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Inflation, supergravity, and superstrings

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The positive potential energy required for inflation spontaneously breaks supersymmetry and in general gives any would-be inflaton an effective mass of the order of the inflationary Hubble parameter thus ruling it out as an inflaton. In this paper I give simple conditions on the superpotential that eliminate some potential sources for this mass, and derive a form for the Kähler potential that eliminates the rest. This reduces the problem of constructing a model of inflation in supergravity to that of constructing one in global supersymmetry with the extra conditions $W = W_\varphi = \psi = 0$ during inflation (where W is the superpotential, the inflaton $\in \varphi$, and $W_\psi \neq 0$). I then point out that Kähler potentials of the required form often occur in superstrings and that the target space duality symmetries of superstrings often contain R parities which would make $W = W_\varphi = 0$ automatic for $\psi = 0$.

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I. INTRODUCTION

The approximate isotropy of the cosmic microwave background radiation implies that the inflation [1,2] that inflated the observable Universe beyond the Hubble radius must have occurred at an energy scale $V^{1/4} \leq 4\times 10^{16}$ GeV [3], and it is thought that physics at energies below the Planck scale is described by an effective $N = 1$ supergravity theory [4]. Thus models of inflation should be constructed in the context of supergravity. However, this immediately leads to a problem. The positive potential energy $V > 0$ required for inflation spontaneously breaks supersymmetry,[1] which would generally be expected to give effective masses $\sim \sqrt{8\pi V}/m_{\rm Pl} \sim H$ to any would-be inflatons. But inflation requires $|V''/V| \ll 1$, i.e., the effective mass of the inflaton must be much less than the inflationary Hubble parameter H.

Natural inflation [5] avoids this problem by assuming the inflaton corresponds to an angular degree of freedom whose potential is kept flat enough by an approximate compact global symmetry. The model of Holman *et al.* [6,2] assumes the form of supergravity that gives minimal kinetic terms[2] and fine-tunes a parameter in the superpotential to eliminate the troublesome mass term. Solutions to this problem which work for ϕ^n chaotic inflation have also been proposed [7,1,8], but they rely on forms for the supergravity Kähler potential that have no independent motivation. In this paper I will propose a solution for inflaton fields which are not purely angular degrees of freedom, which requires no fine tuning, and which uses a well-motivated form for the Kähler potential. Some aspects of this solution have been investigated in [9].

[1] After inflation V disappears and so supersymmetry is restored modulo whatever breaks supersymmetry in our vacuum.

[2] This is no longer regarded as realistic.

II. BASIC FORMULAS AND NOTATION

I will use the following conventions in this paper: $m_{\rm Pl}/\sqrt{8\pi} = 1$, a prime will denote the derivative with respect to the canonically normalized inflaton field σ, a bar will denote the Hermitian conjugate, ϕ will represent a vector whose components ϕ^α are complex scalar fields, and subscript ϕ will denote the derivative with respect to ϕ, so, for example, W_ϕ represents the vector with components $\partial W/\partial\phi^\alpha$.

A. Global supersymmetry

In global supersymmetry [4] the scalar kinetic terms are

$$|\partial_\mu\phi|^2 \,, \tag{1}$$

where $\phi = (\phi^1, \phi^2, \ldots)$ and the ϕ^α are complex scalar fields. The scalar potential is

$$V = |W_\phi|^2 + \frac{1}{2}\sum_a g_a^2 D_a^2 \,, \tag{2}$$

with

$$D_a = \bar\phi Q_a \phi + \xi_a \,, \tag{3}$$

where the superpotential $W(\phi)$ is an analytic function of ϕ, a labels the gauge group generators Q_a, g_a is the gauge coupling constant, and the real constant ξ_a is a Fayet-Iliopoulos term that can be nonzero only if Q_a generates a U(1) gauge group. The first term is called the F term and the second the D term. I will assume that the F term gives rise to the inflationary potential energy density and that the D term is flat along the inflationary trajectory so that it can be ignored during inflation. It may however play a vital role in determining the trajectory and in stabilizing the noninflaton fields.

