Enhanced Symmetries and the Ground State of String Theory

Michael Dine

Santa Cruz Institute for Particle Physics
University of California, Santa Cruz, CA 95064
dine@scipp.ucsc.edu

Yosef Nir and Yael Shadmi

Department of Particle Physics
Weizmann Institute of Science, Rehovot 76100, Israel
ftnir@wicc.weizmann.ac.il, yshadmi@wicc.weizmann.ac.il

The ground state of string theory may lie at a point of “maximally enhanced symmetry”, at which all of the moduli transform under continuous or discrete symmetries. This hypothesis, along with the hypotheses that the theory at high energies has $N = 1$ supersymmetry and that the gauge couplings are weak and unified, has definite consequences for low energy physics. We describe these, and offer some suggestions as to how these assumptions might be compatible.
1. The Problem of Moduli

Moduli, exact and approximate, seem almost ubiquitous in string ($M$) theory. In recent years, they have proven quite useful in elucidating the underlying structure of the theory. But they also pose one of the greatest challenges to the theory. If string theory describes nature, how are they fixed?

The aim of this note is to enumerate, in a very general way, mechanisms by which this problem might be solved, and to explore their consequences for low energy physics. To study this problem, we will adopt a particular framework. We will assume, first, that the theory has a fundamental energy scale comparable to the Planck mass, $M_p$. (This assumption can be relaxed, but is convenient for the initial discussion). Slightly below this scale, we will assume that the theory has $N = 1$ supersymmetry, with some spectrum of light particles, including the gauge and matter fields of the MSSM. We will assume that the standard model gauge couplings at this scale are small and unified, but that there is no weak coupling description of the full theory. Finally, we will assume that supersymmetry breaking is hierarchically small, i.e., it is of order $e^{-8\pi^2/\alpha g^2}$ for some constant $a$.

These are plausible hypotheses, but of course they need not be true. Adopting them, however, we can quickly narrow down possible solutions to the moduli problem. First, a definition: we will use the phrase “true vacuum of string theory”, to describe the putative state of the theory which corresponds to the world we observe. This does not mean that there are not many other vacuum states; for example, we expect there are states with $N = 2, 4, 8$ supersymmetry in four dimensions which are exact non-perturbative solutions of the theory, not to mention higher dimensional ground states with varying degrees of supersymmetry.

There are a number of possibilities one can imagine for the true vacuum. It could be that this state contains approximate moduli. By this we mean that there are fields with potentials that tend to zero in some directions. This definition includes states which have a weakly coupled string theory limit or an $M$-theory limit. Virtually all ideas about the moduli problem which have been discussed to date involve approximate moduli.

Alternatively, it could be that in the true vacuum there are no approximate moduli. One proposal along these lines was made in [1], though that particular example did not
have any supersymmetry. If models exist without approximate moduli and unbroken supersymmetry at a high scale, supersymmetry must be broken by low energy field theory dynamics. To understand why this is so, one merely need note that (small) susy breaking by high energy string effects must appear in the low energy lagrangian through a linear superpotential for some light fields; these fields are necessarily singlets under all of the symmetries. But by assumption, there are no moduli whose $F$-terms can break supersymmetry. While it may be risky to apply conventional notions of naturalness in string theory, it seems highly unlikely that there are light fields which are singlets under all (gauge and discrete) symmetries. So high energy effects cannot generate a supersymmetry breaking superpotential.

Supersymmetry breaking, then, should be a phenomenon of the low energy field theory. In other words, the low energy field theory must contain a set of fields and interactions which break supersymmetry. The scale of this breaking is of order $M e^{-8\pi^2/bg^2}$, where $g$ is the coupling of this theory at the scale $M$ and $b$ is a constant. According to our hypotheses this coupling is small; note that it is not the expectation value of a dynamical field. It is important to stress that we are not assuming that the string coupling is small or that the complete theory is described by perturbation theory; only that there is an effective low energy theory described by weak gauge couplings. The plausibility of this assumption will be discussed further below. Because there are no singlet fields, one cannot generate gaugino masses through non-renormalizable couplings to $W^2_\alpha$, as in conventional supergravity models. So it is probably necessary to generate gaugino masses by gauge mediation. Our hypotheses, then, lead almost inevitably to a picture of dynamical (field theory) supersymmetry breaking and gauge mediation, with the supersymmetry breaking scale anywhere between tens of TeV, as in the models of \cite{6} \cite{7}, and much higher scales, as in “hybrid models” first discussed in \cite{8}.

