String Universality

R. Brustein(1) and S. P. de Alwis(2)

(1) Department of Physics, Ben-Gurion University, Beer-Sheva 84105, Israel
(2) Department of Physics, Box 390, University of Colorado, Boulder, CO 80309.
e-mail: (1) ramyb@bgumail.bgu.ac.il (2) dealwis@pizero.colorado.edu

(January 2000)

Abstract

If there is a single underlying “theory of everything” which in some limits of its “moduli space” reduces to the five weakly coupled string theories in 10D, and 11D SUGRA, then it is possible that all six of them have some common domain of validity and that they are in the same universality class, in the sense that the 4D low energy physics of the different theories is the same. We call this notion String Universality. This suggests that the true vacuum of string theory is in a region of moduli space equally far (in some sense) from all perturbative theories, most likely around the self-dual point with respect to duality symmetries connecting them. We estimate stringy non-perturbative effects from wrapped brane instantons in each perturbative theory, show how they are related by dualities, and argue that they are likely to lead to moduli stabilization only around the self-dual point. We argue that moduli stabilization should occur near the string scale, and SUSY breaking should occur at a much lower intermediate scale, and that it originates from different sources.
We discuss the problems of moduli stabilization and SUSY breaking in currently popular scenarios, explain why these problems are generic, and discuss how our scenario can evade them. We show that String Universality is not inconsistent with phenomenology but that it is in conflict with some popular versions of brane world scenarios.
I. INTRODUCTION

It is well known that many fundamental questions in string theory, such as the question of moduli stabilization, are beyond the reach of perturbative techniques. Recent progress in elucidating the duality relations between the different perturbative theories has not resolved these questions, but has enabled them to be posed in a sharper fashion.

For instance, if in a given perturbative string theory the string coupling constant is small $g < 1$, so that perturbation theory is valid, no potential can be generated for the dilaton, so that the coupling constant which is the vacuum expectation value of the exponential of the dilaton is not fixed. S-duality on the other hand tells us that the strong coupling region $g > 1$, of this string theory also has a perturbative description in terms of the S-dual theory, so, by the same argument, no potential is generated in this theory. Thus one should expect that the string coupling is stabilized in a region which is inaccessible to perturbative

---

$^1$See for example chapter 18 (pp 359-362) of and references therein.
calculation from either of the S-dual theories, in other words, it is fixed at an intermediate value, at or around the self-dual point $g \sim 1$. Similar arguments can be made for the moduli governing the sizes of the compact dimensions (T-moduli), namely that they are fixed at or around the self-dual point under T-duality. This is because if the size of some compact dimensions is much larger than the string scale then the theory is effectively 5 or more dimensional, and no potential is generated for the size moduli, since no potential for the moduli is allowed in supergravity theories in more than four dimensions. If the size of some compact dimensions is smaller than the string scale then one takes the T-dual theory and makes the same argument. Thus internal dimensions should be fixed around the self-dual point.

FIG. 1. “Star” diagram illustrating all string theories, and 11D SUGRA as limits of one theory.

2A similar idea was proposed in and by Veneziano. Our arguments are similar in spirit but different in detail.

3We ignore here the so-called gauged supergravity theories which yield AdS spaces as vacuum configurations since a small 6-space would imply a large 4D cosmological constant. They may however be relevant to brane world scenarios.
Unfortunately, these self-dual points are in regions of minimal computability; they are in some sense equally far away from all the perturbatively computable string theories. The region in which we expect the true vacuum of string theory (with a fixed stable dilaton and compact extra dimensions) to lie, is in the middle of the well known star diagram (Figure 1) illustrating the unity of the different perturbative string theories.

Our key assumption is that the four dimensional low energy theory, which lives in the middle of the star diagram, should be in the universal region of string/M theory. This means that all five string theories and the 11D supergravity must be in the same universality class, and the four dimensional low energy physics of the different theories must be the same. From this point of view, the different perturbative string theories need not be any more than different perturbative (physical) regularization schemes. In particular, we assume that as one goes from the perturbative region at any of the cusps to the center of the star diagram, the infra-red spectrum that is common to all the different starting points is unchanged. For example, there should be a graviton in this region since it exists in all the perturbative string theories.\(^4\) As for the surviving gauge group, this is a more complicated issue since different starting points have different gauge groups, with recent developments, for instance F-theory, giving a wide variety of groups.\(^6\)

Traditionally the view that has been taken is that the real world is described by a

\(^4\)See, for example, [9].

\(^5\)This is similar to the assumption made in [10] in connection with the strongly coupled heterotic string, but [10] was written before the importance of duality was fully realized. What we are proposing here amounts in part to an attempt to extend the arguments of [10] to all perturbative string theories taking the dualities into consideration.

\(^6\)It may be the case that the gauge group that survives in the four dimensional theory in the central region, is the common subgroup of all the starting points. However, we would leave this for the moment as an open question.
single perturbative string model. In other words, that for some unknown reason nature picks one weakly coupled model over the others. Thus for instance up until recently it was thought that the heterotic $E_8 \times E_8$ (HE) theory was the theory that describes the real world. The realization that the different theories are just perturbative descriptions about different points in moduli space has changed that perspective. Nevertheless the traditional belief still survives in a modified form; for example, currently it has become fashionable to use the phenomenology of type I theories (with D-branes) on the grounds that they may be better descriptions of nature than the heterotic string.

One reason for the popularity of type I models is that in weak coupling heterotic theories one gets too small a value for the 4D Planck mass $M_P$, when “experimental” values for $\alpha_{YM}$ at unification are used and additionally the compactification scale $M_C$, is identified with the grand unification scale $M_{GUT}$. The argument was made by Witten \cite{Witten} in the heterotic SO(32) (HO) - type I S-duality context, that one might replace weakly coupled HO string theory $g_{HO} < 1$, by weakly coupled type I theory $g_I < 1$, in which case one could avoid this problem. However our string universality conjecture would give a different interpretation. The physics of the strongly coupled HO theory must be equivalent to the physics of the weakly coupled type I theory, where of course the dilaton cannot be stabilized, so that any comparison to phenomenology is not meaningful. The true vacuum should be around the self-dual coupling $g_{HO} \sim g_I \sim 1$.

In the HE case according to Horava and Witten \cite{Horava}, the strong coupling theory is (in the low energy limit) 11-dimensional supergravity on $S_1/Z_2$ (HW). The naive relations here would seem to give the size of the interval in the eleventh direction $\rho$ to be about 70 times the size of the six volume (and the 11D Planck scale). This has given rise to a picture where the gauge couplings evolve according to the four dimensional gauge theory picture but the gravitational coupling becomes five dimensional at a point well below the unification point \cite{Horava}. There is no possibility of such a picture arising in the HO/I theories where as we have argued above, the discrepancy of scales needs to be resolved with a coupling of order unity and there is no room for a five dimensional scenario. According to our conjecture of
string universality the HE theory should give the same low energy physics. So the naive M theory picture needs to be revised. Our (preliminary) investigations show that with currently available calculations of threshold corrections $g \sim O(1)$ is indeed a viable scenario.

It seems unlikely that the strong coupling picture, which implies that the 11D action is a good starting point for understanding 4D physics, is a description of the real world, since within this picture it is not possible to generate a potential which stabilizes the moduli. Indeed the phenomenology of HW theory is just a reparametrization of that coming from HE theory. In particular, the potentials that have been obtained upon compactification demonstrate the same runaway behavior as in the weak coupling analysis \cite{10,11}.

The calculation of the length of the eleventh dimension needs to revised by allowing for an order unity numerical factor in the relation between the Kaluza-Klein scale and the unification scale, and also to allow for the analog of threshold corrections. Indeed the latter may be taken into account by using a result in Witten’s original paper \cite{8}. If these numerical factors are included and we use a nonstandard embedding then a picture emerges where $\rho$ is of the order of the eleven dimensional Planck scale and the compactification scale. In other words the distance between the boundaries is of the order of the quantum fluctuations and no fifth dimension appears. This would then be compatible with our universality hypothesis since in the HO/I picture there is no room for a five dimensional picture below the string scale. In all this one is assuming that the compactification/unification scale is somewhat below, though close, to the string scale so that the field theory approximation makes sense.

Recently there has been much discussion within the context of so-called brane-world scenarios that the string scale may be as low as 1 TeV \cite{13,14}. While there is as yet no convincing string model that accounts for all the requirements that need to be met to have a

\footnote{These are the corrections that one would identify as threshold corrections (in the context of the field theory they are identified with certain Green-Schwarz anomaly cancelation terms) in the 10D theory as pointed out in \cite{5} and discussed in detail in \cite{12}.}
viable description of this type, the string models that might be considered as candidates for this are T-duals of the type I theory. For example consider compactifying on a six torus (or orbifold) and T-dualizing, in which case we get a IIB orientifold with $2^6$ orientifold planes and 32 D3 branes on which the standard model may live. Gravity however propagates in the bulk as well. If the string scale $l_s^{-1}$ is taken to be around a TeV then in order to get the right value for the Planck mass one needs to have the dimensions transverse to the D3 brane to be very large compared to the string scale (about $10^3 l_s$). Now by our universality hypothesis the same low energy physics must be seen from the U dual (S duality times T duality on a six torus) HO theory. But in the latter gravity and gauge theory propagate on the same space. The low energy theories are compatible only if in the former (i.e. the brane world) the transverse directions are of the order of the string scale. This is of course consistent with the argument above that the compactification scale should be around the string scale. However this would mean that the string scale must be close to the 4D Planck scale. Thus it seems that only the conventional view of the size of the string scale is compatible with our hypothesis. Of course, it is possible that these theories have to be considered in terms of non-standard compactifications (leading to gauged supergravity), but since it is not clear to us whether such a scenario has a viable string description we will not pursue this question further in this paper.

Any of the five string theories and the Horava-Witten theory must lead at low energies and in 4D to the same potential for the moduli, assuming a compactification which results in N=1 SUSY in 4d. The low energy theory should have a SUSY breaking minimum with vanishing cosmological constant, and 4D dilaton stabilized at such a value that the (unified) gauge coupling is weak. On the other hand, the 10D theory must have intermediate coupling so that string perturbation theory has the opportunity to break down, otherwise there is no way that a potential could have been generated. We assume that at intermediate, or strong coupling, the low energy 10D actions are in fact just determined by general covariance supersymmetry and gauge invariance, and the actions for the different string theories are obtained by field redefinitions. Whatever mechanism gives rise to the 4D potential, our string
universality assumption is that it ought to be independent of the particular perturbative string starting point.

General arguments based on PQ symmetries, and how they break due to non-perturbative effects, show that the superpotential must be a sum of exponentials in the moduli. These exponentials could be generated by stringy or field theoretic non-perturbative effects. For instance, the “race-track” mechanism for the stabilization of moduli\(^8\) envisages at least three exponential terms, coming, for instance from gaugino-condensates \(15,16\), which can balance against each other when all of them are small \(17,18\). A constant term in the superpotential while allowed by the symmetries has no natural mechanism for its generation. The one known exception is when the field strength of the antisymmetric two form field acquires a vacuum expectation value but this is quantized in units of the string scale \(19\) and hence yields too large a value for the scale of the gauge theory.

