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PRIMORDIAL TWO-COMPONENT INFLATION

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ABSTRACT

We present a class of inflationable no-scale GUT supergravity models with the following characteristic properties: (i) superheavy ($\sim M_P$) or superlight ($<$ 1 keV) gravitinos avoiding the cosmological gravitino problem; (ii) a natural Higgs triplet-doublet splitting; (iii) proton decay mainly to kaons at an observable level $\tau_p \sim 10^{32\pm 1}$ years; (iv) sufficient cosmological baryon asymmetry due to the decays of superheavy Higgs triplets ($M_{H_3} \gtrsim 10^{15}$ GeV); (v) acceptable inflationary picture, where the physics responsible for the slow roll-over and energy density perturbations is decoupled from the the physics of reheating, which can be as high as $0(10^{16})$ GeV; (vi) rapid, second order GUT phase transition to the SU(3) × SU(2) × U(1) minimum triggered by and happening during inflation, allowing thus for enough GUT magnetic monopole dilution. The inflaton field knocks the GUT-breaking Higgs (guton) field out of its zero minimum. The field responsible for reheating will then be a mixture of the inflaton and the guton.

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1. - INTRODUCTION

The inflatory Universe[1),2)] appears to offer the only natural solution to the horizon and flatness problems of cosmology. By now, it is also evident that the early inflatory models based on grand unified theories (GUTs) with a Coleman-Weinberg-type symmetry breaking [3)] by radiative corrections have severe problems. In de Sitter space, quantum fluctuations cause density perturbations that are scale-invariant when they re-enter the Friedmann-Robertson-Walker horizon[4)]. This is a more than welcome feature, since this is just the Harrison-Zel'dovich[5)] spectrum of fluctuations needed for galaxy formation. Unfortunately, in inflation based on GUT phase transition, these perturbations turn out to be too large by several orders of magnitude[4)]. This problem, together with some fine-tuning problems, can be solved if one divorces inflation from GUTs and moves to supergravity. There inflation is driven by a GUT singlet inflaton field at the scale of the Planck mass M_P, where all fine tunings disappear[6)-8)]. The phenomenological models suitable for this purpose are naturally based on softly broken N = 1 supergravity[9)].

In supersymmetric inflation, the properties directly connected with inflation appear to be in good shape; namely, the roll-over time scale of the inflaton and the magnitude of density perturbations. However, when grand unification, supersymmetry and inflation are considered together, several problems appear. The questions related to grand unification are the production of the observed baryon asymmetry, and how to make the Universe go from the GUT phase to the correct SU(3) × SU(2) × U(1) phase. These difficulties arise because the overall scale of the inflaton potential, which is determined by the density perturbations, must be at most of the order of 10^{16} GeV. If the inflaton is decoupled from matter (except for gravitational interactions), the Universe cannot reheat after inflation to much more than about 10^{11} GeV. To accomplish the phase transition at that temperature, one has to arrange for a very small barrier between the symmetric and broken minima, and resort to strong coupling phenomena[10)]. One postulates that as the Universe supercools in the symmetric phase, the gauge coupling strength becomes 0(1) and the theory confines. This is assumed to lead to the disappearance of the light degrees of freedom of the symmetric phase, and thus the broken phase becomes favourable because of finite temperature effects. This possibility contains many uncertainties, not the least of them being the whole regime of strong coupling. It should also be stressed that the finite temperature behaviour that has been used to determine the onset of phase transition has, in fact, been derived assuming weak couplings.

The low reheating temperature T_R is not so problematic for mere baryogenesis. That can be arranged with low mass ($\sim 10^{10}$ GeV) Higgses. The problem is rather founded on the possibility that these same Higgses will also mediate catastrophically rapid proton decay. The low reheating temperature also has a supersymmetric angle. Namely, if the gauge hierarchy is to be protected by supersymmetry that is broken softly by the gravitino mass $m_{3/2} \sim M_W$, one can show[11)] that there is a stringent bound on T_R. This bound arises because the density of gravitinos, although initially diluted[12)] by inflation, will be regenerated to an unacceptably high value during the reheating unless $T_R \lesssim 10^9 - 10^{10}$ GeV.

A way to avoid the severe bound on T_R imposed by gravitinos is to decouple the scales of local and global supersymmetry breaking. This can be achieved in the no-scale supergravity models[13)-15)] where the mass of the gravitino may be either very large[14)] or very small[15)]. The scale of the necessary global supersymmetry breaking is set by the gaugino masses, which can be determined dynamically to be $O(M_W)$. These models may also explain[15),16)] why the QCD vacuum parameter θ is so small. When no-scale models are coupled to inflation[8)], it is possible to have a rapid, second order phase transition[17)] that actually prefers SU(3) × SU(2) × U(1) to SU(4) × U(1). Recently, also cosmological baryosynthesis, the questions of proton decay rate suppression and the splitting of the GUT and weak scales within the particle multiplets [the $\underline{5}$ of Higgses in SU(5), say] have been investigated[18)] in no-scale models. The result was a SUSY model that can accommodate baryosynthesis and proton lifetime compatible with experiment, together with a successful splitting of GUT and weak scales. In that model, the GUT scale is very close to the Planck mass.

In the present paper, we address the questions of how this large GUT scale can be reconciled with inflation and how the world can undergo a phase transition from the symmetric to the broken phase. We show that all models that are effectively globally supersymmetric at the scale of inflation, as no-scale models are, can lead to a rapid second order phase transition to the broken phase. In addition, we show that then naturally $T_R \sim 10^{16}$ GeV and $M_X \sim O(10^{-1})M_P$, and that despite the relatively high reheating temperature, the density of monopoles is suppressed because part of the inflation takes place in the broken phase.

