CERN-TH.4054/84

COSMIC SUPERGRAVITY

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ABSTRACT

Supergravity seems to play a very important role in particle physics. This paper is devoted to recent cosmological applications of supergravity which seem to indicate that supergravity is of prime importance also in cosmology.

Invited Talk given at the
Bielefeld Conference
"Phase Transitions in the Very Early Universe"
4-8 June, 1984

CERN-TH.4054/84
November 1984

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1. INTRODUCTION

Local supersymmetry (SUSY) or supergravity (SUGRA) seems, more and more, to play a fundamental and key role in particle physics. If this is indeed so, it is interesting to discuss the possible cosmological consequences of SUGRA. This paper is thus devoted to cosmic SUGRA. Section 2 analyses certain potential problems of the standard Big Bang Cosmology (BBC) and of SUGRA models, whilst Section 3 discusses general ways of resolving these problems through inflation, and also sets the rules for its success. In Section 4 the different steps towards realistic primordial (SUGRA) inflation are explained and the construction of healthy cosmic SUGRA models is indicated.

2. COSMO-PARTICLE MYSTERIES

A complete physical understanding of our Universe requires a successful cosmological and particle physics model. On the cosmological front the hot Big-Bang model¹ is the best case at hand, whilst in particle physics $N = 1$ SUGRA Yang-Mills field theories² ($\equiv$ SUGRA models) seem to be the most promising candidates for an effective theory describing all particle physics at energy scales below the (super) Planck scale

$$M = \frac{m_p}{\sqrt{8\pi}} = \frac{1}{\sqrt{8\pi G_{\text{Newton}}}} \approx 2.4 \times 10^{18} \text{ GeV} , \quad (1)$$

which will be henceforth set equal to 1.

The observed universal (Hubble) expansion, the existence of the $T_0 \approx 2.7 \text{ K}$ cosmic microwave background radiation, and the observed abundances of ${}^4\text{He}$, D , etc., serve as the milestones in the hot Big Bang Cosmology (BBC). The credentials of SUGRA models include²: automatic unification of electroweak and strong interactions with gravity; a possible solution of the threefold gauge hierarchy problem³, namely the dynamical origin, the stability at the tree-level⁴, as well as, to all orders in perturbation theory, of the tiny ratio:

$$\frac{m_W}{M} \sim 10^{-16} \quad (2)$$

(m_W denotes the electroweak scale $m_W \approx 100$ GeV); and a possible solution⁵ of the strong CP hierarchy problem

$$\theta_{\text{QCD}} < 10^{-9} . \quad (3)$$

However, neither the BBC model nor the SUGRA model is entirely problem-free. Indeed, the observed homogeneity and isotropy of the present Universe at large scales ($d \gg 100$ Mpc), its inhomogeneity at small scales ($d < 100$ Mpc), its flatness, the predominance of matter over antimatter, and the absence of a cosmological constant are big mysteries in the BBC. The SUGRA models have not only inherited certain problems from (global) SUSY Grand Unified Theories (GUTs), such as the well-known magnetic monopole⁶ and GUT phase transition problems⁷, but they have also created their own problems such as the gravitino⁸ and the Polonyi problems⁹. The gravitino, the spin-3/2 SUSY partner of the graviton (spin-2), is the gauge fermion of local SUSY, whilst the Polonyi field z is the spin-0 field that causes the spontaneous breakdown of local SUSY (super-Higgs effect). The common origin of the gravitino and Polonyi problems lies in their tiny, gravitational couplings to themselves or to ordinary matter. In the case of the gravitino, these gravitational couplings to ordinary matter make the gravitino very long-lived, so that when it decays (and if it is not stable) it causes certain cosmological trouble. For example, it can disturb catastrophically primordial nucleosynthesis, if its mass lies in the region⁸

$$1 \text{ keV} < m_{3/2} < O(10^4 \text{ GeV}) . \quad (4)$$

The lower bound of the forbidden region corresponds to stable gravitinos whose energy density should not exceed ρ_{critic} , whilst the upper bound is coming from the primordial nucleosynthesis constraint. Notice that $m_{3/2} \approx m_W$, which is the gravitino mass in conventional SUGRA models², falls just in the middle of the forbidden range (4). This is the gravitino problem⁸. In the case of the Polonyi field z , its tiny gravitational couplings to itself or to matter are responsible for the extremely long time needed for this field to reach its minimum and to dispose of its energetic content -- again creating cosmic uneasiness. In a typical SUGRA model, the Polonyi field is, even today, still oscillating about its minimum, and its coherent energy density would exceed the present critical energy density by 16 orders of magnitude! This is the Polonyi problem⁹.

A successful cosmo-particle physics model should explain all these cosmo-particle mysteries.

