CERN-TH.4145/85

INFLATION FROM A RIPPLE ON A VANISHING POTENTIAL

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ABSTRACT

We propose a very simple model of inflation having essentially one free parameter, the value of which is fixed by the amplitude and scale independence of energy density fluctuations. The model, based on the maximally symmetric supergravity with SU(n,1) manifold, has asymptotically flat inflaton potential. All inflationary conditions can be satisfied without any fine-tuning and all mass parameters can be $O(M_{P\ell})$.

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CERN-TH.4145/85
April 1985

Cosmological inflation[1] is conceptually a very attractive idea; yet no particle physics model exists in which one can implement both inflation and reheating of the universe in a simple way. The shortcomings of the "new inflation"[2], based on the Coleman-Weinberg type flat potentials, led people to consider supersymmetric model of inflation[3]. Simplest versions, based on the so-called minimal supergravity or trivial Kähler metric, were also shown to suffer from some diseases[4]. In addition, minimal supergravity has a gravitino of mass $m_{3/2} \simeq O(M_W)$, which in the absence of inflation is a cosmological embarrassment[5]. Inflation can dilute[6] any primordial gravitino number density, but unfortunately gravitinos will be regenerated after inflation by $2 \rightarrow 2$ scatterings. Recently, it has been shown[7] that if the regenerated gravitino density is to be suppressed below levels where the cosmic abundances of light elements will not be disrupted by gravitino decays, one finds a very low reheating temperature $T_{RH} \lesssim 10^8$ GeV. Such a low reheating temperature is problematic for particle physics, especially for obtaining a correct baryon asymmetry while suppressing unwanted nucleon decays due to dimension 5 baryon number violating operators[8]. This constraint is an obvious difficulty for models having a $O(10^2)$ GeV gravitino, and although it is possible by clever choices to rotate away[9] the $d = 5$ $\Delta B \neq 0$ operators, the resulting models lack all aesthetic appeal.

The decoupling of the gravitino mass from the scale of SUSY breaking, as felt by the light fields, can be achieved only in "non-minimal" versions of $N = 1$ supergravity with non-trivial Kähler metric. Here the efforts to write down an acceptable model of inflation and particle physics have focused on the no-scale models[10],[11]. These have a maximally symmetric Kähler manifold with an SU(n,1) global symmetry, where n is the number of chiral fields. Such models have flat tree level potentials and can lead to interesting low-energy particle physics as well as to a dynamical determination of all mass scales below the Planck mass. It can also give raise to a novel mechanism[11], whereby the QCD vacuum angle relaxes dynamically to $\Theta = 0$. In the no-scale models the source of the low-energy SUSY breaking is the gaugino masses, whereas the gravitino can be very heavy[10] or very light[11]. The various cosmological and particle physics constraints for a very low mass gravitino in no-scale models have recently been considered in Ref. 11).

The SU(n,1) structure of the Kähler manifold has also recently been advocated[12] in the context of superstrings[13]. The compactification of the effective ten-dimensional theory to $N = 1$ supergravity in four dimensions dictates[14] the compact internal manifold to have SU(3) holonomy. Calabi-Yau manifolds have such a property, but otherwise little is known of them. Therefore it has been suggested[12)] that a simple truncation of ten-dimensional supergravity preserving an SU(3) subgroup of the rotation group of the six-dimensional internal space may serve to mimic the properties of compactification on Calabi-Yau manifold. It is remarkable that the resulting effective theory can be shown to possess SU(n,1) Kähler manifold such as has been studied in Refs. 10) and 11).