2. - PARTICLE PHYSICS AND COSMOLOGY CONSTRAINTS

In this section, we describe in more detail the various particle physics and cosmological constraints relevant for a successful unification of supersymmetric GUTs and inflation. As is well known, the lure of supersymmetric GUTs is based on the fact that supersymmetry can protect the gauge hierarchy from radiative corrections, provided

$$m_B^2 - m_F^2 \simeq m_S^2 \lesssim \theta(M_W^2) \qquad (1)$$

where m_S is the scale of global supersymmetry breaking. The scale of local supersymmetry breaking is set by the gravitino mass $m_{3/2}$. Several phenomenological supersymmetry models[19] have been constructed by identifying $m_{3/2} \sim m_S$. Such a gravitino will usually decay to photons and photinos through d = 5 operators with a typical lifetime

$$\tau \simeq \frac{4M_P^2}{m_{3/2}^3} = 4 \times 10^8 \left(\frac{100\ \text{GeV}}{m_{3/2}}\right)^3 s \qquad (2)$$

Primordal gravitinos will therefore decay after nucleosynthesis, and they could be a possible cosmological embarrassment, unless their abundance can be suppressed. Although inflation dilutes the initial gravitino density[12), reheating produces a density[11)

$$n_{3/2}/n_\gamma \simeq \theta\left(T_R/M_P\right) \qquad (3)$$

The constraints on the density of gravitinos come from the requirements that gravitinos should not disrupt nucleosynthesis, produce too much entropy, distort the microwave background or dissociate light nuclei such as D, ^3He, ^4He and ^7Li. The most stringent bound is obtained from deuterium photodissociation, which implies[11)

$$\frac{T_R}{10^9\,\mathrm{GeV}} \lesssim 2\times 10^{10}\,\delta_B\left(\frac{100\,\mathrm{GeV}}{m_{3/2}}\right) \tag{4}$$

where $\delta_B = n_B/n_\gamma$ at the time of gravitino decay. For $\delta_B \simeq 10^{-9} - 10^{-10}$, one obtains the limit $T_R \lesssim 10^9 - 10^{10}$ GeV.

Such a limit on T_R implies that the mass M of the particles mediating baryon-violating interactions and participating in cosmological baryosynthesis should be very low. Clearly, if these particles are to be produced during reheating, we must have $M < T_R$. Barring some exotic possibilities[20] the only natural candidates for baryosynthesis are the Higgs triplets H_3. However, a light Higgs triplet will cause the proton to decay too rapidly because of d = 5 operators[18]. Indeed, to suppress these operators beyond experimental limits on proton lifetime, one would need[21] $M_{H_3} \gtrsim 10^{15}$ GeV.

As far as particle physics is concerned, a way out of these conflicting requirements was pointed out in Ref. 18). As stressed there, a necessary ingredient of the solution appears to be the decoupling of local and global supersymmetry breaking scales. This is indeed possible in the no-scale supergravity GUTs[13)-15)], where the scale of global supersymmetry breaking is set by the gaugino masses m_V. These can arise at tree level because of the non-minimal kinetic terms found in supergravity. It has been shown that while $m_V \sim O(M_W)$, the gravitino mass may be either very large[14], $m_{3/2} \sim M_P$, or very small[15]; for instance, in one particular example, $m_{3/2} \sim m_V^p/M_P^{p-1}$, $1 < p < 2$. Therefore the severe bound (4) on T_R can be avoided.

In these models, the world is effectively globally supersymmetric at the high mass scale. The scale of global supersymmetry breaking, together with the electroweak scale, is determined dynamically by radiative corrections. In Ref. 18), it was discovered that it is indeed possible to write down such a globally supersymmetric SU(5) model that can give rise to the observed baryon asymmetry while suppressing baryon decay through d = 5 operators. This model employs the missing partner mechanism[22),23)] to achieve a natural doublet-triplet splitting. Its superpotential is

$$f = \frac{\alpha_{ij}}{M_P^4}\,\bar\theta_i\theta_j\Sigma^5 + \frac{\beta_{ij}}{M_P^2}\,\bar\theta_i H_j\Sigma^2 + \frac{\gamma_{ij}}{M_P^2}\,\theta_i\bar H_j\Sigma^2 \tag{5}$$

$$(i,j = 1,2)$$

where $\theta_i \sim \underline{50}$, $H_i \sim \underline{5}$ and $\Sigma \sim \underline{24}$. The form of the superpotential (5) is a consequence of a discrete symmetry imposed on f, under which the fields transform as $\phi_k \to e^{i\alpha_k \pi}\phi_k$ (here $\phi_k = \theta_i, \bar{\theta}_i, H_i, \bar{H}_i, \Sigma$), where α_k is some (irrational) number. The model (5) gives sufficient baryon asymmetry provided the GUT scale $M_X \sim M_P$.

Let us now turn to inflation. In the above model, the masses of Higgs triplets are $O(M_X^2/M_P)$. On the other hand, after inflation, the coherent inflaton field oscillations dominate the energy density of the Universe leading to the reheating temperaure

$$T_R = (\Gamma_\phi M_P)^{\frac{1}{2}} \qquad (6)$$

where $\Gamma_\phi \lesssim H$ is the decay rate of the inflaton, where $H = \dot{R}/R$ is the Hubble constant during inflation. If the inflaton couples to matter only via gravitational interactions, $\Gamma_\phi \sim m_\phi^3/M_P^2$. The overall scale $m_0 \simeq 10^{16}$ GeV of the inflation potential is fixed by the density perturbations. Indeed, it has been shown[24)] that one obtains a stringent upper limit $m_0 << 8.9 \times 10^{16}$ GeV from the quadrupole anisotropy of the microwave background, which is determined by energy density perturbations that re-entered the horizon during the matter-dominated era. Therefore the mass of an inflaton that is decoupled from other fields is small, $m_\phi \simeq m_0^2 / M_P << m_0$. This leads to the low reheating temperature $T_R \simeq 10^{10} - 10^{11}$ GeV, which appears to be at variance with the particle physics constraints discussed above.

As emphasized in the Introduction, the low reheating scale is also associated with the difficulty of effecting the GUT phase transition in the first place. As we will show in the the following sections, a natural solution to these problems is to divorce the period of slow roll-over from the period of reheating by making the mass scales governing the two epochs totally disjointed. Indeed, it should be stressed that, for example, the density perturbation spectrum is not sensitive[25),26)] to the details of reheating, and therefore inflation and reheating are, in fact, two disconnected phenomena. This separation can be accomplished if we assume, as seems natural to us, that the inflation is not decoupled from GUT fields.

