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Brane world susy breaking from string/M-theory

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Abstract: String and M-theory realizations of brane world supersymmetry breaking scenarios are considered in which visible sector Standard Model fields are confined on a brane, with hidden sector supersymmetry breaking isolated on a distant brane. In calculable examples with an internal manifold of any volume the Kähler potential generically contains brane-brane non-derivative contact interactions coupling the visible and hidden sectors and is not of the no-scale sequestered form. This leads to non-universal scalar masses and without additional assumptions about flavor symmetries may in general induce dangerous sflavor violation even though the Standard Model and supersymmetry branes are physically separated. Deviations from the sequestered form are dictated by bulk supersymmetry and can in most cases be understood as arising from exchange of bulk supergravity fields between branes or warping of the internal geometry. Unacceptable visible sector tree-level tachyons arise in many models but may be avoided in certain classes of compactifications. Anomaly mediated and gaugino mediated contributions to scalar masses are sub-dominant except in special circumstances such as a flat or AdS pure five-dimensional bulk geometry without bulk vector multiplets.

Keywords: M-Theory, D-branes, Superstring Vacua, Supersymmetry Breaking.

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1. Introduction

The origin of supersymmetry breaking must ultimately be addressed by any supersymmetric theory of nature. A plethora of scenarios for breaking supersymmetry and for the messenger interactions which must transmit the breaking to the Standard Model fields have been proposed. Brane world supersymmetry breaking (BWSB) scenarios utilize the very

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old idea that the Standard Model is confined to a brane in a higher-dimensional space. In this case supersymmetry breaking may be isolated on a distant hidden sector brane which is not in direct contact with the visible sector Standard Model brane. The messenger interactions which couple the visible and hidden sectors then arise from exchange of fields which reside in the bulk of the higher-dimensional space.

In this paper BWSB is investigated in consistent string and M-theory backgrounds which have a geometric interpretation. The general form of visible sector scalar and gaugino masses which arise in such a scenario are addressed. Tree-level scalar masses comparable to the gravitino mass in general arise from exchange of bulk supergravity fields between the visible and hidden sector branes in addition to universal radiative anomaly mediated soft masses. In almost all cases fields in the minimal bulk supergravity multiplet (required by higher-dimensional bulk supersymmetry) are sufficient to generate the brane-brane couplings which lead to tree-level masses. The scalar masses turn out generally to be non universal. Without additional assumptions about flavor symmetries the squark and slepton mass eigenstates then need not be aligned with the quark and lepton mass eigenstates, and dangerous flavor violating processes can result. This implies that geometric separation of the visible and hidden sector branes within extra dimensions alone is not sufficient to solve the supersymmetric flavor problem. The phenomenology of BWSB then is similar to standard supergravity (SUGRA) scenarios with hidden sector breaking.

In any scenario for supersymmetry breaking, the messenger interactions which couple the visible and hidden sectors largely determine the form of the soft supersymmetry breaking terms and in turn the superpartner mass spectrum. In almost any scenario there are irreducible contributions for both scalar and gaugino masses which are related by supersymmetry to anomalous violations of scale invariance. Anomalous one-loop contributions to gaugino masses were noted some time ago [1], but their origin was not fully understood, and it was thought that these contributions were too small to be phenomenologically interesting. Subsequently, Randall and Sundrum [2] and Giudice, Luty, Murayama and Rattazzi [3] provided a theoretical understanding of these masses in terms of the supersymmetric multiplet of anomalies and regulator dependence of the infrared theory. They showed, systematically, that there are not only anomalous contributions to gaugino masses but, at two loops, to scalar masses as well. The one-loop gaugino and two-loop scalar anomaly mediated masses are given by

$$m_g = -b_0 \frac{g^2}{16\pi^2} \tilde{m}_{3/2}\,, \qquad \tilde{m}_q^2 = \frac{1}{2} c_0 b_0 \left(\frac{g^2}{16\pi^2} \right)^2 |m_{3/2}|^2\,, \tag{1.1}$$

where b_0 and c_0 are the leading beta function and anomalous dimension coefficients, respectively, (for vanishing Yukawa couplings), and $m_{3/2}$ is the gravitino mass. These contributions may be thought of as arising from gauge mediation through the ultraviolet regulator fields for which supersymmetry is broken by a non-vanishing auxillary component for the conformal compensator. Further theoretical insight into the anomaly has been provided by the work of [4, 5]. These authors provided a thorough understanding of the nature of the anomaly, and also gave certain conditions under which the one- and two-loop formulas (1.1) are applicable.

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Anomaly mediated contributions to soft masses are loop suppressed and therefore generally unimportant unless tree-level masses vanish or are suppressed by some mechanism. It has been argued that in BWSB, tree-level masses vanish and that anomaly mediation gives the dominant contributions [2]. The argument is based on the assumption that tree-level brane-brane interactions are absent since the visible and hidden sector branes are not in direct physical contact. However, such contact interactions do in fact arise from exchange of bulk fields. And as shown below, in almost every case fields in the minimal bulk supergravity multiplet are sufficient to generate tree-level brane-brane couplings. Such interactions are suppressed by the volume of the internal manifold or separation scale between the branes, and therefore might naively appear unimportant for large volume. However, the four-dimensional gravitino mass, which sets the scale for supersymmetry breaking effects in the low energy theory, is also suppressed by the volume,

$$m_{3/2}^2 \sim \frac{F^2}{V}\,, \tag{1.2}$$

where F is an auxiliary expectation value on the hidden sector brane. So the magnitude of the tree-level scalar masses are not suppressed relative to the gravitino mass. Said another way, from the perspective of the low energy theory, there is no sense in which the visible and hidden sector branes are far apart.

In order to study BWSB we utilize string and M-theory backgrounds which contain separated branes with world volume fields which model the visible and hidden sectors. While many phenomenological issues of supersymmetry breaking are inaccessible with our present understanding of string and M-theories, one feature which can be addressed with present technology is the form of the Kähler potential which couples the visible and hidden sectors. This determines the form of the visible sector soft masses arising from hidden sector supersymmetry breaking. So string and M-theory backgrounds are well suited to address generic features of the squark and slepton spectra, in particular for BWSB.

In the next section we present simple macroscopic arguements based on extended supersymmetry and tree-level inheritance for the form of the lowest order Kähler potential in a variety of brane world backgrounds, including Hořava-Witten models with end of the world branes, D-brane models, and pure five-dimensional supergravity with end of the world branes. This section includes a review of the results presented in [6]. In general the Kähler potentials are not of the sequestered form, although the no-scale form can be inherited at lowest order in special circumstances. In section 3, Kähler potentials for backgrounds with extended supersymmetry are derived from the microscopic point of view within the high energy theory. Interactions between visible and hidden sector branes contained within the Kähler potentials can be understood in this langauge as arising from exchange of bulk supergravity fields. The volume dependence of brane-brane interactions is illustrated in some D-brane examples. In section 4 we discuss the corrections to the Kähler potential which arise in backgrounds with less than maximal supersymmetry from warping of the internal geometry in both D-brane models as well as Hořava-Witten models with end of the world branes. Flavor violating corrections in the latter models are also shown to have implications for standard heterotic string compactifications. In the D-brane models the

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classical bulk warping corrections are quantum effects from the four-dimensional point of view and demonstrate that the Kähler potential is in general not protected in brane world backgrounds. We also discuss the relation between co-dimension one brane world models with an AdS bulk and a dual boundary field theory description. Section 5 presents the role of moduli masses in the low energy theory. In section 6 the problem of gaugino masses in brane world realizations is addressed. The implications of the general form of the Kaher potential on the brane world mechanism of gaugino mediation is considered in section 7. The implications for the supersymmetric flavor problem in BWSB scenarios is summarized in section 8. In appendix A the form of the inherited Kähler potentials for untwisted states in different classes of orbifold compactifications are presented. In appendix B the soft masses which arise from hidden sector supersymmetry breaking with these Kähler potentials are derived. In many cases the extended supersymmetry of the underlying Kähler potential leads to sum rules for scalar masses which imply the existence of unacceptable tree-level tachyons; the conditions for avoiding these are given.

2. Kähler potentials from simple arguments

The form of the couplings between the visible and hidden sector branes determines the visible sector soft terms arising from hidden sector supersymmetry breaking. These couplings may be studied in a number of different frames. In Einstein frame the coefficient of the gravitational Einstein term is by definition independent of all fields and has canonical normalization. In this frame, couplings between the branes which determine visible sector scalar masses reside in the Kähler potential. We therefore focus on the leading form of the Kähler potential which arises in consistent calculable examples.

Before addressing the form of the Kähler potential in Einstein frame it is instructive to consider the so-called supergravity or conformal frame which is particularly convenient for displaying couplings between the branes. This frame is defined by a Weyl rescaling of the metric from the Einstein frame by $g_{\mu\nu}^{SG} = e^{K/3} g_{\mu\nu}^{E}$ where K is the Einstein frame Kähler potential. The auxiliary F-components are also rescaled as $F^{SG} = e^{-K/6} F^{E}$. The supergravity frame has the advantage that the lagrangian for the auxiliary fields closely resembles the analogous expression in global supersymmetry. There, a visible sector scalar field acquires a tree-level soft mass only if it has a direct D-term coupling with a field which breaks supersymmetry. The same is true in local supersymmetry in the supergravity frame. In the local case the relevant supergravity frame lagrangian is [7]

$$\mathcal{L} = \frac{f}{6}\mathcal{R}_4 - \sum_{i\bar{j}} f_{i\bar{j}} \partial_\mu \varphi_i \partial_\mu \varphi_{\bar{j}}^* - \frac{1}{4f}\left(\sum_i f_i \partial_\mu \varphi_i - \text{h.c.}\right)^2 + \cdots + \\ + \sum_{i\bar{j}} f_{i\bar{j}} F_i F_{\bar{j}}^* + |F_\Phi|^2 f + \sum_i (W_i F_i + f_i F_\Phi^* F_i + 3F_\Phi W + \text{h.c.})\,, \tag{2.1}$$

where f is the field dependent supergravity function which multiplies the four-dimensional Einstein term, and W is the superpotential. Subscripts on scalar fields, φ_i, and auxiliary components, F_i, label the different fields, while subscripts on f and W indicate derivatives

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with respect to the corresponding scalar field. F_Φ is the auxiliary component of the conformal compensator superfield. Note that current-current couplings between matter fields proportional to $(1/f)(f_i\partial_\mu\varphi_i - \text{h.c.})^2$ depend on first derivatives of the f function, while non-derivative couplings between matter fields proportional to $f_{i\bar{j}}F_iF_{\bar{j}}^*$ depend on the field dependence of two derivatives of the f function. Specifically, the latter non-derivative couplings give contributions to the soft masses of visible sector fields Q_i from hidden sector auxiliary expectation values F_{Σ_j} of the form

$$\mathcal{L} \supset f_{Q_i\bar{Q}_{\bar{j}}\Sigma_k\bar{\Sigma}_{\bar{l}}} Q_{\bar{j}}^* Q_i F_{\Sigma_{\bar{l}}}^* F_{\Sigma_k} \tag{2.2}$$

and are the analogs of direct D-term couplings between the visible and hidden sector fields in the globally supersymmetric case. For general f both current-current and non-derivative couplings arise between visible and hidden sector fields. Note that since the current-current coupling in supergravity frame between charged fields depends on $(f_i)^2$, its lowest order form already has four matter fields. Knowing the current-current coupling to this order does not then uniquely determine the higher order non-derivative terms (2.2) which contribute to the soft masses in this frame. The supergravity f function and Einstein frame Kähler potential are related by

$$K = -3\ln\left(-\frac{f}{3}\right). \tag{2.3}$$

A special class of supergravity f functions is of the separable form

$$f(T_i, Q_i, \Sigma_i) = f_{\text{vis}}(Q_i) + f_{\text{hid}}(\Sigma_i) - f_{\text{mod}}\left(T_i + T_i^\dagger\right), \tag{2.4}$$

where Q_i and Σ_i are visible and hidden sector fields, respectively, and T_i are moduli. The imaginary components of moduli often transform non linearly under Peccei-Quinn symmetries. Tree-level invariance under these symmetries then fixes the moduli functional dependence to be $T_i + T_i^\dagger$ in the classical supergravity f function. With a non-vanishing auxiliary expectation value for either hidden sector fields, F_{Σ_i}, or moduli, F_{T_i}, visible sector scalar masses vanish for a supergravity f function of the separable form (2.4). If f were the Kähler potential of global supersymmetry this statement would be obvious since with the separable form there are no non-derivative couplings between the sectors from the $f_{i\bar{j}}F_iF_{\bar{j}}^*$ terms. In supergravity this result is still true even though the different sectors are indirectly coupled through the conformal compensator auxiliary field $F_\Phi \sim m_{3/2}$. This follows from (2.1) with the separable form (2.4) after integrating out F_{Q_i} because there is a cancellation in the expression for the visible sector soft masses between four terms each proportional $|F_\Phi|^2$. This also follows from the fact that at tree-level there is a basis where the conformal compensator only couples to operators in f which are not bilinear in fields, and to operators in the superpotential which are not tri-linear in fields [3]. Note that with the separable form (2.4) current-current interactions between visible and hidden sector fields do exist, even though non-derivative couplings are absent. So even in this special case, the visible and hidden sectors are not actually decoupled.

A supergravity f function of the separable form (2.4) has been referred to as "sequestered". The Kähler potential in Einstein frame associated with this sequestered form is of the no-scale type. With canonical tree-level kinetic terms for the visible and hidden

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sector fields, $f_{\rm vis} = 3\,{\rm tr}\, Q_i^\dagger Q_i$ and $f_{\rm hid} = 3\,{\rm tr}\, \Sigma_i^\dagger \Sigma_i$, and only a single modulus T, the no-scale Kähler potential is

$$K = -3 \ln \left(f_{\rm mod} \left(T + T^\dagger \right) - {\rm tr}\, Q_i^\dagger Q_i - {\rm tr}\, \Sigma_i^\dagger \Sigma_i \right) . \tag{2.5}$$

A non-vanishing auxiliary component for either a hidden sector field, F_{Σ_i}, or modulus, F_T, does not give rise to visible sector scalar masses. In Einstein frame, this vanishing of visible sector scalar masses from (2.5) involves a seemingly miraculous cancellation depending crucially on the functional form of the logarithm and prefactor of "3". However, as discussed above, this follows from the separable form of the supergravity f function in supergravity frame.

For the question of visible sector scalar masses in BWSB, an understanding of the form of the supergravity f function in supergravity frame, or equivalently the Kähler potential in Einstein frame, is clearly crucial. In the brane world scenario, visible and hidden sector fields reside on physically separated branes. It has been argued [2, 8] that the separable form (2.4) might then appear naturally since visible and hidden sector fields are not in direct physical contact in the microscopic theory. In the following subsections we investigate the form of the supergravity f function or equivalently the Kähler potential which arises in string and M-theory models of BWSB. Even at the leading order the no-scale sequestered form is not generally obtained, and unsuppressed non-universal tree-level visible sector scalar masses result from hidden sector supersymmetry breaking. The lowest order supergravity f function is generally not of the separable form and gives unsuppressed non-derivative couplings between branes, even though the branes are physically separated. We also give a simple argument for the form of the leading Kähler potential arising from pure five-dimensional supergravity with end of the world branes. Some corrections to the Kähler potentials are discussed in section 4.

2.1 Hořava-Witten theory

The first brane world model was the Hořava-Witten compactification of M-theory on an $S^1/\mathbb{Z}_2$ interval [9]. In order to cancel gravitational anomalies, E_8 gauge degrees of freedom must be introduced at the two fixed points which bound the interval. The E_8 gauge supermultiplets then reside on end of the world branes separated by a distance R_{11}, and may be thought of as twisted states of this M-theory background. These end of the world branes may be identified with the visible and hidden sector branes. In the limit $R_{11} \ll \ell_{11}$, where ℓ_{11} is the eleven-dimensional Planck length, the Hořava-Witten compactification of M-theory reduces to weakly coupled $E_8 \times E_8$ heterotic string theory in ten dimensions. The strongly coupled limit, $R_{11} \gtrsim \ell_{11}$, provides a realization of the brane world picture.

Compactification of the Hořava-Witten theory to four dimensions on a Calabi-Yau space has been studied extensively [10, 11]. But to explore the structure of four-dimensional Kähler potentials we first consider the simpler case of toroidal compactification on $S^1/\mathbb{Z}_2 \times \mathbb{T}^6$. In this case, all of the states of the $S^1/\mathbb{Z}_2$ Hořava-Witten background, including the gauge supermultiplets, survive compactification to four dimensions. While this is not a realistic phenomenological compactification since it preserves $N = 4$ supersymmetry in

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four dimensions, it demonstrates that even with high degree of symmetry, the sequestered form is not obtained. It also illustrates the origin of certain features of inherited Kähler potentials in more realistic compactifications presented below which preserves only $N=1$ in four dimensions.

Compactification on $S^1/\mathbb{Z}_2 \times \mathbb{T}^6$ can be described by three complex coordinates z_i, $i=1,2,3$. States may be organized in four-dimensional $N=1$ multiplets. At the level of two derivatives the four-dimensional theory has a SU(4) global R-symmetry. In an $N=1$ description only a $\mathrm{SU}(3) \times \mathrm{U}(1)_R$ subgroup is manifest. In this description the SU(3) is a manifest global flavor symmetry of the low-energy theory. The four-dimensional chiral multiplets include geometric moduli $T_{i\bar{j}}$ transforming as $\mathbf{8}_0 \oplus \mathbf{1}_0 \in \mathrm{SU}(3) \times \mathrm{U}(1)_R$ and T_{ij} transforming as $\mathbf{6} \oplus \overline{\mathbf{3}} \in \mathrm{SU}(3) \times \mathrm{U}(1)_R$, and the dilaton S which is invariant. These moduli are related to the geometric parameters of the compactification in eleven-dimensional Planck units by $S+S^\dagger \sim V_6$ and $T \equiv [\det(T_{i\bar{j}}+T_{i\bar{j}})]^{1/3} \sim R_{11}V_6^{1/3}$. The eleven-dimensional gravitational multiplet reduces to the four-dimensional $N=4$ gravitational multiplet plus 6 U(1) $N=4$ vector multiplets. The $E_8 \times E_8'$ gauge multiplets reduce to visible and hidden sector brane fields Q_i and Σ_i transforming as $\mathbf{3}_+ \in \mathrm{SU}(3) \times \mathrm{U}(1)_R$ and $\mathbf{248} \in E_8$, and $\mathbf{248} \in E_8'$, respectively.

The four-dimensional Kähler potential is exact and not renormalized with $N=4$ supersymmetry [12, 13, 14]. The Kähler potential is therefore exact for the case of toroidal compactification. The four-dimensional Kähler potential in the strongly coupled brane world limit, $R_{11} \gtrsim \ell_{11}$, is then identical in this case to that of the weakly coupled heterotic string compactified on $\mathbb{T}^6$. Non renormalization of the Kähler potential can also be obtained by considering a strongly coupled ten-dimensional limit $R_i \gg R_{11} \gtrsim \ell_{11}$, where R_i are the $\mathbb{T}^6$ radii. The Kähler potential is determined by two derivative terms in the low energy theory. It can therefore be obtained from the lowest order ten-dimensional supergravity lagrangian with gauge group $E_8 \times E_8$. This is identical to that obtained in the weakly coupled ten-dimensional limit, $R_i \gg \ell_{11} \gg R_{11}$, since the ten-dimensional supergravity lagrangian is unique at the level of two derivatives. By any of these methods the four-dimensional Kähler potential is [12, 13, 14],

$$K = -\ln\det\left(T_{i\bar{j}} + T_{i\bar{j}}^\dagger - \mathrm{tr}\, Q_i Q_{\bar{j}}^\dagger - \mathrm{tr}\,\Sigma_i \Sigma_{\bar{j}}^\dagger\right) - \ln\left(S + S^\dagger\right), \tag{2.6}$$

where the two traces are over E_8 and E_8' gauge groups, respectively, and the dependence on the T_{ij} moduli is suppressed. Note that this result is explicitly $\mathrm{SU}(3) \times \mathrm{U}(1)_R$ invariant. The microscopic origin of this expression is explained in section 3.1. The supergravity f function associated with the Kähler potential (2.6) is

$$f = -3\left[\left(S + S^\dagger\right)\det\left(T_{i\bar{j}} + T_{i\bar{j}}^\dagger - \mathrm{tr}\, Q_i Q_{\bar{j}}^\dagger - \mathrm{tr}\,\Sigma_i \Sigma_{\bar{j}}^\dagger\right)\right]^{1/3}. \tag{2.7}$$

The Kähler potential (2.6) is not of the no-scale sequestered form and the f function (2.7) is clearly not separable. This is true even ignoring the dilaton. So we see that in this highly symmetric brane world model the sequestered intuition breaks down even without the inclusion of corrections which would generically be present in more realistic models with less supersymmetry.

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The Kähler potential (2.6) or supergravity f function (2.7) imply the existence of non-derivative couplings, proportional to $f_{i\bar{j}}F_iF^*_{\bar{j}}$ in (2.1), between visible and hidden sector fields which reside on the two end of the world branes. These interactions are tree-level and unsuppressed in any way in the four-dimensional theory. The visible sector soft masses arising from hidden sector auxiliary expectation values with the Kähler potential (2.6) are derived in appendix B. Assuming that the moduli are stabilized with vanishing auxiliary components by superpotential interactions and that the cosmological constant vanishes, the eigenvalues of the visible sector scalar mass squared matrix in the $3 \times 3\,\mathrm{SU}(3)$ flavor space in this case are

$$m^2_{Q_i} = m^2_{3/2}(-2, 1, 1)\,. \tag{2.8}$$

These soft masses are proportional to the gravitino mass, $m_{3/2}$, and are not suppressed by the compactification volume or brane separation (relative to the gravitino mass). In addition the tree-level masses are not degenerate. So degeneracy and universality can not be guaranteed simply by separating visible and hidden sectors on different branes. The soft scalar masses squared, in this case, satisfy $\mathrm{Tr}\,m^2 = 0$. As described in appendix B, this special sum rule is the result of the $N = 4$ supersymmetry of the Kähler potential and is not necessarily obtained in more realistic models with less supersymmetry.

Phenomenologically viable models should possess at lowest order only $N = 1$ supersymmetry in four dimensions. The extended supersymmetries of critical string or M-theory must therefore be broken in compactification to four dimensions. This can be achieved by compactifying on a manifold which preserves only 4 supersymmetries. For simplicity we consider compactifications of the Hořava-Witten M-theory background of the form $S^1/\mathbb{Z}_2 \times \mathcal{M}$, where $\mathcal{M}$ is a Calabi-Yau manifold or orbifold, although the conclusions may be applicable to more general compactifications on G_2 manifolds for example.

With only $N = 1$ supersymmetry in four dimensions the lowest order form of the Kähler potential can in general receive corrections. In some cases, however, the corrections to the superpotential and Kähler potential might be argued to be small in the brane world limit, $R_{11} \gtrsim \ell_{11}$. To illustrate this consider first the case in which $\mathcal{M}$ is $K3 \times \mathbb{T}^2$ which preserves 8 supersymmetries which gives $N = 2$ supersymmetry in four dimensions. In this case the Kähler potential is derivable from a holomorphic prepotential. Holomorphy and the classical Peccei-Quinn shift symmetries imply that there are no S or T dependent corrections to the prepotential at any order in perturbation theory, where here T refers generically to any Kähler moduli. However, the shift symmetries allow non-perturbative corrections which are exponentials of functions of $-S$ and $-T$. In the perturbative heterotic string limit these arise from gauge instantons proportional to $\exp(-S)$ and string world sheet instantons wrapping two-cycles proportional to $\exp(h(-S, -T))$, where the function $h(-S, -T)$ depends on the shift symmetries. The perturbative Kähler potential then receives exponentially small corrections in the weakly coupled heterotic string theory limit corresponding to $R_{11} \sim T/S^{1/3} \ll \ell_{11}$, with $S \gg T^3 \gg 1$, where the perturbative string coupling is $g_s^{2/3} \sim R_{11}$. It is then possible to move to the strongly coupled brane world limit, $R_{11} \gg \ell_{11}$ with $T^3 \gg S \gg 1$, while keeping the corrections exponentially small. So the perturbative prepotential receives small corrections also in this strongly coupled

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limit [15]. In this limit the non-perturbative corrections arise from M5-branes wrapping $K3 \times \mathbb{T}^2$ proportional to $\exp(-S)$ and M2-branes wrapping two cycles of $K3 \times \mathbb{T}^2$ and extended in $S^1/\mathbb{Z}_2$ proportional to $\exp(h(-S,-T))$. For $S, T \gg 1$ the lowest order form of the Kähler potential obtained from perturbative heterotic string theory on $K3 \times \mathbb{T}^2$ in this case is then identical, up to exponentially small corrections, to that for M-theory on $S^1/\mathbb{Z}_2 \times K3 \times \mathbb{T}^2$ in the brane world limit. This illustrates that the leading form of the Kähler potential for Hořava-Witten brane world backgrounds can be obtained from weakly coupled heterotic string theory in certain regions of moduli space, at least for compactifications which preserve 8 supersymmetries.

For an $\mathcal{M}$ which preserves only 4 supersymmetries a similar arguement utilyzing holomorphy and symmetries can be made for the lowest order form of the superpotential in a Hořava-Witten brane world limit. The Kähler potential can, however, in principle receive non-holomorphic corrections which are not restricted by holomorphy and symmetries. In fact, continuously connecting the weakly coupled heterotic limit, $R_{11} \ll \ell_{11}$ to the strongly coupled brane world limit, $R_{11} \gg \ell_{11}$, necessarily involves passing through a region in the which the pertrubative string coupling, $g_s^{2/3} \sim R_{11}$, is not small. In this region string non-holomorphic corrections are not suppressed in any way with only 4 supersymmetries. So in this case the weakly and strongly coupled Kähler potentials can not be connected directly by any path through moduli space. However, for the case in which $\mathcal{M}$ is a Calabi-Yau manifold, the Kähler potential at string tree-level to zeroth-order in α' is determined in the large volume limit, $S \gg T^3 \gg 1$ by the classical geometry of the Calabi-Yau manifold. Likewise, in the large volume brane-world limit, $T^3 \gg S \gg 1$ with $S \gg T$, the Kähler potential is also determined by classical geometry up to volume suppressed quantum M-theory corrections. In this region of moduli space $V_6^{1/6} \gg R_{11} \gg \ell_{11}$. Below the scale R_{11}^{-1} the theory is effectively ten-dimensional supergravity on a large Calabi-Yau, and so the Kähler potential is determined by classical geometry. So in this region of moduli space the lowest order form of the Kähler potential obtained from perturbative heterotic string theory on CY is the same, up to power suppressed corrections, as that for M-theory on $S^1/\mathbb{Z}_2 \times \mathrm{CY}$ in the brane world limit.

