LBNL-40830
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Physics Division

September 1998

Submitted to
Physical Review D

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LANCS-TH/9705
LBNL-40830
UCB-PTH-97/48
hep-th/9806157
September 7, 1998

hep-th/9806157 v3 8 Sep 1998

Inflation and flat directions in modular invariant superstring effective theories*

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Abstract

The potential during inflation must be very flat in, at least, the direction of the inflaton. In renormalizable global supersymmetry, flat directions are ubiquitous, but they are not preserved in a generic supergravity theory. It is known that at least some of them are preserved in no-scale supergravity, and simple generalizations of it. We here study a more realistic generalization, based on string-derived supergravity, using the linear supermultiplet formalism for the dilaton. We consider a general class of hybrid inflation models, where a Fayet-Illiopoulos D term drives some fields to large values. The potential is dominated by the F term, but flatness is preserved in some directions. This allows inflation, with the dilaton stabilized in its domain of attraction, and some moduli stabilized at their vacuum values. Another modulus may be the inflaton.

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*This work was supported in part by PPARC grant GR/L40649 and NATO grant CRG 970214, in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098 and in part by the National Science Foundation under grant PHY-95-14797. HM was also supported by Alfred P. Sloan foundation.

I. INTRODUCTION

Cosmological inflation has been regarded as the most elegant solution to the horizon and flatness problems of the standard Big Bang universe. Even though it explains why the current Universe appears so homogeneous and flat in a natural manner, it has been difficult to construct a model of inflation without a small parameter. In fact, one needs a scalar field (inflaton) that rolls down the potential very slowly to successfully generate a viable inflationary scenario [1,2]. This requires the potential to be almost flat in the direction of the inflaton.

Other cosmological considerations may call for a flatness of the potential in non-inflaton directions. For instance, the Affleck–Dine baryogenesis scenario [3] requires a scalar field which carries baryon number to have a large amplitude to start with. To maintain a large amplitude during the rapid expansion of the universe, the scalar potential needs to either be almost flat [3,4], or to have a negative squared mass [5,6].[1]

In general, it appears rather unnatural to impose an almost flat scalar potential in quantum field theory, because such flatness is likely to be destroyed by radiative corrections. Specifically, scalar field masses are generally quadratically divergent and are not protected by any symmetries.[2] Supersymmetry, however, may maintain the flatness of a tree-level scalar potential against radiative corrections due to the nonrenormalization theorem.

A renormalizable, globally supersymmetric theory typically has several directions in which the tree-level potential is very flat. However, because inflation couples the energy density of a scalar potential to gravity to cause a rapid cosmological expansion, supersymmetry has to be made local: supergravity. During inflation, a generic supergravity theory lifts the flat directions of global supersymmetry, generating [7,8] a squared mass at least of order $3H^2 \simeq M_{\rm Pl}^{-2}V$ in magnitude.[3]

This generic result must be evaded for the inflaton field, [8] since in its direction $|V''| \ll M_{\rm Pl}^{-2}V$ is necessary for slow-roll inflation. (V'' is the second derivative of the inflaton

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[1]It should be emphasized, though, that we are looking at the global SUSY flat directions only *during* inflation. After inflation is over one has rapidly oscillating fields and/or thermalized particles, and a separate discussion is required in this much more complicated situation. One might argue [5] that all fields are likely to acquire masses at least of order H (until H falls below their true mass). If that is so, Affleck-Dine baryogenesis may proceed more or less as in [5] or [6], whether or not the directions responsible for it are flat during inflation. However, the thermal effects are exponentially suppressed if the amplitude of the scalar field is larger than the temperature. The effect of the oscillating field on the flat direction depends sensitively on the structure of the non-renormalizable Kähler potential terms. Discussion on these issues is beyond the scope of this paper.

[2]The exception is when the scalar field is a Nambu–Goldstone boson of a spontaneously broken global symmetry.

[3]Here, H is the Hubble parameter defined in terms of the scale factor a by $H = \dot{a}/a$, and $M_{\rm Pl} \equiv (8\pi G)^{-1/2} = 2.4 \times 10^{18}\,{\rm GeV}$ is the reduced Planck mass. The inflaton potential is $V(\phi)$, a prime denotes $d/d\phi$, and we have remembered that $V \simeq 3M_{\rm Pl}^2H^2$ during inflation.

potential, which cannot be much less than the squared mass along the inflaton trajectory.) One would like to understand how this evasion comes about, and whether it occurs for scalar fields other than the inflaton.

The question of whether a scalar potential has a flat direction consistent with inflation requires knowledge not only of the superpotential, which is not renormalized so that a specific form of it is at least technically natural, but also the Kähler potential, which can be renormalized by higher dimension operators with arbitrary coefficients in generic supergravity. In fact, one needs to know the Kähler potential at least up to quartic terms to determine if the potential is flat. Since these contributions are arbitrary in general supergravity, a natural question is whether an underlying quantum theory of gravity, such as superstring theory, determines a specific form of Kähler potential that ensures the flatness of the scalar potential along certain directions in the field space.

This question has been a difficult one to address, because of other cosmological problems in superstring-inspired supergravity theories. The dilaton field exhibits a runaway behavior and it has been difficult to obtain a minimum of the potential, consistent with spontaneously broken supersymmetry (as required by phenomenology) and vanishing cosmological constant (at least on the scale of supersymmetry breaking). Recently, a modular invariant formalism, based on string orbifold compactification, was proposed to study the stabilization of the dilaton by employing the linear multiplet description of the dilaton [9].

In the present paper, we explore the possibilities for inflation in the context of this formalism. Because we are writing down the supergravity Lagrangian after integrating out all the massive string and Kaluza–Klein excitations around a consistent vacuum, the superpotential has a power series expansion in the matter fields. It starts at the cubic order, and higher order terms are allowed with power suppression in the string scale. The effective mass terms and/or linear terms that are presumably necessary for inflation will appear when some of the fields acquire nonzero values (*vev*'s).[4] We suppose, following Stewart [10], that the *vev*'s are generated when a Fayet-Illiopoulos D term is driven to a small value. We show how this can preserve some of the flat directions of global SUSY, generating a potential flat enough for inflation, which will probably be of the hybrid [11,8] type. (During hybrid inflation a non-inflaton field is displaced from the vacuum, and is responsible for most of the potential.)

Our paper is organized as follows. In the next section, we introduce the special form of the supergravity Lagrangian obtained from superstrings in [9]. Section 3 contains the strategy for inflation model-building, and the main discussion on the flatness of the scalar potential. We conclude in Section 4.

