sqrt{3} - x)^2 - 1\} e^{-\frac{x^2}{2}} m_{3/2} \quad (17)$$

and the masses of $\omega = A + iB$ are

$$\begin{aligned} m_A^2 &= e^{x^2/4} \sqrt{3} \sum_n c_n n(n-1)(n-2) x^{n-2} m_{3/2} - \sqrt{3} x(x-\sqrt{3})(x-2\sqrt{3}) m_{3/2}^2 > 0 \\ m_B^2 &= 4 m_{3/2}^2 - m_A^2 > 0 \end{aligned} \quad (18)$$

where

$$m_{3/2} = |f(x)| e^{x^2/2} \quad (19)$$

Just to show how the above scenario can be accomplished in practice, we want to present now a particular realization of it. The equations for getting a scalar potential, as given in Fig. 1, Eqs.(13),(14),(16) and (17), can be written as

$$\begin{aligned} f(0) = f'(0) = f(1) = f'(1) = 0 \\ f'(x) = (\sqrt{3}-x)f(x) \quad , \quad f''(x) = \{(\sqrt{3}-x)^2 - 1\} f(x) \end{aligned} \quad (20)$$

A superpotential satisfying Eq.(20) is given by

$$\begin{aligned} f(\omega) = m_{3/2} \frac{\omega^2(\omega-1)^2}{x^2(x-1)^2} \left\{ 1 + d_1(\omega-x) + \frac{1}{2} d_2(\omega-x)^2 + \right. \\ \left. + \frac{1}{6} d_3(\omega-x)^3 + \frac{1}{24} d_4(\omega-x)^4 \Phi(\omega) \right\} \end{aligned} \quad (21)$$

where $\Phi(\omega)$ is a regular function, chosen in such a way as to avoid supersymmetric minima at small scales

$$\begin{aligned} d_1 = \tau_1 \quad , \quad \tau_1 = -\frac{2}{x} - \frac{2}{x-1} - x + \sqrt{3} \\ d_2 = \tau_2 + \tau_1^2 \quad , \quad \tau_2 = \frac{2}{x^2} + \frac{2}{(x-1)^2} - 1 \\ d_3 = \tau_3 + 3\tau_2\tau_1 + \tau_1^3 \quad , \quad \tau_3 = -\frac{4}{x^3} - \frac{4}{(x-1)^3} + \alpha \end{aligned} \quad (22)$$

α being a constant which gives the mass splitting between m_A and m_B .

The coefficients c_n of the power expansion in ω given in Eq.(12) can be easily derived from Eq.(21). In the minimal case when $\Phi(\omega)$ is chosen to be the unity, one has

$$c_n \simeq 10^{2(n-2)} m_{3/2} \simeq 10^{2(n-10)} \quad , \quad n=2, \dots, 8. \quad (23)$$

The corresponding coefficients a_n of the $\text{tr} \Sigma^n$ expansion of the superpotential, Eq.(8), are trivially related to c_n of Eq.(23) by Eq.(13).

Using Eq.(21) we get the mass of the singlet field which breaks local supersymmetry

$$\begin{aligned} m_A^2 &= \sqrt{3} \alpha m_{3/2}^2 \\ m_B^2 &= (4 - \sqrt{3} \alpha) m_{3/2}^2 \end{aligned} \quad (24)$$

so that $(4/\sqrt{3}) > \alpha > 0$ is the positive (mass)² condition.

On the other hand, we can derive for m_0 and m_1 the mass relations

$$\begin{aligned}\frac{m_0}{m_{3/2}} &\simeq \left(\frac{M}{M_x}\right)^2 \\ \frac{m_1}{m_{3/2}} &\simeq \left(\frac{M}{M_x}\right)^{5+n}\end{aligned}\quad (25)$$

where n is a positive integer, $n > 1$, depending on the particular choice of $\phi(\omega)$.

The height of barriers can be easily understood from Eq. (25). For the first barrier in Fig. 1 we get

$$a \simeq (m_0 m_{3/2}) M_x^2 \simeq (m_{3/2} M)^2 \quad (26)$$

as stated before. On the other hand, for the second barrier we have

$$b \simeq (m_0 m_1) M^2 \simeq a \left(\frac{M}{M_x}\right)^{7+n} \quad (27)$$

so that the ratio $b/a \geq 10^{16}$ which corresponds to the minimal case $n = 1$.

Finally, some comments about the solution of standard cosmological problems in our model are in order. Indeed, there is room for cosmological inflation in our picture. The critical temperature at which the $SU(5)$ phase transition occurs is clearly below the temperature at which inflation would be produced. Moreover, primordial inflation has been recently proposed⁵⁾ as a mechanism suitable to produce enough inflation in the early universe so as to solve standard cosmology problems. This mechanism also works in $N = 1$ supergravity via a singlet field which suffers a phase transition from the de Sitter phase¹⁶⁾ at the Planck temperature and does not break local SUSY. It is clear that this mechanism also works in our picture, producing primordial cosmological inflation. Notice that monopole production is sufficiently suppressed in our model. The critical temperature for the $SU(5)$ phase transition is low enough, 10^{10} GeV, and one obtains⁵⁾ $n_m/n_\gamma \sim 10^{-24}$, which is in agreement with the experimental upper bound.

Notice that, in our model, the mass of the singlet coming from the adjoint representation of $SU(5)$, Eq.(24), is of the order of the gravitino mass. One might think this could destabilize²³⁾ the masses of the low energy theory through the coupling $\lambda H_5 \Sigma H_5$. However, in our case $\lambda \langle \Sigma \rangle$ is of the order of 10^{10} GeV (the mass of "light" triplets) so that no destabilization mechanism⁶⁾ occurs in our model.

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FIGURE CAPTIONS

Fig. 1 Zero temperature potential corresponding to the superpotential [Eq.(21)] with the minimal choice $\phi(\omega) = 1$. The three minima $\Sigma_0 = 0$, $\Sigma_1 \simeq M_x$ and $\Sigma_2 \simeq M$ correspond to the $SU(5)$ supersymmetric, $SU(3) \times SU(2) \times U(1)$ non-supersymmetric and $SU(3) \times SU(2) \times U(1)$ supersymmetric phases, respectively.

Fig. 2 The simplest zero temperature potential exhibiting minima only at Σ_0 and Σ_1 , but not at Σ_2 . It corresponds to the superpotential of Eqs.(21), (22) where terms in $(\omega - 1)$ and $(x - 1)$ have been suppressed.

Fig. 3 The same as in Fig. 2, at the critical temperature $T_c \simeq \Lambda_{SU(5)} \simeq 10^{10}$ GeV.

Fig. 1

Fig. 2

Fig. 3