Cosmological stabilization of moduli with steep potentials

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A scenario which overcomes the well-known cosmological overshoot problem associated with stabilizing moduli with steep potentials in string theory is proposed. Our proposal relies on the fact that moduli potentials are very steep and that generically their kinetic energy quickly becomes dominant. However, moduli kinetic energy redshifts faster than other sources when the universe expands. So, if any additional sources are present, even in very small amounts, they will inevitably become dominant. We show that in this case cosmic friction allows the dissipation of the large amount of moduli kinetic energy that is required for the field to be able to find an extremely shallow minimum. We present the idea using analytic methods and verify with some numerical examples.

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I. INTRODUCTION

Let us start with a brief review of the framework in which the problem of stabilizing string moduli in the perturbative outer region of moduli space, where the string coupling is weak and/or the volume of the compact dimensions is large, is posed.

Realistic models, for example, in flux compactifications of string theory, usually include an effective $N = 1$ supergravity (SUGRA) theory below the string scale $M_s = 10^{-2}$ (our conventions are such that $M_p \equiv \frac{1}{\sqrt{8\pi G_N}} = 1.2 \times 10^{18}$ GeV $= 1$) and supersymmetry (SUSY) breaking in some hidden sector at an intermediate scale $M_I = 10^{-7}$. General arguments based on symmetries show that the moduli superpotential must be a sum of exponentials in the moduli and perhaps an additional constant. These exponentials could be generated by stringy or field theoretic nonperturbative effects.

In this framework, a typical potential for moduli fields σ_i has the $N = 1$ SUGRA form:

$$V_{\rm SUGRA} = e^K(K^{i\bar{j}}F_iF_{\bar{j}} - 3|W|^2), \tag{1}$$

and a typical superpotential has the form:

$$W(\sigma) = \sum_i A_i e^{-\alpha_i\sigma}, \tag{2}$$

where $K^{i\bar{j}}$ is the inverse metric derived from the Kahler potential K and $F_i = \partial_i W + \partial_i K W$ is the Kahler derivative. Here we will consider a situation where all but one of the moduli have been stabilized at the string scale so that we can focus on the dynamics of one light modulus σ as in the recent work of Kachru *et al.* (KKLT)[1].

The Kahler potential of moduli in the perturbative region is typically logarithmic in the fields, so that in terms of a canonically normalized component field ϕ the potential involves exponentials of exponentials. An important feature of such potentials is that they vanish or approach a finite constant when $\phi \to \infty$. This corresponds to decompactification if ϕ is the volume modulus and zero coupling if ϕ is the effective dilaton. This is the well-known Dine-Seiberg problem [2]. Therefore, if the potential has a minimum, it needs to be separated from the asymptotic region by a potential barrier.

Realistic models need some fine-tuning of parameters. SUSY breaking in the observable sector at the TeV scale requires that at the minimum of the potential $F_{\rm min} = M_I^2 = O(10^{-14})$. In addition, if one assumes that the recent cosmological observations indeed indicate that the cosmological constant (CC) is nonvanishing then at the minimum $V_{\rm min} = O(10^{-120}) > 0$[1], so $|F_{\rm min}| = \sqrt{3}|W_{\rm min}| + O(10^{-120})$. A stable minimum further requires tuning of at least four parameters to prevent tachyonic directions. Until recently it was very hard to find a single working model because the framework and parameters were too constrained. Now, with the development of models based on flux compactifications and the understanding of their vast parameter space, the discretuum, it has become possible to find models with minima in the outer region of moduli space [1].

We are interested in estimating the height of the barrier that separates the minimum of the potential from the asymptotic region where the potential vanishes or approaches a constant. Let us assume that at some value $\phi_{\rm min}$ the potential has a true minimum. Since each exponential term is smaller in absolute value for $\phi > \phi_{\rm min}$, we can generically expect that for $\phi > \phi_{\rm min}$, $|F| < |F_{\rm min}|$, and $|W| < |W_{\rm min}|$ so that the height of the separating barrier is at most limited by the intermediate scale $V_{\rm max} \sim M_I^4 = 10^{-28}$. If no further tuning is performed the height of the barrier is much lower than this estimate. Typical moduli potentials are therefore steep and have a

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[1]This assumption will not be very important for us. All our considerations go through almost unmodified if $V_{\rm min} = 0$.

very shallow minimum, when a stable one exists. A typical steep moduli potential with a shallow minimum is shown in Fig. 1.

