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OF SUPERSTRING INSPIRED MODELS

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Inflation and Other Cosmological Aspects of Superstring Inspired Models * †

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Abstract

Investigations of the possibilities for inflation using the four dimensional models inspired by superstring theory is reported and discussed in the context of related work.

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INFLATION AND OTHER COSMOLOGICAL ASPECTS OF SUPERSTRING INSPIRED MODELS

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Investigations of the possibilities for inflation using the four dimensional models inspired by superstring theory is reported and discussed in the context of related work.

Introduction. Attempts to realize a successful inflationary scenario have led to the elaboration[1] of a number of conditions that must be met in the standard field theoretic context where a period of exponential expansion arises from a nonvanishing vacuum energy associated with the vacuum expectation value (vev) of some scalar field ϕ, called the inflaton. Among these conditions are
1) A sufficient number of e-foldings

$$N_e \simeq H\Delta\tau \gtrsim 65 \tag{1}$$

to solve the flatness and horizon problems. In Eq. (1) $\Delta\tau$ is the duration of the inflationary epoch and

$$H = \frac{\dot{a}}{a} = \sqrt{\frac{V(\phi_0)}{3}}/m_{Planck} \tag{2}$$

is the value of the Hubble constant at the start of inflation, and is dominated by the vacuum energy $V(\phi_0)$ stored in the false vacuum $<\phi>_{av} \equiv \phi_0$. 2) Density fluctuations $\delta\rho/\rho$ arising from spatial fluctuations in the scalar field ϕ during inflation should be small enough to be consistent with the observed homogeneity in the microwave background and, ideally, large enough to account for galaxy formation:

$$\delta\rho/\rho \simeq 10^{-5} - 10^{-4}. \tag{3}$$

In the case that there is a post-inflationary source of density fluctuations such as cosmic strings, Eq. (3) can be reinterpreted as an upper limit: $\delta\rho/\rho \lesssim 10^{-4}$.

The inflaton potential $V(\phi)$ is roughly characterized by two parameters: the value σ of the inflaton ϕ at the global minimum of the zero temperature potential,

$$V'(\sigma) = 0 = V(\sigma),\ (V(\phi) \geq 0), \tag{4}$$

where the second equality in (4) is assumed to assure a vanishing cosmological constant today, and the value

$$\mu^4 = V(\phi_0) \tag{5}$$

of the energy density stored in false vacuum at the onset of inflation. Typically, the quantities N_e and $\delta\rho/\rho$ in specific inflationary models are determined by the parameters σ and μ as

$$N_e \simeq \sigma^m/\mu^n, \quad \delta\rho/\rho \simeq \mu^{n'}/\sigma^{m'}, \tag{6}$$

where m, n, m', n' are small positive integers. Since we require large N_e and small $\delta\rho/\rho$, the constraints (1) and (3) are usually met by choosing[1] large σ and small μ:

$$\sigma \sim 1,\ \mu \sim 10^{-4} - 10^{-3} \tag{7}$$

in Planck mass units (that I shall use throughout unless otherwise specified: $m_P = (8\pi G_F)^{-1/2} \equiv 2 \times 10^{18}$ GeV $\equiv 1$), with a potential that is rather flat over a broad region,

$$V(\phi) \sim \mu^4,\ \phi_0 \lesssim \phi \lesssim \phi_e,\ \phi_e - \phi_0 \sim 1, \tag{8}$$

so that the vacuum energy density remains appreciable until the field ϕ rolls to and "end-of-inflation" value ϕ_e after which it rapidly falls (via damped oscillations) to its true ground state value σ. As μ characterizes the self-couplings of the inflaton, the small value (7) of μ implies that it is very weakly coupled.

The required properties of the potential are most easily implemented in the context of $N = 1$ supergravity. As this theory is not renormalizable, the potential is an a priori arbitrary func-

tion of ϕ and so can be adjusted to the desired shape, albeit with considerable fine tuning. In addition, once one introduces a gauge singlet field that couples to ordinary matter only with interactions of gravitational strength, there is a natural mechanism for satisfying a third criterion[1] for a successful inflationary scenarios, namely: 3) Baryogenesis. The reheating temperature after inflation must be sufficient to allow for an out-of-equilibrium distribution of heavy particles that can generate a net baryon number in their decays. If the inflaton has only (generalized) gravitational couplings, it will decay into the heaviest particles that are kinematically available, thus possibly producing the desired result.