In the absence of a constant in the superpotential, if we use the Kahler potential of string perturbation theory, the moduli potential will give only a weak local minimum with zero cosmological constant while the global minimum has negative cosmological constant \(20\). One might think that this conclusion is avoided in the no-scale type models \(21\), but these are valid only at tree level and are destabilized, for example, by threshold corrections. So unless one finds a mechanism for generating a constant in the superpotential the only way out of this is to assume that the Kahler potential is drastically modified from its form in string perturbation theory \(7,22\). One can now speculate as to how such non-perturbative terms might arise (from wrapping of branes on cycles in CY spaces for instance). Combined with constraints coming from duality we find that it is extremely hard to generate significant contributions. We propose therefore that the question of moduli stabilization should be decoupled from SUSY breaking (See also \(4\)).

In the next section we review some known facts perturbative string theories and their

---

\(^8\)See \(6\) for a recent review.
dualities, and set up notation. In section 3 the origin of string non-perturbative effects and their dualities are discussed. Section 4 is devoted to highlighting the problems associated with currently popular scenarios for moduli stabilization and SUSY breaking. In section 4 we propose as an alternative a decoupling of the two issues. We suggest that moduli are all stabilized at or near the self-dual point by string non-perturbative effects while SUSY breaking happens at a much lower scale perhaps as a result of field theoretic non-perturbative effects.

II. PERTURBATIVE STRING THEORIES

For clarity we review some well known issues first. The perturbation expansion of any string effective action is given by

$$\Gamma[\phi, G_{\mu\nu}, B_{\mu\nu},...] = \sum_i e^{-\phi_0 \chi_i} S_i[\tilde{\phi}, G_{\mu\nu}, B_{\mu\nu},...]$$  \hspace{1cm} (1)$$

where the sum is over Riemann surfaces and \(\chi_i = 2 - 2h_i - b_i\) is the Euler character of the surface with \(h_i\) handles and \(b_i\) boundaries and we have split \(\phi = \phi_0 + \tilde{\phi}\) where the first term is the constant part of the dilaton defined so that \(\int \tilde{\phi} = 0\). This form of the effective action clearly restricts the form of the potential for \(\phi\). If one translates \(\phi \rightarrow \phi + \frac{1}{2} \ln t\) then each term in the expansion acquires a factor \(t^{1-h_i-b_i}\) and the only potential that would be allowed is of the form

$$V(\phi) = \sum_i \Lambda_i e^{-\phi \chi_i}. \hspace{1cm} (2)$$

In superstrings with unbroken SUSY (formulated on a flat background) \(\Lambda_i = 0\) for all \(i\)\(^9\) which is of course a necessary consistency condition. But in general such a potential is present in non-supersymmetric string theories (except that \(\Lambda_{S_2}\) is zero) or superstrings with broken SUSY (say by the Scherk-Schwarz mechanism). Assuming it exists, the critical point

\(^9\)Rigorously proven up to \(i = 2\) but expected to be true for all \(i\).
\( \phi = \phi_0 \) (which should be such that \( V(\phi_0) = 0 \)) is generically at \( g = e^{\phi_0} \sim 1 \) since the ratios of coefficients of the perturbation series should be of order one. In general the potential may be written as:

\[
V[\phi] = \sum_i \Lambda_i e^{-\phi \chi_i} + V_{np}[\phi],
\]

The non-perturbative term \( V_{np} \) is expected to depend on the coupling as \( e^{-1/g} \) or \( e^{-1/g^2} \). We will discuss later the contributions to \( V_{np} \) coming from brane-instanton effects.

Let us first list the low energy actions of the different string theories and their relations with each other. We only include the dilaton-gravitational and the gauge couplings since our discussion is going to be confined to the relations between these couplings.

The low energy effective action of type I string theory is the following,

\[
\Gamma_I = \frac{1}{(2\pi)^7 l_I^7} \int_{M_{10}} \left[ e^{-2\phi_I} \sqrt{-G_I (R + 4(\nabla \phi)^2)_I} \right] - \frac{1}{4(2\pi)^7 l_I^6} \int_{M_{10}} e^{-\phi_I} \sqrt{-G_I \text{tr} F^2_I}. \tag{4}
\]

The low energy effective action of heterotic SO(32) string theory (HO) is the following,

\[
\Gamma_{HO} = \frac{1}{(2\pi)^7 l_{HO}^7} \int_{M_{10}} \sqrt{-G_{HO} e^{-2\phi_{HO}}} \left\{ R + 4(\nabla \phi)^2 \right\}_{HO} - \frac{1}{4(2\pi)^7 l_{HO}^6} \int_{M_{10}} \sqrt{-G_{HO} e^{-2\phi_{HO}} \text{tr} F^2_{HO}}. \tag{5}
\]

The fields and parameters of these two theories which are S-dual to each other are related by

\[
\begin{align*}
\phi_I &= -\phi_{HO} \\
G_{\mu\nu,I} &= g_{HO} e^{-\phi_{HO}} G_{\mu\nu,HO} \\
g_I &= e^{<\phi_I>}_0 = e^{-<\phi_H>}_0 = \frac{1}{g_H} \\
l_{I}^2 &= g_{HO} l_{HO}^2.
\end{align*}
\]

\(^{10}V\) will of course depend on other moduli as well but we will ignore this for the time being.
The low energy effective action of heterotic $E8 \times E8$ string theory (HE) is the following,

$$\Gamma_{HE} = \frac{1}{(2\pi)^7 l_{HE}^8} \int_{M_{10}} \sqrt{-G_{HE}} e^{-2\phi_{HE}} \left\{ R + 4(\nabla \phi_{HE})^2 \right\}_{HE}$$

$$- \sum_i \frac{1}{4(2\pi)^7 l_{HE}^6} \int_{M_{10}} \sqrt{-G_{HE}} e^{-2\phi_{HE}} \text{tr} F_i^2.$$  \hspace{1cm} (7)

The HO and HE theories are related by T-duality. So $l_{HE} = l_{HO}$ and $G_{HE} = G_{HO}$ and if HO is compactified on a circle of radius $R$ then the physically equivalent HE is compactified on a circle of radius $l_{HE}^2/R$, with $e^{\phi_{HE}} = \frac{l_{HO}}{R} e^{\phi_{HO}}$.

In the strong coupling limit the HE theory goes over to the Horava-Witten (HW) \footnote{For discussions of the phenomenology of the Horava-Witten theory see the reviews 10,11,23.} theory whose action is given by

$$\Gamma_{HW} = \frac{1}{2\kappa_{11}^2} \int_{\mathcal{M}_{11}} d^{11}x \sqrt{G} R - \frac{1}{8\pi (4\pi k_{11}^2)^{2/3}} \left[ \int_{\mathcal{M}_{10}} d^{10}x \sqrt{G} \text{tr} F_1^2 + \int_{\mathcal{M}_{10}} d^{10}x \sqrt{G} \text{tr} F_2^2 \right].$$  \hspace{1cm} (8)

The gauge fields in HW are the same as in the HE theory and the metric is related by

$$ds_{HW}^2 = e^{-\frac{2}{3} \phi_{HE}} G_{\mu\nu,HE} dx^\mu dx^\nu + e^{\frac{4}{3} \phi_{HE}} dy^2,$$  \hspace{1cm} (9)

where $y$ is the eleventh coordinate. The parameters of the two theories are related by

$$l_{11} = l_{HE} g_{HE}^{1/3}, \quad \rho = l_{HE} g_{HE}$$  \hspace{1cm} (10)

where we put $2\kappa_{11}^2 = (2\pi)^8 l_{11}^8$ and $\int dx^{11} \sqrt{G}_{11} = 2\pi \rho$.

Let us now make the following reasonable assumptions about the low energy effective actions of all string theories and 11D SUGRA. These assumptions are usually made by most authors in superstring theory though they are not always explicitly stated.

- There exists a minimum of the effective action that breaks supersymmetry with zero cosmological constant.

- The low energy effective action can be written in terms of the perturbative spectrum of low energy fields even in the regime where perturbation theory is formally invalid,
except that some fields (moduli) which are not protected by gauge symmetries will acquire a potential and become massive. In particular, there will be massless graviton and gauge fields that will couple exactly as expected from the perturbative calculations because of general covariance and gauge invariance.

The first assumption is certainly non-controversial, but can be posed in two different degrees of severity. The weaker variant that we call the “practical cosmological constant problem”, is the one of ensuring that to a given accuracy within a given model the cosmological constant vanishes. Stated differently, it is the requirement that models should allow a large universe to exist with reasonable probability, and have 4D flat space as a solution. We believe that this requirement should be imposed on any model. The stronger variant is the general question of why the cosmological constant today is so small in natural units. This question is especially interesting in light of the hints that experimentally a small non-vanishing cosmological constant seems to be favored. Of course, the resolution of this issue is extremely important, but is outside the scope of many models, and we believe that it should not be absolutely required from models.

As for the second assumption, one might think that the spectrum of the theory at strong coupling is completely different from the spectrum at weak coupling. This is the case, for instance, in QCD, though even there the fundamental theory is still written in terms of the quarks and gluons. This is a possibility that cannot be ruled out at this stage of development of string theory, but existing indications (that we outline below) suggest that it is at least plausible that there are no phase transitions on the way from weak to strong coupling.

12The latter point appears to call into question the necessity of formulating M theory in terms of fundamental degrees of freedom that are different from the perturbative states of string theory. In the view expressed here all the non-perturbative objects - branes, Black holes etc. - are the analogs of instantons and solitons in field theory.
Perturbation theory establishes the form of the low energy action in the two extreme regions which are related by field redefinitions. For instance if \( \phi_{HO} \) is the HO dilaton then in the region \( \phi_{HO} \ll -1 \) we would have a perturbative heterotic description of the theory with perturbative gravity and gauge dynamics, and a type I description in the region \( \phi_{HO} \gg +1 \), with perturbative gravity and gauge dynamics. The region of coupling in which we cannot be sure that perturbation theory of at least one of the theories I and HO is valid is probably quite small (say, \( 0.7 < g_{HO} < 1.5 \)). In going from one side of this region to the other we get low energy 4D theories that are just field redefinitions of each other. Thus one might expect that at intermediate values of the coupling the nature of the theory is not radically altered. A similar argument holds for the HE/HW duality. Although HW theory has no dilaton (and there is no perturbative analog as in type I) the physical distance \( \rho \) along the transverse eleventh direction plays the role of the dilaton. Whatever the underlying theory is, in the region where all length scales including \( \rho \) are greater than \( l_{11} \) 11D SUGRA should be valid. This is a strong coupling version of the heterotic string effective action but there is no trace of any non-perturbative effects or phase transitions as is evident from the fact that the four dimensional effective action coming from HW is basically the same as that coming from string theory with appropriate field redefinitions. Thus again one might argue that the nature of the intermediate region is not very different from what these boundary regions would suggest except of course that the moduli would have a potential. One would then hope to extract some information about the central region from the physics of the boundary region.