To obtain an acceptable inflationary scenario, the SU(n,1) symmetry must however be broken; otherwise there would be no vacuum energy to drive inflation. A suitable generalization of the SU(n,1) manifold, which can be shown to fulfil all requirements for inflation[15)-18)], is given by the Kähler potential

$$\begin{aligned} G &= -3\ln\left(z+z^* - K(\phi,\phi^*) - \tfrac{1}{3} y^i y_i^*\right) - aK(\phi,\phi^*) + F + F^+ \\ &\equiv G_0 - aK + F + F^+ \end{aligned} \tag{1}$$

Here z is a gauge singlet field responsible for SUSY breaking, y^i are the matter fields, ϕ is a gauge singlet inflaton field and $F = F(\phi, y^i)$ is the superpotential. The function K can be assumed to be a function of $\phi\phi^+$, and a is the parameter responsible for SU(n,1) breaking with $a \to 0$ corresponding to the limit where one recovers the SU(n,1) structure of the Kähler manifold. The choice $a \neq 0$ is thus dictated by the requirement of non-zero vacuum energy. The scalar potential corresponding to (1) can be written as

$$V = \frac{x^3}{3x-a} e^{F+F^+-aK} \frac{|F_\phi - aK_\phi|^2}{K_{\phi\phi^+}} + x^2 e^{F+F^+-aK} F_i^+ F^i + \tfrac{1}{2} \tilde{D}^\alpha \tilde{D}^\alpha \tag{2}$$

where $\tilde{D}^\alpha$ refers to a correctly normalized D-term and $x = \exp G_0/3$ [for a discussion on normalizations of fields and Lagrangian terms in non-minimal supergravities, see Ref.19)]. In the domain of positive kinetic energies where $3x-a > 0$ and $K_{\phi\phi^+} > 0$ the potential (2) is positive semi-definite. The properly normalized kinetic energy terms are obtained after a point transformation

$$\begin{aligned} \phi &\to \left[(3x-a)K_{\phi\phi^+}\right]^{1/2}\Big|_{\substack{\phi=\phi_0 \\ x=x_0}} \phi \\ y^a &\to x^{1/2}\Big|_{\substack{\phi=\phi_0 \\ x=x_0}} y^a \end{aligned} \tag{3}$$

valid at the minimum (ϕ_0, x_0). The minimization of the potential (2) leads to $F_i = \tilde{D}^\alpha = 0$, and at the metastable minimum with non-zero vacuum energy one also has the condition

$$x = a/2 \tag{4}$$

At very high temperatures, all non-singlet fields y^i have been shown[19),20)] to choose their symmetric minima so that for the purposes of inflation, we will consider the case $y^i = 0$, which, together with (4), leads to

$$V = \left(\frac{a}{2}\right)^2 e^{F+F^+ - aK} \left|F_\phi - aK_\phi\right|^2 K^{-1}_{\phi\phi^+} \tag{5}$$

where now $F = F(\phi)$. However, at the global minimum $F_\phi = aK_\phi$, x, and the gravitino mass, is no longer restricted by (4) but only by the positivity condition $x > a/3$.

The potential (5) is attractive for inflationary purposes because of its simple structure and because of its positive definiteness. A model based on this generalization of maximally symmetric Kähler manifold is defined once $F(\phi)$ and $K(|\phi|)$ are given. In previous attempts[15)-19),21)] rather complicated forms were used, leading to a potential suitable for inflation but bearing no relation to the SU(n,1) manifold. Here we will consider the extremely simple possibility of potentials that are globally flat in the sense that $V \to 0$ as $|\phi| \to \infty$, but show local deviations from flatness. Surprisingly enough, although such potentials are very constrained, having only one free parameter, we are able to obtain an inflationable potential in a natural way for a large range of this one parameter.

The aim of this paper is to show that, even in the simplest of the models described by (5), all inflation conditions are satisfied and scale-independent energy density fluctuation with the correct amplitude can be made consistent with all mass parameters being $O(M_{P\ell})$.

We will take the simplest imaginable ansatz for the function K:

$$K = \phi\phi^+ \tag{6}$$

so that the inflaton field has canonical kinetic terms up to an irrelevant numerical factor $(a/2)^{1/2}$. The superpotential can be written as

$$F(\phi) = F(0) + \xi\phi + \frac{1}{2}(a-\xi^2)\phi^2 + \frac{1}{6}\rho\phi^3 + \cdots \qquad (7)$$