We generally set $M_{\rm Pl} = 1$, where $M_{\rm Pl} = (8\pi G)^{-1/2} = 2.3 \times 10^{18}\,\mathrm{GeV}$.

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[4]Other intermediate-scale *vev*'s occurring in Nature (associated say with Peccei-Quinn symmetry, neutrino masses or a GUT) might be related to the inflationary one, or they might be in a different sector of the theory.

II. SUPERSTRING-DERIVED SUPERGRAVITY

We are going to show how to construct models of inflation, in the context of a realistic supergravity theory derived from the weakly coupled superstring [9]. It is based on a class of orbifold compactifications [12,13] with three untwisted three moduli t_I, and contains an effective potential for the dilaton, induced by gaugino condensation. Dilaton stabilization in the true vacuum is achieved by the inclusion of nonperturbative string effects in the Kähler potential, that modify the form of this potential.[5]

An important part of our program is to demonstrate that the dilaton can be stabilized during inflation, by the same nonperturbative string effects that stabilize it in the true vacuum. First though, we look at a simplified model which ignores the dilaton and the gaugino condensate. Then the scalar fields are all in chiral multiplets; they consist of the moduli t_I, and matter fields which we shall denote by ϕ_α.[6]

The tree-level potential has the usual form $V = V_F + V_D$. The F term is

$$V_F = e^K \left[\sum_{nm} K^{n\bar{m}} \left(W_n + K_n W \right) \left(\bar{W}_{\bar{m}} + K_{\bar{m}} \bar{W} \right) - 3|W|^2 \right] . \tag{1}$$

The superpotential W is a holomorphic function of the complex scalar fields, while the Kähler potential K is a function of the fields and their complex conjugates. A subscript n denotes the derivative with respect to the nth field, and $\bar{n}$ the derivative with respect to its complex conjugate. (In this context, n runs over both the both matter fields and the moduli.) The matrix $K^{n\bar{m}}$ is the inverse of the matrix $K_{n\bar{m}}$.

A. The potential ignoring the dilaton

We suppose that the only relevant part of the D term involves a $U(1)$ with a Fayet-Illiopoulos term,

$$V_D = \frac{g^2}{2} \left(\sum_n q_n K_n \phi_n + \xi_D \right)^2 . \tag{2}$$

---

[5]In this paper "modulus" refers to the three untwisted moduli of the class of orbifold compactifications that we consider in explicit examples. In the usual chiral formalism, the dilaton and the universal axion are the real and imaginary parts of of a complex field s. In the linear multiplet formalism used here, the axion is replaced by a two-form potential $b_{\mu\nu}$ related to Ims by a duality transformation that determines the dilaton ℓ in the classical limit as $\ell = (2\mathrm{Re}s)^{-1}$; this relation is modified in the presence of both perturbative and nonperturbative quantum effects.

[6]We include as "matter" the so-called twisted moduli that are Standard Model gauge singlets, but have nonvanishing modular weights. For our purposes, their couplings are no different from those of twisted matter fields that are SM gauge nonsinglets.

In this expression, n runs only over matter fields charged under the relevant $U(1)$. Its gauge coupling is g, and q_n are the charges. As discussed later, weakly coupled string theory predicts that ξ_D will be an order of magnitude or so below the Planck scale.

To warm up, we consider only a single modulus t, corresponding to compactification on a six-torus [14]. Its Kähler potential is $K = -3\ln x$ where $x \equiv t + \bar{t} - \sum_n |\phi_n|^2$, and W is independent of t. This leads to what is termed a 'no-scale' theory [15], in which Eq. (1) becomes simply (for $|\phi_n| \ll 1$)

$$V = \frac{3}{x^2}\sum |W_n|^2 . \tag{3}$$

Instead of $\mathrm{Re}\, t$, one can regard x as a field since this choice too corresponds to approximately canonical normalization. The precise form of the kinetic terms is given for example in [16].

It looks as if x will run away to infinity, but we shall now see that this need not happen if vevs for the matter fields are generated from a Fayet-Illiopoulos D term.[7]

Suppose that there is only one W_n, say W_3, which comes from a term $\lambda\phi_1\phi_2\phi_3$ in the power series expansion of W, and occurs because the D term lifts the flat 1 and 2 directions. With matter fields $|\phi_n| \ll 1$ one has $K_{n\bar{m}} \simeq x^{-1}\delta_{nm}$ so this will generate a vev

$$|\phi_1\phi_2|^2 = cx^2\xi_D^2 , \tag{4}$$

where c is a constant of order 1. This will give

$$V \simeq \frac{3}{x^2}|W_3|^2 \simeq 3c\lambda^2\xi_D^2 . \tag{5}$$

All of the flat directions are preserved except for $n = 1, 2$ and 3, and the potential is also flat in the direction t. Slopes in these directions can be generated from nonrenormalizable terms, departures from the no-scale assumption, gaugino condensation or loop effects. When they are included the inflaton (corresponding to the direction of steepest descent in the space of the flat directions) might turn out to be any combination of the flat directions and t.

Earlier authors working with the potential Eq. (3) supposed instead that x was fixed, either by an ad hoc functional form for $K(x)$ [17–19], or by a loop correction [16]. Then all of the flat directions are preserved, and $\mathrm{Im}\, t$ and x are also flat.

For the rest of this paper, we invoke three moduli fields t_I. Also, we allow W to have a dependence on the moduli, that is determined by the modular invariance of the theory. The matter fields are divided into twisted fields ϕ_A and untwisted fields ϕ_{AI}. For the most part we suppose that the twisted fields vanish during inflation. Ignoring both them and the dilaton, the Kähler potential is

$$K = -\sum_{I=1}^{3} \ln x_I, \tag{6}$$

---

[7]The possibility of generating vevs from such a term has been considered in two previous works [8,10], but these failed to notice the crucial effect of the nontrivial kinetic term K_n in Eq. (2).

where

$$x_I = t_I + \bar{t}_I - \sum_A |\phi_{AI}|^2. \tag{7}$$

The potential Eq. (1) becomes

$$V = e^K \left[\sum_I \left(x_I \sum_A |W_{AI} + \bar{\phi}_{AI} W_I|^2 + |x_I W_I - W|^2 \right) - 3|W|^2 \right]. \tag{8}$$

In this expression, $W_I \equiv \partial W / \partial t_I$.

