PRIMORDIAL TWO-COMPONENT MAXIMALLY SYMMETRIC INFLATION

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ABSTRACT

We propose a two-component inflation model, based on maximally symmetric supergravity, where the scales of reheating and the inflaton potential at the origin are decoupled. This is possible because of the second-order phase transition from $SU(5)$ to $SU(3) \times SU(2) \times U(1)$ that takes place when $\phi \approx \phi_c < \phi_0$, when $\phi_0 \approx 0(M)$ is the value of the inflaton at the global minimum, and leads to a reheating temperature $T_R \approx (10^{15} - 10^{16})$ GeV. This makes it possible to generate baryon asymmetry in the conventional way without any conflict with experimental data on proton lifetime. The mass of the gravitinos is $m_{3/2} \approx 10^{12}$ GeV, thus avoiding the gravitino problem. Monopoles are diluted by residual inflation in the broken phase below the cosmological bounds if $\phi_c \lesssim 0.3M$.

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1. INTRODUCTION

The inflationary Big Bang cosmological model [1,2] is at present the most appealing solution to such standard cosmological puzzles as the curvature and the horizon problem. Furthermore, the existence of an inflationary epoch predicts the scale-independent density perturbations [3] needed to explain the galaxy formation [4]. The incapability of inflationary models based on conventional field theories to explain the scale of energy density perturbations ($\delta\rho/\rho \simeq 10^{-4}$), and accordingly the isotropy of the microwave background radiation ($\delta T/T \simeq \delta\rho/\rho$) led people to consider inflationary models based on supersymmetric theories [5,6].

In the framework of supergravity, the flatness of the inflationary potential, which is needed to generate such small density perturbations, remains stable under quantum corrections. However, even in supergravity the necessary final step from the inflationary era to the $SU(3) \times SU(2) \times U(1)$ symmetric Universe has proven to be difficult. Firstly, one has to effect somehow the phase transition breaking the grand unified symmetry and secondly, one has to produce the observed baryon asymmetry of the Universe during reheating. Because the overall scale of the inflaton potential usually dictates the reheating temperature T_R to be very low, $T_R \leq T_H$, one will run into trouble with the $d = 5$ operators capable of mediating catastrophically rapid proton decay.

Recently, a new mechanism has been proposed [7,8] to solve both these problems. There the inflation causes the symmetry-conserving high-temperature minimum of adjoint Higgs field to become unstable during the period of inflation, and a second-order phase transition takes place. In fact, then the period of inflation, produced by the inflaton field ϕ , can be decoupled from the period of reheating in the sense that the scales governing these two periods are completely different. In this two-component inflation the

reheating temperature can be as high as $T_R \simeq V(\varphi = 0)^{1/4} \simeq (10^{15} - 10^{16}) \text{ GeV}$. Reheating is good because the field responsible for it is a mixture of the inflaton and the adjoint Higgs and couples directly to matter. Baryon asymmetry will then be generated through heavy Higgs triplets ($\simeq 10^{16} \text{ GeV}$) in the conventional way [9,10], while $d = 5$ operators predict a proton lifetime around or larger than the present experimental bounds.

In the present paper we propose a supergravity model of inflation that: produces sufficient inflation; gives rise to correct energy density perturbations; has a second-order phase transition to the $SU(3) \times SU(2) \times U(1)$ phase during inflation; has good reheating; generates baryon asymmetry in the conventional way without contradicting data on proton lifetime; can suppress the density of monopoles even below the very stringent limit coming from the monopole-catalysed baryon decay in neutron stars; avoids the cosmological problems with light gravitinos; and allows for a dynamical determination of the electroweak scale. The model is based on the maximally symmetric Einstein-Kähler manifold [11-13] of the scalar fields.

2. MAXIMALLY SYMMETRIC NO-SCALE SUPERGRAVITY AND INFLATION

Two different problems find a common solution in the context of maximally symmetric no-scale supergravity models [13]: the inflation scale problem and the gravitino problem. In particular, the requirement of energy density perturbations consistent with galaxy formation and also with the observed isotropy of the background radiation implies a vacuum energy at the metastable minimum of the order of $V(0)^{1/4} \simeq (10^{15} - 10^{16}) \text{ GeV}$, i.e. a Hubble constant $H \simeq (10^{12} - 10^{13}) \text{ GeV}$. On the other hand, in supergravity any De Sitter vacuum breaks supersymmetry and defines a gravitino mass term equal to $V(0)^{1/2}/M$. This can easily be seen from the general structure of the supergravity potential [14]

$$V = e^G \left[G_i (G^{-1})^i_j G^j - 3 \right] + \frac{1}{2} D^2. \quad (1)$$

Now, if a minimum has $V(0) \geq 0$, with $D(0) = 0$, we conclude that

$$V(0) \leq e^G G_i (G^{-1})^i_j G^j \simeq e^G = m_{3/2}^2 M^2, \quad (2)$$

where we assume that $G_i (G^{-1})^i_j G^j \simeq 0(1)$, i.e. that the coupling constant are of the order of unity. This simple observation shows that the gravitino mass term at the local minimum is fixed by the correct scale of energy density fluctuations: in our particular case we have $m_{3/2}(0) \simeq (10^{12} - 10^{13})$ GeV. This large value of $m_{3/2}(0)$ is what we call the inflation scale problem.