The bad news for supergravity models is their failure, in general, to fulfill a fourth criterion, namely to insure initial conditions conducive to an inflationary epoch, that is,

4) The thermal constraint. In renormalizable theories the high temperature potential is modified by quantum corrections that add a term $A\phi^2T^2$, $A \geq 0$, to the effective potential. In a non-renormalizable model the leading high temperature correction is of the form $f(\phi)T^2$ where $f(\phi)$ can be an arbitrary function of the inflaton field since the functional form of the tree Lagrangian is arbitrary. In many supergravity models the leading high temperature correction to the potential takes the form:[2]

$$\Delta V(\phi,T) \simeq \frac{T^2N}{12}\Big(V(T=0)+0(1/N)\Big),$$

where N is the total number of chiral supermultiplets ($N \gtrsim 50$ for a realistic model), so the minimum at high temperature is not significantly shifted from the true vacuum.

There are other problems[1] that must be avoided in a successful inflationary scenario, for example, an excessive production of gravitinos during inflation. In this talk I will focus primarily on points 1, 2, and 4. I will describe work[3] in collaboration with Binétruy, and related work[4-6] that has addressed the question as to whether the locally supersymmetric models suggested by string theories[7-9] have the right ingredients to satisfy these criteria. I will also comment on the problem of baryogenesis.

Superstring inspired models. These models possess several of the desired features for a successful inflationary cosmology. They contain[10,11] a number of gauge-singlet scalars that couple to ordinary matter with gravitational strength. Among these are a dilaton ϕ and the "breathing mode" e^σ which, together with axion-like fields β and h, form the complex scalars

$$\begin{aligned} S &= \phi^{-3/4}e^{3\sigma} + 3i\sqrt{2}h \\ T &= \phi^{3/4}e^{\sigma} - i\sqrt{2}\beta + k\,|y|^2 \end{aligned} \qquad (9)$$

of two chiral supermultiplets. In (9) $|y|^2 = \sum |y_i|^2$, $i = 1,\cdots N$, where the y_i are the scalar components of N chiral (gauge non-singlet) matter supermultiplets. Among the possible gauge singlets, the fields S and T in (9) have received the most attention in studies of phenomenology.

Superstring inspired models further embody naturally flat directions in the spin-zero field space. The scalar potential[10] is a sum of terms of the form

$$V(\phi) \ni \frac{1}{(S+S^*)}\frac{1}{(T+T^*-k\,|y|^2)^n}G(y) \qquad (10)$$

where $G(y) \sim y^4$ + higher order terms, and n is an integer: $n = 1-3$. $G(y)$ is either a D-term, implied by supersymmetry from gauge couplings, or an F-term, derived from a superpotential,[10] $W(y) \sim y^3$+ higher order. $G(y)$ is positive-semi-definite, so the minimum of the potential occurs for $y_i = 0$, with S and T undetermined.

In addition, the E_6 models[9] inspired by the heterotic[8] string include among the matter fields y singlets N of the observed gauge group $SU(3)_c \times SU(2)_L \times U(1)$ that have flat directions at least in the F-terms of the potential because it is derived from a superpotential $W(y)$ that does not allow terms $W \ni N^3$; such terms are forbidden by extra $U(1)$'s of the low energy gauge group.[9]

In a realistic model the vacuum degeneracy associated with these flat directions must be lifted because supersymmetry and any additional gauge $U(1)$'s must be broken. Supersymmetry could be broken by non-perturbative effects. In the heterotic string inspired models there is a hidden sector described by a pure $N=1$ Yang Mills theory with E_8 or some sub-group thereof as gauge group. The hidden sector is infrared strong and its gauginos acquire a vev $< \bar{\lambda}\lambda > \neq 0$ at some scale $\Lambda_c < m_{Plank}$; this vev is characterized by a parameter h. In addition the 10-dimensional theory includes a totally antisymmetric tensor that can have a nonvanishing vev $< H_{mnp} > \neq 0$ (where m,n,p are Lorentz indices in the compactified 6-space), characterized by a constant c that is in fact a quantized number that depends on the topology of the compactified manifold.