If W does not depend on the moduli t_I, we have simply

$$V = e^K \sum_{I,A} x_I |W_{AI}|^2 \tag{9}$$

This is the same as in global SUSY, except for the factors $e^K x_I = x_I / x_1 x_2 x_3$.

As in the previous case, a Fayet-Illiopolos D term may prevent the runaway of the x_I, even if W has no dependence on x_I.

As pointed out in [16], the preservation of flat directions for the toy models (3) and (9) is a consequence of the Heisenberg invariance [20] of the Kähler potential which depends on the scalar fields of these models only through the invariant combinations x and x_I of the moduli t or t_I and the (untwisted) matter fields ϕ_A or ϕ_{IA},

B. The full potential from orbifold compactification

To construct a realistic string-derived model, we have to include the dilaton, the Green-Schwarz term needed to cancel the modular anomaly induced by field theory loop corrections, a superpotential for the twisted sector fields, and the effective potential for the dilaton that is induced by gaugino condensation. The last two terms break Heisenberg invariance through an explicit dependence on the moduli (although the last will be considered negligible in much of our discussion). The Kähler potential and the G-S term are not completely known. In addition to imposing modular invariance we will assume that they are Heisenberg invariant. This is equivalent to the, at least plausible, assumption that the Kähler potential and the G-S term involve the untwisted scalar fields only through the radii R_I of compactification of the three tori: in string units $1/2R_I^2 = t_I + \bar{t}_I - \sum_A |\phi_{AI}|^2$. We will indicate the modifications that occur if this assumption is relaxed.

To incorporate the dilaton, we turn to the model of [9], which arguably reflects string constraints more faithfully than any so far. It was realized some time ago [21–23] that the usual chiral formalism for gaugino condensation, that uses an unconstrained chiral multiplet as in interpolating field for the gaugino condensate, is inconsistent with the Bianchi identity for the Yang-Mills chiral superfield. The most straightforward way to introduce a chiral field with the correct constraint is to identify it with the chiral projection of a vector superfield whose components also contain a those of a linear supermultiplet, interpreted as the dilaton supermultiplet in this formalism. In fact, the chiral multiplet for the dilaton that is commonly used is obtained after a duality transformation from the components of a

linear supermultiplet that are remnants of the dilaton, dilatino, and a three-form field of ten dimensional supergravity. There is increasing evidence [24] that the linear supermultiplet is the correct formulation for the dilaton in the context of superstring theory. Although it has been argued [21] that the linear and chiral multiplets are equivalent through a duality transformation even at the quantum level and in the presence of nonperturbative effects, the implementation of the correct constraint leads to considerable complication in the chiral formalism. Using the linear formalism for gaugino condensation [21,22], together with constraints from string theory, including perturbative modular invariance [25] and matching conditions [26,27] at the string scale, as well as a parameterization [28,29] of nonperturbative string effects [30] to stabilize the dilaton, it was shown [9] that the moduli are naturally stabilized at their self dual points in realistic theories [12,13] from orbifold compactification. Since the dilaton is stabilized at weak coupling, one expects the result for moduli *vev*'s to hold up to possible small corrections that could arise if the nonperturbative string corrections to the dilaton Kähler potential are moduli-dependent [31].

We take the Kähler potential K and the Green-Schwarz term V_{GS} to be

$$K = G + \ln V + g(V), \quad G = \tilde{G} + \sum_A X_A, \quad V_{GS} = b\tilde{G} + \sum p_A X_A,$$

$$\tilde{G} = \sum_I \tilde{G}_I, \quad \tilde{G}_I = -\ln(T_I + \bar{T}_I - \sum_A |\Phi_{AI}|^2), \quad X_A = \exp\left(\sum_I q_I^A \tilde{G}^I\right)|\Phi_A|^2, \tag{10}$$

where $g(V)$ parameterizes nonperturbative string effects, $b = 30/8\pi^2$, V is a vector superfield whose scalar component $V_{\theta=\bar{\theta}=0} = \ell$ is the dilaton. The T_I are the chiral multiplets containing the moduli. The Φ_{AI} are untwisted sector chiral multiplets, Φ_{AJ} having modular weight $q_I^{AJ} = \delta_I^J$, and the Φ_A are twisted sector chiral multiplets with modular weights $q_I^A > 0$ (typically less than 1).[8] The coefficients p_A are unknown, but can plausibly be either zero or equal to $b \approx .38$. If the twisted sector fields decouple from the GS term we have $p_A = 0$, while if the GS term is simply $V_{GS} = bK$, we have $p_A = b$. These unknown couplings determine the soft SUSY breaking parameters at the vacuum and also during inflation, and therefore may be relevant to the issue of maintaining flat directions during inflation.

To obtain the K that is to be used in calculating the potential, one replaces V, T_I, Φ_A and Φ_{AI} by the corresponding scalar fields ℓ, t_I, ϕ_A and ϕ_{AI}. (The matter fields will be denoted collectively by ϕ_α, with $\alpha = A$ or AI.) As before, we define $x_I \equiv t_I + \bar{t}_I - \sum_A |\phi_{AI}^2|$, and we also define

$$X_A \equiv \left(\prod_I x_I^{-q_I^A}\right)|\phi_A|^2 . \tag{11}$$

Then

$$K = \ln(\ell) + g(\ell) - \sum_I \ln x_I + \sum_A X_A , \tag{12}$$

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[8]The group of modular transformations on the moduli is generated by $t_I \to 1/t_I$ and $t_I \to t_I \pm i$. Under this group, a matter field ϕ_α transforms like $\prod_I \eta^{-2q_I^\alpha}(t_I)$, where η is the Dedekind function and q_I^α are the weights of the field.

The choices for K and V_{GS} are consistent with what is known [32] from string theory. The Lagrangian for the untwisted sector of the theory can be obtained by direct compactification of ten-dimensional supergravity. Therefore, we know the part of K that depends only on the *untwisted* fields, which will allow us to keep their masses under control and hence preserve some untwisted flat directions. To obtain information about twisted sector couplings, an expansion of the S-matrix as power series in matter fields has been used to obtain the moduli-dependence of the coefficient of Φ_A^2. In writing (10) we made an additional assumption of Heisenberg invariance, and dropped higher order terms in the twisted sector fields. The latter cannot affect masses if twisted sector fields vanish during inflation. Under these assumptions, the masses of the twisted sector fields during inflation can be determined: $m_A^2 = F(q_I^A, b_A)V$, where $|F(q_I^A, b_A)| \sim 1$. For example, if we take $W_A = W_I = W_{A1} = W_{A2} = 0$ and assume $|W_{A3}| >> |W|$ (as in the explicit model introduced below), we get twisted sector masses $m_B^2 = V[1 - (1 + p_B\ell)q_3^B]$. If Heisenberg invariance is not preserved by K and V_{GS}, there can be additional terms of the form $X_A|\phi_{B_I}|^2/(t_I + \bar{t}_I)$; these can modify the masses of the twisted fields by coefficients of order one.