3. - TWO-COMPONENT INFLATION (I)

We will consider for definiteness an SU(5) SUSY GUT with a single adjoint field Σ. It is sufficient to consider the effective potential only in the direction of vanishing D-terms. We assume that, as in no-scale models, at the onset of inflation the effective potential for the inflaton ϕ and the $SU(3) \times SU(2) \times U(1)$ singlet σ of 24 can be written as

$$V_{eff} = E(|\phi|,\sigma)\left[|F_\phi|^2 + \frac{1}{2}\kappa^2(|\phi|)|F_\sigma|^2\right] \tag{7}$$

Here $F = F(\phi,\Sigma)$ is the superpotential, and E and κ are some positive scaling functions which can be present if the effective theory has been derived from local supersymmetry with non-minimal kinetic terms. When evaluated at the global minimum, they will give rise to rescaling of the (effective) superpotential. For instance, the no-scale model for inflation[8),17)] given by the Kähler potential

$$G = -3\ln\left(z + z^* + h(|\phi|)\right) + g(|\phi|) + |\sigma|^2 + \ln|F|^2$$

will result in

$$E = -\frac{1}{4}\frac{g_{\phi\bar\phi}}{h^3_{\phi\bar\phi}}\exp\left(g + |\sigma|^2\right)$$

$$\kappa^2 = -\frac{1}{2} g_{\phi\bar\phi}$$

Here $g_{\phi\bar\phi} = (\partial^2 g/\partial\phi\partial\phi^*) < 0$, $h_{\phi\bar\phi} < 0$, and z is the field breaking supersymmetry. Apart from the scaling functions E and κ^2, our effective potential is globally supersymmetric. Of course, it may turn out that in the correct model $E = 1$ and $\kappa^2 = 2$. In any case, in the following we will make the natural assumption that below M_P, E and κ are smooth, slowly-varying functions. For most purposes, these scaling functions can be taken to be constants, and the following considerations are not sensitive to their precise forms.

The global minima of the potential (7) are given by

$$F_\phi = F_\sigma = 0 \tag{8}$$

In addition, there should exist a local minimum with a positive cosmological constant at $\phi = \sigma = 0$. Inflation takes place when ϕ starts to roll out from this minimum, and reheating happens when fields oscillate about the global minimum (8) where the cosmological constant $\Lambda = 0$.

The effective superpotential $F(\phi,\sigma)$ is a priori an arbitrary analytic function. Because σ is an SU(5) non-singlet, high temperature effects will force σ to be at $\sigma = 0$ when $T \sim M_p$, independently of the form of $F(\phi,\sigma)$. Therefore it is natural to expand $F(\phi,\sigma)$ about this point to obtain

$$F(\phi,\sigma) = A(\phi) + B(\phi)\sigma^2 + C(\phi)\sigma^3 + \cdots \tag{9}$$

assuming that in terms of Σ, $F(\phi,\Sigma)$ can be expanded as $F = \sum_{n=2} a_n(\phi)\mathrm{Tr}\Sigma^n$. Here, A, B and C are unknown functions. The decoupling of inflation from the GUT field would be achieved if B and C were constants. We see no physical reason why this should happen. Therefore we take B and C [and other higher order coefficients in (9)] to be reasonably smooth and well-behaving functions of ϕ. With the expansion (9), the effective potential (7) can be written as

$$V_{eff} = E\Big[\,|A_\phi|^2 + 2\,\mathrm{Re}(A_\phi^* B_\phi \sigma^2) + 2\kappa^2|B|^2|\sigma|^2 \tag{10}$$
$$+ 6\kappa^2|\sigma|^2\,\mathrm{Re}(B^* C\sigma) + 2\,\mathrm{Re}(A_\phi^* C_\phi \sigma^3) + \cdots\Big]$$

Here, $A_\phi = \partial A/\partial\phi$, etc.

Higher order terms in expansion (9) are not essential for our discussion and we do not write them down explicitly. At high temperatures $\sigma = 0$, and V_{eff} collapses to

$$V_0 = E|A_\phi|^2 \tag{11}$$

It is evident that the form of V_0 must be such that a sufficient amount of inflation results. This means that the scale of $V_0 = m_0^4$ must be $m_0 \simeq 10^{16}$ GeV to produce correct density fluctuations. We take it for granted that $A(\phi)$ can be chosen in such a way that it gives successful inflation. The various constraints for achieving enough inflation have been extensively discussed in the literature[7),8),25)], and we assume that these conditions can be met.

Suppose now that ϕ starts to roll into the real ϕ direction and that F is a real function. This is no essential restriction, but it simplifies the following discussion considerably. Defining $\sigma = (1/\sqrt{2})(\sigma_R + i\sigma_I)$, the effective potential is then given by

$$V_{eff} = E\Big[|A_\phi|^2 + (\kappa^2 B^2 + A_\phi B_\phi)\sigma_R^2 + (\kappa^2 B^2 - A_\phi B_\phi)\sigma_I^2 + \frac{1}{\sqrt{2}}(3\kappa^2 BC + A_\phi C_\phi)\sigma_R^3 + \frac{1}{\sqrt{2}}(3\kappa^2 BC - 3A_\phi C_\phi)\sigma_R\sigma_I^2 + O(\sigma^4)\Big] \quad (12)$$

We assume that the only σ-dependence in E is of the form $E = E_0(|\phi|)e^{\xi|\sigma|^2}$, where $\xi = 0$ or 1. This is a reasonable assumption if σ is to have canonical kinetic terms. The masses of σ_R and σ_I are then

$$\begin{aligned} m_R^2 &= E_0\left[\kappa^2 B^2 + \xi A_\phi^2 + A_\phi B_\phi\right] \\ m_I^2 &= E_0\left[\kappa^2 B^2 + \xi A_\phi^2 - A_\phi B_\phi\right] \end{aligned} \quad (13)$$

We can assume without loss of generality that $A_\phi \leq 0$ in the interval $0 \leq \phi \leq \phi_0$, where ϕ_0 is the point where the inflaton would roll to in the absence of σ. However, this need not be a global minimum, i.e., it may be that $|A_\phi(\phi_0)| > 0$. Although it may be that $\phi_0 = \infty$, we usually take $\phi_0 \simeq O(M_P)$. In Fig. 1, we show a schematic illustration of the inflaton potential in the ϕ direction.