Either $c \neq 0$ or $h \neq 0$ breaks supersymmetry. When both of these parameters are nonvanishing,

the additional term induced in the potential takes the form of a squared quantity such that the S-field (which enters because it couples to all gauge and gaugino fields) can acquire a vev such that the (tree level) potential still vanishes at its minimum. Specifically, the effective tree potential in four dimensions acquires a term:

$$V \ni \frac{1}{(S+S^*)}\frac{1}{(T+T^*-k\,|y|^2)^3}\times$$
$$\left|W(y)+c+h(1+\alpha)^{-\alpha/2}e^{-i\beta/2}\right|^2 \qquad (11)$$

where

$$\alpha = 3ReS/b_0, \qquad \beta = 3ImS/b_0 \qquad (12)$$

and b_0 determines the β-function of the hidden gauge group. The term (11) fixes the vev of S, but T remains undetermined at tree level, and so also the gravitino mass, proportional[13] to $(T+T^*)^{-3/2}$.

I will first assume that ReT, and the gravitino mass $m_{\tilde{G}}$, are somehow fixed and study the potential as a function of S. I will discuss later the possible determination of ReT. Setting T at its presumed ground state value and $y =< y >= 0$, we obtain the potential for the S-field.

$$V(\alpha,\beta) = \mu^4\frac{\alpha_0}{\alpha}\left[1+\frac{1}{\hat{c}^2}(1+\alpha)^2e^{-\alpha}\right.$$
$$\left.-\frac{2}{\hat{c}}(1+\alpha)e^{-\alpha/2}\cos\beta/2\right], \qquad (13)$$

where $\alpha_0 = 3/4\pi b_0\alpha_{GUT}$ and $\beta_0 = 4\pi n$ are the ground state values of α,β, Eq. (12), and we have chosen $c<0$ to avoid CP violation through the coupling[13] (ImS)$F\tilde{F}$ where F is the gauge field strength. $< ReS >$ determines the gauge coupling α_{GUT} at the scale of compactification through its coupling to the gauge-gaugino supermultiplet. The parameter $\hat{c}$ in (13) is the ratio of vevs:

$$\hat{c} = -c/h = (1+\alpha_0)e^{-\alpha_0/2} \qquad (14)$$

and the energy that might be stored in a false vacuum is governed by the scale μ:

$$\mu/m_P = \left[\frac{m_{\tilde{G}}}{m_P}\frac{1+\alpha_0}{\alpha_0}\right]^{1/2} \qquad (15)$$

which has the desired order of magnitude $\mu \sim 10^{-3}$ for $m_{\tilde{G}} \sim 10^{12}$ GeV and $\alpha_{GUT} \simeq 0.3$. This is a rather larger gravitino mass than in conventional SUSY and SUGRA models such as those considered in the previous talk by Ovrut, and one might worry about the stability of scalar masses. It has been shown[14,15] that in these models scalars remain massless at one loop. Gauginos apparently[16] acquire masses $m_{\tilde{g}} \sim m_{\tilde{G}}^3/16\pi^2m_P^2$, suggesting that scalars should acquire similar masses at the 2-loop level that can be below a TeV for $m_{\tilde{G}}$ as high as 10^{12} GeV.

Is β the inflaton? Setting $\alpha = \alpha_0$, the potential (13) for β becomes

$$V(\alpha_0,\beta) = 4\mu^4\sin^2\beta. \qquad (16)$$

Finite temperature quantum corrections to (17) do not shift the position of its minimum[3], so the thermal constraint cannot be satisfied. If one assumes that some other mechanism stabilizes field β at an initial value $\beta_0 = 4\pi n + 2\pi$, one finds[3] for the number of e-foldings during which the vacuum energy density dominates ($\alpha_0 \geq 1$, see fig. 1):

$$N_e = -\frac{4}{\alpha_0^2}\ln\sin\left(\frac{\alpha_0}{\sqrt{6}}\frac{\mu^2}{m_P^2}\right) << 65.$$