B. Supergravity

The scalar fields in a supergravity theory are the coordinates of a Kähler manifold. The metric on a Kähler manifold is $K_{\bar{\phi}\phi}$ where the Kähler potential $K(\phi, \bar{\phi})$ is a real function of ϕ and its Hermitian conjugate $\bar{\phi}$. The scalar kinetic terms are

$$\partial_\mu \bar{\phi} \, K_{\bar{\phi}\phi} \, \partial^\mu \phi \,, \tag{4}$$

the F-term part of the scalar potential is

$$V_F = e^K \left[\left(W_\phi + W K_\phi \right) K^{-1}_{\bar{\phi}\phi} \left(\bar{W}_{\bar{\phi}} + \bar{W} K_{\bar{\phi}} \right) - 3|W|^2 \right] , \tag{5}$$

and the D-term part is

$$V_D = \frac{1}{2} \sum_{a,b} \left(\mathrm{Re}\, f_{ab} \right)^{-1} D_a D_b \,, \tag{6}$$

with

$$D_a = K_\phi Q_a \phi + \xi_a \,, \tag{7}$$

where $f_{ab}(\phi)$ is an analytic function of ϕ transforming as a symmetric product of two adjoint representations of the gauge group. Only the combination

$$G(\phi, \bar{\phi}) = K + \ln |W|^2 \tag{8}$$

is physically relevant and we are always free to make a Kähler transformation:

$$\begin{aligned} K(\phi, \bar{\phi}) &\to K(\phi, \bar{\phi}) - F(\phi) - \bar{F}(\bar{\phi}), \\ W(\phi) &\to e^{F(\phi)} W(\phi) \,. \end{aligned} \tag{9}$$

C. Inflation

I assume the effective action during inflation [1,2] to have the form

$$S = \int d^4x \sqrt{-g} \left[-\frac{1}{2} R + g^{\mu\nu} \partial_\mu \bar{\phi} \, K_{\bar{\phi}\phi} \, \partial_\nu \phi - V(\phi, \bar{\phi}) \right] , \tag{10}$$

and make the usual flat Robertson-Walker ansatz

$$ds^2 = dt^2 - a(t)^2 d\mathbf{x}^2 , \;\; \phi = \phi(t) \,. \tag{11}$$

The Hubble parameter H is defined as $H \equiv \dot{a}/a$. Inflation requires[3] $-\dot{H}/H^2 \ll 1$, or equivalently $3H^2 \simeq V$, i.e., the energy density of the Universe should be dominated by the scalar potential. The dynamics of the scalar fields then rapidly approaches the slow-roll equations of motion

$$\dot{\phi} = -3 H K^{-1}_{\bar{\phi}\phi} V_{\bar{\phi}} \,, \tag{12}$$

and I will assume that they have been attained for all epochs of interest. The canonically normalized inflaton σ is defined by

$$\frac{1}{2} d\sigma^2 = d\bar{\phi} \, K_{\bar{\phi}\phi} \, d\phi \,. \tag{13}$$

The conditions necessary for inflation can be expressed in terms of the potential as

$$\left(\frac{V'}{V} \right)^2 \ll 1, \;\; \left| \frac{V''}{V} \right| \ll 1 \,, \tag{14}$$

where a prime denotes the derivative with respect to σ.

III. THE PROBLEM

At any point in the space of scalar fields we can make a holomorphic field redefinition such that $\phi = 0$ and the scalar fields have canonical kinetic terms at that point. Any purely holomorphic terms in the Kähler potential can then be absorbed into the superpotential using a Kähler transformation. Then, in the neighborhood of that point, the Kähler potential will be

$$K = |\phi|^2 + \cdots \,, \tag{15}$$

where the ellipsis stands for higher-order terms. Therefore, from Eq. (5),

$$\begin{aligned} V = {} & \exp\left(|\phi|^2 + \cdots \right) \\ & \times \left\{ \left[W_\phi + W \left(\bar{\phi} + \cdots \right) \right] (1 + \cdots) \left[\bar{W}_{\bar{\phi}} + \bar{W} \left(\phi + \cdots \right) \right] \right. \\ & \left. -3|W|^2 \right\} \end{aligned} \tag{16}$$