3. COSMIC INFLATION

A first step towards a partial resolution of our cosmo-particle mysteries is clearly the use of Cosmic Inflation¹⁰. Inflation -- a period of sufficiently long exponential expansion of the cosmic scale factor R in the very early Universe -- automatically resolves, in principle, the problems of large-scale homogeneity and isotropy, of flatness, and, in some cases, of the magnetic monopole and the gravitino¹¹. A sufficient exponential growth of the volume of the Universe during very early times makes it a possibility that the whole present observable Universe was once, in the far distant past, contained in a causally connected region, so that the observed homogeneity and isotropy may be explained in terms of processes of causal microphysics. In addition, such an exponential growth of the cosmic scale factor R completely diminishes the importance of the curvature term k/R^2 so that, independently of the initial value of k (+1: closed; 0: flat; -1: open Universe), we live today in an effectively flat ($k \approx 0$) Universe, thus resolving the flatness problem. Furthermore, such a blow-up of the cosmic scale factor R can potentially dilute completely any topological defects such as magnetic monopoles, or any cosmologically undesirable particles such as gravitinos¹¹. Needless to say, extra care should be taken so that after inflation, gravitinos should not be abundantly regenerated¹²⁻¹⁴, whilst the classic GUT cosmic-baryon asymmetry mechanism¹⁵ should be still applicable. In any case, cosmic inflation¹⁰ seems to be an unavoidable ingredient of our cosmo-particle model, and we show next how to implement it.

A period of exponential expansion of the cosmic scale factor R corresponds to a period of constant cosmic energy density ρ , which is realizable if the Universe, during the very early stages of its existence, was trapped for a length of time in some local, metastable minimum where it suffered sufficient supercooling, before it reached its ground state (global minimum). Such phase transitions are endemic in spontaneously broken gauge theories at finite temperatures and they are best described by the study of the scalar potential V (free energy) as a function of the scalar field ϕ , causing the phase transition which in the case of inflation has been called¹⁶ the inflaton.

The first attempts¹⁰ which, justifiably enough, identified the inflaton field ϕ with the guton field Σ responsible for the spontaneous breaking of the GUT group G down to $SU(3) \times SU(2) \times U(1)$, failed badly. If the phase transition from $G \rightarrow SU(3) \times SU(2) \times U(1)$ is of the first order and occurs through nucleation, growth, and coalescence of bubbles of the $SU(3) \times SU(2) \times U(1)$ phase, the phase transition will never end and the Universe will never reheat back to high temperatures, since all the steps discussed above occur at 'normal' (adiabatic expansion) rates much smaller than the environmental exponential

expansion rate. This is the 'graceful exit' problem. If, though, we concentrate on one bubble, and extra care is taken so that because of the flatness of the scalar potential the inflaton takes a long time to reach its global minimum, it may happen that one bubble may be blown up sufficiently to encompass all our present Universe at that time. The 'graceful exit' problem has thus been automatically evaded. In this 'New Inflation' scenario¹⁷ the reheating of the Universe is due to the radiation coming from the oscillations of the inflaton field about its global minimum. Furthermore, quantum fluctuations of the inflaton field ϕ during inflation are responsible for the creation of energy density perturbations¹⁸ that are scale-independent when they enter their own Friedman horizon, exactly as needed¹⁹ for galaxy formation. The small-scale ($d < 100$ Mpc) inhomogeneity problem is, in principle, resolved in this scenario. It seems that i) sufficient reheating, so that conventional GUT baryogenesis occurs, and ii) GUT magnetic monopole dilution (there may exist only one magnetic monopole in the whole Universe!) ask for the identification of the inflaton field ϕ with the guton field Σ .

Alas, such identification, once more, has several catastrophic consequences for the New Inflation scenario¹⁷. The implied relations between scalar self-couplings and gauge couplings are of the wrong type^{20,21} for inflation. Indeed, uncontrollable quantum fluctuations of the inflaton field not only terminate inflation too soon^{20,21} to be of any cosmological use, but in addition they create scale-independent energy density perturbations of the wrong magnitude¹⁸ ($\delta\rho/\rho \sim 10$ - 100) instead of the needed ones ($\delta\rho/\rho \sim 10^{-4}$). Furthermore, the form of the implied scalar potential is such that eventually we will end up in some wrong vacuum and not in the correct $SU(3) \times SU(2) \times U(1)$ one²².

