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Phenomenology and cosmology of weakly coupled string theory *†

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Abstract

The weakly coupled vacuum of $E_8 \otimes E_8$ heterotic string theory remains an attractive scenario for phenomenolgy and cosmology. The particle spectrum is reviewed and the issues of gauge coupling unification, dilaton stabilization and modular cosmology are discussed. A specific model for condensation and supersymmetry breaking, that respects known constraints from string theory and is phenomenologically viable, is described.

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Disclaimer

1 Introduction

It has been suggested that weakly coupled string theory has serious difficulties [1, 2]. Specifically, arguments have been made that *i*) dilaton stabilization is not possible, *ii*) the prediction for coupling constant unification is incompatible with experiment, and *iii*) there are serious cosmological problems. In this talk I address these issues, and show that in the context of a specific model, consistent with the known constraints from string theory, a weakly coupled vacuum presents a viable scenario.

First recall the reasons for the original appeal of the weakly coupled $E_8 \otimes E_8$ heterotic string theory [3] compactified on a Calabi-Yau (CY) manifold [4] (or a CY-like orbifold [5]). The zero-slope (infinite string tension) limit of superstring theory [6] is ten dimensional supergravity coupled to a supersymmetric Yang-Mills theory with an $E_8 \otimes E_8$ gauge group. To make contact with the real world, six of these ten dimensions must be compact – of size much smaller than distance scales probed by particle accelerators, and generally assumed to be of order of the reduced Planck length, 10^{-32}cm. If the topology of the extra dimensions were a six-torus, which has a flat geometry, the 8-component spinorial parameter of $N = 1$ supergravity in ten dimensions would appear as the four two-component parameters of $N = 4$ supergravity in ten dimensions. On the other hand, a Calabi-Yau manifold leaves only one of these spinors invariant under parallel transport; for this manifold the group of transformations under parallel transport (holonomy group) is the $SU(3)$ subgroup of the maximal $SU(4) \cong SO(6)$ holonomy group of a six dimensional compact space. This breaks $N = 4$ supersymmetry to $N = 1$ in four dimensions. As is well known, the only phenomenologically viable supersymmetric theory at low energies is $N = 1$, because it is the only one that admits complex representations of the gauge group that are needed to describe quarks and leptons. For this solution, the classical equations of motion impose the identification of the affine connection of general coordinate transformations on the compact space (described by three complex dimensions) with the gauge connection of an $SU(3)$ subgroup of one of the E_8's: $E_8 \ni E_6 \otimes SU(3)$. As a consequence the gauge group in four dimensions is $E_6 \otimes E_8$. Since the early 1980's, E_6 has been considered the largest group that is a phenomenologically viable candidate for a Grand Unified Theory (GUT) of the Standard Model (SM). Hence E_6 is identified as the gauge group of the "observable sector", and the additional E_8 is attributed to a "hidden sector", that interacts with the former only with gravitational strength couplings.

Orbifolds, which are flat spaces except for points of infinite curvature, are more easily studied than CY manifolds, and orbifold compactifications that closely mimic the CY compactification described above, and that yield realistic spectra with just three generations of quarks and leptons, have been found [7]. In this case the surviving gauge group is $E_6 \otimes G_o \otimes E_8,\ G_o \in SU(3)$.

The low energy effective field theory is determined by the massless spectrum, that is the spectrum of states with masses very small compared with the scales of the string tension and of compactification. Massless bosons have zero triality under an $SU(3)$ which is the diagonal of the $SU(3)$ holonomy group and the (broken) $SU(3)$ subgroup of one E_8. The ten-dimensional vector fields $A_M,\ M = 0, 1, \ldots 9$, appear in four dimensions as four-vectors $A_\mu,\ \mu = M = 0, 1, \ldots 3$, and as scalars $A_m,\ m = M - 3 = 1, \cdots 6$. Under the decomposition $E_8 \ni E_6 \otimes SU(3)$, the E_8 adjoint contains the adjoints of E_6 and $SU(3)$, and the representation $(\mathbf{27}, \mathbf{3}) + (\overline{\mathbf{27}}, \overline{\mathbf{3}})$. Thus the massless spectrum includes gauge fields in the adjoint representation of $E_6 \otimes G_o \otimes E_8$ with zero triality under both $SU(3)$'s, and scalar fields in $\mathbf{27} + \overline{\mathbf{27}}$ of E_6, with triality ± 1 under both $SU(3)$'s, together with their fermionic superpartners. The number of $\mathbf{27}$'s and $\overline{\mathbf{27}}$ chiral supermultiplets that are massless depends on the detailed topology of the compact manifold. The important point for phenomenology is the decomposition under $E_6 \to SO(10) \to SU(5)$:

$$(\mathbf{27})_{E_6} = (\mathbf{16} + \mathbf{10} + \mathbf{1})_{SO(10)} = (\{\overline{\mathbf{5}} + \mathbf{10} + \mathbf{1}\} + \{\mathbf{5} + \overline{\mathbf{5}}\} + \mathbf{1})_{SU(5)} \,. \qquad (1)$$