Here $F(0) = \log m_0$ and $\xi \neq 0$ are related to the Hubble constant during the inflation and the particular form of the coefficient of the quadratic term is dictated by the condition of having a stationary point at the origin. Here we do not consider higher order than cubic terms. The curvature along the real and imaginary directions at the origin is given by

$$\begin{aligned} m_R^2 &= m_0^2\,\xi\,(\rho - 2\xi^3) \\ m_I^2 &= m_0^2\,(4a^2 - 6a\xi^2) - m_R^2 \end{aligned} \qquad (8)$$

If $\rho \neq 0$ the global minimum is not unique: there are two complex-conjugate global minima with

$$\begin{aligned} u &= \frac{\sqrt{2}}{\xi}\,\rho^{-1}(\xi^2 - 2a) \\ v^2 &= \frac{4\xi}{\rho}\left(1 + \frac{2a-\xi^2}{\rho}\right) + \left(\frac{2a-\xi^2}{\rho}\right)^2 \frac{2}{\xi^2} \end{aligned} \qquad (9a)$$

where $\phi = 1/\sqrt{2}(u+iv)$. If $\rho < 1/2\xi^3$, there are two real (v=0) minima given by the solutions of

$$1 - \frac{\xi}{\sqrt{2}}u + \frac{\rho}{4\xi}u^2 = 0 \qquad (9b)$$

The most economical assumption is, however, that $\rho = 0$. In that case one finds a unique real minimum at

$$\phi_0 = \sqrt{2}/\xi \qquad (10)$$

The potential will now be flat at large Reϕ, independently of ξ, and globally flat in the whole complex plane if $\xi^2 < 2a$. Moreover, if $\xi^2 < a$, the origin becomes a minimum along the imaginary direction and the path followed by the inflaton will be along the real direction.

Let us now examine the conditions on inflation assuming for the moment that the initial distribution of the inflaton field is such that at the beginning of inflation $\phi \simeq 0$. We will return to this question at the end of the paper.

A slow roll-over of the inflaton field can be achieved if[22)]

$$|V''(\phi)| \lesssim 9H^2(\phi) \tag{11a}$$

$$|V'(\phi)/V(\phi)| \lesssim \sqrt{6} \tag{11b}$$

where the Hubble parameter is

$$H^2(\phi) = \frac{1}{3}\left[V(\phi) + \frac{1}{2}\dot{\phi}^2\right] \tag{12}$$

Equation (11a) is satisfied for $0 < \phi < \phi_0$ provided $\xi < 3/2$. Inflation ends when $|V''(\phi_f)| \simeq 9H^2(\phi_f)$ which translates in our case to

$$3\left[1 - \frac{\xi}{\sqrt{2}}\phi_f\right]^2 = \xi^2 \left|2 - \sqrt{2}\xi\phi_f - \frac{7}{2}\xi^2\phi_f^2 + 2\sqrt{2}\xi^3\phi_f^3 - \frac{1}{2}\xi^4\phi_f^4\right| \tag{13}$$

During the slow roll-over, the semi-classical equation of motion of the inflaton

$$\ddot{\phi} + (3H + \Gamma)\dot{\phi} + V'(\phi) = 0 \tag{14}$$

and the Hubble parameter (12) can be approximated by

$$3H\dot{\phi} + V'(\phi) = 0 \tag{15a}$$

$$H^2 = \frac{1}{3}V(\phi) \tag{15b}$$

In that case, the number of e-foldings of the scale factor R(t) as ϕ rolls from ϕ_a to ϕ_b can be calculated exactly and is

$$N(a,b) = \int_a^b H(t)\,dt = \frac{1}{2\xi^2} \log \frac{\phi_b\left(1 - \frac{\xi}{2\sqrt{2}}\phi_b\right)}{\phi_a\left(1 - \frac{\xi}{2\sqrt{2}}\phi_a\right)} \tag{16}$$