III. STRING/M THEORY NON-PERTURBATIVE EFFECTS

String universality suggests that not only are different perturbative effective actions related, but also non-perturbative induced terms in these actions are universal. After all, much of the low energy physics of any string theory is determined by non-perturbative interactions. A natural source of such universal non-perturbative effects are BPS brane-
Motivated by string universality we propose that string/M theoretic non-perturbative effects (SNP) originate from supersymmetric BPS branes. The branes that we consider are objects which are point like in Euclidean four space and are obtained by wrapping the extended directions of the brane (including its Euclidean time direction) around some cycle in the compact space. Their action is the product of brane tension and the volume of the cycle around which its world volume is wrapped. Since under dualities, BPS branes transform into BPS branes, SNP in one string theory (or 11d SUGRA) transform into SNP in the other theory. The matching between BPS states in theories dual to each other is complete, and is actually considered one of the central pieces of evidence for the “correctness” of dualities. Therefore, we argue that we can make a complete list of SNP, by taking into account the known BPS branes, and that there’s no room for additional SNP which are sometimes assumed in some models. Such complete matching is obviously compatible with non-perturbative string universality.

We focus for simplicity on two moduli, the volume of the 6 compact dimensions $V$, and the string coupling $g$ (or the size of the 11d interval $\rho$ in Horava-Witten theory). In accordance with the general arguments about the dependence of SNP on moduli, we find that SNP depend exponentially on $V$ and on $1/g$ (or $1/g^2$). In the weak coupling, large volume region of moduli space, which we call the “boundary region” of moduli space, dominant SNP come from two or sometimes three BPS branes. The dominant SNP come from the brane-instantons whose actions have the weakest dependence on the string coupling $g$, and compactification radius $R$, or their product. Therefore for drawing conclusions about the boundary region of moduli space it is enough to consider only very few contributions.

\[\text{13Recently Sen has discussed non BPS branes, which deserve a separate discussion. Recently it has been argued that these correspond to sphelarons in field theory. But we do not expect our qualitative conclusions on moduli stabilization to be significantly modified by these effects.}\]
For each theory we compare the 4d Yang-Mills coupling \( \alpha_{YM} = \frac{g_{YM}^2}{4\pi} \), and Newton’s constant \( G_N \), and express the two couplings in terms of the string coupling and string scale of each string theory. We formally define the volume of the 6D compact manifold, \( V = R^6 \), and assume for simplicity that Euclidean time extension is also given by \( R \). Obviously more complicated possibilities can be considered, but we leave them to future work. In this section we will omit most numerical factors, \( \pi \)’s etc., since they will not be important for our purposes. The necessary numerical factors have appeared in previous sections or can be computed in a straightforward manner. SNP are simply given \( e^{-\text{action}} \) (ignoring the prefactor and questions related to fermion zero modes). In the following the branes are denoted by standard notation, \( F/D \) stands for a fundamental/Dirichlet brane\(^{14}\), followed by the number of spatial dimensions of the brane.

A. Type I vs. heterotic \( SO(32) \)

In type I and HO, the 4d Yang-Mills coupling \( \alpha_{YM} = \frac{g_{YM}^2}{4\pi} \), and Newton’s constant \( G_N \) are given by table I. The \( S \)-duality relations between the two I-HO couplings are given in

| \( \alpha_{YM} \) | Type I | HO |
|---|---|---|
| \( g_I \left( \frac{l_I}{R} \right)^6 \) | | \( g_{HO}^2 \left( \frac{l_{HO}}{R} \right)^6 \) |
| \( G_N \) | \( g_I^2 \left( \frac{l_I}{R} \right)^6 l_I^2 \) | \( g_{HO}^2 \left( \frac{l_{HO}}{R} \right)^6 l_{HO}^2 \) |

**TABLE I.** \( \alpha_{YM} \) and \( G_N \) in type I/HO.

\(^{14}\)For uniformity of notation we have denoted all branes which are conventionally called NS, because they are sources for NS-NS fields, by the letter \( F \) (conventionally used only for the NS 1-brane). After all D branes are not called “R branes”.
equations (6). In table II we list the type of branes and their actions in the two S-dual theories, type I, and HO. In each row the branes are related by duality, and their actions are related by (6). In type I there are only D-branes, and in HO only the heterotic string

| Type I | HO |
|-------|----|
| **brane** | **brane** | **action** | **action** |
| D5 | $\frac{1}{g_I} \left( \frac{R_{lI}}{R} \right)^6$ | F5 | $\frac{1}{g_{lHO}} \left( \frac{R_{lHO}}{R} \right)^6$ |
| D1 | $\frac{1}{g_I} \left( \frac{R_{lI}}{R} \right)^2$ | F1 | $\left( \frac{R_{lHO}}{R} \right)^2$ |

TABLE II. BPS branes and actions in type I/HO.

and 5-brane.

### B. M on $R^{10} \times S^1/Z_2$ vs. heterotic $E_8 \times E_8$

In HE and HW theories, the 4d Yang-Mills coupling $\alpha_{YM} = \frac{g^2_{YM}}{4\pi}$, and Newton’s constant $G_N$ are given by table III. The duality relations between the two HE-HW couplings are

| | HE | HW |
|---|---|---|
| $\alpha_{YM}$ | $g^2_{HE} \left( \frac{l_{HE}}{R} \right)^6$ | $\left( \frac{l_{HE}}{R} \right)^6$ |
| $G_N$ | $g^2_{HE} \left( \frac{l_{HE}}{R} \right)^6 \frac{l^2_{HE}}{4\pi}$ | $\frac{1}{\rho} \left( \frac{l_{HE}}{R} \right)^6 \frac{l^3_{11}}{4\pi}$ |

TABLE III. $\alpha_{YM}$ and $G_N$ in HE/HW.

\begin{align*}
    l_{HE} \ g_{HE} &= \rho \\
    l_{HE}^{\frac{1}{3}} g_{HE} &= l_{11}
\end{align*}  

(11)
\[ l_{HE}^2 = \frac{l_{11}^3}{\rho}, \]

where only two are independent.

In table IV we list the type of branes and their actions in the two theories, HE, and HW. In each row the branes are related by duality, and their actions are related by \((\mathbb{I})\). In HE there are only the F string and its magnetic dual the F 5-brane. \(M\) stands for an \(M\)-theory brane, followed by the number of spatial dimensions of the brane. We need to discuss two possible orientations of each type of brane, one where one direction is longitudinal to the 11th dimension and the other where all brane directions are transverse to it. Only one possible orientation of each \(M\)-brane has a dual and therefore (we conjecture) is allowed. The other two possibilities, \(M5\) longitudinal and \(M2\) transverse\(^{15}\), are (we believe) absent (i.e. their pre-factors should vanish) since in the S-dual HE theory they are absent. In the HW theory this happens at the boundaries because the fields that they couple to are odd under the \(Z_2, x_{11} \rightarrow -x_{11}\) and are projected out there.

| HE | action | HW | action |
|----|--------|----|--------|
| \(F5\) | \(\frac{1}{g_{HE}^2 \left( \frac{R}{l_{HE}} \right)^6}\) | \(M5\) transverse | \(\left( \frac{R}{l_{11}} \right)^6\) |
| \(F1\) | \(\left( \frac{R}{l_{HE}} \right)^2\) | \(M2\) longitudinal | \(\left( \frac{R}{l_{11}} \right)^2 \frac{\rho}{l_{11}}\) |

TABLE IV. BPS branes and actions in HE/HW.

\(^{15}\)These would correspond to \(D_4\) and \(D_2\) branes in the string theory limit and are present in IIA but are absent in the heterotic theories.
C. \( M \) on \( R^{10} \times S^1 \) vs. type IIA

In \( M \) on \( R^{10} \times S^1 \) (MS1) and type IIA theories, the comparison of the 4d Yang-Mills coupling is meaningless since there are no perturbative gauge fields in these theories. Newton’s constant \( G_N \) is given by table V. The duality relations between the two IIA, HW couplings

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|}
\hline
 & IIA & MS1 \\
\hline
\( G_N \) & \( g_{IIA}^2 \left( \frac{l_{11}}{R} \right)^6 l_{IIA}^2 \) & \( \frac{1}{\rho} \left( \frac{l_{11}}{R} \right)^6 l_{11}^3 \) \\
\hline
\end{tabular}
\caption{\( G_N \) in IIA/M.}
\end{table}

are

\[
\begin{align*}
l_{IIA} g_{IIA} &= \rho \\
l_{IIA} g_{IIA}^{1/3} &= l_{11} \\
l_{IIA}^2 &= \frac{l_{11}^3}{\rho},
\end{align*}
\]

(12)

where only two are independent.

In table VI we list the type of branes or states and their actions in the two \( S \)-dual theories, type IIA, and HW. The branes are denoted by standard notation, \( F/D \) stands for a fundamental/Dirichlet brane, followed by the number of spatial dimensions of the brane. The notation \( KK \) denotes Kaluza-Klein states. In each row the branes or \( KK \) states are related by duality, and their actions are related by (12).

D. Type IIA vs. type IIB

In type IIA and type IIB theories, the 4d Yang-Mills coupling \( \alpha_{YM} = \frac{g_{YM}^2}{4\pi} \), and Newton’s constant \( G_N \) are given by table VII. For one single \( T \)-duality the relations between the IIA, IIB couplings are
TABLE VI. BPS branes and actions in IIA/MS1. Dominant contributions are marked by *.

| brane | action | brane | action |
|-------|--------|-------|--------|
| * D0  | $\frac{1}{g_{\text{IIA}}} \frac{R}{l_{\text{IIA}}}$ | KK graviton | $\frac{R}{\rho}$ |
| * F1  | $\left( \frac{R}{l_{\text{IIA}}} \right)^2$ | M2 longitudinal | $\left( \frac{R}{l_{\text{II}}} \right)^2 \frac{\rho}{l_{\text{II}}}$ |
| D2    | $\frac{1}{g_{\text{IIA}}} \left( \frac{R}{l_{\text{IIA}}} \right)^3$ | M2 transverse | $\left( \frac{R}{l_{\text{II}}} \right)^3$ |
| D4    | $\frac{1}{g_{\text{IIA}}} \left( \frac{R}{l_{\text{IIB}}} \right)^5$ | M5 longitudinal | $\left( \frac{R}{l_{\text{IIB}}} \right)^5 \frac{\rho}{l_{\text{IIB}}}$ |
| F5    | $\frac{1}{g_{\text{IIA}}} \left( \frac{R}{l_{\text{IIB}}} \right)^6$ | M5 transverse | $\left( \frac{R}{l_{\text{IIB}}} \right)^6$ |

TABLE VII. $\alpha_Y$ and $G_N$ in IIA/IIB.

| | IIA | IIB |
|---|-----|-----|
| $\alpha_Y$ | $g_{\text{IIA}}^2 \left( \frac{l_{\text{II}}}{l_{\text{IIA}}} \right)^6$ | $g_{\text{IIB}}^2 \left( \frac{l_{\text{IIB}}}{l_{\text{IIB}}} \right)^6$ |
| $G_N$ | $g_{\text{IIA}}^2 \left( \frac{l_{\text{II}}}{l_{\text{IIA}}} \right)^6 l_{\text{IIA}}^2$ | $g_{\text{IIB}}^2 \left( \frac{l_{\text{IIB}}}{l_{\text{IIB}}} \right)^6 l_{\text{IIB}}^2$ |

In table VIII we list the type of branes or states and their actions in the two $T$-dual theories, type IIA, and IIB. The notation $KK$ denotes Kaluza-Klein states. In each row the branes or $KK$ states are related by duality, either in a transverse (denoted by $T$) or a longitudinal direction (denoted by $L$), and their actions are related by (13). $KK$ MM stands

$$l_{\text{IIA}} = l_{\text{IIB}} \equiv l_{\text{II}}$$

$$R_{\text{IIA}} = \frac{l_{\text{II}}^2}{R_{\text{IIB}}}$$

(13)

$$g_{\text{IIA}} = \left( \frac{l_{\text{II}}}{R_{\text{IIB}}} \right) g_{\text{IIB}}.$$

In table VIII we list the type of branes or states and their actions in the two $T$-dual theories, type IIA, and IIB. The notation $KK$ denotes Kaluza-Klein states. In each row the branes or $KK$ states are related by duality, either in a transverse (denoted by $T$) or a longitudinal direction (denoted by $L$), and their actions are related by (13). $KK$ MM stands
for Kaluza-Klein momentum mode. The size of a direction that is not T-dualized is denoted \( R \). Unlike the previous cases of S-duality, we have to make a distinction between different radii, since even if their sizes were equal to begin with, the T-duality changes them.