From the above expressions we find

$$K_\alpha = \left(\prod_I x_I^{-q_I^\alpha}\right)\bar{\phi}_\alpha\,, \tag{13}$$

and near the origin

$$K_{\alpha\beta} = \left(\prod_I x_I^{-q_I^\alpha}\right)\delta_{\alpha\beta}\,. \tag{14}$$

Supersymmetry is supposed to be broken by gaugino condensation. The condensates have masses larger than the condensation scale $\Lambda_c = |u|^{\frac{1}{3}} \sim 10^{13}\,\mathrm{GeV}$, where u is the *vev* of the gaugino condensate. Below this scale, we can integrate them out to get the effective theory.

The potential including matter fields was not given explicitly in [9]. Assuming that the D term vanishes, it is

$$\begin{aligned} V = &\frac{1}{16\ell^2}\left(\ell g'(\ell) + 1\right)\left|u(1 + b_a\ell) - 4\ell W(\phi)e^{K/2}\right|^2 \\ &-\frac{3}{16}\left|b_a u - 4W(\phi)e^{K/2}\right|^2 + \hat{G}^{n\bar{m}}Y_n\bar{Y}_{\bar{m}}\,, \end{aligned} \tag{15}$$

where b_a is one third the β-function coefficient for the confined hidden gauge group (*e.g.*, $b_a = b$ for E_8, $b_a = n/8\pi^2$ for pure $SU(n)$ Yang-Mills, *etc.*). The subscript n takes on the values I (corresponding to t_I), and $\alpha = A$ or AI. Roughly speaking the first and last terms correspond, respectively, to the F-terms for the dilaton and for the other fields in the usual chiral formulation, while the middle term corresponds to the usual $-3e^K|W|^2$.

The factor $\hat{G}^{n\bar{m}}$ is the inverse of the matrix $\hat{G}_{n\bar{m}}$, with

$$\hat{G}_{n\bar{m}} = G_{n\bar{m}} + \ell V_{n\bar{m}}^{\mathrm{GS}} \tag{16}$$

$$= (1 + b\ell)\tilde{G}_{n\bar{m}} + \sum_A [1 + p_A\ell]\,(X_A)_{n\bar{m}} \tag{17}$$

and $\xi = \eta'/\eta$ is the logarithmic derivative of the Dedekind function η. The factors Y_n are given by

$$Y_n = e^{K/2}(W_n + G_n W) + \frac{u}{4}(b - b_a)\tilde{G}_n$$
$$+\frac{u}{4}\sum_A (p_A - b_a)(X_A)_n - \frac{u}{2}(b - b_a)\xi(t_I)\delta_{nI}\,. \tag{18}$$

Because we are dealing with a linear multiplet, the superpotential W is independent of the dilaton. This is in contrast with the case for the chiral multiplet formulation, and is an important simplification. The superpotential has a power series expansion in the matter fields which we take to be

$$W = \sum_m \lambda_m \prod_\alpha \phi_\alpha^{n_m^\alpha} \prod_I \eta(t_I)^{2(\sum_\alpha n_m^\alpha q_I^\alpha - 1)} \tag{19}$$

with the n_m^α positive integers or zero. The t_I dependence of each coefficient is dictated by modular invariance, which requires that W transforms like $\prod_I \eta^{-2(t_I)}$ (up to a modular-invariant holomorphic function, which we do not consider because it would have singularities). Using this expression one sees that

$$\frac{\partial W}{\partial t_I} \equiv W_I = 2\xi(t_I)\left(\sum_\alpha q_I^\alpha \phi_\alpha W_\alpha - W\right). \tag{20}$$

Putting all this together, the potential is

$$V = \frac{1}{16\ell^2}\left(\ell g'(\ell) + 1\right)\left|u(1 + b_a\ell) - 4\ell W e^{K/2}\right|^2 - \frac{3}{16}\left|b_a u - 4W e^{K/2}\right|^2$$
$$+\sum_A \left(\prod_I x_I^{q_I^A}\right)\frac{|Y_A|^2}{1 + p_A\ell} + \sum_I \frac{1}{1 + b\ell + \sum_B (1 + p_B\ell) q_I^B X_B} \times$$
$$\left[\left|A_I\left(2\xi(t_I)x_I + 1\right) - e^{K/2}\sum_A \phi_{AI} W_{AI}\right|^2\right.$$
$$\left. + x_I \sum_A \left|W_{AI} e^{K/2} + 2\xi(t_I) A_I \bar{\phi}_{AI}\right|^2\right]. \tag{21}$$

In this expression,

$$A_I = e^{K/2}\left(\sum_\alpha q_I^\alpha \phi_\alpha W_\alpha - W\right) - \frac{u}{4}(b - b_a), \tag{22}$$

or equivalently

$$2\xi(t_I)A_I = e^{K/2}W_I - 2\xi(t_I)\frac{u}{4}(b - b_a). \tag{23}$$

Also,

$$Y_A = \frac{1}{\phi_A}\left\{e^{K/2}\left[\phi_A W_A + X_A W\right] + \frac{u}{4}(p_A - b_a)X_A\right\} \tag{24}$$
$$= e^{K/2}\left[W_A + K_A W\right] + \frac{u}{4}(p_A - b_a)K_A \tag{25}$$

where $K_A = \left(\prod_I x_I^{-q_I^A}\right) \bar{\phi}_A$.

This potential has degenerate vacua with broken supersymmetry, at $t_I = 1$ and $t_I = e^{i\pi/6}$ (with all matter fields equal to zero in both cases).[9] Only the former was considered in [9], but the qualitative features are the same if one takes $t_I = e^{i\pi/6}$.