A sufficient amount of inflation is obviously obtained for values of ξ that are less than O(1). The most stringent constraint on the parameters of inflationary potentials comes, however, from the magnitude of the energy density fluctuations, $\delta\rho/\rho \simeq 10^{-4}$ on galactic scales. The amplitude of such perturbations, when they re-enter the Friedman horizon after the inflationary epoch, is given by

$$\frac{\delta\rho}{\rho} = \frac{1}{\sqrt{2}\,\pi^{3/2}} H^2(\phi_*)\,\dot{\phi}_*^{-1} \tag{17}$$

where $\phi_* = \phi_*(t_*)$ and t_* is the time of the first horizon crossing during inflation. We estimate the number of e-foldings between ϕ_* and ϕ_f, the end of inflation, to be

$$N_* \simeq 43 + \frac{1}{3}\log M/M_\odot \tag{18}$$

where $M_\odot$ is the solar mass, M is the mass contained in the comoving volume, and a reheating temperature $T_{RH} \simeq 5\times10^{15}$ GeV, consistent with our further results, has been used. For the scales $M_1 = 10^{12}M_\odot$ (galaxy mass), $M_2 = 10^{15}M_\odot$ (supercluster mass) and $M_3 = 10^{22}M_\odot$ (present horizon mass), N_* takes the values 52, 55 and 60, respectively. Therefore the deviation from exact scale invariance is very small but will in general depend on ξ. We can obtain $\dot{\phi}_*$ from the equations of motion (15a), whereas ϕ_* can be expressed as a function of ϕ_f and N_* by using (16). Then energy density perturbations can be written as

$$\frac{\delta\rho}{\rho} = \frac{\left[1 - \sqrt{2}\,\xi\phi_f\left(1 - \frac{\xi}{2\sqrt{2}}\phi_f\right)e^{-2\xi^2 N_*}\right]^{1/2} H_*}{(2\pi)^{3/2}\xi^2\,\phi_f\left(1 - \frac{\xi}{2\sqrt{2}}\phi_f\right)}\, e^{2\xi^2 N_*} \tag{19}$$

where $H_* = H(t_*)$ is the Hubble parameter at the time of the first horizon crossing.

From now on, we will proceed in the following way: for a given value of ξ, we obtain the value of the inflaton field at the end of inflation ϕ_f, solving Eq. (13), and at the time of horizon crossing, ϕ_*, using (16) with $N_* = N_3 = 60$. The Hubble parameter at t_* can be deduced from (19) and using the condition (17). From the very definition of $H(\phi)$ and the values of H_*, ϕ_* and ϕ_f, we can deduce $H_0 \equiv H(\phi)$ and $H_f \equiv H(\phi_f)$ and thus the energy $V_0^{1/4}$ and $V_f^{1/4}$. Finally the total amount of e-foldings which the cosmic scale undergoes during the whole inflationary period can be read off from (16) as

$$N_t = \frac{1}{2\xi^2} \log \frac{\phi_f\left(1 - \frac{3}{2\sqrt{2}}\phi_f\right)}{H_0} \tag{20}$$

The deviation of $\delta\rho/\rho$ from scale invariance is given by

$$\Delta \equiv \frac{(\delta\rho/\rho)_2}{(\delta\rho/\rho)_1} = \left[\frac{1 - \sqrt{2}\xi\phi_f\left(1 - \frac{3}{2\sqrt{2}}\phi_f\right)e^{-2\xi^2 N_2}}{1 - \sqrt{2}\xi\phi_f\left(1 - \frac{3}{2\sqrt{2}}\phi_f\right)e^{-2\xi^2 N_1}}\right]^{1/2} e^{-2\xi^2(N_2 - N_1)} \tag{21}$$

In Fig.1 we show ϕ_m, ϕ_f, ϕ_*, $V_f^{1/4}$, $V_0^{1/4}$, N_t and Δ for $M_1 = 10^{12}M_\odot$ and $M_2 = 10^{22}M_\odot$ as functions of ξ. We see that $N_t > 60$, which is the minimum amount of inflation needed to solve the cosmological problems of horizon and flatness, for the whole considered interval of ξ. However, the deviation from scale invariance is $\Delta < O(3)$, which is probably allowed, for $\xi \lesssim 0.24$ which is, as can be seen from Fig.1, the interesting range. On the other hand ϕ_* becomes $\ll M_{P\ell}$ for $\xi \gtrsim 0.2$. This can be understood analytically as follows: for $\xi^2 \gtrsim 1/2N_*$, i.e., $\xi \gtrsim 0,1$, ϕ_* can be approximated by

$$\phi_* = \phi_f\left(1 - \frac{3}{2\sqrt{2}}\phi_f\right)e^{-2\xi^2 N_*} \tag{22}$$

which approaches zero asymptotically as ξ^2 increases.