A $\overline{\mathbf{5}} + \mathbf{10} + \mathbf{1}$ contains one generation of quarks and leptons of the Standard

Model, a right-handed neutrino and their scalar superpartners; a $\mathbf{5}+\overline{\mathbf{5}}$ contains the two Higgs doublets needed in the supersymmetric extension of the Standard Model and their fermion superpartners, as well as color-triplet supermultiplets. Thus all the states of the Standard Model (and its minimal supersymmetric extension) are present.

On the other hand, there are no scalar particles in the adjoint representation of the gauge group. In conventional models for grand unification, these (or one or more other representations much larger than the fundamental one) are needed to break the GUT group to the Standard Model. In string theory, this symmetry breaking can be achieved by the Hosotani, or "Wilson line", mechanism [8] in which gauge flux is trapped around "holes" or "tubes" in the compact manifold, in a manner reminiscent of the Arahonov-Bohm effect. The vacuum value of the trapped flux $< \int d\ell^m A_m >$ has the same effect as an adjoint Higgs, without the complications of having to construct a potential for large Higgs representations that can actually reproduce the properties of the observed vacuum [9]. When this effect is included, the gauge group in four dimensions is

$$\begin{aligned}&\mathcal{G}_{obs}\otimes\mathcal{G}_{hid}, \quad \mathcal{G}_{obs}=\mathcal{G}_{SM}\otimes\mathcal{G}'\otimes\mathcal{G}_o, \quad \mathcal{G}_{SM}\otimes\mathcal{G}'\in E_6, \quad \mathcal{G}_o\in SU(3),\\ &\mathcal{G}_{hid}\in E_8, \quad \mathcal{G}_{SM}=SU(3)_c\otimes SU(2)_L\otimes U(1)_w. \end{aligned} \tag{2}$$

There are of course many other four dimensional string vacua besides the class of vacua described above.

• The gauge group in four dimensions may be larger. Its rank r can be greater than the rank 16 of $E_8\otimes E_8$, if some Kaluza-Klein vector fields $g_{\mu m}$ are massless. In weakly coupled string theory $r\leq 22$, but the group can become arbitrarily large in the strongly coupled regime [10]. These scenarios, however, seem to lead further away from observation.

• The above scenario corresponds to affine level $k_a=1$, where a refers to the gauge group $\mathcal{G}_a$, and the affine levels k_a are the coefficients of Schwinger terms in the current algebra on the string world sheet. For nonabelian groups k_a is a positive integer. If $k_a>1$, it is possible to have adjoint Higgs multiplets in the low energy theory, offering the possiblity of a conventional GUT below the Planck scale (with the problems alluded to above). Models with $k_a \neq 1$ have proven difficult to construct, although some examples exist.

The attractiveness of the picture described above is that the requirement of $N = 1$ supersymmetry (SUSY) naturally results in a phenomenologically viable gauge group and particle spectrum. Moreover, the gauge symmetry can be broken to a product group embedding the Standard Model without the necessity of introducing large Higgs representations.

Supersymmetry is broken in nature. It is well known that spontaneous breaking of global supersymmetry in the observable sector is incompatible with the observed low energy mass spectrum. This fact led Arnowitt and Nath, among others, to the formulation [11] of spontaneously broken supergavity by "hidden sector" interactions that communicate with the observable sector *via* gravitational strength couplings that induce soft SUSY breaking terms in the effective low energy theory, assumed to be a supersymmetric extension of the Standard Model. The $E_8 \otimes E_8$ string theory provides us with the needed hidden sector.

More specifically, if some subgroup $\mathcal{G}_a$ of $\mathcal{G}_{hid}$ is asymptotically free, with a β-function coefficient $b_a > b_{SU(3)}$, defined by the renormalization group equation (RGE)

$$\mu \frac{\partial g_a(\mu)}{\partial \mu} = -\frac{3}{2} b_a g_a^3(\mu) + O(g_a^5), \tag{3}$$

confinement and fermion condensation will occur at a scale $\Lambda_c \gg \Lambda_{QCD}$, and hidden sector gaugino condensation $< \bar{\lambda}\lambda >_{\mathcal{G}_a} \neq 0$, may induce [12] supersymmetry breaking.