$$= V|_{\phi=0} + V|_{\phi=0} \, |\phi|^2 + \text{other terms} \,. \tag{17}$$

Thus at $\phi = 0$ the exponential term gives a contribution V to the effective mass squared of *all* scalar fields. Therefore,

$$\frac{V''}{V} = 1 + \text{other terms} \,, \tag{18}$$

where the prime denotes the derivative with respect to the canonically normalized inflaton field. But $|V''/V| \ll 1$ is necessary for inflation to work. So a successful model of inflation must arrange for a cancellation between the exponential term and the terms inside the curly brackets. This will require fine-tuning unless a symmetry is used to enforce it. Natural inflation [5] uses an approximate compact global symmetry. I will use a combination of a discrete R symmetry and a noncompact global symmetry.

[3] Strictly speaking $-\dot{H}/H^2 < 1$ is all that is required. However realistic models satisfy $-\dot{H}/H^2 \ll 1$. See [10] for a more detailed discussion.

IV. A SOLUTION

Divide the vector of scalar fields ϕ into two separate vectors φ and ψ with the inflaton contained in φ:

$$\phi = (\varphi, \psi), \quad \text{inflaton} \in \varphi\,. \tag{19}$$

Then, during inflation, ψ is constant and so without loss of generality we can set

$$\psi = 0\,. \tag{20}$$

When we want to distinguish noninflaton φ fields from the inflaton we will denote them by χ:

$$\chi \subset \varphi, \quad \text{inflaton} \notin \chi\,. \tag{21}$$

Note that χ is constant during inflation.

Now the key point of the solution is to assume that

$$W = W_\varphi = 0 \;\text{ and } W_\psi \neq 0 \tag{22}$$

during inflation. Then the scalar potential Eq. (5) simplifies to

$$V = e^K W_\psi K^{-1}_{\bar\psi\psi} \bar W_{\bar\psi}\,. \tag{23}$$

Now it becomes possible to choose a form for the Kähler potential that cancels the inflaton dependent corrections to the global supersymmetry potential in a natural way.

For simplicity, assume that

$$K_{\bar\varphi\psi}|_{\psi=0} = 0\,, \tag{24}$$

so that φ and ψ have no mixed kinetic terms during inflation. Then, using a Kähler transformation to remove any remaining terms linear in ψ, and expanding about $\psi = 0$, we get

$$K = A(\varphi,\bar\varphi) + \bar\psi\, B(\varphi,\bar\varphi)\,\psi + O\left(\psi^2, \bar\psi^2\right)\,, \tag{25}$$

where A is a real function and B is a positive definite Hermitian matrix.

Note that Eqs. (22), (24), and (25) become automatic if we impose the symmetry (an R parity)

$$\psi \to -\psi, \quad \varphi \to \varphi, \quad W \to -W, \quad K \to K\,, \tag{26}$$

which also helps to stablize ψ at 0 because $V_\psi = 0$ is also automatic.[4]

From Eqs. (23) and (25),

$$V = e^A W_\psi B^{-1} \bar W_{\bar\psi}\,, \tag{27}$$

and so to eliminate the inflaton dependent corrections to the global supersymmetry potential we require

$$B^{-1} = f(\varphi,\bar\varphi)\, C^{-1}(\chi,\bar\chi)\,, \tag{28}$$

and

$$A = -\ln f(\varphi,\bar\varphi) + g(\chi,\bar\chi)\,, \tag{29}$$

where f and g are real functions, and C is a positive definite Hermitian matrix. This gives the inflationary potential

$$V = e^{g(\chi,\bar\chi)} W_\psi\, C^{-1}(\chi,\bar\chi)\, \bar W_{\bar\psi}\,, \tag{30}$$

and the Kähler potential is required to have the general form

$$K = -\ln f(\varphi,\bar\varphi) + \frac{\bar\psi\, C(\chi,\bar\chi)\,\psi}{f(\varphi,\bar\varphi)} + g(\chi,\bar\chi) + O\left(\psi^2,\bar\psi^2\right) \tag{31}$$

$$= -\ln\left[f(\varphi,\bar\varphi) - \bar\psi\, C(\chi,\bar\chi)\,\psi\right] + g(\chi,\bar\chi) + O\left(\psi^2,\bar\psi^2\right)\,. \tag{32}$$