Clearly enough, controllable quantum fluctuations cry out²¹ for supersymmetry, whilst insurmountable fine-tunings are avoided if $\langle\phi\rangle \sim M$, which together with the 'wrong vacuum' problems requires the inflaton field to be a GUT-singlet and thus unidentifiable with the guton field Σ . Supersymmetric inflation with a GUT-singlet inflaton field ϕ with $\langle\phi\rangle \sim M$ has been called Primordial Inflation²³. Notice that since we are approaching the (super) Planck scale M , where gravitational interactions cannot be neglected, the correct manner^{16,24} of treating inflation is to use SUPERGRAVITY models². It is highly remarkable that not only in particle physics were we led almost uniquely to consider SUGRA models², but also the construction of a successful cosmological model seems to need^{16,23,24} SUGRA.

Before moving to SUGRA inflation let us recall the rules of the game. A successful inflationary potential should have the following properties^{16,23-25}:

- i) A local (metastable) minimum at the 'origin', stable under high-temperature corrections (*thermal stability*)²⁶ and of correct value [$V(0) \sim (10^{15} \text{ GeV})^4$] such that the scale-independent energy density perturbations have the correct magnitude:

$$\left. \frac{\delta \rho}{\rho} \right|_{\lambda=ct} \approx 0(10^{-4}) . \quad (5)$$

- ii) Sufficient flatness and *rigidity* (no 'holes')²⁷ so that there is enough inflation for cosmological purposes:

$$H \cdot \Delta\tau > 0(65) , \quad (6)$$

where $H(\approx \sqrt{V(0)}/M)$ is the Hubble parameter during inflation, whilst $\Delta\tau$ denotes the time-period that inflation lasts.

- iii) A very steep global minimum, with *naturally* vanishing vacuum energy (zero cosmological constant), so that sufficient reheating is feasible, meaning that

$$T_R \geq 0(10^{15} \text{ GeV}) , \quad (7)$$

if we insist on conventional GUT baryogenesis¹⁵.

It should be stressed here that in the case of Primordial Inflation ($\phi \neq \Sigma$), extra care should be taken about the dilution of magnetic monopoles and the GUT phase transition.

4. PRIMORDIAL (SUGRA) INFLATION

The form of the scalar potential V in an $N = 1$ SUGRA Yang-Mills theory is given by²⁸

$$V = e^G [G_i (G^{-1})^i_j G^j - 3] + \text{'D-terms'} , \quad (8)$$

where G is a real gauge-invariant function of the scalar components x_i ($i = 1, 2, \dots$) of the chiral multiplets χ_i , $G_i \equiv \partial G / \partial x^i$, $G^i \equiv \partial G / \partial x_i$, etc., and the contributions from the gauge interactions are generically presented as the 'D-terms'. The scalar kinetic terms are given by

$$G^i_j (\partial_\mu x^j) (\partial^\mu x_i) , \quad (9)$$

which have a form characteristic of SUSY non-linear σ -models. The scalar fields x_i in $N = 1$ SUGRA span a Kähler manifold with G^i_j as its metric or G as its Kähler potential²⁹. It is apparent from the potential (8) that, unlike the global case, spontaneously broken SUSY does not imply $\langle V \rangle > 0$. This is fortunate since we can now obtain spontaneous breakdown of local SUSY (the super-Higgs effect) in Minkowski space ($\langle V \rangle = 0$). In this case, the gravitino mass is given by

$$m_{3/2} = \left\langle \exp \frac{G}{2} \right\rangle . \quad (10)$$

Let us now see what form of G-function may lead to an acceptable cosmo-particle model.

The most naïve choice for the function G would be the one that provides canonical scalar terms,

$$G_{ij}^j = \delta_{ij}^j, \quad (11)$$

corresponding to a flat Kähler manifold. In such a case we write

$$G = zz^* + \phi\phi^* + y_\alpha y_\alpha^* + \ln |F(z, \phi, y_\alpha)|^2, \quad (12)$$

where we have separated the Polonyi or super-Higgs field z and the inflaton field ϕ from the 'matter' fields y_α which include all GUT chiral multiplets, and where $F(z, \phi, y_\alpha)$ is an analytic, gauge-invariant function of the indicated superfields, the so-called superpotential. However, the minimal choice (12) has some very unpleasant cosmo-particle consequences. On the particle physics front, we may succeed in getting $\langle V \rangle = 0$, but only thanks to some unbearable fine-tuning of the parameters in F . Furthermore, since in this case all scalar boson masses are proportional to the gravitino mass, and any low-energy fermion-boson mass splitting is bounded by the gauge hierarchy to be at most of $O(m_W)$, we deduce that³⁰

$$m_{3/2} \approx O(m_W). \quad (13)$$

Notice that in this case the relation (13) is introduced by hand, and any hope for dynamical determination disappears.