To discuss supersymmetry breaking in more detail, we need the low energy spectrum resulting from the ten-dimensional gravity supermultiplet that consists of the 10-d metric g_{MN}, an antisymmetric tensor b_{MN}, the dilaton ϕ, the gravitino ψ_M and the dilatino χ. For the class of CY and orbifold compactifications described above, the massless bosons in four dimensions are the 4-d metric $g_{\mu\nu}$, the antisymmetric tensor $b_{\mu\nu}$, the dilaton ϕ, and certain components of the tensors g_{mn} and b_{mn} that form the real and imaginary parts, respectively, of complex scalars known as moduli; the number of moduli is related to the number of particle generations (# of $\mathbf{27}$'s $-$ # of $\overline{\mathbf{27}}$'s). (More precisely, the scalar components of the chiral multiplets of the low energy theory are obtained as functions of the scalars ϕ, g_{mn} while the pseudoscalars b_{mn} form axionic components of these supermultiplets.) Typically, in a three generation orbifold model there are three moduli t_I; the *vev*'s $< \mathrm{Re}t_I >$ determine the radii of compactification of the three tori of the compact space. In some compactifications there are three other moduli u_I; the *vev*'s $< \mathrm{Re}u_I >$ determine the ratios of the two *a priori* independent radii of each torus. These form chiral multiplets with fermions χ^t_I, χ^u_I obtained from components of ψ_m. The 4-d dilatino χ forms a chiral multiplet with with a complex scalar field s whose *vev*

$$< s >= g^{-2} - \frac{i}{8\pi^2}\theta \tag{4}$$

determines the gauge coupling constant and the θ parameter of the 4-d Yang-Mills theory. The "universal" axion Ims is obtained by a duality transformation [13] from the antisymmetric tensor $b_{\mu\nu}$: $\partial_\mu \mathrm{Im}s \leftrightarrow \epsilon_{\mu\nu\rho\sigma}\partial^\nu b^{\rho\sigma}$.

Because the dilaton couples to the (observable and hidden) Yang-Mills sector, gaugino condensation induces [14] a superpotential for the dilaton superfield[1] S:

$$W(S) \propto e^{-S/b_a}. \tag{5}$$

The vacuum value

$$< W(S) >\propto \left\langle e^{-S/b_a} \right\rangle = e^{-g^{-2}/b_a} = \Lambda_c, \tag{6}$$

is governed by the condensation scale Λ_c as determined by the RGE (3). If it is nonzero, the gravitino $\tilde{G}$ acquires a mass $m_{\tilde{G}} \propto< W >$, and local supersymmetry is broken.

[1]Throughout I use capital Greek or Roman letters to denote a chiral superfield, and the corresponding lower case letter to denote its scalar component.

2 Gauge coupling unification

Precision data on the three gauge couplings of the Standard Model is often construed as indirect evidence for supersymmetry. Using the measured values of these couplings at the Z mass, the RGE equations applied to the SM give approximate, but not exact unification at a scale around 10^{15}GeV, whereas in the minimal supersymmetric extension of the SM (MSSM), the data are consistent with exact unification at the scale $\Lambda_G \approx 2 \times 10^{16}$GeV, with a value of the fine structure constant $\alpha_G = g_G^2/4\pi \approx 1/25$. As discussed above, string theory is not necessarily a GUT, but all the coupling constants are determined – at a scale μ_0 characteristic of the underlying string theory – by the *vev* (4). Allowing for affine levels $k_a \neq 1$, the prediction is

$$g_a^{-2}(\mu_0) = k_a < \mathrm{Re}s > . \tag{7}$$

There are several possibilities within the general context of weakly coupled superstring theory:

• $k_a = 1 \;\forall\; a$, in which case SM unification is predicted[2] as in conventional GUTs, but the theory above the unification scale is not a GUT field theory.

• $k_a \neq 1$ in which case there are two distinct scenarios.

i) Since Higgs superfields in the adjoint representation can appear in the effective field theory if $k_a > 1$, this theory may be a GUT that is broken to the SM by a Higgs *vev* which is determined by the dynamics of the field theory. One recovers the conventional SUSY GUT scenario (with its conventional difficulties) and string theory provides no additional constraint.

ii) The theory may not be a GUT, and the RGE prediction is modified if the k_a of the three SM gauge groups are not the same.