Now the tree-level Kähler potential for heterotic string theory on a Calabi-Yau manifold to zero-th order in α' and expanding in the fluctuating matter fields takes the general form

$$K = -\ln\left(S + S^\dagger\right) + K\left(T, T^\dagger\right) + Z_{i\bar{j}}\varphi_i\varphi_{\bar{j}}^\dagger + 12 Z_{i\bar{j}k\bar{l}}\varphi_i\varphi_{\bar{j}}^\dagger\varphi_k\varphi_{\bar{l}}^\dagger + \cdots, \tag{2.9}$$

where here T refers generically to any of the $(1,1)$ Kähler moduli. The matter fields $\varphi_i = Q_i$ or Σ_i have wave functions $Z_{i\bar{j}} = Z_{i\bar{j}}(T)$ and quartic couplings $Z_{i\bar{j}k\bar{l}} = Z_{i\bar{j}k\bar{l}}(T)$ which at tree-level are general functions of the $(1,1)$ moduli. For $(2,2)$ Calabi-Yau compactifications with the spin connection embedded in the gauge connection the tree-level form of the Kähler potential is exact to all orders in perturbation theory, while for $(0,2)$ compactifications with torsion general corrections are allowed so that at the perturbative level the matter field wave functions and quartic couplings acquire S dependence, $Z_{i\bar{j}} = Z_{i\bar{j}}(S,T)$ and $Z_{i\bar{j}k\bar{l}} = Z_{i\bar{j}k\bar{l}}(S,T)$. The supergravity f function for the Kähler potential (2.9) contains

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a coupling between the visible and hidden sector fields

$$f \supset \left(S+S^\dagger\right)^{1/3} e^{-K(T,T^\dagger)/3} \left(Z_{Q_i\bar{Q}_{\bar{j}}\Sigma_k\bar{\Sigma}_{\bar{l}}} - \frac{1}{3} Z_{Q_i\bar{Q}_{\bar{j}}} Z_{\Sigma_k\bar{\Sigma}_{\bar{l}}}\right) Q_i Q_{\bar{j}}^\dagger \Sigma_k \Sigma_{\bar{l}}^\dagger \,. \tag{2.10}$$

As long as $Z_{Q_i\bar{Q}_{\bar{j}}\Sigma_k\bar{\Sigma}_{\bar{l}}} \neq \frac{1}{3} Z_{Q_i\bar{Q}_{\bar{j}}} Z_{\Sigma_k\bar{\Sigma}_{\bar{l}}}$ non-derivative couplings exist between the visible and hidden sectors and non-vanishing visible sector soft masses arise from hidden sector supersymmetry breaking. The matter wave function $Z_{i\bar{j}}(T)$ can be calculated to all orders in perturbation theory from the tree-level result for classes of $(2,2)$ compactifications and are generally moduli dependent. The quartic couplings $Z_{i\bar{j}k\bar{l}}(T)$ have not been calculated but should also reasonably be expected to be moduli dependent. The combination $Z_{Q_i\bar{Q}_{\bar{j}}\Sigma_k\bar{\Sigma}_{\bar{l}}} - \frac{1}{3} Z_{Q_i\bar{Q}_{\bar{j}}} Z_{\Sigma_k\bar{\Sigma}_{\bar{l}}}$ is then likely to be non vanish except perhaps at isolated points on moduli space. This should also generally be the case for $(0,2)$ compactifications. So in the region of moduli space in which the corrections to the zeroth-order values of the wave and quartic couplings functions in (2.9) are small in the Hořava-Witten brane world limit, the visible sector soft masses arising from hidden sector supersymmetry breaking are very likely to be non vansihing on a generic Calabi-Yau manifold. Outside these regions of moduli space there will in general be further corrections to these couplings and therefore further corrections to the soft masses. The leading corrections to the zeroth-order Kähler potential in the strongly coupled brane world limit are proportional to $(T+T^\dagger)/(S+S^\dagger)$, and are discussed in section 4.2. These corrections generally violate flavor, so even if the lowest order Kähler potential is flavor conserving, at generic points on moduli space flavor violation can arise.

The Kähler potential simplifies considerably for $(2,2)$ Calabi-Yau compactifications with $h^{1,1}=1$ and therefore only one Kähler modulus T. In this case the Kähler potential including the T modulus, dilaton, and matter fields is fixed by the extra world sheet supersymmetry to be the no-scale form [14]

$$K = -3 \ln\left(T+T^\dagger - \mathrm{tr}\, Q^\dagger Q\right) - \ln\left(S+S^\dagger\right). \tag{2.11}$$

However, since these compactifications have the spin connection embedded entirely in the visible sector gauge connection, there are no hidden sector matter fields. So this class of models is not useful for hidden sector BWSB, but might be applicable to dilaton or moduli dominated scenarios.

Another possibility for $S^1/\mathbb{Z}_2 \times \mathcal{M}$ brane world compactifications of M-theory is for $\mathcal{M}$ an orbifold. The general conditions for consistent singular M-theory backgrounds are not known, although in specific examples evidence for consistent backgrounds with orbifold singularities and possibly M2- and M5-branes can be found [16]. However, for $S^1/\mathbb{Z}_2 \times \mathcal{M}$ compactifications in the weakly coupled heterotic string limit, $R_{11} \ll \ell_{11}$, the consistency conditions based on moduli invariance of the perturbative string description are well established, and reviewed in appendix A. Although we can not demonstrate that a generic consistent orbifold compactification of heterotic string theory lifts to M-theory compactifications, it seems reasonable that this is in fact the case, and that chiral symmetry protects all chiral states which are massless in the heterotic string theory from gaining a mass in the strongly coupled brane world limit. So with these caveats we consider orbifold compactifications of Hořava-Witten theory of this type.

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In an orbifold construction the states which survive in the low energy theory are invariant under the orbifold action. In general this action is non trivial in both compact geometric directions as well as in the gauge group of the underlying theory. In addition, twisted states which reside at orbifold fixed points also appear in the low energy theory. For M-theory backgrounds $S^1/\mathbb{Z}_2 \times \mathcal{M}$, a subset of the $E_8 \times E_8'$ gauge supermultiplets which reside on the end of the world visible and hidden sector branes survive in the low energy theory, $Q_i \subset \mathbf{248}_i \in E_8$ and $\Sigma_i \subset \mathbf{248}_i' \in E_8'$, respectively, where $i = 1, 2, 3$ labels the internal complex coordinates of $\mathcal{M}$. From the weakly coupled heterotic string point of view these fields are untwisted states, and will be referred to as such below. The lowest order tree-level Kähler potential for these states is inherited directly from the $N = 4$ Kähler potential (2.6) by simply removing non-invariant states. This can be obtained from the eleven-dimensional supergravity solution for bulk fields as discussed in section 3.1.1. There are in general additional twisted states in the low energy theory which reside at fixed points of $\mathcal{M}$. The dependence of the Kähler potential on these twisted states is not restricted by extended symmetries since $\mathcal{M}$ preserves only $N = 1$ supersymmetry. Here we focus only on the Q_i and Σ_i visible and hidden sector fields which do inherit a lowest order Kähler potential from the ten-dimensional theory.

For an orbifold which preserves 4 supersymmetries, the inherited Kähler potential for the untwisted states receives corrections even in the weakly coupled heterotic limit. It will also receive additional corrections in the brane world limit. Unlike the case of a Calabi-Yau manifold in which the latter corrections can be studied perturbatively in certain regions of moduli space as described above, the general form of the corrections are unknown in this case. So the lowest order form of the Kähler potentials given below may not literally be applicable to the brane world limit in these backgrounds. However, since corrections are very likely to give yet additional contributions to soft masses, the inherited Kähler potentials are instructive in indirectly illustrating couplings between the visible and hidden sector branes and the associated lowest order irreducible contribution to the scalar masses.

The form of the inherited Kähler potential for an $S^1/\mathbb{Z}_2 \times \mathcal{M}$ compactification which preserves $N = 1$ supersymmetry in four dimensions is restricted by the moduli and gauge quantum numbers of the matter fields. The simplest case arises with $T_{i\bar{j}}$ moduli and three generations of visible and hidden sector matter, Q_i and Σ_i, which then have identical gauge quantum numbers for each i, where $i = 1, 2, 3$ are the complex coordinates of $\mathcal{M}$ and also the flavor index. A simple example of this type with $\mathcal{M}$ a symmetric $\mathbb{Z}_3$ orbifold is given in the appendix A. In this class of compactifications the off-diagonal combinations of fields $Q_i Q_{\bar{j}}^\dagger$ and $\Sigma_i \Sigma_{\bar{j}}^\dagger$ are gauge invariant and can appear in the Kähler potential. With all the $T_{i\bar{j}}$ moduli the inherited Kähler potential and supergravity f function are then identical to (2.6) and (2.7), respectively. Note that the inherited Kähler potential has a SU(3) flavor symmetry in this case. As discussed in appendix B, with any supersymmetry breaking hidden sector auxillary expectation values, $F_{\Sigma_i} \neq 0$, and assuming the moduli are stabilized with vanishing auxillary expectation values, the non-universal tree-level visible sector masses given in (2.8) are obtained.

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Another class of $S^1/\mathbb{Z}_2\times\mathcal{M}$ compactifications have only diagonal $T_{i\bar{i}}$ moduli and visible and hidden sector matter, Q_i and Σ_i, which have different gauge quantum numbers for each i. In this case off-diagonal combinations are not guaranteed to be gauge invariant, and in fact are not if a given representation under the unbroken subgroup arises only once. A simple example of this type with $\mathcal{M}$ a $\mathbb{Z}_6$ orbifold which does not respect any permutation symmetries is given in appendix A. In this class of compactifications, since the off-diagonal combinations $Q_iQ_{\bar{j}}^\dagger$ are not gauge invariant they can not appear in the Kähler potential. The tree-level Kähler potential inherited from (2.6) in this case is a sum of logarithms

$$K=-\sum_i \ln\left(T_i+T_i^\dagger-\mathrm{tr}\,Q_i^\dagger Q_i-\mathrm{tr}\,\Sigma_i^\dagger\Sigma_i\right)-\ln\left(S+S^\dagger\right), \tag{2.12}$$

where the traces are over gauge qauntum numbers and are in general different for each i. This Kähler potential is invariant under an $S_3\times\mathrm{U}(1)_R$ global symmetry. The supergravity f function associated with the Kähler potential (2.12) is

$$f=-3\left[\left(S+S^\dagger\right)\prod_i\left(T_i+T_i^\dagger-\mathrm{tr}\,Q_iQ_i^\dagger-\mathrm{tr}\,\Sigma_i\Sigma_i^\dagger\right)\right]^{1/3}. \tag{2.13}$$

The Kähler potential (2.12) is not of the no-scale sequestered form and the f function (2.13) is not seperable even ignoring the dilaton.

The visible sector scalar masses arising from hidden sector supersymmetry breaking with the Kähler potential (2.12) are discussed in appendix B. Assuming the moduli are stabilized with vanishing auxillary expectation values the scalar masses squared for each i are

$$m_i^2=m_{3/2}^2(1-x_i)\,, \tag{2.14}$$

where $x_i=2\,\mathrm{Re}\,T_i|F_i|^2/|W|^2$. Vanishing of the cosmological constant implies $\sum_i x_i=3$ as shown in appendix B. Again in this case the soft masses are proportional to the gravitino mass and are non universal with the detailed spectrum depending on the x_i. The condition on the x_i in this case implies that the sum of the three mass squared eigenvalues vanish

$$m_1^2+m_2^2+m_3^2=0\,. \tag{2.15}$$

However, since the multiplicities for each i are not necessarily the same, $\mathrm{Tr}\,m^2\neq 0$ in general. The condition (2.15) implies the existence of unacceptable tree-level visible sector tachyons. However, in more realistic compactifications these can be avoided by, for example, projecting out the dangerous states. A special case of the condition (2.15) is obtained for both hidden sector fields with diagonal auxillary expectation values, $F_i=F$, and diagonal moduli expectation values, $T_i=T$. Vanishing cosmological constant and unbroken S_3 then implies $x_i=1$ and the scalar masses vanish, $m_i^2=0$. It is important to note in this case, however, the vanishing masses result from an unbroken S_3 flavor symmetry rather than the form of the Kähler potential.

As a final example consider a class of $S_1/\mathbb{Z}_2\times\mathcal{M}$ orbifold compactifications which have moduli $T_{i\bar{j}}$ for $i=1,2$, and $T_{3\bar{3}}$, and visible and hidden sector chiral matter Q_i, Σ_i. Further, suppose there is a S_2 permutation symmetry for $i=1,2$, so that the low-energy

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theory has two generations in $Q_{i=1,2}$ and $\Sigma_{i=1,2}$ (Q_i and Σ_i are charged under different groups), and states Q_3 and Σ_3 which have different quantum numbers from the first two generations. A class of orbifold $\mathbb{Z}_6$ orbifold examples leading to this spectrum is provided in the appendix A. In this class of compactifications the off-diagonal elements $Q_i^\dagger Q_3$ for $i = 1,2$ are not gauge invariant and do not appear in the Kähler potential. The lowest order inherited tree-level Kähler potential is then

$$K = -\ln \det_{i=1,2}\left(T_{i\bar{j}} + T^\dagger_{i\bar{j}} - \operatorname{tr} Q_i^\dagger Q_{\bar{j}} - \operatorname{tr} \Sigma_i^\dagger \Sigma_{\bar{j}}\right) - \ln\left(T_{3\bar{3}} + T^\dagger_{3\bar{3}} - \operatorname{tr} Q_3^\dagger Q_{\bar{3}} - \operatorname{tr} \Sigma_3^\dagger \Sigma_3\right). \tag{2.16}$$

Assuming that only $T_{i\bar{i}}$ acquire vevs and that all the $T_{i\bar{j}}$ moduli are stabilized with vanishing auxillary components, then with hidden sector supersymmetry breaking the soft masses are

$$m^2 = m^2_{3/2}(1, -2 + x_3, 1 - x_3)\,, \tag{2.17}$$

where $x_3 \equiv 2\,\mathrm{Re}\, T_3 |F_{\Sigma_3}|^2/|W|^2 \leq 3$. It is important to note that the first two states have the same gauge quantum numbers. Inspecting the mass eigenvalues indicates that with hidden sector supersymmetry breaking the two generations are not generically degenerate, and one may be tachyonic. Degeneracy occurs with unbroken or approximate SU(2) flavor symmetry in the hidden sector, and is lifted for $x_3 = 3 - \epsilon$. For small enough ϵ the first two generations are approximately degenerate, but results from an approximate SU(2) flavor symmetry.

So we see that in general without unbroken flavor symmetries tree-level non-universal scalar masses arise from Hořava-Witten BWSB.

2.2 D-branes

Large classes of brane world models can in principle be constructed using perturbative string theory D-branes on compact manifolds. Examples of this type are instructive since, unlike compactifications of the Hořava-Witten theory, the co-dimension of the internal manifold can be larger than one.

In order to illustrate a simple D-brane world model consider first a toroidal compactification of type-I string theory with gauge group SO(32) on $\mathbb{T}^6$. This preserves $N = 4$ supersymmetry in four dimensions, and is therefore of course not phenomenologically realistic but again illustrates that even in this highly symmetric case the sequestered intuition breaks down. In the absence of Wilson lines, the four-dimensional Kähler potential in $N = 1$ notation is determined by $N = 4$ supersymmetry to be

$$K = -\ln\det\left(T_{i\bar{j}} + T^\dagger_{i\bar{j}} - \operatorname{tr} \varphi_i^\dagger \varphi_{\bar{j}}\right) - \ln\left(S^\dagger + S\right), \tag{2.18}$$

where $T_{i\bar{j}}$ are the geometric moduli, φ_i are the $N = 1$ chiral matter fields arising from the compactification of the ten-dimensional gauge supermultiplets, and the trace is over SO(32) indices. This result may also be obtained directly by compactification of the ten-dimensional type-I theory with $T_{i\bar{j}} = g^{\mathrm{I}}_{i\bar{j}}/\lambda_{\mathrm{I}}$ and $S = V_6^{\mathrm{I}}/\lambda_{\mathrm{I}}$, where $g^{\mathrm{I}}_{i\bar{j}}$ and V_6^{I} are the $\mathbb{T}^6$ complex metric and volume in type-I string frame, and λ_{I} is the type-I string coupling. Alternately it can also be obtained from the heterotic SO(32) string theory. This theory

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compactified on $\mathbb{T}^6$ has the Kähler potential (2.18) with $T_{i\bar{j}} = g^h_{i\bar{j}}$ and $S = V^h_6/\lambda^2_h$ where λ_h is the heterotic string coupling. Under type-I-heterotic duality the string couplings and metrics of the ten-dimensional lagrangians are related by $\lambda_h = 1/\lambda_{\rm I}$ and $g^h_{i\bar{j}} = g^{\rm I}_{i\bar{j}}/\lambda_{\rm I}$.

A brane world model in type-II theory with D-branes may be obtained by a T-duality transformation of the type-I theory. Under T-duality on all the $\mathbb{T}^6$ directions

$$R_i \longrightarrow \frac{1}{R_i}\,, \qquad \lambda_I \longrightarrow \frac{\lambda_I}{V}\,. \tag{2.19}$$

The $N = 4$ supersymmetry and the invariance of the low energy theory under this transformation uniquely determines the Kähler potential in the type-I$'$ theory compactified on $\mathbb{T}^6$ to be identical to (2.18).

The type-IIB theory has (including images) 32 D3-branes and 16 O3 orientifold planes in a 1/2 BPS configuration which preserves 16 supersymmetries in the four-dimensional theory. Motions of the D3-branes away from one of the orientifold planes correspond to Wilson lines in the type-I description which break the SO(32) gauge symmetry to a product group. Separating the D3-branes into two groups provides a model of visible and hidden sector branes in the type-I$'$ description. For example, 16 D3-branes (including images) at each orientifold plane gives SO(16)$\times$SO(16) gauge group. In this case, the four-dimensional matter fields break up into visible and hidden sector fields, $\varphi_i = Q_i$ or Σ_i, which reside on each group of D-branes. The Kähler potential (2.18), in this case, then has the same form (2.6) as that of the Hořava-Witten theory as dictated by $N = 4$ supersymmetry. As discussed in the previous subsection, even in this highly symmetric case the Kähler potential is not of the no-scale sequestered form. Supersymmetry breaking isolated in matter fields on the hidden sector brane would in general then lead to unsuppressed tree-level visible sector scalar masses of order the gravitino mass. Note that the branes in this example are co-dimension six, unlike the Hořava-Witten example of the previous subsection in which the end of the world branes are co-dimension one.

The Kähler potential with $N = 4$ supersymmetry is uncorrected. So the inherited Kähler potential (2.18) is exact for compactification of type-I or -I$'$ theories on $\mathbb{T}^6$. However, in a more realistic model compactified on a manifold $\mathcal{M}$ which preserves only $N = 1$ supersymmetry, the inherited Kähler potential is not protected and in general receives corrections in the low energy theory. In the type-I picture there potentially are one-loop quantum corrections proportional to inverse powers of the Wilson line which breaks to the product gauge group. Alternately, in the type-I$'$ picture these one-loop quantum corrections may be understood as arising from integrating out massive string states which stretch between the visible and hidden sector branes. In the closed string channel this one-loop quantum amplitude amounts to tree-level interaction of the visible and hidden sector branes through exchange of bulk closed string states. In the limit of large separation this is dominated by exchange of bulk supergravity fields. The explicit form of Kähler potential corrections in a 1/4 BPS D-brane configuration which preserves 8 supersymmetries is illustrated in section 4.1. For general compactifications with only $N = 1$ supersymmetry, the corrections should be expected to give further contributions to the tree-level masses.

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As in the heterotic Hořava-Witten example, visible sector scalar masses arising from supersymmetry breaking on the hidden sector brane in D-brane realizations of BWSB are generally already present at tree-level, are non universal, and are parameterically of the same order as the gravitino mass, $m_Q^2 \sim m_{3/2}^2$. This would hold generally for D-brane models of any co-dimension.

2.3 Pure five-dimensional supergravity

A BWSB situation which can not at present by analyzed directly within the framework of string or M-theory backgrounds is pure five-dimensional supergravity on an interval with end of the world branes on which the visible and hidden sectors reside. The pure five-dimensional example turns out to be very special for a number of reasons. First, the minimal five-dimensional supergravity multiplet contains only a single gauge field. The only other higher-dimensional case in which this arises is that of eleven dimensions. Geometric compactifications of M-theory generally result in a number of vector and antisymmetric tensor fields as well as scalar moduli and therefore cannot lead to the pure five-dimensional supergravity theory. A more promising possibility for obtaining a five-dimensional BWSB model might be to consider asymmetric orbifold compactifications in type-II string theory since these generally possess far fewer moduli. Still, these compactifications always contain a dilaton. At best one could hope that the pure five-dimensional theory is the strongly coupled description of an asymmetric orbifold compactification to four dimensions, with the dilaton equivalent to the radion of the five-dimensional theory. At present, however, no realization of this possibility has been constructed [17].

The most important feature of pure five-dimensional supergravity is that in compactification of five to four dimensions, the single real radius modulus, R, for the compact volume turns out not to possess a kinetic term. This is not the case for compactification of higher dimensions down to four dimensions. The imaginary component partner of the four-dimensional radius modulus is the periodic Wilson line for the single U(1) gauge boson of the five-dimensional supergravity multiplet. The absence of a kinetic term for the real component, along with the classical Peccei-Quinn symmetry for the imaginary component, implies that the Einstein frame tree-level four-dimensional Kähler potential for the volume modulus without any brane matter is

$$K = -3 \ln \left(T + T^\dagger \right), \tag{2.20}$$

or equivalently that $f_{\text{mod}} = -3(T + T^\dagger)$ in supergravity frame, where $T + T^\dagger = 2R$. As an aside, note also that the supergravity frame, in this case, happens to be the same as the geometric frame obtained by simply dimensionally reducing the five-dimensional theory (or in string theory language the string frame) since $f = -6R$, which again is special to the case of compactification of five to four dimensions.

Compactification of pure minimal five-dimensional supergravity on S^1 gives pure $N = 2$ supergravity in four dimensions. This theory possess an $\mathrm{SL}(2, \mathbb{R})$ non-compact symmetry at the classical level which acts on the radius modulus and implies that the moduli space metric is quaternionic. The Kähler potential (2.20) is consistent with this requirement. For

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the orbifold compactification $S^1/\mathbb{Z}_2$, with projections to give $N = 1$ in four dimensions, the radius modulus (which necessarily survives the orbifold projection) inherits at lowest order the tree-level Kähler potential (2.20).

For a brane world realization, compactification from five to four dimensions on an $S^1/\mathbb{Z}_2$ interval allows the introduction of four-dimensional $N = 1$ chiral multiplet matter fields on end of the world branes at the orbifold fixed points which bound the interval. The end of the world branes may be interpreted as the visible and hidden sector branes. The four-dimensional tree-level Kähler potential including hidden and visible sector brane matter cannot be derived from consistency of the low-energy four-dimensional theory alone. In principle it could be derived microscopically from an underlying string or M-theory background as in the Hořava-Witten and D-brane examples of the previous subsections if a compactification to pure five-dimensional supergravity (plus perhaps a multiplet which contains the dilaton) on the interval $S_1/\mathbb{Z}_2$ were known. However, if an example of this type exists in which the visible and hidden sector brane matter are remnants of $N = 4$ states in the underlying theory, then as described in appendix B, the lowest order tree-level Kähler potential inherited from the $N = 4$ form (2.6) would presumably contain

$$K \supset -3 \ln \left(T + T^\dagger - \mathrm{tr}\, Q_i^\dagger Q_i - \mathrm{tr}\, \Sigma_i^\dagger \Sigma_i \right) . \tag{2.21}$$

This is the no-scale sequestered form of the Kähler potential. The associated supergravity f function, in this case, is of the separable form

$$f = -3 \left(T + T^\dagger - \mathrm{tr}\, Q_i^\dagger Q_i - \mathrm{tr}\, \Sigma_i^\dagger \Sigma_i \right) . \tag{2.22}$$

Even with the separable form (2.22) the branes are not decoupled since there exist current-current couplings in (2.1) proportional to $(1/f)(f_{Q_i}\partial_\mu Q_i - \mathrm{h.c.})(f_{\Sigma_j}\partial_\mu \Sigma_j - \mathrm{h.c.})$ which couple fields on the two branes [8]. As mentioned in section 2, the lowest order form of the current-current couplings in supergravity frame are not sufficient to fix the non-derivative couplings in this frame. The no-scale form can therefore not be derived from pure five-dimensional supergravity with end of the world brane matter by simply matching the lowest order supergravity frame current-current couplings to the effective four-dimensional theory without additional assumptions about the bulk-brane couplings. Here, the Kähler potential (2.21) and f function (2.22) follow at leading order from inheritence of the underlying theory with flat Kähler metric.

With supersymmetry breaking isolated on the hidden sector brane, the no-scale Kähler potential (2.21) gives rise to vanishing tree-level visible sector masses, $m_i^2 = 0$. The separable form for the radion modulus supergravity f function, or equivalently the Kähler potential (2.20) and the presumed inherited form of the Kähler potential with end of the world brane matter (2.21) is a direct result of the fact that the volume modulus for compactification from five to four dimensions does not possess a kinetic term. This occurs only for compactification from five to four dimensions. Therefore, vanishing tree-level soft masses arising with generic hidden sector supersymmetry breaking from the no-scale Kähler potential (2.21) should clearly be thought of as a property of a particular hypothetical model rather than a general feature of brane world supersymmetry breaking.

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With $N = 1$ supersymmetry in four dimensions the Kähler potential is not protected from corrections. It has been argued that with pure five-dimensional supergravity BWSB corrections to the no-scale form of the Kähler potential (2.21) occur only at one-loop and are suppressed by additional powers of the brane separation [2, 8]. This is only true if the brane tensions vanish, which seems a rather strong additional assumption. In the general case there are tree-level corrections to the Kähler potential from bulk warping and therefore to the tree-level soft masses, as discussed in section 4.3. This also occurs in the Hořava-Witten theory compactified on a Calabi-Yau three-fold discussed in section 4.2. In contrast, a brane world model with bulk warping but a sequestered Kahler potential occurs in the supersymmetric Randall-Sundrum model with two branes (and pure anti-de Sitter bulk) [27]. This sequestering can be understood from the AdS/CFT correspondance [30]. We return to this example in section 4, where we speculate that the sequestered form can be spoiled by the presence of additional bulk vectors.

So we conclude from the above simple arguments for the form of Kähler potentials that the sequestered intuition and associated vanishing tree-level scalar masses are not generally realized in BWSB. Anomaly mediated supersymmetry breaking therefore does not seem to be a generic or robust feature of BWSB scenarios, but might be obtained in very special models, or as the result of unbroken flavor symmetries.

3. Locality and the low energy effective action

Couplings between branes which are physically separated within a compact manifold manifest themselves as contact interactions in the low energy theory below the compactification scale. As mentioned in section 1, such interactions do not in fact violate notions of bulk locality since brane-brane interactions can arise already at tree-level from exchange of bulk fields between the branes. We have seen that these brane-brane tree-level interactions in general violate the naive sequestered expectation that the supergravity f function is separable or equivalently that the Kähler potential is of the no-scale form. In this section the microscopic origin of brane-brane interactions from exchange of bulk fields is presented in a number of calculable, controlled examples. We begin with examples with $N = 4$ supersymmetry in four dimensions. The Kähler potentials for both the heterotic and type-I or -I$'$ examples derived in the previous section from general arguments are shown to arise directly from exchange of bulk supergravity fields. The existence of these bulk fields is guaranteed by higher-dimensional supersymmetry. Even though these $N = 4$ examples are not phenomenologically realistic they have the virtue that explicit computations are straightforward, and already the sequestered intuition breaks down at tree-level. As discussed in section 2 these features are inherited by the lowest order tree-level Kähler potential of general $N = 1$ compactifications. Additional tree-level corrections to the brane Kähler potential from exchange of bulk fields in examples with less supersymmetry are discussed in section 4.