Coming now to the cosmological front we notice first that the unavoidably deduced relation (13) seems to lie in the middle of the forbidden gravitino-mass region (4). Of course, inflation may be of help here¹¹, but we have to make sure that regeneration of gravitinos by scattering processes will not be disastrous. Recent detailed calculations have shown¹²⁻¹⁴ that in the case where (13) is valid, the maximal reheating temperature allowed is¹²⁻¹⁴

$$T_R \leq O(10^{9 \pm 1} \text{ GeV}), \quad (14)$$

where the 10^8 GeV case is presumably more probable¹⁴. Clearly, the bound (14) means trouble because it makes conventional baryogenesis¹⁵ problematic³¹. Thermal production of colour Higgs triplets (H_3), the conventional baryogenerators, would imply masses of these H_3 obeying the upper bound (14) and thus leading to catastrophically short proton lifetime through dimension 5 ($m_{H_3} > 10^{15}$ GeV) or dimension 6 ($m_{H_3} > 10^{12-13}$ GeV) operators³¹. Actually, the presently known³² lower bound on the proton lifetime ($\tau_p > 10^{32}$ years) seems to be fatal also for non-thermal production of H_3 's through decays of the inflaton field, because it entails inflaton masses much above those allowed by inflation. A dead end. It is interesting to note that insistence on the 'minimal' choice of G as given by (12) creates also inflationary problems. True enough, the 'minimal'

G choice (12) was the starting point of SUGRA inflation^{24,16} and made it possible to establish the whole framework^{16,23,24} of discussing SUGRA inflation. However, subsequent detailed studies^{27,33} show that independently of the particular choice of $F(\phi)$ in (12), one can never satisfy simultaneously the *thermal stability* and the *rigidity* criteria discussed at the end of Section 3. So it seems that which ever way you look (cosmic or particle physics) at the 'minimal' G (12), it is no good.

The next interesting development³⁴ came from our desire for a naturally vanishing cosmological constant, with no fine-tunings, at least at the tree level. Indeed, it was noticed³⁴ that the G-function

$$G = -3 \ln(z+z^*) + \phi\phi^* + y_\alpha y_\alpha^* + \ln |F(\phi, y_\alpha)|^2 \quad (15)$$

has some really remarkable properties. The scalar potential V vanishes automatically at the tree level

$$V(z, z^*) = 0, \quad \forall z; \quad (16)$$

in other words we get³⁴ a flat-zero potential, i.e., zero cosmological constant. It is important to note that local SUSY is spontaneously broken, that is $m_{3/2} \neq 0$, but it is *undetermined* at the tree level. Of course, non-gravitational (electroweak, strong) radiative corrections will dynamically determine $m_{3/2}$, thus realizing the so-called no-scale dream. No-scale models³⁵⁻³⁸ are particle physics models where there is only one scale (acting as the unit) presumably to be identified with the (super) Planck scale M , and every other scale is dynamically determined in terms of this basic scale M . While using (15) one can construct very interesting particle physics models³⁵, (13) is still obeyed (derived and not introduced by hand), thus leading to all kinds of cosmological disasters discussed above³¹. In the case of inflation, it does not even allow for the construction of a local minimum³⁹: Some further modification is needed³⁹, at least for cosmological purposes.

Actually, most of the problem comes from the validity of (13), which can be traced back to the use of 'minimal', canonical kinetic terms for the 'matter' fields both in (12) and in (15). Furthermore, we know that global SUSY theories are characterized by semipositive definite potentials, so it may be interesting to employ them in order to ensure the *rigidity* constraint automatically [see discussion before (6)]. So we are³⁹ after some modification of (15) with non-minimal kinetic terms for all fields, and resulting in some kind of an effective, globally SUSY, scalar potential V . Such a form of G does exist³⁶ and can be parametrized as follows^{36,37}:

$$G = -3 \ln \left[f(z, z^*) - \frac{\phi\phi^*}{3} - \frac{y_\alpha y_\alpha^*}{3} \right] + \ln |F(\phi, y_\alpha)|^2, \quad (17)$$

which under certain general assumptions about $f(z, z^*)$ gives^{36,37}

$$V = |F_i|^2 + \text{'D-terms'} , \quad (F_i \equiv \partial F / \partial x^i) , \quad (18)$$

the classic semi-positive definite potential of (effective) global SUSY, thus automatically satisfying the *rigidity* constraint. Notice that in this case (13) is no longer valid, since there is no connection between gravitino and scalar masses, and the gravitino mass can be anything. For example, we can choose³⁷ $f(z, z^*)$ such that $m_{3/2}$ may be as big as M , or any other value we wish³⁸, presumably outside the forbidden range (4) and thus evading the gravitino problem. It should be stressed that in this gravitino mass-liberation movement^{37,38}, the burden of global SUSY-breaking shifts³⁶ to the gauginos, the superpartners of the gauge bosons. The gauginos may get tree-level global SUSY-breaking masses through non-trivial choices of an analytic function $f_{ab}(x_i) = f_{ba}(x_i)$, related to the Y-M part of the Lagrangian²⁸

$$- \frac{1}{4} (\text{Re } f_{ab}) F_{\mu\nu}^a F^{\mu\nu b} + \frac{i}{4} (\text{Im } f_{ab}) F_{\mu\nu}^a \tilde{F}^{\mu\nu b} \quad (19)$$