V. SIMPLE EXAMPLES OF SUITABLE KÄHLER POTENTIALS

A. SU$(m,1)$/[SU(m) × U(1)]

The simplest example of Eq. (31) is

$$K = -\ln X, \quad X = 1 - |\phi|^2\,. \tag{33}$$

The corresponding Kähler manifold is SU$(m,1)$/[SU(m) ×U(1)], where m is the number of components of ϕ. It is a maximally symmetric space with constant Riemannian curvature. Such coset spaces form the basis of no-scale supergravity [11], though it is important to note that the Kähler potential in Eq. (33) only corresponds to part of one sector of a no-scale model. Now

$$K_{\bar\phi\phi} = \frac{1}{X^2}\left(X + \phi\bar\phi\right)\,, \quad K^{-1}_{\bar\phi\phi} = X\left(1 - \phi\bar\phi\right)\,, \tag{34}$$

and, from Eq. (5),

$$V = |W_\phi|^2 - |W_\phi\phi - W|^2 - \frac{2}{X}|W|^2\,. \tag{35}$$

Let $\phi = (\varphi,\psi)$, and assume $W = W_\varphi = \psi = 0$. Then the kinetic terms are

$$\frac{1}{X^2}\,\partial_\mu\bar\varphi\,(X + \varphi\bar\varphi)\,\partial^\mu\varphi\,, \tag{36}$$

and the potential is

$$V = |W_\psi|^2\,. \tag{37}$$

Thus for $|\varphi| \ll 1$ we have canonical kinetic terms and the potential has the global supersymmetry form, though with the additional requirements $W = W_\varphi = \psi = 0$.

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[4]In fact, any unbroken discrete (or continuous) R symmetry of the form $W \to e^{i\theta_0}W$, $\psi \to e^{i\theta_0}\psi$, $\chi^\alpha \to e^{i\theta_\alpha}\chi^\alpha$, $\theta_\alpha \neq \theta_0$ would suffice.

B. SO$(m,2)$/[SO(m) × SO(2)]

Another example of Eq. (32) is

$$K = -\ln\left(1 - \sum_{\alpha=1}^{m} \phi^\alpha \bar{\phi}^\alpha + \frac{1}{4}\left|\sum_{\alpha=1}^{m} \phi^\alpha \phi^\alpha\right|^2\right) . \qquad (38)$$

The corresponding Kähler manifold is SO$(m,2)$/[SO(m) × SO(2)]. For example, if $m = 2$, $\varphi = \phi^1$, and $\psi = \phi^2$ we get

$$K = -\ln\left[\left(1 - \frac{1}{2}|\varphi|^2\right)^2 - |\psi|^2\right] + O\left(\psi^2, \bar{\psi}^2\right) , \qquad (39)$$

or if $m = 3$, $\varphi = \left(\phi^1 + i\phi^2\right)/\sqrt{2}$, $\chi = \left(\phi^1 - i\phi^2\right)/\sqrt{2}$, and $\psi = \phi^3$ we get

$$K = -\ln\left[\left(1 - |\varphi|^2\right)\left(1 - |\chi|^2\right) - |\psi|^2\right] + O\left(\psi^2, \bar{\psi}^2\right) . \qquad (40)$$

VI. MODEL BUILDING

The solution described in Sec. IV suggests a natural strategy for inflationary model building. Construct a globally supersymmetric model which gives rise to inflation and satisfies Eqs. (20) and (22), at least to some approximation—see Sec. VI B. Then extend to supergravity by choosing a Kähler potential of the form of Eq. (32). However, as we shall see in Sec. VI C, it may not even be necessary for the globally supersymmetric model to give rise to inflation.

