Towards a Solution of the Moduli Problems of String Cosmology

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Abstract

There are at least two serious moduli problems in string cosmology. The first is the possibility that moduli dominate the energy density at the time of nucleosynthesis. The second is that they may not find their minima all together. After reviewing some previously proposed solutions to these problems, we propose another: all of the moduli but the dilaton sit at points of enhanced symmetry. The dilaton has a potential similar to those of racetrack models; it is very massive and its dynamics do not break supersymmetry. The dilaton is able to find the minimum of its potential because the energy is dominated by non-zero momentum modes. This energy need not be thermal. The effective potential for the dilaton is quite different from its flat space form. If certain conditions are satisfied, the dilaton settles into the desired minimum; if not, it is forced to weak coupling.
1 Introduction: Moduli in Cosmology

Moduli are ubiquitous in string models, but they are almost certainly not present in nature. Most ideas for fixing the moduli still leave over approximate moduli. In pictures where supersymmetry is relevant to understanding the hierarchy, there are some approximate moduli which are associated with supersymmetry breaking (e.g. in models of gluino condensation\[1\], Kahler stabilization\[2\], or the racetrack picture\[3, 4, 5\]). In scenarios with large or warped dimensions, there are approximate moduli associated with the size of the extra dimensions. These approximate moduli are simultaneously interesting and dangerous for cosmology. They are interesting because they are candidate inflatons\[6, 7\], and could play a role in generating the baryon asymmetry\[8\]. They are problematic because they can easily carry too much energy\[9\], or simply fail to find their minima altogether\[10\]. The first of these difficulties is usually called the “cosmological moduli problem,” the second the “Brustein-Steinhardt problem.”

Several solutions to the cosmological moduli problem have been proposed. The moduli might simply be much more massive than one might have expected from considerations of hierarchy; if their masses are of order 10’s of TeV, or their interactions somewhat stronger than expected from naive considerations, their decays can restart nucleosynthesis. The baryon number can be produced either directly in their decays, or through the Affleck-Dine mechanism\[11\]. Late inflation\[12\] (and thermal inflation\[13\]) have been discussed as solutions.

One particularly simple possibility is that of “maximally enhanced symmetry”\[14, 15, 16\]. This is the idea that all of the moduli transform non-trivially under unbroken (or slightly broken) symmetries. The main difficulty with this idea is connected with what we will refer to as the dilaton, $S$, that modulus (or more generally moduli) which controls the values of the standard model gauge couplings. It seems unlikely that the dilaton would have an enhanced symmetry at a point where the gauge couplings are small (some ideas for how a weak coupling might emerge were discussed in \[17\]). So one might consider, instead, the possibility that the dilaton does not sit at an enhanced symmetry point; this point is determined by supersymmetry-preserving dynamics at some high scale. Any other moduli sit at enhanced symmetry points. Provided that the mass of the dilaton is greater than a few 10’s of TeV (in the mechanisms to be discussed below the natural mass scale is much larger), this would solve the cosmological moduli problem: the dilaton would decay early; the other moduli could naturally start out near their minima\[5\].

This still leaves the Brustein-Steinhardt problem, and this will be the focus of our attention in what follows. The question is: why should the dilaton end up anywhere near the correct
vacuum. In most pictures for the origin of the dilaton potential, the potential is extremely steep. This applies, in particular, to both the racetrack models\[3, 5\] and to models of Kahler stabilization\[2\]. As a result, if one assumes that the zero mode of the dilaton dominates the energy density, one finds that the field inevitably overshoots the minimum, except for very special initial conditions.

This problem is troubling, but a number of mechanisms have been suggested which might mitigate it. Horne and Moore\[17\] noted that the dilaton moduli space has finite volume. As a result, there is a finite probability that the system will start out with a suitable initial coordinate and velocity so as to end up in the correct vacuum. Banks et al\[18\] have made several other points concerning the Brustein-Steinhardt problem. First, they note that there may be modifications to the Kahler potential which affect the steepness of the potential. As we will see, this is relevant to both the Kahler stabilization and the racetrack schemes. On the other hand, we will see that the required modifications of the Kahler potential are not very plausible. Essentially, it is necessary to fine tune the Kahler potential over a significant volume of the field space.