If $B_\phi \geq 0$ in the interval $0 \leq \phi \leq \phi_0$, we see that $m_I^2 \geq 0$ always. However, as B is basically an arbitrary function, it is natural to assume that it can have one (or several) zero(s) ϕ_* with $0 \leq \phi_* \leq \phi_0$. But then, $m_R^2 < 0$ if

$$\xi A_\phi^2(\phi_*) + A_\phi(\phi_*)B_\phi(\phi_*) < 0 \quad (14)$$

signalling the breaking of the SU(5) symmetry and facilitating a rapid phase transition to the SU(3) × SU(2) × U(1) phase. Note that if $\xi = 0$, (14) is satisfied automatically. The value of ϕ_* depends on the functional form of $B(\phi)$, and thus the phase transition can happen during any stage of inflation, not only at the end of it, as in Ref. 17). This means that it is possible that the Universe inflates partly in the broken phase. This will effectively suppress the density of the monopoles produced by the phase transition, as in the old Coleman-Weinberg type inflation.

The condition (14) for symmetry breaking can be satisfied whenever $B_\phi(\phi_*) \gtrsim \left|A_\phi(\phi_*)\right|$. By naturalness, one would argue that this is possible in <u>all</u> models having the effective potential (7). Indeed, this means that in <u>all</u> models having an effective global supersymmetry at the scale of inflation, the GUT phase transition happens during the inflation, is rapid and of the second order.

4. - TWO-COMPONENT INFLATION (II)

As has been repeatedly stressed in this paper, the scale of inflation or V_0 [see Eq. (11)] is fixed by density perturbations to be roughly $m_0^4 \simeq$ $\simeq (10^{16}\ \mathrm{GeV})^4$. In the present approach, there are two natural possibilities for achieving this:

(i) $A \sim B \sim C \sim \mathcal{O}(1)\ ;\ E_0 \sim m_0^4$

(ii) $A \sim m_0^2\ ;\ B \sim C \sim E_0 \sim \mathcal{O}(1)$

(in natural units $M_P/\sqrt{8\pi} = 1$). We take the above to be overall scales multiplying the functions in question. Otherwise, we expect these functions to be typically polynomials with coefficients of the order of 1. The mass of the σ-field as a function of ϕ for both cases is depicted in Fig. 2. Although for inflatory purposes these two cases are equivalent, they will lead to very different physics at the end of inflation. In both cases, finite T effects push σ to origin, where inflation starts, and in both cases $B(\phi) \simeq 0$ makes SU(5) breaking phase transition possible, after which the fields will roll to the broken global minimum (8). It is also evident that at that minimum $\sigma_0 = M_X \sim 0(1)$ and $\phi_0 \simeq \sigma_0$.

Because of the overall scale m_0, the masses of ϕ and σ in case (i) are necessarily of the order of $m_0^2/M_P \sim 10^{13}$ GeV. It is not probable that Higgs triplets could be this light. Therefore, in case (i), the question of baryogenesis remains open. On the other hand, in case (ii), the scale $m_0 << 1$ can be neglected at the global minimum. Therefore, case (ii) gives rise to <u>a true two-component description of inflation, where the periods of slow roll-over and reheating are governed by two totally different scales</u>. In case (ii), the field oscillating about the global minimum is not the inflaton but a mixture of ϕ and σ. The mass eigenstates at the global minimum (8) are easily calculated and are given for case (ii) as

$$M_{1,2}^2 = \frac{1}{2} F_{\phi\sigma}^2 \left[1 + 2x \pm (1+4x)^{1/2} \right] \tag{15}$$

where we have assumed a real minimum and where $x \equiv (F_{\sigma\sigma}/F_{\phi\sigma})^2$. We expect $F_{\phi\sigma}$ and x to be $O(1/10)$-$O(1)$, or $M_{1,2} \sim O(1/100)$-$O(1)$. In this case, the Universe will reheat to $T_R \simeq (HM_P)^{1/2} \simeq m_0$. As we will show later in the context of a concrete model, this reheating is sufficient to generate the cosmological baryon asymmetry.

Let us examine more closely the period during which the phase transition takes place. We expand the functions A and B about the zero ϕ_* of B as

$$B = b(\phi - \phi_*)$$

$$A = a_0 + a_1(\phi - \phi_*)$$

In case (i), all coefficients are naturally of the order 1, whereas in case (ii), $b \sim O(1)$ and $a_0 \sim a_1 \sim O(m_0^2)$. Inserting the above form into Eq. (13), and solving $m_R^2 = 0$ for $\phi - \phi_*$, one finds in case (i)

$$|\phi - \phi_*| \sim O(1)$$

That is, the period during which the SU(5) phase transition is possible is very long, of the order of M_P. Therefore, the phase transition would occur very rapidly after ϕ has tunnelled out of the metastable minimum at $\phi = 0$, and most of the inflation should occur in the broken phase. Although this would be very effective in diluting the density of magnetic monopoles, the low reheating temperature makes case (i) unrealistic.

In case (ii), one finds

$$|\phi - \phi_*| \simeq \frac{m_0}{\kappa b^{1/2}} \simeq 10^{-2} - 10^{-3} \tag{16}$$

In this case, the period when $m_R^2 < 0$ is very short compared to the roll-over scale. One may even wonder whether it is too short for field fluctuation to trigger the beginning of the phase transition. Indeed, the semiclassical evolution of the inflaton field is given by

$$\ddot{\phi} + 3H\dot{\phi} + \Gamma\dot{\phi} + V'(\phi) = 0 \tag{17}$$

During the slow roll-over phase, $\ddot{\phi}$ and $\Gamma\dot{\phi}$ can be neglected. The time scale of the period when $m_R^2 < 0$ can then be evaluated from Eqs. (16) and (17) to be

$$\Delta t \sim \frac{3H|\phi-\phi_*|}{V'(\phi_*)} \sim m_0^{-1}$$

This is much shorter than the time scale of thermal fluctuations due to the Hawking temperature $T_H = H/2\pi$, which is $(\Delta t)_{th} \simeq H^{-1} \simeq m_0^{-2}$. Therefore the symmetry breaking opening to the $\sigma \neq 0$ direction is too narrow for thermal fluctuations to "see" it. However, because the curvature of the potential at $\phi \simeq \phi_*$ in the σ-direction is, from (13), $O(m_0^2)$, the σ-field is subject to quantum fluctuations that occur with a typical time scale $(\Delta t)_q \sim m_0^{-1}$. These quantum fluctuations are sufficiently rapid for triggering the beginning of the GUT phase transition. In effect, in addition to ϕ, also σ starts rolling.