A. An example of a suitable globally supersymmetric model

Consider the following globally supersymmetric model:

$$W = \lambda_1 \varphi \chi_1 \psi_1 + \lambda_2 \chi_2^n \psi_2 , \qquad (41)$$

and

$$D = \Lambda^2 - |\chi_1|^2 - |\chi_2|^2 + |\psi_1|^2 + n|\psi_2|^2 . \qquad (42)$$

This model is invariant under the R parity

$$\psi_1 \to -\psi_1, \quad \psi_2 \to -\psi_2, \quad W \to -W , \qquad (43)$$

and the U(1) gauge symmetry[5]

$$\chi_1 \to e^{-i\theta}\chi_1, \quad \chi_2 \to e^{-i\theta}\chi_2, \quad \psi_1 \to e^{i\theta}\psi_1,$$

$$\psi_2 \to e^{in\theta}\psi_2 . \qquad (44)$$

[5]For example an "anomalous" U(1) often appears in string theory [12] with $\Lambda \sim 10^{17} - 10^{18}$ GeV *if* the dilaton is fixed near its usual value.

To obtain the effective potential during inflation we minimize the potential [Eq. (2)] for fixed φ as follows. For $|\chi_1|^2 + |\chi_2|^2 \le \Lambda^2$ the potential is minimized for

$$\psi_1 = \psi_2 = 0 , \qquad (45)$$

and so the R parity ensures that

$$W = W_\varphi = W_{\chi_1} = W_{\chi_2} = 0 . \qquad (46)$$

Therefore

$$V = |W_{\psi_1}|^2 + |W_{\psi_2}|^2 + \frac{1}{2}g^2 D^2 \qquad (47)$$

$$= \lambda_1^2 |\varphi|^2 |\chi_1|^2 + \lambda_2^2 |\chi_2|^{2n} + \frac{1}{2}g^2\left(\Lambda^2 - |\chi_1|^2 - |\chi_2|^2\right)^2 . \qquad (48)$$

Now if

$$|\varphi|^2 \ge \frac{g^2}{\lambda_1^2}\left(\Lambda^2 - |\chi_2|^2\right) , \qquad (49)$$

the potential is minimized for

$$\chi_1 = 0 . \qquad (50)$$

Then

$$V = \lambda_2^2 |\chi_2|^{2n} + \frac{1}{2}g^2\left(\Lambda^2 - |\chi_2|^2\right)^2 . \qquad (51)$$

Now if

$$\frac{\lambda_2 \Lambda^{n-2}}{g} \ll 1 , \qquad (52)$$

the potential is minimized for

$$|\chi_2|^2 \simeq \Lambda^2 - \frac{n\lambda_2^2 \Lambda^{2n-2}}{g^2} , \qquad (53)$$

and so

$$V \simeq \lambda_2^2 \Lambda^{2n} . \qquad (54)$$

Thus, from Eqs. (49) and (53), for

$$|\varphi| \ge \frac{\sqrt{n}\,\lambda_2 \Lambda^{n-1}}{\lambda_1} , \qquad (55)$$

we have a positive potential energy density and a flat potential for the inflaton field φ.

The above globally supersymmetric model satisfies the conditions Eqs. (20) and (22) [Eqs. (45) and (46)] and so if the Kähler potential is of the form of Eq. (32) then the supergravity corrections will not spoil the flatness of the inflaton's potential (which is exactly flat in this case but there are many possible sources for a small slope for the inflaton's potential—see the next section). Also, it is easy to check that the supergravity corrections do not spoil the stability of the model.

Alternatively to Eq. (52), if $n = 1$ and $g\Lambda/\lambda_2 < 1$ then the potential is minimized for $\chi_2 = 0$ and

$$V = \frac{1}{2} g^2 \Lambda^4 . \tag{56}$$

In this case the inflationary potential energy density is dominated by the D-term part of the scalar potential which might provide an alternative solution to the problem discussed in Sec. III.

