RAPID PHASE TRANSITIONS IN PRIMORDIAL INFLATION
WITH FLAT SUPERGRAVITY POTENTIALS

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ABSTRACT

We demonstrate that the recently proposed naturally flat $N = 1$ supergravity potentials not only provide a working primordial inflation scenario, but they also considerably facilitate the GUT phase transition. Towards the late stages of inflation, a rapid (second order) phase transition occurs naturally to the $SU(3) \times SU(2) \times U(1)$ phase, thus gracefully and simultaneously ending the reign of the inflationary as well as GUT eras.

The success of the inflatory Universe^{1),2)} in explaining the many outstanding problems of the standard big bang model is by now well known. These problems include the puzzles of homogeneity and isotropy of the present-day Universe, the origin of galaxy formation and the absence of magnetic monopoles. However, GUT potentials of the Coleman-Weinberg type, which are the most promising candidates for producing inflation in GUTs²⁾, can do the job only through a severe fine-tuning of the scalar masses³⁾. However, although they give rise to scale-independent energy density fluctuations needed for galaxy formation⁴⁾, these come out to be too large by four orders of magnitude. This has led several people⁵⁾⁻⁷⁾ to consider supersymmetric primordial inflation, where inflation is driven by a GUT singlet field (the inflaton) at the scale of the Planck mass M_P , where all fine-tunings disappear. This is done most naturally in the context of models based on simple $N = 1$ supergravity⁸⁾.

In primordial inflation, the inflatory period is divorced from the GUT phase transition. At the end of inflation, the fields will usually find themselves in the symmetric minimum^{5),6),9)}, which is separated by a potential barrier from the broken minima [e.g., $SU(3) \times SU(2) \times U(1)$ or $SU(4) \times U(1)$]. To accomplish the GUT phase transition, one has to arrange for a small barrier height and resort to strong coupling phenomena¹⁰⁾. The idea is that, as the Universe cools in the symmetric phase, the gauge coupling becomes strong, leading to confinement and disappearance of the light degrees of freedom from the particle spectrum of the symmetric phase. Finite temperature effects will then lower the broken minima below the symmetric one¹⁰⁾ and phase transition becomes possible.

Although strong coupling phenomena may well be responsible for the GUT phase transition, the very notion of strong coupling involves many unknowns. In this respect, the possibility of a rapid second order phase transition after inflation would be more than welcome. In the present paper, we will show that this is indeed what will happen if inflation is considered within the recently proposed class of supergravity models with a naturally flat potential^{11),12)}, and consequently a vanishing cosmological constant.

The flatness of the potential can be obtained if the theory has a global non-compact $SU(1,1)$ symmetry broken by the gravitino mass term¹²⁾. It is interesting that these theories may provide us a theory with no other

intrinsic mass scale than the Planck mass. The $SU(1,1)$ symmetry forces the tree level scalar potential

$$V = e^G \left[G_a (G^{-1})^{\bar{a}b} G_{\bar{b}} - 3 \right] + \frac{1}{2} D^\alpha D_\alpha \quad (1)$$

to be flat with an undetermined gravitino mass. Here $G = G(\phi, \phi^*)$ is the Kähler potential which depends on a set of chiral superfields ϕ , and

$$G_a = \frac{\partial}{\partial \phi_a} G \quad ; \quad G_{\bar{a}} = \frac{\partial}{\partial \phi_a^*} G \quad ; \quad G_{a\bar{b}} = \frac{\partial^2}{\partial \phi_a \partial \phi_b^*} G \quad (2)$$

Throughout the paper the units are such that $M_P/\sqrt{8\pi} = 1$. The flatness of (1) is ensured by choosing a non-trivial Kähler potential G for the hidden fields z , which are responsible for the supersymmetry breaking. This is in fact imposed on us by the $SU(1,1)$ symmetry, which leads to an Einstein space with¹²⁾

$$R_{z\bar{z}} = \partial_z \partial_{\bar{z}} \ln G_{z\bar{z}} = \frac{2}{3} G_{z\bar{z}} \quad (3)$$

where $R_{z\bar{z}}$ is the curvature of the Kähler manifold. Up to field redefinitions, (3) is solved uniquely by

$$G = -3 \ln (z + z^*) \quad (4)$$

The gravitino mass $m_{3/2}$ is generated dynamically through radiative corrections, and $m_{3/2} \sim M_W \sim M_P \exp(-O(1)/\alpha)$.