If one considers time-dependent solutions, one encounters a cosmological version of the Dine-Seiberg problem, first discussed by Brustein and Steinhardt [3]. If the moduli start from a generic point on the potential they are expected to reach the outer region of moduli space by classically rolling towards the asymptotic region. If they start to the left of the minimum, they will roll over the shallow barrier, and if they start to the right of the barrier, they will never reach the minimum and roll to the asymptotic outer region. To avoid this without additional sources, such as radiation, the initial position of the field, ϕ_0, has to be such that the initial height of the field, $V(\phi_0)$, is at most an order of magnitude larger than the height of the barrier. Thus getting a bound solution requires fine-tuning of initial conditions to a very high accuracy. In addition, the steep potential inhibits the possibility of inflation while the moduli are rolling, since moduli kinetic energy tends to dominate the energy budget of the universe. The consideration of time-dependent solutions thus leads to additional criteria of stability of moduli beyond the standard static stability criteria, and therefore leads to additional requirements and constraints beyond the standard criteria which guarantee the absence of tachyonic directions. Such arguments have led us to propose previously that the most likely place for moduli stabilization is the central region of moduli space [4].

Several previous attempts to resolve the cosmological stability of moduli in general, and, in particular, the cosmological overshoot problem, have been made. Banks *et al.* have emphasized the role of the nonzero modes [5] and have noticed that they redshift slower than zero modes. Dine [6] noticed the possible role of such slower redshift in stabilizing moduli. Barriero *et al.* [7] discovered that the presence of additional sources helps to

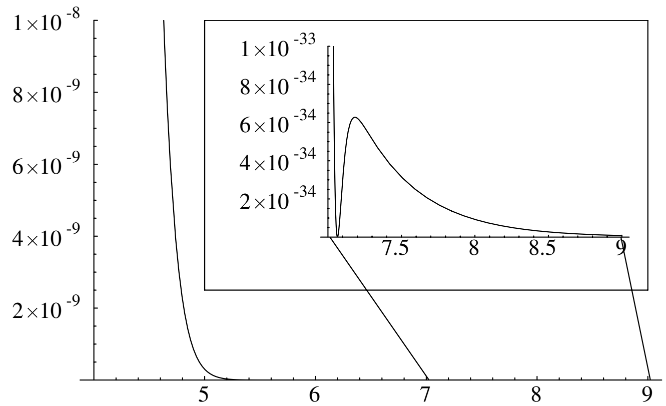

FIG. 1. A typical moduli potential. The region of the shallow minimum had to be magnified by 26 orders of magnitude so that it can be seen. The vertical axis is in units of M_p^4, and the horizontal axis is in units of M_p.

relax the problem. Furthermore Huey *et al.* [8] determined that if the moduli receive some specific temperature corrections then the other sources (i.e., radiation) could be much more effective. We will comment on this further below. Of course all these discussions were made in the context of then known string models which were much more constrained than the models that we have now.

Our proposed resolution relies on the existence of other sources, for example, a gas of relativistic particles: radiation. The nature of the additional sources and their energy density is not particularly important to us. A key fact that is crucial to the scenario that we propose is that kinetic energy (KE) in an expanding universe redshifts at the fastest rate of all known sources. As the field rolls on its steep potential its KE builds up and very quickly leads to KE dominance. Since KE redshifts faster than the other sources, the additional sources will eventually become the dominant energy component. While these sources dominate, they create a large amount of cosmic friction which dissipates a large amount of energy, allows the field to gently land into the basin of attraction of the shallow minimum, and to eventually settle down at the minimum. While these additional sources dominate, the field moves only a finite amount.[2, 3]

We do not rely on additional temperature dependent coupling beyond cosmic friction. An example of such coupling is $T^2\phi^2$. In fact, in the particular example that we use to illustrate our idea we assume that they are absent. If such couplings exist (as assumed by Huey *et al.* [8]) they could further help in relaxing the constraints on moduli evolution. However, it has been argued that high temperature effects do not modify the form of the potential of string theoretic moduli [12]. More recently, extending previous calculations to higher order, and considering heterotic moduli that are not in thermal equilibrium, it was shown in [13] that corrections proportional to T^4 could change the potential in the vicinity of the minimum in a significant way. These corrections affect the evolution of the moduli mostly near the shallow minimum.