B. The slope of the inflaton's potential

The solution described in Sec. IV is unlikely to hold exactly in realistic models. Small deviations from it lead to small contributions to the slope of the inflaton's potential. In some cases these could dominate the contributions coming from the globally supersymmetric model and so effectively determine the slope of the inflaton's potential. For example, if $W = W_\varphi = \psi = 0$ but $K = K_0 + \delta(\varphi, \bar{\varphi})$ where K_0 is of the form of Eq. (32), then we get $V = e^\delta V_0$ where V_0 is the potential which would have been obtained if $K = K_0$ had been used. We thus get a contribution of δ' to V'/V and of δ'' to $V''/V - (V'/V)^2$. See [9] for an explicit example. Also, $W \neq 0$ or $W_\varphi \neq 0$ would typically give a contribution to V''/V of order $|W|^2/|W_\psi|^2$ or $|W_\varphi|^2/|W_\psi|^2$, respectively.

C. Inflation without inflation in the global supersymmetry limit

Another example of a globally supersymmetric model satisfying Eqs. (20) and (22) is

$$W = \lambda\, f(\varphi)\, \chi^n \psi \,, \tag{57}$$

and

$$D = \Lambda^2 - |\chi|^2 + n\, |\psi|^2 . \tag{58}$$

For $\lambda \Lambda^{n-2} |f(\varphi)|/g \ll 1$ it gives a potential

$$V \simeq \lambda^2 \Lambda^{2n} \, |f(\varphi)|^2 . \tag{59}$$

This will not give rise to inflation in the global supersymmetry limit[6] for a generic function $f(\varphi)$. However, in the supergravity theory with the Kähler potential discussed in Sec. V A, the kinetic terms are noncanonical and diverge as $|\varphi|$ approaches one,[7] but the φ dependence of the potential is unchanged from the global supersymmetry limit. Therefore, transforming to the canonically normalized inflaton field stretches out the potential and so, assuming that $f(\varphi)$ does not diverge as $|\varphi| \to 1$, we will get a flat potential.

To illustrate this, consider the following simple example. For the case of only one φ field Eq. (36) reduces to

$$\frac{1}{X^2}\, |\partial \varphi|^2 . \tag{60}$$

For simplicity assume the phase of φ is constant. Then the canonically normalized inflaton σ is given by

$$|\varphi| = \tanh \frac{\sigma}{\sqrt{2}} . \tag{61}$$

Now during inflation $\sigma \gg 1$ and so

$$|\varphi| \simeq 1 - 2e^{-\sqrt{2}\,\sigma} . \tag{62}$$

Therefore

$$V \simeq V|_{|\varphi|=1} - 2\, \frac{dV}{d|\varphi|}\bigg|_{|\varphi|=1} e^{-\sqrt{2}\,\sigma} . \tag{63}$$

The coefficient of the exponential can be absorbed by the redefinition

$$\tilde{\sigma} = \sigma - \frac{1}{\sqrt{2}} \ln \frac{2\frac{dV}{d|\varphi|}}{V}\Bigg|_{|\varphi|=1} , \tag{64}$$

to give the inflationary potential[8]

$$V = V_1 \left(1 - e^{-\sqrt{2}\,\tilde{\sigma}}\right) , \tag{65}$$

which has only one free parameter $V_1 = V|_{|\varphi|=1}$ and that is determined by the Cosmic Background Explorer (COBE) normalization to be $V_1^{1/4} = 6 \times 10^{15}$ GeV. It is also straightforward to calculate the spectral index of the density perturbations produced during inflation [13,10],

$$n = 1 - 3\left(\frac{V'}{V}\right)^2 + 2\frac{V''}{V} \simeq 1 - \frac{2}{N} \simeq 0.96 , \tag{66}$$

which is the same as ϕ^2 chaotic inflation. However, the ratio of gravitational waves to density perturbations is [13,10]

$$R = 6\left(\frac{V'}{V}\right)^2 = \frac{3}{N^2} \sim 10^{-3} , \tag{67}$$

compared with $R = 6/N = 0.1$ for ϕ^2 chaotic inflation.

It is interesting that these results are quite robust, at least for a single inflationary degree of freedom. For example, if we had instead chosen the Kähler potential of Eq. (39), we would have got the inflationary potential $V = V_1 (1 - e^{-\sigma})$ which also gives $n = 1 - 2/N \simeq 0.96$ but slightly larger $V_1^{1/4} = 7 \times 10^{15}$ GeV and $R = 6/N^2 \sim 10^{-2.5}$.