Is α the inflaton? This possibility was briefly considered by Maeda and Pollock[4] who concluded that the potential did not satisfy the slow rollover[1] criteria. In ref. 3 the potential (13) was considered in more detail. Setting $\beta = \beta_0$, the potential (13) for α takes the form

$$V(\alpha,\beta_0) = \mu^4\frac{\alpha_0}{\alpha}[1-\frac{1}{\hat{c}}(1+\alpha)e^{-\alpha/2}]^2 \qquad (17)$$

which is plotted in Fig. 1 for representative values of $\hat{c}$. The potential (17) vanishes for $\alpha \to \infty$, a general feature[17] of potentials generated by nonperturbative SUSY breaking. Then the initial value of α could fall to the right of the maximum α_2 and subsequently roll off to infinity. One possibility for avoiding[18,5] this disaster is the existence of Q-balls[19] in the early universe, where Q is the charge associated with the nonhomogeneous transformation $s \to s+i\gamma$. Q-balls are spatial volumes with nonvanishing $< \partial_\mu\beta\partial^\mu\beta^\dagger >$ and hence net charge Q. Because of the non-canonical form of the kinetic energy:

$$\mathcal{L}_{KE} = \frac{1}{(S+S^*)^2}\partial_\mu S\partial^\mu S^* = \frac{1}{4\alpha^2}[(\partial\alpha)^2+(\partial\beta)^2] \qquad (18)$$

that couples α to the kinetic term for β, this vev induces[18] an effective potential for α, proportional to $Q(S+S^*)^2$ that can drive α to zero. This mechanism might suffice to fix the initial value at $\alpha_i > \alpha_0$.

The conventional mechanism for establishing initial conditions conducive to inflation is through quantum corrections to the finite temperature po-

tential. In the model considered here, the finite temperature corrections to the potential are positive definite and of the form:[3]

$$\Delta V_T(\alpha,T) = \frac{T^2}{24}[(2N+N_G+16)V(\alpha,T=0)+0(1)] \tag{19}$$

where N_G is the number of gauginos. Since N and N_G are both large $(0(10^2))$ numbers the minimum of the potential cannot be significantly shifted at high temperature.

In the context of other inflationary supergravity models, Holman et al.[20] have stressed that scalars with interactions of only gravitational strength are thermally decoupled below the Planck mass down to temperatures well below that at which the energy of the false vacuum should dominate. They argue that such a system – in the present case S, T – should be treated as a closed system isolated from the rest of matter. If, following this line of reasoning, we compute the high temperature corrections neglecting the gauge and chiral degrees of freedom (although it seems difficult to argue that they would not reappear through precisely the quantum corrections we are considering) we obtain the result of eq. (19) with $N = N_G = 0$, in which case the 0(1) term can be competitive with the term proportional to $V(\alpha, T = 0)$. It turns out[3] that this truncated potential tends to stabilize the initial value α_i of α near the inflection point (Fig. 1b) or second local minimum (Fig. 1c), α_1 for the parameter range $0.937 \lesssim \hat{c} < 1$.

It is clear from inspection of Fig. 1 that, whatever the origin of the initial value α_i, the only chance for an acceptable inflationary scenario is for just this range of parameters with $\alpha_i \simeq \alpha_1$. If instead we assume $\alpha_i \simeq \alpha_2$, where α_2 is the position of the maximum $\alpha_2 > \alpha_0$ for any $\hat{c}$, domains of both $\alpha < \alpha_2$ and $\alpha > \alpha_2$ will form as the universe cools; the latter will expand indefinitely and ultimately dominate. For similar reasons chaotic[21] initial conditions would not appear to induce inflation for the potentials of Fig. 1.

We therefore assume

$$\hat{c} = .937 + \delta\hat{c} < 1, \quad \alpha_i = \alpha_1. \tag{20}$$

As we let $\hat{c} \to 1$, the energy stored in the false vacuum becomes too small to be useful and the barrier to be tunnelled across becomes dangerously large. The analysis of Hawking and Moss[22] shows that the effective mass $m_\alpha^2(\alpha_1) = \partial^2 V(\alpha)/\partial\alpha^2\Big|_{\alpha=\alpha_1}$ should be limited to

$$m_\alpha^2(\alpha_1) < 2H^2 = 2V(\alpha_1)/3m_P^2. \tag{21}$$

When the bound (21) is exceeded tunnelling is too slow for nucleation of the bubbles of true vacuum to occur and one encounters the difficulties of "old" inflationary models.[1] For $\hat{c} = 1, \alpha_1$ is an inflection point: $m_\alpha^2(\alpha_1) = 0$. As we turn on $\delta\hat{c} > 0, m_\alpha^2(\alpha_1) > 0$ and the bound (21) is satisfied only for

$$\delta\hat{c} < 3 \times 10^{-4} \tag{22}$$

which requires strong fine tuning of the potential. In fact $\hat{c}$ is not a parameter to be tuned. It depends on the topology of the compact manifold (c) and the on the hidden gauge group (h); Eq. (22) is the statement that successful inflation requires very specific values for these parameters.