If $V \gg u^2$ (restoring the Planck mass, $V^{1/4} \gg 10^{11}$ GeV), then u is presumably negligible and we obtain

$$V = e^K \Bigg\{ (\ell g'(\ell) + 1)\,|W|^2 - 3|W|^2 + \sum_A \left(\prod_I x_I^{q_I^A}\right) \frac{|W_A + K_A W|^2}{1 + p_A \ell} + \sum_I \frac{1}{1 + b\ell + \sum_B (1 + p_B \ell) q_I^B X_B} \times \left[\left| x_I W_I - W + \sum_A q_I^A \phi_A W_A \right|^2 + x_I \sum_A \left| W_{AI} + W_I \bar{\phi}_{AI} \right|^2 \right] \Bigg\}. \tag{26}$$

This corresponds to Eq. (8), with the dilaton and twisted-sector fields now included.

III. BUILDING A MODEL OF INFLATION

Guided by references [8,10], we suppose that during inflation the following conditions hold.

1. The gaugino condensate u is negligible, corresponding to $V^{1/4} \gg \sqrt{u} \sim 10^{11}$ GeV.
2. Every term in the expansion (19) of W vanishes (i.e., at least one of the fields in each term vanishes).
3. All derivatives of W with respect to the fields vanish, except for W_{C3} which is fixed during inflation corresponding to a single untwisted matter field $\alpha = C3$.[10]
4. All matter field values are $\ll 1$.

Somewhat less specific conditions would have essentially the same effect, but these have the virtue of simplicity, and we shall later be making a specific proposal for achieving them.

Since the twisted fields ϕ_A are $\ll 1$, the terms X_A defined by Eq. (11) are also $\ll 1$. Both ϕ_A and X_A can be ignored in Eq. (26), which becomes simply

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[9]These vacua correspond to $2\xi(t_I)x_I + 1 = 0$. Points in field space that are obtained from these vacua by a modular transformation do not represent physically distinct vacua, since the group of modular transformations is a gauge discrete symmetry as opposed to a global one.

[10]The choice $I = 3$ is arbitrary, and one could allow nonvanishing W_α for more untwisted fields from the *same* sector without changing anything. As we shall see later, one might also allow nonvanishing W_α for fields from two or all three of the untwisted sectors, but for the moment we insist on just one.

$$V = \frac{\ell e^{g(\ell)}}{(1+b\ell)x_1 x_2}|W_{C3}|^2 . \tag{27}$$

A. A simple possibility

In order to proceed, we need to know the dependence of $|W_{C3}|$ on the moduli and the dilaton. Following reference [8], let us first suppose that W_{C3} (considered as a function of the matter fields and the t_I) is independent of t_3, but has dependence on t_1 and t_2 which stabilizes the potential. Then flat directions are preserved for matter fields in the $I = 3$ sector (except for any which are spoiled by coupling to fields that are displaced from the origin) and the t_3 direction is also flat. This is because the terms in Eq. (26) that give zero contribution to V, also give zero contribution to the squared masses of these fields. Flat directions in the $I = 1$ and $I = 2$ sectors are not preserved because of the factors x_1 and x_2 in Eq. (27). As noted earlier, the masses of twisted sector fields obtained from Eq. (26) could be modified by unknown coefficients of order one, if the assumption of Heisenberg invariance of the Kähler potential is dropped. The inflaton trajectory could be any combination of t_3 and the flat $I = 3$ directions (excluding ϕ_{C3} which is supposed to be fixed).

These conditions would be achieved [8] if W_{C3} came from a term $\Lambda^2\phi_{C3}$, with Λ independent of the matter fields. Then, modular invariance would imply that $\Lambda^2 \propto \eta^{-2}(t_1)\eta^{-2}(t_2)$, and

$$V \propto \left[|\eta(t_1)\eta(t_2)|^4 x_1 x_2\right]^{-1} . \tag{28}$$

To discuss the stability of the moduli, we can set the matter fields equal to zero so that $x_I = t_I + \bar{t}_I$. As shown in [8], V is stabilized at $t_1 = t_2 = e^{i\pi/6}$ up to modular transformations. The squared masses of the canonically normalized t_1 and t_2 turn out to be precisely V, which presumably hold them in place during inflation.

The value $t_I = e^{i\pi/6}$ corresponds to a fixed point[11] of the modular transformations. Since it must be an extremum of the potential, it is not particularly surprising to find that it represents the minimum during inflation. As we noted earlier it also represents a possible true vacuum. As a result, the moduli stabilized at this point during inflation will remain there, and, as has often been noted before, would not be produced in the early Universe.

To complete this simple model, note that $|W_{C3}|$ has no dependence on the dilaton. The Eq. (27) gives

$$V = \lambda\frac{\ell e^g}{1+b\ell}, \quad V' = \frac{V}{\ell}\left(1+\ell g' - \frac{b\ell}{1+b\ell}\right),$$
$$V'' = \frac{V}{\ell^2}\left(\ell^2 g'' - 1 + \frac{b^2\ell^2}{(1+b\ell)^2}\right) + \frac{V'}{\ell}\left(1+\ell g' - \frac{b\ell}{1+b\ell}\right). \tag{29}$$

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[11]The other fixed point in the fundamental domain, namely $t_I = 1$, is a saddle point of potential (28); see e.g. [33].

We require $V' = 0,\ V'' > 0$ for stabilization, which means

$$\ell g' + 1 = \frac{b\ell}{1+b\ell}, \quad \ell^2 g'' > 1 - \frac{b^2\ell^2}{(1+b\ell)^2}. \tag{30}$$

The function $f(\ell)$ is related to $g(\ell)$ by a differential equation (which assures a canonical form for the Einstein term in the Lagrangian):

$$g'\ell = f - f'\ell. \tag{31}$$

An example (taken here for calculational simplicity) of a choice for f that reflects string nonperturbative effects [30], and stabilizes the dilaton at weak coupling $[\alpha(m_{str}) \approx .17]$ and vanishing cosmological constant, is

$$f = 4.25 e^{-1/\sqrt{b\ell}}(1 - .53/\sqrt{b\ell})\,. \tag{32}$$

With this parameterization the potential (27) has a local minimum at $\ell = 4.2$, which is in the domain of attraction and roughly satisfies our initial assumption that $\ell = O(1)$ during inflation.

B. More general possibilities

Contrary to what was stated in [8], one cannot argue that Eq. (28) always holds, because in general W is an arbitrary expression of the form Eq. (19). Although W_{A3} transforms like $(\eta(t_1)\eta(t_2))^{-2}$, this is automatically satisfied by Eq. (19) and it does not in general determine the dependence of W_{A3} on the t_I at (say) fixed values of the matter fields. More specifically, since we are working in the context of string theory, we need to justify the emergence of a superpotential term linear in a matter field, since the effective Lagrangian from string theory contains terms of cubic and higher order. Thus a linear term can arise only from some fields acquiring *vev*'s.