[6]Which requires $|\varphi| \ll 1$ for consistency.

[7]Note that φ is defined only for $|\varphi| < 1$.

[8]Note that this is the potential during inflation ($\tilde{\sigma} \gg 1$). When inflation ends ($\tilde{\sigma} \sim 1$), the neglected, model dependent (i.e., f dependent) terms become important.

VII. SUPERSTRING EXAMPLES

A. Orbifold compactifications

The Kähler potential of the untwisted sector of the low-energy effective supergravity theory derived from orbifold compactification of superstrings always contains [14]

$$K = -\ln\left(S+\bar{S}\right) - \sum_{i=1}^{3} \ln\left(T_i + \bar{T}_i - |\phi_i|^2\right) , \quad (68)$$

where S is the dilaton, T_i are the untwisted moduli associated with the radii of compactification, and ϕ_i are the untwisted matter fields associated with T_i. Now if we divide the scalar fields into φ, ψ, and χ fields as

$$T_1 \in \varphi , \quad (69)$$

$$\phi_1 \in \psi , \quad (70)$$

$$S, T_2, T_3, \phi_2, \text{ and } \phi_3 \in \chi \subset \varphi , \quad (71)$$

then we get a Kähler potential of the required form [Eq. (32)]

$$K = -\ln\left(\varphi + \bar{\varphi} - |\psi|^2\right) + g(\chi, \bar{\chi}) , \quad (72)$$

and the target space duality symmetries [15],

$$T_i \to \frac{a_i T_i - i b_i}{i c_i T_i + d_i}, \quad \phi_i \to \frac{\phi_i}{i c_i T_i + d_i}, \quad a_i d_i - b_i c_i = 1 , \quad (73)$$

contain the desired R parity [Eq. (26)] on setting $b_i = c_i = 0$, $a_1 = d_1 = -1$, and $a_2 = a_3 = d_2 = d_3 = 1$.

B. More orbifold compactifications

A Kähler potential of the form

$$K = -\ln\left[\left(T+\bar{T}\right)\left(U+\bar{U}\right) - \left(B+\bar{C}\right)\left(C+\bar{B}\right)\right] , \quad (74)$$

often occurs in orbifold compactifications [14,16,17]. In particular, it can arise in orbifolds with continuous Wilson lines, in which case T corresponds to one of the T_i moduli of Sec. VII A, U is a (1,2) modulus, and B and C are continuous Wilson line moduli [17]. Now if we divide the fields as

$$T \text{ and } U \in \varphi , \quad (75)$$

$$B \text{ and } C \in \psi , \quad (76)$$

then we get a Kähler potential of the required form [Eq. (32)]

$$K = -\ln\left[\left(\varphi^1 + \bar{\varphi}^1\right)\left(\varphi^2 + \bar{\varphi}^2\right) - |\psi|^2\right] + O\left(\psi^2, \bar{\psi}^2\right) , \quad (77)$$

and the target space duality symmetries [17] contain the desired R parity [Eq. (26)].

C. Fermionic four-dimensional string models

The Kähler potential of the untwisted sector of the revamped flipped SU(5) model [18] is [19]

$$\begin{aligned} K = &-\ln\left(1 - |\Phi_1|^2 - |\Phi_{23}|^2 - |\Phi_{\overline{23}}|^2 - |h_1|^2 - |h_{\bar{1}}|^2 + \frac{1}{4}\left|\Phi_1^2 + 2\Phi_{23}\Phi_{\overline{23}} + 2h_1 h_{\bar{1}}\right|^2\right) \\ &-\ln\left(1 - |\Phi_2|^2 - |\Phi_{31}|^2 - |\Phi_{\overline{31}}|^2 - |h_2|^2 - |h_{\bar{2}}|^2 + \frac{1}{4}\left|\Phi_2^2 + 2\Phi_{31}\Phi_{\overline{31}} + 2h_2 h_{\bar{2}}\right|^2\right) \\ &-\ln\left(1 - |\Phi_4|^2 - |\Phi_5|^2 - |\Phi_3|^2 - |\Phi_{12}|^2 - |\Phi_{\overline{12}}|^2 - |h_3|^2 - |h_{\bar{3}}|^2 \right. \\ &\left. \qquad + \frac{1}{4}\left|\Phi_4^2 + \Phi_5^2 + \Phi_3^2 + 2\Phi_{12}\Phi_{\overline{12}} + 2h_3 h_{\bar{3}}\right|^2\right) , \end{aligned} \quad (78)$$