At this stage we can proceed in two different ways:

i) To enforce [15] a phase transition of $m_{3/2}(\varphi)$ from $m_{3/2}(0) \simeq 10^{12}$ GeV to $m_{3/2}(\varphi_0) \simeq 10^2$ GeV, where φ_0 is the global minimum of the inflation potential. This phase transition needs some additional conditions to be imposed and, in particular, a fine-tuning on the parameters of the superpotential. Furthermore, such a light gravitino (inside the forbidden region [16] $1 \text{ keV} < m_{3/2} < 10^4 \text{ GeV}$) is sufficiently stable to live beyond the primordial nucleosynthesis epoch, but its mass density invalidates conventional calculations of the abundance of light elements. It was initially proposed [17] that an inflationary epoch could suppress sufficiently the primordial gravitino abundance in such a way that gravitinos produced after inflation, but before nucleosynthesis, would not lead to cosmological troubles. However, recent calculations [18] show that one can only tolerate a reheating temperature of the Universe after inflation $T_R < (10^8 - 10^9) \text{ GeV}$ since, otherwise, the regenerated gravitinos may change drastically the primordial production of light elements and distort the isotropy of cosmological background radiation. However, such a low reheating temperature is

unable to generate the observed cosmological baryon asymmetry of the Universe. This is the so-called gravitino problem.

ii) The way to solve simultaneously the inflation scale and the gravitino problems [19] is to let $m_{3/2}(\varphi_0) \simeq m_{3/2}(0) \simeq 10^{12} \text{ GeV} \gg 10^4 \text{ GeV}$, which is the most natural solution. There is no need for any phase transition for $m_{3/2}(\varphi)$ and the gravitino problem is non-existent. Note that such a large value for the gravitino mass term is only permitted when it does not appear, or is strongly suppressed, in the low-energy spectrum of the theory; otherwise it will destabilize the electroweak scale.

In particular, for all the supermultiplets (B_i, F_i) which couple "strongly" (i.e. through gauge interactions), electroweak stability requires

$$|M_{B_i}^2 - M_{F_i}^2| \leq \mathcal{O}(500 \text{ GeV})^2, \quad (3)$$

while for gauge singlet supermultiplets, such as the (inflaton-inflatino)-supermultiplet, which may couple gravitationally with conventional non-singlet fields, we only need

$$\eta^2 |M_B^2 - M_F^2| \leq \mathcal{O}(500 \text{ GeV})^2, \quad (4)$$

where η represents the strength of interaction. For instance, if $\eta \simeq 10^{-12}$, then $|M_B^2 - M_F^2|^2$ can be as large as $(10^{13} \text{ GeV})^2$ without destabilizing the gauge hierarchy. These features can only be accommodated in supergravity models where the super-Higgs mechanism takes place in a maximally symmetric way [13].

An inflationary model based on maximally symmetric supergravity has recently been proposed in Ref. [19]. In that model the reheating is produced by inflaton oscillations around its global minimum φ_0 . Since the

inflaton field is "weakly" coupled to matter, the reheating is poor so that the reheating temperature is $T_R \simeq (10^{10} - 10^{11})$ GeV and (light) triplets have to be around this scale to generate the cosmological baryon asymmetry. In that case the $d = 5$ operators for proton decay have to be rotated away by using a complicated Higgs sector in the theory.

In the present paper we show that such complications are, in fact, not necessary. In our model reheating is good and consequently there are no problems with baryon asymmetry or $d = 5$ operators.

3. THE INFLATON POTENTIAL

The Kähler potential of our model is chosen to be

$$G = G_0 - aK(\phi, \phi^+) + \log |F(\phi, y^a)|^2, \quad (5)$$

where

$$G_0 = -3 \log \left[z + z^+ - \frac{1}{3} y^a y_a^+ - K(\phi, \phi^+) \right]. \quad (6)$$

Here F is an analytic function (the superpotential) of the matter fields y^a and the gauge singlet inflaton field ϕ , $K(\phi, \phi^+)$ is a real function, a a positive constant, and z the field responsible for supersymmetry breaking. The scalar potential corresponding to Eq. (5) can be written as [19]

$$V = \frac{x^3}{3x-a} e^{-aK} \frac{|F_\phi - aK_\phi F|^2}{K_{\phi\phi^+}} + x^2 e^{-aK} |F_y|^2 + \frac{1}{2} x^2 D^2, \quad (7)$$

where $x = \exp(G_0/3)$, and D is the gauge part of the potential (D-term), which in our case takes the usual form $D = g_{\text{GUT}} y_a^+ T_b^{ab}$.

The existence of a metastable minimum with positive cosmological constant requires the positivity of a . In fact, since x is a function only of the properly normalized real field $A_D = G_0/\sqrt{6}$, the dilaton field, the potential (7) has to be minimized with respect to that direction. This minimization condition ($V_x = 0$) is automatically satisfied for any global minimum with zero cosmological constant for all values of A_D lying in the physical domain

$$\begin{aligned} 3x - a &> 0 \\ K_{\phi\phi^+} &> 0 \end{aligned} \quad (8)$$

However, in the metastable minimum case, $F_\psi - aK_\psi F \neq 0$, $V_x = 0$ implies $x = a/2$ and $V_y = 0$ implies $F_y = 0$ with $D = 0$, which can only be satisfied for $a > 0$. The simplest choice, which will be adopted hereafter, is $a = 2$ so that $x = 1$ ($A_D = 0$) and there is no necessity of rescaling matter fields during the inflationary epoch. Also, this would be the case for inflaton field if $K_{\phi\phi^+} \simeq 1$ during the inflation; thus we choose

$$K(\phi, \phi^+) = \phi\phi^+ - \frac{1}{4}\beta(\phi\phi^+)^2 \quad (9)$$

For $\beta > 0$, there is a domain of positive kinetic energy; at the boundary of this domain the potential becomes infinite so that the fields will indeed be confined in this domain. For $\beta \leq 0$ the domain of positive kinetic energy covers the whole of the complex plane. Although in what follows we will use a certain value of β as an example, we do not commit ourselves to either of the above possibilities.