Since our main goal here is to explain our idea and show that it can be realized rather than to explore in a general and systematic way the various possibilities and caveats, we concentrate on some explicit examples that allow us to discuss the basic argument. We will present a systematic search in a forthcoming publication. In Sec. II we use approximate analytic solutions to study modular cosmology with sources, and in Sec. III we verify numerically that our approximations and quantitative estimates are valid. Sec. IV contains our conclusions and a brief description of possible extensions of our scenario.

[2]This fact was noticed by previous authors in a different context [9,10].

[3]Similar phases of moduli evolution were discussed in a different context in [11].

II. MODULAR COSMOLOGY IN THE PRESENCE OF SOURCES

We will discuss a specific model to expose the idea on which our proposal is based and to explain its basic ingredients. We consider a cosmology with a single field that has a potential of the form shown in Fig. 1. The field is assumed to start in the steep region of the potential. Some amount of radiation in a thermal state (with constant entropy so that $T \sim 1/a$) is also assumed to be present $\rho_{\rm rad} = CT^4 = \frac{c^2}{a^4}$ where C and c are constants, T is the temperature and a is the Friedman-Robertson-Walker (FRW) scale factor. We further assume for simplicity a spatially flat universe. As we have explained, our proposal is not particularly sensitive to the nature of the additional sources. In fact all that is required is some additional source that redshifts slower than kinetic energy (see below). Here we will just incorporate radiation—in a later paper we will consider more general sources in detail. The qualitative features that we wish to illustrate will however remain the same.

We do not discuss here the evolution prior to the epoch in which we can treat the effective dynamics as a single scalar field in a cosmological background, or whether the universe starts in a quantum region. In both of these cases the universe would arrive at a starting point that is similar to the one that we assume. Such initial conditions can be arrived at in different ways. For example, a short period of inflation, a phase of so-called pre-big bang evolution, or nucleation from nothing.

The equations of motion that we need to solve are therefore

$$H^2 = \frac{1}{3}\left[\frac{\dot{\phi}^2}{2} + V(\phi) + \frac{c^2}{a^4}\right], \qquad \ddot{\phi} + 3H\dot{\phi} + \frac{\partial V}{\partial \phi} = 0. \tag{3}$$

Here a dot represents a derivative with respect to time, and H is the Hubble parameter $H = \dot{a}/a$.

Our scenario includes four distinct epochs that are described below. Each epoch starts with some specific initial values for the Hubble parameter and the field and its derivatives. The end values of the previous epoch provide initial values for the subsequent epoch. For definiteness, at the end of stage i we denote the values of the relevant variables by the subscript i so for example t_2 is the time at which epoch number two ends.

(a) *Epoch one*: Potential domination
The field starts with zero initial velocity and a large potential energy on the steep part of the potential. The universe expands at a fast rate. If a substantial amount of radiation (or other sources) is initially present, the radiation energy density quickly redshifts as a^{-4}, and the main source of energy becomes the potential energy of the field.
The equations of motion can be approximated by

$$\ddot{\phi} + \frac{\partial V}{\partial \phi} = 0 \qquad H = \sqrt{\frac{V}{3}}. \tag{4}$$

The energy $E = \frac{\dot{\phi}^2}{2} + V(\phi)$ is conserved in this epoch so all the potential energy is converted into the field's kinetic energy (KE).
The solution of the approximate Eqs. (4) with the initial conditions at $t = 0$, $\phi = \phi_0$, $V(\phi_0) = V_0$, and $\dot{\phi} = 0$ is the following,

$$\dot{\phi} = \sqrt{2[V_0 - V(\phi)]} \qquad t = \int \frac{d\phi}{\sqrt{2[V_0 - V(\phi)]}}$$
$$\ln a = \frac{1}{\sqrt{6}} \int_{\phi_0}^{\phi} \frac{d\phi \sqrt{V(\phi)}}{\sqrt{V_0 - V(\phi)}}. \tag{5}$$

The velocity of the field at the end of this epoch

$$\dot{\phi}_1 = \sqrt{2[V_0 - V(\phi_1)]}, \tag{6}$$

can be quite large if the potential is steep $V_0 \gg V(\phi_1)$, and the KE of the field becomes the dominant energy component.