all three parts of which are of the form of Eq. (38).[9] Furthermore, if we divide the fields as

$$\Phi_4 \text{ and } \Phi_5 \in \varphi , \quad (79)$$

$$\Phi_3, \Phi_{12}, \Phi_{\overline{12}}, h_3 \text{ and } h_{\bar{3}} \in \psi , \quad (80)$$

$$\Phi_1, \Phi_2, \Phi_{23}, \Phi_{\overline{23}}, \Phi_{31}, \Phi_{\overline{31}}, h_1, h_{\bar{1}}, h_2 \text{ and } h_{\bar{2}} \in \chi \subset \varphi , \quad (81)$$

then we get a Kähler potential of the required form [Eq. (32)]

$$K = -\ln\left(1 - |\varphi|^2 + \frac{1}{4}\left|\varphi^{\mathrm{T}}\varphi\right|^2 - |\psi|^2\right) + g(\chi, \bar{\chi}) + O\left(\psi^2, \bar{\psi}^2\right) , \quad (82)$$

and the target space duality symmetries [19] contain the desired R parity [Eq. (26)].

[9] Up to trivial redefinitions.

D. Calabi-Yau compactifications

Here I give the Calabi-Yau manifold discussed in Sec. 4.3 of [20] as another example of a compactification of superstrings that can give Kähler potentials of the form discussed in Sec. V. At a particular point in the moduli space of the above-mentioned Calabi-Yau manifold, the low-energy gauge group includes an extra $\mathrm{U}(1)^4$ factor, a subgroup of which may be preserved on subspaces of the moduli space that pass through that point. The Kähler potential on a subspace of the moduli space that preserves a $\mathrm{U}(1)^3$ subgroup of the extra $\mathrm{U}(1)^4$ gauge symmetry is [20]

$$K = -\ln\left(1 - |N|^2 - |C|^2\right) + O\left(C^2, \bar{C}^2\right) , \qquad (83)$$

where $N = (N_1, N_2)$ are the neutral[10] (1,2) moduli that span the subspace, and C is a vector whose components are 63 of the 99 charged (1,2) moduli and their associated matter fields. Also the Kähler potential on a subspace of the moduli space that preserves a $\mathrm{U}(1)^2$ subgroup of the extra $\mathrm{U}(1)^4$ gauge symmetry is [20]

$$K = -\ln\left(1 - |N|^2\right) , \qquad (84)$$

where N is a vector whose components are the 12 neutral (1,2) moduli that span the subspace.

[10]With respect to the unbroken $\mathrm{U}(1)^3$.

VIII. CONCLUSIONS

A globally supersymmetric model of inflation (see, for example, [9,21]) will not work in a generic supergravity theory because the higher-order, non-renormalizable supergravity corrections destroy the flatness of the inflaton's potential. In this paper I have derived a form for the Kähler potential which eliminates these corrections if $W = W_\varphi = \psi = 0$ during inflation (where W is the superpotential, the inflaton $\in \varphi$, and $W_\psi \neq 0$). It is encouraging that Kähler potentials of the required form often occur in superstrings and that the target space duality symmetries of superstrings often contain R parities which would make $W = W_\varphi = 0$ automatic for $\psi = 0$.

Also, I have shown that supergravity theories with Kähler potentials of this form may give rise to inflation even if the corresponding globally supersymmetric theory does not. The simplest examples of this new idea for inflation give a spectral index $n = 1 - 2/N \simeq 0.96$ for the density perturbations and negligible gravitational waves, though more complicated examples lose this predictive power.

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