Second, and most importantly, Banks et al point out that it is not consistent to focus simply on the zero modes. We will review this argument, and pursue its implications. The energy stored in non-zero modes might be thermal or not. In either case, there is an effective potential for the dilaton, which one can think of as arising from the usual coupling constant dependent corrections to the free energy and/or the kinetic terms of the various fields. If the system is thermal, the thermodynamic free energy is a function of the couplings, \(\Omega(T, g^2)\). If the coupling is dynamical, and if the dilaton mass is small compared to the temperature (which it is for temperatures below \(M_p\)), this is just the dilaton potential. Even if the energy is non-thermal, and carried by non-zero (or zero) modes of some field \(\Phi\) (which might be the dilaton itself), there will be corrections to the energy of the system in powers of the coupling constant, which again constitute a potential for the dilaton. In either case, at weak coupling, this potential is likely to be far larger than the non-perturbative potential. The minimum of this potential will not coincide with \(S_0\), that of the zero temperature, flat space theory, so the dynamics is quite different from that considered in \[10\]. The zero coupling limit can be a minimum or a maximum of the potential. If it is a minimum, the system may be driven to weak coupling. If it is a maximum (i.e. if the potential tends to zero from below) the system may well be set gently into the ground state. As we will see, even if the weak coupling point is a minimum of the potential, under plausible conditions, the system finds its way to the true minimum. If these conditions are not satisfied, the system is driven to weak coupling.

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If the system is truly in thermal equilibrium, one can be rather definite, and we will consider this case as well. At very weak coupling, one can compute the potential. However, one expects the true minimum to lie at a point where weak coupling methods are not reliable. For gauge groups typical of the racetrack scheme, we will see that perturbation theory is not reliable, and we need to make some assumptions about the form of the potential. With plausible assumptions, the system does indeed find its ground state. But it is also quite possible for the system to move continually to ever weaker coupling.

In non-supersymmetric scenarios, such as the large radius and warped compact dimension pictures which have been developed recently [19, 20, 21], it is harder to make definite statements. One doesn’t have quite as detailed a picture of the underlying theory in this case (e.g. one doesn’t usually know much about the moduli which determine the gauge couplings). We will make some comments on these ideas in our conclusions.

In sum, we conclude that the usual formulation of the Brustein-Steinhardt problem is not appropriate. The answer to the question: do the moduli end up close to their true minima, depends on the behavior of non-zero momentum modes. With quite plausible assumptions about the potentials and about the conditions of the early universe, the system can find its stable ground state. Actually determining the course of events requires far greater knowledge than we have at present about the stabilization of the moduli and about the initial conditions in the early universe.

2 Why is the Dilaton Special?

In the weakly coupled picture of the heterotic string theory, the gauge couplings are controlled by a field, $S$. The universality of the $S$ couplings can lead to coupling unification. Given that we don’t expect this weak coupling picture to hold in any detailed way (if at all), what field might be singled out? Consider, for example, the strong coupling limit of the heterotic string. Here, the fields $S$ and $T$ are both large, and $T$-dependent corrections to the gauge couplings are important and can potentially spoil the prediction of unification [22, 23].

In the enhanced symmetry picture which we are proposing, it is natural to suppose that all but one linear combinations of moduli are fixed and of order one in the appropriate fundamental units. One linear combination is not. Call this combination $S$. Suppose it couples to both a “hidden sector” gauge group, as well as to the ordinary gauge groups. Then, as we will discuss below, one can consider various mechanisms which stabilize $S$ at a large value, giving rise to
weak couplings. Whether couplings are unified is a separate question. This depends on whether it is reasonable that the $S$ couplings are universal. Similarly, whether this picture is realized in any known approximate moduli space is a question which requires investigation. In the naive Horava-Witten picture, for example, both the $S$ and $T$ moduli are large, but for special backgrounds there is unification. These issues are under study and will be discussed elsewhere.

For now, we will simply assume that it is sensible to focus on one particular modulus, as suggested by the enhanced symmetry picture we have proposed. Our cosmological remarks below are likely to be applicable to models in which several moduli suffer from the Brustein-Steinhardt difficulty.

3 Moduli Stabilization

Two ideas for stabilizing the moduli have been seriously pursued. Both focus on the puzzle: how can a theory with no small parameter generate a small gauge coupling. Both assume that the resolution to this question lies in holomorphy. More precisely, some dynamics generates a minimum for the dilaton (where we use the term in the sense of the previous section) at large $S$. Corrections to the superpotential are exponentially small in $S$, but corrections to the Kahler potential may be large.