(b) *Epoch two*: Kinetic energy domination
The equations of motion in this epoch can be approximated by

$$\ddot{\phi} + 3H\dot{\phi} = 0 \qquad H^2 = \tfrac{1}{6}\dot{\phi}^2. \tag{7}$$

In epoch two the equation of motion for ϕ can be implicitly solved $\dot{\phi} = \dot{\phi}_1 a_1^3/a^3$, implying that $KE = \frac{1}{2}\dot{\phi}^2$ redshifts quickly $KE = \frac{1}{2}\dot{\phi}_1^2 a_1^6/a^6$.
The solution of the approximate Eqs. (7) with the initial conditions at $t = t_1$, $\phi = \phi_1$, and $\dot{\phi} = \dot{\phi}_1 = \sqrt{2[V_0 - V(\phi_1)]}$, given in Eq. (6) is the following,

$$\phi - \phi_1 = \sqrt{\frac{2}{3}} \ln\left[\sqrt{\frac{3}{2}}\dot{\phi}_1 (t - t_1) + 1\right]$$
$$a(t) = \left[\sqrt{\frac{3}{2}} a_1^3 \dot{\phi}_1 (t - t_1) + a_1^3\right]^{1/3}. \tag{8}$$

We can relate the amount of KE dissipation to the displacement in the field during this epoch

$$\ln\left(\frac{\dot{\phi}_1^2}{\dot{\phi}^2}\right) = \sqrt{6}(\phi - \phi_1), \tag{9}$$

which means that to dissipate a few tens of orders of magnitude in KE, as required, the field needs to move quite a bit in (reduced) Planck units to the extreme outer region of moduli space. If the field moves a several Planck lengths, as is the case in models that we have considered, then only a few orders of magnitude of KE are dissipated, approximately 1 order of magnitude per planck length displacement of the field.

As the universe expands, both the KE and the radiation energy density redshift. Since the radiation redshifts at a slower pace, it will eventually "catch up" with the KE no matter how small its initial value. A possibility is that the potential energy will become dominant before the radiation. Whether this happens or not depends on the details of the model. In the successful cases (in the case of steep potentials) the radiation becomes dominant first.

(c) *Epoch three*: Radiation domination
Epoch two leads to a radiation dominated epoch when the radiation "catches up" with the KE, $(\rho_{\rm rad})_2 \gg (KE)_2$,

$$\frac{c^2}{a_2^4} \gg \frac{\dot{\phi}_1^2}{2}\frac{a_1^6}{a_2^6}. \qquad (10)$$

In epoch three the KE continues to redshift as $K = \frac{\dot{\phi}_1^2}{2}\frac{a_1^6}{a^6}$, and the expansion of the universe is faster, hence in this epoch KE is dissipated in a more efficient way.
The equations of motion in epoch three can be approximated by

$$\ddot{\phi} + 3H\dot{\phi} = 0 \qquad H^2 = \frac{1}{3}\frac{c^2}{a^4}. \qquad (11)$$

The solution of Eqs. (11) is given by

$$a = \sqrt{\frac{2c}{\sqrt{3}}(t - t_2) + a_2^2}$$
$$\phi = \phi_2 + \frac{\sqrt{3}\dot{\phi}_1 a_1^3}{c a_2} - \frac{\sqrt{3}\dot{\phi}_1 a_1^3}{c}\frac{1}{\sqrt{\frac{2c}{\sqrt{3}}(t - t_2) + a_2^2}}. \qquad (12)$$

The displacement of the field during this epoch can be expressed after some algebra as follows:

$$\phi - \phi_2 = \sqrt{6}\sqrt{\frac{(KE)_2}{(\rho_{\rm rad})_2}}\left(1 - \frac{a_2}{a}\right) \le \sqrt{6}. \qquad (13)$$

Therefore, even if this epoch continues indefinitely the field will only move a finite distance [9,10]. KE on the other hand, continues to redshift as $1/a^6$, so contrary to epoch 2, the field can dissipate a lot of energy while staying almost constant.