Assuming optimistically that $\hat{c}$ is very close to the critical value .937, the number of e-foldings[1]

$$N_e = -\int_{\alpha_i}^{\alpha_e} \frac{3H^2(\alpha)}{2\alpha^2 V'(\alpha) m_P^2} d\alpha \tag{23}$$

can be calculated. Taking $\delta\hat{c} = 0$ and $\alpha_i = \alpha_1$ gives[3] $N_e > 65$ for $\mu/m_P \lesssim 0.14$, which is not a stringent constraint on μ. However, a successful inflationary scenario is still not established. Since $V(\alpha_1) > V(\alpha_2)$ in Fig. 1b, (or 1c if $\delta\hat{c} < 3\times10^{-4}$), there is a danger that α would overshoot the barrier at α_2 if the friction due to the expansion of the universe is insufficient to prevent the classically rolling field from reaching the value α_2.

Finally, we must consider the initial conditions that determine the shape of the zero-temperature potential itself. The gaugino condensate of the strongly coupled hidden gauge sector forms only below a critical temperature $T \simeq \Lambda_c$. On the other hand the parameter c can take only discreet values that can change by vacuum tunnelling. In Ref. 3 we followed Rohm and Witten[23] who argue that as long as $< \lambda\bar{\lambda} > \propto h = 0$, the H-vacuum chooses the lowest energy SUSY preserving configuration $< H >= 0$, and at $h \neq 0, < H >$ follows by tunneling so as to again achieve a vanishing vacuum energy. Thus in Ref. 3 it was assumed that $h = c = 0$ above the critical temperature $T_c \simeq \Lambda_c$ and $h, c \neq 0$ below T_c. When the gaugino condensate first forms there will be energy stored in the false H-vacuum, but presumably tunnelling will occur too rapidly to allow for significant inflation; otherwise one could encounter the nucleation and inhomogeneity problems of "old" inflation.

Ellis et al.[5] assume that $< H >$ is fixed at its present non-zero value above the critical tem-

perature T_c and treat $< \bar{\lambda}\lambda >$ as a temperature dependent quantity. Then $|\hat{c}| = |c/\lambda|$ decreases with temperature so the potential has a time dependent shape that for $\hat{c}(T) < 1.21$ runs backwards from Fig. 1f towards Fig. 1a, ending at the potential for $\hat{c}(0) = \hat{c}$. They argue that for $c(T) = 1.21$ α would be stabilized at α_0' (Fig. 1f) and evolve through α_0' (Fig.1e) to α_1 (Fig. 1c) where it could be stabilized with nonvanishing vacuum energy before tunnelling à la Hawking and Moss[22] to the zero temperature ground state value α_0. Their scenario also requires $\hat{c}$ close to 0.937, with a similar fine tuning condition.

Is ReT the inflaton? To consider this possibility requires knowledge of the effective potential for the T-field that determines its ground state value and thus the gravitino mass. The vacuum degeneracy is lifted[14,15] by one loop quantum corrections which must be regularized by the introduction of a cut-off of order of the condensate scale Λ_c. In ref. 15 the effective one-loop potential was calculated neglecting terms of order $(m_i^2(\phi)/\Lambda_c^2)$ where $m_i^2(\phi)$ are field-dependent masses. This potential has a global minimum at a finite value of ReT which determines the relevant scales of the theory as a function of α_{GUT} and b_0. For a range of these parameters such that $m_{GUT} < m_{Planck}$, one finds[15] $m_{\tilde{G}} > 10^{15}$ GeV. The results of Ref. 16 suggest that this would entail gaugino masses $m_{\tilde{g}} \sim m_{\tilde{G}}^3/16\pi^2 \simeq 10^{-11} m_P \simeq 10^7$ GeV and hence a Higgs scalar mass $m_H \sim m_{\tilde{G}}\sqrt{\alpha_{weak}/4\pi} \gtrsim 10^6$ GeV. In addition, the parameter (15) that governs the scale of the false vacuum has in this case a value $\mu/m_P \gtrsim 0.05$ for $\hat{c} = .937 (\alpha_0 = 2.8)$ which would lead to excessive density fluctuations for the S-inflaton scenario. It was subsequently argued[24] that, since the values $m(\phi)$ at the minimum are comparable to Λ_c, one should keep the full cut-off dependence in the effective potential. It turns out that the $O(m^2/\Lambda_c^2)$ terms destabilize* the potential in the direction $ReT \to 0$, $m_{\tilde{G}} \to \infty$. Since when these corrections are included there is no longer even a local minimum in the region $m(ReT) < \Lambda_c$, the potential for ReT must be considered as unknown at present.