3.1 The strongly coupled heterotic theory in the ten-dimensional limit

The microscopic origin of brane-brane interactions from the exchange of bulk fields is well illustrated by the $S^1/\mathbb{Z}_2$ Hořava-Witten M-theory orbifold background. The $S^1/\mathbb{Z}_2$ M-

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theory interval has length R_{11} with E_8 gauge supermultiplets localized on end of the world brane boundaries. Consider first the theory at energy scales below the inverse interval length R_{11}^{-1}. The action of the low energy ten-dimensional supergravity theory coupled to $E_8 \times E_8'$ super Yang-Mills at these scales includes a kinetic term for the NS field strength proportional to

$$\int d^{10}x e^{-2\phi_{10}} H^2 = \int d^{10}x e^{-2\phi_{10}} \left(dB + \omega_{E_8} + \omega_{E_8'} - \omega_L \right)^2 , \tag{3.1}$$

where ω_{E_8} and $\omega_{E_{8'}}$ are the Chern-Simons terms associated with each E_8, and where ω_L is the Lorentz Chern-Simons form. From an eleven-dimensional point of view, E_8 and E_8' lie on different end of the world branes, separated by the distance R_{11}. At low energies this theory must reduce to the ten-dimensional theory above. This low energy theory clearly has tree-level contact interactions which couple the Chern-Simons forms on each brane.

The way in which the brane-brane contact interaction comes about was explained in [18]. In eleven dimensions, a non-vanishing Chern-Simons term on one of the brane serves as a source for the three-index antisymmetric tensor potential, C_{ABC}, in the bulk. This appears as a modification of the Bianchi identity for the four-form field strength G_{ABCD} [9],

$$(dG)_{ABCD11} = \lambda \left(J^{\text{hid}}_{ABCD} \delta(y_{11} - y_{\text{hid}}) + J^{\text{vis}}_{ABCD} \delta(y_{11} - y_{\text{vis}}) \right) , \tag{3.2}$$

where

$$J^i = \text{tr} \left(F^i \wedge F^i \right) - \frac{1}{2} \text{tr}(R \wedge R) , \tag{3.3}$$

and $\lambda \equiv (\kappa_{11}/4\pi)^{2/3}/(2\sqrt{2}\pi^2)$. This implies that G_{ABCD} is in general non vanishing in the bulk [18]

$$G_{ABC11} = 3dB_{[ABC]} + \frac{\lambda}{2} \left(\omega_{E_8} + \omega_{E_8'} - \omega_L \right)_{ABC} , \tag{3.4}$$

where here $x^{11} \in [0, 1]$. Note that the four-form field strength is constant in the bulk between the branes. Now the eleven-dimensional lagrangian contains a kinetic term for the field strength proportional to

$$\int d^{11}x \sqrt{g}\, G_{ABC11} G^{ABC11} . \tag{3.5}$$

Inserting the expression (3.4) for the four-form field strength into this action and integrating over the eleventh dimension gives precisely the ten-dimensional action (3.1).

Two features are important to the form of the brane-brane tree-level interaction in this case. First, and most crucially, fields on the branes are sources for bulk fields. This is a generic feature of the all the examples addressed here. Second, the bulk field strength sourced by the brane fields in this geometry is constant and does not fall with distance. This feature is special to a bulk with co-dimension one, and as described below, is not crucial to the existence of unsuppressed brane-brane interactions in examples with larger co-dimension.

Next consider the further reduction to four dimensions. The Chern-Simons terms discussed above play a crucial role in yielding the Kähler potentials described in the previous

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section. Consider first the special compactification $\mathbb{T}^2 \times \mathbb{T}^2 \times \mathbb{T}^2$ where each $\mathbb{T}^2$ is a symmetric torus with both radii given by R_i for $i = 1, 2, 3$. This compactification preserves a complex structure with complex metric $g_{i\bar{j}}$. After compactification to four dimensions, but before Weyl rescaling to Einstein frame, the four-dimensional kinetic terms for visible and hidden sector scalars, $\varphi_i = Q_i$ or Σ_i which arise from the ten-dimensional gauge boson kinetic terms in string frame are

$$\mathcal{L}_{\rm kin} = \sum_i e^{-2\phi_{10}} \frac{V_6}{R_i^2} \operatorname{tr} |\partial_\mu \varphi_i|^2 \,, \tag{3.6}$$

where $V_6 = \det(g_{i\bar{j}})^{1/2}$ is the volume of the compact space, ϕ_{10} is the ten-dimensional dilaton, and the trace is over the gauge and flavor indices. After Weyl rescaling to the Einstein frame, where the string and Einstein frame metrics are related by

$$g_{\mu\nu}^{ST} = \frac{e^{2\phi_{10}}}{V_6} g_{\mu\nu}^{E} \tag{3.7}$$

the kinetic terms become

$$\mathcal{L}_{\rm kin} = \sum_i \frac{1}{R_i^2} \operatorname{tr} |\partial_\mu \varphi_i|^2 \,. \tag{3.8}$$

In Einstein frame the moduli terms for the radii and four-dimensional dilaton are

$$\mathcal{L}_{\rm rad} = \sum_i \frac{1}{R_i^2} (\partial_\mu R_i)^2 + (\partial\phi)^2 = \sum_i \frac{1}{4R_i^4} \left(\partial_\mu R_i^2\right)^2 + (\partial_\mu \phi)^2 \,, \tag{3.9}$$

where the four-dimensional dilaton is a combination of the ten-dimensional dilaton and the volume modulus

$$\phi = \phi_{10} - \frac{1}{4} \ln\left(\det g_{i\bar{j}}\right) = \phi_{10} - \frac{1}{2} \ln V_6 \,. \tag{3.10}$$

In addition, the four-dimensional Einstein frame lagrangian contains a term coming from the reduction of the Chern-Simons squared terms (3.1) which arise from the eleven-dimensional theory by integrating out $G_{\mu ab11}$,

$$\mathcal{L}_{\rm cs} = \sum_i \frac{1}{R_i^4} \left(\partial_\mu a_i - \frac{i}{\sqrt{2}} \operatorname{tr}(\varphi_i^* \partial_\mu \varphi_i - \varphi_i \partial_\mu \varphi_i^*) \right)^2 \,, \tag{3.11}$$

where a_i is the pseudoscalar partner of radii, $a_i = B_{i\bar{i}} = C_{11 i\bar{i}} \equiv \operatorname{Im}(T_i)$. The trace is over the gauge degrees of freedom which includes a sum over *both* visible and hidden sector fields, $\varphi_i = Q_i$ or Σ_i, which from an eleven-dimensional point of view reside on different boundaries. Note that the radii, R_i, do not have derivative couplings to the brane matter fields.

Now in order to write the entire action in a manifestly supersymmetric form in terms of $N = 1$ supermultiplets, it is necessary to define the scalar component of the chiral multiplet fields T_i which contain the radii moduli as [19]

$$T_i = R_i^2 + i a_i + \frac{1}{2} \operatorname{tr} Q_i^* Q_i + \frac{1}{2} \operatorname{tr} \Sigma_i^* \Sigma_i \,, \tag{3.12}$$

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and the scalar component of the dilaton chiral multiplet as

$$S = e^{-2\phi} + i\sigma\,, \tag{3.13}$$

where σ is the model independent axion. With these definitions, the Kähler potential for these four-dimensional $N = 1$ chiral multiplets is

$$K = -\ln\left(S + S^\dagger\right) - \sum_i \ln\left(T_i + T_i^\dagger - \mathrm{tr}\,\Sigma_i^\dagger\Sigma_i - \mathrm{tr}\,Q_i^\dagger Q_i\right). \tag{3.14}$$

The appearance of the logarithm in (3.14) might at first sight seem puzzling since this non-linear function might be expected to generate, first, an infinite series of high dimension couplings between the brane fields, and second, derivative couplings between R_i and the brane fields. In contrast, the classical ten-dimensional lagrangian discussed above has operators which involve at most six fields (arising from the Chern-Simons squared interactions), and does not involve derivative couplings of R_i to matter fields. However, the field redefinitions (3.12) and (3.13) with the Kähler potential (3.14) gives precisely the four-dimensional lagrangian terms identified above including the Chern-Simons squared terms, and no others. The supergravity f function is obtained by a Weyl rescaling from the four-dimensional Einstein frame in which supersymmetry is manifest, and not by a matching of the original four-dimensional Einstein frame with current-current couplings (3.11) directly to supergravity frame.

It should be noted that the Einstein frame Chern-Simons lagrangian (3.11) is of the form of a current-current interaction, and does not contain non-derivative interactions between the branes. The non-derivative couplings appear only after the field redefinitions (3.12) to obtain a manifestly supersymmetric action in terms of $N = 1$ supermultiplets. Since these field redefinitions are so important in obtaining in the correct form of the Kähler potential and supergravity f function it is instructive to consider them in more detail. First, the field redefinitions (3.12) indicate that the arguments of the logarithms in the Kähler potential (3.14) depend only on the geometric radii, and not the brane fields, $T_i + T_i^\dagger - \mathrm{tr}\,\Sigma_i^\dagger\Sigma_i - \mathrm{tr}\,Q_i^\dagger Q_i = 2R_i^2$. The corresponding supergravity f function then also only depends on the compact volume and not the brane fields, $f = -3(R_1R_2R_3)^{2/3}(S + S^\dagger)^{1/3} = -3V_6^{1/3}(S + S^\dagger)^{1/3}$. This in fact must be the case for direct compactification from ten dimensions since the coefficient of the Einstein term in the geometric frame just depends on the compact volume V_6 and the dilaton, independent of any brane matter. A Weyl rescaling from geometric to Einstein to supergravity frame then results in the f function given above which depends only on V_6 and S.

The field redefinitions (3.12) can be derived in this case from the underlying theory. The fermionic partners of the scalar components of the T_i moduli are components of the ten-dimensional gravitino. In the four-dimensional theory the supersymmetric variation of a fermion is proportional to the derivative of its scalar partner. The ten-dimensional local supersymmetric variation of the gravitino is [20]

$$\delta_\eta \psi_M = D_M\eta + \frac{\sqrt{2}}{32}\eta e^{-\phi_{10}/2}\left(\Gamma_M^{\ NPQ} - 9\delta_M^N\Gamma^{PQ}\right)H_{NPQ}\,, \tag{3.15}$$

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where a Γ matrix with n indices is the antisymmetric product of n Dirac matrices that satisfy the (field-dependent) ten-dimensional Clifford algebra. The variation (3.15) gives rise to two types of contributions in the four-dimensional theory. The first is from the covariant derivative appearing in the first term. The covariant derivative includes derivatives of R_i^2. The second type of term arises from the the three-form field strength, which contains the Chern-Simons form. Upon compactification to four dimensions, the Chern-Simons form includes derivatives of two matter fields as described above. In this way the supersymmetric variation of ψ_{T_i} is the derivative of a sum of two terms, the geometric radius squared and a composite of matter fields. The scalar partner of ψ_{T_i} is then precisely this combination of scalar fields, as given in (3.12). The four-dimensional superpotential and Kähler potential are then functions of this combination of fields, rather than the original geometric variables.

For a general $\mathbb{T}^6$ compactification which preserves $N = 4$ and includes off-diagonal fields, $T_{i\bar{j}}$ with $i \neq \bar{j}$, we may use the SU(3) symmetry to infer from (3.14) the full result

$$K = -\ln\left(S + S^\dagger\right) - \ln\det\left(T_{i\bar{j}} + T^\dagger_{i\bar{j}} - \operatorname{tr}\Sigma_i\Sigma^\dagger_{\bar{j}} - \operatorname{tr}Q_i Q^\dagger_{\bar{j}}\right), \tag{3.16}$$

which agrees with the known result for $N = 4$. As discussed in section 2, this Kähler potential is not of the sequestered form. In supergravity frame the supergravity function f of the low energy four-dimensional theory is not separable and there are unsuppressed tree-level interactions between the visible and hidden sector branes. In an S_3 preserving $N = 1$ orbifold the lowest tree-level Kähler potential for the untwisted states is the same as in (3.16) but where the non-invariant states are projected out. With hidden sector supersymmetry breaking this in general gives rise to unsuppressed non-degenerate visible sector soft scalar masses which are of order the gravitino mass.

An interesting feature of the $N = 4$ Kähler potential as well as those obtained by inheritance is that the overall scale for the associated soft masses are independent of the moduli and can not be be made parametrically small in some region of moduli space, as discussed in appendix B. This feature can be easily seen in the strongly coupled Hořava-Witten limit derivation given here with the rescalings between frames given above. The coefficient of the eleven-dimensional Einstein term $\frac{1}{2}\int d^{11}x\sqrt{g}\mathcal{R}$ reduced to geometric frame in four dimensions is proportional the total internal volume $R_{11}V_6 \sim TS^{2/3}$, where $R_{11} \sim T/S^{1/3}$ and $V_6 \sim S$. The coefficient of the ten-dimensional Hořava-Witten brane gauge fields $\frac{1}{4}\int d^{10}x\sqrt{g}F_{\mu\nu}F^{\mu\nu}$ reduced to geometric frame in four dimensions gives brane matter field kinetic terms proportional to $V_6 g^{i\bar{i}} \sim S^{2/3}$ where $S \sim V_6$ is the coefficient of the four-dimensional gauge coupling and the inverse metric $g^{i\bar{i}} \sim 1/V_6^{1/3} \sim 1/S^{1/3}$ arises from contraction of the brane gauge fields indices on the internal space. Weyl rescaling from geometric to Einstein frame then gives brane field kinetic terms proportional to $1/T$. Now the coefficient of the brane-brane quartic couplings arising form integrating out the bulk four-form field strength kinetic term $\int d^{11}x\sqrt{g}G_{ABC11}G^{ABC11}$ with the solution (3.4) reduced to geometric frame in four dimensions is proportional to $R_{11}V_6 g^{11}g^{i\bar{i}}g^{j\bar{j}} \sim V_6^{1/3}/R_{11} \sim S^{2/3}/T$, where $R_{11}V_6$ is the internal volume, $g^{11} = 1/R_{11}^2$ is the inverse metric with the $x^{11} \in [0, 1]$ coordinates of (3.4), and the inverse metric $g^{i\bar{i}}g^{j\bar{j}} \sim 1/V_6^{2/3}$ arises from contraction of the brane gauge field indices on the internal space. A Weyl rescaling to Einstein frame then

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gives four-dimensional matter quartic couplings proportional to $1/T^2$. Finally, in order to obtain the physical quartic couplings and physical soft masses the matter fields must rescaled to canonical normalization by multiplying each matter field by $T^{1/2}$. This gives quartic couplings and soft masses which are independent of both T and S, and therefore also of R_{11} and V_6 holding a four-dimensional scale such as the gravitino mass fixed. All these rescalings are of course implicitly contained in the above Kähler potentials and eleven-dimensional derivation given above.

3.1.1 Inheritance for M-theory Orbifolds

The microscopic derivation of the Kähler potential for $N = 4$ compactifications from Hořava-Witten theory can be modified to provide a justification for the inheritance principle determining the lowest order tree-level Kähler potential for the untwisted states in an $N = 1$ orbifold compactification $\mathbb{T}^6/\Gamma \times S_1/\mathbb{Z}_2$ in the brane world limit. General reasons were presented in section 2.1 arguing that the lowest order Kähler potential for these states is determined from the $N = 4$ expression by deleting the non-invariant states. This inheritance can justified in the large volume brane world limit, $T^3 \gg S \gg 1$ by inspecting the explicit supergravity solution for the four-field strength (3.4), which is valid at generic points in the compact space in this limit up to small power suppressed quantum M-theory corrections.

If the fluctuating fields appearing on the right-hand side of (3.4) are restricted o those massless states that survive the orbifold projection, then the Chern-Simons form

$$\omega_{\mu i \bar{j}} = \mathrm{tr}\, Q_j^\dagger \partial_\mu Q_i - \mathrm{h.c.} + \cdots \tag{3.17}$$

contains some states Q_i^A and Q_j^A transforming under a representation $\{A\}$ of the unbroken gauge group. The ellipses denote terms involving other representations and also terms involving three fields. The $i = 1, 2, 3$ denotes the SU(3) label, where under an abelian orbifold action the complex coordinate of the torus $z_i \to \gamma_i z_i$ with $\gamma_i^{|\Gamma|} = 1$. The orbifold is also embedded into the gauge group, so that each field transforms under the orbifold group as

$$Q_i^A \longrightarrow e^{2\pi i r_i} e^{-2\pi i \beta_A} Q_i^A = Q_i^A \,, \tag{3.18}$$

where the last equality follows from restricting to invariant states. More details on the notation can be found in appendix A.

The argument for inheritance appearing in the subsequent paragraph utilizes the following fact: if after the orbifold projection the Chern-Simons terms in (3.17) appear, then necessarily $\gamma_i \equiv \gamma_j$. This is almost obvious, since the states are invariant and it naively appears that the gauge transformation appearing in (3.18) cancels between Q_i and Q_j. But there is a loophole that could allow two states Q and Q' with the same quantum numbers to have different transformations under the embedding of the orbifold into the gauge group, namely $e^{2\pi i \beta} \neq e^{2\pi i \beta'}$. This could in principle occur if in the decomposition of a represention $\mathcal{R}$ of $\mathcal{G}$ into representations $\{\mathcal{R}'\}$ of a regular subalgebra $\mathcal{G}'$ the same representation $\mathcal{R}'$ appeared more than once. Inspecting the branching tables appearing in Slanksy [21] indicates that this does occur but rarely and only for very high-dimensional

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representations $\mathcal{R}$. But for an adjoint representation this does not occur. This is because the U(1) charges of a charged root of $\mathcal{G}$ are enough to completely specify the root. Thus two states Q and Q' with the same quantum numbers must have the same orbifold gauge transformation, $\beta = \beta'$, since in the higher-dimensional theory the quantum numbers of these states originate from the adjoint representation of E_8. This combined together with the requirement that $r_i \equiv \beta_A \mod |\Gamma|$ for invariant states is enough to imply that if gauge invariant terms appear in (3.17), then necessarily $\gamma_i = \gamma_j$. These arguments can also be directly applied to $\omega_{\mu ij}$ with the conclusion that if these terms appear then $\gamma_i = \gamma_j^*$.

With the above information it is now possible to to argue that (3.4) restricted to invariant states remains a solution in the orbifold background at tree-level. The point is that the solution for the four-form field strength on the covering space must respect the transformatation properties specified by the orbifold group. Thus (3.4) is still the correct solution at lowest order provided it satisfies the boundary conditions

$$G_{\mu i\bar{j}11}(x_{\text{torus}}, x_{11}) = q_{\mathbb{Z}_2}\gamma_i\gamma_j^* G_{\mu i\bar{j}11}(\mathcal{U}x_{\text{torus}}, -x_{11}), \tag{3.19}$$

where $\mathcal{U}$ is the orbifold rotation acting on the internal coordinates of the torus and $q_{\mathbb{Z}_2}$ is the M-theory orbifold charge. There is a similar condition for $\bar{j} \to j$. To describe the massless untwisted states we restrict the fluctuating fields appearing on the right-hand side of (3.4) to those massless states that survive the orbifold projection. Then the Chern-Simons term appearing in (3.4) is necessarily constant across the $\mathbb{T}^6$ orbifold. Thus in order that (3.4) restricted to the invariant states remains a solution, the boundary conditions for $G_{i\bar{j}\mu 11}$ and $G_{ij\mu 11}$ given above must allow for constant solutions across the $\mathbb{T}^6$ orbifold. This will be the case provided that $\gamma_i\gamma_j^* \equiv 1$ and $\gamma_i\gamma_j \equiv 1$, respectively. The argument of the previous paragraph demonstrated that this is always the case. The remaining arguments leading from (3.4) to the Kähler potential (3.14) then follow through except that now in (3.14) or in (3.16) the non-invariant states are projected out. Thus the solution (3.4) for the toroidal compactification can be trivially extended to obtain the lowest order tree-level Kähler potential for the untwisted states of the the orbifold by simply deleting the non-invariant states. This provides a supergravity justification for the use of the inheritance principle in the large volume brane world limit. The expression (3.14) is then appropriate for an abelian orbifold that does not preserve any permutation symmetry, whereas (3.16) applies for an abelian orbifold that preserves an S_3 permutation symmetry.

3.2 The strongly coupled heterotic theory in the five-dimensional limit

The Hořava-Witten theory compactified on a generic six-manifold preserving at least $N = 1$ supersymmetry in four dimensions leads, as discussed in section 2.1, to unsuppressed non-universal tree-level visible sector scalar masses comparable to the gravitino mass, even with supersymmetry breaking isolated on a hidden sector brane. The microscopic origin of these masses in the ten-dimensional limit in which the M-theory direction is smaller than any of the compact six-manifold directions, $R_i \gg R_{11}$, was explained in the previous subsection. The required contact terms arise from exchange in the M-theory direction of the bulk four-form field strength between the branes. In this subsection we consider the

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five-dimensional limit in which the length scales of the six-dimensional compact manifold are much smaller than the M-theory direction, $R_{11} \gg R_i$. This yields a five-dimensional theory below the compactification scales of the six-dimensional compact manifold. An $S_1/\mathbb{Z}_2$ orbifold projection of the M-theory direction in this five-dimensional theory then yields the five-dimensional analog of the Hořava-Witten theory with end of the world branes just as in eleven dimensions. If the compact manifold is $\mathbb{T}^6$ then $N = 4$ supersymmetry is preserved in the low energy four-dimensional theory obtained by the $S^1/\mathbb{Z}_2$ orbifold of the five-dimensional theory. The Kähler potential of a four-dimensional $N = 4$ theory is not renormalized. Because of this, the Kähler potential in the ten-dimensional limit is identical to the Kähler potential in the five-dimensional limit. Non-universal tree-level scalar masses then also arise in this limit. The non-universal nature of the tree-level soft masses survives generic projections of the $\mathbb{T}^6$ compactification which preserve only $N = 1$ supersymmetry in four dimensions.

The origin of the non-universal visible sector scalar masses is simple to understand in the five-dimensional limit. In particular, the five-dimensional theory obtained by compactifying from eleven dimensions contains a number of U(1) gauge bosons which couple non-universally to the matter living on the boundaries. Visible and hidden sector matter on the branes are neutral under the bulk U(1)'s, but appear as sources in a modification of the Bianchi identities for the bulk gauge bosons. Integrating out these gauge bosons to obtain the four-dimensional effective theory precisely generates the Chern-Simons squared contact interactions (3.11) and Kähler potentials described in the previous sections.

To see the origin of the non-universal brane-brane interactions in this limit, note that the five-dimensional decomposition of the eleven-dimensional three-form potential with two compact indices and one non-compact index gives five-dimensional vector bosons. The number of these gauge bosons depends on properties of the compact six-manifold. For toroidal compactification on $\mathbb{T}^6$ the low energy five-dimensional theory has an $\mathrm{SO}(6) \cong \mathrm{SU}(4)$ symmetry at leading order in the low energy derivative expansion. The five-dimensional vector bosons arising from the three-form potential in this case transform as $\mathbf{15} \in \mathrm{SU}(4)_R$. Under the $N = 1$ decomposition of the $N = 4$ R-symmetry the vector bosons transform as $\mathbf{8}_0 \oplus \mathbf{1}_0 \oplus \mathbf{3}_+ \oplus \overline{\mathbf{3}}_- \in \mathrm{SU}(3) \times \mathrm{U}(1)_R \subset \mathrm{SU}(4)_R$. For definiteness we focus on the $\mathbf{8} \oplus \mathbf{1} \in \mathrm{SU}(3)$ gauge bosons with four- and five-dimensional components

$$A_\mu^{(i\bar{\jmath})} \equiv C_{i\bar{\jmath}\mu}\,, \qquad A_{11}^{(i\bar{\jmath})} \equiv C_{i\bar{\jmath}11}\,, \tag{3.20}$$

where $i, \bar{\jmath}$ are complex coordinate indices on the compact six-manifold, and where the M-theory direction (indicated by a subscript 11) is now the fifth dimension. These U(1) gauge bosons are the ones which would survive the $\mathbb{Z}_3$ orbifold projection which preserves $N = 1$ described in section 2.1. The SU(3) global symmetry is a flavor symmetry under which matter on both the visible and hidden sector end of the world branes transform. The bulk five-dimensional U(1) gauge bosons then also transform under this global flavor symmetry. Most importantly, the field strengths for these gauge bosons have a modified Bianchi identity that is inherited from the eleven-dimensional Bianchi identity (3.2) for the

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four-form field strength [9]

$$\left(dF^{(i\bar{j})}\right)_{\mu\nu 11} = \delta(x_{11} - x_{11,\mathrm{hid}}) J^{(\mathrm{hid})}_{\mu\nu i\bar{j}}(x^\mu) + \delta(x_{11} - x_{11,\mathrm{vis}}) J^{(\mathrm{vis})}_{\mu\nu i\bar{j}}(x^\mu)\,. \tag{3.21}$$

The sources appearing on the right-hand side of the Bianchi identity depend on the $i\bar{j}$ flavor for both the visible and hidden sector fields and also transform under the global flavor symmetry as $\mathbf{8} \oplus \mathbf{1} \in \mathrm{SU}(3)$. The solution to (3.21) is the same as eq. (3.4) for the four-form field strength in the previous section, but where here the components $G_{\mu i\bar{j}11}$ are relevant. The dimensional reduction of the action from eleven to five dimensions contains a term

$$\int dx_{11} d^4x V_6 e^{-2\phi_{10}} \frac{\left(F^{(i\bar{j})}_{\mu 11}\right)^2}{R_i^2 R_j^2}\,. \tag{3.22}$$

Integrating out $F^{i\bar{j}}_{\mu 11}$ in the five-dimensional theory then yields precisely the brane-brane contact interactions discussed above. In this limit the SU(3) flavor symmetries on each brane are coupled by exchange of the bulk gauge bosons. This arises in the five-dimensional theory by integrating out the Kaluza-Klein tower of five-dimensional bulk gauge boson states even though the four-dimensional zero modes of these fields are projected out by the $S^1/\mathbb{Z}_2$ orbifold in the M-theory direction. So in this limit the brane-brane contact interactions are generated by exchange of massive bulk fields without any (exponential) suppression of the coupling.