In Refs. [15] and [19] it is shown that when $\langle y^a \rangle = 0$, all conditions of inflation and correct energy density perturbations ($\delta\rho/\rho \simeq 10^{-4}$) fix the inflaton superpotential to be

$$\log F(\phi, y^a=0) = \log m_0 + F(\phi), \quad (10)$$

$$F(\phi) = \xi\phi + \frac{1}{2}(2-\zeta^2)\phi^2 + \frac{1}{6}\left(2\zeta^3 + \frac{1}{\xi}\frac{m_R^2}{m_0^2} - \beta\zeta\right)\phi^3 + O(\phi^4),$$

where the parameters ξ and m_0 determine the value of the potential at $\phi = 0$, $V(0) = \xi^2 m_0^2 M^2$, and m_R^2 is the curvature of the inflaton potential along the real axis at $\phi = 0$.

Inflation then occurs within a large range of the parameters ξ , β , and m_0 . For example, taking $\xi = -1/\sqrt{10}$ and $\beta = 1/50$, as in Ref. [15], and dropping all $O(\phi^4)$ -terms from Eq. (10), we get from inflation and energy density fluctuations, $m_0 = 11 \times 10^{12}$ GeV, which corresponds to $H = 2.0 \times 10^{12}$ GeV and $V_0 = (2.9 \times 10^{15} \text{ GeV})^4$, and $m_R = 3.1 \times 10^{10}$ GeV. Of course, other values of ξ and β are possible, and we have explicitly checked that, for example, negative values of β will also lead to an acceptable inflationary model. Since the curvature of the inflaton potential along the imaginary direction is larger than the curvature along the real axis [10^6 times for $\phi = O(H)$ and 10^2 times near the global minimum], it is sufficient to consider the inflaton potential along the real direction ϕ . With our choice of parameters it takes the form

$$V(\phi) = v(\phi) m_0^2 M^2 \exp \left[2y - y^2 + (0.6 + \frac{1}{3}\epsilon)y^3 + y^4 \right], \quad (11)$$

where

$$y = \text{Re}(-\phi/\sqrt{10}), \quad \epsilon = 100 \left(\frac{m_R}{m_0} \right)^2 \simeq 8 \times 10^{-4}$$

and

$$v(\phi) = \frac{1}{10 - 2y^2} \left[1 - y + (0.9 + \frac{1}{2}\epsilon)y^2 + 2y^3 \right]^2 \quad (12)$$

The potential (11) exhibits, inside the physical domain $|\varphi| < 1/\sqrt{\beta} = 7.07M$ a unique global minimum (in the complex φ -plane) at $\varphi_0 = 3.817M = 0.92 \times 10^{19}$ GeV (see Fig. 1).

4. GUT PHASE TRANSITION AND REHEATING

In this section we will consider a mechanism whereby the grand unified symmetry will become broken during inflation. We will consider an SU(5) GUT for definiteness. A second-order phase transition from SU(5) to SU(3) $\times$ SU(2) $\times$ U(1) can naturally be accomplished with the following superpotential:

$$\log (F(\phi, \Sigma, y^a=0)) = F(\phi) + \log (m_0 + f(\phi, \Sigma)), \quad (13a)$$

where

$$f(\phi, \Sigma) = -\frac{1}{2} C_2(\phi) \text{Tr} \Sigma^2 + \frac{1}{3} C_3(\phi) \text{Tr} \Sigma^3 \quad (13b)$$

Here Σ is the adjoint Higgs (24) and the coefficients of the superpotential are generalized to be φ -dependent, as proposed in Ref. [8]. This is a very natural assumption, because φ is a singlet and can therefore have all kinds of (usually non-renormalizable) interactions with $\text{Tr} \Sigma^n$. We take now the SU(3) $\times$ SU(2) $\times$ U(1) invariant direction σ of Σ and the effective superpotential is

$$f(\phi, \sigma) = -\frac{1}{2} a_2(\phi) \sigma^2 + \frac{1}{3} a_3(\phi) \sigma^3, \quad (13c)$$

where the coefficients $a_i(\varphi)$ are trivially expressed in terms of $c_i(\varphi)$. All other fields can for our purposes be taken to lie in their minima $y^a = 0$.

It is worth stressing here that the functions $a_i(\varphi)$ in Eq. (13c) contain direct couplings between the φ and Σ sectors. In particular, if we expand $a_i(\varphi)$ around the global minimum $\tilde{\varphi} = \varphi - \varphi_0$, we can write in the superpotential (13c): $a_2(\varphi) \sigma^2 = a_2(\varphi_0) \sigma^2 + a'_2(\varphi_0) \tilde{\varphi} \sigma^2 + O(\tilde{\varphi}^2 \sigma^2)$ and $a_3(\varphi) \sigma^3 = a_3(\varphi_0) \sigma^3 + O(\tilde{\varphi} \sigma^3)$. This shows that $a'_2(\varphi_0)$ measures the strength of renormalizable interactions that couple φ to σ which is represented by the parameter η in Eq. (4).