Maeda et al.[6] nevertheless assumed an effective potential

$$V(\phi) = V_{Tree}(\phi) + V_{BG}(\phi) + \text{constant}. \qquad (24)$$

*However, contributions from loop momenta $\Lambda_c^2 \lesssim |p|^2 \lesssim m_{GUT}^2$ may restabilise the potential[25].

where V_{BG} is the one loop contribution calculated in Ref. 15 that has a negative cosmological constant which is cancelled by the constant added in (24). Maeda et al. argue that as long as $S \propto (\alpha, \beta) \neq (\alpha_0, \beta_0)$, $V_{Tree}(\phi)$ is nonvanishing and will drive (c.f. eq. (11)) ReT towards infinity. Once $< S >$ has settled at its ground state value, ReT will roll to its ground state. They find that the conditions for successful inflation are satisfied for a broad range of the parameters c and h.

However, as the authors of Ref. (6) note, the ansatz (24) does not satisfy the general result[17] that the potential should vanish as $ReT \to \infty$.

Is a matter field (N) the inflaton? In the heterotic string inspired models matter fields fill fundamental 27 representations of E_6 that include two singlets (N and N' or ν^c that I will generically call N) of the observed gauge group. The renormalizable F-terms have flat directions for $N \neq 0$, but the D-terms are not flat in the $N \neq 0$ direction if all matter fields transform as 27's. However, the effective low energy theory may contain 27 + *overline*27 pairs in addition to the 27's that comprise the standard families. The direction $N = \bar{N} \neq 0$ is a flat direction for the D-term of the potential as well as of the F-term. This degeneracy should be lifted by effective nonrenormalizable terms in the superpotential that arise[26] when one integrates out the heavy degrees of freedom of the theory. These give an effective potential of the form:

$$\begin{aligned} V_{eff} &\sim m^2(|N|^2 + |\bar{N}|^2) \\ &+ \lambda_{eff}^2(|N\bar{N}|)^{2n-2}(|N|^2 + |\bar{N}|^2) \\ &+ D - \text{terms} + \cdots \end{aligned} \qquad (25)$$

with $n \geq 2$. The potential (25) includes a SUSY breaking mass term for scalars that is probably[16] of order $m \sim m_{\tilde{G}}^3/16\pi^2 m_P^2$. The squared mass parameter has a logarithmic dependence on the field strength $|N|$ and, for suitable N-couplings, approaches a negative fixed point $-m_0^2$, triggering a vev for $N(\bar{N})$. Since one expects all scalars to have comparable SUSY breaking masses, we require $m_0^2 \simeq (m_W^2$ to $(1\text{ TeV})^2)$ if the Higgs mass is to be in the range required for the observed Higgs vev without strong couplings. This means that a $U(1)$ gauge invariance is broken at a critical temperature $T_c \sim (m_W - 1\text{ TeV})$, while quantum corrections stabilize the field at $N = 0$, for $T > T_c$. The onset of inflation occurs when the vacuum energy density dominates over the radia-

tion density, i.e. for $T_i \sim (V_0)^{1/4} \simeq \mu_0$, and inflation continues until the critical density T_c for the symmetry breaking phase transition is reached. Since the temperature drops exponentially during inflation ($T \simeq T_i e^{-Ht}$) the number of e-foldings is given simply by [3]

$$N_e \simeq \ln(T_i/T_c) \simeq \ln(\mu_0/m_W) \simeq 6-16 \qquad (26)$$

for $n_0 \geq 2$, which is insufficient to solve the flatness and horizon problems. Another possibility is that the flat directions are lifted only by non-perturbative effects on the string world sheet[27] in which case the effective coupling λ_{eff} in (25) is of order $\exp(-2 < ReT >)$, increasing[3] the number of e-foldings (26) by about $1/4 < \mathrm{Re}\,T >$ which is still insufficient for plausible values of $< \mathrm{Re}\,T >$.