The reheating temperature T_{RH} is related to the available energy at the end of the inflation $V_f^{1/4}$, in the case of good reheating, by

$$T_{RH} = \left(\frac{30}{\pi^2 g_*}\right)^{1/4} V_f^{1/4} \tag{23}$$

where g_* is the effective number of relativistic degrees of freedom. Typically $g_* = O(10^2)$ and $T_{RH} \simeq (0.4\text{–}0.5)V_f^{1/4}$. In this way one can obtain 10^{15} GeV $< T_{RH} \lesssim 10^{16}$ GeV for $\xi \lesssim 0.2$. This range overlaps with the range of scale-invariant density fluctuations, and it is the interesting one.

For very small values of ξ, H_* and H_f behave like

$$\begin{aligned} H_* &\simeq 2.4 \times 10^{14}\ \text{GeV} \\ H_f &\simeq 1.3 \times 10^{13}\ \text{GeV} \\ V_f^{1/4} &\simeq 7.3 \times 10^{15}\ \text{GeV} \end{aligned} \tag{24}$$

while H_0 and $V_0^{1/4}$ increase as

$$\begin{aligned} H_0 &\simeq 1.35 \times 10^{13}\ \text{GeV}\ \xi^{-1} \\ V_0^{1/4} &\simeq 7.49 \times 10^{15}\ \text{GeV}\ \xi^{-1/2} \end{aligned} \tag{25}$$

For the sake of completeness we have tabulated, in the Table, the numerical values of $\phi_i = H_0$, ϕ_*, ϕ_f, ϕ_m, $V_f^{1/4} = (\pi^2 g_*/30)^{1/4} T_{RH}$, Δ and N_t corresponding to a typical set of values of the parameter ξ.

It is astonishing that we really can obtain a successful description of inflation by using a one-parameter model (with m_0 being only an over-all scale factor which, as we will see, can be set equal to $M_{P\ell}$) and that the condition on ξ, as dictated by Δ, is only the mild $\xi \lesssim 0.2$.

Therefore, having now established the feasibility of inflation and scale-independent energy density fluctuations with the correct amplitude for small ξ, we now come to the value of the over-all mass scale m_0. In most inflationary models m_0 has to be adjusted to a very small value $m_0 \simeq O(10^{-5}\text{–}10^{-6})$ dictated by $\delta\rho/\rho \simeq 10^{-4}$. However in our model, using the expression (25) for small values of ξ, we get

$$m_0 \simeq 10^{-5}\,\xi^{-2} \tag{26}$$

so that m_0 gets small values only for values of ξ ruled out by scale independent energy density fluctuations; but $m_0 \simeq O(1)$ for $\xi \simeq 3\times10^{-3}$.

In conclusion, we have considered the simplest model of inflation based on maximally symmetric supergravity. The model has two free parameters: ξ and m_0. All requirements needed for a successful inflationary scenario are satisfied without any fine tuning of parameters. Energy density fluctuations are scale-independent for $\xi < 0.24$. The correct amplitude $\delta\rho/\rho \simeq 10^{-4}$ determines the Hubble constant at the end of the inflationary period H_f [and hence the reheating temperature for the case of good reheating as in the two-component inflationary model proposed in Ref. 16)] and at the origin H_0. The further requirement of large reheating temperature ($T_{RH} \gtrsim 10^{15}$ GeV) - needed to generate, after the inflation, the cosmological baryon asymmetry by out-of-equilibrium decays of heavy ($M_H \gtrsim 10^{16}$ GeV) colour triplets - leads to the mild constraint $\xi < 0.2$.