5. - A TOY MODEL

The evolution of ϕ and σ, induced by the potential (10), is in general quite complicated. Exact quantitative statements would require computer simulations. However, some aspects of the behaviour of the effective potential can be studied using simplified models. In particular, we wish to show that it is indeed possible for the fields to roll to the global minimum, without being trapped in some local minimum with a positive cosmological constant.

We consider two-component inflation with the superpotential of the form

$$F(\phi,\sigma) = A(\phi) + b(\phi-\phi_1)\sigma^2 + c(\phi-\phi_2)\sigma^3 \qquad (18)$$

where b, c, ϕ_1 and ϕ_2 are real constants of the order of 1, and $A \sim m_0^2$. As explained in Section 3, it is sufficient to consider real σ only, provided we take $b > 0$. We restrict ourselves to real ϕ, and make a further simplification by taking the scaling functions to be $E(\phi,|\sigma|) = 1$ and $\kappa^2(|\phi|) = 2$. This means that our effective potential is globally supersymmetric.

The history of the Universe starts in this toy model from $\phi = \sigma = 0$. Eventually, we expect the Universe to land at the global minimum $F_\phi = F_\sigma = 0$.

From (18) it is given by

$$\sigma_0 = M_X \simeq -\frac{b}{c}$$
$$\phi_0 = 3\phi_2 - 2\phi_1 \tag{19}$$

At the global minimum, the scale m_0 has disappeared from the potential, and therefore reheating is independent of the slow roll-over period.

The path to the global minimum is, however, rather complicated. Let us take for simplicity $c = -1$, $\phi_2 = 2\phi_1$. Then the effective potential for real σ and ϕ is

$$V_{eff} = A_\phi^2 + 2\left[A_\phi M_X + 2M_X^2(\phi-\phi_1)^2\right]\sigma^2 - 2\left[A_\phi + 6M_X(\phi-\phi_1)(\phi-2\phi_1)\right]\sigma^3$$
$$+ \left[M_X^2 + 9(\phi-2\phi_2)^2\right]\sigma^4 - 2M_X\sigma^5 + \sigma^6 \tag{20}$$

When $\phi \simeq \phi_1$, the mass of σ becomes negative (we take $A_\phi < 0$). With $A_\phi \simeq -m_0^2$, the corresponding minimum in the σ-direction is at $\phi = \phi_1$

$$\sigma_1 \simeq \left[\frac{M_X}{M_X^2 + 9\phi_1^2}\right]^{\frac{1}{2}} m_0 + O(m_0^2) \tag{21}$$

with positive vacuum energy of the order of m_0^4. However, this minimum is not stable in the ϕ-direction:

$$V_\phi(\phi_1, \sigma_1) \simeq 12 M_X \phi_1 m_0^3 + O(m_0^4)$$

Note that this means that ϕ actually starts to roll back towards $\phi = 0$. The curvature is $V_{\phi\phi} \simeq 0(m_0^2)$, so that the potential is not flat enough for additional inflation.

Let us now consider the region $|\phi - \phi_1| >> m_0^2$. There we can neglect $F_{\phi\phi} = A_{\phi\phi}$, and the extrema are given by

$$V_\sigma = F_{\phi\sigma}F_\phi + F_{\sigma\sigma}F_\sigma = 0$$
$$V_\phi = F_{\phi\sigma}F_\sigma = 0 \tag{22}$$

In our toy model, $F_{\phi\sigma} = F_\sigma = 0$ cannot be satisfied simultaneously, and hence there is no corresponding stationary point. The case $F_{\sigma\sigma} = F_{\phi\sigma} = 0$ corresponds to a stationary point at $\phi = 3\phi_1$, $\sigma = (2/3)M_X$. However, calculating the determinant of the mass matrix, one finds

$$\det M^2 = -V_{\phi\sigma}^2 < 0$$

That is, this extremum is a maximum in some direction and not stable. This leaves us with only one stable extremum, $F_\phi = F_\sigma = 0$, which is just the global minimum. Therefore, whatever adventures the fields may have in between, in the present model they will finally end at the global minimum. It is also evident that there will be some additional inflation in the broken phase; exactly how much is very model dependent.

A possible trajectory of the inflaton field, which is a superposition of ϕ and the guton σ, is shown in Fig. 3. A birds' eye view of this trajectory will generally resemble a meandering river that flows through valleys overlooked by mountains.

Let us stress that the temperature at the time of the phase transition is the Hawking temperature $T_H = H/2\pi$. Therefore, monopole to photon ratio is suppressed by a factor $T_H^3 \sim 10^{-21}$ if $m_0 \sim 10^{-3}$. Additional suppression will be obtained while the Universe inflates in the broken phase. Hence, it is possible to satisfy the limits on the monopole abundance coming from the overall density of the Universe[27)] and even the stringent limits coming from the more uncertain monopole catalyzed proton decays in neutron stars[28)]. There is no need to invoke a very low reheating temperature ($T_R \lesssim 10^9$ GeV) to suppress monopoles. Note also that in the present case, the possibility exists that there is a monopole flux that could be observable in the near future.

If we take $\phi_1 \simeq 1$, the mass eigenstates of the ϕ,σ mass matrix are given from (15) by

$$M_1^2 \simeq 72 M_X^2 \simeq M_P^2$$
$$M_2^2 \simeq \frac{3}{8} M_X^4 \simeq m_0^2 \qquad (23)$$

if $M_X \simeq 0.1$. If we adopt also the model given in Eq. (5), the baryon asymmetry will be[18)] proportional to $(\alpha/\beta)^4 M_X^{12} \times 0(10^{-2})$, which will then be of correct magnitude if $\beta \sim 0.1$, $\alpha \sim 1$. The inflaton, being now a superposition of ϕ and σ, will couple to Higgs triplets through the superpotential couplings (5). Note that it does not couple to matter fermions. Therefore, the lightest eigenstate which will be produced by the coherent field oscillations can decay predominantly to Higgs triplets.