If the inflaton field decays rapidly after the end of inflation the fractional entropy increase is related to the number of e-foldings by $s_f/s_i = e^{3N_e} \simeq 10^8 - 10^{20}$ for the range (26). However, since with the above choice of parameters N is in fact weakly coupled to ordinary matter, the entropy release can in fact be considerably larger.[28] Such a late ($T \lesssim T_c \sim m_W$) entropy release has raised the issue[28,29] as to whether the phase transition associated with an intermediate scale may cause serious problems for standard cosmology.

It has recently been pointed out[28,30] that the phase transition will actually occur at a much higher temperature. The one-loop quantum corrections to the finite temperature potential can be expressed as $\Delta V(\phi,T) = \sum_i v_i(m_i(\phi),T)$ where the index i runs over all spin and internal symmetry degrees of freedom and $m_i(\phi)$ is the appropriate field-dependent mass. For $T_i^2 >> m_i^2(\phi)$ this expression reduces to the usual high temperature approximation

$$v_i = -\frac{T^4}{90}\pi^2(1-\frac{1}{8}[1-(-)^{F_i}]) + \frac{T^2}{24}m_i^2(\phi), \quad (27)$$

where F_i is fermion number. The point is[28,30] that (27) is an approximation valid only in the high temperature limit, while the exact expression falls rapidly to zero as soon as $m_i^2(\phi) > T^2$. For fixed T the potential possesses a barrier enclosing the false vacuum (the origin for the N-fields) and coincides with the zero temperature potential for values of the scalar fields ϕ such that $m_i^2(\phi) >> T$. As the temperature decreases the position of the barrier moves closer to the field origin. As soon as it is close enough that thermal fluctuations are sufficiently large ($\Delta\phi = 0(\phi_{barrier})$) to "spill" the field ϕ over the barrier with a significant probability, domains of true vacuum will form. If this "spill-over" occurs during an inflationary epoch (caused presumably by some other scalar sector; in this case $\Delta\phi \sim H$) domains of true N-vacuum could grow large enough to encompass our present observable universe with no need for a later phase transition.[30]

Other gauge-singlet fields? The models inspired by superstrings contain additional structurerelated gauge singlet chiral supermultiplets namely[23] b_{11} fields $T^{(k)} = \alpha^{(k)} + +i\beta^{(k)}$ and b_{21} fields $C^{(\alpha)}$ where b_{11} and b_{21} are topological numbers. One of the b_{11} $T^{(k)}$ fields is the T-field discussed above for which there is no superpotential at tree level and, in fact, all the $\beta^{(k)}$ are axion-like objects, implying that, at tree level at least, there can be no superpotential for the $T^{(k)}$ fields. On the other hand, the ground state values of the $C^{(\alpha)}$ are determined by $< H >$ which fixes the cohomology class of the the complex structure of the compact manifold, so there is a tree level potential for these fields. In the absence of a model for the shape of the potential, this possibility cannot be analyzed further.

Baryogenesis. In superstring inspired models, the baryon number violating X-bosons, responsible for baryon decay in standard GUTs, have masses of the order of the Planck scale. They cannot play any role in generating a net baryon number for the observed universe if there is inflation in the epoch $T < m_P$ where the effective four-dimensional theories are valid. These four-dimensional theories allow two distinct scenarios: one with an "intermediate scale" associated with lifting the flat direction in N-space, and one in which there is no intermediate scale, with $< |N| >$ of the order (albeit higher by a factor of at least 2-3) of the electroweak scale.