(d) *Epoch four*: Potential domination
Since the radiation energy density redshifts quickly, eventually, the potential energy density no matter how small, will come to dominate when $V(\phi) \gg \frac{c^2}{a_2^4}$. In this epoch the equations of motion can again be approximated by

$$\ddot{\phi} + \frac{\partial V}{\partial \phi} = 0 \qquad H = \sqrt{\frac{V}{3}}. \qquad (14)$$

At this point the solution of this system depends on the value of ϕ, and whether it ended up in the basin of attraction of a minimum, beyond the last minimum, or still at a high point. If the transition from KE dominance to RD occurs at a point that is close enough to the minimum (and to the left of the maximum), then the field will be trapped even in a shallow minimum since it has lost a huge amount of KE.

The conclusion of this analysis is that as the field acquires a large amount of KE the latter seeds the elements of its own destruction by the lurking radiation. In the KE domination epoch the field does not move much per order of magnitude of dissipated KE, and in the radiation dominated epoch its motion is clearly bounded. Thus if the potential is steep enough, so that enough kinetic energy is acquired in the initial stage, and friction becomes important early on, and the difference in elevation between the starting point and the minimum is large enough, then with generic initial conditions the field would be bound.

Obviously the detailed realization of this scenario depends on the values of the parameters in the potential and the nature of the specific sources.

III. NUMERICAL EXAMPLES

To verify that our approximations are meaningful, and to check their range of validity, we have made a series of numerical investigations. We report here only about some of them to illustrate some features of the scenario.

For the numerical investigation we have used the following potential for the volume modulus used in [1]:

$$V(\sigma) = \frac{aAe^{-a\sigma}}{2\sigma^2}\left(\frac{1}{2}\sigma aAe^{-a\sigma} + W_0 + Ae^{-a\sigma}\right) + \frac{d}{\sigma^3}, \qquad (15)$$

where $\sigma \equiv e^{\sqrt{2/3}\phi}$. This potential was derived using the superpotential (coming from flux contributions and nonperturbative effects)

$$W(\sigma) = W_0 + Ae^{-a\sigma}, \qquad (16)$$

and the classical tree level Kahler potential. The last term in Eq. (15) comes from the contribution of an anti-Dbrane (Dbar brane). We stress that we use this particular potential merely for the sake of illustration of our mechanism. Similar results follow for potentials which incorporate alternatives to the Dbar brane term.

All our numerical investigations were done in the conventions that $M_p^2 \equiv \frac{1}{8\pi G_N} = 1$. We present our results in Figs. 2–5. To construct the numerical examples in these figures we used the following values for our parameters: $a = 0.1$, $A = 1.0$, $d = 3 \times 10^{-26}$, $W_0 = -2.96 \times 10^{-13}$. For these parameters the potential has a true minimum at $\phi = 7.06$. At the minimum the value of the potential (the CC) is 6.35×10^{-42} and the barrier separating the minimum from the asymptotic region is