We must still ascertain that the produced Higgs triplets will be out of equilibrium. From (5), we find that the ΣH^2 coupling has a strength M_X^3. This, then, will be the order of magnitude for the ϕH^2 coupling also. The triplets will decay out of equilibrium if

$$\frac{T_D}{M_{H_3}} \simeq \frac{\Gamma_\phi^{1/2}}{M_X^2} << 1 \tag{24}$$

where T_D is the decay temperature of the inflaton, and $M_{H_3} \simeq M_X^2$. From (23), we find that $\Gamma_\phi \simeq M_X^8$ and therefore the bound (24) is satisfied if $M_X \simeq 0.1$. With a heavy Higgs mass $M_{H_3} \sim 10^{15}-10^{16}$ GeV, the produced baryon asymmetry cannot be erased[28)] by the potentially dangerous $2 \leftrightarrow 2$ scattering. As stressed earlier, such heavy Higgs triplets are sufficient for rendering the dangerous dimension five operators harmless. The proton decay then proceeds in the standard way[21)], producing mainly kaons with a lifetime $\tau_p \simeq 10^{32\pm 1}$ years, which can hopefully be observed in the near future.

6. - CONCLUSIONS

In the present paper, we have argued that a solution to discrepancies between cosmological inflation and constraints coming from particle physics is to separate the scale of the slow roll-over period from the scale of reheating. We showed that this separation occurs naturally in no-scale models of supergravity[13)-15)], where the scales of global and local supersymmetry breaking are decoupled. This decoupling allows one to avoid the potentially disastrous limit on the reheating temperature[11)] obtained in the case of the gravitino with mass of the order of the weak scale. Indeed, in no-scale models, the gravitino may be either very heavy[14)] or very light[15)]. As a consequence, we were led to consider primordial inflation in the context of effectively global supersymmetry.

We have demonstrated how inflation is responsible for triggering a rapid second order phase transition from SU(5) to SU(3) × SU(2) × U(1) during inflation. After that, the inflaton is, in fact, a superposition of the SU(5) breaking Higgs (guton) field and the old inflaton field. Part of the inflation takes place in the broken phase giving an effective suppression of monopole abundance. At the end of inflation, the scale characterizing the slow roll-over epoch has disappeared, and consequently the Universe may reheat to $T_R \sim 0(10^{16})$ GeV. This is essentially a two-component picture of inflation.

Furthermore, we considered a toy model where we showed that fields can indeed roll to the global minimum, where SU(5) is broken, without being trapped in any local minimum. In general, we expect the trajectory of the inflaton field from the origin to the global minimum to be a very complicated one. Our toy model is a proof of existence for such a trajectory.

We feel that two-component inflation is a promising candidate for unifying all the successful features of GUTs and cosmological inflation. It can lead to a reheating high enough to be compatible with both the limits on proton stability and the generation of the cosmological baryon asymmetry by superheavy ($M_{H_3} > 10^{15}$ GeV) Higgs triplets. Indeed, in two-component inflation, one can accommodate GUTs, supersymmetry and early cosmological inflation in one single consistent picture.

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FIGURE CAPTIONS

Fig. 1 : A schematic form of the inflaton potential. There need be no minimum of the potential for $\phi > \phi_*$, where ϕ_* is the location of the zero of the coefficient B [see Eq. (9)].

Fig. 2 : The mass squared $m_R^2(\sigma)$ of the SU(5) adjoint field as a function of ϕ: a) in case (i), b) in case (ii) (see text).

Fig. 3 : A possible "snake river" trajectory of the inflaton field in the (ϕ,σ) plane. The fields roll from $\phi = \sigma = 0$ to the global minimum at $\sigma = M_X$, $\phi = \phi_0$.