To implement inflation into supergravity models with the above desirable properties, the Kähler potential (4) needs to be generalized to include the inflaton field ϕ and the observable matter fields y_a . A suitable solution is⁷⁾

$$G = -3 \ln (z + z^* + h(\phi, \phi^*)) + g(\phi, \phi^*, y_a, y_a^*) \quad (5)$$

This leads to a semi-positive potential⁷⁾

$$V = e^G \left[\frac{|g_\phi|^2}{(g_{\phi\bar{\phi}} + G_Z h_{\phi\bar{\phi}})} + |g_{y_a}|^2 \right] + \frac{1}{2} D^\alpha D_\alpha \quad (6)$$

where $g_\phi = (\partial/\partial\phi)g$, etc., and a sum over a is implied. The form of the potential (6) ensures that there are no potentially dangerous supersymmetric minima with a negative cosmological constant. The global minima of (6) are given by the conditions

$$g_\phi = g_{y_a} = 0 \quad (7)$$

At the tree level, the gravitino mass is again undetermined. Its value and the scale of SUSY breaking will be determined by radiative corrections.

In addition to the global minima (7), there can exist local minima. As was shown in Ref. 7), when $g_{\phi\bar{\phi}} < 0$ and $h_{\phi\bar{\phi}} < 0$, the local minimum can have a positive cosmological constant, thus giving rise to inflation. When $g_\phi \neq 0$, the minimizing condition $V_Z = 0$ leads to

$$G_Z = -\frac{3}{2} g_{\phi\bar{\phi}} / h_{\phi\bar{\phi}} ; \quad g_{y_a} = 0 \quad (8)$$

The gravitino mass $m_0 = e^{G/2}$ at that point is related to the scale of inflation. At the end of inflation, when fields arrive at their global minima (7), the gravitino mass relaxes dynamically to the value $m_{3/2} \sim M_W$.

To consider GUT phase transition in primordial inflation with flat supergravity potential, it is useful to separate from $g(\phi, \phi^*, y_a, y_a^*)$ the part depending on the total superpotential $F_T(\phi, y_a)$:

$$g = K(|\phi|) + |y_a|^2 + \ln |F_T|^2 \quad (9)$$

where we have chosen a minimal kinetic term for the observable matter fields y_a . As in Ref. 7), $h(\phi, \phi^*)$ and $K(|\phi|)$ are chosen to be

$$K = -2|\phi|^2 + \frac{1}{2}\beta|\phi|^4 \quad (10)$$

$$h = K / 2\alpha$$

where we take $\alpha = 1$ for simplicity. In the following we will assume that the total superpotential $F_T(\phi, y_a)$ can be written as a sum of the superpotentials of the inflaton and matter sectors:

$$F_T(\phi, y_a) = F_1(\phi) + F_2(y_a) \quad (11)$$

The inflaton superpotential $F_1(\phi)$ can be written as

$$F_1(\phi) = \left(\frac{\phi - \phi_0}{\phi_0} \right)^2 e^{P(\phi)} m^3 \quad (12)$$

where m^3 gives the scale of inflation and ϕ_0 is the value of the inflaton field at the global minimum (7). $P(\phi)$ is a polynomial such that all the necessary conditions for inflation are satisfied. This entails a definite expansion for $P(\phi)$ around $\phi = 0$ and has been considered in detail in Ref. 7). Sufficient inflation was found to arise with