The total scalar potential can then be written as

$$V(\phi, \sigma) = \frac{e^{F+F^*-2K}}{K_{\phi\phi^*}} \left\{ |(F_{\phi} - 2K_{\phi}) (m_{\sigma} + f) + f\phi|^2 + K_{\phi\phi^*} |f\sigma|^2 \right\} \quad (14)$$

For simplicity, and to avoid large mixing between φ and Σ sectors around the global minima induced by radiative corrections, we will choose the coefficients $a_i(\varphi)$ such that

$$\frac{da_i}{d\phi} \leq 10^{-12} \quad (15)$$

around the global minimum φ_0 and at the inflationary region. This is required to ensure the stability of the electroweak gauge hierarchy. In fact, since the inflaton-inflatino mass splitting is $|M_{\varphi}^2 - M_{\tilde{\varphi}}^2| \simeq [(10^{12} - 10^{13}) \text{ GeV}]^2$ we need the parameter $\eta \simeq (da_2/d\varphi)(\varphi_0)$ in Eq. (4) to satisfy Eq. (15) just to have an effective breaking of supersymmetry in the σ -sector, due to the gravitino contribution, $\mu_{3/2} \simeq O(1-10) \text{ GeV}$. Even if this effective breaking were negligibly small [although we can control its strength by the functions $a_i(\varphi)$], supersymmetry would be broken at low energies by the gaugino mass and the dynamical determination of the electroweak scale would proceed as in ordinary maximally symmetric no-scale models [13]. We may thus neglect f_{φ} in both regions, and write the potential (in these regions) as follows:

$$V(\phi, \sigma) = e^{F+F^\dagger - 2K} \tilde{V}(\phi, \sigma) \quad (16)$$

where

$$\tilde{V}(\phi, \sigma) = v(\phi) |m_0 + f|^2 + |f\sigma|^2, \quad (17)$$

and the function $v(\phi)$ has been defined in Eq. (12). It is a monotonic function from $\phi = 0(H)$, where $v \simeq 1/10$, to the global minimum. As in Ref. [19] the SU(5) phase transition occurs when $m_0 v(\phi_c) = a_2$. This also shows that for the phase transition to take place the mass of the σ -field around the SU(3) $\times$ SU(2) $\times$ U(1) minimum is $M_\sigma < (1/10) m_0 \simeq 10^{12}$ GeV, and $a_2 < m_0 / (10 \langle \sigma \rangle) \lesssim 10^{-4}$, which is unable to produce a sufficiently large reheating if the coefficients $a_i(\phi)$ are constant functions along the whole range of ϕ .

In the general case, where $a_2(\phi)$ is a function of ϕ , the SU(5) phase transition only requires (see Fig. 1)

$$a_2(\phi) \underset{\phi \simeq \phi_c}{\simeq} m_0 v(\phi_c) \quad (18)$$

locally and does not imply any further restriction for the value of $a_2(\phi_0)$ which can easily be orders of magnitude larger than $a_2(\phi_c)$. For instance, in the case where $a_2(\phi_0) \simeq 10^4 a_2(\phi_c)$, then the mass of the σ -field, $M_\sigma = a_2(\phi_0) \simeq 10^{16}$ GeV, is sufficiently large to reheat the Universe up to a temperature $T_R \sim V(0)^{1/4} \simeq (10^{15} - 10^{16})$ GeV by the decays of σ into ordinary matter fields through SU(5) gauge interactions.

We now follow the history of the Universe as dictated by our considerations. At high temperatures, $T > V(0)^{1/4} \simeq 10^{16}$ GeV, all fields sit at their symmetric minima and, in particular, the inflaton field ϕ chooses its metastable vacuum ($\phi = 0$). For temperatures $T \lesssim 10^{16}$ GeV the inflaton is trapped at the metastable de Sitter vacuum owing to the existence of a

small barrier (supercooling) in the inflaton potential and the energy density of the Universe is dominated by the cosmological constant. By tunnelling à la Hawking-Moss [20] (the Coleman-de Luccia mechanism is not important in our model because [21] $m_R \ll H$) the inflaton jumps the barrier and rolls slowly for 10^4 Hubble times, $H \Delta t \simeq 10^4$, producing the required amount of inflation. The present inflaton potential gives rise to a correct spectrum of energy density perturbations at the first horizon crossing, as was explained in detail in Refs. [15] and [19]. During this period, the SU(5) phase transition will take place. We can arrange it to happen during the inflationary epoch by choosing the critical value of the inflation, ϕ_c , at that position, as indicated in Fig. 1. We choose the function $a_2(\phi)$ so as to satisfy Eq. (15) for $\phi < \phi_1$, with $\phi_c < \phi_1 < \phi_0$, and $a_2(\phi) = M_X$ for $\phi > \phi_1$. For that purpose we can choose any analytic function that approaches the step function, as for instance the error function E

$$a_2(\phi) = m_0 V(\phi_c) + M_X E \left[\left(\frac{\phi - \phi_1}{M} \right)^n \right] \quad (19)$$

with n being some integer number such that the derivative of Eq. (19) satisfies the requirement (15) for $|\phi - \phi_1|/M > 1$. This happens very easily owing to the exponential behaviour of the derivative of the error function. For instance, if $\phi - \phi_1 \simeq 2-3$ it is enough to take $n \leq 2$. For the function $a_3(\phi)$ we can make the following two simple choices

$$\frac{a_2(\phi)}{a_3(\phi)} = M_X \quad (\text{Model I}) \quad (20a)$$

$$a_3(\phi) = a_3 = \text{constant} = \mathcal{O}(1) \quad (\text{Model II}). \quad (20b)$$

The global shape of the potential (17) for $\varphi < \varphi_1$ has been studied in Ref. [19] and we present here its main features. The potential $\tilde{V}$ has three stationary points along the real axis; they are given by

$$\sigma_1(\phi) = 0 \quad (21a)$$

$$\sigma_2(\phi) = \frac{a_2 - V(m_0 + f(\sigma_2))}{2a_3} \underset{\phi < \phi_1}{\simeq} \frac{a_2 - Vm_0}{2a_3} \quad (21b)$$

$$\sigma_3(\phi) = \frac{a_2(\phi)}{a_3(\phi)} \quad (21c)$$

For $\varphi < \varphi_c$, $\sigma_1 = 0$ is a saddle point (maximum along the real and minimum along the imaginary direction) and σ_3 is also a saddle point (minimum along the real and maximum along the imaginary direction), given in model I by $\sigma_3^I = M_x$ (Fig. 2a) and by $\sigma_3^{II} = [v(\varphi_c)/a_3] m_0$, in model II (Fig. 3a). Since the potential has the symmetry $\text{Im } \sigma \rightarrow -\text{Im } \sigma$ there are two global minima in the complex σ -plane around σ_3 and symmetric with respect to the real axis. As φ rolls these complex minima move towards σ_3 as $\varphi \rightarrow \varphi_c$, and collapse to it at $\varphi = \varphi_c$. On the other hand, σ_2 is a local minimum in the σ -plane which also moves towards σ_1 as $\varphi \rightarrow \varphi_c$. For $\varphi = \varphi_c$ the real direction at $\sigma_1 \equiv \sigma_2$ and the imaginary direction at σ_3 become flat (Figs. 2b and 3b). The σ_3 becomes global minimum of the potential and the system suffers a second-order phase transition.