Note that since the bulk gauge bosons exist in the five-dimensional theory, their mass is protected by a gauge invariance. It is then not possible to suppress the couplings they generate by giving them a large mass in the $S^1/\mathbb{Z}_2$ projection to the four-dimensional theory. Eliminating the brane-brane contact interactions in this limit would require projecting out of the five-dimensional theory all the bulk U(1) gauge bosons which have brane sources through Bianchi identities of the form (3.21). In the compactification from eleven to five dimensions this would require lifting all of the three-form potentials with indices in the compact directions. This is in fact not possible with a symmetric abelian orbifold projection which does not reduce the rank and necessarily leaves invariant at least three five-dimensional vector bosons with diagonal internal indices $i = \bar{j}$. These three bulk gauge bosons have Bianchi identity brane sources for the $\mathrm{U}(1)^3 \subset \mathrm{SU}(3) \times \mathrm{U}(1)_R$ preserved by any abelian orbifold. Removing all vector bosons from the five-dimensional theory (except of course the one in the gravitational supermultiplet) might in principle be possible with asymmetric and/or non-abelian projections which reduce the rank of the global symmetry.

Since the flavor symmetries on each brane are coupled through exchange of bulk gauge bosons, non-universal visible sector tree-level scalar masses arise for generic supersymmetry breaking on the hidden sector brane. Universal vanishing tree-level visible sector masses require flavor symmetric hidden sector supersymmetry breaking. However, in this case, universality is the result of unbroken flavor symmetries and not simply from physically separating the visible and hidden sector branes.

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3.3 Open strings

The Kähler potential for D-brane world models of BWSB are generally not of the no-scale sequestered form, as described in section 2.2. This implies that brane-brane contact interactions exist in the low energy four-dimensional theory. Microscopically, bulk locality implies that these interactions must arise from tree-level exchange of bulk fields between the branes. In this subsection we illustrate the origin of these interactions in simple D-brane world models which preserve $N = 4$ supersymmetry. While not phenomenologically realistic, these models illustrate that the sequestered intuition breaks down even in situations with a high degree of supersymmetry. There is no reason to think that things will be different in models with less supersymmetry to protect the form of the Kähler potential, and we will illustrate such effects in examples with $N = 2$ and $N = 1$ supersymmetry later. D-brane models are also instructive in illustrating the dependence on the internal volume in examples with brane co-dimension larger than one. The volume dependence of brane-brane interactions leads to unsuppressed visible sector soft scalar masses from hidden sector supersymmetry breaking as compared with the gravitino mass, just as in the co-dimensions one cases.

The simplest example of a D-brane world model is provided by type-I string theory with gauge group SO(32) in ten dimensions compactified on a circle of radius R. Interaction terms between visible and hidden sector fields are easily exhibited, as above, in Chern-Simons squared couplings. The low energy nine-dimensional theory at energy scales below R^{-1} contains, among other terms, the dimensional reduction of the Chern-Simons term proportional to $R\omega^2_{\rm MNO}$, where upper case Latin indices denote the non-compact dimensions. In particular, the nine-dimensional action includes a term proportional to

$$\int d^9x R \left(dC_{(2)} - \omega_{\rm CS}\right)^2 , \tag{3.23}$$

where $C_{(2)}$ is the type-I RR two-form potential. With Wilson lines turned on the Chern-Simons form factorizes into a sum over the Chern-Simons forms of the unbroken subgroups. Thus (3.23) includes quadratic terms involving the product of different Chern-Simons forms.

Now consider the type-I$'$ description obtained by T-duality on the compact direction. In this theory there are 2 O8 planes separated by a distance $R' = 1/R$ and (including images) 32 D8-branes transverse to the interval between the O8 planes in a 1/2 BPS configuration which preserves 16 supersymmetries. A brane world model may be obtained by arranging the D8-branes in two physically separated groups along the interval. For example, if (including images) 16 D8-branes are placed on each O8 plane an SO(16)×SO(16) gauge theory is obtained. The bulk between the branes is co-dimension one in this example. Generalization to other co-dimensions will be analyzed below. The T-dual of the low energy nine-dimensional action (3.23) may be written

$$\int d^9x \left(\frac{1}{R'}\left(dC_{(2)}\right)^2 - \frac{2}{R'} dC_{(2)} \cdot \omega_{\rm CS} + \frac{\omega^2_{\rm CS}}{R'}\right) , \tag{3.24}$$

where dependence on the type-I RR two-form field has been retained. An important observation is that for separated branes the Chern-Simons form (as opposed to its square)

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appearing in (3.24) is a sum of the individual Chern-Simons forms from each of the groups of D-branes. This follows by T-duality from the type-I description with Wilson lines turned on.

We have focused on terms where all the indices are in the non-compact direction and they are raised and lowered using the non-compact metric. Since the low-energy nine-dimensional theory is invariant under T-duality, these interactions must also be present in the type-I$'$ description. Cross terms between fields localized on the visible and hidden sector D-branes in the type-I$'$ description must then arise microscopically from exchange of bulk fields between the visible and hidden sector D-branes. Of importance are the volume suppressed Chern-Simons squared interactions.

The first two terms in (3.24) may be easily understood in the type-I$'$ D-brane picture. This theory has a RR three-form potential $C_{(3)}$ in the bulk which couples to gauge fields propagating on the branes in a manner described below. The first term in (3.24) is just the dimensional reduction of the bulk kinetic term for this three-form

$$\int d^{10}x \left|dC_{(3)}\right|^2 \longrightarrow \int d^9x R' g^{99} \left|dC_{(3)\mu\nu 9}\right|^2 . \tag{3.25}$$

Under T-duality the type-I RR two-form potential becomes $C_{(2)\mu\nu} \to C_{(3)\mu\nu 9}$ in the type-I$'$ theory. Then with $g^{99} = 1/R'^2$ the kinetic term (3.25) agrees with the first term of (3.24) in terms of type-I fields. The second term in (3.24) is obtained in the type-I$'$ description from Wess-Zumino couplings between the brane fields and the bulk RR potentials [14],

$$S_{\mathrm{WZ}} = \int_9 \mathrm{Tr}\left(e^{2\pi\alpha' F}\right) \wedge \sum_q C_{(q)} , \tag{3.26}$$

where the sum is over all the RR potentials. The existence of these interactions may be understood from the observation that instantons on the brane world volume act as sources for RR forms of lower rank [22]. Note that the interactions (3.26) do not depend on either the dilaton or volume of the compact space. The expansion of the exponential in the type-I$'$ theory gives a coupling of brane instanton number to the bulk five-form potential

$$S_{\mathrm{WZ}} \supset \int_9 F \wedge F \wedge C_{(5)} . \tag{3.27}$$

With $d\omega_{\mathrm{CS}} = F \wedge F$, this interaction may be written as

$$\int_9 \omega_{\mathrm{CS}} \wedge dC_{(5)} . \tag{3.28}$$

Using Poincaré duality $dC_{(5)}$ can be written as a four-form field strength, $F_{(4)} =^* dC_{(5)} = dC_{(3)}$. However, Poincaré duality utilizes the Levi-Civita ϵ-tensor, which, since one of its indices is in the compact direction, introduces a factor of $1/R'$. To see this, note that the above coupling is proportional to

$$\omega_{(\mathrm{CS})[\mu_1\mu_2\mu_3} F_{(6)\mu_4\cdots\mu_9]} \tag{3.29}$$

and the Poincaré duality gives $F_{(6)\mu_4\cdots\mu_9} = \epsilon_{\mu_4\cdots\mu_9}{}^{a_1\cdots a_4} F_{(4)a_1\cdots a_4}$. Now all of $\mu_1, \ldots, \mu_9$ are along the D8-brane directions, so one of the components of $F_{(4)}$ and an upper component of the ϵ-tensor is in the compact direction. If the metric is used to express the ϵ-tensor with

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mixed components in terms of an ϵ-tensor with all lower components, then three of the indices of $F_{(4)}$ are contracted with the indices of the Chern-Simons form, with one index left out. Since the non-zero value of the ϵ-tensor with all lower indices is $\pm\sqrt{-g}$, where g is the determinant of the bulk metric, the dependence of this coupling on the metric is $g^{99}\sqrt{-g} = 1/R'$. After dimensional reduction the remaining component of $F_{(4)}$ in the compact direction cannot involve a derivative, so it must also be a component of $C_{(3)}$.

Putting this together, the Wess-Zumino interactions, together with Poincaré duality, imply a coupling of the brane gauge fields to the bulk three-form of the form

$$\int d^9x \frac{1}{R'} \omega_{\rm CS} \cdot dC_{(3)} \,, \tag{3.30}$$

with $C_{(3)} = C_{(3)\mu\nu 9}$, and the three indices of $\omega_{\rm CS}$ are contracted with the three non-compact indices of dC_3 using the lower-dimensional metric. The generalization to any number of tori or circles transverse to the brane implies that this interaction is in general suppressed by the volume of the compact space transverse to the brane. Notice that the above dependence on the volume and the dilaton agrees with the second term (3.24) since $C_{(2)\mu\nu} \to C_{(3)\mu\nu 9}$ under T-duality.

Now the full action of the compactified nine-dimensional theory should be invariant under T-duality. But since the low-energy lagrangian in the type-I theory contains the $R\omega_{\rm CS}^2$ interaction in (3.23), we infer from the invariance under T-duality that the low energy type-I$'$ D-brane world theory contains not only the first two terms of (3.24) discussed above from the point of view of brane field interactions with bulk RR fields, but also the third term of (3.24) which is the volume suppressed contact interaction

$$\int d^9x \frac{1}{R'} \omega_{\rm CS}^2 \,, \tag{3.31}$$

which couples branes that are physically separated in the microscopic type-I$'$ theory. More generally, in any number of dimensions, this coupling is suppressed by a factor of the compact volume transverse to the brane as discussed below.

In the closed string channel general brane-brane interactions are obtained from integrating out the Kaluza-Klein tower of RR fields. In the D8-brane example above, brane gauge instanton number, $F \wedge F$, appears as a source in the bulk five-form potential equation of motion. By using Poincaré duality the sources appear in a modified Bianchi identity for the four-form field strength. The contact brane-brane interaction (3.31) is then generated from integrating out the Kaluza-Klein tower of the three-form potential, much as for the one-form in the five-dimensional limit of the Hořava-Witten model discussed in the previous subsection. This can be generalized to branes of lower dimension, but the modified Bianchi identity that follows is now more intricate and a solution is not presented here.

In the open string channel the brane-brane interactions are generated by the one-loop amplitude of open string states stretching between the branes. By modifying the RR amplitude for the force between two parallel groups of D-branes [14] to include the gauge boson vertex operators, it is possible to obtain the brane-brane interaction (3.31) directly with correct dependence on the volume. Further work on this open string perspective is in progress [23].

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The general volume dependence of brane-brane interactions for brane co-dimensions larger than one and resulting Kähler potentials in four dimensions may be illustrated by additional T-duality transformations of the type-I theory. Consider first type-I theory on a symmetric $\mathbb{T}^2$ of radius R_1. The T-dual description in the $\mathbb{T}^2$ directions of this background is type-IIB string theory with (including images) 32 D7-branes and 4 O7 orientifold planes on $\mathbb{T}^2/\mathbb{Z}_2$ with radius $R' = 1/R$. This configuration is 1/2 BPS and preserves 16 supersymmetries. Following arguments similar to those preceding (3.24) and (3.31) the low energy eight-dimensional action contains the Chern-Simons squared volume suppressed contact interactions between D7-brane fields

$$\int d^8x \frac{\omega_{\rm CS}^2}{R'^2} \,. \tag{3.32}$$

By T-duality, similar transverse volume suppressed brane-brane contact interactions arise for type-IIB configurations with 32 D5-branes and 8 O5 planes or with 32 D3-branes and 16 O3 planes. The Kähler potential derived below applies to these cases also.

Now consider the further toroidal compactification of the eight-dimensional type-IIB configuration with 32 D7-branes and four O7 planes on a product of two symmetric tori $\mathbb{T}^2 \times \mathbb{T}^2$ with radii R_2 and R_3, and we relabel $R' \to R_1$. The low energy four-dimensional theory has $N = 4$ supersymmetry. Before Weyl rescaling the lagrangian in geometric frame obtained by direct compactification contains the couplings

$$\begin{aligned}\int d^4x \Bigg\{ & V_6 e^{-2\phi_{10}} \mathcal{R}_{(4)} + e^{-\phi_{10}} \left(R_2^2 R_3^2\right) \frac{1}{4} \,{\rm tr}\, F_{\mu\nu}^2 + \\ & + R_2^2 R_3^2 e^{-\phi_{10}} \left({\rm tr}\, \frac{(\partial_\mu \varphi_2)^2}{2R_2^2} + \frac{e^{\phi_{10}}}{8R_1^2 R_2^4} ({\rm tr}\, \varphi_2^* \partial_\mu \varphi_2 - {\rm h.c})^2 + \right. \\ & \left. + \frac{{\rm tr} (\partial_\mu \varphi_3)^2}{2R_3^2} + \frac{e^{\phi_{10}}}{8R_1^2 R_3^4} ({\rm tr}\, \varphi_3^* \partial_\mu \varphi_3 - {\rm h.c})^2 \right) \Bigg\} ,\end{aligned} \tag{3.33}$$

where $V_6 = R_1^2 R_2^2 R_3^2$ is the compact volume, ϕ_{10} is the type-IIB ten-dimensional dilaton, and φ_2 and φ_3 are complex matter fields that are the zero modes of the brane gauge field with components along $\mathbb{T}^2 \times \mathbb{T}^2$ directions. The lagrangian (3.33) includes the Einstein-Hilbert action, the gauge kinetic function, and the matter action obtained from the dimensional reduction of the higher-dimensional gauge kinetic terms localized on the branes, and also the volume suppressed dimensionally reduced Chern-Simons squared interactions. The dependence on the radii arises from the internal metric, except for the quartic terms which contain an additional volume suppression R_1^2. After a Weyl rescaling to Einstein frame the above action for the matter fields becomes

$$\int d^4x \frac{e^{\phi_{10}}}{R_1^2} \left(\frac{(\partial_\mu \varphi_2)^2}{2R_2^2} + \frac{e^{\phi_{10}}}{8R_1^2 R_2^4} ({\rm tr}\, \varphi_2^* \partial_\mu \varphi_2 - {\rm h.c})^2 + (2 \to 3) \right) . \tag{3.34}$$

It is important to note that now the dilaton and radii dependence of the coefficient of the quartic term is the square of the coefficient of the quadratic term. Finally, there is also an adjoint scalar from the collective coordinates of the D7-branes. The kinetic terms for

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these collective coordinates come from the lowest terms in the Nambu-Goto action for the D7-branes

$$T_7 \int d^8x e^{-\phi_{10}} \sqrt{h_{||}} h_{\perp ij} \frac{1}{2} \operatorname{tr} \partial_\mu x^i \partial_\mu x^j \longrightarrow \int d^4x R_1^2 R_2^2 R_2^3 e^{-\phi_{10}} \operatorname{tr} |\partial_\mu \varphi_1|^2 , \tag{3.35}$$

where T_7 is the string frame D7-brane tension, $h_{||}$ is the determinant of the induced longitudinal metric, and $h_{\perp ij}$ is the induced transverse metric with $i, j = 8, 9$. In the second expression in (3.35) the D7-branes are compactified on $\mathbb{T}^2 \times \mathbb{T}^2$ in geometric frame and the collective coordinate scalars properly normalized in geometric string frame are

$$\varphi_1 = \sqrt{T_7} \frac{\sqrt{h_\perp}\,(x^8 + ix^9)}{\sqrt{2} R_1} . \tag{3.36}$$

After the same Weyl rescaling to Einstein frame as above these kinetic terms become

$$\int d^4x e^{\phi_{10}} \operatorname{tr} |\partial_\mu \varphi_1|^2 . \tag{3.37}$$

In order to write the action in a manifestly supersymmetric form it is necessary to redefine the scalar components of the four-dimensional chiral supermultiplets in terms of the geometric radii, ten-dimensional dilaton, and brane matter fields. With the field redefinitions for the scalar components of the four-dimensional dilaton

$$S = e^{-\phi_{10}} R_2^2 R_3^2 , \tag{3.38}$$

and the other moduli

$$\begin{aligned} T_1 &= e^{-\phi_{10}} + \frac{1}{2} \operatorname{tr} \varphi_1^* \varphi_1 , \qquad T_2 = e^{-\phi_{10}} R_1^2 R_2^2 + \frac{1}{2} \operatorname{tr} \varphi_2^* \varphi_2 , \\ T_3 &= e^{-\phi_{10}} R_1^2 R_3^2 + \frac{1}{2} \operatorname{tr} \varphi_3^* \varphi_3 , \end{aligned} \tag{3.39}$$

the lagrangian terms described above are obtained from the Kähler potential

$$K = -\ln\left(S + S^\dagger\right) - \sum_i \ln\left(T_i + T_i^\dagger - \operatorname{tr} \varphi_i^\dagger \varphi_i\right) . \tag{3.40}$$

For a general $\mathbb{T}^6$ compactification, inclusion of the off-diagonal moduli and interactions yields the $N = 4$ Kähler potential

$$K = -\ln\left(S + S^\dagger\right) - \ln\det\left(T_{i\bar{j}} + T_{i\bar{j}}^\dagger - \operatorname{tr} \varphi_i \varphi_{\bar{j}}^\dagger\right) . \tag{3.41}$$

In order to obtain a brane world model, the D-branes may be separated into two groups which model the visible and hidden sectors, $\varphi_i = Q_i$ or Σ_i. Due to $N = 4$ supersymmetry there are no corrections to the Kähler potential as one turns on the branes (i.e. turns on expectation values for certain scalar fields on the brane. One can check that integrating out massive fields does not generate such couplings at tree level). The four-dimensional scalar fields describing the D-brane separation may then be replaced by their expectation values, and the brane fields in (3.41) break up into a sum of visible and hidden sector fields,

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$\mathrm{tr}\,\varphi_i\varphi^\dagger_{\bar{j}} = \mathrm{tr}\,Q_iQ^\dagger_{\bar{j}} + \mathrm{tr}\,\Sigma_i\Sigma^\dagger_{\bar{j}}$. This is then of course the same Kähler potential given above for an $N = 4$ theory with visible and hidden sectors.

The Kähler potential (3.41) is identical to the one found in the toroidal compactification of heterotic theory in ten dimensions since both of these low-energy theories preserve $N = 4$. Since the ten-dimensional lagrangian for the dilaton and metric are identical in both the type-I and the heterotic theories, it may then appear strange that the type-I′ definitions (3.38) and (3.39) of the moduli in terms of the string coupling and torus radii are different than the heterotic definitions and (3.12) and (3.13). This is, however, not a puzzle since these two sets of definitions are in fact mapped into each other by performing T-duality in the $\mathbb{T}^2$ direction with radius R_1 and a subsequent type-I-heterotic duality, both of which are symmetries of the low energy theory.

For a compactification which preserves only $N = 1$ in four dimensions and for which all the off-diagonal moduli and interactions are lifted, such as one of the $\mathbb{Z}_6$ orbifold compactifications discussed in appendix A the lowest order tree-level inherited Kähler potential would be of the sum of logarithms form (3.40). For a compactification preserving a S_2 or S_3 permutation symmetry, such as one of the $\mathbb{Z}_6$ and $\mathbb{Z}_3$ orbifolds discussed in appendix A having 2 or 3 generations in the untwisted sector, their lowest-order inherited Kähler potential would be related to (3.41) Again, a parallel brane world model can be constructed by simply separating the D-branes into two groups. In this case, $N = 1$ supersymmetry does not forbid corrections to the Kähler potential (3.40). The origin of some leading corrections are described in the next section. However, the lowest order tree-level Kähler potential is indeed inherited in this case from (3.40) or (3.41).

In either case, with generic hidden sector supersymmetry breaking with stabilized moduli, the Kähler potentials (3.40) and (3.41) lead to unsuppressed non-degenerate tree-level visible sector scalars and potentially undesirable tachyons. As suggested above, the brane-brane contact interactions contained in these Kähler potentials arise microscopically from exchange of bulk RR fields between visible and hidden sector brane fields. The interactions are suppressed by the compact volume transverse to the branes. In terms of the underlying theory with the microscopic Planck scale held fixed this leads to four-dimensional visible sector scalar masses suppressed by the total internal volume $m_Q^2 \sim F^2/V$, where F is a hidden sector auxiliary expectation value. However, the four-dimensional gravitino mass is determined by the four-dimensional Newton constant, which is also volume suppressed, $m_{3/2}^2 \sim F^2/V$. So the visible sector scalar masses are unsuppressed with respect to the gravitino mass, $m_Q^2 \sim m_{3/2}^2$. The volume dependence for both the brane-brane interactions and four-dimensional Newton constant is of course implicitly contained within the Kähler potentials derived above.

Finally, it is worth noting that all the brane world models discussed here with $N = 4$ supersymmetry in four dimensions are actually related by dualities. The Hořava-Witten M-theory background is the strongly coupled limit of type-IIA on $S_1/\mathbb{Z}_2$ which is T-dual to to the type-I description, and this is in turn related back to the Hořava-Witten background by type-I-heterotic duality in the strongly coupled limit. Since these symmetries survive compactification to four dimensions these dualities are also sufficient to show that the Kähler potentials in these examples are identical. These are also sufficient to show that

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the brane-brane interactions of these examples for any co-dimension are suppressed by one power of the internal volume, just as the four-dimensional gravitino mass. The analysis here exhibits the different microscopic mechanisms which explain the seeming non locality in each of these pictures.

4. Corrections to the inherited Kähler potential

In many theories the visible and hidden sector matter fields are remnants of extended supermultiplets of the high energy theory. In these cases the lowest order four-dimensional Kähler potential which couples these sectors is inherited from the form dictated by the extended supersymmetry of the microscopic theory. As described from a number of points of view in the previous sections, this occurs for many brane world backgrounds of string and M-theories. This generally leads to model dependent relations among the visible sector tree-level soft masses, as discussed in section 2.1 and appendix B. However, with only $N = 1$ supersymmetry in four dimensions, the Kähler potential is not protected from generic corrections. The inherited Kähler potential, and therefore the tree-level relations among the visible sector scalar masses, is then generally modified in the full low-energy theory. It is interesting to address what form and magnitude these corrections may take in BWSB scenarios.

With standard hidden sector supersymmetry breaking the Kähler potential is believed to obtain generic corrections. However, BWSB differs in that the visible and hidden sector fields are physically separated in the microscopic theory. It has been argued that this feature implies that the separable form of the supergravity f function is stable, that the leading corrections are suppressed by additional powers of the compact volume, and are therefore controllably small [2]. The argument starts with the observation that corrections which are quantum one-loop from the bulk point of view are suppressed by two powers of the internal volume. Since the four-dimensional gravitino mass is suppressed by only one power of the internal volume, these corrections are small in the large volume limit. In addition, quantum corrections arising from point-like bulk fields are ultraviolet finite since the four-dimensional brane-brane correlators are point split (and therefore regulated) in the higher-dimensional theory. These features were used to argue that the sequestered hypothesis was natural and stable in brane world realizations [2]. As we have seen, the sequestered hypothesis breaks down even at tree level in an expansion in powers of the brane separation, and that there are generically brane-brane contact interactions which give rise to tree-level soft masses. However, one might wonder if the above arguments could be modified to demonstrate that the inherited form of the Kähler potential in brane world models is stable with controllably small corrections. In this section we show that in fact this is not the case. There is generally no sense in which the lowest order form of the Kähler potential is protected from generic corrections in BWSB models, although the corrections may be small in some corners of moduli space.

The origin of corrections to the inherited Kähler potential in brane world theories is easy to understand. From the bulk point of view, brane-brane contact interactions responsible for visible sector soft scalar masses coming from hidden sector supersymmetry

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breaking are generated by tree-level exchange of bulk supergravity fields. With $N = 4$ supersymmetry in four dimensions these interactions are fixed to have a particular form, as discussed in the previous sections. With compactifications which preserve only $N = 1$ supersymmetry these interactions take a more general form. This in turn gives rise to more general tree-level Kähler potentials with additional contributions to scalar masses beyond those of the lowest order inherited Kähler potential. In addition, warping of the internal space by non-vanishing brane tensions which generally arise even with $N = 2$ supersymmetry. This also modifies the form of the Kähler potential. If moduli acquire auxiliary components, which commonly occurs in mechanisms which stabilize the moduli, additional contributions to tree-level masses arise.

The leading corrections to the inherited Kähler potential may also be understood in a slightly different manner in weakly coupled D-brane models. In this case couplings between the visible and hidden sectors are generated quantum mechanically by integrating out heavy states which are charged under both the hidden and visible gauge groups. In a brane world model there are no localized point-like states which couple these sectors directly since they are physically separated. However, there are massive string states which stretch between the hidden and visible sector branes. The amplitude for integrating out these states at one-loop is just the open-string channel of the annulus diagram with boundaries on the visible and hidden sector branes. In the closed string channel this amplitude is simply tree-level exchange of bulk closed-string modes between the branes. For large volume this is dominated by exchange of massless bulk supergravity fields, and is therefore only suppressed by one power (rather than two as for point-like quantum amplitudes) of the internal volume. This gives rise to unsuppressed corrections to visible sector soft masses from hidden sector supersymmetry breaking. We see that the low-energy local effective field theory reasoning described above breaks down in this picture because of the existence of physically extended states in the full underlying theory.

In the next subsection corrections to the inherited Kähler metric are illustrated in simple 1/4 BPS D-brane models which preserve only 8 supersymmetries. In the following subsections warping of the internal space by non-vanishing brane tensions which result in generic $N = 1$ compactifications of the Hořava-Witten and pure five-dimensional supergravity examples are shown to generally give corrections to the lowest order inherited Kähler potential. These in turn generally give additional unsuppressed tree-level contributions to scalar masses which are not small except in corners of moduli space.

4.1 Kähler potential corrections in the D-brane picture

It is instructive to consider how corrections to the lowest order inherited Kähler potential arise in D-brane world models. This is easily illustrated in simple D-brane configurations. In BPS configurations which preserve 16 supersymmetries, specifically parallel Dp-branes, the Kähler metric is flat and the Kähler potential is exact. But in configurations with less supersymmetry the Kähler metric is modified by brane-brane interactions.

To illustrate the modification of the metric in more general D-brane configurations consider type-II uncompactified string theories in ten dimensions with a source Dp'-brane and a probe Dp-brane. We follow and slightly extend an analysis due to Brodie [24]. The

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source Dp'-brane may be thought of as the hidden sector and the Dp-brane probe as the visible sector. The metric line element and dilaton backgrounds of the source Dp'-brane at distances large compared to the string scale are

$$ds^2 = f(r)^{-1/2} dx_{\|}^2 + f(r)^{1/2} dx_{\perp}^2\,, \qquad e^{-2\phi} = f(r)^{\frac{p'-3}{2}}\,, \tag{4.1}$$

with

$$f(r) = 1 + g_s \left(\frac{\sqrt{\alpha'}}{r} \right)^{7-p'} . \tag{4.2}$$

On the visible sector probe Dp-brane world volume, these background bulk fields yield possible corrections to the potential and visible sector kinetic terms. Evaluating the Dp-brane Dirac-Born-Infeld action in these background fields

$$S_p = -T_p \int d^{p+1}x e^{-\phi} \sqrt{\det\left(h_{\mu\nu} + F_{\mu\nu}\right)}, \tag{4.3}$$

where $h_{\mu\nu}$ is the induced metric, yields

$$S_p \sim -T_p \int d^{p+1}x f(r)^{\frac{p'-3}{4}} f(r)^{-\left(\frac{p+1}{4}\right)} \left[1 + f(r) \left(\frac{1}{2} \partial_\mu X^i \partial^\mu X_i + \frac{1}{4} F_{\mu\nu} F^{\mu\nu} \right) + \cdots \right] \tag{4.4}$$

with indices now raised and lowered using the Minkowksi metric. For $p = p'$, the D-branes are in a 1/2 BPS configuration and preserve 16 supersymmetries. In this case, from (4.4), it is apparent that the Dp-brane world volume kinetic terms receive no corrections. The Kähler metric is flat and the Kähler potential is exact as required with 16 supersymmetries. (The correction to the potential term in (4.4) is canceled by the exchange of the RR p-form antisymmetric tensor field for $p = p'$). For $p = p' - 4$, the configuration is 1/4 BPS and preserve 8 supersymmetries. In this case, the dilaton contribution cancels the gravitational contribution to the potential, but there is a correction to the Dp-brane world volume kinetic terms. The flat inherited Kähler metric for an isolated Dp-brane, which in the absence of other D-branes would preserve 16 supersymmetries, is modified by the background fields generated by the Dp'-brane. The inherited Kähler potential is therefore modified by brane-brane interactions in the configuration with only 8 supersymmetries.