$$P(\phi) = -\xi\phi + \frac{1}{2}(2 - \xi^2)\phi^2 - \frac{1}{3}\xi^3\left(1 - \frac{\beta}{2\xi^2} + \frac{\epsilon}{2}\right)\phi^3 + O(\phi^4)$$

where $\xi = 0.32$, $\epsilon = 5 \times 10^{-4}$, and β was set equal to $(2\phi_0^2)^{-1}$ with $\phi_0 = 5$. The scale of inflation was given by $m^3 = 3 \times 10^{-6}$. We will also choose these values for definiteness, although the GUT phase transition is not sensitive to them.

The condition (8) for a local minimum is obviously satisfied for gauge non-singlets at $y_a = 0$. As high temperature effects prefer gauge conserving minima¹³⁾, inflation will begin while the gauge non-singlet fields sit at this minimum. However, as the inflaton field tunnels through the barrier and starts rolling towards the global minimum (7), the shape of the potential in the y_a direction may and indeed will change. As we will show, the potential

barrier separating the gauge conserving minimum at $y_a = 0$ from gauge breaking minima at $y_a = \langle y_a \rangle \neq 0$ can disappear during inflation, thus facilitating a rapid phase transition to a broken phase.

We will consider an SU(5) model with an adjoint Higgs field Σ . All other fields are assumed to be at their minima $y_a = 0$ and are of no consequence to the following analysis. Consider therefore the superpotential

$$F_2(\Sigma) = \sum_{n \geq 2} c_n \text{Tr} \Sigma^n \quad (13)$$

The broken minima may lie either in the SU(3)×SU(2)×U(1) preserving direction or in the SU(4)×U(1) preserving direction:

$$\begin{aligned} \langle \Sigma \rangle &= \sqrt{\frac{2}{15}} \sigma \text{diag} \left(1, 1, 1, -\frac{3}{2}, -\frac{3}{2} \right) \text{ for } \text{SU}(3) \times \text{SU}(2) \times \text{U}(1) \\ \langle \Sigma \rangle &= \frac{1}{\sqrt{20}} \sigma \text{diag} (1, 1, 1, 1, -4) \text{ for } \text{SU}(4) \times \text{U}(1) \end{aligned} \quad (14)$$

As we will show later, the SU(3)×SU(2)×U(1) direction will be favoured over SU(4)×U(1). Meanwhile, it suffices to note that in both cases we are led to consider an effective superpotential $F_2(\sigma)$ for the SU(5) singlet field σ :

$$F_2(\sigma) = \sum_{n \geq 2} a_n^{(k)} \sigma^n \quad (15)$$

where $k = 1, 2$ refers to SU(3)×SU(2)×U(1) and SU(4)×U(1) cases, respectively. The coefficients $a_n^{(k)}$ are related to c_n by

$$a_n^{(1)} = 3 \left(\frac{2}{15} \right)^{\frac{n}{2}} \left[1 - \left(-\frac{3}{2} \right)^{n-1} \right] c_n$$

$$a_n^{(2)} = 4 \left(\frac{1}{20} \right)^{\frac{n}{2}} \left[1 - (-4)^{n-1} \right] c_n \quad (16)$$

Let us focus on the configurations in the field space which lie along the minimizing curve $G_z = (-3/2)g_{\phi\phi}/h_{\phi\phi}$ [see Eq. (8)] but for which $g_\sigma \neq 0$. We will examine the behaviour of the potential (6) in the σ direction. We restrict ourselves to real ϕ only. Although the path of inflaton does not necessarily follow the line $\text{Re } \phi$, the v.e.v. of ϕ can be chosen to be real. Since, as we will see later, the GUT phase transition can happen only in the very late stages of inflation, when $\text{Im } \phi$ must be small, this restriction to real ϕ is justified. The potential can then be expanded around $\sigma = 0$, using Eqs. (9), (11) and (15), to yield