The first of these mechanism is known as Kahler stabilization. Gaugino condensation or some similar non-perturbative effect is supposed to generate a superpotential for the dilaton, which falls to zero exponentially rapidly for weak coupling. The stabilization of the dilaton arises because of properties of the Kahler potential. It has been argued that this is plausible, since one might expect large corrections to the Kahler potential from its weak coupling value even when the coupling is rather small.

The second of these mechanisms is known as the racetrack model. Here, the idea is that the stabilization occurs through properties of the superpotential. For example, if one has two gaugino condensates, one might expect the superpotential to look as:

$$W = Ae^{-S/N_1} - Be^{-S/N_2}. \tag{1}$$

This superpotential has a stationary point at

$$S = \frac{N_1 N_2}{(N_1 - N_2)} \ln(B/A). \tag{2}$$
So if $N_1$ and $N_2$ are large, the coupling may be small. In order that any systematic analysis be possible, it is necessary that $N_1 \approx N_2$ \cite{5}. In addition, in general, one does not expect to be able to compute the Kahler potential in this scheme (though this may be possible under certain circumstances).

An appealing version of the racetrack scenario has been put forth in \cite{4}. These authors consider models with $R$ symmetries, which can give rise to unbroken supersymmetry with $W = 0$ at the minimum. From the perspective of the cosmological issues which we have discussed in the introduction, this is a particularly attractive possibility. One might imagine that the coupling is fixed at a small value, but in such a way that the dilaton is quite massive, and rather harmless in cosmology. In this model, one has a set of singlets, $X$, coupled to two gauge groups with quantum modified moduli spaces, and nearly identical beta functions. The couplings include

$$\mathcal{L} = X(A Q_1 Q_1 - B Q_2 Q_2)$$

(3)

where $X$ is one of the singlets and $Q_1$ and $Q_2$ denoting matter in the two different gauge groups. The potential, then, has the form

$$V = |A e^{-S/N_1} - B e^{-S/N_2}|^2.$$  

(4)

This has a supersymmetry-conserving minimum with $S$ given by eqn. \cite{4} above.

4 The Brustein-Steinhardt Problem

Both the Kahler Stabilization and Racetrack scenarios are characterized by potentials which fall exponentially with the dilaton field. At weak coupling, the Kahler potential for the dilaton is

$$K(S, S^\dagger) = -\ln(S + S^\dagger)$$

(5)

so the canonical field, $\phi$, is related to $S$ by

$$S = e^\phi.$$  

(6)

This means that the potential behaves for very weak coupling as:

$$V(S) \sim \exp(-A e^\phi)$$

(7)

i.e. it is an extremely steep function of $\phi$ for large $\phi$.  

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Figure 1: A typical racetrack potential (used in calculations below). Minimum is a small dimple.

In both of these proposals, the potential is already assumed to be quite small at the true minimum. So one has the situation illustrated in fig. The potential is extremely steep almost everywhere, with a tiny dimple at the minimum. It is natural to expect that the system overshoots. To get some feeling for the problem, first simply ignore the potential, and suppose that $\phi$ dominates the energy density of the universe. In this case $p = \rho$ so the scale factor grows as $R(t) \sim t^{1/3}$. $\phi$ obeys the equation:

$$\ddot{\phi} + \frac{1}{t} \dot{\phi} = 0,$$

with solution

$$\dot{\phi} = \frac{c}{t}, \quad \phi = \ln t + d.$$  \hfill (9)

If one plugs this into the expression for $V$, one sees that $V$ falls of exponentially fast with time, while the kinetic energy falls off only as a power. So the neglect of the potential is consistent. By the time the system reaches the minimum, the potential energy is a tiny fraction of the kinetic energy, and the system overshoots. Numerical study readily verifies that this is the case.
5 Modification of the Kahler Potential

The analysis above relied on the weak coupling form of the Kahler potential. In the case of Kahler stabilization, however, by definition the Kahler potential is significantly different from its weak coupling form near the minimum. Similarly, in the case of the racetrack scenario, one also expects significant corrections to the Kahler potential. So it is natural to ask whether one could avoid the Brustein-Steinhardt problem if the Kahler potential is significantly different from its weak coupling form. In this section we explore this possibility.