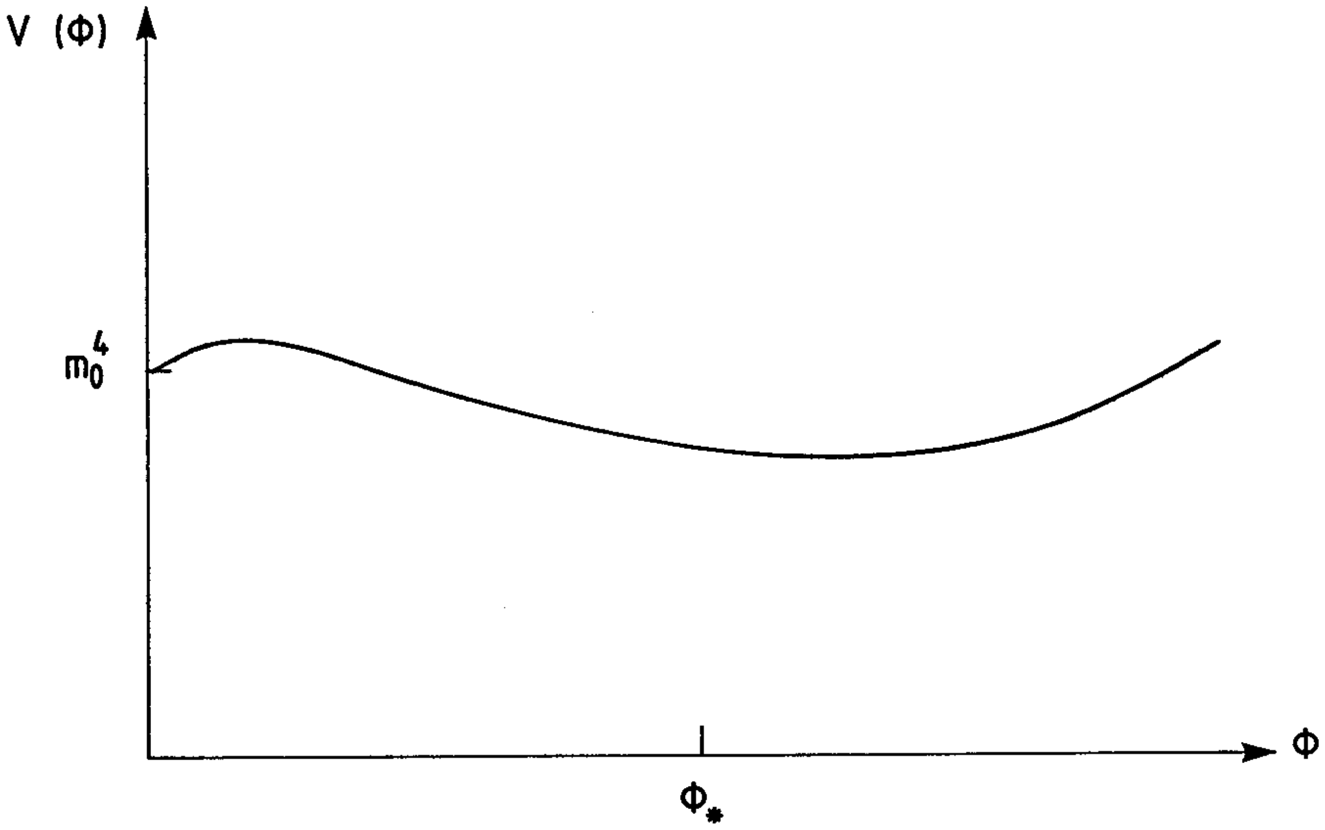

Fig. 1

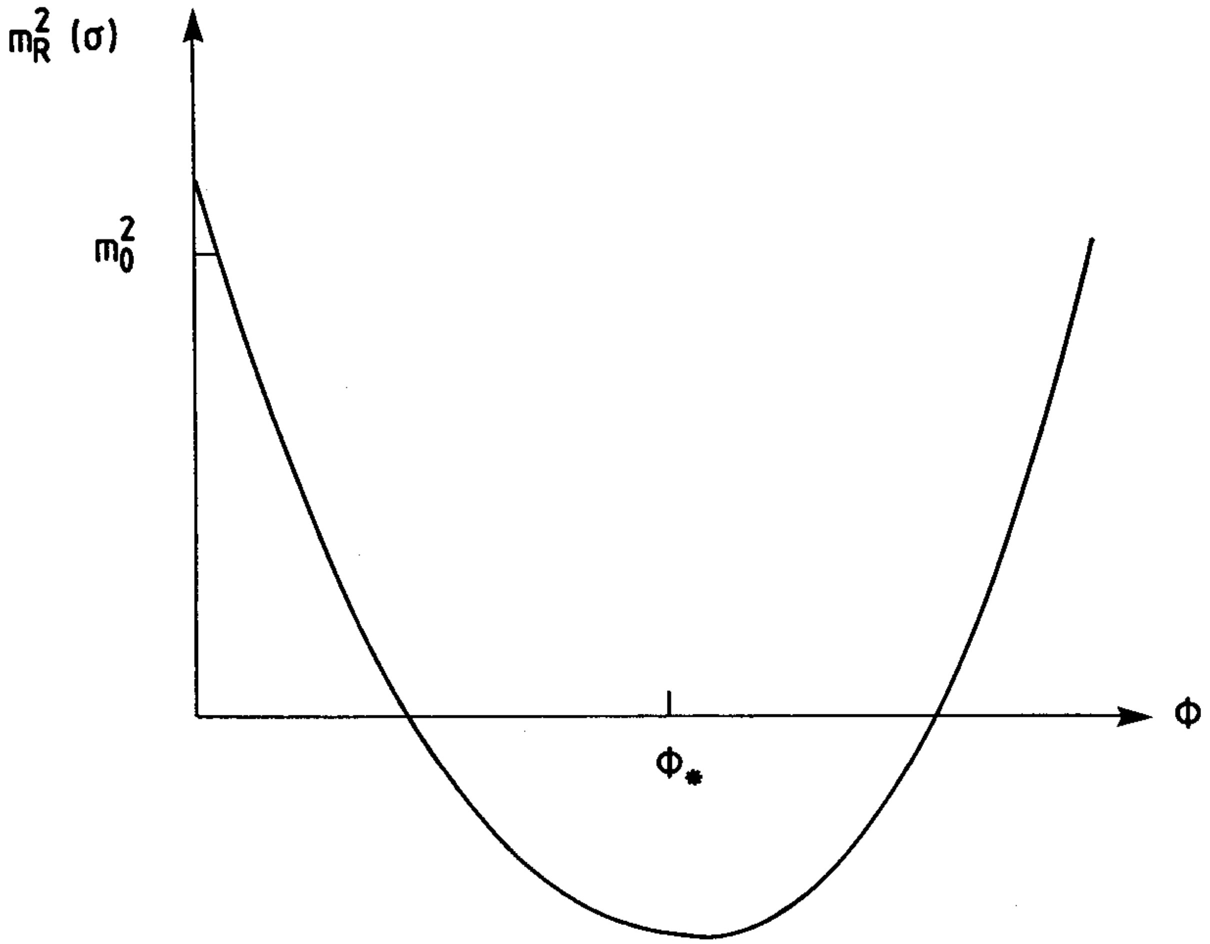

Fig. 2a