(a and b are indices of the adjoint representation of the gauge group G). For $f_{ab} = f(x_i) \delta_{ab}$ we get³⁸

$$m_{\tilde{V}} = m_{3/2} \frac{\langle 0 | f' | 0 \rangle}{\langle 0 | f | 0 \rangle} \frac{G'}{G} , \quad (20)$$

with primes indicating z differentiation. Clearly, in such types of no-scale models³⁶⁻³⁸ the previous role of the gravitino is played by the gaugino, which eventually will fix its mass dynamically through non-gravitational radiative corrections³⁶⁻³⁸. If we insist on dynamical determination of everything³⁸ -- $m_{3/2}$, $m_{\tilde{V}}$, etc. -- then we have to identify

$$f(z, z^*) = z + z^* \quad (21)$$

and proceed from there. Actually, then, something very interesting occurs: a natural solution³⁸ to the strong CP-hierarchy problem (3). One simply notices from (19) that non-minimal gauge-boson kinetic terms, needed for non-vanishing gaugino masses (20), give rise to

$$\theta_{\text{QCD}} = \frac{\langle \text{Im } f_{ab} \rangle}{\langle \text{Re } f_{ab} \rangle} . \quad (22)$$

Then, since there are *two* dynamical degrees of freedom ($\text{Re } z$, $\text{Im } z$), it is evident that these will dynamically determine two physical parameters. Since the gravitino is decoupled, we may naturally choose $m_{\tilde{V}}$ and θ_{QCD} as the two dynamically determined quantities. Although Standard Model [$\text{SU}(3) \times \text{SU}(2) \times \text{U}(1)$] radiative corrections fix the gaugino mass $m_{\tilde{V}}$, the only non-trivial dynamical dependence on θ_{QCD} comes from non-perturbative QCD effects which favour $\theta_{\text{QCD}} = 0$. The dynamical freedom accorded us by the $\text{Im } f_{ab}$ field enables us to 'relax' to $\theta_{\text{QCD}} = 0$, thus solving³⁸ the strong CP-hierarchy problem (3). In a class of

no-scale models (p-type models), the gravitino mass is determined to be³⁸

$$m_{3/2} = 0 \left(\frac{m_V^p}{M^{p-1}} \right) \ll m_{\tilde{V}} \approx 0(m_W) , \quad (23)$$

with $1 < p < 2$, in order to avoid too singular couplings to ordinary matter, and more accurately $1.2 < p < 1.4$ in order to be consistent with a large number of cosmological, astrophysical, and particle physics constraints. Amazingly enough, the above-mentioned 'singular' couplings to matter, or self-couplings, may help the Polonyi field to 'relax' more quickly to its minimum and thus possibly evade³⁸ also the Polonyi problem.

It is rather apparent that we are on the right track, so let us next try to construct inflationary models based on a certain generalization of (17) that is suitable for our purposes³⁹⁻⁴²

$$G = -3 \ln \left[f(z, z^*) - K(\phi, \phi^*) - \frac{y_\alpha y_\alpha^*}{3} \right] - \alpha K(\phi, \phi^*) + \ln |F(\phi, y_\alpha)|^2 . \quad (24)$$

Once more, the potential V emerging from (24) is semi-positive definite, for $\alpha \geq 0$ and so the *rigidity* constraint is automatically satisfied. In the case $K(\phi, \phi^*) = \phi\phi^*/3$ and $\alpha = 0$ we obviously fall back on (17). There are two general classes of inflationary models using (24) as their starting point.

Class I⁴²

Choose $f(z, z^*)$ in such a way that second and higher derivatives, at least up to fourth order, vanish at the physical minimum z_0 and $|f_z| = \text{constant}$, e.g.

$$f(z, z^*) = (z + z^*) + A(z + z^*)^m , \quad m > 4 \quad (z_0 = 0) , \quad (25)$$

and $K(\phi, \phi^*) = \phi\phi^*/3$ and $\alpha = 0$. In such a case it is a straightforward matter to check⁴² that a simple superpotential F , like

$$F(\phi, y_\alpha = 0) = m_0 M^2 \left(\phi - \frac{\phi^4}{4} \right) , \quad (26)$$

with $m_0 \approx (10^{-5} - 10^{-6})M$, satisfies the inflation constraints [(5) and (6)]. Furthermore, a detailed calculation⁴² of high-temperature effects^{26,43} shows that indeed $\phi = 0$ is not only a local minimum of the $T = 0$ potential $V(T = 0)$, but that it remains a local minimum also of the $V(T \neq 0)$ potential. *Thermal stability* is ensured. Notice that in this class of models the gravitino mass is determined by hand [through $f(z, z^*)$] at the tree level, once and for all at a rather superhigh value [$m_{3/2} \approx 0(m_0)$].