$$V(\sigma, \phi) = J(\phi) \left[f^2(\phi) + U(\phi, \sigma) \right] \quad (17)$$

where

$$J(\phi) = \frac{e^K}{1 - \beta\phi^2} \quad (18)$$

Here $f(\phi)$ is the generalized derivative defined by

$$f(\phi) = F_{1\phi}(\phi) + K_\phi(\phi)F_1(\phi) \quad (19)$$

where F_1 and K are given by Eqs. (10) and (11). The expansion of $U(\phi, \sigma)$ around $\sigma = 0$ is

$$\begin{aligned}
 U(\phi, \sigma) = & \left[f^2(\phi) + (1 - \beta\phi^2)(F_1^2(\phi) + 4a^2) \right] |\sigma|^2 \\
 & + \left[2f(\phi)K_\phi a + 4(1 - \beta\phi^2)aF_1(\phi) \right] \text{Re}\sigma^2 \\
 & + \left[2f(\phi)K_\phi + 6F_1(\phi)(1 - \beta\phi^2) \right] b \text{Re}\sigma^3 \\
 & + 12ab(1 - \beta\phi^2) |\sigma|^2 \text{Re}\sigma + \mathcal{O}(\sigma^4)
 \end{aligned} \tag{20}$$

where we have denoted $a_2^{(k)} = a$, $a_3^{(k)} = b$.

With the choice $\beta = (2\phi_0^2)^{-1}$, $\phi_0 = 5$, one readily sees from Eq. (10) that $K_\phi < 0$ in the interval $0 < \phi \leq \phi_0$. Since $K_\phi(\phi=0) = 0$, from Eq. (20) one finds that around $\phi \approx 0$ the minimum $\sigma = 0$ is protected by a potential barrier in the whole complex plane. As ϕ keeps on rolling towards ϕ_0 and $F_1(\phi) \rightarrow 0$, this situation will change, and as we will see, the potential barrier will vanish. However, before we embark on a detailed analysis, let us note that it suffices to consider real σ only. When $\sigma \rightarrow 0$,

$$\begin{aligned}
 U(\phi, \sigma) \approx & \left[f^2 + 2K_\phi f a + (1 - \beta\phi^2)(2a + F_1)^2 \right] (\text{Re}\sigma)^2 \\
 & + \left[f^2 - 2K_\phi f a + (1 - \beta\phi^2)(2a + F_1)^2 \right] (\text{Im}\sigma)^2 + \mathcal{O}(\sigma^3)
 \end{aligned}$$

Since all other terms but $2fK_\phi a$ are positive definite, we can always choose the sign of a in such a way that the curvature of the potential is positive in the $\text{Im } \sigma$ direction. From now on, we choose $a < 0$ and restrict ourselves to real σ only. This results in a potential $U(\phi, \sigma)$ which is given by

$$U(\phi, \sigma) = A(\phi)\sigma^2 + B(\phi)\sigma^3 + \mathcal{O}(\sigma^4) \tag{21}$$

where

$$A(\phi) = f^2 + 2K_\phi f a + (1-\beta\phi^2)(2a+F_1)^2 \quad (22)$$

$$B(\phi) = 2b [K_\phi f + 3(1-\beta\phi^2)(2a+F_1)]$$

The potential barrier confining σ to $\sigma = 0$ will disappear during inflation if $A(\phi)$ has any zeros in the interval $0 < \phi < \phi_0$. As is evident, this can happen only when $f \sim 0(a)$, or when $\phi \rightarrow \phi_0$; that is, at the very end of the inflatory period. But then, by virtue of (12),

$$|F_{1\phi}| \sim 2|\phi - \phi_0| \gg F_1 \sim (\phi - \phi_0)^2$$

This means that $f \sim F_{1\phi} < 0$. Therefore, as $\phi \rightarrow \phi_0$, $A(\phi)$ can be written as