This completes the basic relations between the theories at the edges of the star diagram [I]. There are additional relations between the perturbative string theories, but they are generated by the relations we have listed.

### E. Discussion

Our conjecture has been that the all the possible string non-perturbative effects are accounted for by the D and F instantons (in 4D) given in the above tables. We note that all expected SNP in each theory appear, and that there’s no room for exotic SNP. For example, suppose that in the HO theory we look for SNP of strength \( e^{-\frac{1}{g_{HO}}(\frac{R}{R_{HO}})^6} \) as might be required for Kahler stabilization [II]. By S-duality this would require an effect in type I of

|   | IIA |   | T/L | IIB |
|---|-----|---|-----|-----|
| * | \( D0 \) | \( \frac{1}{g_{IIA}} \frac{R_{IIA}}{t_{II}} \) | L | \( D-1 \) | \( \frac{1}{g_{IIB}} \) |
| * | \( D0 \) | \( \frac{1}{g_{IIA}} \frac{R}{t_{II}} \) | T | \( D1 \) | \( \frac{RR_{IIB}}{g_{IIB}t_{II}} \) |
| * | \( F1 \) | \( \frac{RR_{IIA}}{t_{II}} \) | L | \( KK \) MM | \( \frac{R}{R_{IIB}} \) |
| * | \( F1 \) | \( \frac{R^2}{t_{II}} \) | T | \( F1 \) | \( \frac{R^2}{t_{II}} \) |
| | \( D2 \) | \( \frac{1}{g_{IIA}} \frac{R^2_{IIA}}{t_{II}} \) | L | \( D1 \) | \( \frac{1}{g_{IIB}} \left( \frac{R}{t_{II}} \right)^2 \) |
| | \( D2 \) | \( \frac{1}{g_{IIA}} \left( \frac{R}{t_{II}} \right)^3 \) | T | \( D3 \) | \( \frac{1}{g_{IIB}} \frac{R_{IIB}R^3}{t_{II}} \) |

**TABLE VIII.** BPS branes and actions in IIA/IIB. Dominant contributions are marked by *. 

---

[1]: Reference for the star diagram
[II]: Reference for Kahler stabilization
the form $e^{-g_1^{-\beta} \left( \frac{\mu}{\eta} \right)^{2\beta}}$, which surely cannot exist. Similarly we would like to argue that in the HW theory the correction to the Kahler potential coming from transverse membranes \cite{28} actually vanishes since (see table \[IV\] and the related discussion) this has no analog in the S-dual HE theory. Indeed it would be surprising (and contrary to the spirit of String Universality) if such an NP contribution to the HW theory were present since there appears to be no way this could arise in the HO theory or its S-dual type I theory. Another piece of supporting evidence for this is that in the IIA/MS1 case (see table \[VI\]) all contributions in the strong coupling theory have a weak coupling analog so it would be strange if this were not the case in the HW theory.

Considering only the leading exponential dependence it is clear that moduli are unstable, since they have runaway potentials which force them towards free 10d (11d) theories. We therefore conclude that moduli stabilization cannot occur when either the inverse coupling or the volume are parametrically large.

**IV. MODULI STABILIZATION AND SUSY BREAKING**

In this section we will discuss possible mechanisms for moduli stabilization, and their likely values. We have already argued that stabilization of moduli in the boundary region of moduli space is unlikely. Here we explicitly show why this is so. We then argue that stabilization near the self-dual point, motivated by string universality, is plausible, and that SUSY breaking is likely to occur at a much lower scale than moduli stabilization scale. To discuss moduli stabilization in more detail we need to take into account supersymmetry in a more quantitative way. Moduli are chiral superfields of $D = 4$, $N = 1$ SUGRA, which means that their interactions are constrained by the general form of the $N = 1$ SUGRA Lagrangian.

The potential for moduli $\phi_i$ in $N = 1$ supergravity is given by \cite{29} the following

\[^{16}\text{A convenient source for the formulae in this section is volume II chapter 18 of [1] which should}\]
expression \footnote{17}

\[ V = e^\theta (G_i \tilde{G}^{i\bar{j}} \tilde{G}^j - 3) = F_i \tilde{G}^{i\bar{j}} \tilde{F}^j - 3 e^\theta, \]  
\[ (14) \]

and \( G \) is usually written as

\[ G = K(\phi_i, \bar{\phi}_i) + \ln |W(\phi_i)|^2, \]  
\[ (15) \]

where the real analytic function \( K \) is the Kahler potential, the holomorphic function \( W \) is the superpotential and \( F_i \equiv e^{\frac{i}{2}} \tilde{G}_i \) is an order parameter for SUSY breaking. Also

\[ G^{ij} = G_{ij}^{-1}, \quad G_i = \frac{\partial G}{\partial \phi_i}, \quad G_{ij} = \frac{\partial^2 G}{\partial \phi_i \partial \phi_j}. \]  
\[ (16) \]

The potential \( V \) can be expressed in terms of \( K \) and \( W \) as follows,

\[ V = e^K \left[ D_i W K^{\bar{j}} D_j W - 3 |W|^2 \right] \]  
\[ (17) \]

where \( D_i W = \partial_i W + \partial_i K W \). The coupling to the gauge sector modifies the \( F \)-terms into

\[ F_i = e^{\frac{i}{2}} \tilde{G}_i + \frac{1}{4} f_i \lambda \]  
\[ (18) \]

(where \( f_i = \frac{\partial f}{\partial \phi^i}, f \) being the gauge coupling function) and the potential can then still be written in the form given in the second equality of (14).

Let us consider string theory compactified on a 6-manifold which preserves \( N = 1 \) SUSY. Conventionally a four dimensional complex modulus \( S \) (related to the model independent 10D dilaton/axion) and three complex moduli \( T_i \) (related to the size and shape of the compact manifold) are defined. In particular, we have

\[ \text{Re} S = e^{-2\phi} \sqrt{\det G_c} \]  
\[ (19) \]

where \( \phi \) is the 10D dilaton and the last factor on the RHS is the measure on the compact manifold. Then in perturbative string theory we have at tree level,

\footnote{17} In this section we will put \( \kappa = M_P^{-1} = 1 \) for convenience.
\[ \mathcal{G} = - \sum_{i=1}^{3} \ln (T_i + \bar{T}_i - C^a \bar{C}^a) - \ln (S + \bar{S}) + \ln |W(C)|^2, \]  

(20)

where \( C^a \) are matter fields coming from the 10D gauge fields tangent to the compact directions. The first term comes from the metric of the 6 manifold (CY space or orbifold) which is parametrized in terms of the three complex \( T_i \) moduli. This form of the function \( \mathcal{G} \) is (a generalization of) the so-called no-scale model \([21]\) which leads to a positive definite potential. Up to field redefinitions (involving the \( S \) and \( T \) fields) it is also obtained from the compactification of the HW theory \([10,11]\).

We would like at this point to emphasize the appearance and possible implications of the “practical cosmological problem”, which as we explained is not the same as the “cosmological constant problem”, whose solution should not be a requirement of phenomenological SUSY breaking models. In this particular context the “practical cosmological problem” takes the following form. To allow models of supersymmetry breaking to predict reliably the structure of soft-supersymmetry-breaking terms it is essential that the absolute value of the potential at its minimum does not exceed\(^{18}\) the order of \( m_{3/2}^4 \), where \( m_{3/2} \) is the gravitino mass which determines a typical size of soft breaking terms. In most supergravity models (not of the no-scale type) the value of the potential at the minimum has a tendency to be negative and of order \( M_P^2 m_{3/2}^2 \). In addition to the loss of predictive power this feature has obvious disastrous cosmological consequences, to be avoided. Most non-perturbative SUSY breaking models suffer from this problem and special measures have to be taken to remedy the situation.

We discuss three different classes of models:

- no-scale type models, which solve automatically the practical cosmological constant problem (at least at tree level), but have trouble in stabilizing moduli and predicting non-vanishing gaugino masses.

\(^{18}\)Assuming that there are no additional phase transitions at intermediate scales in between the SUSY breaking scale and the scale at which it becomes explicit.
• “race-track” type models which have problems breaking SUSY, stabilizing moduli and solving the practical cosmological constant problem.

• our suggestion for stabilization around the self-dual point.

A. No scale models

No scale models are models for which (20) holds\(^{19}\). Let us assume now for simplicity that the three \(T\) moduli are identified. At string tree level the gauge coupling function \(f = f_S S\) so that, \(F_T = e^{\frac{1}{2}G} G_T\), \(F_C = e^{\frac{1}{2}G} G_C\), \(F_S = e^{\frac{1}{2}G} G_S + \frac{1}{4} f_S (\lambda \lambda)\) and the potential becomes,

\[
V = F_s G_S \bar{G} S + e^K \frac{|\partial_C W|^2}{3(T + \bar{T} - |C|^2)^2}.
\]

(21)

This is a positive definite potential whose minimum is at \(F_S = 0\) and \(\partial_C W = 0\). SUSY may still be broken if \(F_T \neq 0\) and/or \(F_C \neq 0\). This requires that the superpotential at the minimum be non-vanishing. In the tree level compactification \(W = d_{ijk} C^i C^j C^k\) where \(d_{ijk}\) are coupling constants. At the minimum we must have \(\partial_l W = d_{ijl} C^i C^j + d_{ilj} C^i C^j + d_{lij} C^i C^j = 0\), but this implies that \(W = 0\) at the minimum and no SUSY breaking. So to have broken SUSY in this scenario one must require a constant in the superpotential \(W_0 = \Lambda^3\) (generated by some yet unknown non-perturbative mechanism). Furthermore \(\Lambda\) must be around \(10^{13} GeV\) in order to get the required scale of SUSY breaking. It is not at all clear why such a scale should arise but if it does exist\(^{20}\) then (taking into account gaugino condensation) the dilaton is stabilized with SUSY broken at an acceptable scale.