In model II, $\sigma_3^{II} \simeq m_0 \simeq H$ is comparable to the thermal fluctuations so that they will trigger the beginning of phase transition. In model I, $\sigma_3^I \simeq M_x \gg H$ so that the transition will proceed slowly since the σ -potential is very flat. This follows because, in both models, the curvature of the σ -potential for $\varphi < \varphi_1$ around the global minimum σ_3 is $M_\sigma^2 \simeq H^2$. For

$\varphi_c < \varphi < \varphi_1$ (Figs. 2c and 3c) σ_2 becomes a saddle point (maximum along the real and minimum along the imaginary direction) and begins to construct a barrier between σ_1 and σ_3 . In model II the system rolls down mainly along the φ -direction so that energy density perturbations $\delta\rho/\rho \simeq 10^{-4}$ are produced as in the absence of a σ -field. In model I the system rolls down following an intermediate direction in the φ - σ plane and, since the potential along the σ -direction is smoother than along the φ -direction ($M_\sigma < m_0$), energy density perturbations $\delta\rho/\rho \lesssim 10^{-4}$ are expected. At $\varphi \simeq \varphi_1$, $a_2(\varphi)$ takes the value M_X . For model I, σ_3 does not change, but the curvature (and with it the barrier between minima) increases from $M_\sigma \simeq m_0$ to $M_\sigma \simeq M_X$ (Fig. 2d). In model II, σ_3 changes from m_0 to M_X and also M_σ (the curvature) changes as in model I (Fig. 3d). In Fig. 4 we show the schematic three-dimensional forms of the potential in model I (Fig. 4a) and in model II (Fig. 4b).

In both models the inflaton near its global minimum is really a mixture of the fields φ and σ . The two components have very different decay rates due to their interaction strengths. The σ component has gauge interactions (short lifetime), whereas the φ component interacts only through gravitational interactions, having thus a long lifetime. In this respect, the system (σ, φ) resembles the familiar (K_S^0, K_L^0) -system. In particular, the decay rate of the σ component is

$$\Gamma_\sigma \simeq \alpha_{\text{GUT}} M_X > H.$$

The whole system enters the reheating regime when σ decays to ordinary matter through SU(5) gauge interactions [8], and the Universe is reheated to the temperature $T_R \simeq V(0)^{1/4} \simeq (10^{15} - 10^{16})$ GeV.

At $\varphi = \varphi_0$ there are two degenerate minima, the one with $\sigma = 0$ and the other with $\sigma = M_X$. Both have zero vacuum energy. However, they are separated

by a barrier with height M_X^4 , situated at $1/2 M_X$ (see Figs. 2e and 3e). Therefore the Universe will be trapped in the broken phase because the energy of the inflaton field fluctuations is not high enough to drive it over the barrier. When the SU(5) phase becomes unstable there will be other equally favourable directions in the group space besides SU(3) x SU(2) x U(1). We do not consider this a problem. Roughly speaking, this means only that in 50% of the inflated universes the symmetry will be SU(4) x U(1). Ours belong to the other 50% with SU(3) x SU(2) x U(1) symmetry.

In both models I and II the density of magnetic monopoles will be diluted by the residual inflation that takes place after phase transition (but before reheating), i.e. when $\phi_c < \phi < \phi_1$. This is more than welcome, since the thermal dilution of the monopole density, which is given by $(T_H/M)^3$, is not sufficient to suppress the monopole-to-photon ratio below the cosmological [22] upper bound 10^{-27} . In our model, when $\phi \gtrsim 10^2 H \simeq 10^{-4}$, the potential is dominated by the cubic term, and we can solve the equations of motion for $\phi \gtrsim \phi_c$. If we denote $V \simeq -\gamma H^2 \phi^3$, where γ is a constant (in our case $\gamma = 0.39$), we find that the number of e-foldings that the scale factor of the Universe will undergo as ϕ rolls from ϕ_c to ϕ_1 is

$$H \Delta t = \frac{1}{\gamma \phi_c} \left[1 - \frac{\phi_c}{\phi_1} \right]. \quad (22)$$

With $\phi_1 \simeq 1$ we satisfy the cosmological bound on the monopole density if $\phi_c \lesssim 0.34$ (in our case). This is a rather mild constraint. In fact, to satisfy the five orders of magnitude more severe bound from neutron stars [23], we still would need only $\phi_c \lesssim 0.22$. Therefore the additional dilution of the monopole density occurs naturally in our model. However, we cannot exclude the possibility that this dilution is not too effective and that there might be an observable number of monopoles in the present Universe.

5. DOUBLET-TRIPLT SPLITTING AND BARYON ASYMMETRY GENERATION

The baryon-to-entropy ratio will be generated by out-of-equilibrium decays of heavy coloured triplets ($M_H \simeq 10^{16} - 10^{17}$ GeV) into ordinary matter fields, as in Ref. [10]. To discuss this baryon asymmetry generation, we must now specify the interactions of all other fields besides φ and σ .