In this paper we have not addressed the question of the initial pre-inflationary conditions. Two different mechanisms have been so far proposed: the high-temperature phase transition[23)] and the chaotic inflation[24)] scenario. The feasibility of a high-temperature phase transition depends on the value of the typical relaxation time ($\Delta t \sim 1/m_0 T$) as compared to the age of the (radiation-dominated) Universe ($t_U \sim 1/T^2$). In fact, as has been argued by Linde[24)], if $\Delta t > t_U$ for $T \gtrsim T_c$, then any initial distribution of the inflaton field may remain effectively frozen and the theory of temperature phase transitions does not apply. In that case one should consider some chaotic initial distribution for ϕ and assume that in some domain of the Universe the field ϕ was initially near the origin $0 < \phi \lesssim H_0$.

On the other hand, if $\Delta t < t_U$ then any initial distribution will have the time to relax dynamically to the origin at T_c and temperature corrections are relevant. In our model, small values of ξ, i.e., large values of m_0, seem to favour this possibility. Not only that, recently Mazenko et al.[25)] have argued that large field fluctuations would form domains where the inflaton lies at the global minimum $\phi = \phi_0$ and no slow roll-over would take place. However, as was pointed out by Albrecht and Brandenberger[26)], the domain formation rate will depend on the strength of the inflaton interaction rate as compared to the rate of the Hubble expansion which will red-shift large fluctuations. According to Albrecht and Brandenberger the correct initial state will be obtained for small enough coupling of the inflaton, which entails that the position of the global minimum should be larger than $O(M_{P\ell})$. Also our simple model with small ξ will favour the formation of a correct initial state: in particular for $\xi \simeq 3\times10^{-3}$ we get $\phi_0 \simeq 10^2\ M_{P\ell}$. The problem of initial conditions in inflationary models based on maximally symmetric supergravity is presently under investigation[27].

ACKNOWLEDGEMENTS

One of us (K.E.) gratefully thanks the Academy of Finland for financial support.

ξ	$\phi_i/M_{P\ell}$	$\phi^*/M_{P\ell}$	$\phi_f/M_{P\ell}$	$\phi_m/M_{P\ell}$	$(\pi^2 g_*/30)^{1/4} T_{RH}$	Δ	N_t
0.35	2.8×10^{-11}	1.7×10^{-7}	0.66	0.81	2.2×10^{13} GeV	11.5	95
0.30	1.3×10^{-9}	8.8×10^{-6}	0.79	0.94	1.3×10^{14} GeV	6.1	109
0.25	2.8×10^{-8}	2.8×10^{-4}	0.97	1.13	5.6×10^{14} GeV	3.5	133
0.20	3.6×10^{-7}	5.7×10^{-3}	1.25	1.41	1.8×10^{15} GeV	2.2	183
0.15	2.2×10^{-6}	5.9×10^{-2}	1.70	1.86	3.9×10^{15} GeV	1.6	281
0.10	$8\ \times10^{-6}$	0.44	2.63	2.79	6.2×10^{15} GeV	1.3	588
0.05	2.7×10^{-5}	2.74	5.46	5.62	7.2×10^{15} GeV	1.14	2.5×10^{3}

TABLE Numerical values of $\phi_i = H_0$, ϕ_*, ϕ_f, ϕ_m, $(T_{RH})_{max}$, Δ and N_t for a sample of values of ξ.

REFERENCES

1) A.H. Guth. Phys.Rev. D23 (1981) 347;
D. Kazanas, Ap.J. 241 (1980) L59;
K. Sato, Mon.Not.Roy.Astr.Soc. 195 (1981) 467; Phys.Letters 99B (1981) 66.

2) A.D. Linde, Phys.Letters 108B (1982) 289;
A. Albrecht and P.J. Steinhardt, Phys.Rev.Letters 48 (1982) 1220.

3) J. Ellis, D.V. Nanopoulos, K.A. Olive and K. Tamvakis, Phys.Letters 118B (1982) 335; 120B (1983) 331; Nuclear Phys. B221 (1983) 524;
D.V. Nanopoulos, K.A. Olive, M. Srednicki and K. Tamvakis, Phys.Letters 123B (1983) 41;
D.V. Nanopoulos, K.A. Olive and M. Srednicki, Phys.Letters 127B (1983) 30;
G.B. Gelmini, D.V. Nanopoulos and K.A. Olive, Phys.Letters 131B (1983) 53.