What distinguishes string theory from conventional GUTs is that the scale μ_0 of unification is not an arbitrary parameter, but is determined in terms of the Planck scale by one or more scalar *vev*'s, since the theory contains only

[2]I use the GUT normalization for $U(1)$, not the SM normalization which is sometimes used in the literature.

one fundamental scale: the string tension m_s^{-2}, related to the reduced Planck mass $m_P = (8\pi G_N)^{-\frac{1}{2}} \approx 2\times 10^{18}$GeV by $m_s^2 = m_P^2/ < \mathrm{Re}s >= g^2 m_P^2$. An educated guess [13] would be to identify μ_0 with the scale of compactification:

$$\Lambda_{comp} = \left\langle (\mathrm{Re}t)^{-\frac{1}{2}} \right\rangle m_s = \left\langle (\mathrm{Re}t\mathrm{Re}s)^{-\frac{1}{2}} \right\rangle m_P, \tag{8}$$

where t is the geometric mean of the moduli t_I. Then comparison with the data assuming affine level one would yield $< \mathrm{Re}s >\approx 2$, $< \mathrm{Re}t >\approx 50$. It has been argued by Kaplunovski [15] that such a large value of $< t >$ is not possible; he concluded that consistency requires $t \sim 1$, $\mu_0 \sim m_s$. This argument was revisited [16] in the light of recent progress in strongly coupled string theory (M-theory [17]). The conclusion was that large t (small radius of compactification) is excluded in most string vacua, a notable exception being a particular strongly coupled limit [18] of M-theory in which there is an eleventh dimension much larger than the ten-dimensions on which string theory lives.

However, it is possible to be more precise, particularly in the case of well-understood orbifold compactifications. The field theory loop corrections must be calculated using a regularization procedure that respects supersymmetry. Using a Pauli-Villars (PV) regularization[3] one obtains for the loop-corrected gauge couplings [20] (in the sense of the "Wilson coefficient" of the Yang-Mills field strength operator $-\frac{1}{4}F^a_{\mu\nu}F_a^{\mu\nu}$):

$$\begin{aligned} g_a^{-2} &= k_a\mathrm{Re}s + \sum_I \frac{\ln(t_I+\bar t_I)}{16\pi^2}\left[C^a - \sum_\alpha (1-2q^I_\alpha)C^a_\alpha\right] + \frac{\ln(s+\bar s)}{16\pi^2}\left(C^a - \sum_\alpha C^a_\alpha\right) \\ &\equiv k_a\mathrm{Re}s - 3C^a\frac{\ln\Lambda_a^2}{16\pi^2} + \sum_\alpha C^a_\alpha \frac{\ln\Lambda_\alpha^2}{16\pi^2}, \quad b_a = \frac{1}{8\pi^2}\left(C^a - \frac{1}{3}\sum_\alpha C^a_\alpha\right), \end{aligned} \tag{9}$$

where q^I_α is a "modular weight", C^a is the quadratic Casimir operator in the adjoint representation of the gauge group $\mathcal{G}_a$ and $C^a_\alpha = \mathrm{Tr}(T^a_\alpha)^2$, where T^a_α represents a generator of $\mathcal{G}_a$ on the matter chiral superfield Φ^α. On the RHS

[3] The cancellations present in supersymmetric theories allow one to regulate gauge loops with Pauli-Villars fields in chiral supermultiplets [19], and BRST invariance is maintained.

the Λ's are the PV masses that act as effective cut-offs (which determine the reference value μ_0 in the RGE equations) for the different sectors of the theory in a one-loop calculation. For the matter fields of the "untwisted sector" in orbifold compactification, $\alpha = (AI)$, $q^I_\alpha = q^I_{AJ} = \delta^I_J$, the effective cut-off turns out to be precisely (8) if $< t_I >=< t >$. For the gauge sector, the effective cut-off is $\Lambda_a = g^{-\frac{2}{3}}\Lambda_{comp}$; the factor $g^{-\frac{2}{3}} =< (\mathrm{Re}s)^{\frac{1}{3}} >$ corresponds to a two-loop correction that appears automatically in a one-loop calculation if the PV masses are chosen so as to respect supersymmetry. For the remaining ("twisted sector") contributions, different powers of the moduli appear in the cut-off, but the anticipated result that the effective cut-off for the low energy field theory is determined by the moduli is indeed borne out.