The background fields generated by the Dp'-brane which modify the Dp-brane Kähler metric are a condensate of closed string states. In terms of brane-brane interactions this amounts at the perturbative level to tree-level exchange of closed string states between the branes. This classical exchange amplitude may also be interpreted in the crossed channel as a one-loop quantum amplitude over open string states which stretch between the branes. In the open string language this corresponds to a one-loop correction to the Dp-brane kinetic terms. This is not surprising since one-loop corrections to the Kähler potential are allowed with 8 supersymmetries.

Modification of the Kähler metric and associated Kähler potential in the above example illustrates that corrections to the effective Kähler potential generally exist in D-brane configurations with less than 16 supersymmetries. From the expression for the modification of the kinetic terms in (4.4) there is clearly a direct coupling between the

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brane-brane separation r and the brane world volume fields. If this modulus acquires an F-component, then barring a cancellation, this contact interaction leads to tree-level visible sector scalar masses.

The breakdown of the sequestered intuition for the leading decoupling of separated branes may be understood in general D-brane world models as illustrated above in two equivalent ways. First, in the closed string channel tree-level couplings between the branes can in fact give rise to non-derivative brane-brane interactions. Second, in the open string channel, the quantum one-loop amplitude which involves extended string states stretching between the branes is not part of the low-energy local effective field theory description. The low energy arguments for volume suppression of point-particle quantum amplitudes therefore do not apply. In either channel, the leading inherited Kähler potential is also seen not to be protected from unsuppressed corrections in generic configurations.

In the closed string language warping of the background metric and dilaton also affects the propagation of the RR or vector potentials that generate the brane-brane contact interactions discussed in section 3. In theories with 8 or fewer supersymmetries this provides yet another source of modification to the Kähler potential. An example of this effect is presented in the next subsection.

4.2 The strongly coupled heterotic theory On a Calabi-Yau space

The Hořava-Witten background of M-theory provides a brane world model with end of the world branes with transverse co-dimension one. As discussed in section 2.1, for a Hořava-Witten compactification of M-theory on $S^1/\mathbb{Z}_2 \times \mathbb{T}^6$ the 16 unbroken supersymmetries fix the form of the Kähler potential. This Kähler potential includes brane-brane couplings which at low energy lead to non-derivative contact interactions between matter fields residing on the two end of the world branes. The origin of these brane-brane couplings may be understood in, for example, the eleven-dimensional limit as arising from exchange of the M-theory bulk four-form field strength, as described in section 3.1, or in the five-dimensional limit as arising from the exchange of vector bosons as described in section 3.2. In this description, it is reasonable to expect that in backgrounds with less supersymmetry brane-brane couplings arising from exchange of bulk fields persists in general. And since the Kähler potential is not as constrained with less supersymmetry the specific form of the brane-brane couplings may be more general. In the brane world picture this can arise because the form of the background and fluctuating bulk fields and metric are also less constrained with less supersymmetry.

In this subsection the leading form of the Kähler potential for a class of Calabi-Yau fibrations over $S^1/\mathbb{Z}_2$ M-theory backgrounds which preserve only 4 supersymmetries is illustrated. Calculable (in principle) flavor violating modifications to the lowest order Kähler potential and tree-level soft masses arise and may be traced to warping and distortion of the bulk space between the branes. From the bulk point of view these corrections are tree-level effects and not suppressed in all regions of moduli space where the brane world limit is obtained. So even if the leading tree-level contributions to the soft scalar masses happen to vanish in specific models, tree-level flavor violation generally arises at generic

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points on moduli space. And as discussed below these effects are not likely to be small in a Hořava-Witten model of nature.

For simplicity we consider Calabi-Yau compactifications in which the spin connection is embedded in the SU(3) $\subset E_8$ visible sector gauge connection [10]. This leaves an unbroken $E_6 \times E_8'$ gauge symmetry with $h^{1,1}(h^{2,1})$ generations(antigenerations) of $\mathbf{27}(\overline{\mathbf{27}}) \in E_6$ chiral matter on the visible sector brane and pure E_8' super Yang-Mills on the hidden sector brane. Hořava-Witten compactifications of this type have been extensively studied [11, 18, 25]. Since the hidden sector does not contain any chiral matter, this class of compactifications is not a viable model of BWSB arising from hidden sector matter. However, these examples are instructive in illustrating corrections to the lowest order form of the Kähler potential in brane world backgrounds with only 4 supersymmetries. Warping of the bulk metric in these examples induces corrections to the tree-level Kähler potential at order T/S. These corrections are also interesting in that they do not arise at any order in perturbation theory in the weakly coupled heterotic string limit.

More realistic models with hidden sector chiral matter can arise in Calabi-Yau fibrations over $S^1/\mathbb{Z}_2$ in which the spin connection is not embedded solely in the visible sector gauge connection. This requires in general turning on a hidden sector gauge connection and bulk four-form field strength. Corrections to the Kähler potential appearing at order T/S are also expected in these classes of models, as explained below.

For toroidal compactification of the Hořava-Witten background, the $S^1/\mathbb{Z}_2$ M-theory bulk interval separating the branes is flat. This is guaranteed by the 16 unbroken supersymmetries and remains exact in the full interacting theory. Physically, this follows from the fact that the end of world branes in a background with this much supersymmetry do not have any tension and therefore do not distort the bulk. For compactifications with less supersymmetry the brane tensions need not vanish and the bulk is not guaranteed to be flat. If the brane background sources are small, a systematic expansion for the bulk fields may be developed by expanding about the solutions for a flat interior [10].

For the class of Calabi-Yau compactifications considered here the brane tensions and background sources for the bulk four-form field strength vanish to zeroth order in an expansion in powers of $\kappa^{2/3}R_{11}/V_6^{2/3} \sim T/S$, where κ^2 is the eleven-dimensional Newton's constant and V_6 is the Calabi-Yau volume. At this order the $S^1/\mathbb{Z}_2$ interval is flat and the internal space is a direct product $S^1/\mathbb{Z}_2 \times \mathrm{CY}$. At higher orders in the $\kappa^{2/3}R_{11}/V_6^{2/3}$ expansion, non-vanishing brane fields act as sources for the metric and bulk four-form field strength through the modified Bianchi identity (3.2). In order to determine the four-dimensional Kähler potential it is important to keep both background (B)-brane fields resulting from the gauge connection embedding as well as fluctuating (F)-brane fields which represent visible sector fields. For the standard embedding, $\mathrm{tr}(F\wedge F)^{(\mathrm{B})}_{E_8} = \mathrm{tr}(R\wedge R)$ and $\mathrm{tr}(F\wedge F)^{(\mathrm{B})}_{E_8'} = 0$, the solution for the bulk three-form potential associated with the four-form field strength [11] with both background and fluctuating brane sources using the unperturbed flat metric is

$$C_{ABC} = \frac{\lambda}{12}\left[\frac{1}{2}\omega^{(\mathrm{B})}_{E_8} + \left(1-x^{11}\right)\left(\omega^{(\mathrm{F})}_{E_8} - \frac{1}{2}\omega^{(\mathrm{F})}_{L}\right) - x^{11}\left(\omega^{(\mathrm{F})}_{E_8'} - \frac{1}{2}\omega^{(\mathrm{F})}_{L}\right)\right]_{ABC}, \qquad (4.5)$$

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where $x^{11} \in [0,1]$ and $\mathrm{tr}(F \wedge F)_{E_8} = d\omega_{E_8}$ and $\mathrm{tr}(R \wedge R) = d\omega_L$. Note that the background brane gauge fields only induce a constant three-form potential in the bulk at this order. In fact with this embedding the background brane sources do not induce a background four-form field strength in the bulk at any order in the expansion, although the fluctuating fields do give rise to a fluctuating field strength.

The background brane fields do, however, lead to finite brane tensions. Four-dimensional $N=1$ supersymmetry along with a vanishing bulk four-form field strength resulting from the standard embedding imply that one-brane has positive tension while the other has negative tension of equal magnitude. This gravitational source warps the internal space which becomes a fibration of Calabi-Yau over the $S^1/\mathbb{Z}_2$ interval. At the first non-trivial order the induced perturbation of the eleven-dimensional metric grows linearly between the boundaries. The solution for the metric can be expressed in terms of R_{11}^0 and the $h^{1,1}$ Kähler moduli appearing in the low-energy theory. The metric for the Calabi-Yau fiber can be expanded in a fixed basis $\{\omega^n_{i\bar{j}}\}$ for the $(1,1)$ Kähler forms as

$$g_{i\bar{j}} = V_6^{1/3} \sum_{n=1}^{h^{1,1}} b^n \omega^n_{i\bar{j}} \,, \tag{4.6}$$

where $V_6 = V_6(x^{11})$ and the $b^n = b^n(x^{11})$ are the x^{11} dependent volume and Kähler moduli, related to the standard normalization by $\mathrm{Re}\, S = V_6$ and $\mathrm{Re}\, T_n = R_{11} b^0_n$. The $h^{1,1}$ moduli b^n have a normalization which is independent of V_6, R_{11}, and x^{11}, and are constrained to satisfy $d_{ijk} b^i b^j b^k = 6$ (and therefore only describe $h^{1,1} - 1$ independent degrees of freedom) where d_{ijk} are the Calabi-Yau intersection numbers. In terms of these moduli the metric for the Calabi-Yau fiber is given by (4.6) with moduli and volume varying along x^{11} as [11, 25]

$$\begin{aligned} b^i &= b^i_0 - \frac{R^0_{11}}{\sqrt{2}\,\mathrm{Re}\, V^0_6} \left(\gamma_i - \frac{2}{3} b^i_0 \sum_{k=1}^{h^{1,1}} \gamma_k b^k_0 \right) \left(x^{11} - \frac{1}{2} \right) , \\ V_6 &= V^0_6 \left(1 + \frac{3}{2} \sum_{n=1}^{h^{1,1}} \alpha_n \frac{\mathrm{Re}\, T_n}{\mathrm{Re}\, S} \left(x^{11} - \frac{1}{2} \right) \right) , \end{aligned} \tag{4.7}$$

where V^0_6 and b^i_0 are the unperturbed values and α_n and $\gamma_n = 3\alpha_n/2\sqrt{2}$ are numerical constants that depend topological data of the Calabi-Yau and may be found in [11, 25] and $x^{11} \in [0,1]$. The four-dimensional metric in geometric frame, and $S^1/\mathbb{Z}_2$ interval metric along the x^{11} direction, are given by

$$\begin{aligned} g_{\mu\nu} &= \left(1 + \sum_{n=1}^{h^{1,1}} \alpha_n \frac{\mathrm{Re}\, T_n}{\mathrm{Re}\, S} \left(x^{11} - \frac{1}{2} \right) \right) \eta_{\mu\nu} \,, \\ g_{11,11} &= \left(1 + 2 \sum_{n=1}^{h^{1,1}} \alpha_n \frac{\mathrm{Re}\, T_n}{\mathrm{Re}\, S} \left(x_{11} - \frac{1}{2} \right) \right) \left({R^0}_{11} \right)^2 . \end{aligned} \tag{4.8}$$

Because the brane tensions are equal in magnitude and opposite in sign the metric perturbations averaged over the $S^1/\mathbb{Z}_2$ M-theory direction vanish at this order for this class

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of compactifications with standard embedding. The zero-mode components of the four-dimensional moduli are therefore equal to the unperturbed values, $\langle V_6 \rangle = V_6^0$ and $\langle b^i \rangle = b_0^i$. It is important to note that the bulk warping not only modifies the Calabi-Yau volume along the M-theory direction, but also distorts its geometry along this direction in a manner which depends on the four-dimensional moduli.

The magnitude of the bulk warpings (4.7) and (4.8) induced by the brane tensions depends on the ratios of moduli $\epsilon_i = T_i/S$. This moduli dependence is easily understood from the typical magnitude of the brane sources $R \wedge R \sim 1/V_6^{2/3}$. Since for co-dimension one the magnitude of the perturbation grows with the brane separation, the total effect is proportional to $R_{11}/V_6^{2/3} = T/S$. Since this dependence is a consequence of co-dimension one, warpings for more general compactifications with unequal brane tensions resulting from the spin connection not embedded entirely in the visible sector gauge connection should have the same parametric dependence.

In the Hořava-Witten theory unification of four-dimensional gauge and gravity couplings is obtained for an average Calabi-Yau volume of order $S \simeq (3 \times 10^{16}\,\mathrm{GeV})^{-6}$. This, along with the value of the unified gauge coupling, imply that numerically $S^{1/6} \simeq 2\kappa^{2/9}$ and that the brane separation is of order $R_{11} \simeq 8\kappa^{2/9}$ [10, 15]. As advocated by Witten, one might hope that, given these parameters, a supergravity approximation to the bulk physics is roughly valid. The expansion parameter $\epsilon = T/S$ for the overall volume modulus $T = R_{11}V_6^{1/3}$ is then of order one [15]. The other moduli may generally be expected to have similar magntidue; and therefore none of the expansion parameters may be particularly small in a Hořava-Witten model with gauge coupling unification. Aside from the T_i/S dependence the numerical factors in (4.7) and (4.8) are all order one numbers which depend on the Calabi-Yau topology. So distortion of the Calabi-Yau and the associated modification of the inherited Kähler potential discussed below, is likely to be a significant effect in the Hořava-Witten brane world theory.

The four-dimensional Kähler potential for this class of compactifications with standard embedding can in principle be obtained from the brane field action arising in the perturbed background given above, just as for the unperturbed case discussed in section 3.1. There are two classes of corrections to the Kähler potential which are important in determining the visible sector soft masses arising from hidden sector supersymmetry breaking. First there are corrections to the chiral matter kinetic term wave functions, parameterized by $Z_{i\bar{j}}$ in (2.9), arising from the perturbations of the metric. These have been computed from the warping deformations (4.7) of the Calabi-Yau metric (4.6) for the class of standard embeddings considered here [11, 25]. There are also corrections to terms involving four chiral matter fields, parameterized by $Z_{i\bar{j}k\bar{l}}$ in (2.9). These arise from integrating out the bulk four-form in the warped background metric, and to our knowledge have not been computed.

The warping of the Calabi-Yau metric (4.6) at the position of the visible sector brane modifies the normalization of the brane field kinetic terms and therefore corrects the Kähler potential. Since the magnitude of the metric perturbations averaged over the $S_1/\mathbb{Z}_2$ M-theory direction vanishes to $\mathcal{O}(\epsilon_i)$ with standard embedding, to this order Weyl rescaling from geometric frame (obtained from compactification to four dimensions) to Einstein

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frame is the same as in the unperturbed flat bulk case. As discussed in section 3.1 the coefficients of the brane matter kinetic terms in the flat case are proportional to $1/T$ in four-dimensional Einstein frame. So warping of the internal metric at the position of the visible sector brane proportional to T/S modifies the visible sector brane matter kinetic terms in Einstein frame by an amount proportional to $1/S$. Since this is independent of R_{11}, this implies that corrections to the quadratic terms in the Einstein frame Kähler potential must be proportional to $1/S$ and involve ratios of T_j/T_i. In fact, Lukas, Ovrut, Stelle and Waldram find corrections [25] of just this form for the brane matter wave functions

$$Z_{i\bar{j}}(S,T_k) = Z^0_{i\bar{j}}(T_k) + \delta Z_{i\bar{j}}\left(S, \frac{T_k}{T_l}\right), \qquad \delta Z_{i\bar{j}}\left(S, \frac{T_k}{T_l}\right) = \frac{1}{S+S^\dagger} h_{i\bar{j}}\left(\frac{T_k}{T_l}\right), \tag{4.9}$$

where $Z^0_{i\bar{j}}(T_k)$ is the wave function in the unperturbed case without warping and is a function of the Kähler moduli but not the dilaton as discussed in section 2.1. The function $h_{i\bar{j}}(T_k/T_l)$ is an order one function determined by the lowest order Kähler potential of the $(1,1)$ moduli to be a function of ratios of these moduli. The explicit forms for these functions for the class of compactifications considered here have been computed in terms of topological data α_n and intersection numbers d_{ijk} of the Calabi-Yau fiber [25]. Note that since the wave function corrections do not depend on R_{11} in four-dimensional Einstein frame they do not depend on the overall volume modulus $T = R_{11}/V_6^{1/3}$.

A determination of the visible sector soft masses arising from hidden sector supersymmetry breaking also requires the corrections to the quartic terms in the warped background, $\delta Z_{i\bar{j}a\bar{b}}$, appearing in the Kähler potential (2.9) where i,j refer to visible sector and a,b refer to hidden sector. In principle these couplings are determined by integrating out the four-form field strength with brane Yang-Mills Chern-Simons sources localized on the branes. However, since the background bulk four-form field strength vanishes exactly for compactifications in which the spin connection is embedded in the gauge connection the quartic couplings can arise only from integrating out the fluctuating part of the bulk four-form field strength sourced by brane field fluctuations. This involves integrating over the internal metric. But since with standard embedding the metric perturbations averaged over the $S^1/\mathbb{Z}_2$ interval vanish to $\mathcal{O}(\epsilon_i)$ the quartic couplngs are identical to those of the unperturbed case at this order. Corrections could however arise at $\mathcal{O}(\epsilon_i^2)$. For a compacitification with more general embedding with unequal brane tensions and non-vanishing background bulk four-form field strength, corrections to the lowest order quartic couplings would however be expected to arise at order T_i/S since the metric perturbations do not average to zero in this case. As discussed in section 3.1 the coefficients of the brane-brane quartic Kähler potential couplings in the flat case are proportional to $1/T^2$ in four-dimensional Einstein frame. So for a compactification with general embedding the brane-brane interactions in Einstein frame should be modified by an amount proportional to $1/TS$. For a standard embedding these corrections would vanish at this order, and the next order corrections would modify the quartic couplings by an amount proprotional to $1/S^2$. For a general embedding the Einstein frame quartic Kähler potential couplings should then be modified

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by warping and distortion of the internal metric by an amount

$$\begin{aligned} Z_{i\bar{j}a\bar{b}}(S,T_k) &= Z^0_{i\bar{j}a\bar{b}}(T_k) + \delta Z_{i\bar{j}a\bar{b}}\left(S,\frac{T_k}{T_l}\right), \\ \delta Z_{i\bar{j}a\bar{b}}\left(S,\frac{T_k}{T_l}\right) &= \frac{1}{(S+S^\dagger)(T+T^\dagger)} j_{i\bar{j}a\bar{b}}\left(\frac{T_k}{T_l}\right), \end{aligned} \tag{4.10}$$

where $Z^0_{i\bar{j}a\bar{b}}(T_k)$ is the Kähler potential quartic coupling in the unperturbed case without warping and is a function of the Kähler moduli but not the dilaton as discussed in section 2.1. The function $j_{i\bar{j}a\bar{b}}(T_k/T_l)$ is determined by the lowest order Kähler potential of the $(1,1)$ moduli to be a function of ratios of these moduli. For general embeddings it should be an order one function which could in principle be determined from topological data of the Calabi-Yau manifold. For standard embedding $j_{i\bar{j}a\bar{b}}(T_k/T_l)$ includes at least an explicit factor of ϵ. To our knowledge the lowest order quartic terms even without warping have not been calculated for a Calabi-Yau fibration with any embedding.

With hidden sector supersymmetry breaking, the visible-hidden quartic couplings in Einstein frame for the general Kähler potential (2.9) to lowest order in the expansions described above for a general Calabi-Yau fibration over $S^1/\mathbb{Z}_2$ are, from the associated supergravity f function (2.10), proportional to

$$\frac{1}{3}\left(Z^0_{i\bar{j}}Z^0_{a\bar{b}} + \frac{1}{S+S^\dagger}\left(h_{i\bar{j}}Z^0_{a\bar{b}} + Z^0_{i\bar{j}}h_{a\bar{b}}\right)\right) - \left(Z^0_{i\bar{j}a\bar{b}} + \frac{1}{(S+S^\dagger)(T+T^\dagger)}j_{i\bar{j}a\bar{b}}\right). \tag{4.11}$$

In order to obtain the physical mass squared matrix the scalar kinetic term wave functions factors

$$Z^0_{i\bar{j}} + \frac{h_{i\bar{j}}}{S+S^\dagger}, \qquad Z^0_{a\bar{b}} + \frac{h_{a\bar{b}}}{S+S^\dagger} \tag{4.12}$$

must be rescaled to give canonically normalized fields. As discussed in section 3.1, in four-dimensional Einstein frame the coefficients of the brane matter kinetic terms in the flat case are proportional to $1/T$ while the quartic couplings are proportional to $1/T^2$. To lowest order, extracting this dependence and rescaling the fields by a compensating amount, the dependence on the expansion parameter $(T+T^\dagger)/(S+S^\dagger)$ is apparent. In order to obtain the physical masses, the mass squared matrix must be diagonalized and canonically normalized, which in general requires a unitary rotation and rescaling. In a basis which both the visible and hidden sector are diagonalized at zero-th order in the expansion, the mass squared matrix is proportional to

$$\left(\frac{1}{3}\delta_{i\bar{j}} - \tilde{Z}_{i\bar{j}a\bar{b}}X_{a\bar{b}}\right) + \frac{T+T^\dagger}{S+S^\dagger}\left[\frac{1}{3}\tilde{h}_{i\bar{j}} + \left(\frac{1}{3}\tilde{h}_{a\bar{b}} - \tilde{j}_{i\bar{j}a\bar{b}}\right)X_{a\bar{b}}\right], \tag{4.13}$$

where the tilde functions are related to the previous ones by rotation and rescaling, and $X_{a\bar{b}} = F_aF^*_{\bar{b}}/\,\mathrm{tr}_{c\bar{d}}(F_cF^*_{\bar{d}})$. Even if the lowest order masses arising from the first terms in paranthesis are universal, as would occur with a no-scale Kähler potential, they are not universal in general. The leading $\mathcal{O}(\epsilon_i)$ effects arise by diagonalization at this order through rescalings which lead to non-degenerate masses at the same order. Given that the functions inside the square paranthesis are unrelated in any simple way, there is no

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reason for the non-degenerate contributions from these terms to cancel. The upshot is that generically there are irreducible non-degenerate contributions to the visible sector masses appearing at $\mathcal{O}(\epsilon_i)$. However, without a theory of flavor this rescaling need not be aligned with the quark and lepton mass eigenstates, and would in general introduce dangerous supersymmetric flavor violation, since the expansion parameters T_i/S are not particularly small in Hořava-Witten models. Geometrically the flavor violation arises because the bulk warping (4.7) distorts the Calabi-Yau geometry (4.6) at the position of the brane in a manner depending on the four-dimensional moduli and is therefore in general not aligned with the lowest order metric or matter field wave functions.

For Calabi-Yau manifolds with the standard embedding of the spin connection in the visible sector gauge connection with $h^{1,1} = 1$ the Kähler potential including the lowest order effects of warping can be obtained in a simple form [11]

$$K = -\ln\left(S + S^{\dagger} - \epsilon Q^{\dagger} Q\right) - 3\ln\left(T + T^{\dagger} - Q^{\dagger} Q\right). \tag{4.14}$$

where ϵ is the fixed background value $(T+T^{\dagger})/(S+S^{\dagger})$ evaluated on the unperturbed background, while S and T appearing (4.14) including fluctuating pieces. With the standard embedding the corrections to the quartic terms vanishes to this order as discussed above, and the only warping effect is the overall normalization of the visible sector matter wave function. The lowest order no-scale form is modified at first non-trivial order. Although with standard embedding this class of backgrounds is not useful for hidden sector BWSB since the hidden sector does not contain any chiral matter, this modification of the Kähler does demonstrate that the lowest order Kähler is modified in a non-trivial way by warping as expected.

The tree-level warping and distortion corrections to the Kähler potential also have implications for a dilaton dominated scenario where only the dilaton auxillary component is non vanishing (or dominant). To lowest order the visible sector scalar fields acquire from the Kähler potential (2.9) with (4.9) a universal soft mass. But to next leading order the second term in (4.9) with non-vanishing dilaton auxillary component generates a general matrix in flavor space proportional to $h_{i\bar{j}}|F_S|^2$ and this is only suppressed only by T_i/S. Although the kinetic terms given by $Z_{i\bar{j}}$ are also not diagonal, the point is that generically these two matrices are not simultaneously diagonalizable. Again this flavor violation is not small. So we see in the Hořava-Witten brane world theory without additional assumptions about flavor symmetries, supersymmetric flavor violation is not particularly suppressed. Because the bulk is co-dimension one the warping of the Calabi-Yau grows with transverse distance and these effects become larger if the brane separation R_{11} is increased, but become smaller if V_6 is increased.

This again illustrates that brane world realizations of hidden sector supersymmetry breaking alone do not provide a solution of the supersymmetric flavor problem. In the absence of assumptions about flavor symmetries, supersymmetric flavor violation is not necessarily suppressed over much of moduli space, just as for standard hidden sector supersymmetry breaking scenarios. In addition, the flavor violating effects are tree-level effects from the bulk point of view. So the lowest order tree-level Kähler potential is not

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protected in any way from recieving tree-level flavor violating corrections over much of moduli space. In the Hořava-Witten theory with moduli expectation values which allow standard gauge coupling unification these effects are significant.

Finally, it should be emphasized that the lowest order form of the brane-brane couplings discussed here as well as the corrections arising from warping and distortion of the internal geometry are obtained in the eleven-dimensional supergravity limit. In this limit the order T_i/S corrections result from the finite M-theory interval separating the branes. Now effects associated with the M-theory direction do not appear at any order in perturbation theory in the heterotic string theory limit. So this class of non-holomorphic corrections to the Kähler potential are not visible in perturbative heterotic string theory. As long as the brane separation R_{11} and Calabi-Yau volume V_6 and inverse curvature are large in eleven-dimensional Planck units, the supergravity approximation should be good and quantum M-theory corrections should be at least power law suppressed in these parameters. There are in additional semi-classical non-perturbative effects due to M2- and M5-brane instantons which are exponentially suppressed. These are visible in the perturbative heterotic limit as world sheet and gauge theory instantons as discussed in section 2.1.