As is well known, no scale type models have several problems even if one assumed that the form of the \(G\) and the tree level result for \(f\) survived string loop corrections and that

\(^{19}\)The discussion in this subsection is included for reasons of clarity and completeness. All the material given here can be found in the literature.

\(^{20}\)In available examples constants in the superpotential are quantized in units of the string scale.
there is a constant in the superpotential resulting in SUSY breaking of the desired magnitude. We list them here for completeness.

- a) The T-moduli are not determined.

- b) Since SUSY breaking is not dilaton dominated generically there will be flavor changing neutral currents.

- c) The gaugino masses are given by the formula $m_\lambda = \langle f_i G_i \bar{F}_j \rangle$ [29]. If $f_i = f S$ as in tree level string theory $f_i = fS$ and since $F_S = 0$ at the minimum this mass vanishes (note that $G^{ST} = 0$ in these models).

Of course the gauge coupling function is corrected by string loop effects and one should write $f = fS + fT$. But in this case there are two additional terms in the potential. i.e. defining $Q_i = \frac{1}{4} f_i \lambda \lambda$

$$Q_T \bar{G}^{TT} Q_T + 2ReQ_T \bar{G}^{TT} e^\frac{i}{4} \bar{G} T.$$

The second term clearly spoils the positive definiteness of the potential. Alternatively if one works with an effective Lagrangian after gaugino condensation one has a superpotential that is now T-dependent and that will spoil the no-scale property. Thus even if one assumed that the general no-scale form of the Kahler potential [20] is unaffected by quantum corrections one cannot preserve the positive definiteness of the potential and get non-zero gaugino masses.

Since the no-scale field is expected to be quite light, and its interaction are of gravitational strength, a particularly severe potential problem for this class of models is the amount of energy stored in the no-scale scalar field and the implications of this energy on late cosmological evolution. This is a manifestation of the so-called moduli problem [30].

**B. “Race track” models: Stabilization near the boundary of moduli space**

An alternative to no-scale models is that the parameters of the superpotential and the Kahler potential are chosen in such a way that results in a SUSY breaking minimum with
vanishing cosmological constant. This has been the class of models of choice for most previous works on string phenomenology, and it was recently revived in a slightly different context [32]. It is usually assumed that stabilization and SUSY breaking occur at the same scale, which has to be much below the string scale. This leads to known problems, which we recall for emphasis, and further argue that they are generic to such models.

We will argue shortly that it is very unlikely that a minimum of vanishing cosmological constant which breaks SUSY is found at weak coupling and large compactification volume (which we called “the boundary region of moduli space”). But even if we assume that such a minimum exists then there is always a deep minimum with a large negative cosmological constant towards weaker coupling and larger volume [20]. In addition, there’s always a supersymmetric minimum at vanishing coupling and infinite volume. This multi-minima structure brings into focus the barriers separating them. If these barriers are high enough one may argue that flat space is a metastable state with a large enough life time. Generically, however, this is not the case, and classical or quantum transitions between minima are quite fast. In the context of gaugino-condensation race-track models this was discussed in [20]. In particular, in a cosmological setup it was shown [20] that classical roll-over of moduli towards weak coupling and large volume are generic, and occur for a large class of moduli initial conditions. Later it was shown that cosmic friction can somewhat improve the situation [33], and recently it was argued that finite temperature effects drastically improve the situation [34].

We would like to show that the problems arising in this class of models are rooted at their basic assumptions, and that they cannot be remedied by choosing different parameters or playing with numbers. Our conclusion is that, at the very least, it is inconsistent to consider, in this class of models, only \( T \) and \( S \) moduli, and parametrize SUSY breaking in terms of a complex vector in \((F_T, F_S)\) plane as suggested in [35,36]. This conclusion was in fact already recognized in gaugino-condensation models by [36], where an additional ad-hoc chiral superfield was included, and later also in [37,38].

Let us now examine more closely the possibility of stabilizing moduli near the bound-
aries of moduli space. We consider generic moduli chiral superfield, which we denote by $S$, which could be either the dilaton $S$-modulus, or the $T$-modulus. We assume that its Kahler potential is given by $K = -\ln(S + S^*)$, and that $ReS > 0$, corresponding to having a well defined compactification volume and gauge coupling. The generic feature of the superpotential $W(S)$ near the boundaries of moduli space is its steepness. This has to be so, because we insist that a new scale, much lower than the string scale is generated dynamically, and the ratio of this new scale to the Planck scale has to be reached within about Planck distance in moduli space. This requires that derivatives of the superpotential are large. In mathematical terms, the steepness property of the superpotential is expressed as follows,

$$\left| \frac{(S + S^*)\partial_S^{(n+1)}W}{|\partial S W|} \right| \gg 1 \quad n = 0, 1, 2, 3.$$  \hspace{1cm} (23)

This property certainly holds for all “gaugino-condensation” superpotentials, but as explained, it is generic to all models of stabilization around the boundaries of moduli space. The typical example of a superpotential satisfying (23) is a sum of exponentials $W(S) = \sum_i e^{-\beta_i S}$, with $Re\beta_i \gg 1$, in the region $|S \beta_i| \gg 1$. In this example the “boundary region of moduli space” is simply the region $ReS \gg 1$, but in general, the precise definition will depend on the details of the model. It is good to keep this example in mind while going through the following arguments, but we will not use any particular specific form for $W$.

Inequality (23) holds as a functional relation and can be, of course, violated at isolated points. Obviously (see eq. (26)), it is violated at extrema. But the violations of (23) are only in some of the relations between the derivatives, for example the first and second derivatives, while for the rest, the rule that the higher the derivative, the larger it is, still holds.

The potential $V$, is given in terms of the superpotential $W$, and its derivatives

$$V(S, S^*) = (S + S^*)F(S, S^*)F^*(S, S^*) - \frac{3}{S + S^*}W(S)W^*(S^*),$$  \hspace{1cm} (24)

where $F(S, S^*) = \partial_S W(S) - \frac{1}{S + S^*}W(S)$. The first derivatives of the potential are give by

$$\partial_S V = (S + S^*)\partial_S^2 W(S)F^* - \frac{2}{S + S^*}FW^*$$
$$\partial_{S^*} V = (S + S^*)\partial_{S^*}^2 W^*(S^*)F^* - \frac{2}{S + S^*}F^*W.$$  \hspace{1cm} (25)
An extremum is determined by solutions of $\partial S V = 0$, that is,

$$(S + S^*)^2 \partial S^2 W(S) F^* = 2 F W^*; \quad (26)$$

which can be satisfied in two ways,

$$F \neq 0 \quad \text{and} \quad (S + S^*)^2 \partial S^2 W(S) F^* = 2 F W^* \quad (27)$$

$$F = 0. \quad (28)$$

At an extremum with a vanishing $F$ as in (28), SUSY is unbroken and the cosmological constant is $-3|W|^2$, which is generically too large to obey our requirement that it solves the practical cosmological constant problem. If $W$ is also tuned to zero at the minimum, then one encounters the problems associated with the appearance of an additional deep minimum with negative cosmological constant and those arising from the multi-minima structure which were alluded to previously. Only an extremum as in (26) breaks SUSY, but as we will show shortly it is never a minimum for steep potentials.

To determine whether the extrema are minima, maxima or saddle points we need to analyze the matrix of second derivatives of the potential at the extrema. Using the following expressions for the derivatives of $F$

$$\partial S F(S, S^*) = \partial S^2 W(S) - \frac{1}{S + S^*} F(S), \quad (29)$$

and

$$\partial S^* F(S, S^*) = \frac{1}{(S + S^*)^2} W(S), \quad (30)$$

and similar expressions for derivatives of $F^*$, we can compute second derivatives of $V(S, S^*)$,

$$\partial^2_{SS} V = -\frac{2}{(S + S^*)^2} WW^* - \frac{2}{(S + S^*)} FF^* + (S + S^*) \partial S^2 W \partial S^2 W^* \quad (31)$$

$$\partial^2_ S V = \frac{4}{(S + S^*)^2} F W^* - \frac{1}{S + S^*} \partial S^2 WW^* + \partial S^2 WF^* + (S + S^*) \partial S^3 WF^*. \quad (32)$$

Expressions (31,32) take the following form,

Case (i): $F = 0$; $\partial S V = 0$ at the extremum $S_0$, then at $S_0$
\[
\partial_{SS}^2 V = -\frac{2}{(S + S^*)^3} WW^* \quad \text{and} \quad \partial_{SS}^2 W \partial_{S}^2 W^*, \quad (33)
\]

and

\[
\partial_{S}^2 V = -\frac{1}{S + S^*} \partial_{S}^2 WW^*. \quad (34)
\]

Case (ii): \( F \neq 0; \partial_S V = 0 \) at the extremum \( S_0 \), then at \( S_0 \)

\[
| (S + S^*)^2 \partial_{S}^2 W | = |2W|, \quad (35)
\]

and therefore

\[
\partial_{SS}^2 V = \frac{2}{(S + S^*)^3} WW^* - \frac{2}{S + S^*} FF^*, \quad (36)
\]

and

\[
\partial_{S}^2 V = (S + S^*) \partial_{S}^2 W F^* - \frac{1}{S + S^*} \partial_{S}^2 WW^* + \frac{6}{(S + S^*)^2} FW^*. \quad (37)
\]

From these expressions we can calculate the various partial derivatives with respect to \( S_R = Re(S) \) and \( S_I = Im(S) \). Using

\[
\frac{\partial^2 V}{\partial S_R \partial S_R} = + \frac{\partial^2 V}{\partial S \partial S} + \frac{\partial^2 V}{\partial S^* \partial S^*} + 2 \frac{\partial^2 V}{\partial S \partial S^*},
\]

\[
\frac{\partial^2 V}{\partial S_I \partial S_I} = - \frac{\partial^2 V}{\partial S \partial S} - \frac{\partial^2 V}{\partial S^* \partial S^*} + 2 \frac{\partial^2 V}{\partial S \partial S^*},
\]

\[
\frac{\partial^2 V}{\partial S_R \partial S_I} = + \frac{\partial^2 V}{\partial S \partial S} - \frac{\partial^2 V}{\partial S^* \partial S^*}. \quad (38)
\]

The relevant quantity is the determinant of the matrix of second derivatives,

\[
H = \frac{\partial^2 V}{\partial S_R \partial S_R} \frac{\partial^2 V}{\partial S_I \partial S_I} - \left( \frac{\partial^2 V}{\partial S_R \partial S_I} \right)^2 = -4 (\partial_{SS} V \partial_{S^*S^*} V - \partial_{SS} V \partial_{S^*S} V). \quad (39)
\]

We can now substitute Eqs. \((33, 34)\) and \((36, 37)\) into \((39)\) and obtain expressions for \( H \) at the extrema. First, Case (i),

\[
H = 4 \left( (S + S^*)^2 (|\partial_{S}^2 W|^2)^2 - \frac{5}{(S + S^*)^2} |\partial_{S}^2 W|^2 |W|^2 + \frac{4}{(S + S^*)^3} (|W|^2)^2 \right), \quad (40)
\]
where we have used eqs. (33, 34). Expression (40) for $H$ can be written as

$$H = 4 \left( (S + S^*) |\partial_S^2 W|^2 - \frac{2}{(S + S^*)^3} |W|^2 \right)^2 - \frac{4}{(S + S^*)^2} |\partial_S^2 W|^2 |W|^2. \quad (41)$$

This means that the extremum of type (i) is either a local minimum or a local maximum. To check which of the two, it is enough to choose an arbitrary direction and check if the second derivative is positive or negative. For example, choose the $S_R$ direction. From eq. (38) we obtain

$$\partial_{S_R S_R} V = 2(S + S^*) |\partial_S^2 W|^2 - \frac{2}{S + S^*} Re(\partial_S^3 W W^*) - \frac{4}{(S + S^*)^3} |W|^2. \quad (42)$$

Using Eq. (23), we see that if $\partial_S^3 W \neq 0$ at the extremum of case (i), then from (42)

$$\partial_{S_R S_R}^2 V \approx 2(S + S^*) |\partial_S^2 W|^2 > 0. \quad (43)$$

The conclusion is that a generic extremum of type (i) is a minimum. If additional conditions are imposed on the superpotential it may be possible to have a local maximum at an extremum of type (i).