4) B.A. Ovrut and P.J. Steinhardt, Phys.Letters 133B (1983) 161.

5) H. Pagels and J.R. Primack, Phys.Rev.Letters 48 (1982) 223;
S. Weinberg, Phys.Rev.Letters 48 (1982) 1303.

6) J. Ellis, A.D. Linde and D.V. Nanopoulos, Phys.Letters 118B (1982) 59.

7) M.Yu. Khlopov and A.D. Linde, Phys.Letters 138B (1984) 265;
J. Ellis, J.E. Kim and D.V. Nanopoulos, Phys.Letters 145B (1984) 181;
J. Ellis, D.V. Nanopoulos and S. Sarkar, CERN preprint TH.4057 (1984).

8) For a recent review, see:
K. Enqvist and D.V. Nanopoulos, CERN preprint TH.4066 (1984).

9) D.V. Nanopoulos and K. Tamvakis, Phys.Letters 114B (1982) 235;
A. Masiero, D.V. Nanopoulos, K. Tamvakis and T. Yanagida, Z.für Phys. C17 (1983) 33;
J. Kim, A. Masiero and D.V. Nanopoulos, Phys.Letters 145B (1984) 187.

10) J. Ellis, C. Kounnas and D.V. Nanopoulos, Nuclear Phys. B241 (1984) 406; B247 (1984) 373; Phys.Letters 143B (1984) 410.

11) J. Ellis, K. Enqvist and D.V. Nanopoulos, Phys.Letters 147B (1984) 99; 151B (1985) 357.

12) E. Witten, Princeton preprint (1985);
M. Dine, R. Rohm, N. Seiberg and E. Witten, Princeton preprint (1985).

13) M.B. Green and J.H. Schwarz, Phys.Letters 149B (1984) 117;
D.J. Gross, J.A. Harvey, E. Martinec and R. Rohm, Phys.Rev.Letters 54 (1985) 502.

14) P. Candelas, G. Horowitz, A. Strominger and E. Witten, Santa Barbara preprint (1984).

15) G.B. Gelmini, C. Kounnas and D.V. Nanopoulos, Nuclear Phys. B250 (1985) 177.

16) K. Enqvist and D.V. Nanopoulos, Phys.Letters 142B (1984) 349; Nuclear Phys. B252 (1985) 508.

17) C. Kounnas and M. Quiros, Phys.Letters 151B (1985) 189;
J. Ellis, K. Enqvist, D.V. Nanopoulos, K. Olive and M. Srednicki, Phys.Letters 159B (1985) 175.

18) K. Enqvist, C. Kounnas, D.V. Nanopoulos and M. Quiros, CERN preprint TH.4027 (1984).

19) K. Enqvist, C. Kounnas, D.V. Nanopoulos and M. Quiros, CERN preprint TH.4048 (1984).

20) K. Enqvist, D.V. Nanopoulos, K. Olive, M. Quiros and M. Srednicki, Phys.Letters 152B (1985) 181.

21) A.S. Goncharov and A.D. Linde, Class.Quant.Grav. 1 (1984) L25.

22) P.J. Steinhardt and M.S. Turner, Phys.Rev. D29 (1984) 2162.

23) S.W. Hawking and I.G. Moss, Phys.Letters 110B (1982) 35;
G. Gelmini and L. Mizrachi, CERN preprint TH.4047 (1984).

24) A.D. Linde, Phys.Letters 129B (1983) 177; 132B (1983) 317; Rep.Prog.Phys. 47 (1984) 925.