4.3 Kähler potential corrections with pure five-dimensional supergravity

Pure five-dimensional supergravity with a flat unwarped interval separating end of the world branes can give a low energy four-dimensional no-scale Kähler potential of the separable sequestered form [2, 8], at least in the case of inheritance from the $N = 4$ form, as argued in section 2.3. A flat bulk is only obtained, however, for vanishing brane and bulk tensions. Non-vanishing brane tension will generically warp the bulk metric. The presence of a stabilizing mechanism will also contribute to the warping due to non-zero bulk stresses. So in general the four-dimensional metric is expected to be a function of the coordinate, y, transverse to the branes

$$g_{\mu\nu}(y) = (1 + \xi(y))\eta_{\mu\nu} \,, \tag{4.15}$$

where $\xi(y)$ is a model dependent function. The warping of the metric (4.15) at the positions of the branes gives a universal modification of the brane field kinetic terms. For perturbations about a flat interior the warping will typically grow linearly with the brane separation due to the one-dimensional geometry. This warping induced modification will be reflected in a modification of the four-dimensional Kähler potential. The lowest order no-scale Kähler potential (2.22) inherited from the five-dimensional theory is then generally modified. This is consistent with the fact that a general warped compactification only preserves $N = 1$ supersymmetry in four dimensions, so the inherited form of the Kähler potential is not protected in any way. Corrections to the inherited no-scale Kähler potential are small if the brane tensions are small.

The pure five-dimensional supergravity with end of the world branes scenario is very similar to the five-dimensional limit of the Hořava-Witten theory discussed in section 3.2. Here, however, there is by assumption no dilaton. In addition there is only a single bulk vector boson which is part of the five-dimensional gravitational supermultiplet, whereas in the Hořava-Witten theory on a generic background there are a number of additional

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vector bosons. In pure five-dimensional supergravity, exchange of the single vector boson between the branes can be understood as giving rise to brane-brane current-current couplings contained in the no-scale Kähler potential [8], while in the Hořava-Witten theory exchange of the additional vector bosons can be understood as giving rise to the more general non-sequestered form of the inherited Kähler potential which contains non-derivative brane-brane contact interactions, as discussed in section 3.2. Just as in the Hořava-Witten theory, warping of the bulk geometry should lead to corrections to the lowest order inherited form of the Kähler potential. And as described below, in the pure five-dimensional case in fact warping generally does leads to visible sector tree-level scalar masses, even though the masses vanish in the flat case.

In the Hořava-Witten theory with bulk warping discussed in the previous subsection, it was possible to argue for the form of the four-dimensional Kähler potential to first order in perturbation about a flat background by inspecting the form of the brane field interactions in the underlying theory. In this pure five-dimensional supergravity scenario this is not possible without a consistent underlying theory. We therefore assume an *ansatz* for the four-dimensional Kähler potential with bulk warping of the form

$$\begin{aligned} K = -3\ln\Big(f_{\rm mod}\left(T+T^\dagger\right) - \left(1+h_{\rm vis}\left(T+T^\dagger\right)\right)\,{\rm tr}\,Q_i^\dagger Q_i - \\ - \left(1+h_{\rm hid}\left(T+T^\dagger\right)\right)\,{\rm tr}\,\Sigma_i^\dagger \Sigma_i,\Big) \end{aligned} \tag{4.16}$$

where $h_{\rm vis}$ and $h_{\rm hid}$ are assumed to be universal factors for each field localized on a given brane, and $f_{\rm mod}$ is the total compact volume, and T is the radion modulus. These are model dependent functions of the brane tensions and bulk stress-energy momentum tensor. For a flat bulk $h_{\rm vis}(x) = h_{\rm hid}(x) = 0$ and $f_{\rm mod}(x) = x$, resulting in the no-scale form. Kähler potentials of this type have been obtained [26, 27] in the supersymmetric version of the Randall-Sundrum model with two-branes [28], where the warping is obviously an important effect.

The Kähler potential (4.16) contains a direct coupling between the radion or volume modulus, T, and visible sector brane fields Q_i. In the geometric frame obtained by direct compactification to four dimensions this follows from the warping induced dependence of the visible sector kinetic terms on the compact volume. This survives unchanged in the supergravity frame since these frames coincide for compactification from five to four dimensions. If the radion obtains an auxiliary expectation value, $F_T \neq 0$, this direct coupling leads to a tree-level visible sector scalar mass. As discussed below in more detail this is in fact the case in most scenarios for stabilizing the bulk geometry. From (4.16) the ratio of visible sector scalar masses to the gravitino mass for both radion and hidden sector auxiliary fields, $F_T \neq 0$ and $F_\Sigma \neq 0$, is

$$\frac{m_Q^2}{m_{3/2}^2} = \frac{(T+T^*)^2}{9(1+h_{\rm vis})\left(f_{\rm mod}'^2 - f_{\rm mod}f_{\rm mod}''\right)^2}\left(\frac{(h_{\rm vis}')^2}{1+h_{\rm vis}} - h_{\rm vis}''\right)\left(\frac{|F_T|^2}{|W|^2}(T+T^*)^2\right), \tag{4.17}$$

where here $' \equiv \partial/\partial\,{\rm Re}\,T$.

The expression for the soft masses (4.17) has a number of interesting features. First the masses are universal since by assumption of the ansatz the warping modification of

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kinetic terms is the same for all fields on the visible sector brane. Universality would, however, not be preserved for visible sector fields propagating on different branes. And the h functions in general need not even be universal for fields propagating on the same brane. Second, the overall magnitude of the tree-level mass (4.17) is clearly model dependent, but can be sizable. The vanishing cosmological constant condition gives an upper limit on the last term in parenthesis of $2\,\mathrm{Re}\,T|F_T|/|W| \leq 3$. The masses do, however, vanish even with $F_T \neq 0$ if the condition condition $(h'_{\mathrm{vis}})^2 = (1 + h_{\mathrm{vis}})h''_{\mathrm{vis}}$ is satisfied. This arises for a flat interior $h_{\mathrm{vis}}(x) = 1$ as well as for an AdS intertior $h_{\mathrm{vis}}(x) = e^{-x} - 1$ (a boundary interpretation of the latter case is presented in subsection 4.4). However, the requisite dynamics which stabilizes the T modulus will generically create stresses which lead to warping of the bulk. In general this will not preserve the flat interior metric or result in a pure AdS warping. Third, even though the auxillary component for the T modulus, F_T, contributes directly to the scalar masses, auxillary hidden sector auxiliary components, F_Σ, do not give rise directly to tree-level visible sector masses, as a consequence of the assumed ansatz for the Kähler potential (4.16). This is apparent in supergravity frame where the supergravity f function associated with (4.16) is separable and therefore does not give rise to non-derivative brane-brane couplings between visible and hidden sector fields. So in this sense the hidden sector does remain sequestered with the Kähler potential ansatz (4.16), and it is only the indirect effects of warping and a non-vanishing F_T which lead to scalar masses. Finally, the masses are directly proportional to the radion auxillary expectation value, F_T.

Since the auxillary component of the T modulus, F_T, is crucial in determining the scale for the resulting soft scalar masses it is instructive to consider its magntidue including the effects of stabilization and vanishing cosmological constant. To see this consider the supergravity potential including the effects of hidden sector supersymmetry breaking. The derivative of the potential with respect to T, assuming a canceled cosmological constant is

$$0 = V' = e^K \Big(K^{TT^*} F_T^* F_T' + K^{TT^*} F_T K_{TT^*} W^* + K^{TT^*\prime} F_T^* F_T + \\ + K^{\Sigma\Sigma^*\prime} |F_\Sigma|^2 - 3W^* W' \Big) , \tag{4.18}$$

where $\langle\Sigma\rangle \ll 1$ is assumed, and $' = \partial/\partial T$, and the superpotential is assumed to factorize as $W = W(T) + W(\Sigma)$. An extremum of the potential with vanishing radion auxiliary component would require $K^{\Sigma\Sigma^*\prime}|F_\Sigma|^2 = 3W^*W'$. Combining this with the condition for vanishing cosmological constant, $K^{\Sigma\Sigma^*}|F_\Sigma|^2 = 3|W|^2$, implies $K^{\Sigma\Sigma^*\prime}/K^{\Sigma\Sigma^*} = -K'$ which is not satisfied for the no-scale potential with the prefactor 3. So one finds that $F_T = 0$ is not an extremum of the potential with hidden sector supersymmetry. So assuming vanishing cosmological constant with $F_\Sigma \neq 0$, then F_T must be non vanishing in the stable ground state.

The overall magnitude of the radion auxiliary component depends on whether the radion is stabilized in the globally supersymmetric limit or only including supergravity effects. First consider the case in which it is stabilized in the globally supersymmetric limit in the absence of hidden sector supersymmetry breaking. In this case $\langle W'\rangle = 0$ and $F_T = 0$ in this limit. So near the ground state minimum $\langle T\rangle = T_0$ the superpotential can

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be expanded as $W = \frac{1}{2}M(T - T_0)^2 + \cdots$. One can then verify that with hidden sector supersymmetry breaking turned on $W' = 0$ is no longer a solution. This is clear from the following equivalent expression for V' obtained by substituting for F'_T

$$0 = V' = e^K \left(K^{TT^*} F^*_T \left(W'' + K_{TT} W + K_T W' \right) + K^{TT^*} K_{TT^*} W^* F_T + \right.$$
$$\left. + K^{TT^*\prime} |F_T|^2 + K^{\Sigma\Sigma^*\prime} |F_\Sigma|^2 - 3W^* W' \right). \tag{4.19}$$

Now suppose that $W' = 0$ which implies $F_T = K'W$. Then with $W \sim \mathcal{O}(m_{3/2})$, all the remaining terms in this equation are $\mathcal{O}(m^2_{3/2})$ except the one involving W'' which is $\mathcal{O}(m_{3/2})$. For $M \gg m_{3/2}$ (which is the case if globally supersymmetric dynamics dominates the stabilization of the radion) this term cannot possibly cancel the other terms. So $W' \neq 0$ which implies that $F_T \neq 0$ with the superpotential above. In this limit with $M \gg m_{3/2}$ one then expects that an expression for the extremum can be written as an expansion in terms of the of the radion auxiliary component, derivatives of the superpotential, and $\Delta \equiv (T - T_0)/T_0$ in powers of $m_{3/2}$. Having established that F_T and W' are both non zero, one might guess that $F_T \sim \mathcal{O}(m_{3/2})$. An inspection of (4.19) indicates that this is not possible for essentially the same reasons that W' could not vanish. Namely, all the terms are manifestly $\mathcal{O}(m^2_{3/2})$ except the term $F^*_T W''$ which would be $\mathcal{O}(m_{3/2})$. In fact, solving (4.19) for Δ implies $T - T_0 = 3\langle W \rangle / 2M \,\mathrm{Re}(T) + \cdots$ which leads to a vanishing radion auxiliary component at lowest order. At next order one finds $F_T \sim \mathcal{O}(m^2_{3/2}/M)$. Thus radion stabilization in the global supersymmetric limit leads to a non zero but negligible radion auxiliary component. So in this case the soft scalar masses resulting from warping would be insignificant.

Another possibility is that the radion is stabilized only in the presence of supersymmetry breaking and only including supergravity interactions. Examples of this type have been studied [8, 29]. Inspecting (4.19) one finds that a consistent solution with the superpotential $\langle W \rangle$, its derivatives, $\langle W' \rangle$, and the radion auxiliary component F_T, all order $\mathcal{O}(m_{3/2})$ at the extremum is possible. Of course, the dependence of the radion auxiliary component on the volume is model-dependent. In the previous example with the radion stabilized in the supersymmetric limit the obstruction to finding such a solution appears to be that there $W'' \sim M$. An example of a model in which the radion is stabilized by the supersymmetry breaking and supergravity interactions with $F_T \sim \mathcal{O}(m_{3/2})$ is provided by both bulk and boundary gaugino condensation which in the presence of hidden sector supersymmetry breaking stabilizes the radion [29]. In models of this type warping of the bulk geometry generally induces tree-level visible sector masses through the radion auxiliary component.

4.4 Conformal sequestering from geometry

The question of what classes of models have vanishing or suppressed tree-level soft masses even with hidden sector supersymmetry breaking is an interseting one. In these cases if the scalar masses do in fact vanish then anomaly mediation can be important [2, 3]. As we have seen, vanishing tree-level masses do not generally arise from BWSB, but can be achieved in specific models. In particular, pure five-dimensional supergravity with a flat interior separating end of the world branes gives vanishing tree-level masses [2, 8],

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at least by inheritance if this case can be obtained by a string or M-theory orbifold as argued in section 2.3. Vanishing masses are not protected once the effects of warping and radion stabilization are included, and tree-level masses generally result, as discussed in the previous subsection. However, the soft masses (4.17) for the pure five-dimesional case do in fact vanish for arbitrary F_T if the special condition

$$\left(h'_{\rm vis}\right)^2 = (1+h_{\rm vis})h''_{\rm vis} \tag{4.20}$$

is satisfied, where $h_{\rm vis}$ is the function parameterizing the bulk warping in the Kähler potential ansatz (4.16). This condition is satisfied not only for a flat internal metric, but also for an internal metric which is a section of AdS space with exponential warping

$$m_Q^2 = 0\,, \qquad \text{for } 1+h_{\rm vis}\left(T+T^\dagger\right) = \begin{cases} 1\,, & \text{flat}\,, \\ e^{-k(T+T^\dagger)}\,, & \text{AdS}\,. \end{cases} \tag{4.21}$$

An exponential warping of this form occurs in the supersymmetric version of the Randall-Sundrum model with two branes [28, 27, 26].

The vanishing visible sector scalar masses which result with a co-dimension one AdS bulk can probably be traced to the existence in the five-dimensional bulk of Anti-de Sitter supersymmetry which is related in four dimensions to superconformal symmetry. This relation can be made precise through the AdS/CFT correspondence between a boundary four-dimensional conformal field theory and a five-dimensional anti-de Sitter space. This correspondence suggests that there may be a boundary description entirely in terms of four-dimensional superconformal field theory dynamics for the vanishing of scalar masses with a five-dimensional pure AdS bulk [30]. Of course a precise correspondence between a boundary field theory and large low curvature bulk is only good at large g^2N. But it is still interesting to explore the heurtistic correspondence between the geometric brane world example considered here and four-dimensional field theories based solely on approximate conformal symmetry.

Two classes of field theory models which are approximately superconformal have been suggested for suppressing visible sector scalar masses [31, 30]. Both rely on the observation that in global supersymmetry the soft mass squared for the scalar component of a chiral multiplet may be represented as the highest component of the wave function factor treated as a vector superfield. $m^2 = Z|_{\theta^2\bar\theta^2}$. Now field theory dynamics which is approximately conformal over some range of energies and couples to the wave function factor results in an approximately constant anomalous dimension γ for the wave function over this range, $Z = Z_0(\mu/\mu_0)^\gamma$, where μ is the renormalization group scale and μ_0 refers to the scale below which the theory is approximately conformal. As [31] demonstrates, if all non-$U(1)_R$ symmetries are explicitly broken, then the highest component, namely the soft mass, likewise also has an approximately constant anomalous dimension over this range. This results in a potentially sizeable suppression of the soft mass.

The first class of field theory models proposed by Nelson and Strassler utilize strong dynamics which is approximately conformal over some range and couples directly at tree-level to some of the visible sector scalars [31]. The second class of models proposed by Luty and Sundrum assume a hidden sector which is strongly interacting and conformal

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down to a low scale, and couples only to the weakly coupled visible sector fields through non-renormalizable operators represented by the wave function factor in the low energy theory of the hidden sector [30].

The AdS/CFT correspondence suggests a five-dimensional bulk geometric interpretation of these models [30]. In the AdS/CFT correspondence, five-dimensional anti-de Sitter space with exponential warping corresponds to a conformal field theory in the four-dimensional boundary theory. In the present case the end of the world branes break the full anti-de Sitter symmetry which in the boundary description correspond to UV and IR mass scales of the theory. In this interpretation the end of the world branes may be thought of as the UV and IR ends points of the conformal renormalization group flow. Interactions which couple the visible and hidden sectors in the boundary field theory, $(1/M^2)\int d^4\theta\; Q^\dagger Q \Sigma^\dagger \Sigma$ in global supersymmetry, are suppressed by a large mass scale M and so can be thought of as generated by physics residing on the UV-brane. The mass scale M then determines the position of the UV-brane.

In the Nelson-Strassler models the strongly coupled approximately conformal dynamics which couples directly to the standard model corresponds in the bulk to the AdS warping. In the Luty-Sundrum models it is the strongly coupled approximately conformal hidden sector which corresponds in the bulk to the AdS warping. In both cases the exponential warping in the AdS bulk between the UV- and IR-branes corresponds to the exponential suppression of the soft mass operators $m^2 = Z|_{\theta^2\bar\theta^2}$ in the boundary theory between the UV and IR limits of the conformal renormalization group flow. Since the boundary theories are only approximately conformal over some range the corresponding bulk geometry should only be approximately AdS with some residual warping near the UV and IR walls where conformal invariance is violated. This deviation from AdS would presumably result in only exponentially suppressed masses from (4.17) in the bulk geometric picture as required by the bulk/boundary correspondence.

In the Luty-Sundrum conformal suppression mechanism there is potentially a puzzle in translating the anomalous dimension effects of the hidden sector wave functions in the boundary theory written in globally supersymmetric language as a D-term quartic interaction between the visible and hidden sectors, into an operator written in locally supersymmetric language. In particular, the boundary theory suppression of D-term quartic interactions between the visible and hidden sectors in global supersymmetry applies in the locally supersymmtric theory to the suppression of quartic operators in the IR in which frame — supergravity or Einstein frame? Since the supergravity frame is manifestly supersymmetric this seems the most natural choice. And as discussed in section 2, vanishing of non-derivative quartic couplings in this frame results in vanishing scalar masses. A suppression in the IR of the couplings appearing in the Kähler potential would with supersymmetry breaking in the *visible* sector lead to universal masses in the *hidden* sector of order the gravitino mass. Since this would not give a picture consistent with the results of [31], the suppression occurs in the supergravity frame.

In the case of a flat five-dimensional bulk discussed in section 3.2, vanishing tree-level scalar masses are only obtained if the only vector boson in the bulk is that of the minimal supergravity multiplet. Additional bulk vector multiplets coupled to brane

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matter Cherns-Simons forms through a modified Bianchi identity do give rise to tree-level masses from hidden sector supersymmetry breaking. Although an analogous calculation for the case of a five-dimensional AdS with additional vector multipets coupled to brane matter as above has not been completed, it should be expected that also in this case tree-level masses arise from integrating out the bulk vectors (more on this below).

In the case of a flat interior the existence of bulk gauge symmetries under which the five-dimensional vector bosons transforms is closely related to the existence of unbroken global symmetries which act on the brane matter and vice versa. This is at least true in orbifold compactifications for which bulk five-dimensional gauge symmetries arise microscopically from higher-dimensional diffeomorphisms are in one to one correspondence with global symmetries on the D-brane which have the same origin. And this correspondence is likely to hold also for more general compactifications. So in the flat interior case vanishing tree-level masses require no gauge symmetries in the bulk (beyond that related through the minimal supergravity multiplet) or equivalently that the D-brane matter not possess any global symmetries.

Now in the AdS/CFT correspondence, on very general grounds and by construction, global symmetries in the boundary theory are realized as local gauge symmetries in the bulk. In the boundary field theory description the existence of an unbroken non-anomalous global symmetry has an important effect on the mechanism of conformal suppression of soft masses [31]. The anomalous dimension for an exactly conserved global current vanishes. So for every abelian global symmetry in the conformal field theory (which commutes with supersymmetry — see below) there is one eigenvector (determined by the global charges) of the (visible or hidden) soft mass squared matrix which is not suppressed by the conformal suppression mechanism. In the AdS/CFT correspondence such global symmetries in the conformal field theory are realized by local symmetries in the bulk with concomitant vector bosons.

So in the bulk AdS interpretation of the field theory conformal suppression mechanism, vanishing scalar masses require the absence of global symmetries acting on the conformal sector matter. In the field theory interpretation, if the goal is vanishing tree-level scalar masses even with hidden sector supersymmetry breaing, the model building problem is to find theories with the conformal suppression mechanism but no global symmetries. In the bulk language the problem is to find compactifications which reduce to pure five-dimensional Anti-de-Sitter supergravity with end of the world branes and no bulk vector multiplets. Which is easier may depend on the model builder.

As a final comment about the bulk-boundary correspondence between these pictures of conformal sequestering, note that superconformal invariance of an $N = 1$ four-dimensional boundary theory requires the existence of an unbroken $\mathrm{U}(1)_R$ symmetry. This current is related by superconformal symmetry to the dilation current, and since it does not commute with supersymmetry it does not leave an eigenvector of the scalar mass squared matrix unsuppressed. In the bulk picture, anti-de Sitter supersymmetry requires the existence of a single U(1) gauge boson in the minimal gravity multiplet, and couples on the boundary to the $\mathrm{U}(1)_R$ current. And (at least in the flat interior case) does not give non-derivative

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couplings between the branes which would lead to tree-level soft masses. So in both pictures the one abelian symmetry which is required by supersymmetry does not give rise to soft masses, but additional ones would.

5. Massive moduli

The lowest order inherited Kähler potentials in all the brane world scenarios presented here (with the exception of pure five-dimensional supergravity) are not sequestered and contain brane-brane non-derivative contact interactions which give rise to unsuppressed tree-level visible sector scalar masses from hidden sector supersymmetry breaking. All these contain, in the absence of supersymmetry breaking, four-dimensional massless moduli such as the $T_{i\bar{j}}$. Although it seems a priori unlikely, one might wonder if stabilizing these moduli in the four-dimensional effective theory could modify the form of the inherited Kähler potential in such a way that tree-level masses are not induced and that the sequestered intuition is recovered. In this section we show that the opposite is the case. Namely, without stabilizing the moduli the tree-level masses vanish. This feature is not restricted to the special sequestered Kähler potential but applies the more general class of no-scale Kähler potentials and appears to be a consequence of the auxillary equations of motion of the moduli. When the moduli are stabilized however, it is seen that the most naive expectation that with hidden sector supersymmetry breaking the moduli can simply be replaced with scalar expectation values is correct.

There are two possibilities for the scales associated with stabilizing the moduli. The first is that the moduli gain masses well below the compactification scale. The analysis of the previous sections within the low energy four-dimensional theory is then unaltered since the moduli are part of this theory, and the Kähler potential and resulting patterns of tree-level scalar masses are obtained assuming moduli stabilization as detailed below. Alternately one might hope that the moduli gain a mass much larger than the compactification scale. In this case, a different low energy theory might in principle result. This latter option seems unlikely, however, at least for moduli which arise from dimensional reduction of supermultiplets which are part of the higher-dimensional theory and appear in the effective theory above the compactification scale. These moduli must be present in the bulk space and are protected from gaining a mass parameterically larger than the compactification scale by gauge or supersymmetries of the underlying theory. Even for other moduli which do not arise directly from extended supermultiplets, stabilization or projection mechanisms generally give masses at most of order the compactification scale. At best it might be possible to consider compactifications with more than one geometric compactification scale. The minimal set of moduli in the low energy four-dimensional theory are then just those that arise from extended supermultiplets in the effective theory just above the lowest compactification scale. Other moduli (or vector multiplets) could gain masses at the higher compactification scales and not appear in the five-dimensional effective theory. The most extreme possible example of this type is a compactification of a fundamental theory of gravity to pure five-dimensional supergravity at a mass scale well above the compactification scale from five to four dimensions. Since by assumption all

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moduli which might be present in the five-dimensional theory are lifted by the higher mass scale compactification, the bulk contains no additional vector or hyper multiplets, and the only modulus which survives in the four-dimensional theory is the radion supermultiplet which arises from fields in the five-dimensional supergravity multiplet.

In order to investigate the influence of a moduli stabilizing potential consider for simplicity the inherited Kähler potentials of the types discussed above with only diagonal moduli. These Kähler potentials may be written in the form

$$K = -\sum_{i=1}^{n} p_i \ln \left(T_i + T_i^\dagger - \operatorname{tr} Q_i^\dagger Q_i - \operatorname{tr} \Sigma_i^\dagger \Sigma_i \right) , \tag{5.1}$$

where the sum of logarithms form (2.12) is obtained for $n = 3$ and $p_i = 1$, while the no-scale sequestered form (2.21) is obtained for $n = 1$ and $p = 3$. Also assume that the moduli are stabilized by superpotential interactions of the form

$$W = W(\Sigma) + W(T_i) \, . \tag{5.2}$$

Together with the Kähler potential this superpotential is assumed to give a mass to both the real and imaginary components of T_i. This may or may not require supersymmetry breaking. In the Einstein frame the scalar potential is given by the standard formula

$$V = e^K \left(K^{i\bar{j}} F_i F_{\bar{j}}^* - 3|W|^2 \right) = e^K \left(V_0(T_i, \Sigma_j, W) + \sum_i \lambda_i(T_j, W) Q_i^\dagger Q_i + \cdots \right) , \tag{5.3}$$

where in the second expression the dependence on the visible scalars to second order in the fields has been isolated, and where the coefficients λ_i are given below. Up to a wavefunction renormalization the coefficients of $Q_i^\dagger Q_i$ in the potential (5.3) are the scalar masses for those fields.

There is an important subtlety here because the field redefinition for T_i (3.12) also contains a dependence on the scalars, and one might worry whether this also contributes to their masses. When the moduli are massive this concern is spurious, and it is straightforward to see that this dependence does not contribute to the scalar masses. The worrisome field redefinition is

$$T_i = \langle T_i \rangle + \delta T_i \, , \tag{5.4}$$

where $\langle T_i \rangle$ is the expectation value of the moduli at the minimum of the potential (if it exists), including possible shifts due to supersymmetry breaking, and

$$\delta T_i = g_i(\delta R_j) + i a_i + \frac{1}{2} Q_i^\dagger Q_i \, , \tag{5.5}$$

where g_i is a polynomial function of the geometric moduli, as for example in (3.12) and (3.39). Since moduli stabilization and a vanishing cosmological constant is assumed, the action for T_i begins at $\mathcal{O}((T_i - \langle T_i \rangle)^2)$. But a contribution to the scalar mass from the field redefinition (5.5) could only come from $(\partial V_0 / \partial T_j)(T_j - \langle T_j \rangle)$ which vanishes. Note

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that it is the chiral component T_i which is stabilized, instead of the geometric variable R_i. This is the only assumption consistent with $N = 1$ $d = 4$ supersymmetry. Thus to compute the soft mass with a stabilizing potential for the moduli one may simply freeze the moduli at their expectation values. Using the usual supergravity potential, the Kähler potential in (5.1), the superpotential in (5.2), and assuming $\langle \Sigma_i \rangle$ is small compared with the four-dimensional Planck scale, then with stabilized moduli the scalar masses are given by

$$m_i^2 = m_{3/2}^2 \left(1 - \frac{2 \operatorname{Re} T_i}{p_i^2} \frac{|F_{\Sigma_i}|^2}{|W|^2} - \frac{4 (\operatorname{Re} T_i)^2}{p_i^2} \frac{|F_{T_i}|^2}{|W|^2} \right) , \tag{5.6}$$

where the gravitino mass in Einstein frame is $m_{3/2}^2 = e^K |W|^2$, and where here the unconventional notation $F_I \equiv \partial_I W + (\partial_I K) W$ is employed which differs from the standard notation by a factor of $e^{K/2} K^{i\bar{j}}$. Note that with vanishing auxillary component for the moduli (5.6) agrees with the formulae found in section 2.1 and appendix B where for the Kähler potential (5.1) the moduli were simply replaced by their expectation values. But before concluding that the masses are non vanishing when the moduli are stabilized, it is necessary to check that (5.6) is not identically zero once the cosmological constant is canceled. The condition for the cosmological constant to vanish is found to be

$$0 = \sum_i \left(\frac{2 \operatorname{Re} T_i}{p_i} |F_{\Sigma_i}|^2 + \frac{4 (\operatorname{Re} T_i)^2}{p_i} |F_{T_i}|^2 \right) - 3 |W|^2 \tag{5.7}$$

and this is not in general equivalent to (5.6). However, there are some special situations in which (5.7) and (5.6) equivalent, some of which have been detailed in section 2.1 and appendix B. By including an auxillary component for the diagonal moduli the results presented here slightly extend the discussion found in these other sections. In the five-dimensional supergravity brane world model there is only the single radion modulus, T, for which $p = 3$. In this case, the inherited no-scale Kähler potential without warping gives vanishing visible sector scalar masses even for $F_T \neq 0$. For three moduli T_i, with $i = 1, 2, 3$, it is seen by inspection that the scalar masses continue to vanish if both the stabilization and supersymmetry breaking preserve an S_3 symmetry. For generic patterns of moduli values and supersymmetry breaking on the distant brane, however, the equation for the scalar mass does not vanish when the cosmological constant vanishes.