Class II³⁹⁻⁴¹

In this class of models we choose

$$f(z, z^*) = z + z^* ; \quad K(\phi, \phi^*) = \frac{\phi\phi^*}{3} - \frac{\beta}{4} (\phi\phi^*)^2 \dots ; \quad \alpha > 0 . \quad (27)$$

In such a case we write

$$\ln F(\phi, y_\alpha=0) = \ln m_0 + F(\phi) , \quad (28)$$

with

$$F(\phi) = \xi\phi + \frac{1}{2}(2-\xi^2)\phi^2 + O(\phi^3) , \quad (29)$$

and easily satisfy the inflation constraints (5) and (6) by again choosing, for example, $m_0 \approx (10^{-5})M$ and reasonable values of ξ and β belonging to the range $(10^{-2}-1)$. Then straightforward but tedious calculations⁴⁴ show that the *thermal stability* constraint is trivially satisfied if $\xi > 0$ -- a rather innocuous assumption. Notice that in this case ($\xi > 0$), the local minimum is moved to the 'left' of $\phi = 0$ at high temperatures, so that at lower temperatures it moves forwards and stabilizes at $\phi = 0$. In this class³⁹⁻⁴¹ of inflationary models the gravitino mass, which is of order m_0 during inflation, could be again undetermined after inflation because of the form of $f(z, z^*)$ (27). Thus, such inflationary models can be tagged on to the p-type models³⁸ (23), automatically resolving several problems.

Additionally, we have yet to discuss the connection with the GUT breaking, the magnetic monopole problem, and the reheating constraint (7). Recently, a novel, simple mechanism has been proposed⁴⁵ for a simultaneous resolution of the aforementioned problems. It is actually possible to decouple the scale of the slow roll-over, fixed by $\delta\rho/\rho$ (5) to be $H \approx m_0 \approx (10^{-5}-10^{-6})M$, from the scale present during the end of inflation. This happens automatically if the inflaton is coupled to the GUT non-singlets in a non-trivial way⁴⁵. Indeed, rather than impose a complete decoupling of the inflaton and GUT sector, it is more natural to consider the most general linked superpotential of the form⁴⁵

$$F(\phi, \Sigma, y_{\alpha \neq \Sigma} = 0) = A(\phi) + B(\phi)\sigma^2 + C(\phi)\sigma^3 + \dots , \quad (30)$$

where σ is the $SU(3) \times SU(2) \times U(1)$ singlet in the Σ , and $\sigma^n = \text{Tr } \Sigma^n$. *A priori*, A , B , C , etc., are arbitrary analytic functions of ϕ which represent the (in general non-renormalizable) couplings of ϕ to σ . Care must be taken to ensure that these couplings do not perturb our hard-won satisfaction with the thermal stability constraint^{42,44}. Because of these couplings the curvature of the potential in the σ -direction changes as ϕ rolls, and can become negative for some range of ϕ values, triggering a second-order phase transition which breaks the GUT group (see figure). After the phase transition, the dominant scale of the potential is $M_{\text{GUT}} \gg m_0 (\approx H)$. In effect, the inflaton becomes a mixture of ϕ and σ , and can give adequate reheating (7) through the decays of the σ -oscillations^{41,45}. Furthermore, magnetic monopoles, created during the GUT phase transition, may, in principle, be easily diluted by the residual inflationary period after the GUT phase transition, as quantitative estimates

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 the line $\phi = \phi_0$ and in the (σ, ϕ) plane for simplicity. The dotted line represents the trajectory of the inflaton field.

show^{41,45}. The amount of magnetic monopole dilution is model-dependent, but the possibility remains that an experimentally detectable magnetic monopole number density does survive^{41,45}. As promised, this novel mechanism⁴⁵ (30) does provide a simultaneous solution to the GUT phase transition, magnetic monopole, and reheating problems, and can be used in any primordial inflation model. Two-component primordial inflation models^{41,45} seem to share all the 'goodies' of primordial inflation^{16,23,24} around the local minimum (ϕ is really ϕ) and the 'goodies' of 'New Inflation'¹⁷ around the global minimum (ϕ becomes really σ). We seem to get the best of two pictures⁴¹. It is thus fair to call such an inflationary approach: *NEW PRIMORDIAL INFLATION*.