Let us turn now to the question of theories with approximate moduli. It is usually

\footnote{Such models presumably must be discovered in some self-consistent scheme, since by assumption supersymmetry is broken (albeit by a small amount) and there is no small parameter. However, it seems quite possible that such theories exist.}

\footnote{Some related observations about theories without moduli appear in \cite{2}.}
assumed that the vacuum lies at some random point in the moduli space. But work on string dualities in the last few years suggests a more compelling possibility. We now know that there are often points in the moduli space with “maximally enhanced symmetry”, i.e., points where all of the moduli are charged under some symmetry (which may be continuous or discrete). For example, compactifying the heterotic string on $T_6$, there are many points where all of the geometrical moduli are charged under enhanced gauge symmetries. Because these compactifications exhibit $S$-duality, at the self-dual point one has a $Z_2$ symmetry under which $S$ transforms.

This example is unrealistic, since it has $N = 4$ supersymmetry, and $\frac{g_4^2}{4\pi} = 1$ at the self-dual point, which contradicts our hypotheses. But we now know of vast numbers of string compactifications with high degrees of symmetry. It is plausible that among these are states with maximally enhanced symmetry and small, unified gauge couplings. In general, maximally enhanced symmetry points are of interest for at least two reasons:

1. They are automatically stationary points of the effective action, and thus candidate minima. Tadpoles for the moduli are forbidden; it is a detailed question whether all of the moduli have positive masses.

2. They don’t suffer from the cosmological moduli problem\[5\]. In thinking about string cosmology, it is usually assumed that the minimum of the moduli potential lies at a random point in the moduli space, i.e., a point with no special properties. The universe is unlikely to start out at such a point. Finite temperature and/or curvature effects (say during inflation) are likely to leave the universe somewhere else. Assuming the moduli potential is of order the scale of supersymmetry breaking, the universe remains at this point until the Hubble parameter is of order the curvature of the potential, after which the system oscillates. The moduli typically dominate the energy density of the universe when they decay, leading to catastrophic consequences. However, if the vacuum is a point of maximally enhanced symmetry, it is quite natural for the universe to start out in this state \[9\]. For example, during inflation, even though the potential for the moduli is modified, it quite naturally may prefer the symmetric state. Finite temperature effects also tend to prefer states of higher symmetry. But even without any specific picture of the universe, since maximally enhanced symmetry points are
special, it is quite plausible that they will be singled out in the early universe.

Let us first explore the consequences of the assumption that the true vacuum lies at an enhanced symmetry point. Again, high energy string effects cannot be responsible for supersymmetry breaking, since the symmetries forbid any linear term in the moduli superpotential. In addition, gluino condensation cannot yield supersymmetry breaking in such a vacuum, since couplings $\mathcal{M}W_{\alpha}^2$ are forbidden. So the situation is very similar to that in theories without moduli; even if the moduli acquire substantial $F$-components, because of the symmetries they cannot couple to $W_{\alpha}^2$, and thus cannot contribute to gaugino masses. Note that field theories with dynamical supersymmetry breaking similarly do not contain fields that are singlets under all symmetries. So, again, we are led to a picture of dynamical supersymmetry breaking in the effective low-energy field theory, almost surely involving some sort of hidden sector and gauge mediation.