$$A(\phi) \simeq f^2 + 2K_\phi f a + 4a^2(1-\beta\phi^2) \quad (23)$$

Near $\phi \simeq \phi_0$, K_ϕ can be approximated by $K_\phi \simeq -3\phi_0/2$, and $1-\beta\phi^2 \simeq \frac{1}{2}$ [see Eq. 10)]. Therefore a good approximation of (23) is

$$A(\phi) \simeq f^2 - 3fa\phi_0 + 2a^2 \quad (24)$$

Since $f < 0$ and a was also chosen to be negative, it is now clear that there will be a period during inflation when $A(\phi) < 0$. This occurs when

$$15a \lesssim f(\phi) \lesssim 0.13a \quad (25)$$

Therefore, during this time a rapid phase transition to an $SU(5)$ breaking minimum is possible^{*)}.

Because the potential barrier reappears again when $f \sim 0.13a$, there exists the possibility that the full $SU(5)$ symmetry becomes restored again. However, it turns out that the fields will become trapped in the broken phase. This follows from the fact that when $\phi \rightarrow \phi_0$, $B(\phi)$ is approximately [see Eq. (22)]

$$B(\phi) \simeq 3b [2a - f] \quad (26)$$

^{*)} The scale of a is thus basically given by $a^2 \sim M_W M_P$.

Recalling that $f < 0$ and $a < 0$, one observes that for $b > 0$, B becomes negative when $f > 2a$. It is remarkable that this happens during the period (25) when $A < 0$. This makes it possible for σ to remain in the $SU(5)$ breaking minimum also after the potential barrier around $\sigma = 0$ has reappeared. The existence of such a minimum depends on the magnitude of the σ^4 -term. This in turn can have contributions from the non-renormalizable part of the superpotential (15). We assume that these can be arranged in such a way that a minimum results. Then there will be no restoration of the $SU(5)$ symmetry. Instead, the inflaton field ϕ will roll to the global minimum defined by $g_\phi = 0$ and

$$g_\sigma = F_{2\sigma}(\sigma_0) + \sigma_0 [F_1(\phi) + F_2(\sigma_0)] = 0 \quad (27)$$

For $\sigma_0 \ll 1$, and recalling that $|F_1(\phi)| \ll |a|$, (27) is satisfied roughly by $F_{2\sigma}(\sigma_0) \approx 0$, or

$$\sigma_0 \approx - \frac{2a}{3b} \quad (28)$$

The condition $g_\phi = 0$ yields

$$f(\phi) \approx \frac{1}{2} a \phi_0 \sigma_0^2 \quad (29)$$

where the fact that $F_2(\sigma_0) \approx a\sigma_0^2/3$ has been used. Since the order of magnitude of $f(\phi)$ is $(\phi - \phi_0)$, the new minimum is reached only after the reappearance of the potential barrier around $\sigma = 0$.

If the potential barrier separating the minima $\sigma = 0$ and $\sigma = \sigma_0$ is not large enough, field fluctuations could drive the Universe back to the symmetric phase where there is the least amount of entropy. The energy of the inflaton field fluctuations is of the order of the inflaton mass $m_\phi \approx 10^{13}$ GeV. As this is also the scale of $f(\phi)$ at the end of inflation, (25) dictates the scale a to be of the same order of magnitude. Accordingly, this is also the scale of the barrier separating the broken and unbroken minima, and therefore field fluctuations are not of sufficient strength to force the Universe back to the symmetric phase.