All in all, SUGRA seems to be a fundamental ingredient in the construction of realistic particle physics as well as cosmological models. Gravitational

interactions, as contained in SUGRA, seem to be needed not only (as usual) at very long distances and times, but they are badly needed also at very small distances and times.

ACKNOWLEDGEMENTS

I would like to express my appreciation to R. Baier and S. Satz for organizing such an exciting conference on a very interesting theme. Many thanks are also due to Kari Enqvist for useful and enjoyable discussions on many topics of this talk and for his critical reading of the manuscript.

REFERENCES

- 1) For a detailed description of the standard hot Big-Bang Cosmology see, for example, S. Weinberg, *Gravitation and Cosmology* (Wiley, New York, 1972), Chapter 15.
- 2) For recent reviews of physical supergravity see, for example, J. Ellis, preprint CERN-TH.3802 (1984);
H.P. Nilles, *Phys. Rep.* 110C (1984) 2;
D.V. Nanopoulos, preprints CERN-TH.3699 (1983) and CERN-TH.3995 (1984).
- 3) E. Witten, *Nucl. Phys.* B188 (1981) 513.
- 4) A. Masiero, D.V. Nanopoulos, K. Tamvakis and T. Yanagida, *Phys. Lett.* 115B (1982) 380.
B. Grinstein, *Nucl. Phys.* B206 (1982) 387.
C. Kounnas, D.V. Nanopoulos, M. Srednicki and M. Quirós, *Phys. Lett.* 127B (1983) 82.
- 5) J. Ellis, S. Ferrara and D.V. Nanopoulos, *Phys. Lett.* 114B (1982) 231.
- 6) For a review see, for example, J. Preskill, in: *The very early Universe*, eds. S.W. Hawking, G.W. Gibbons and S. Siklos (Cambridge University Press, Cambridge, UK, 1983), p. 119.
- 7) D.V. Nanopoulos and K. Tamvakis, *Phys. Lett.* 110B (1982) 449.
M. Srednicki, *Nucl. Phys.* B202 (1982) 327, and B206 (1982) 132.
D.V. Nanopoulos, K.A. Olive and K. Tamvakis, *Phys. Lett.* 115B (1982) 15.
- 8) H. Pagels and J.R. Primack, *Phys. Rev. Lett.* 48 (1982) 223.
S. Weinberg, *Phys. Rev. Lett.* 48 (1982) 1303.
- 9) G.D. Coughlan, W. Firschler, E.W. Kolb, S. Raby and G.G. Ross, *Phys. Lett.* 131B (1983) 59.
D.V. Nanopoulos and M. Srednicki, *Phys. Lett.* 133B (1983) 287.
- 10) A.H. Guth, *Phys. Rev.* D23 (1981) 347.
D. Kazanas, *Ap. J.* 241 L59 (1980).
K. Sato, *Mon. Not. R. Astron. Soc.* 195 (1981) 467 and *Phys. Lett.* 99B (1981) 66.
- 11) J. Ellis, A.D. Linde and D.V. Nanopoulos, *Phys. Lett.* 118B (1982) 59.