One way to analyze this problem is to note that, if the Brustein-Steinhardt problem is to be avoided, one should be in an approximately slow-roll regime. Let us first formulate the conditions for slow-roll in the context of a generic superpotential $W$ and Kahler potential $K$. For a Lagrangian of the form

$$L = f(\phi, \phi^\dagger) \partial_\mu \phi^\dagger \partial^\mu \phi - V(\phi, \phi^\dagger),$$

the condition for slow-roll reads

$$\frac{d}{d\phi} \left( \frac{1}{f} \frac{dV}{d\phi} \right) \ll 3V. \quad (11)$$

In the supergravity framework, we have

$$V = \left| \frac{\partial W}{\partial X} \right|^2 g_{X\bar{X}}^{-1} e^K, \quad f = K''.$$

Let us denote $\bar{V} = \left| \frac{\partial W}{\partial X} \right|^2$. We assume that it depends on the dilaton field $S$ and that $g_{X\bar{X}} = 1$. The condition (11) implies then the following three conditions for $\bar{V}$ and $K$:

$$\frac{\bar{V}''}{\bar{V}} \ll K'', \quad (13)$$

$$\frac{\bar{V}'}{\bar{V}} \ll \left( \frac{2K'}{K''} - \frac{2K'''}{(K'')^2} \right)^{-1}, \quad (14)$$

$$1 + \frac{(K')^2}{K''} - \frac{K'K'''}{(K'')^2} \ll 1. \quad (15)$$

Let us first investigate the case of the racetrack superpotential. Equations (13)-(14) can be rewritten then as conditions on the Kahler potential:

$$K'' \gg 1/N^2, \quad (16)$$
\[
\frac{(K'')^2}{2K'K'' - K'''} \gg \frac{1}{N}. \tag{17}
\]

At the minimum, \( S = \mathcal{O}(N) \). It is straightforward then to see that with the weak coupling form, \( K = -\ln(S + S^\dagger) \), the condition (13) is satisfied but the two conditions (14) and (17) are not. Allowing strong modifications of the Kähler potential, such as \( K = S^n \) for any \( n \geq 2 \) or \( K = e^{aS} \) with \( a \) positive or negative, we still cannot satisfy the three conditions simultaneously.

One can understand that the problem of satisfying (13), (16) and (17) simultaneously is generic by Taylor expanding around the minimum \( S_0 \). Take
\[
K = a_0 + a_1 (S - S_0) + \frac{1}{2} (S - S_0)^2 + \frac{1}{6} (S - S_0)^3. \tag{18}
\]
We are interested in finding a solution to the slow-roll condition for \( S \) that is within a few \( N \) away from \( S_0 \), that is, \( \delta \equiv (S - S_0)/N = \mathcal{O}(1) \). The \( a_m \) coefficients of eqn. (18) should then have the following \( N \) dependence:
\[
a_m = \alpha_m / N^m, \tag{19}
\]
where \( \alpha_m \) are \( N \)-independent. One can write the three conditions (13), (16) and (17) in terms of the three \( \alpha_m \) and \( \delta \). It becomes clear that fine tuning of the \( \alpha_m \) parameters is required to satisfy these conditions for a given \( \delta \). Even if we manage to find a solution for a given \( \delta = \mathcal{O}(1) \), moving a distance \( \Delta \delta = \mathcal{O}(1) \) away from this point will reintroduce strong violations of the slow-roll conditions.

We conclude that it is impossible to satisfy the slow-roll condition for a racetrack superpotential and any non-singular Kähler potential for a finite range of \( S \) that is a few \( N \) away from the minimum.

Second, we would like to investigate the case of Kähler stabilization. The superpotential is of the form
\[
W = e^{-aS}. \tag{20}
\]
Assume that \( K \) can be expanded around the minimum \( S_0 \) of the scalar potential \( V \) as
\[
K = K_0 + K'(S - S_0) + \frac{1}{2} K''(S - S_0)^2 + \frac{1}{6} K'''(S - S_0)^3. \tag{21}
\]
The scalar potential is then of the following form:
\[
V = e^K e^{-2aS} \left[ -a + K' + K''(S - S_0) + \frac{1}{2} K'''(S - S_0)^2 \right]^2 \frac{1}{K'' + K'''(S - S_0)} - 3 \right]. \tag{22}
\]
Requiring that $V(S_0) = 0$ and $V'(S_0) = 0$, we get

$$K'' = \frac{1}{3}(K' - a)^2, \quad (23)$$

$$K''' = \frac{2}{9}(K' - a)^3. \quad (24)$$

Thus, setting $K' - a$, we get $K''$ and $K'''$. Combining analytical and numerical searches, we were unable to find a solution of the slow-roll equations for $\delta \equiv (S - S_0)/N = \mathcal{O}(1)$ that is not fine-tuned.