Let us finally address the question of the direction of the $SU(5)$ breaking. Returning to (17), we proceed to estimate the value of the potential at the minimum σ_0 for $SU(3) \times SU(2) \times U(1)$ and $SU(4) \times U(1)$. The

parameter a is the same for both vacua [see Eq. (16)]. However, the parameter b is not. This gives rise to a deeper minimum in the $SU(3) \times SU(2) \times U(1)$ direction. Using (17), (21) and (28) we find that

$$\Delta \equiv V(\phi, \sigma_0) \Big|_{SU(3) \times SU(2) \times U(1)} - V(\phi, \sigma_0) \Big|_{SU(4) \times U(1)} = \frac{100}{729} \mathcal{J}(\phi) \frac{A^3}{B_{SU(3)}^2} < 0 \quad (30)$$

where $B_{SU(3)}$ denotes $B(\phi)$ as calculated in the $SU(3) \times SU(2) \times U(1)$ phase. By virtue of (30), we can argue that the $SU(3) \times SU(2) \times U(1)$ direction is naturally preferred over the $SU(4) \times U(1)$ direction.

Up to now, we have neglected all finite temperature effects¹³⁾. This has been justified if the scale a of the superpotential (15) is not much smaller than the Hawking temperature $T_H = H/2\pi \sim 10^{-7}$. In that case, finite temperature effects cannot lower the minimum in the $SU(4) \times U(1)$ direction below the minimum in the $SU(3) \times SU(2) \times U(1)$ direction, independently of the exact light field content of the full theory. The finite T corrections add to the potential a term $\Delta V_T = -(\pi^2/90) N(T) T^4$, where $N(T) = N_B + 7N_F/8$ measures the number of light degrees of freedom. In general possible differences induced by different light field contents of the $SU(3) \times SU(2) \times U(1)$ and the $SU(4) \times U(1)$ vacua are thus of the order of $\Delta V_T \approx 0(10)(\pi^2/90)T_H^4 \sim T_H^4$. Taking $A \sim a \sim T_H$ as a characteristic magnitude of A , with $\sigma_0 \approx 10^{-2}$ and assuming that the phase transition becomes possible when $\phi \approx 4$, one estimates from (30) that

$$\Delta \approx 10^{-24} \gg T_H^4$$

That is, finite T effects can be neglected, and we can conclude that the Universe will indeed go to the correct $SU(3) \times SU(2) \times U(1)$ preserving phase.

The exact moment of the phase transition depends on the behaviour of $P(\phi)$ around $\phi \approx \phi_0$, which is not directly connected to the inflatory epoch. It seems likely that the phase transition will happen during the reheating period, when the inflaton field is oscillating about the global minimum of the potential.

Since the temperature at which the GUT phase transition takes place is the Hawking temperature, magnetic monopoles will be diluted¹⁴⁾ by a factor T_H^3 . However, in contrast to previous considerations, the oscillations of the inflaton field take now place in the broken phase. As the scale of inflation, m^3 , is still present, the Universe will feel an effective non-zero cosmological constant. This will cause an additional dilution of the density of magnetic monopoles. After ϕ has settled to its new minimum, the generation of the baryon asymmetry can proceed in the usual way^{5),15)}.

Let us conclude by presenting the early history of the Universe as it has emerged from the present analysis. (i) Inflation begins in the SU(5) symmetric phase with $\sigma = 0$; (ii) ϕ approaches the global minimum ϕ_0 , and $\sigma = 0$ becomes unstable (A becomes negative); the phase transition into the broken minimum occurs rapidly with a natural preference for SU(3)×SU(2)×U(1) over SU(4)×U(1); (iii) ϕ keeps on moving towards the new minimum, while the minimum of the broken phase steepens (B becomes negative), and eventually a new barrier appears around $\sigma = 0$ (A becomes positive again); (iv) due to the barrier, σ remains trapped in the new minimum, while ϕ oscillates and finally settles into its new global minimum, completing the transition. It seems to us that the ease of the GUT phase transition is yet another attractive feature of the supergravity models with a flat potential.

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FIGURE CAPTION

A schematic form (not to scale) of the potential $V(\phi, \sigma)$, showing the barrier between the broken and unbroken minima at the global minimum ϕ_0 . The potential has been cut along line $\phi = \phi_0$ and in the (σ, ϕ) plane for simplicity. The dotted line represents the trajectory of the inflaton field.

FIG. 1