- 12) M.Yu. Khlopov and A.D. Linde, Phys. Lett. 138B (1984) 265.
- 13) J. Ellis, J.E. Kim and D.V. Nanopoulos, Phys. Lett. 145B (1984) 181.
- 14) J. Ellis, D.V. Nanopoulos and S. Sarkar, preprint CERN-TH.4057 (1984).
- 15) For recent reviews of the interplay between unified theories and cosmology see, for example, K.A. Olive, Fermilab preprint FERMILAB Conf.-84/59 A (1984); D.V. Nanopoulos, in: Proc. 1st ESO-CERN Symposium on Large-Scale Structure of the Universe, Cosmology and Particle Physics, eds. G. Setti and L. Van Hove (ESO-CERN Publication, Geneva, 1984), p. 349.
- 16) D.V. Nanopoulos, K.A. Olive and M. Srednicki, Phys. Lett. 127B (1983) 30.
- 17) A.D. Linde, Phys. Lett. 108B (1982) 389.
A. Albrecht and P.J. Steinhardt, Phys. Rev. Lett. 48 (1982) 1220.
- 18) S.W. Hawking, Phys. Lett. 115B (1982) 295.
A.H. Guth and S.Y. Pi, Phys. Rev. Lett. 49 (1982) 1110.
J.M. Bardeen, P.J. Steinhardt and M.S. Turner, Phys. Rev. D28 (1983) 679.
- 19) E.R. Harrison, Phys. Rev. D1 (1970) 2726.
Ya.B. Zel'dovich, Mon. Not. R. Astron. Soc. 160 (1972) 18.
- 20) A.D. Linde, Phys. Lett. 116B (1982) 335.
- 21) J. Ellis, D.V. Nanopoulos, K.A. Olive and K. Tamvakis, Phys. Lett. 118B (1982) 335.
- 22) J. Breit, S. Gupta and A. Zaks, Phys. Rev. Lett. 51 (1983) 1007.
- 23) J. Ellis, D.V. Nanopoulos, K.A. Olive and K. Tamvakis, Nucl. Phys. B221 (1983) 224.
- 24) D.V. Nanopoulos, K.A. Olive, M. Srednicki and K. Tamvakis, Phys. Lett. 123B (1983) 41.
- 25) P.J. Steinhardt and M. Turner, Phys. Rev. D29 (1984) 2162.
- 26) G.B. Gelmini, D.V. Nanopoulos and K.A. Olive, Phys. Lett. 131B (1983) 53.
- 27) B.A. Ovrut and P.J. Steinhardt, Phys. Lett. 133B (1983) 161.
- 28) E. Cremmer, B. Julia, J. Scherk, S. Ferrara, L. Girardello and P. van Nieuwenhuizen, Phys. Lett. 79B (1978) 231 and Nucl. Phys. B147 (1979) 105.
E. Cremmer, S. Ferrara, L. Girardello and A. Van Proeyen, Phys. Lett. 116B (1982) 231 and Nucl. Phys. B212 (1983) 413.
- 29) B. Zumino, Phys. Lett. 87B (1979) 203.
- 30) J. Ellis and D.V. Nanopoulos, Phys. Lett. 116B (1982) 283.
- 31) J.E. Kim, A. Masiero and D.V. Nanopoulos, Phys. Lett. 145B (1984) 187.
J. Ellis, K. Enqvist, G. Gelmini, C. Kounnas, A. Masiero, D.V. Nanopoulos and A.Yu. Smirnov, Phys. Lett. 147B (1984) 27.

- 32) For a recent review see M. Koshiha, Experiments on baryon number non-conservation, Plenary talk given at the 22nd Int. Conf. on High-Energy Physics, Leipzig, 1984 (to appear in the Proceedings).
- 33) A.D. Linde, Phys. Lett. 132B (1983) 317.
B.A. Ovrut and P.J. Steinhardt, Phys. Rev. Lett. 53 (1984) 732 and Univ. Pennsylvania preprints UPR-0264 and UPR-0266T (1984).
R. Holman, P. Ramond and G.G. Ross, Phys. Lett. 137B (1984) 343.
- 34) E. Cremmer, S. Ferrara, C. Kounnas and D.V. Nanopoulos, Phys. Lett. 133B (1983) 61.
- 35) J. Ellis, A.B. Lahanas, D.V. Nanopoulos and K.A. Tamvakis, Phys. Lett. 134B (1984) 429.
J. Ellis, C. Kounnas and D.V. Nanopoulos, Nucl. Phys. B241 (1984) 406.
- 36) J. Ellis, C. Kounnas and D.V. Nanopoulos, preprint CERN-TH.3824 (1984).
- 37) J. Ellis, C. Kounnas and D.V. Nanopoulos, Phys. Lett. 143B (1984) 410.
- 38) J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys. Lett. 147B (1984) 99 and preprint CERN-TH.4036 (1984).
- 39) G.B. Gelmini, C. Kounnas and D.V. Nanopoulos, preprint CERN-TH.3777 (1983).
- 40) C. Kounnas and M. Quirós, preprints LPTENS 84/21 (1984).
- 41) K. Enqvist, C. Kounnas, D.V. Nanopoulos and M. Quirós, preprint CERN-TH.4027 (1984).
- 42) J. Ellis, K. Enqvist, D.V. Nanopoulos, K.A. Olive and M. Srednicki, preprint CERN-TH.4039 (1984).
- 43) K.A. Olive and M. Srednicki, Finite temperature corrections in $N = 1$ supergravity, Univ. California, Santa Barbara, preprint (1984).
P. Binétruy and M.K. Gaillard, Univ. California, Berkeley, LBL preprint UCB-PTH-84/19/LBL-18072 (1984).
- 44) K. Enqvist, D.V. Nanopoulos and M. Quirós, preprint CERN-TH.4040 (1984).
K. Enqvist, C. Kounnas, D.V. Nanopoulos and M. Quiros, preprint CERN-TH.4048 (1984).
- 45) K. Enqvist and D.V. Nanopoulos, Phys. Lett. 142B (1984) 349 and preprint CERN-TH.3935 (1984).