Our model will be completely defined by the total superpotential $F(\varphi, \Sigma, y^a)$, which includes all matter superfields. We choose

$$F(\phi, \Sigma, y^a) = F(\phi, \Sigma, y^a=0) + F_{DT}(H^i, \bar{H}^i, \theta^i, \bar{\theta}^i, \Sigma) + F_{Yuk}(y^a), \quad (23)$$

where F_{Yuk} contains all ordinary (trilinear) Yukawa couplings,

$$F_{Yuk} = C_{ijk} Q_i Q_j^c H_k. \quad (24)$$

Here Q_i are ordinary matter superfields [quark and lepton supermultiplets in $\bar{5}$ and 10 representations of $SU(5)$], H_K are fiveplet Higgses ($5 + \bar{5}$ representations), and F_{DT} is superpotential responsible for the doublet-triplet splitting of H_K through the missing partner mechanism [24]

$$F_{DT} = \lambda_1^{ij} \bar{\theta}_i \Sigma'^n \theta_j + \lambda_2^{ij} H_i \Sigma'^2 \bar{\theta}_j + \lambda_3^{ij} \bar{H}_i \Sigma'^2 \theta_j + \delta^{ij} \bar{H}_i \Sigma' H_j, \quad (25)$$

where θ_i are superfields in the $\underline{50}$ representation, and

$$\det \delta_{ij} = 0 \quad (26)$$

so that at least one linear combination of doublet fields remains massless and can break $SU(2) \times U(1)$ at low energy. In the case that only one doublet

stays massless, and the other heavy ($\sim M_X$), only one pair of $\theta, \bar{\theta}$ is needed in Eq. (25).

To have Eq. (26) automatically satisfied, as a consequence of a global symmetry, then $n = 5$ is required, as in Ref. [10], and $\delta_{ij} = 0$. In that case the Γ' field in Eq. (25) cannot be identified with the Γ field of the superpotential (13b). If we insist on the identification of both fields, the presence of that global symmetry will also imply $a_n = 0$ and, therefore, no SU(5) breaking induced by the inflaton field, with the hard-won good reheating, exists. One possibility would be to introduce two adjoint fields Γ and Γ' . The Γ would produce the reheating after inflation (by gauge interactions) and Γ' the doublet-triplet splitting, while the correct scale of SU(5) breaking would be $M_{\text{GUT}}^2 = \langle \Gamma \rangle^2 + \langle \Gamma' \rangle^2$. In fact, along the SU(3) $\times$ SU(2) $\times$ U(1) direction the potential is flat for the Γ' field so that $\langle \Gamma' \rangle$ will be fixed by non-gravitational radiative corrections and is expected to be $O(M)$. The mass of triplets is now $O(\cos^2 \theta_\sigma M_{\text{GUT}}^2)$, where $\sin \theta_\sigma = \langle \Gamma \rangle / M_{\text{GUT}}$, while the mass of the reheating fields σ is fixed to be $M_\sigma = a_2(\varphi_0) \simeq (10^{15} - 10^{16})$ GeV and the vacuum expectation value of Γ is now liberated from the bound $M_T > O(10^{16} - 10^{17})$ GeV. However, we believe that temperature corrections will keep probably $\langle \Gamma' \rangle = 0$ forever.

The most economical alternative, from the point of view of the number of fields, is to identify $\Gamma' = \Gamma$, to keep only one pair of θ 's and one pair of light doublets, abandoning thus the globally symmetric solution $n = 5$. The proposed doublet-triplet superpotential is chosen to be ($n = 2$):

$$F_{DT} = \lambda_1 \bar{\theta} \Sigma^2 \theta + \lambda_2 H_1 \Sigma^2 \bar{\theta} + \lambda_3 \bar{H}_1 \Sigma^2 \theta + \delta \bar{H}_2 \Sigma^2 H_2, \quad (27)$$

with the coefficients λ_i and δ of the order of 1. In this minimal model, the requirement that the reheating temperature be $T_R \simeq (10^{15} - 10^{16})$ GeV implies that the masses generated by the Γ -superpotential, i.e. the masses of fields in the unbroken direction of the effective theory $SU(3) \times SU(2) \times U(1)$, are $M_\sigma \simeq M_{(8,1)} \simeq M_{(1,3)} \simeq 10^{16}$ GeV, so that all superheavy degrees of freedom are massive enough to guarantee the perturbative calculations up to scales close to the Planck mass, $\alpha_i(M) < O(1)$. This permits a slight modification of that minimal model by replacing the 24 Γ -field in Eq. (13) by the 75 Ξ -field and using the conventional doublet-triplet [25] superpotential

$$F_{DT} = \lambda_1 \bar{\theta} \Xi^n \theta + \lambda_2 H_1 \Xi \bar{\theta} + \lambda_3 \bar{H}_1 \Xi \theta + \delta \bar{H}_2 \Xi H_2 \quad (28)$$

with $n = 0, 1$, which does not contain dimensions larger than three terms. Whether one uses the 24 or the 75 representation, the effective superpotential relevant for baryon asymmetry generation and proton decay is the same, and can be written as $\mu_1 \bar{H}_1 H_2 + \mu_2 \bar{H}_2 H_2$, where $\mu_i \simeq M_T \simeq (M_X^2 \text{ or } M_X)$, depending on the model chosen (Γ^2 or Ξ). The present experimental limits for proton decay require [26] $M_T \gtrsim 10^{16}$ GeV, which implies

$$\begin{aligned} M_X &\gtrsim 10^{16} \text{ GeV} \quad \text{in the } \Xi\text{-model} \\ M_X &\gtrsim 10^{17} \text{ GeV} \quad \text{in the } \Sigma^2\text{-model.} \end{aligned} \quad (29)$$