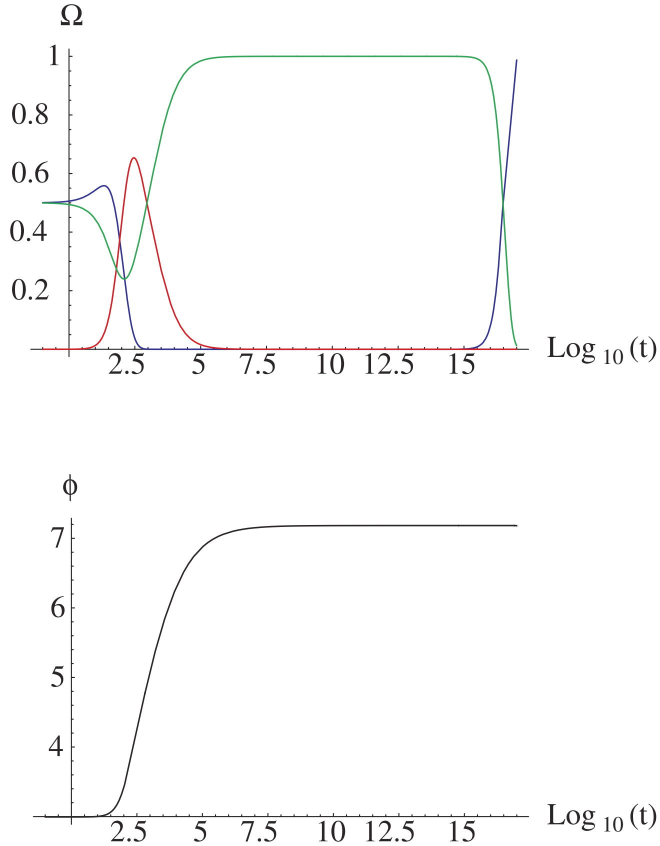

FIG. 2 (color online). A bound case. Shown in the top panel are fractional densities of potential energy (blue, dark), kinetic energy (red, medium) and radiation energy (green, light) as a function of time. KE becomes dominant and then the radiation. Shown in the bottom panel is the evolution of the field as a function of time ending in the shallow minimum of the potential.

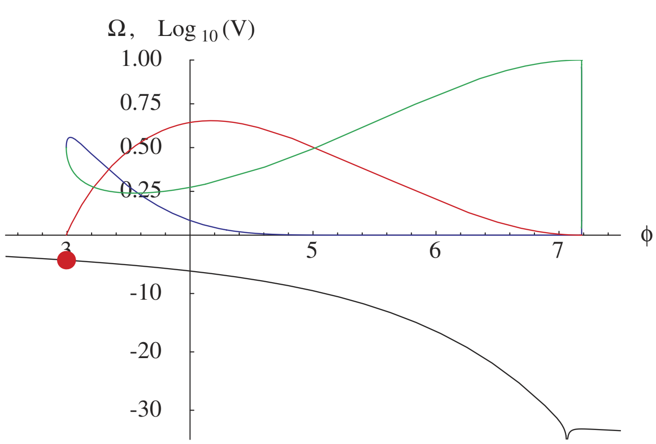

FIG. 3 (color online). A bound case. Shown in the top panel are the various energy densities (color coded as in Fig. 2) as a function of the scalar field position. Shown in the bottom panel is the potential and the starting position of the field. Note that the scale is logarithmic and that the difference in potential energy between the starting point and the minimum is about 30 orders of magnitude.

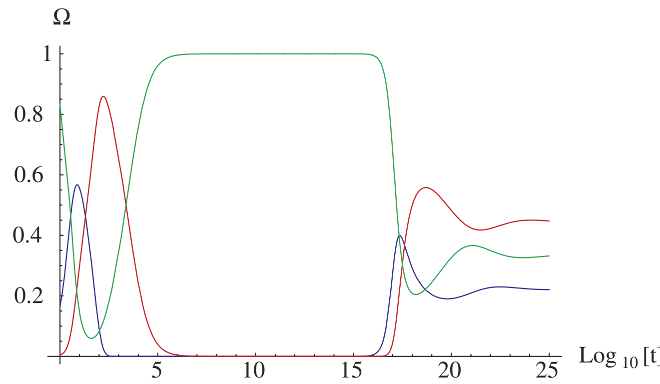

FIG. 4 (color online). An unbound case: various energy densities (color coded as in Fig. 2) as a function of time. The only difference from the bound case is where the field ends up at the end of the radiation dominated epoch. In this case it lands to the right of the shallow minimum and continues to run on the potential in a "tracking" solution.

located at $\phi = 7.18$. The height of the barrier is 6.28×10^{-34}.