Returning to the issue of moduli stabilization, we note that the expression for the soft masses above are dramatically modified if the moduli do not have any superpotential interactions. In [8] it was demonstrated that for the sequestered Kähler potential the absence of superpotential interactions for the radion implies that the conformal compensator always vanishes. It is not too difficult to extend their results to include the more general no-scale Kähler potentials (5.1). To see this, note that without stabilizing the moduli there are additional contributions to the soft masses from the term $\partial V_0 / \partial T_j (T_j - \langle T_j \rangle)$. Operationally it is then more convenient to substitute the field redefinition (5.5) into the scalar potential. Now for instance, the inverse metric for the hidden sector fields is independent of $Q_i^\dagger Q_i$. Visible scalar masses only come from terms involving $F_{T_i} = -p_i W_0 / R_i$ where $R_i = 2 \operatorname{Re} \langle T_i \rangle$

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with $\langle T_i \rangle$ defined in (5.5), and terms involving $F_{Q_i} = p_i |W| Q^* / R_i$. In this case with vanishing superpotential interactions for the moduli and the no-scale Kähler potentials (5.1) the scalar potential for the visible sector fields to quadratic order is proportional to

$$\frac{1}{p_i} \begin{pmatrix} F_{T_i} & F_{Q_i} \end{pmatrix} \begin{pmatrix} R_i^2 + Q_i^\dagger Q_i R_i & Q_i^\dagger R_i \\ Q_i R_i & R_i \end{pmatrix} \begin{pmatrix} F_{T_i}^\dagger \\ F_{Q_i}^\dagger \end{pmatrix}, \tag{5.8}$$

where there is no sum on repeated indices. Using the auxillary components given above it is readily confirmed that the soft scalar masses vanish. This demonstrates that without superpotential interactions or additional Kähler corrections for the moduli, general non-vanishing auxillary components do not lead to soft scalar masses. But as argued in the previous paragraph, these conclusions do not apply when there are superpotential interactions that stabilize the moduli.

Finally, one might wonder why (5.6) does not contain any contributions suppressed by the radion masses. The moduli couple to both hidden and visible sectors, and in integrating out these states one might expect corrections to the Kähler potential suppressed by these masses. But the operators that are generated are not of the form that would lead to tree-level masses after supersymmetry breaking. This was demonstrated in a model with a single volume modulus [8]. But here we see more generally that integrating out any number of moduli in the low energy theory does not affect the form of the tree-level visible sector masses. Instead the operators which are generated include a number of supersymmetric covariant derivatives which, with hidden sector supersymmetry breaking, do not contribute directly to scalar masses. This form for the effective operators which are generated by integrating out moduli is readily confirmed by either a component or supergraph calculation.

6. Gaugino masses

Gaugino masses require the breaking of both supersymmetry and R-symmetry. Both these symmetries are broken by a gauge kinetic function auxiliary expectation value

$$\int d^2\theta \mathcal{F}\left(S, T, \frac{\mu}{(\Phi M_{\rm reg})}\right) W^\alpha W_\alpha\,, \tag{6.1}$$

where μ is the infrared renormalization scale. The gauge kinetic function $\mathcal{F}(S, T, \mu/(\Phi M_{\rm reg}))$ depends on the dilaton S, the various T moduli, and at the loop-level on the conformal compensator superfield, Φ, through the regulator mass scale $M_{\rm reg}$. If hidden sector brane fields Σ are charged under global or gauge symmetries then the visible sector gauge kinetic function is at least bilinear in these fields. If the hidden sector scalar expectation values are small compared with the fundamental Planck scale, any hidden sector auxiliary expectation values give rise directly to visible sector gaugino masses which are suppressed by powers of hidden sector scalar expectation values over the fundamental scale, and therefore unimportant. In this case, only the auxiliary components for the dilaton, S, moduli, T, or conformal compensator, Φ, can give important direct contributions to the visible sector gaugino masses.

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The ratio of gaugino to scalar masses depends crucially on the origin of the auxiliary component within the gauge kinetic function. One-loop anomaly mediated contributions to the gaugino masses arise from the conformal compensator auxiliary expectation value, F_Φ, which, as discussed in section 2, can be induced indirectly by hidden sector supersymmetry breaking. If supersymmetry breaking is isolated in matter fields on the hidden sector brane, then as argued above, tree-level scalar masses in general arise from interactions between the branes, except in very special circumstances. In the scenario with supersymmetry breaking isolated on the hidden sector brane the gaugino masses are then a loop factor smaller than the scalar masses. Although model dependent, avoiding experimental bounds on gaugino masses of $\mathcal{O}(50-100\,\mathrm{GeV})$ from direct searches then generally requires some tuning to obtain radiative electroweak symmetry breaking. This occurs because the scalar masses, including that of the scalar Higgs, are somewhat larger than the electroweak scale as implied by the large ratio between scalar and gaugino masses. This is also the case for moduli dominated supersymmetry breaking, $F_T \neq 0$, which gives rise to tree-level scalar masses and one-loop gaugino masses through threshold effects. Dilaton dominated supersymmetry breaking, $F_S \neq 0$, however, gives rise to tree-level visible sector masses for both gauginos and scalars. From the criterion of avoiding tuning of electroweak symmetry breaking by obtaining both scalar and gaugino masses at the same order, dilaton dominated supersymmetry breaking, or a combination of dilaton, moduli, and hidden sector breaking, then seems most natural.

7. Gaugino mediation

Gaugino mediation is a variant of BWSB in which Standard Model gauge multiplets reside in (at least a subspace of) the bulk, while standard model quarks and leptons are confined to a visible sector brane [32, 33]. Supersymmetry breaking takes place on a hidden sector brane which is physically separated from the visible brane in a compact manifold. Since the gauge field multiplets reside in the bulk of the compact manifold and are in direct physical contact with the hidden sector supersymmetry breaking brane, the gauginos feel supersymmetry breaking directly at tree level and obtain masses at the compactification scale of order [32]

$$m_{1/2}^2 \sim \frac{F}{V_g} \sim \frac{V^{1/2}}{V_g} m_{3/2}^2 \,, \tag{7.1}$$

where V_g is the volume of the subspace within the compact manifold in which the gauge field multiplets reside, F is the supersymmetry breaking auxiliary expectation value on the hidden sector brane, and from (1.2) the gravitino mass is $m_{3/2} \sim F/V^{1/2}$. Note that for $V_g \sim V$ the gaugino masses are parameterically smaller than the gravitino mass. If the Higgs multiplets also reside in the bulk and are in contact with hidden sector brane, Higgs sector mass squared parameters receive direct tree level contributions at the compactification scale of order

$$m_h^2 \sim \frac{F^2}{V_h} \sim \frac{V}{V_h} m_{3/2}^2 \,, \tag{7.2}$$

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where V_h is the volume of the subspace within the compact manifold in which the Higgs multiplets reside. The scalar squarks and sleptons confined on the visible sector brane do not couple directly to the hidden sector brane and therefore have been argued to receive masses squared only radiatively from bulk gauge multiplets at one loop [32, 33]. At the compactification scale these contributions to scalar masses are

$$m^2 \sim \frac{g_4^2 m_{1/2}^2}{16\pi^2} \sim \frac{V}{V_g^2} \frac{g_4^2 m_{3/2}^2}{16\pi^2}\,, \tag{7.3}$$

where g_4 are the four-dimensional gauge couplings. In order to obtain the observed values of the four-dimensional gauge couplings the compactification volume should be not much larger than the fundamental scale. For one extra dimension a compactification radius of order the inverse GUT scale is generally assumed.

The gaugino mediated contributions (7.3) to the scalar squark and slepton radiative masses at the compactification scale are suppressed by a loop factor relative to the gaugino masses. This leads to the expectation that, from a low energy point of view, gaugino mediation amounts effectively to no-scale boundary conditions in which the squark and slepton masses nearly vanish at the messenger scale while gaugino masses and perhaps Higgs sector parameters are non vanishing [32, 33]. If this is the case, the dominant contribution to the squark and slepton masses squared comes from renormalization group evolution between the compactification and electroweak scales. This yields $m^2 \sim m_{1/2}^2$ at the electroweak scale. Theses contributions do not violate flavor.

However, as discussed above, with brane world supersymmetry breaking scalar masses squared generically arise at tree level from volume suppressed couplings between the visible and hidden sector branes,

$$m^2 \sim \frac{F^2}{V} \sim m_{3/2}^2\,. \tag{7.4}$$

A specific realization of gaugino mediation has not yet been presented in a string or M-theory background. String or M-theory models in which the gauge groups arise from ADE singularities on a sub-space with visible and hidden sector matter fields localized at physically separated singularities within this sub-space might give a realization of gaugino mediation within a consistent theory which contains gravity. The expectation that tree-level scalar masses (7.4) generically arise should hold in such models. These contributions are not suppressed compared with the gravitino mass and parameterically dominate over the gaugino mediated contributions (7.3) by a loop factor and at least one power of the volume. In addition, as discussed above these contributions are generally not universal. So gaugino mediation through the bulk alone is not sufficient to solve the supersymmetric sflavor problem. Additional assumptions are required to solve the sflavor problem such as flavor symmetries. Dominance of the gaugino mediated contributions to scalar masses might might arise if one had a model such as a flat five-dimensional realization of string or M-theory with very small or vanishing tensions (so as to avoid warping of the internal geometry). Dominance of gaugino mediation would also occur in models where the volume of the subspace in which the gauge multiplets propagate is much smaller than the volume of the gravitating space.

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8. The supersymmetric flavor problem

Virtual squark and slepton processes can in principle generate low energy quark and lepton flavor violating processes at levels much larger than allowed by current experimental bounds. This seems to imply that sufficent structure exists in the squark and slepton soft mass matrices to suppress these potentially dangerous supersymmetric contributions to flavor violation. Within any senario for transmitting supersymmetry breaking to the squark and slepton fields it is then worth addressing what assumptions, if any, are required to avoid excessive sflavor violation.

Within a given supersymmetry breaking mediation scenario, if squarks and sleptons with the same gauge quantum numbers acquire non-universal masses at the messenger scale, then there is potentially a sflavor problem. This is because the squark and slepton eigenvectors associated with the non-universal masses define directions in flavor space. If these are not aligned or proportional to the eigenvectors defined by the quark and lepton masses to a sufficient degree dangerous sflavor violation can arise. Without additional assumptions about underlying flavor symmetries this is of course not in general quaranteed.

In the string and M-theory BWSB models presented here, such as the orbifold examples, fields with the same gauge quantum numbers do generally acquire non-universal masses. So without additional assumptions about flavor symmetries in the underlying theory, BWSB scenarios alone do not appear to be free of the supersymmetric flavor problem. However, it is worth noting that, as discussed in section 2.3, if a direct compactification of string/M-theory to pure five-dimensional supergravity with end of the world branes exists, it would presumably inherit at lowest order the no-scale form of the Kähler potential for the untwisted states from the underlying $N = 4$ Khaler potential, and give vanishing tree-level masses. This should probably be considered an interesting property of a given model, rather than a general feature of BWSB.

In most of the models presented here multiple copies of chiral matter arise as the result of discrete or continuous internal geometric symmetries. Copies of visible sector chiral matter with the same gauge quantum numbers are by definition different flavor generations. So in these models flavor is closely related to internal geometric symmetries. The appearance of soft masses which do not commute with flavor can be traced to the existence of fields in the bulk which in the underlying theory necessarily also transform under these geometric symmetries.

From the low energy effective field theory point of view it might seem that including bulk fields which transform under flavor is the only way in which to obtain non-universal masses. However, as the models presented here also illustrate, the soft masses depend not only on which bulk fields are present, but also on the form of the bulk-brane matter couplings which are suppressed by the fundamental scale of theory. From an effective field theory point of view these couplings are restricted only by four-dimensional $N = 1$ supersymmetry (and possibly anomaly cancellation). So even with the minimal supergravity multiplet in the bulk, an effective field theory analysis of BWSB requires additional assumptions about flavor symmetries at the fundamental scale in order to avoid dangerous sflavor violation in general.

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9. Conclusions

Brane worlds provide an interesting scenario in which to realize hidden sector supersymmetry breaking. As discussed here, even though supersymmetry breaking can be isolated on a hidden sector brane which is not in direct physical contact with the visible sector brane, unsuppressed tree-level squark and slepton masses are generally obtained. This is illustrated in a number of models. In those models with more than one generation the masses are generically non universal. In addition, these models typically contain tachyons. These can be removed by a projection, but this requires a correlation between the projection and the pattern of hidden sector supersymmetry breaking.

The origin of the non-derivative interactions which couple the branes and gives rise to tree-level visible sector masses is easy to understand in the underlying theory as arising from exchange of bulk supergravity fields between the branes. It is important to note that for total space-time dimensions greater than five, bulk fields in the minimal supergravity multiplet are sufficient to give rise to non-derivative unsuppressed brane-brane couplings even in flat backgrounds. Higher-dimensional supersymmetry guarantees the existence of these bulk states which therefore can not simply be ignored by fiat. In addition, corrections to the leading tree-level masses determined by the form of the inherited Kähler potential are not suppressed by additional powers of the compactification volume as has been claimed and can be significant. This is consistent with the expectation that the inherited Kähler potential is not protected in any way in the low energy theory.

Since the tree-level soft masses are not generally degenerate, additional assumptions about flavor symmetries and their breaking are required to ensure proportionality or alignment in order to avoid dangerous sflavor changing effects. So contrary to previous expectations, BWSB alone does not give a solution to the supersymmetric flavor problem. In this sense the compactification model building and phenomenology of BWSB is similar to standard hidden sector supersymmetry breaking. Since tree-level scalar masses are generic, brane world realizations do not in general provide a robust rationale for anomaly mediated supersymmetry breaking. It may be possible, however, to construct specific models in which tree-level masses vanish and anomaly mediation is important. The most natural setting for this would seem to be compactifications of a consistent microscopic theory to pure five-dimensional supergravity with end of the world branes, although none are known at present. Another possibility is a Hořava-Witten compactification of M-theory on a Calabi-Yau fibration with $h^{1,1} = 1$. The lowest order inherited Kähler potential for these cases is the sequestered no-scale form which gives vanishing tree-level soft masses. For more general compactifications in which the inherited Kähler potential is the sum of logarithms form, unbroken flavor symmetries in the hidden sector can also give vanishing tree-level masses; but this of course requires additional assumptions about flavor symmetries. In either case, corrections to the inherited Kähler potential, which generally give rise to tree-level soft masses, would have to be tuned to be small by moving to some corner of moduli space and/or ensuring that the brane tensions vanish, presumably by some discrete symmetry. On top of this, additional interactions have to be assumed which lift the tachyonic right-handed slepton of anomaly mediation (although these might be provided by the

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corrections to the inherited Kähler potential — however in this case anomaly mediation does not give the dominant contribution to scalar masses). All of these model assumptions required to achieve anomaly mediation could also likely be achieved in a standard hidden sector supersymmetry scenario without reference to a brane world picture.

Scenarios for transmitting supersymmetry breaking to the visible sector seem most natural if the scalars and gauginos receive masses which are of the same order without any tuning of parameters. In the brane world realization tree-level masses for scalars are naturally obtained. Tree-level masses for gauginos require an auxiliary component for the dilaton. So the most natural BWSB scenario seems to be either dilaton domination, or one in which the dilaton acquires an auxiliary expectation value comparable to those in the hidden sector. Dilaton domination gives, at leading order, universal tree-level contributions to visible sector scalar masses. However, even in this case which is flavor blind at lowest order, as we have seen, flavor violating corrections to the dilaton Kähler potential are not necessarily small, in particular in the Hořava-Witten theory with standard gauge coupling unification. So additional assumptions about flavor symmetries seem to be required. This also has implications for the standard dilaton domination scenario in the language of weakly coupled string theory. Phenomenological analyses [34] have assumed that scalar degeneracy and therefore alignment is not likely to be better than $\mathcal{O}(\alpha_{\rm GUT})$. This is just barely large enough to explain the suppression of flavor violating processes in the kaon system. However, in the strongly coupled Hořava-Witten limit of heterotic string theory the non-universal flavor violating couplings of the dilaton arise classically at $\mathcal{O}(T/S)$. On both phenomenological and theoretical grounds, one expects that T/S is not much smaller than unity. So even though these corrections are non perturbative from the heterotic string point of view, they are not likely to be small numerically. And therefore again, additional assumptions about flavor symmetries seem to be required.

In a top-down approach to models of nature, one could hope that large classes of models or vacua within an underlying theory might have generic features which can eventually be confronted experimentally at low energies. If so, then these classes, and perhaps ultimately the underlying theory, can be considered to be predictive and testable. Even though there are many problems that string/M-theory can not address with our present level of understanding, certain questions regarding the mechanism of supersymmetry breaking and resulting patterns of squark and sleptons masses can be addressed with present technology. One might have hoped that BWSB vacua within string/M-theory might have provided an example of a class of models with universal generic features by predicting universal flavor conserving squark and slepton masses as the result of physically separating the hidden and visible sectors within an internal manifold. Unfortunately, for a rather large class of models this does not appear to be the case. The superpartner spectrum and magnitude of sflavor violation appears to be very model dependent. So just as for standard perturbative string theory with hidden sector supersymmetry breaking, one is reduced to presenting predictions (in principle) for specific models rather than for broad classes. The latter would clearly be preferable. Perhaps future scenarios for realizing supersymmetry breaking within string/M-theory will provide general patterns or features which are predictive and testable.

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A. M-theory orbifold compactifications and their inherited Kähler potentials

On possibility for obtaining brane world models with $N = 1$ supersymmetry from M-theory is to consider Hořava-Witten orbifold backgrounds $S^1/\mathbb{Z}_2 \times \mathcal{M}$ where $\mathcal{M}$ is a six-dimensional $N = 1$ preserving compact orbifold of $\mathbb{T}^6$. If the $S^1/\mathbb{Z}_2$ interval length is large, $R_{11} \gtrsim \ell_{11}$, a brane world background with end of the world branes results. In the weakly coupled limit, $R_{11} \ll \ell_{11}$, modular invariance of one-loop string amplitudes in the resulting perturbative heterotic string theory gives certain consistency conditions on $\mathcal{M}$ discussed below, and also in general requires the existence of additional twisted states which reside at fixed points of $\mathcal{M}$. In the absence of an underlying theory of M-theory orbifolds, we content ourselves with orbifold compactifications of this type which are modular invariant in the weakly coupled heterotic limit. We will assume that some of the perturbative heterotic models exist at strong coupling with a geometric description. Then for these models the lowest order Kähler potential for the untwisted states is inherited from the $N = 4$ Kähler potential obtained from compactification on $\mathbb{T}^6$. As discussed in section 3.1 a low energy supergravity analysis implies that this lowest order inherited form for the untwisted states should receive small corrections in the brane world limit as long as the compact volume is large in eleven-dimensional Planck units. Our discussion below applies to these classes of models, again assuming they are consistent compactifications of M-theory.

In an orbifold construction the states which survive in the low energy theory are invariant under the orbifold action. In general this action is non trivial in both compact geometric directions as well as in the gauge group of the underlying theory. In addition, twisted states which reside at orbifold fixed points also appear in the low energy theory. For M-theory orbifold backgrounds $S^1/\mathbb{Z}_2 \times \mathcal{M}$, a subset of the $E_8 \times E_8'$ M-theory twisted gauge supermultiplets which reside on the end of the world visible and hidden sector branes survive in the low energy theory, $Q_i \subset \mathbf{248}_i \in E_8$ and $\Sigma_i \subset \mathbf{248}'_i \in E_8'$, respectively, where $i = 1, 2, 3$ labels the internal complex coordinates of $\mathcal{M}$. From the weakly coupled heterotic string point of view these fields are untwisted states, and will be referred to as such below. The lowest order tree-level Kähler potential for these states is inherited directly from the $N = 4$ Kähler potential (2.6) by simply removing non-invariant states. There are in general additional twisted states in the low energy theory which reside at fixed points of $\mathcal{M}$. The dependence of the Kähler potential on these twisted states is not restricted by extended symmetries since $\mathcal{M}$ preserves only $N = 1$ supersymmetry. Here we focus only on the

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untwisted Q_i and Σ_i visible and hidden sector fields which do inherit a lowest order Kähler potential from the ten-dimensional theory.

For simplicity we consider orbifold backgrounds $\mathcal{M}$ which are symmetric abelian orbifolds of $\mathbb{T}^6$. An abelian orbifold group Γ acts on the three complex planes of $\mathbb{T}^6$ by

$$z_i \longrightarrow e^{2\pi i r_i} z_i\,. \tag{A.1}$$

Modular invariance of one-loop string amplitudes requires that [14]

$$|\Gamma| \sum_i r_i = 0 \bmod 2\,, \tag{A.2}$$

where $|\Gamma|$ is the order of the orbifold group. The existence of an unbroken $N = 1$ supersymmetry requires that

$$\sum_i r_i = 0 \tag{A.3}$$

for which the modular invariance condition (A.2) is then automatically satisfied.

To describe the action of Γ in the gauge groups in the present case, it is convenient to consider the subgroup $\mathrm{SU}(8) \subset \mathrm{SO}(16) \subset E_8$ and likewise for E_8'. The actions on $\lambda \in \mathrm{SU}(8)$ and $\lambda' \in \mathrm{SU}(8)'$ are

$$\begin{aligned} \lambda_a &\longrightarrow e^{2\pi i \beta_a} \lambda_a\,, \\ \lambda'_a &\longrightarrow e^{2\pi i \beta'_a} \lambda'_a\,, \end{aligned} \tag{A.4}$$

where $a = 1, \ldots, 8$. In the weak coupling limit, modular invariance requires that [14]

$$|\Gamma| \sum_a \beta_a = |\Gamma| \sum_a \beta'_a = 0 \bmod 2 \tag{A.5}$$

and

$$|\Gamma| \left[\sum_a \left(\beta_a^2 + \beta_a'^2 \right) - \sum_i r_i^2 \right] = \begin{cases} 0 \bmod 2\,, & |\Gamma| \text{even}\,, \\ 0 \bmod 1\,, & |\Gamma| \text{odd}\,. \end{cases} \tag{A.6}$$

The orbifold action (A.4) breaks the gauge groups to a subgroup, but does not reduce the rank for abelian Γ.

One of the simplest orbifold constructions begins with three $\mathbb{T}^2$ tori, each of which preserves a $\mathbb{Z}_3$ symmetry $z_i \to e^{2\pi i/3} z_i$. A $\mathbb{Z}_3$ orbifold twist consistent with this symmetry, modular invariance, and preserving $N = 1$ supersymmetry is $r_i = (1, 1, -2)/3$. This orbifold leaves invariant an S_3 geometric symmetry since the twist on each plane is actually identical in this case, $z_i \to e^{2\pi i/3} z_i$. The so-called standard embedding involves also a gauge twist by an element of $\mathrm{SU}(3) \subset E_8$ given by $\beta_a = r_i \delta_{ia}$, and $\beta'_a = 0$. This leaves an unbroken subgroup $\mathrm{SU}(3) \times E_6 \times E_8'$. The visible sector untwisted matter fields, Q_i, which are invariant under the orbifold projection transform as $(\mathbf{3}, \mathbf{27})_i \in \mathrm{SU}(3) \times E_6$ for $i = 1, 2, 3$. The untwisted sector then has 3 generations of $(\mathbf{3}, \mathbf{27}) \in \mathrm{SU}(3) \times E_6$. There are additional twisted states which cancel gauge anomalies. For this orbifold there are no hidden sector matter fields since all the Σ_i are non invariant. So this particular orbifold is not a useful model for hidden supersymmetry breaking. However, because of the S_3

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symmetry it does provide a simple example in which there are states in different sectors $i \neq j$ with the same gauge quantum numbers. So off-diagonal combinations of fields, $Q_i Q_{\bar{j}}^{\dagger}$, are gauge invariant and can appear in the Kähler potential. The lowest order tree-level Kähler potential inherited from (2.6) for the untwisted states of this orbifold is then [14]

$$K = -\ln\det\left(T_{i\bar{j}} + T_{i\bar{j}}^{\dagger} - \mathrm{tr}\, Q_i Q_{\bar{j}}^{\dagger}\right) - \ln\left(S + S^{\dagger}\right), \tag{A.7}$$

where the trace is over the gauge quantum numbers. Note that even though the orbifold only leaves invariant an $S_3 \times \mathrm{U}(1)_R \subset \mathrm{SU}(4)$ R-symmetry, the inherited Kähler potential possess an accidental $\mathrm{SU}(3) \times \mathrm{U}(1)_R$ global symmetry since it depends only on two derivative terms in the ten-dimensional theory. Higher order corrections to the full Kähler potential in the four-dimensional theory would of course break the accidental continuous flavor symmetry to the discrete $S_3 \subset \mathrm{SU}(3)$.

In the perturbative string limit this orbifold may be deformed to a smooth Calabi-Yau manifold by turning on blow up modes which resolve the fixed points, giving a Calabi-Yau with $h^{1,1} = 36$ and $h^{2,1} = 0$. The blow up modes break the SU(3). There are in total 36 generations of $\mathbf{27} \in E_6$, of which 27 come from the twisted sector. It would be useful to understand the M-theory geometric lift of this compactification.