How would the moduli be fixed in such a picture? One can imagine several possibilities. The first question, though, is to identify the moduli. At the renormalizable level, the MSSM possesses several flat directions. Some of these could, in fact, be exact at the level of the superpotential. For example, there is a set of flat directions labeled by chiral fields $\bar{u}\bar{d}\bar{d}$. These would be exactly flat if one had a (discrete) $R$ symmetry under which the $\bar{u}$, $\bar{d}$ and $\bar{e}$ fields were neutral. It is also possible to find flat directions for which all of the standard model gauge symmetries are broken. (One such direction has non-zero vev’s for the $Q$ and $L$ fields, and discrete symmetries; these directions are described by 12 parameters, i.e., they could correspond to as many as 12 moduli.) So it could be that the moduli are in fact linear combinations of the ordinary squarks, sleptons and Higgs fields. The enhanced symmetries might just be the symmetries of the standard model!

If this is the case, it is easy to imagine how the moduli are fixed. In gauge mediation, for example, the squarks and sleptons gain mass from Kähler potential corrections; for a sensible model, the curvature of the potential must be positive at the origin, i.e., at the enhanced symmetry point. This is, indeed, a realization of the “Kähler stabilization” of [12].

---

3 Models such as those of Intriligator and Thomas [10] or Nelson [11] are only natural if the various singlets transform under discrete symmetries.
Alternatively, some or all of the moduli might play a role in the hidden sector dynamics. The models of \([10]\) and \([11]\), for example, have classical moduli. These moduli may be stabilized near the origin, i.e., near the enhanced symmetry point. In this case, the masses of the moduli could easily be 10’s of TeV or larger.

Finally, some or all of the moduli might be singlets both with respect to the standard model gauge group and with respect to any hidden sector gauge group. Assuming that supersymmetry is broken by the \(F\)-components of some hidden sector fields, \(Z\), these fields could gain mass through Kähler potential terms such as

\[
K = \mathcal{M}^\dagger \mathcal{M} Z^\dagger Z .
\] (1.1)

Such couplings respect the symmetries, and if the sign is suitable, stabilize the moduli at the symmetry point. The moduli fields would be quite light. However, because of the symmetries, their interactions with ordinary matter would be highly suppressed. For example, typical axion or Brans-Dicke couplings would be forbidden. Couplings to quarks and leptons would also be suppressed by extra powers of the large mass scale, the amount of suppression depending on the charges of the quarks and leptons.

2. The Cosmological Constant Problem

One obvious question in this picture is: can the cosmological constant vanish? We have no new wisdom to offer as to the detailed mechanisms by which the cosmological constant vanishes. In particular, we do not know whether the vacuum energy should vanish or nearly vanish in the effective action. Still, we can give conditions under which cancellation of the vacuum energy at the level of the effective action is at least possible. In conventional supergravity model building, one starts with the potential

\[
V = e^K \left[ \left( \frac{\partial W}{\partial \phi_i} + \frac{\partial K}{\partial \phi_i} \right) K^{ij} \left( \frac{\partial W^*}{\partial \phi_j^*} + \frac{\partial K}{\partial \phi_j^*} W^* \right) - 3|W|^2 \right] .
\] (2.1)

One then notes that the superpotential can include a constant, and this constant is adjusted to give zero potential at the minimum.

We wish, then, to ask whether there can be a constant of the correct order of magnitude, in the context of models with maximally enhanced symmetry. Consider, first,
theories with no moduli. In such theories, one would expect that $W \sim 1$, unless there is a symmetry. This is presumably a disaster, but such a term can be forbidden by a discrete $R$ symmetry. In order that there be a small constant term in $W$, one then requires a spontaneous breaking of the discrete symmetry in the low energy theory. This can occur, for example, if there is a sector containing a pure gauge theory without matter. In the absence of moduli, such a sector can generate a non-zero, constant value for $W$. This can be of the correct order of magnitude (this depends on the $\beta$-functions of the various sectors).

In the case where there are approximate moduli, there are similar issues. It is possible, as we will describe below, that $e^{-S}$ is numerically small. Alternatively, it may be forbidden by discrete symmetries. However, in either case, obtaining a constant of the correct order of magnitude may require a sector like that described above to generate the necessary constant. Of course, in any of these pictures, the precise cancellation of the cosmological constant looks miraculous from the perspective of the effective field theory.