We conclude that it is highly unlikely that there exists a form of the Kähler potential that would both stabilize the dilaton and satisfy the slow-roll conditions over a finite range of $(S - S_0) = \mathcal{O}(N)$.

6 A Possible Solution: The Role of Non-Zero Modes

In the previous section, we concluded that modification of the Kahler potential is unlikely to resolve the Brustein-Steinhardt problem. On the other hand, the authors of [18] made an observation which both sharpens the problem and is likely crucial to its resolution. They noted that the focus on zero momentum modes alone is inconsistent. Assuming the behavior $R \propto t^{1/3}$ described above, the equation for the non-zero momentum modes is:

$$\ddot{\phi} + \frac{1}{t} \dot{\phi} + \frac{k^2}{R^2} \phi = 0. \quad (25)$$

This equation has the solution

$$\phi = \frac{1}{t^{1/3}} \cos \left(\frac{3}{2}t^{2/3}\right). \quad (26)$$

As a result, the energy density of the non-zero modes falls off as $t^{4/3} \sim R^4$, i.e. more slowly than that of the zero modes (which falls off as $1/t^2$), and just like that of radiation.

We have little insight as to what might be appropriate initial conditions, but this result suggests some possibilities. One is to suppose that the field $\phi$ initially has a roughly thermal distribution, and dominates the energy density. This is consistent with the $1/R^4$ falloff (of course, the distribution does not have to be thermal to justify these statements; it is only necessary that the energy density be dominated by the non-zero modes of $\phi$). This provides additional damping, which, as we will see, can appreciably slow the motion of the field. Just
as important, it also means that there is an additional potential for the dilaton, a potential which can be much larger than that due to the non-perturbative effects associated with gluino condensation. The point is that the action for the moduli has coupling constant dependent corrections, i.e.

\[ \mathcal{L}_{\text{kin}} = (\partial_{\mu} \Phi)^2 (1 + a g^2 + b g^4 + \ldots). \] (27)

Averaging \((\partial_{\mu} \Phi)^2\) over the background, generates a potential for the dilaton:

\[ V(S) = \langle (\partial_{\mu} \Phi)^2 f(g^2) \rangle \] (28)

This potential vanishes for large \(S\) (small coupling). Its behavior at large coupling, where we expect \(S_0\) lies, is unknown, and we can imagine many possibilities.

To see that the assumption of a non-zero momentum background significantly alters the picture, first ignore the potential all together, as we did earlier. The zero mode now obeys the equation

\[ \ddot{\phi} + \frac{3}{2} \dot{\phi} = 0. \] (29)

This has solution

\[ \phi = \alpha t^{-1/2} + \beta. \] (30)

In other words, the field creeps to some particular point. Including the potential, it is reasonable to hope that the system will track the potential, and eventually settle into the correct minima.

Now consider the problem with the potential, including the large effects associated with the non-zero energy. Suppose that the potential tends to zero from below. As a model, take the energy to have the form

\[ V_{\text{eff}} = b g^2 + V_{np}(g^2). \] (31)

Here \(T\) is simply to be thought of as a parameter which characterizes the energy of the system; it is not necessary that the system be in thermal equilibrium, just that its energy redshift like radiation. If \(b < 0\), the potential has a minimum for non-zero coupling, resulting from the competition of the non-perturbative piece and the additional, finite energy contribution. If initially the coupling is not too small, the finite energy contribution dominates. As this energy redshifts, the minimum moves to weaker coupling; eventually, it will lie near \(S_0\). Provided \(b\) is
Figure 2: In the case that $b < 0$, the system settles nicely to its minimum. Not too small, the system tracks this minimum. This can be seen in fig. 2, where we have taken $b = -1$, and for $V_{np}$ we have taken

$$V_{np} = \left(e^{-\frac{sg^2}{16\pi^2}} - 4e^{-\frac{sg^2}{16\pi^2}}\right)^2$$

We have used the weak coupling form for the Kahler potential in studying the motion of the field. In the figure, one can see that the coupling evolves to its minimum, and then oscillates about it, as expected.

In the case that $b > 0$, it is more difficult for the system to find its minimum. We find for this simple model that if $b > 0.02$, the system typically overshoots the minimum. If $b$ is smaller, however, the damping is adequate, and the system, as in fig. 2, settles into the minimum near $S_o$ (in this case located at $S = 1.91$). This is compatible with crude estimates obtained by using the slow roll approximation.