To illustrate the effect of radiation on the evolution of a moduli field we present examples of a bound and an unbound case. These two examples differ only in initial conditions. To create these two examples we set the initial conditions as follows. For the bound case: $V_0 = (\rho_{\rm rad})_0 = 5.20 \times 10^{-5} M_p^4$. The corresponding initial value of the field is $\phi_0 = 2.99$, and the velocity of the field vanishes initially. For the unbound case: $V_0 = (\rho_{\rm rad})_0 = 3.78 \times 10^{-2} M_p^4$. The fractional energy densities created by our bound solution are shown in Figs. 2 and 3. The first

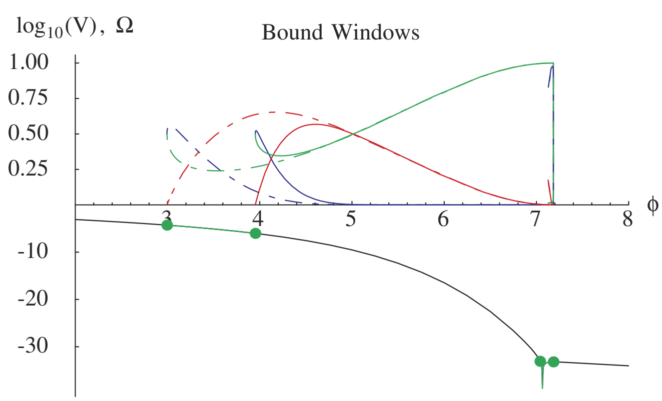

FIG. 5 (color online). Window of allowed initial conditions for bound solutions. Shown in the top panel are the various energy densities as a function of scalar field position (color coded as in Fig. 2) for two bound solutions with two different initial conditions, depicted by solid and dashed curves. In the lower panel we show the potential as a function of the scalar field. The dots on the potential are the end points of the two regions of initial conditions that lead to bound solutions. One of the regions is around the minimum, and the other is way up on the potential.

depicts the value of the field and the fractional energy densities as a function of time, while the second shows the same fractional energy densities and the potential as a function of the field. The four epochs mentioned previously in this paper are clearly visible in both representations. The fractional energy densities created by our unbound solution are shown in Fig. 4. Note that a significant fourth epoch is not present in this example. The field lands to the right side of the barrier, where a scaling solution quickly takes over.

The addition of radiation to our model created an extra window, (2.99, 3.95), of initial conditions which lead to bound solutions. The other interval, (7.04, 7.18), remains relatively unchanged with the addition of radiation. The two intervals of bound solutions can be seen in Fig. 5.

IV. DISCUSSION AND OUTLOOK

We have proposed a scenario for resolving the cosmological moduli stabilization problem. We believe that we have identified the basic ingredients of the required solution. The application of the idea may vary according to the detailed models of string compactifications.

Obviously this preliminary investigation needs to be followed by detailed and quantitative exploration, including various sources, a systematic study with a wider range of parameters and potentials, and a quantitative analysis of stable potentials. In particular, the dependence of the range of bound initial conditions on the parameters of the potential such as the width and height of the minimum and barrier.

Having found cosmologically stable models, it is clear that string models become less constrained. This provides a new perspective on the approach to string model building. Previously, because the theory was highly constrained, the hope was that only a single model, or a single class of models will satisfy the necessary conditions. Now the theory needs to be constrained as much as possible by data and phenomenological bottom-up constraints.

As a final comment we wish to point out that the presence of the fourth epoch (of potential domination) does not necessarily imply slow roll inflation; however, constructing models in this framework which produce a sufficient amount of primordial inflation is straightforward. First, one must set the parameters of the potential so that the barrier is very flat (i.e., $\frac{V''}{V} \ll 1$) at the top. Such an example has been produced recently using the imaginary component (the axion) of the volume modulus, ϕ [14]. Second, one must set the initial conditions of the model so that the radiation places the field close to the top of the barrier. It has not yet been possible to find a model which is not so severely fine tuned so that it may survive quantum corrections. We are currently investigating this issue and hope to report our results in a future publication. Other models of inflation from a flat maximum of the potential of stabilized moduli were introduced in [15,16]. Additional regions in the discretuum may allow the construction of other types of inflationary models.

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