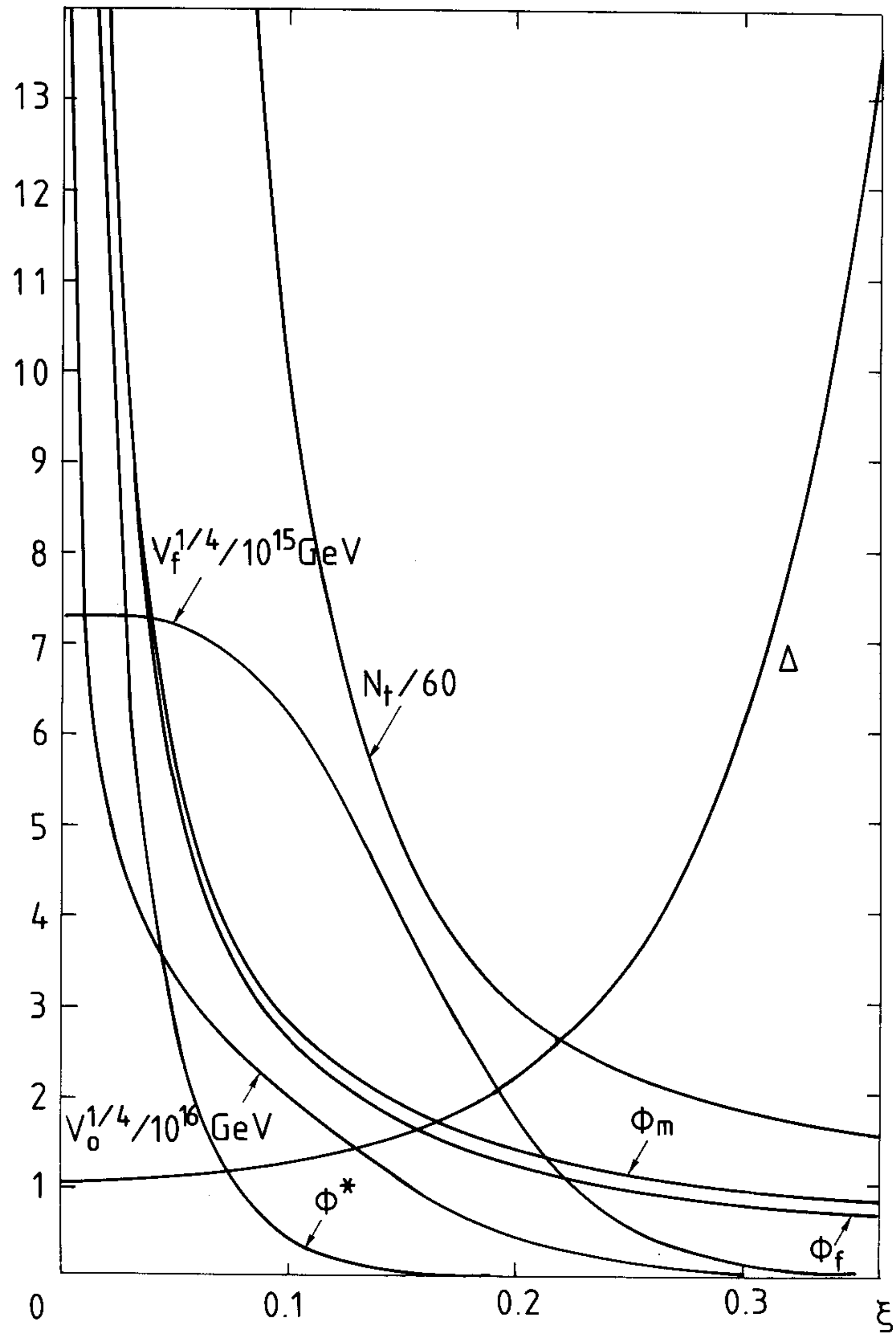
25) G. Mazenko, W. Unruh and R. Wald, Phys.Rev. D31 (1985) 273.

26) A. Albrecht and R. Brandenberger, Phys.Rev. D31 (1985) 1225.

27) K. Enqvist, C. Kounnas, D.V. Nanopoulos and M. Quiros, paper in preparation.

FIGURE CAPTION

Graphs of ϕ_m, ϕ_f, ϕ_*, $V_f^{1/4}$, $V_0^{1/4}$, N_t and Δ as functions of ξ (for definitions, see text); ϕ_m, ϕ_f and ϕ_* are in units of $M_{P\ell} = 1.2\times10^{19}$ GeV, and the total roll-over scale N_t is expressed in units of $(N_t)_{min} \simeq 60$.



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Δ
Φ_m
Φ_f
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