The analysis in case (ii) is a little bit more complicated but is essentially the same analysis. For case (ii), rather than write the full complicated expressions, we analyze them using eq. (23). If $\partial_S^3 W \neq 0$ at the extremum then

$$H \approx -4(S + S^*)^2 |\partial_S^2 W|^2 |F|^2, \quad (44)$$

so $H < 0$ and the extremum is necessarily a saddle point. If $\partial_S^3 W = 0$ at the extremum, then using Eq. (33), one finds that $H \approx +\frac{16}{(S + S^*)^2} (|F|^2)^2$ so $H > 0$ and the extremum is either a maximum or a minimum. Checking, for example, $\partial_{S_R S_R}^2 V \approx -\frac{4}{S + S^*} |F|^2$ reveals that this is a local maximum. No further simplification can occur and therefore the analysis is complete. The conclusion is that, under the conditions of Eq. (23), an extremum of case (ii) is either a saddle point or a maximum, but never a minimum.

So far we have used the moduli perturbative Kahler potential $K = -\ln(S + S^*)$ in our discussion. We expect the results of our analysis to be similar if this Kahler potential receives
some moderate corrections. If, as expected, corrections to the Kahler potential preserve the steepness of the potential so that superpotential derivatives are larger than derivatives of the Kahler potential, an analog of (23) exists, in which powers of \((S + S^*)\) are replaced, where appropriate, by \(e^K\), \(\partial_S K\) or \(\partial_{S^*} K\). The subsequent analysis follows through, since the essential ingredient that we have used was a classification of the largest terms in the equations, which should still be valid. We do not consider the case when corrections to the Kahler potential are larger than its original perturbative value, since this would mean that perturbation theory is badly broken in the outer region of moduli space, an unlikely situation in contradiction with available information.

To summarize, we have found that it is not possible to find a SUSY breaking minimum in the region where the superpotential is a steep function. This is a very general conclusion which shows how hard it is to make “race-track” models work. This conclusion has been reached previously using different arguments in the context of gaugino condensation models \([20,36,38]\).

C. Stabilization around the self-dual point

We propose that moduli stabilization around the self-dual point is a plausible scenario, allowing the possibility of SUSY breaking while evading the problems we discovered for the other scenarios. Our discussion will be qualitative, postponing the quantitative analysis to a dedicated project \([39]\).

Our suggestion, motivated by string universality is the following:

- Moduli are stabilized at a scale below, but not much below, the 4D Planck scale, which is similar to the string scale. All flat directions are lifted at this stage, leaving no light moduli. A possible source of moduli potential can be SNP. SUSY is unbroken at that scale, and the cosmological constant at the minimum is parametrically smaller than \(M_P^4\). The dynamics which results in this must be intrinsically stringy.
• SUSY is broken at a much lower intermediate scale, by additional small non-perturbative effects, which do not spoil the stabilization of moduli. The dynamics here probably can be understood in field theoretic terms, for example gaugino-condensation in the hidden sector.

The moduli stabilizing superpotential should be of order $M_P^3$, and all the F-terms, and the superpotential itself should vanish at the minimum. SNP are likely to obey these requirements, which are nothing but the consistency requirement that flat space is stable to SNP. Since BPS branes are solutions of string/M-theory in a flat background, there is no reason to expect that they will destabilize flat-space.

The SUSY breaking superpotential should be of order

$$\delta W \sim \frac{m_3}{2} M_P^2,$$

(45)

and the resulting $F$ term is of order

$$\delta F \sim \frac{m_3}{2} M_P,$$

(46)

which is not expected to destabilize a minimum with curvature of the order of $M_P^2$, since the new location of the minimum is shifted by a small amount proportional to $m_3/2$.

We argue that the scenario that we are proposing will not suffer from the problems that the other scenarios were inflicted with. First, a minimum with broken SUSY as in case (ii) (eq.35) can be found, since the arguments against such a possibility depended on the steepness of the superpotential, but in the current situation the superpotential is not a steep function. The practical cosmological constant problem can be solved provided the superpotential satisfies the following additional condition,

$$(S + S^*)^2 |F|^2 = 3 |W|^2 / M_P^2.$$

(47)

That eq.(47) can be satisfied is clear from comparing eqs. (15) and (16). The problems associated with the multi-minima structure become as benign as possible, since the height of the barrier between the SUSY breaking minimum around the self-dual point and the
SUSY preserving minimum at infinity is of order $M^4_P$ and its width is of order $M_P$, which are the best values that we can hope for.

The above arguments are not a proof that our proposed scenario works, but they show that it is plausible, and that it may not suffer from the problems that existing ideas for moduli stabilization and SUSY breaking.

Previously, Kaplunovsky and Louis [40] have proposed in the context of $F$-theory a scenario which has some of the ingredients that we are proposing, namely, stabilization at a high scale and SUSY breaking at a lower scale. However, as pointed out in [6] their proposal suffered from drawback that a field theory argument is inconsistent at the string/Planck scale. What we are proposing on the other hand is that the moduli are stabilized at the string scale by intrinsically stringy effects while the SUSY breaking could be field theoretic, in the spirit of [2].

V. THE CONSISTENCY OF STRING UNIVERSALITY WITH PHENOMENOLOGY

In this section we would like to show that if indeed moduli are stabilized near the self-dual point, then the phenomenology of the effective low energy theories derived from the various string theories is consistent with each other and with expectations. Of course, this analysis is done using perturbative theories in the boundary region of moduli space, but we hope to demonstrate that approaching the central region from different directions gives a consistent picture of the physics there.

A. The value of $\rho$ in Horava-Witten Theory

In string theories I, HO the point $g = 1$ is the one where one would expect perturbation theory to break down. The phenomenology of these theories is not inconsistent with this value (assuming that the perturbative phenomenology can be extrapolated with suitable assumptions about stabilization of moduli). In particular, with threshold corrections there
is no inconsistency with the value of the Planck mass. On the other hand the HW phenomenology would seem to yield a large value of the eleventh dimension and hence a large value of the string coupling. Let us therefore examine this question in some detail.

In eq.(8) we take \( M_{11} = R^{10} \times S^1 \) and work with fields that are reflection symmetric under \( x^{11} \rightarrow -x^{11} \). In the zero'th order calculation one assumes that \( M^{11} \) can be compactified as a direct product space \( R^4 \times CY_3 \times S^1 \). Defining the volume of \( CY \) \( \int \sqrt{G} \, d^6x = V \) and putting \( \int dx^{11} \sqrt{G_{11}} = 2\pi \rho \) as before we may identify,

\[
G_N = \frac{\kappa_{11}^2}{16\pi^2 V \rho}, \quad \alpha_{GUT} = \frac{(4\pi \kappa_{11}^2)^{2/3}}{2V}.
\] (48)

If we identify \( V = \frac{a^6}{M_{GUT}^3} \) where \( a \sim O(1) \), then we obtain,

\[
\rho = a^3 \frac{(2\alpha_{GUT})^{3/2}}{64\pi^3} \frac{M_P^2}{M_{GUT}^3} \sim a^3 \frac{1.8}{M_{GUT}}
\] (49)

\[
\kappa_{11}^2 = \frac{(2\sqrt{\alpha_{GUT}})^{1/6}}{(4\pi)^{1/6}} = 0.5 a \frac{1}{M_{GUT}}.
\] (50)

In the above \( M_P^2 = G_N^{-1} = 1.2 \times 10^{19} GeV \) and we have put \( \alpha_{GUT} = 1/25, \quad M_{GUT} = 3 \times 10^{16} GeV \). For the low energy M-theory expansion to make sense we should have \( \frac{\kappa_{11}^2}{V} < 1 \) and \( \frac{\kappa_{11}^2}{V^{2/3}} < 1 \). With the above numbers the former is about 0.02 but the latter is 0.3\( a^2 \) so that we must have \( a \sim O(1) \) and seems to rule out values close to \( 2\pi \). We define the eleven dimensional Planck length by the relation, \( 2\kappa_{11}^2 = (2\pi)^8 l_{11}^9 \) as before. This and the radius of the M theory circle are related to the HE string length and coupling constant by eq. (10). So we may summarize this naive comparisons of scales as follows,

\[
\frac{\pi \rho}{2\pi l_{11}} \approx 9a^2
\] (51)

\[
\frac{\pi \rho}{\sqrt[6]{V}/6} \approx 6a^2
\] (52)

\[
\frac{V^{1/6}}{2\pi l_{11}} \approx 1.6.
\] (53)

Thus we see that with \( a \sim 1 \), the extra (eleventh) dimension is about an order of magnitude larger than the other length scales of the theory. However it should be noted that although consistency of the arguments require that \( a \) has the upper bound given above,
there is no lower bound. Indeed the last ratio is independent of \( a \) but larger than one as required.

In order to compare with perturbative HE string theory (see \( \text{[III]} \)) we make the identification \( \rho = g_{HE}^{2/3} l_{HE} \). Then we get from the above \( g_{HE} = 76a^3 \) which must certainly be considered a value that is in the strong coupling region if \( a \sim 1 \). However, \( a \) is cubed in this relation and with \( a \sim 0.25 \) we would obtain \( g \sim O(1) \). It is not unreasonable to expect a numerical coefficient of this magnitude in the relation between the observed unification scale and the Kaluza-Klein scale, but we will presently argue that if one also takes into account threshold effects one can indeed get values of \( g \sim 1 \) even with \( a \sim 1 \).

If the phenomenology does indeed require large values of \( \pi \rho \) or equivalently of \( g \), then the idea of string universality is not viable. But then we would have to say that the HE/HW theory is special and is the only one that gives the low energy energy physics of the real world. The evolution of the physical couplings would be as discussed in \( \text{[I]} \) and illustrated in Fig 18.2 of \( \text{[I]} \) according to which a fifth dimension opens up a little above \( 10^{15} GeV \) so that the dimensionless gravitational coupling constant starts evolving as in a 5D theory to meet the other three couplings at the GUT scale. So as one increases the energy the world appears to go from being 4D to 5D and finally to 11D. Such a picture clearly cannot hold in any other string theory where one should only see a transition from 4D to 10D at the string scale of around \( 10^{18} \). In this picture the dimensionless gravitational coupling does not meet the other three at the GUT scale, one merely has a determination of its value at that scale.