In models inspired by the heterotic string left handed components of matter fields fall into 27-plets of E_6: $27_L = (\bar{5}+10+5+\bar{5}+1+1)_L$ under $SU(5)$. The $\bar{5}+10$ represent ordinary quarks and leptons and their sfermion partners; $5+\bar{5}$ contain the SUSY standard model Higgs doublets $H, \bar{H}$ plus their color triplet, $SU(2)_L$ singlet, counterparts $(g,\bar{g})$, and two SU(5) singlets (N). The superpotential must contain a term $W(\phi) \ni g\bar{g}N$ that gives a mass to the color triplets g when N acquires a vev. In addition, there are E_6 allowed terms $W(\phi) \ni gqq + gq\ell + \cdots$ that when combined can generate proton decay. If there is no intermediate scale, $m_g \sim m_W$, discreet symmetries in the

effective low energy renormalizable theory must be invoked to forbid couplings which would allow fast nucleon decay. Then one finds[31] that observable proton decay can proceed only if there are extra $\overline{27}$'s as well as yet more gauge singlets[11] (remnants of the E_8 gauge fields). This picture allows little prospect for post-inflationary baryogenesis via the standard mechanism of decaying color-triplet, super-heavy bosons.

If there is an intermediate scale with large vevs for the N-fields yielding very massive $g, \bar{g}$, their baryon/lepton number violating couplings can be allowed at tree level. Then the scalar members of these supermultiplets can play much the same role as the color-triplet Higgs of ordinary (or SUSY) GUTs, except that their couplings are not constrained by low energy data. Unless the couplings responsible for nucleon decay are somehow suppressed, limits on the nucleon lifetime probably require that they be too heavy to have been produced in sufficient abundance for efficient baryogenesis during postinflationary reheating. The reheating temperature is bounded by $T_R \lesssim \mu \lesssim 10^{-3}_P \simeq 10^{15}$ GeV, and is expected to be considerably lower for a very weakly coupled inflaton whose oscillations are strongly red-shifted before it can decay.

One can forbid nucleon decay by imposing lepton number conservation. Then baryon number violating processes could be mediated by objects sufficiently light to be copiously produced during post inflationary reheating. The problem is then to generate the out-of-equilibrium abundance necessary for their decays to generate a baryon asymmetry. A possible mechanism noted recently[32] is the dependence of the g-mass on the value of $< N >$. For $N = 0$ the g's are massless, with a mass that increases as N rolls towards its true ground-state value, and oscillates as the field N oscillates and decays. If the N-phase transition occurs after or near the end of inflation, an out of equilibrium g-distribution could be generated.

The difficulties for efficient baryogenesis have led some authors to consider alternatives that do not require explicit baryon number violating couplings. One approach[33] is based on the observation that instanton effects related to the anomalous baryon number and lepton number currents in the standard model are unsuppressed at high temperature and could have been important in the hot early universe. Another suggestion[34] is that immediately after inflation squarks and leptons may have had large vevs; a significant baryon asymmetry could be generated as this system evolves to the true vacuum state.

Baryogenesis could proceed according to conventional mechanisms if inflation took place at a temperature $T \gtrsim m_P$, prior to or during compactification. This possibility has been studied[4] but a convincing mechanism has not yet been found within the superstring context. If there is not an inflationary epoch after compactification, one is faced with the usual graviton, gravitino, monopole and domain wall (in particular those topological defects arising from compactification) problems.

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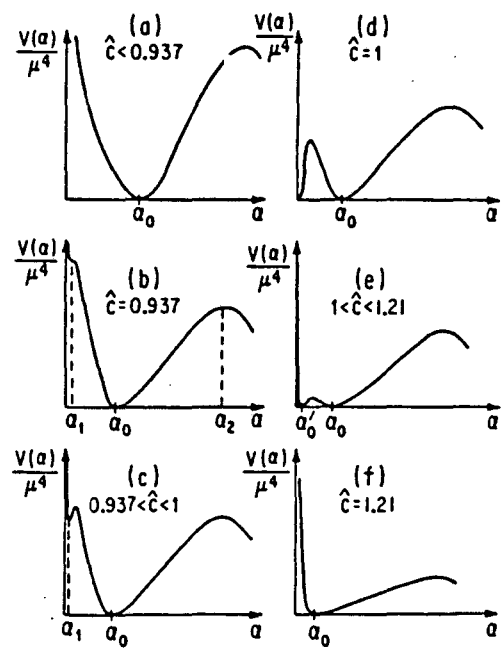

Fig. 1: Shape[6] of the potential of Eq. (18) for different values of $\hat{c}$, Eq. (15).

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