### 3. Why are the couplings small and unified?

One of the great puzzles of string theory is: why are the gauge couplings we observe small and unified, if the underlying theory is strongly coupled. The weakly coupled heterotic string provides a simple picture of coupling unification, but it is not clear why this should hold at stronger coupling, when $e^{-S}$ cannot be neglected. The picture we have developed of maximally enhanced symmetries provides, perhaps, a plausible answer. At such points, the gauge couplings will be pure numbers, independent of any moduli, determined, perhaps, by topological considerations. One can certainly conjecture that these numbers are sometimes small, and occasionally equal. This seems more plausible than equality in theories where the string coupling is large but the moduli are simply random.

Alternatively, in [12], a scenario for coupling unification at small gauge coupling was proposed, along with a mechanism for stabilizing the dilaton. It was argued that perhaps $e^{-S}$ is small, so there are no large corrections to the superpotential, $W$, and the gauge coupling function, $f$, but that corrections to the Kähler potential might be large, and that
these might be responsible for the stabilization of the dilaton. The point is that while one expects that string perturbation theory is less convergent than the perturbation expansion of field theory \[13\], holomorphy severely restricts the form of the corrections to \( W \) and \( f \).

We can ask whether we can account for weak coupling and unification in a similar way in the case of maximally enhanced symmetry. In the simple example of maximally enhanced symmetry discussed in the introduction, \( \frac{g^2}{4\pi} \) is equal to one. Indeed, we might expect this behavior is generic. Consider starting with some weak coupling description, and varying the dilaton. At weak coupling, the superpotential is independent of \( S \), up to terms of order \( e^{-S} \). However, if a symmetry is to be restored, we would expect that the corrections to the superpotential must be large, so that couplings which violate the symmetry vanish, and perhaps that massive particles become massless. This requires that \( e^{-S} \) not be negligible at the enhanced symmetry point. In such a view, no matter what the numerical value of the coupling, there would be no sense in which one could treat \( e^{-\frac{g^2 S}{2}} \) as small in the full string theory, and coupling unification would have to be a consequence of a particular choice of strongly coupled ground state.

The possible loophole to this argument is, again, the Kähler potential. Could the Kähler potential be responsible for symmetry restoration? If this is to be the case, \( K \) presumably must be singular at the symmetry point. While we have no theory where we can currently exhibit such an effect in a meaningful calculation, in the following section we offer a toy model for this phenomenon.

### 4. A Model for Symmetry Enhancement Through Kähler Potential Effects

In \[12\], in a weak coupling language, it was suggested that perhaps the moduli are stabilized in a region where corrections to the classical Kähler potential are large, but \( e^{-S} \) is very small. Based on the argument above, we would not expect enhanced symmetries at such a point, unless the Kähler potential were singular. The question is: what sort of singularities are required and how plausible is such singular behavior for the Kähler potential?

\[4\text{ Ref. [12] actually argued in the other direction: that symmetries which are unbroken at weak coupling can only be broken at stronger coupling if the Kähler potential becomes singular.}\]
Ref.\textsuperscript{14} discussed a possibility for stabilizing the dilaton involving Fayet-Iliopoulos terms. It was argued that if the $D$-term vanishes somewhere on the moduli space, it is natural that the potential has a minimum nearby. A little thought indicates that a vanishing $D$-term also provides a model for enhanced symmetries. In models where a gauge anomaly is cancelled by the axion, there is a Fayet-Iliopoulos $D$-term, which, for general points on the moduli space, is given by

$$\xi^2 = -\frac{\delta_{GS}}{2} \frac{\partial K}{\partial S}, \tag{4.1}$$

where

$$\delta_{GS} = \frac{1}{192\pi^2} \sum_i q_i. \tag{4.2}$$

At generic points in the moduli space, the Green-Schwarz term is non-zero; typically the $D$-term is cancelled by the expectation value of some scalar field(s), $N$. This scalar field expectation value may break discrete or continuous symmetries in addition to the anomalous $U(1)$. Now if the derivative were to vanish at some point in the moduli space, the expectation value would vanish, and symmetries might be restored. This symmetry restoration is controlled by a non-holomorphic quantity. So there is the possibility of enhanced symmetry, and small holomorphic coupling.