In fact, something like Eq. (28) may be applicable under rather general circumstances. Let us a assume that at the string scale, W includes a term of the form

$$W = \lambda\phi_{C3}\,[\eta(t_1)\eta(t_2)]^{-2} \prod_\alpha \phi_\alpha \prod_I [\eta(t_I)]^{2q_I^\alpha}\,, \tag{33}$$

where the product over α does not contain ϕ_{C3}. We suppose that during inflation, there are nonzero vevs $|\phi_\alpha|^2$, with the modular invariant form

$$|\phi_\alpha|^2 \prod_I x_I^{-q_I^\alpha} = c_\alpha \ell^{d_\alpha} \prod_I \left[x_I|\eta(t_I)^4|\right]^{p_I^\alpha}\,, \tag{34}$$

where c_α is a constant. As we shall see, *vev*'s of this form can indeed be generated from a Fayet-Illiopoulos D term. Such vevs will drive ϕ_{C3} to zero as required. Once the fields with these *vev*'s have been integrated out to give an effective theory relevant to the scale of inflation, the moduli-dependence of the potential (27) takes the form

$$V \propto \prod_I \left[|\eta(t_I)|^4 x_I\right]^{n_I}, \quad n_{I=1,2} = \sum_\alpha (p_I^\alpha + q_I^\alpha) - 1, \quad n_3 = \sum_\alpha (p_3^\alpha + q_3^\alpha) . \tag{35}$$

Bearing in mind our earlier discussion, we want one or more of the n_I to vanish, providing flat directions suitable for inflation. Any remaining n_I should be negative, which as we noted after Eq. (28) will ensure that the corresponding moduli are stabilized at $t_I = e^{i\pi/6}$. (Positive n_I are excluded, because the potential would be driven to zero in the direction $t_I \to \infty$ (or 0).) Inflaton candidates are the modulus (or moduli) with $n_I = 0$, and matter fields in the corresponding untwisted sector(s) which correspond to flat directions.

In contrast with the earlier situation, any or all of the n_I can vanish. If they all vanish, one could generalize Eq. (33) to be a sum over terms, with ϕ_{C3} is replaced by fields from *different* untwisted sector.

The dilaton-dependence of V is

$$V \propto \frac{e^{g(\ell)}\ell^d}{1+b\ell}, \quad d = 1 + \sum_\alpha d_\alpha, \tag{36}$$

for which the minimization conditions (32) become

$$\ell g' + d = \frac{b\ell}{1+b\ell}, \quad \ell^2 g'' > d - \frac{b^2\ell^2}{(1+b\ell)^2}. \tag{37}$$

The condition that the potential be positive definite requires [9,28] $\ell g' > -1$ and since $0 \leq b\ell/(1+b\ell) < 1$, stabilization can occur only if $d < 2$. With the parameterization introduced above, there is a minimum within the domain of attraction (*i.e.*, with $\ell \gtrsim 1.4$) for $-3.3 \lesssim d \lesssim 1.4$.

Taken literally, this model gives an exactly flat inflaton potential, and no mechanism for ending inflation. There are many possibilities for generating a slope. It can come from small departures from assumptions 2–4, from the gaugino condensate or loop corrections. Also, if the *vev*'s are generated by a D term that term will be driven to a small but nonzero value; this will generate a slope from the inflaton-dependence of the factors $K_{n\bar{n}}$ in Eq. (2). With the slope in place, a generalization of the model exhibited in reference [10] allows inflation to end by the hybrid inflation mechanism.

Provided that no matter fields charged under the strongly coupled hidden gauge group acquire large *vev*'s during inflation, the condensate potential V_c will indeed be present. Let us estimate the mass it generates for the moduli with $n_I = 0$. As mentioned above, the model is viable only if the moduli are close to their vacuum values – *i.e.* within the domain of attraction – during inflation. In this case the mass of moduli with $n_I = 0$ is $m_{t_I}^2 \approx V_c(\ell_i)$, where ℓ_i is the value of the dilaton field during inflation. The magnitude of V_c is governed by the value of the string-scale gauge coupling $g^2(\ell_i) = 2\ell_i/[1+f(\ell_i)]$, and the condition that the vacuum energy vanishes in the true vacuum assures that this is a slowly varying function near its vacuum value ℓ_v. For the parameterization used above with $d = 1$, $(g^{-2}(\ell_v), g^{-2}(\ell_i)) = (.44, .13)$, $u(\ell_i) \approx 10^2 u(\ell_v)$, $m_{t_I}(\ell_i) \approx 5m_t(\ell_v) \approx 100\,\text{TeV} \ll V^{\frac{1}{2}}$. (Note that $V_c(\ell_i)$ is of order $(10^{11}\,\text{GeV})^4$ as one would expect.)

Flat directions in the corresponding untwisted sector are lifted by mass terms of order $|m_{\phi_{AI}}|^2 \sim m_{t_I}^2$ as long as $|\phi_{IA}|^2 \lesssim \text{Re}t_I$. This contribution is negative if $|\phi_{IA}|^2 \lesssim .2\text{Re}t_I$, and

is much smaller than that induced [16] by loop effects ($-m_\phi^2 \sim 10^{-2}V$). If either of these gives the dominant slope, the spectral index $n = 1 + 2m^2/V$ is very close to 1.

Finally, we note that the "moduli problem" [34,35] encountered in generic supergravity/superstring inflationary scenarios, is considerably alleviated in the model studied here. While the dilaton is stabilized at a value shifted from its vacuum value by an amount of $O(1)$, its mass is about $10^6 GeV$ [9], and its decay does not contribute to the moduli problem.[12] The moduli masses are about 20 TeV, which is sufficient to evade the moduli problem of [34] only if R-parity is violated [35] (or the moduli abundance is diluted by thermal inflation [37]). If R-parity is conserved, the problem is still evaded for those moduli that are stabilized at the vacuum value $t_I = e^{i\pi/6}$. It is possible that the requirement that the remaing moduli (e.g., t_3 in the above example) be in the domain of attraction is sufficient to avoid the problem altogether.

C. Generating *vev*'s with the D term

In many models *vev*'s of the form (34) with $p_I^\alpha = 0$, $d_\alpha = 1$ arise from a Fayet-Illiopoulos D term, whose contribution to the potential is given by Eq. (2). In string models it arises as a GS counter term, introduced [38] to cancel a $U(1)$ gauge anomaly of the effective field theory (with no corresponding string theory anomaly), analogous to the the GS term introduced in III.B to cancel the modular anomaly. Orbifold models with an anomalous $U(1)$ and supersymmetric vacua have been found in [12,13], and the GS D term has been used in various applications to phenomenology.