Both these scenarios cannot be true, and one has to pick one over the other. Such a situation seems very strange to us. It would mean that the other string theories have no role to play in nature. The existence of duality relation among them and with the HE theory however seems to argue against such an interpretation. It seems much more natural to have a situation of string universality as discussed in the introduction. Let us therefore reexamine the phenomenology of the HE/HW theory. One possibility was mentioned above, i.e. one may have \( a < 1 \). Below we will argue that if threshold corrections are included in the relations \( \text{[RS]} \) then it is possible (for non-standard embeddings) to obtain a coupling
\( g \sim O(1) \) so that in the equivalent HW picture one has a eleventh (fifth) dimension whose size is of the order of the string length \( l_{11} \sim l_{HE} \).

In the string theory this is of the order of stringy fluctuations and does not have the interpretation of an extra dimension. In other words the picture that emerges is that of an intermediate coupling HE/HW theory that lies at the boundary of the region accessible to both. The coupling unification picture that will emerge from this analysis will be the same as that which comes out of any of the other phenomenologically viable string constructions.

It is easiest to discuss the threshold corrections from the strong coupling HW end following Witten’s arguments [8] (for reviews of work done since this original paper see [10,11,23]). It should however be stressed that this is completely equivalent to the weak coupling string calculation extrapolated to strong (or at least intermediate) coupling, and for large compactification volume. [10]

As pointed out in [8] the volume of the CY space depends on the value of \( x^{11} \). Thus we have

\[
V_H \equiv V(\pi \rho) = V_O + 2\pi^2 \rho \left( \frac{\kappa}{4\pi} \right)^{2/3} \int_{X_O} \omega \wedge (\text{tr} F \wedge F - \frac{1}{2} \text{tr} R \wedge R) \bigg|_{X_O} \equiv V(0). \tag{54}
\]

Calling one of the walls the observable one with CY volume \( V_O \) and the other one the hidden wall with CY volume \( V_H \) we may write

\[
V_{O,H} = V(1 \mp \epsilon), \tag{55}
\]

where the threshold correction \( \epsilon \), is given by the integral over the CY space \( X_O \) at the observable wall,

\[
\epsilon = \frac{2\pi^2 \rho}{2V} \left( \frac{\kappa}{4\pi} \right)^{2/3} \frac{1}{8\pi^2} \int_{X_O} \omega \wedge (\text{tr} F \wedge F - \frac{1}{2} \text{tr} R \wedge R). \tag{56}
\]

and \( V = \langle V \rangle = \frac{V_O + V_H}{2} \). \( \epsilon \) is negative for the standard embedding but may be positive or negative for non-standard embeddings. Also the requirement that \( V_{O,H}, V > 0 \) implies that \( |\epsilon| < 1 \).
Now we identify $\alpha_O = \alpha_{GUT} \simeq \frac{1}{25}$ and $V_O^{1/6} = \frac{a}{M_{GUT}}$. But the important point is that the expression for Newton’s constant involves the average volume $V$. Redoing the calculations leading to (49,50) by taking into account this variation of the volume of the CY space, we get

$$\rho = \frac{(2\alpha_O V_O)^{3/2}}{64\pi^3 G_N V_O} (1 - \epsilon) = a^3 1.8 M_{GUT}^{-1} (1 - \epsilon)$$

$$k_{11}^{2/9} = \frac{(2V_O \alpha_O)^{1/6}}{(4\pi)^{1/9}} = 0.5 \frac{a}{M_{GUT}}.$$

Then the relations of (51,52) acquire a factor $(1 - \epsilon)$ on their right hand sides so that even if $a = 1$, for a non-standard embedding with $(1 - \epsilon_O) \sim O(10^{-1})$ we would get an eleventh dimension whose size is of the order of the eleven dimensional Planck scale (or the string scale) and hence in the ten dimensional theory a coupling $g$ of order unity.

We should stress again that this calculation was by no means an attempt to show that $g_{HE} = 1$. It is merely an argument to demonstrate that the conclusion that $g_{HE} >> 1$ is unwarranted. After all we have been arguing that along with $g = 1$ the $T$ moduli should be stabilized close to the string scale (which is the same as $l_{11}$ if $g = 1$), but the above argument was a large volume one. This is in the spirit of our whole discussion where we approach the the “middle of moduli space” from different (computable) directions to get hints on the nature of this region.

### B. Type I/IIB orientifold compactifications and Brane worlds

Recently there has been much excitement about the possibility that the string scale is close to the weak scale ($\sim 1$ TeV) [13][14]. Within the space of possible string theoretic formal constructs the sort of situation envisaged in [13] may be modeled by compactifying type I string theory on a 6 torus (or orbifold) and T-dualizing in all 6 compact directions. The resulting theory is a type IIB orientifold with $2^6$ orientifold planes and 32 D3 branes.

The compactified type I theory has the following 4d terms,

$$\Gamma_1 = \frac{V_6}{(2\pi)^7 g^{2/5} l_6^6} \int d^4 x \sqrt{G_4 R} + \frac{V_6}{4(2\pi)^7 g l_6^6} \int d^4 x \sqrt{G_4 \text{tr} F^2}.$$

38
In this picture we write the compactification volume as \( V_6 = (2\pi R)^6 \) and the theory has Kaluza-Klein (KK) modes and winding modes with masses \( M_{KK} = \frac{n}{R} \) \( n \in \mathbb{Z} \) and \( M_w = \frac{\omega R}{l_I^2} \) \( w \in \mathbb{Z} \) respectively. For simplicity we have taken all radii equal and we have only kept the constant mode of the dilaton.

The T-dual effective action is obtained by the transformations \( R \rightarrow R' = \frac{l_I^2}{R} \) and \( g \rightarrow g' = \left( \frac{l_I^2}{R} \right)^6 g \). We get

\[
\Gamma' = \frac{V'_6}{(2\pi)^6 l_I^6 g'^2} \int d^4x \sqrt{G}R + \frac{1}{2\pi g'} \int d^4x \frac{1}{4} \text{tr} F^2.
\]

(59)

It should be noted that the second term is just a 4 dimensional integral because the D9 brane of type I has been transformed into a D3 brane and \( V' = (2\pi R')^6 \). The two pictures above must be physically equivalent as string theories though not necessarily as low energy field theories. For instance even though in the second (D3 brane) picture the gauge theory modes are just four dimensional while in the first picture they are ten dimensional it should be recalled that the KK modes in the first picture get replaced by winding modes in the second.

From the type I picture we have

\[
G_N = \frac{1}{8} l_I^8 g^2, \quad \alpha_{YM} = \frac{1}{2} \frac{l_I^6}{R^6} g.
\]

(60)

Now if the string scale is as low as 1 TeV then we would find from the above \( g \simeq 10^{-30} \) which is unnaturally small. On the other hand putting \( \alpha_{YM} \sim O(10^{-2}) \) we have \( \frac{R}{l_I} \simeq 10^{-5} \).

But when the compactification scale is smaller than the string scale this picture should not be used and the physics should be analyzed in the T-dual situation. In this case we have,

\[
G_N = \frac{1}{8} l_I^8 g'^2, \quad \alpha_{YM} = \frac{1}{2} g'.
\]

(61)

Now \( \frac{R}{l_I} \sim 10^5 \) and it would be very difficult to understand how such large volume stabilizations could be achieved\(^\text{21}\). Also the string coupling, although not unnaturally small, is

\(^\text{21}\)It is perhaps also worth pointing out that using a warp fact or in the metric with the transverse
still $O(10^{-1})$ and one would expect that perturbation theory is valid and it is difficult to understand why the dilaton should be stabilized.

Note that the conclusion that $g' = \frac{1}{2} \alpha_Y M$ and is therefore small in the second picture, is independent of the string scale size or the compactification scale since these factors do not enter into the above relation. In other words regardless of the size of the string scale the IIB orientifold picture leads to a weakly coupled string theory in which it is hard to understand how the dilaton is stabilized. To avoid this situation we must argue that the original picture is the correct one to use with $g \sim O(1)$ and $R'_I \sim (2\alpha_Y M)^{-1/6} \sim 1.5$ so that it is still reasonable to argue that the compactification scale is close to the string scale and is thus presumably stabilized by some SNP mechanism. It is not meaningful to go to the D3 brane picture in this case since there $R' < l_I$ and the original picture is the appropriate one.

The picture of string universality that we have presented would then require that the region of moduli space where the real world lies cannot be approached from this particular part of the boundary of moduli space. This is not too surprising in any case. The theory given in (59) is not of the same status as the original 10D theories. It can be obtained by compactifying the type I theory on a torus (or orbifold) and then T-dualizing. Now according to our hypothesis if the original type I theory has a compactification that can approximate the real world in the central region of moduli space, just as its S-dual heterotic theory can from the another side, this compactification is likely to be a point in the moduli space of Ricci flat Kahler manifolds that is more complicated than a torus or a orbifold and this may not necessarily be connected directly to the region of moduli space that yields the type IIB orientifold. In other words a boundary region theory of IIB orientifolds with D-branes may be connected to the central region that we are interested in only via type I coordinates allowed to take arbitrarily large values or even be non-compact as is done in [41] does not solve this hierarchy problem.
theories to which they connected through torus/orbifold compactifications and T-dualities. Otherwise we run into the problem of explaining why the moduli are stabilized with small 10 D string coupling.

Now the scenario that we have considered here for a brane world picture is a fairly conventional one and as we’ve argued above it is hard to see how such a picture could work in that it requires moduli to be stabilized in the weak coupling region. An alternative possibility is to have anti-branes as well as branes at orbifold fixed points. Such an analysis has been carried out in [42,43]. In this scenario SUSY is broken at the string scale which therefore has to be taken to be an intermediate one. The authors have argued that there is possibly a mechanism for stabilizing the moduli even at weak coupling. Even if this were the case these models generically have a cosmological constant at the string scale (there is no mechanism for generating a small number in the theory and since everything is in principle calculable there is no room for fine tuning either), and therefore would not solve the practical cosmological constant problem.

It might be that the brane-world scenario requires a non-standard compactification such as those which lead to gauged supergravity. This is an issue which needs to be investigated further.

VI. SUMMARY AND CONCLUSIONS

Practically all the calculations in string theory have been done in the framework of one or other supersymmetric weakly coupled theory in the boundary region of moduli space. The problem that needs to be addressed is what if anything these calculations tell us about the real world which has:

- Four large space-time dimensions
- No (or broken) supersymmetry.
- Vanishing, or very small, cosmological constant.
• A weak scale hierarchically smaller than the Planck scale $M_P/M_Z \sim 10^{16}$, and perhaps comparable to the size of soft supersymmetry breaking terms in the low energy theory.

• Small gauge couplings.