But, as presented, there is a problem with this idea. In order that the derivative of the Kähler potential vanish, it is necessary that the second derivative, $\frac{\partial^2 K}{\partial S^2}$, change sign. This means, however, that the kinetic term for the dilaton does not have the correct sign. In \textsuperscript{14}, it was asserted that this would not be a problem once kinetic mixing with other moduli was considered. But this is not the case; it is easy to show that for a general symmetric matrix, if a diagonal component of the matrix is negative, at least one eigenvalue must be negative. So this description of the problem cannot be correct.

However, there is no reason why the description of the theory in terms of various light fields, the dilaton multiplet, and the massive $U(1)$ gauge multiplet, should make sense for such values of the coupling. At such points, the gauge multiplet has mass of order the other heavy fields in the theory. It also has a very large width. So it is not clear that we should be keeping only this set of degrees of freedom. Instead, we should focus on the
light degrees of freedom, and ask what features of the Kähler potential are responsible for symmetry restoration.

It will prove instructive to first consider the $D$-term model at weak coupling. Loosely speaking, this limit already provides us with an example of symmetry restoration through Kähler potential effects. Let us suppose that there is a single field, $N$, with charge $q_N < 0$ (we take $\delta_{GS} > 0$), which cancels the $D$-term (in particular, the direction with non-zero $N$ is “$F$-flat”). The weakly coupled theory is described by a Kähler potential

$$K = f(S + S^\dagger - \delta_{GS} V) + N^\dagger e^{2q_N V} N + \sum \phi_i^\dagger e^{2q_i V} \phi_i + \ldots ,$$

(4.3)

where $\phi_i$ are light, chiral fields with $U(1)$ charge $q_i$, which, we assume, do not develop vevs. The $D$-term is given in eqn. (4.1), so

$$N^\dagger N = \frac{\delta_{GS}}{2q_N} \frac{\partial K}{\partial S} .$$

(4.4)

Since the $N$ field is massive (the light modulus at weak coupling is principally $S$), it is appropriate to integrate it out in this region. The theory will also contain, in general, superpotential terms such as (suppressing couplings)

$$W = N\phi_i \phi_j + N\phi_i \phi_j \phi_k + \ldots .$$

(4.5)

The expectation value of $N$ then leads to masses, Yukawa couplings, and so on in the low energy theory. Because the $N$ expectation value is a non-holomorphic function of the dilaton, so are these couplings. Moreover, the $N$ contribution to the masses leads to non-holomorphic corrections to the gauge couplings. Taking $S \to \infty$, the $N$-vev tends to zero, so that various masses and couplings in eqn. (4.3) approach zero as well. Thus, even at weak coupling we have a model of “symmetry restoration”. In the low energy theory, it must be possible to interpret the $S$-dependence of the couplings in terms of the Kähler potential for the light fields.

In order to understand the structure of the low energy theory, first note that the heavy scalar field, which is eaten by the $U(1)$ gauge field, is of the form, $H \sim \delta_{GS} K'' S - 2q_N |\langle N \rangle|^2 \ln N$, and is predominantly $N$, while the light field, which parametrizes the flat
direction, is the gauge-invariant combination $S + \frac{\delta_{GS}}{2qN} \ln N$, which, at weak coupling, is approximately $S$. It is therefore convenient to choose a gauge in which $N$ is fixed, $N = n_0$, for some fixed reference $n_0$. Note that this is a holomorphic gauge condition. In this gauge the Kähler potential can be written as

$$K = f(\Phi + \Phi^\dagger - \delta_{GS}X) + |n_0|^2 e^{2qN X} + \sum \phi_i^\dagger e^{2q_i X} \phi_i + \ldots , \quad (4.6)$$

where $\Phi = S + \frac{\delta_{GS}}{2qN} \ln \frac{N}{n_0}$. Integrating out the heavy field $X$ we have

$$2qN |n_0|^2 e^{2qN X} - \delta_{GS} K'(\Phi + \Phi^\dagger - \delta_{GS}X) = 0 , \quad (4.7)$$