Since we do not want the analysis of the inflaton potential that we made in Section 3 to be modified by matter fields (including Γ -fields) we will usually require, besides Eq. (15), that the gravitino mass

$$m_{3/2}^2 = e^{\langle F+F^\dagger-2K \rangle} \left| m_0 + \frac{\langle f(\phi, \sigma) \rangle}{M^2} \right|^2 \chi^3 \quad (30)$$

at the global minimum ϕ_0 remains unchanged with respect to its value at the metastable minimum $\phi = 0$. This requirement, $m_{3/2}(\phi) \simeq m_{3/2}(0) \simeq 10^{12}$ GeV, which is the most natural solution to the inflation scale problem, is more than welcome to avoid other minima of the scalar potential, Eq. (14), in the (ϕ, σ) -plane. For instance, the solutions to the equation $m_0 + f(\phi, \sigma) = 0$ would correspond to global minima (with zero vacuum energy) in the complex ϕ -plane. Thus we impose the absolute upper bound

$$\langle f(\phi, \sigma) \rangle \leq m_0 M^2, \quad (31)$$

which translates into

$$a_2 \leq \theta \left[m_0 \left(\frac{M}{M_X} \right)^2 \right] = \begin{cases} 10^{16} \text{ GeV}, & M_X = 10^{16} \text{ GeV} \\ 10^{15} \text{ GeV}, & M_X = 10^{17} \text{ GeV} \end{cases} \quad (32)$$

bounds which permit maximal reheating temperature $T_R \simeq V(0)^{1/4} \simeq (10^{15} - 10^{16})$ GeV, since indeed $\Gamma_\sigma > H$ and thus $T_R \simeq (HM)^{1/2}$.

In this case, the stability of the electroweak scale against non-gravitational radiative corrections

$$\delta m_{\text{scalars}} \simeq a'_2(\phi_0) m_{3/2} \lesssim 10^2 \text{ GeV} \quad (33)$$

or

$$a'_2(\phi_0) \lesssim 10^{-16} m_{3/2}^{-1} \simeq 10^{-10} - 10^{-11} \quad (34)$$

is guaranteed by the condition (15).

Baryon asymmetry in our model is generated by the out-of-equilibrium decays of heavy coloured triplets. The mechanism is similar to that discussed in Ref. [10], where the amount of baryon asymmetry was calculated in the context of exact supersymmetry. In our case, there is no substantial suppression factor due to small $H\bar{H}$ -mixing in the doublet-triplet superpotentials (27) and (28). Therefore, the baryon asymmetry per H decay is given simply by

$$\Delta B \simeq \left(\frac{\lambda g_Y}{4\pi} \right)^2 \theta, \quad (35)$$

where g_Y is the Yukawa coupling of the b -quark at M_X , λ is a generic superpotential coupling in Eqs. (27) and (28), and θ is a CP-violating angle possibly $O(1)$. If $\lambda \sim O(1)$, we find $\Delta B \simeq 10^{-6} - 10^{-7}$; the exact number depends on the details of the electroweak radiative breaking, which determines the relative magnitudes of the vacuum expectation values of the two Higgs fields.

However, there is an additional entropy production due to the reheating produced by decays of ϕ -fields and gravitinos. In both cases the amount of entropy production is estimated to be [16]

$$\left(T_R / T_D \right)^3 \simeq m^{-1/2} N(M_P)^{-1/4} N(T_D)^{-1/2}, \quad (36)$$

where $m = (m_\phi, m_{3/2})$, and T_D and T_R are the corresponding decay and reheating temperatures. For the models described by the superpotential [(23), (27), and (28)] $N(M_P) = 423.75$, the decay temperature is

$$\begin{aligned} T_{D_\phi} &\simeq 1.3 \times 10^{10} \text{ GeV} \\ T_{D_{3/2}} &\simeq 3 \times 10^8 \text{ GeV}, \end{aligned} \quad (37)$$

where we have taken $m_\phi \simeq 10^{13}$ GeV, $m_{3/2} \simeq 10^{12}$ GeV, and $N(T_D) = 243.75$.

Thus the amount of entropy production due to ϕ -decays is

$$\left(T_{R_\phi} / T_{D_\phi} \right)^3 \simeq 5 \quad (38)$$

and to gravitino decays

$$\left(T_{R_{3/2}} / T_{D_{3/2}} \right)^3 \simeq 14 \quad (39)$$

and the total amount of entropy production is

$$\left(T_R / T_D \right)_T^3 \simeq 70 \quad (40)$$

We can conclude then that in our model the baryon-to-entropy ratio is

$$10^{-11} < \frac{n_B}{s} < 10^{-10} \quad (41)$$

Note that in this model ϕ -decays cannot produce baryon asymmetry because $m_\phi \ll m_H$. Baryon asymmetry is therefore generated by decays of heavy triplets (as in Ref. [10]) produced out-of-equilibrium after the reheating of the Universe by the σ -fields.

Both models we have just presented (the Σ^2 and the Ξ -model) have the same effective low-energy theory so that it is not possible to distinguish them by low-energy requirements. It is just a question of taste and we leave the reader to choose between them. However, we have not been able to find a symmetry argument to forbid the appearance of δ -terms in Eq. (25) without introducing an additional Σ' -field.

6. CONCLUSION

We have presented a model of inflation and reheating based on maximally symmetric supergravity. It is a genuine two-component model of inflation, where the scale of reheating does not depend on the scale of the inflaton potential at $\varphi \simeq 0$. As we have shown, this attractive possibility can be accomplished when the mass of the adjoint Higgs field σ depends on the inflaton field. Indeed, since φ is a gauge singlet, it can be argued that such a direct coupling between σ and φ is very natural. Then it becomes possible to have an SU(5) breaking phase transition during inflation, which in turn leads to a high reheating temperature $T_R \simeq (10^{15} - 10^{16})$ GeV. We have presented a detailed discussion on such a model with two realistic examples. In both cases we found the high reheating temperature necessary for baryon asymmetry generation and suppression of $d = 5$ operators.