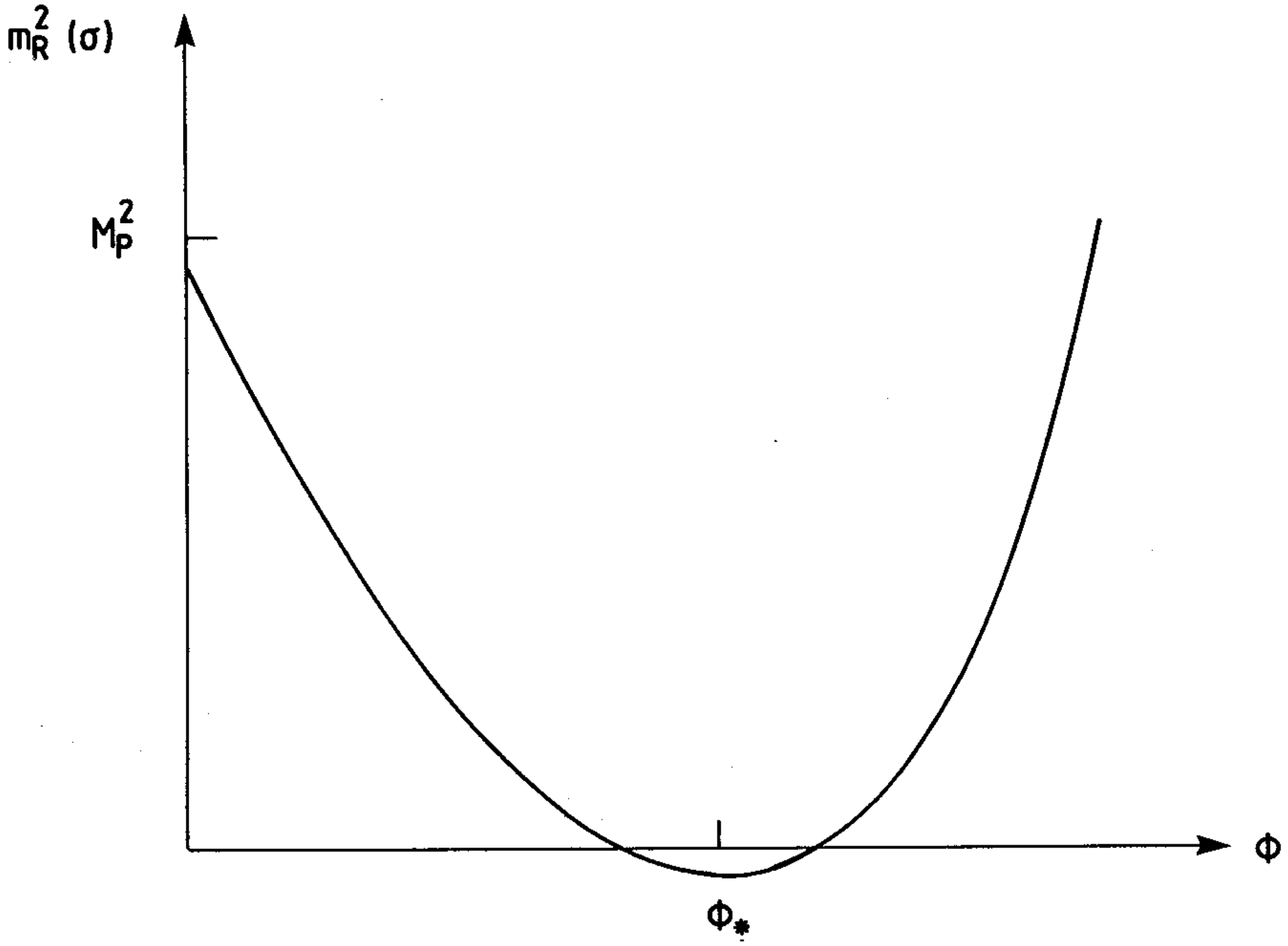

Fig. 2b

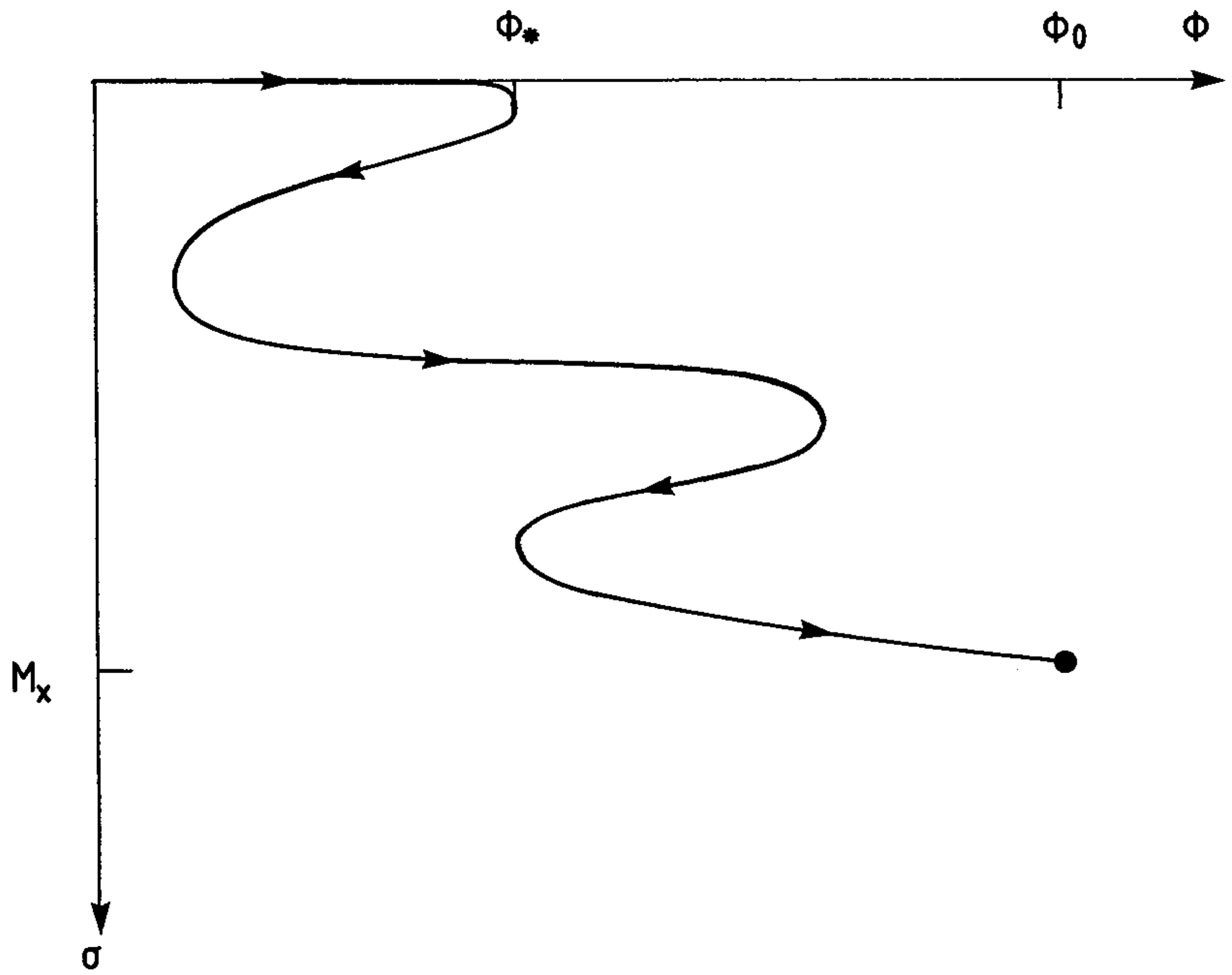

Fig. 3