so that, for large $\Phi$, $X \sim \frac{1}{2qN |n_0|^2} \ln(\frac{\delta_{GS}}{2qN} \frac{1}{\Phi + \Phi^\dagger})$. At low energy we then have the fields $\Phi, \phi_i$ with the Kähler potential

$$K = - \ln(\Phi + \Phi^\dagger) + \sum Z_i \phi_i^\dagger \phi_i + \ldots , \quad (4.8)$$

with

$$Z_i = \left( \frac{\delta_{GS}}{2qN} \frac{1}{|n_0|^2} \frac{1}{\Phi + \Phi^\dagger} \right)^{\frac{q_i}{2qN}} = \frac{|\langle N \rangle|^{\frac{q_i}{2qN}}}{|n_0|^{\frac{q_i}{2qN}}} . \quad (4.9)$$

From the point of view of this low energy theory, (which does not have $N$ in the superpotential), symmetry restoration occurs since the Kähler potential becomes singular as $\Phi \to \infty$. Note that the singularity is just a consequence of our insistence of writing the theory in a holomorphic fashion. Physical quantities are not singular.

The question, then, is whether such singularities can occur on the “interior” of the moduli space, as we vary the gauge coupling. We have seen that the notion of a vanishing $D$-term at an interior point is, at best, a caricature of what might be going on; because of the change of sign of the kinetic term, such a description cannot be complete. Still, it is a suggestive model. Suppose, for a moment, that it is a suitable description near the enhanced symmetry point. We could imagine repeating the previous exercise near the point $K' = 0$. In order to smoothly interpolate between the weak and strong coupling regions, it is again useful to choose the gauge $N = n_0$. From a physical point of view, this is a rather peculiar choice, however, as the modulus, $\Phi$, in this regime is predominantly $N$. It is helpful, indeed, to consider first the gauge $S = s_0$, and then gauge transform to
the gauge with constant $N$, through the gauge transformation $i\Lambda = \ln(N/n_0)$. Proceeding in this way, near the enhanced symmetry point $\langle N \rangle = 0$,

$$Z_i = \left(\frac{N}{n_0}\right)^{q_i/qN} \left(\frac{N}{n_0}\right)^{\dagger q_i/qN} = \left(e^{\Phi-\Phi_0}\right)^{q_i} e^{\Phi-\Phi_0} \right)^{2q_i} \frac{2q_i}{\gamma_{GS}}, \quad (4.10)$$

where $\Phi_0$ is just the constant $s_0$. As expected, the $Z_i$ become singular as $N \to 0$ (or $\Phi \to \infty$), leading to symmetry restoration. Note that in this gauge, at the enhanced symmetry point, the $Z$’s are products of a holomorphic and an antiholomorphic factor and so we can make a holomorphic redefinition of the fields to remove the $Z_i$’s from the Kähler potential. This redefinition puts $N$, now interpreted as a light, dynamical field, in the appropriate places in the superpotential. Indeed, this redefinition is just the gauge transformation in reverse, so we recover the $S = s_\sigma$-gauge description. One should point out that in this model, there are not necessarily new massless particles at the enhanced symmetry point. If there are, one can contemplate more complicated singularities.

It is interesting to note in this description how the gauge coupling function respects the symmetry. After all, one of the main points in our discussion of enhanced symmetries is that there cannot be couplings of the charged modulus to the gauge fields, or linear terms in the charged modulus in the superpotential. In the gauge $N = n_0$, the gauge coupling function is equal to $\Phi$, which does transform under the symmetry. However, the redefinitions of the $\phi_i$’s, required to make their kinetic terms canonical, precisely cancel the $N$-dependent term in $\Phi$, leaving $s_0$ in front of the gauge kinetic terms.