In the linear multiplet formalism, the gauge coupling constant g (defined at the string scale) which appears in Eq. (2) is given by

$$g^2 = \frac{2\ell}{f(\ell)+1}\,. \tag{38}$$

The scale ξ_D of the Fayet-Illiopoulos term is given in this formalism by

$$\xi_D = \frac{2\ell \mathrm{Tr}(T)}{192\pi^2}\,, \tag{39}$$

where T is the generator of the anomalous $U(1)$, whose trace $\mathrm{Tr}(T) = \sum q_n$ is perhaps [39,13] of order 100.

Using Eq. (13), one sees that *vev*'s generated by the D term will be of the form Eq. (34).[13] If this were the only source of *vev*'s, d would just be the dimension of the the superpotential

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[12]Even though there is no constraint from the Big-Bang Nucleosynthesis because of its high mass and hence early decay, it still dilutes the baryon asymmetry by a factor of roughly 10^{-12}. A very efficient mechanism, such as Affleck–Dine mechanism, can generate enough baryon asymmetry to withstand the dilution [36].

[13] In [10], $K_{n\bar{m}}$ was set equal to δ_{nm} in the D-term. Including the nontrivial K will in general affect the flatness, as the present discussion demonstrates.

in (33) with $d \geq 3$. The potential (36) would be driven to zero in the direction of vanishing gauge coupling $\ell \to 0$. However *vev*'s induced by a D term can induce other *vev*'s with a different ℓ-dependence through superpotential terms. If, for example, there is a gauge invariant, modular covariant superpotential term

$$w = F \prod_{\beta=1}^{3} \phi_\beta \prod_I [\eta(t_I)]^{2q_I^\beta}, \quad F = \prod_I \eta^{-2}(t_I), \tag{40}$$

the superpotential

$$W(w) = F \sum_n c_n (w/F)^n, \tag{41}$$

is allowed by all the symmetries. Now suppose the D term induces modular invariant *vev*'s for ϕ_2 and ϕ_3:[14]

$$| < \phi_\beta > |^2 = |v_\beta|^2 = c_\beta \ell \prod_I (t_I + \bar{t}_I)^{q_I^\beta}, \quad \beta = 2, 3. \tag{42}$$

Then if c_1 is nonzero, $W_1(w)$ does not vanish unless $< \phi_1 >= v_1 \neq 0$. Solving $W_\beta = 0$, $\beta = 1, 2, 3$ gives a single equation: $\phi^\beta W_\beta = \sum_n n c_n (w/F)^n = 0$, which is solved by $w/F = \sqrt{c_1} =$ constant. Then

$$|v_1|^2 = \frac{c_1}{c_2 c_3} \ell^{-2} \prod_I x_I^{q_I^1} \left[x_I |\eta(t_I)|^4 \right]^{p_I^1}, \quad p_I^1 = \sum_{\beta=1}^{3} q_I^\beta. \tag{43}$$

It is easy to satisfy the condition $d < 2$ by including such fields in the superpotential (33). For example, if $w = \phi_{A1}\phi_{B2}\phi_{C'3}$ in (40) and $W = \phi_{C3}\phi_{A1}\phi_{B'2}\phi_{C''3}\eta(t_3)^2$ in (33), with D term induced *vev*'s for $\phi_{B2}, \phi_{C'3}, \phi_{B'2}, \phi_{C''3}$, one recovers precisely the behavior in (27) and (28).

The magnitude of the potential will be of the form

$$V = \lambda \Lambda^{-2n} \xi_D^{2(2+n)}. \tag{44}$$

In this expression, λ is a ratio of dimensionless couplings in the superpotential (times a coefficient of order 1), $3 + n$ is the dimension of the term Eq. (33) of the superpotential, and Λ is scale of nonrenormalizable terms in the superpotential. Using Eq. (39), and the perturbative superstring estimate $\Lambda^2 = M_{\rm str}^2 \equiv g_{\rm str}^2 M_{\rm Pl}^2$, this gives

$$V^{1/4} \sim \lambda^{1/4} g_{\rm str} \left(\frac{{\rm Tr}(T)}{192\pi^2} \right)^{(2+n)/4} M_{\rm Pl}. \tag{45}$$

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[14]The ratio of ϕ_2 and ϕ_3 *vev*'s can be fixed, for instance, by gauging a non-anomalous U(1) symmetry under which they have the opposite and equal charges.

One expects that $g_{\rm str}$ will be at most an order of magnitude below unity, [15] but the other factors can be smaller. With reasonable values like $n = 1$ or $n = 2$, one can easily achieve the result $V^{1/4} \lesssim 10^{-2} M_{\rm Pl}$ required by the COBE normalization.

IV. CONCLUSION

We have shown how to construct a general class of inflation models, with some very desirable properties. The scale of inflation is set by a Fayet-Illiopoulos term, derived from the superstring. Some of the flat directions of global SUSY are preserved, and the potential is also flat in the direction of at least one of the moduli. All of these flat directions are candidates for the inflaton field. The dilaton is stabilized within its domain of attraction, and the remaining moduli are stabilized at or near their vacuum values. The models may avoid the usual moduli problem.

These inflation models are constructed within the framework of a specific, modular invariant, model of supersymmetry breaking in the true vacuum, that invokes string nonperturbative effects to stabilize the dilaton. It is based on a class of orbifold compactifications with just three untwisted moduli t_I, and the dilaton is described by the linear supermultiplet formalism.

Although the specific model contains a definite mechanism for stabilizing the dilaton in the true vacuum, this is actually irrelevant for our proposed models of inflation because they make the inflationary energy scale much bigger than the scale of SUSY breaking in the true vacuum ($V^{1/4} \gg 10^{11}$ GeV). If, for instance, SUSY breaking in the true vacuum is gauge-mediated, our models of inflation still work provided that the dilaton is described by the assumed linear-multiplet formalism.

Let us emphasize that what we have given is only a strategy for model-building. We have shown how to achieve a sufficiently flat inflationary potential, but we have not exhibited a complete model. Such a model would define the slope of the inflationary potential, and would include a mechanism for ending inflation. To achieve this last objective we would presumably need a hybrid inflation model, along the lines of the one given in reference [10]. To exhibit such a model would seem to be a worthwhile project, and might be quite nontrivial.