A variant of the above $\mathbb{Z}_3$ orbifold that does possess hidden sector matter involves the same spacetime and visible sector gauge twists, $r_i = (1, 1, -2)/3$ and $\beta_a = r_i \delta_{ia}$, but with the hidden sector gauge twist $\beta'_a = (2, 2, 2, 0^5)/3$, where the exponent indicates the multiplicity of the component. This is modular invariant, preserves $N = 1$ supersymmetry, and has an S_3 symmetry. This choice of gauge twists is not unique, since the same spacetime twist can be accompanied by other modular invariant gauge twists in both the hidden and visible sectors. Orbifolds based upon this $\mathbb{Z}_3$ spacetime twist thus represent a class of models. The form of the Kähler potential for the untwisted states will be common to this class, and all of these models will have in the untwisted sector 3 (or 0) generations because of the S^3 symmetry. For the choice of gauge twists above the visible sector matter content is identical to the previous example and in particular the untwisted sector has 3 generations. The hidden sector unbroken gauge group with this gauge twist is $E'_6 \times \mathrm{SU}(3)' \subset E'_8$. And in this case there is hidden sector matter $(\mathbf{3}, \mathbf{27})_i \in E'_6 \times \mathrm{SU}(3)' \subset E'_8$ for $i = 1, 2, 3$. Because of the S_3 symmetry, off-diagonal combinations of both visible and hidden sector fields, $Q_i Q_{\bar{j}}^{\dagger}$ and $\Sigma_i \Sigma_{\bar{j}}^{\dagger}$, are gauge invariant and can appear in the Kähler potential. The lowest order tree-level inherited Kähler potential or supergravity f function for the untwisted states of this class of $\mathbb{Z}_3$ orbifolds is then just given by (2.6) or (2.7), respectively,

$$K = -\ln\det\left(T_{\bar{i}j} + T_{\bar{i}j}^{\dagger} - \mathrm{tr}\, Q_{\bar{i}}^{\dagger} Q_j - \mathrm{tr}\, \Sigma_{\bar{i}}^{\dagger} \Sigma_j\right) - \ln\left(S + S^{\dagger}\right). \tag{A.8}$$

In this example each sector i has only one chiral matter field appearing in the visible and hidden untwisted sectors. For more generic choices of the gauge twists, however, there will appear in each sector i states with different gauge quantum numbers according to which matter visible sector states $Q_i \subset \mathbf{248}_i \in E_8$ and hidden sector states $\Sigma_i \subset \mathbf{248}'_i \in E'_8$ survive the orbifold projection. In these more generic examples the trace appearing inside the determinant in the inherited Kähler potential includes a sum over these states. But we

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remphasize that the form of the Kähler potential and the existence of 3 (or 0) generations is determined by the spacetime twist, with the choice of gauge twist determining the representation content.

The Kähler potential (A.8) is not of the sequestered form and has non-derivative interactions between the visible and hidden sector branes. It gives rise to tree-level masses from hidden sector supersymmetry breaking as discussed in appendix B. However, the lowest order visible sector scalar masses arising from hidden sector auxiliary expectation values with classes of orbifolds such as the ones given above which preserve an S_3 geometric symmetry inherit the $\mathrm{Tr}\, m^2 = 0$ sum rule from the $N = 4$ Kähler potential discussed in appendix B. The associated visible sector tree-level tachyon(s) imply that such S_3 symmetric orbifolds can not be phenomenologically viable in the absence of large corrections which would lift the tree-level tachyon(s). Even though this particular class of orbifolds are not realistic for the reasons given above, they do illustrate that the sequestered intuition is not generally valid for brane world models which preserve $N = 1$ supersymmetry.

It is also possible to consider orbifolds that reduce the form of the inherited Kähler potential to a sum of logarithms rather than a logarithm of a determinant. This is illustrated in a class of $\mathbb{Z}_6$ orbifolds with the space-time twist $r_i = (1, 2, -3)/6$. This is modular invariant, preserves $N = 1$ supersymmetry, but is not invariant under any discrete symmetry which interchanges the three planes. Because there is no geometric symmetry between the planes, only the diagonal moduli, $T_i \equiv T_{i\bar{i}}$ survive the orbifold projection. For the gauge twists, consider the case $\beta_a = (3^2, 2, 1^2, 0^3)/6$ and $\beta'_a = (3^2, 2, 1^4, 0)/6$, where again the exponents indicate the multiplicity of the component. These twists satisfy the modular invariant conditions. This choice of gauge twists is not unique, since the same spacetime twist can be accompanied by other modular invariant gauge twists in both the hidden and visible sectors. Orbifolds based upon this $\mathbb{Z}_6$ spacetime twist thus represent a class of models. With the gauge twists given above the unbroken gauge groups are $\mathrm{SU}(4) \times \mathrm{SU}(2) \times \mathrm{SO}(4) \times \mathrm{U}(1)^2 \subset E_8$ and $\mathrm{SU}(4)' \times \mathrm{SO}(4)' \times \mathrm{U}(1)'^3 \subset E'_8$. With these twists there is matter in both the visible and hidden sector since some of the Q_i and Σ_i are invariant. However, because there is no discrete geometric symmetry, off-diagonal combinations of fields $Q_i Q_j^\dagger$ and $\Sigma_i \Sigma_i^\dagger$ are not guaranteed to be gauge invariant since states in different $i \neq j$ sectors generally have different gauge quantum numbers. In fact no gauge invariant off-diagonal terms exist if a given representation under the unbroken subgroup arises only once. This is always the case if the representations in the low energy theory are obtained by projecting an adjoint representation of a larger group in the higher-dimensional theory. One may explicitly verify this, although it is important to keep track of all the U(1) charges to verify this. In this model the untwisted sector has many different gauge quantum numbers but only one generation of each.

The lowest order tree-level Kähler potential for the untwisted states of the class of $\mathbb{Z}_6$ orbifolds given above without gauge invariant off-diagonal combinations of fields inherited from (2.6) is then

$$K = -\sum_{i=1}^{3} \ln\left(T_i + T_i^\dagger - \mathrm{tr}\, Q_i^\dagger Q_i - \mathrm{tr}\, \Sigma_i^\dagger \Sigma_i\right) - \ln\left(S + S^\dagger\right), \tag{A.9}$$

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where the traces are over both gauge quantum numbers and are different for each i. This Kähler potential is invariant only under $\mathrm{U}(1)_R$ since the visible and hidden sector representations and multiplicities are distinct for each i. This Kähler potential is not of the sequestered form. The non-derivative couplings between the visible and hidden sector branes implied by (A.9) give rise to unsuppressed tree-level scalar masses as discussed in the appendix B.

The previous class of examples has many untwisted states, but due to the spacetime twist there is only one generation in the untwisted sector for each gauge quantum number. Next we consider a class of orbifolds that is a hybrid of the previous two examples such that in the untwisted sector there are two generations. This is illustrated in a $\mathbb{Z}_6$ orbifold with spacetime twist $r_i = (1,1,-2)/6$, which is modular invariant and preserves $N=1$. In addition we consider gauge twists in the hidden and visible sectors chosen to be modular invariant and provide untwisted hidden and visible sector chiral matter. Since the spacetime twist preserves a S_2 permutation symmetry, states in the untwisted sector with Q_i for $i=1,2$ are guaranteed to have the same gauge quantum numbers. This class of models then has two generations in the untwisted sector. Consequently, off-diagonal contributions $Q_i^\dagger Q_{\bar{j}}$ for $i,j=1,2$ are guaranteed to be gauge invariant. The lowest order tree-level Kähler potential for the untwisted states inherited from (2.6) is then

$$K = -\ln \det_{i=1,2}\left(T_{i\bar{j}} + T_{i\bar{j}}^\dagger - \mathrm{tr}\, Q_i^\dagger Q_{\bar{j}} - \mathrm{tr}\, \Sigma_i^\dagger \Sigma_{\bar{j}}\right) - \ln\left(T_{3\bar{3}} + T_{3\bar{3}}^\dagger - \mathrm{tr}\, Q_3^\dagger Q_{\bar{3}} - \mathrm{tr}\, \Sigma_3^\dagger \Sigma_3\right). \tag{A.10}$$

This is also not of the sequestered form. The soft mass spectrum for the two generations with hidden sector supersymmetry breaking is generically non universal, as discussed in appendix B.

B. Soft scalar masses from inherited Kähler potentials

The spectrum of visible sector soft scalar masses arising from hidden sector supersymmetry breaking depends on the four-dimensional Kähler potential. For any theory geometrically embedded in a higher-dimensional supersymmetric theory, the lowest order tree-level soft masses depend on the form of the Kähler potential inherited from the underlying theory which necessarily has extended supersymmetry. In many cases the lowest order inherited tree-level Kähler metric for the matter fields is either quaternionic if the underlying microscopic theory has 8 supercharges or flat if it has 16 supercharges. At lowest order the inherited Kähler potential for the untwisted states in an orbifolded theory is obtained directly from that of the underlying theory with extended supersymmetry by simply truncating to the light fields which survive in the four-dimensional theory, as discussed in appendix A.

The tree-level spectra of visible sector soft masses arising from the inherited Kähler potentials possess some special features as described below. In general, unacceptable visible sector tachyons arise with hidden sector supersymmetry breaking, but are absent in certain classes of compactifications. These classes however depend on the pattern of hidden sector supersymmetry breaking, and cannot be chosen a priori. In models with more

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than one generation in the untwisted sector the soft masses are generically non universal. For definiteness we focus on microscopic theories which have the matter content of a weakly coupled string theory in ten dimensions in a background which preserves at lowest order 16 supercharges. However, since the form of the inherited Kähler potential only depends on the existence of extended supersymmetries of the underlying background, the results are more generally applicable, to, for example, the untwisted states appearing in Hořava-Witten orbifold backgrounds or D-brane theories on an orbifold.

Consider first the case in which all the visible and hidden sector fields, geometric moduli, and dilaton survive in the four-dimensional theory. While this is not a realistic compactification since the surviving states form $N = 4$ multiplets, it is instructive in illustrating the origin of relations among the visible sector scalar spectrum in more realistic examples discussed below with fewer surviving states. The low-energy four-dimensional theory at the level of two-derivative terms has an SU(4) global symmetry which is inherited from the R-symmetry of the microscopic theory. The fermions in both the visible and hidden sector transform as $\mathbf{4} \in \mathrm{SU}(4)$, and the scalars as $\mathbf{6} \in \mathrm{SU}(4)$. In the $N = 1$ low energy description only an $\mathrm{SU}(3) \times \mathrm{U}(1)_R \subset \mathrm{SU}(4)$ is manifest. Under this subgroup the visible and hidden sector $N = 1$ chiral multiplets transform as complex $\mathbf{3}_+ \in \mathrm{SU}(3)\times\mathrm{U}(1)_R$, the $T_{i\bar{j}}$ moduli chiral multiplets transform as complex $\mathbf{8}_0 \oplus \mathbf{1}_0 \in \mathrm{SU}(3) \times \mathrm{U}(1)_R$, the T_{ij} moduli chiral multiplets transform as $\mathbf{6}\oplus\overline{\mathbf{3}}$, and the dilaton chiral multiplet is invariant. In the $N = 1$ description the SU(3) subgroup is a manifest global flavor symmetry of the low energy theory at the two-derivative level. The four-dimensional inherited Kähler potential, in this case, is fixed by the $N = 4$ supersymmetry to be [12, 13, 14]

$$K = -\ln\det\left(T_{i\bar{j}} + T^{\dagger}_{i\bar{j}} - \mathrm{tr}\, Q_i Q^{\dagger}_{\bar{j}} - \mathrm{tr}\, \Sigma_i \Sigma^{\dagger}_{\bar{j}}\right) - \ln\left(S^{\dagger} + S\right), \tag{B.1}$$

where the traces are over the visible and hidden sector gauge groups, respectively, i and $\bar{j}$ are global SU(3) flavor fundamental and anti-fundamental indices, respectively, and the dependence on the T_{ij} moduli is suppressed. Note that this is manifestly invariant under the $\mathrm{SU}(3) \times \mathrm{U}(1)_R$ global flavor symmetry.

In order to exhibit a simple expression for the visible sector scalar masses arising from the Kähler potential (B.1) consider the case in which only diagonal scalar moduli acquire an expectation value, $\langle T_{i\bar{i}}\rangle \neq 0$, while the off-diagonal moduli vanish, $\langle T_{i\bar{j}}\rangle = \langle T_{ij}\rangle = 0$. If all the $T_{i\bar{j}}$ moduli are stabilized by superpotential interactions and have vanishing auxiliary components, $\langle F_{T_{i\bar{j}}}\rangle = 0$, the visible sector scalar masses arising from hidden sector auxiliary components, $F_i \equiv \langle F_{\Sigma_i}\rangle$, are easily computed from the relevant terms in the $N = 1$ supergravity potential

$$V \supset \exp(K)\left(F_i K^{i\bar{j}} F_{\bar{j}} - 3|W|^2\right), \tag{B.2}$$

where $K_{i\bar{j}} \equiv \partial_i \partial_{\bar{j}} K$, $K^{i\bar{j}}$ is the inverse metric, and $F_i \equiv \partial_i W + K_i W$. The canonically normalized visible sector scalar mass squared matrix in the 3×3 SU(3) flavor space, in this case, is

$$m^2_{i\bar{j}} = \frac{m^2_{3/2}}{|W|^2}\left(|W|^2 \delta_{i\bar{j}} - 2\sqrt{\mathrm{Re}\, T_i\, \mathrm{Re}\, T_j} F_i F^*_{\bar{j}}\right) \tag{B.3}$$

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with no sum on repeated indices, and where $2\,\mathrm{Re}\,T_i \equiv \langle T_{i\bar{i}} + T^\dagger_{i\bar{i}} \rangle$, and $F_i F^*_{\bar{j}} \equiv \mathrm{tr}\,F_i F^*_{\bar{j}}$ for each $i\bar{j}$. In deriving this expression the cosmological constant is assumed to vanish by the relation

$$\sum_i 2\,\mathrm{Re}\,T_i |F_i^2| - 3|W|^2 = 0 \tag{B.4}$$

and the gravitino mass squared with this condition is

$$m^2_{3/2} \equiv \exp(K)|W|^2\,. \tag{B.5}$$

The eigenvalues of (B.3) are in general non vanishing and non degenerate. The trace of the mass squared matrix (B.3) vanishes,

$$\mathrm{Tr}\,m^2 = 0\,, \tag{B.6}$$

as the result of the vanishing cosmological constant condition (B.4) and SU(3) flavor invariance which implies that the Q_i multiplicities in the visible sector are identical for each i, and likewise for the Σ_i in the hidden sector. Note that the Tr is only over visible sector scalars. This of course implies that there is at least one visible sector tree-level tachyon for any non-vanishing hidden sector auxiliary component.

The form of the visible sector scalar spectrum arising from the Kähler potential (B.1) is illustrated in the simplest case where only one of the hidden sector fields has an auxiliary expectation value. The mass squared eigenvalues for $F_1 \neq 0$, $F_2 = F_3 = 0$, but still allowing for arbitrary $T_{i\bar{i}}$, are

$$m^2_{Q_i} = m^2_{3/2}(-2, 1, 1)\,. \tag{B.7}$$

This pattern of breaking preserves a SU(2) $\subset$ SU(3) flavor symmetry which enforces the two-fold degeneracy of the eigenvalues. Another simple case is the situation in which hidden sector fields acquire a SU(3) diagonal auxiliary component, $F_i = F$, and that the moduli have an SU(3) diagonal value, $T_{i\bar{i}} = T$ and are stabilized with vanishing auxillary components. Then from (B.3) and (B.4) the visible sector mass squared matrix, in this case, is

$$m^2_{i\bar{j}} = m^2_{3/2} \begin{pmatrix} 0 & -1 & -1 \\ -1 & 0 & -1 \\ -1 & -1 & 0 \end{pmatrix}. \tag{B.8}$$

This $S_3 \subset$ SU(3) symmetric matrix has eigenvectors which transform as $\mathbf{2} \oplus \mathbf{1} \in S_3$ with the same eigenvalues as (B.7). This may be understood by noting that the SU(3) flavor symmetry of the Kähler potential (B.1) may be used to rotate the three F_i auxiliary components into one component, F_1 say, which is the case previously discussed just above.

Consider next the case in which the only states which survive compactification in the low-energy four-dimensional theory are the diagonal geometric moduli, $T_i \equiv T_{i\bar{i}}$, the dilaton, S, and some of the visible and hidden sector fields, Q_i and Σ_i. In addition, restrict attention to cases in which the off-diagonal combinations of visible and hidden sector fields, $Q_i Q^\dagger_{\bar{j}}$ and $\Sigma_i \Sigma^\dagger_{\bar{j}}$, are not gauge invariant, and therefore do not appear in the Kähler potential. This situation arises in some of the orbifold compactifications discussed

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in section 2.1 and appendix A. The lowest order tree-level Kähler potential inherited from (B.1) for this truncated set of states is

$$K = -\sum_{i=1}^{3} \ln \left(T_i + T_{\bar{i}}^\dagger - \operatorname{tr} Q_i Q_{\bar{i}}^\dagger - \operatorname{tr} \Sigma_i \Sigma_{\bar{i}}^\dagger \right) - \ln \left(S^\dagger + S \right) , \tag{B.9}$$

where again the traces are over visible and hidden sector gauge groups, respectively. Note that for a generic compactification the visible and hidden sector gauge groups, matter representations, and multiplicities are not necessarily the same. So the Kähler potential is then only invariant under a $\mathrm{U}(1)_R \subset \mathrm{SU}(4)$. Allowing for arbitrary hidden sector auxiliary components $F_i \equiv F_{\Sigma_i}$, the canonically normalized visible sector scalar masses squared arising from (B.9) with vanishing cosmological constant and stabilized moduli with vanishing auxillary components are

$$m_i^2 = \frac{m_{3/2}^2}{|W|^2} \left(|W|^2 - 2 \operatorname{Re} T_i |F_i|^2 \right) \tag{B.10}$$

with no sum on the repeated index and where $|F_i|^2 \equiv \operatorname{tr} F_i^* F_i$ is the sum of the hidden sector auxiliary components squared for given i. The scalar masses (B.10) do not vanish in general. Note that all the visible sector fields Q_i with the same i are degenerate with mass squared m_i^2, but are not generically degenerate with the Q_j fields for $i \neq j$. In fact, in this case, there are just three possible eigenvalues, $m_{i=1,2,3}^2$. The multiplicities of each eigenvalue, however, will not be the same in general since the multiplicities of the visible sector fields Q_i are not necessarily the same for each i. Assuming there is at least one matter field Q_i for each i — that is, there are no empty sectors — a sum rule for these masses similar to that found in the $N = 4$ case discussed above in which all fields survive is obtained. In this case, the sum of the three eigenvalues (B.10) vanishes as the result of the vanishing cosmological constant condition (B.4)

$$m_1^2 + m_2^2 + m_3^2 = 0 \,. \tag{B.11}$$

However, since the multiplicities for each i are not necessarily the same, $\operatorname{Tr} m^2 \neq 0$ in general. The condition (B.11) implies that in this case there are also generically tree-level tachyons in the visible sector. Tachyons are avoided in compactifications in which all visible sector matter fields Q_i in the i-th sector(s) with negative mass squared eigenvalue(s) are projected out of the low-energy theory. Since tachyonic states occur in i-th sectors where Σ_i has an auxillary component $|F_{\Sigma_i}|^2 > |W|^2/(2 \operatorname{Re} T_i)$, avoiding tachyons requires a correlation between the orbifold projection and the direction of supersymmetry breaking. This could be achieved in an orbifold compactification where the projections in the visible sector and hidden sector are anti-correlated. That is, empty visible sectors i correspond to non-empty hidden sectors and vice versa.

A special case of the condition (B.11) resulting from the Kähler potential (B.9) is obtained for hidden sector fields which acquire diagonal auxiliary components, $F_i = F$, and scalar moduli which have an S_3 symmetric expectation value, $T_i = T$. In this case, the

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vanishing cosmological constant condition (B.4) reduces to $2\,\mathrm{Re}\,T|F|^2 - |W|^2 = 0$. This implies that all the scalar masses squared eigenvalues (B.10) vanish,

$$m_i^2 = 0\,. \tag{B.12}$$

This differs from the spectrum (B.7) obtained from the Kähler potential (B.1) with the same diagonal moduli and auxiliary components. The difference traces to the fact that unlike (B.1), the Kähler potential (B.9) does not preserve a SU(3) flavor symmetry. The hidden sector auxiliary expectation values can therefore not be rotated into a single flavor. For the case of diagonal hidden sector auxiliary components the mass eigenvalues (B.10) are necessarily identical. The condition (B.11) then in turn implies that the scalar masses in fact vanish in this case. The lowest order vanishing of the visible sector masses may also be understood as a remnant of the underlying $N = 4$ Kähler potential (B.1) since the diagonal elements of (B.8) vanish, and by assumption the off-diagonal elements have been projected out of the low-energy four-dimensional theory.

Consider next the case where the only states that survive in the low enery theory are moduli $T_{i\bar{j}}$ for $i = 1, 2$, $T_{3\bar{3}}$, and visible and hidden sector chiral matter Q_i, Σ_i. Further, suppose there is a S_2 permutation symmetry for $i = 1, 2$, so that the low-energy theory has two generations of Q_i and Σ_i for $i = 1, 2$, and states Q_3 and Σ_3 which have different quantum numbers from the first two generations. The visible and hidden sector gauge groups do not have to be identical in general. A class of orbifold examples leading to this type of matter content is presented in the appendix A. The lowest order tree-level Kähler potential inherited from (B.1) by this set of states is

$$K = -\ln \det_{i=1,2}\left(T_{i\bar{j}} + T^\dagger_{i\bar{j}} - \mathrm{tr}\, Q_i^\dagger Q_{\bar{j}} - \mathrm{tr}\,\Sigma_i^\dagger \Sigma_{\bar{j}}\right) - \ln\left(T_{3\bar{3}} + T^\dagger_{3\bar{3}} - \mathrm{tr}\, Q_3^\dagger Q_{\bar{3}} - \mathrm{tr}\,\Sigma_3^\dagger \Sigma_3\right). \tag{B.13}$$

Assuming that only $T_{i\bar{i}}$ acquire expectation values and that all the $T_{i\bar{j}}$ moduli are stabilized with vanishing auxillary components, then with hidden sector supersymmetry breaking the soft masses are

$$m^2_{i\bar{j}} = m^2_{3/2}\left(\delta_{i\bar{j}} - 2\sqrt{\mathrm{Re}\,T_{i\bar{i}}\,\mathrm{Re}\,T_{j\bar{j}}}\,\frac{F_i F_{\bar{j}}^*}{|W|^2}\right) \tag{B.14}$$

for $i, j = 1, 2$, and

$$m_3^2 = m^2_{3/2}\left(1 - 2\,\mathrm{Re}(T_3)\frac{|F_{\Sigma_3}|^2}{|W|^2}\right). \tag{B.15}$$

The mass eigenvalues satisfy the sum rule $\mathrm{Tr}\, m^2 = 0$, and are given by

$$m^2 = m^2_{3/2}(1, -2 + x_3, 1 - x_3)\,, \tag{B.16}$$

where $x_3 \equiv 2\,\mathrm{Re}\,T_3|F_{\Sigma_3}|^2/|W|^2$ is in the range $0 \le x_3 \le 3$. It is important to note that the first two states have the same gauge quantum numbers. Inspecting the mass eigenvalues indicates that with hidden sector supersymmetry breaking the two generations are not degenerate unless hidden sector supersymmetry breaking is isolated in the third hidden generation. If the visible sector soft masses are non-vanishing one of the scalars in the first two generations is tachyonic.

As a final example consider the restrictive class of examples in which visible and hidden sector gauge groups, matter representations, and multiplicities are the same for each i sector. The inherited Kähler potential is then a special case of the sum of logarithms form (B.9) given above. Each logarithm term in the sum in the Kähler potential is then identical. The Kähler potential (B.9) is then invariant under a $S_3 \times \mathrm{U}(1)_R \subset \mathrm{SU}(4)$ global symmetry. If the discrete S_3 symmetry is gauged, then in addition to the dilaton, only a single overall volume modulus $T = \frac{1}{3}\sum_i T_i$, and a single set of visible and hidden sector matter fields, $Q = \frac{1}{3}\sum_i Q_i$ and $\Sigma = \frac{1}{3}\sum_i \Sigma_i$ survive in the truncated theory. The lowest order tree-level Kähler potential for the untwisted states inherited from (B.1) in this case is

$$K = -3\ln\left(T + T^\dagger - \mathrm{tr}\, QQ^\dagger - \mathrm{tr}\,\Sigma\Sigma^\dagger\right) - \ln\left(S^\dagger + S\right). \tag{B.17}$$

This is the sequestered no-scale Kähler potential. It follows from the special diagonal case of the sum of logarithm form discussed above that non-vanishing hidden sector auxiliary components, $F_\Sigma \neq 0$, give rise to vanishing tree-level visible sector scalar masses

$$m^2 = 0\,. \tag{B.18}$$

The lowest order vanishing of the visible sector soft masses, in this case, is also a result of inheritance from the $N = 4$ Kähler potential (B.1) of the underlying microscopic theory since only the S_3 symmetric combination of the diagonal elements of the mass squared matrix (B.8) remain in the low-energy four-dimensional theory.

The lowest order tree-level visible sector scalar spectra which arises from the inherited Kähler potentials discussed above have a number of interesting features. For generic compactifications and hidden sector supersymmetry breaking, visible sector tachyons in general arise at lowest order. In the absence of large corrections which lift these phenomenologically unacceptable tachyons, specific compactifications may project out these dangerous visible sector states. For orbifold compactifications of string theories in ten dimensions this generally requires completely projecting visible sector matter fields out of one or two planes of the six-dimensional internal manifold. In addition, the supersymmetry breaking must occur in hidden sector fields with internal components along the same planes in order to not give tachyonic masses to the remaining visible sector states. This provides an important previously overlooked criterion for compactification model builders. Another feature is that visible sector scalars are not generically degenerate. This was illustrated in models with 2 and 3 generations. In some cases degeneracy does result from hidden sector auxiliary expectation values which are invariant under an unbroken flavor symmetry. But, in this case, it is the unbroken flavor symmetry which of course enforces degeneracy, rather than the specific form of the inherited Kähler potential. In the no-scale sequestered case, degenerate, and in fact vanishing, lowest order scalar masses do result. But, in this case, degeneracy follows simply from the fact that by definition only flavor singlet states survive in the low-energy theory. Note also that the lowest order no-scale Kähler potential arises in only a very restrictive and special class of backgrounds, and is therefore not a generic feature of standard compactifications or brane world scenarios.

The tree-level sum rules for scalar masses squared for the inherited Kähler potentials can be traced to the fact that the Kähler metric of the underlying parent theory with

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16 supersymmetries is necessarily flat. In more general examples in which the relevant underlying microscopic theory has only 8 supersymmetries the parent quaternionic Kähler metric would imply less restrictive sum rules for the inherited Kähler potentials.

Finally, it is important to note that the inherited form of the Kähler potentials are only valid at lowest order tree-level in the truncated low-energy theory and are not protected in any way from corrections. Properly integrating out heavy states in general modifies the full Kähler potential of the low-energy theory. This is true of both standard compactifications in which quantum corrections can be important, as well as brane world scenarios in which tree-level bulk interactions between branes can be important.

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