We can thus describe the phenomena which occur in this model entirely in terms of the Kähler potential of the light fields. The crucial features are that at the enhanced symmetry point, the Kähler potential for the modulus, $\Phi$, behaves as

$$K_\Phi = |n_0|^2 \left(\left(e^{\Phi-\Phi_0}\right)^{\dagger} e^{\Phi-\Phi_0}\right)^{\frac{2q_N}{\gamma_{GS}}}, \quad (4.11)$$

and $Z_i(\Phi)$ is given in eqn. (4.10), whereas at weak coupling $K_\Phi$ and $Z_i(\Phi)$ are as given in eqns. (4.8)-(4.9). This example makes clear that Kähler potential effects can, in principle, restore a symmetry. While phrasing the discussion in terms of the light modulus $\Phi$ seems to avoid the issue of $K'' < 0$, it should be noted that the Kähler potential for $\Phi$ is not
single-valued; $\Phi \to \infty$ both in the weak-coupling limit and in the limit that the symmetry is exactly restored. We do not know how to assess the plausibility of such behavior of the Kähler potential in real string theory.

Another toy example of symmetry restoration through Kähler effects, that does not involve a negative kinetic term, has two anomalous $U(1)$’s. Several anomalous $U(1)$’s can appear in non-perturbative heterotic vacua (or in perturbative type I vacua) where there can be more than one dilaton field to cancel the $U(1)$ anomalies [15] [16]. Consider then a model in which the Fayet-Iliopoulos terms are canceled, at generic points, by the vevs of two fields, $N_1$ and $N_2$, of charges $(q_a^a, q_b^b)$, $\alpha = 1, 2$, under $U(1)_a, U(1)_b$. It is easy to see that the $N_1$ vev vanishes for

$$q_2^b \delta_{GS}^a \frac{\partial K}{\partial S_a} = q_2^b \delta_{GS}^b \frac{\partial K}{\partial S_b}. \quad (4.12)$$

If $N_1$ transforms under an additional symmetry, this symmetry is restored along the curve (4.12). Note that such a curve exists even if $K$ is perturbative. (We do not know whether this is realized in actual $F$-theory examples.) Thus there may be a region of moduli space that contains curves of enhanced symmetry, but no negative kinetic terms or branched behavior. Here again, in the low energy theory, the enhanced symmetry is associated with singular wave-function renormalizations for some of the light fields. This model, however, does not lead to maximal symmetry restoration along $N_1 = 0, N_2 \neq 0$.

5. Conclusions

In considering the problem of supersymmetry breaking in string theory, it has almost invariably been assumed that the true string vacuum contains approximate moduli, and that they are fixed at some random value. In this note, we have considered two alternative possibilities: either the vacuum has no approximate moduli, or the moduli sit at or very near a point of maximally enhanced symmetry. In both cases, we have argued that supersymmetry breaking must occur in the low-energy field theory, and that superpartner masses are generated, at least in part, by gauge mediation.

In the latter case, we have pointed out that the moduli might actually be the superpartners of ordinary quarks and leptons. The same mechanisms which give them mass
(e.g. gauge mediation) then also fix their expectation values at the symmetry point. Alternatively, the moduli might be very light, and extremely weakly coupled, or they might be fields which participate in a supersymmetry-breaking hidden sector.

We have noted that the moduli problem of cosmology is readily solved in such a framework. This is simply because the zero energy ground state is a special point. Finite temperature or curvature effects can naturally favor such enhanced symmetry points.

We have not seen a clear, low energy signal which distinguishes between the no moduli/approximate moduli alternatives. Since the MSSM contains approximate flat directions, only measurements of high dimension operators determine whether the directions are truly flat. One could imagine, however, a string theory program in which one would enumerate, e.g. $F$-theory compactifications with suitable gauge groups, particle content, symmetries and couplings [17]. Choosing such a state would then determine the low energy gauge group and the detailed pattern of supersymmetry breaking.

It should be noted that it is possible to relax a number of our assumptions. For example, one could imagine that there are a variety of scales, as in [18] [19] [20] [21]. In these cases, one might want to understand not only why the gauge couplings are small, but also why certain “geometrical moduli” are also large or small. This might be relevant to understanding flavor structure (as might be the enhanced symmetries).

In any case, we believe that these ideas establish a framework for considering the problem of the moduli, which can be pursued both at the level of phenomenological model building and of more fundamental string theory.