Acknowledgements

MKG and HM were supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098 and in part by the National Science

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[15] $g^2 \simeq 0.5$ would correspond to the value $\alpha_{\rm str} \equiv g_{\rm str}^2/4\pi \simeq 1/25$, which with naive running of the couplings is suggested by observation at a scale of order 10^{16} GeV. In fits (including additional particles) to the string theory constraint that unification occurs at $\mu^2 = (2e)^{-1} g^2 M_{\rm Pl}^2$, the value of g^2 is somewhat larger.

Foundation under grant PHY-90-21139. HM was also supported by the Alfred P. Sloan Foundation. DHL is supported by PPARC grant GR/L40649 and NATO grant CRG 970214, and acknowledges hospitality at LBNL where this work was started. DHL acknowledges a long-term collaboration with Ewan Stewart, and useful correspondence with Chris Kolda and Toni Riotto. MKG thanks Yi-Yen Wu for discussions.

REFERENCES

[1] A. R. Liddle and D. H. Lyth, Phys. Rep. **231**, 1 (1993).
[2] D. H. Lyth, hep-ph/9609431 D. H. Lyth and A. Riotto, to be submitted to Phys. Rep..
[3] I. Affleck and M. Dine, Nucl. Phys. **B249**, 361 (1985).
[4] B. A. Campbell, M. K. Gaillard, H. Murayama and K. A. Olive, hep-ph/9805300.
[5] M. Dine, L. Randall, and S. Thomas, Nucl. Phys. **B458**, 291 (1996).
[6] E. D. Stewart, M. Kawasaki and T. Yanagida, Phys. Rev. **D54**, 6032 (1996).
[7] G. D. Coughlan, R. Holman, P. Ramond and G. G. Ross, Phys. Lett. **140B**, 44 (1984); M. Dine, W. Fischler and D. Nemeschansky, Phys. Lett. B **136**, 169 (1984).
[8] E. J. Copeland, A. R. Liddle, D. H. Lyth, E. D. Stewart and D. Wands, Phys. Rev. **D49**, 6410 (1994).
[9] P. Binétruy, M. K. Gaillard and Y. Wu, Nucl. Phys. **B493** (1997) 27 and Phys. Lett. **412** (1997) 228.
[10] E. D. Stewart, Phys. Rev. D **51**, 6847 (1995).
[11] A. D. Linde, Phys. Lett. **B259**, 38 (1991).
[12] L.E. Ibañez, H.-P. Nilles and F. Quevedo, Phys. Lett. **B187**, 25 (1987).
[13] A. Font, L. Ibanñez, F. Quevedo and A. Sierra, Nucl. Phys. **B331**, 421 (1990).
[14] E. Witten, Phys. Lett. **B155**, 151 (1985).
[15] A.B. Lahanas and D.V. Nanopoulos, Phys. Rept.. **145** , 1 (1987).
[16] M. K. Gaillard, H. Murayama and K. A. Olive, Phys. Lett. **B355**, 71 (1995).
[17] A. D. Linde, *Particle Physics and Inflationary Cosmology*, Harwood Academic, Switzerland (1990).
[18] K. A. Olive, Phys. Rep. **190**, 307 (1990).
[19] H. Murayama et al., Phys. Rev. D **50**, 2356 (1994).
[20] See [16] and references therein.
[21] C.P. Burgess, J.-P. Derendinger, F. Quevedo and M. Quirós, Phys. Lett. **B348** (1995) 428.
[22] P. Binétruy, M.K. Gaillard and T.R. Taylor, Nucl. Phys. **B455** (1995) 97.
[23] P. Binétruy and M.K. Gaillard, Phys. Lett. **B365** (1996) 87.
[24] W. Siegel, Phys. Lett. **211B** (1988) 55; J. de Boer and K. Skenderis, Nucl. Phys. **B 481** (1996) 129; S.J. Gates, private communication.
[25] A. Giveon, N. Malkin and E. Rabinovici, Phys. Lett. **B220** (1989) 551; E. Alvarez and M. Osorio, Phys. Rev. **D40** (1989) 1150.
[26] M.K. Gaillard and T.R. Taylor, Nucl. Phys. **B381** (1992) 577.
[27] V.S. Kaplunovsky and J. Louis, Nucl. Phys. **B444** (1995) 191.
[28] P. Binétruy, M. K. Gaillard and Y. Wu, Nucl. Phys. **B481** 109 (1996).
[29] J.A. Casas, Phys. Lett. **B384**, 103 (1996).
[30] S.H. Shenker, in *Random Surfaces and Quantum Gravity*, Proceedings of the NATO Advanced Study Institute, Cargese, France, 1990, edited by O. Alvarez, E. Marinari, and P. Windey, NATO ASI Series B: Physics Vol.262 (Plenum, New York, 1990).
[31] E. Silverstein, preprint hep-th/9611195.
[32] L.J. Dixon, V.S. Kaplunovsky and J. Louis, *Nucl. Phys.* **B329**, 27 (1990); S. Ferrara, D. Lüst, and S. Theisen, *Phys. Lett.* **233B**, 147 (1989); L.J. Dixon, V.S. Kaplunovsky and

J. Louis, *Nucl. Phys.* **B355:** 649 (1991); I. Antoniadis, K.S. Narain and T.R. Taylor, *Phys. Lett.* **267B:** 37 (1991).
[33] A. Font, L. Ibanñez, D. Lust and F. Quevedo, Phys. Lett. **B245**, 401 (1990).
[34] G.D. Coughlan, W. Fischler, E.W. Kolb, S. Raby and G..G. Ross, Phys. Lett. **B131** (1983) 59.
[35] M. Kawasaki, T. Moroi, and T. Yanagida, Phys. Lett. **B370**, 52 (1996).
[36] T. Moroi, M. Yamaguchi and T. Yanagida, Phys. Lett. **B342**, 105 (1995).
[37] D. H. Lyth and E. D. Stewart, Phys. Rev. **D53**, 1784 (1996).
[38] M. Dine, N. Seiberg and E. Witten, Nucl. Phys. B289, 589 (1987); J. Attick, L. Dixon and A. Sen, Nucl. Phys. B292, 109 (1987); M. Dine, I. Ichinose and N. Seiberg, Nucl. Phys. B293, 253 (1988).
[39] T. Kobayashi and H. Nakano, Nucl. Phys. **B496**, 103 (1997).

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Prepared for the U.S. Department of Energy under Contract No. DE-AC03-76SF00098