A second possibility is that the system really is in thermal equilibrium. In particular, the hidden sector gauge fields might be in thermal equilibrium. We need to ask what is the effective action for $S$ in a thermal background. But this is easy: the effective potential is precisely the free energy as a function of coupling. To understand this, note that it is appropriate to think of the dilaton as changing adiabatically on the scale $T^{-1}$. As a result, we can integrate out the
Figure 3: For small, positive $b$, the system also settles to its minimum. Here $b = 0.02$.

fast modes – the gauge bosons, gauginos, etc., to obtain an effective action for $S$. The potential just corresponds to setting $S = 1/g^2$ to a constant value, and computing the free energy of the gauge system. In other words, the potential for $S$ is just the free energy of the gauge system.\footnote{This is different than the results of \cite{24}, who argue that the potential behaves as $1/g^2$, i.e. that it blows up at weak coupling. Their form is correct for coupling to certain homogeneous scalar figure configurations. Many of our other observations are similar to theirs.}

The form of the free energy for such a system is known. For gauge group $SU(N)$ (without matter fields, for simplicity) one has, for example\cite{26}

\[
\Omega(g, T) = -\frac{(N^2 - 1)\pi^2}{24}T^4(1 - \frac{3}{8}\frac{g^2N}{\pi^2} + \frac{g^3N^{3/2}}{\sqrt{2}\pi^3} + \ldots).
\]  \hspace{1cm} (33)

As expected, it tends to zero for small coupling. This formula does not exhibit a minimum for larger coupling. However, for $N = 10$ or so, as expected in the racetrack picture, the perturbation expansion has broken down at interesting values of the coupling ($g \sim 1/2 - 1$), and we might speculate on possible behaviors.

The system might have no minimum at all, except at zero coupling. In this case, the problem is similar to that we considered above, with a potential proportional to $g^2$ with a positive coefficient. If one takes the perturbative formula literally, this is roughly like $b \sim 1/3,$
for \( N = 10 \), and so the system is likely to overshoot the minimum. On the other hand, given that the corrections are large, one might be lucky, and the coefficient, effectively, might be much smaller, giving the behavior found for small \( b \) above. Alternatively, the system might have a local minimum at some coupling, \( g_0 \). (Note that in any case, the high temperature calculation is not valid for \( T < \Lambda(g) \), at which point the low energy non-perturbative potential takes over.)

In this case, one might expect that the system will roughly track this local minimum, if it starts out to the left (at stronger coupling), essentially following adiabatically. In this case, one might hope that the system will be gently set in the true minimum. We have written toy potentials with these features, and verified that indeed this does happen.

These results suggest that the extent to which the Brustein-Steinhardt problem is a problem is quite sensitive to dynamics which are not known as present, as well as aspects of the initial conditions of the universe, which are not well understood.

7 Conclusions

There has been much discussion through the years of the moduli problems of string cosmology. These problems have often been viewed as providing a serious challenge to the viability of string theory itself. The main lesson of the studies here is that understanding these problems requires both a much better understanding of moduli stabilization and of the initial conditions of early universe cosmology than we have at present. With quite plausible assumptions about how moduli are stabilized, the moduli need not dominate the energy density of the universe. With equally plausible assumptions about the conditions of the early universe, the Brustein-Steinhardt problem is readily avoided.

We have not offered a complete history of the early universe. We have not committed ourselves, for example, to the question of whether we are describing a period before or after inflation, in part because we suspect that moduli themselves may have something to do with inflation. There are many issues which one would have to consider in a complete model. For example, the authors of [8] have argued that production of non-zero momentum modes of moduli during inflation, and that this might even endanger the enhanced symmetry picture we have advocated here. On the other hand, as they note, provided the effective masses of the moduli are of order \( H \) during inflation, the production is modest. Similarly, given the large mass of the dilaton, it is likely to decay very early, avoiding the usual moduli problem. It is difficult to address these questions without a detailed model of inflation. Unfortunately, we
have not yet seen a simple way to fit inflation into this picture, without fine tuning. From the results presented here, however, we hope to have made clear that the moduli problems, as usually formulated, have plausible, robust solutions.

We know that non-perturbatively there are many string ground states with unbroken supersymmetry. It seems likely that, if there are stable string ground states with broken supersymmetry and vanishing cosmological constant, there are many of them. The results reported here suggest that, as issues involving moduli stabilization are better understood, cosmology will likely be crucial to understanding to understanding how the universe finds itself in the state we see – and not in one of the myriad other possibilities.

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