Acknowledgements

We thank David Kutasov, Jan Louis, Eva Silverstein, Nathan Seiberg, Adam Schwimmer and Scott Thomas for conversations. We are especially grateful to Tom Banks for critical comments on an early version of this manuscript. This work was supported in part by the United States – Israel Binational Science Foundation (BSF). The work of M.D. was supported in part by the U.S. Department of Energy. Y.N. is supported in part by the Israel Science Foundation and by the Minerva Foundation (Munich).
References

[1] M. Dine and E. Silverstein, “New M Theory Vacua With Frozen Moduli,” hep-th/9712166.
[2] C.P. Burgess, A. de la Macorra, I. Maksymyk and F. Quevedo, “Constant Versus Field Dependent Gauge Couplings in Supersymmetric Theories,” hep-ph/9805285.
[3] E. Poppitz and S. Trivedi, “Dynamical Supersymmetry Breaking,” hep-th/9803104.
[4] I. Affleck, M. Dine and N. Seiberg, “Dynamical Supersymmetry Breaking and Its Phenomenological Implications,” Nucl. Phys. B256 (1985) 557.
[5] T. Banks, D. Kaplan and A.E. Nelson, “Cosmological Implications of Dynamical Supersymmetry Breaking,” Phys. Rev. D49 (1994) 779.
[6] M. Dine, A. Nelson, and Y. Shirman, “Dynamical Supersymmetry Breaking Simplified,” Phys. Rev. D51 (1995) 1362, hep-ph/9408384.
[7] M. Dine, A. Nelson, Y. Nir, and Y. Shirman, “New Tools for Low Energy Dynamical Supersymmetry Breaking,” Phys. Rev. D53 (1996) 2658, hep-ph/9507378.
[8] E. Poppitz and S.P. Trivedi, “New Models of Gauge and Gravity Mediated Supersymmetry Breaking,” Phys. Rev. D55 (1997) 5508, hep-ph/9609529.
[9] M. Dine, L. Randall and S. Thomas, “Baryogenesis From Flat Directions of the Supersymmetric Standard Model,” Nucl. Phys. B458 (1996) 291, hep-ph/9507453.
[10] K. Intriligator and S. Thomas, “Dynamical Supersymmetry Breaking on Quantum Moduli Spaces,” Nucl. Phys. B473 (1996) 121, hep-th/9603158.
[11] A.E. Nelson, “A Viable Model of Dynamical of Supersymmetry Breaking in the Hidden Sector,” Phys. Lett. B369 (1996) 277, hep-ph/9511350.
[12] T. Banks and M. Dine, “Coping with Strongly Coupled String Theory,” Phys. Rev. D50 (1994) 7454, hep-th/9406132.
[13] S. Shenker, talk at Cargese workshop on Random Surfaces, Quantum Gravity and Strings, France (1990).
[14] T. Banks and M. Dine, “Quantum Moduli Spaces of N=1 String Theories,” Phys.Rev. D53 (1996) 5790, hep-th/9508071.
[15] G. Aldazabal, A. Font, L.E. Ibanez and G. Violero, “D=4, N=1, Type IIB Orientifolds,” hep-th/9804026.
[16] L.E. Ibanez, “New Perspectives in String Phenomenology from Dualities,” hep-ph/9804236.
[17] V. Kaplunovsky and J. Louis, “Phenomenological Aspects of F-Theory,” Phys. Lett. B417 (1998) 45, hep-th/9708049.
[18] E. Witten, “Strong Coupling Expansion of Calabi-Yau Compactification,” Nucl. Phys. B471 (1996) 135, hep-th/9602070.
[19] T. Banks and M. Dine, “Couplings and Scales in Strongly Coupled String Theory,” Nucl. Phys. B479 (1996) 173, hep-th/9605136.
[20] P. Horava, “Gluino Condensation in Strongly Coupled Heterotic String Theory,” Phys. Rev. D\textbf{54} (1996) 7561, \texttt{hep-th/9608019}.

[21] N. Arkani-Hamed, S. Dimopoulos and G. Dvali, “The Hierarchy Problem and New Dimensions at a Millimeter,” \texttt{hep-ph/9803315}.