A fundamental theory ought to be able to explain these major features of particle physics phenomenology. Unfortunately, at this point we have no idea how string theory might explain them. As we have reviewed at some length in the introduction, the source of the problem appears to be that the true vacuum of string theory is in the central region of moduli space where there are no calculational techniques available.

Given this state of affairs, we have proposed a notion of string universality which may provide a way of extracting information about the central region from the information that we have in the boundary region. The idea is based on the several assumptions. Of our assumptions we believe that the first three are generally accepted by string theorists:

(i) One underlying theory (M-theory?) exists, and has a connected moduli space.

(ii) The moduli space of this theory has boundary regions where the theory approaches the various weakly coupled 10D string theories and 11D supergravity.

(iii) Non-perturbative effects in the theory generate a mechanism for stabilizing moduli, and thus fixing all parameters. This provides a unique prediction of where the true vacuum associated with the real world is situated in moduli space.

To these three assumptions we have added two more:

(iv) The region in which moduli are stabilized is the central region in which the string coupling is of order unity and the compactification scale is of order the string scale. This is motivated by the duality relations between the different corners of moduli space and it is
the region that is equally far in some sense from these boundary regions\textsuperscript{22}. 

(v) Some of the physics in this central region can be extracted by approaching it from different calculable directions and identifying the intersection of compatible predictions from these extrapolations.

In section 1 and 2 of this paper we have discussed these assumptions at some length. In section 3 we argued that the origin of string non-perturbative effects are BPS brane instantons, estimated their actions and discussed their dualities. In Section 4 we highlighted the problems associated with currently popular scenarios for moduli stabilization and SUSY breaking, and proposed as an alternative that the two issues be decoupled. We propose that moduli are all stabilized at or near the self-dual point by string non-perturbative effects while SUSY breaking happens at a much lower scale perhaps as a result of field theoretic non-perturbative effects. In section 5 we have showed that stabilization near the self dual point can be made consistent with phenomenology.

We have concluded that moduli stabilization cannot occur when either the inverse coupling or the volume are parametrically large, and therefore at the very least, it is inconsistent to consider, in “race-track” models, only $T$ and $S$ moduli, and parametrize SUSY breaking in terms of a complex vector in $(F_T, F_S)$ plane. We believe that this is a very general conclusion which shows how hard it is to make “race-track” models work.

We have concluded that moduli stabilization around the self-dual point evades the problems of other scenarios. However, although our ideas are not in conflict with phenomenology we have not yet subjected them to a stringent test by constructing a concrete model and verifying in a quantitative way that they are consistent.

\textsuperscript{22}This assumption is similar to that made independently by G. Veneziano \cite{Veneziano}.  

43
Finally, we would like to emphasize that we have presented a new idea: String Universality, which then leads naturally to many consequences, of which we have explored only some. Perhaps the most important consequence of String Universality is that the true vacuum of string theory is around the self-dual point with string coupling of order unity and compact volume of the order of, but below the string volume. We now need to confront this and find new methods and models to describe the physics of string theory.

VII. ACKNOWLEDGMENTS

We wish to thank D. Eichler and G. Kane for discussions, and M. Dine and Y. Shadmi for comments on the manuscript. We thank the CERN TH group for hospitality during the time this project was initiated. This work is partially supported by the Department of Energy contract No. DE-FG02-91-ER-40672.
REFERENCES

[1] J. Polchinski "String Theory", Vol I and II Cambridge 1998.

[2] M. Dine, Y. Nir and Y. Shadmi, “Enhanced symmetries and the ground state of string theory,” Phys. Lett. B438, 61 (1998) [hep-th/9806124].

[3] M. Dine, L. Randall and S. Thomas, “Supersymmetry breaking in the early universe,” Phys. Rev. Lett. 75, 398 (1995) [hep-ph/9503303].

[4] G. Veneziano, “Is the vacuum self-dual?”, unpublished.

[5] T. Banks and M. Dine, “Couplings and Scales in Strongly Coupled Heterotic String Theory,” Nucl. Phys. B479, 173 (1996) [hep-th/9605136].

[6] M. Dine, “Some Reflections on Moduli, Their Stabilization and Cosmology,” [hep-th/0001157].

[7] T. Banks and M. Dine, “Coping with strongly coupled string theory,” Phys. Rev. D50, 7454 (1994) [hep-th/9406132].

[8] E. Witten, “Strong Coupling Expansion Of Calabi-Yau Compactification,” Nucl. Phys. B471, 135 (1996) [hep-th/9602070].

[9] P. Horava and E. Witten, “Heterotic and type I string dynamics from eleven dimensions,” Nucl. Phys. B460, 506 (1996) [hep-th/9510209]; P. Horava and E. Witten, “Eleven-Dimensional Supergravity on a Manifold with Boundary,” Nucl. Phys. B475, 94 (1996) [hep-th/9603142].

[10] H.P. Nilles, “On the Low-energy limit of string and M theory,” Lectures given at Theoretical Advanced Study Institute in Elementary Particle Physics (TASI 97): Supersymmetry, Supergravity and Supercolliders, Boulder, CO, 1-7 Jun 1997.

[11] B.A. Ovrut “N=1 Supersymmetric Vacua in Heterotic M-Theory” Lectures presented at the APCTP Third Winter School on "Duality in Fields and Strings", February, 1999,
Cheju Island, Korea, hep-th/9905115.

[12] H. P. Nilles and S. Stieberger, “String unification, universal one-loop corrections and strongly coupled heterotic string theory,” Nucl. Phys. B499, 3 (1997) [hep-th/9702110].

[13] J. D. Lykken, “Weak Scale Superstrings,” Phys. Rev. D54, 3693 (1996) [hep-th/9603133].

[14] I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. Dvali, “New dimensions at a millimeter to a Fermi and superstrings at a TeV,” Phys. Lett. B436, 257 (1998) [hep-ph/9804398].

[15] J. P. Derendinger, L. E. Ibanez and H. P. Nilles, “On The Low-Energy D = 4, N=1 Supergravity Theory Extracted From The D = 10, N=1 Superstring,” Phys. Lett. B155, 65 (1985).

[16] M. Dine, R. Rohm, N. Seiberg and E. Witten, “Gluino Condensation In Superstring Models,” Phys. Lett. B156, 55 (1985).

[17] N. V. Krasnikov, “On Supersymmetry Breaking In Superstring Theories,” Phys. Lett. B193 (1987) 37.

[18] T. R. Taylor, “Dilaton, Gaugino Condensation And Supersymmetry Breaking,” Phys. Lett. B252, 59 (1990); B. de Carlos, J. A. Casas and C. Munoz, “Supersymmetry breaking and determination of the unification gauge coupling constant in string theories,” Nucl. Phys. B399, 623 (1993) [hep-th/9204012].

[19] R. Rohm and E. Witten, “The Antisymmetric Tensor Field In Superstring Theory,” Annals Phys. 170, 454 (1986).

[20] R. Brustein and P. J. Steinhardt, “Challenges for superstring cosmology,” Phys. Lett. B302, 196 (1993) [hep-th/9212049].
[21] E. Cremmer, S. Ferrara, C. Kounnas and D. V. Nanopoulos, “Naturally Vanishing Cosmological Constant In N=1 Supergravity,” Phys. Lett. B133, 61 (1983).

[22] T. Barreiro, B. de Carlos and E. J. Copeland, “On non-perturbative corrections to the Kaehler potential,” Phys. Rev. D57, 7354 (1998) [hep-ph/9712443].

[23] C. Munoz, “Effective supergravity from heterotic M-theory and its phenomenological implications,” Corfu Summer Institute on Elementary Particle Physics, (1998) hep-th/9906152.

[24] S. Weinberg, “The Cosmological Constant Problem,” Rev. Mod. Phys. 61, 1 (1989).

[25] E. Witten, “String theory dynamics in various dimensions,” Nucl. Phys. B443 (1995) 85 hep-th/9503124.

[26] A. Sen, “Non-BPS States and Branes in String Theory”, APCTP winter school lectures hep-th/9904207; A. Lerda and R. Russo, “Stable non-BPS states in string theory: a pedagogical review”, hep-th/9905000.

[27] J. A. Harvey, P. Horava and P. Kraus, hep-th/0001143.

[28] K. Choi, H. B. Kim and H. Kim, “Moduli stabilization in heterotic M-theory,” Mod. Phys. Lett. A14, 125 (1999) hep-th/9808122.

[29] E. Cremmer, S. Ferrara, L. Girardello and A. Van Proeyen, “Yang-Mills Theories With Local Supersymmetry: Lagrangian, Transformation Laws And Superhiggs Effect,” Nucl. Phys. B212, 413 (1983).

[30] T. Banks, D. B. Kaplan and A. E. Nelson, “Cosmological implications of dynamical supersymmetry breaking,” Phys. Rev. D49, 779 (1994) hep-ph/9308292.

[31] B. de Carlos, J. A. Casas, F. Quevedo and E. Roulet, “Model independent properties and cosmological implications of the dilaton and moduli sectors of 4-d strings,” Phys.
Lett. B318, 447 (1993) [hep-ph/9308325].

[32] M. Dine and Y. Shirman, “Remarks on the racetrack scheme,” [hep-th/9906246].

[33] T. Barreiro, B. de Carlos and E. J. Copeland, “Stabilizing the dilaton in superstring cosmology,” Phys. Rev. D58, 083513 (1998) [hep-th/9805005].

[34] G. Huey, P. J. Steinhardt, B. A. Ovrut and D. Waldram, “A cosmological mechanism for stabilizing moduli,” [hep-th/0001112].

[35] V. S. Kaplunovsky and J. Louis, “Model independent analysis of soft terms in effective supergravity and in string theory,” Phys. Lett. B306, 269 (1993) [hep-th/9303040].

[36] A. Brignole, L. E. Ibanez and C. Munoz, “Towards a theory of soft terms for the supersymmetric Standard Model,” Nucl. Phys. B422, 125 (1994) [hep-ph/9308271].

[37] R. Brustein, “The Role of the superstring dilaton in cosmology and particle physics,” 29th Rencontres de Moriond: Electroweak Interactions and Unified Theories, Les Arcs, France, 12-19 Mar 1994, [hep-th/9405066].

[38] J. A. Casas, “The generalized dilaton supersymmetry breaking scenario,” Phys. Lett. B384, 103 (1996) [hep-th/9605180].

[39] R. Brustein and S. P. de Alwis, in preparation.

[40] V. Kaplunovsky and J. Louis, “Phenomenological aspects of F-theory,” Phys. Lett. B417, 45 (1998) [hep-th/9708049].

[41] L. Randall and R. Sundrum, “A Large Mass Hierarchy from a Small Extra Dimension” Phys. Rev. Lett. 83 (1999) 3370 [hep-ph/9905221].

“An alternative to compactification,” Phys. Rev. Lett. 83, 4690 (1999) [hep-th/9906064].

[42] I. Antoniadis, E. Dudas and A. Sagnotti, “Brane supersymmetry breaking,” Phys. Lett. B464 (1999) 38 [hep-th/9908023].
[43] G. Aldazabal, L. E. Ibanez and F. Quevedo, “Standard-like models with broken supersymmetry from type I string vacua,” JHEP 0001, 031 (2000) [hep-th/9909172].