In the present model the SU(5) phase transition modifies the overall potential only little before the onset of reheating. Therefore the potential is very flat also in the broken phase, and one can easily achieve a huge suppression of the present monopole density; on the other hand, the critical field φ_c is a free parameter in our model and consequently the number of monopoles can also remain not far below the present observational bounds.

An interesting open question is the extent of effects due to thermal corrections on the two-component maximally symmetric inflation. We do not expect them to affect the main features of inflation or reheating, although they may restrict the allowed range of parameters of the model. A detailed survey of thermal corrections in our maximally symmetric inflation will be published elsewhere [27].

The dynamical determination of the electroweak breaking follows lines similar to those discussed elsewhere [19,28]. In particular, the top quark mass depends on the effective symmetry breaking at low energy induced by

the inflaton field, which in turns depends on the coupling between the φ and σ fields. In the limit where these fields are decoupled the top quark mass is predicted to be [13] $m_t \simeq 45$ GeV, in agreement with preliminary experimental results. On the other hand, if φ and σ are coupled, so that the effective breaking is comparable to the gaugino mass, then $m_t \simeq 65$ GeV, as in minimal coupling models [29] if there is no $m_{\tilde{t}} \bar{H} H$ mixing in the superpotential for Weinberg-Salam Higgs doublets, but it can be lowered easily to the expected range $30 \text{ GeV} < m_t < 50 \text{ GeV}$ if $m_{\tilde{t}} \neq 0$, as was shown elsewhere [30].

Finally, we would like to emphasize again that the only scale of the inflaton potential is fixed by the scale of energy density fluctuations $\delta\phi/\phi \simeq 10^{-4}$. It provides $m_0 \simeq H \simeq m_{3/2}(\varphi = 0) \simeq m_{3/2}(\varphi_0) \simeq 10^{12}$ GeV. It is striking that this value is also required by the observed baryon-to-entropy ratio. In fact, inflaton and gravitino decays produce a total amount of entropy production estimated as $m_{3/2}^{-1} N(M_p)^{-1/2} N(T_D)^{-1} \simeq 10^{-4} m_{3/2}^{-1}$, and the baryon-to-entropy ratio is $n_B/s \simeq (10^{-4} - 10^{-5}) m_{3/2}$. Thus the experimental value $n_B/s \simeq 10^{-10} - 10^{-11}$ leads again to $m_{3/2} \simeq 10^{12}$ GeV. The electroweak gauge hierarchy is stable under non-gravitational radiative corrections because of the very small renormalizable mixing between the inflaton field and matter fields.

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Figure captions

- Fig. 1 : Schematic view of the inflaton potential, Eq. (12), along the real axis. All of its points correspond to minima along the imaginary axis. $a_2(\varphi)$ is the coefficient of the adjoint Higgs superpotential, Eq. (13c), which will change in the regime around $\varphi \simeq \varphi_1$. The GUT phase transition will take place at $\varphi = \varphi_C$.
- Fig. 2 : Schematic view of the SU(5) phase transition as given by the superpotential (13a) for model I in the text. For $\varphi < \varphi_C$, (a), there is a minimum at the real negative axis and saddle points at $\sigma = 0$ and at $\sigma = M_X$. For $\varphi = \varphi_C$, (b), the minimum along the real axis and the saddle point at $\sigma = 0$ collapse and the SU(5) phase transition begins whilst the saddle point at M_X becomes a minimum with a curvature m_0^2 along the real axis and a flat direction along the imaginary axis. For $\varphi_C < \varphi < \varphi_1$, (c), a barrier begins to develop between the minima at the origin and at M_X , and the σ -field begins to roll down to the broken minimum. For $\varphi = \varphi_1$, (d), $a_2(\varphi)$ goes to M_X and both the curvature at $\sigma = M_X$ and the barrier separating the minima at 0 and M_X increase rapidly to the values M_X^2 and M_X^4 , respectively. The reheating by σ -field starts now. At $\varphi = \varphi_0$, (e), the global minimum of the inflaton field, both minima, at 0 and M_X , are degenerate and have zero vacuum energy but the system is at the broken minimum, $\sigma = M_X$, and reheats the Universe by gauge interactions.

Fig. 3 : Schematic view of the SU(5) phase transition for model II in the text. For $\varphi < \varphi_1$, (a), (b) and (c), the saddle point becoming a minimum (global) at the positive real axis is at $\sigma \simeq H$ so that phase transition will be easily driven by thermal fluctuations. Thus at $\varphi_c < \varphi < \varphi_1$, (c), the system is trapped at the "broken" ($\sigma \simeq H$) minimum. At $\varphi = \varphi_1$, (d), the location of the minimum is shifted from $\sigma \simeq H$ to $\sigma \simeq M_X$ so that the system is driven into the normal broken minimum. At the same time the curvature and the barrier between minima jump into M_X^2 and M_X^4 , as in Fig. 2, and the reheating starts. Finally, at the global minimum of the inflaton potential the σ -field oscillates and reheats through its decays and the configuration of the potential is identical to that of Fig. 2e.

Fig. 4 : A schematic three-dimensional representation (not to scale) of the potential (14) in (a) model I and (b) model II [see Eqs. (20)]. The wide, black line is a possible trajectory of the inflaton field.

Fig. 1

Fig. 2a

Fig. 2b

Fig. 2c

Fig. 2d

Fig. 2e

Fig. 3a

Fig. 3b

Fig. 3c

Fig. 3d

Fig. 3e

Fig. 4a

Fig. 4b