However, (9) cannot be the correct answer. The result is not invariant under the group of modular transformations generated by

$$T_I \to T_I^{-1}, \quad T_I \to T_I + i \tag{10}$$

(together with q^I_α-dependent transformations on Φ^α), that is known [21] to be an exact symmetry of string perturbation theory. The effective field theory has a conformal anomaly (due to the noninvariance of the cut-off) and a chiral anomaly (due to the chiral transformation of fermion fields implicit in the superfield transformations) under (10). These anomalies form a supermultiplet, a constraint that fixes [20] the cut-offs Λ in (9), since the chiral anomaly is unambiguously determined at one loop in quantum field perturbation theory.

The upshot is that one must add counterterms to the effective field theory to restore modular invariance. In general two different counterterms provide the anomaly cancellation. Some or all of the variation under (10) of the field theory loop contributions to $g^{-2}(\mu_0)$ can be canceled if S is not modular invariant, resulting in a compensating variation of the tree level contribution [the first term on the RHS of (9)]. This is the so-called Green-Schwarz counterterm [22] which is model independent; in the chiral formulation for the dilaton supermultiplet it has the effect of modifying the Kähler potential for the dilaton by a T_I-dependent term so as to maintain modular invariance of the Kähler potential. In addition there are model-dependent threshold corrections [23] that arise from integrating out the heavy string and Kaluza-Klein modes; these generate terms in the loop-corrected value of $g^{-2}(\mu_0)$ involving the modular invariant function $|\eta(t_I)|^4\mathrm{Re}t_I$, where $\eta(t_I)$ is the Dedekind eta-function. Matching [20] field theory loop calculations to string loop calculations gives the boundary condition for the running coupling constants in the $\overline{MS}$ renormalization scheme:[4]

$$\begin{aligned} g_a^{-2}(\mu_0) &= \frac{k_a}{2\ell} - \frac{C_a}{8\pi^2}\ln 2 - \frac{1}{16\pi^2}\sum_I b_a^I \ln\left(2\mathrm{Re}t_I|\eta(t_I)|^4\right), \\ \mu_0^2 &= < e^{-1}\ell >= \frac{g^2}{2e} = \frac{m_s^2}{2e} \end{aligned} \tag{11}$$

in reduced Plank units $m_P = 1$, where ℓ is a modular invariant function of the dilaton and the moduli:

$$\ell = \left[2\mathrm{Re}s - \frac{C_{E_8}}{8\pi^2}\sum_I \ln(2\mathrm{Re}t_I)\right]^{-1}. \tag{12}$$

The same result was obtained in [25] using a different regularization procedure for field theory loops (and hence a different definition of the Wilson coefficient). The parameters

$$b_a^I = C_{E_8} - \sum_\alpha (1 - 2q_I^\alpha)\, C_a^\alpha \tag{13}$$

vanish [23] for a large class of orbifolds, in particular the Z_3, Z_7 orbifolds that appear to yield realistic models [7].

The scalar field ℓ and the two-form $b_{\mu\nu}$ described in the introduction are the bosonic components of a linear supermultiplet [26] L, that is dual to the chiral multiplet S. The introduction of the Green-Schwarz counterterm is most naturally implemented within this formalism [27]. There is increasing

[4] Not included here is an additional moduli-dependent contribution from $N = 2$ sectors that is independent of the gauge group [24].

evidence [28] that this is the appropriate formulation for the string dilaton. This formulation is used in the explicit model [29] for gaugino condensation and dilaton stabilization to be discussed in the next section. Although it has been argued [30] that the two formalisms (linear and chiral) are equivalent even in the presence of nonperturbative effects like gaugino condensation, the condensate action is more simply expressed in the linear multiplet formalism.

The bottom line of the present analysis is that for affine level one models without significant threshold corrections [the last term in the RHS of (11)], the unification scale μ_0 is related to the value of the common coupling constant at that scale by $\mu_0 \approx gM_P$, as originally found by Kaplunovski [31]. While the result (11) has been derived only for orbifold compactifications, its large t_I limit agrees with the behavior found in the large t_I limit of Calabi-Yau compactification. If one compares the MSSM fit to the data of the unification scale ($\mu_0 \approx 2 \times 10^{16}$GeV) with the value obtained using the string prediction and the fit value of the coupling at unification [$\alpha \approx 1/25,\ g^2 \approx 0.5 \Rightarrow \mu_0 \approx (2e)^{-\frac{1}{2}} g m_P \approx 6 \times 10^{17}$GeV], there is a mismatch of a factor of about 30 in μ_0 [or a factor of about 3.4 in $\ln(\mu_0)$, which is the quantity that enters in the RGE].

The resolution of this discrepancy has been addressed by a number of authors [32]. The options studied include

• String threshold corrections, *i.e.,* the last term in (11). As noted above, these are absent in many orbifolds. In addition they are small if $< t_I >\sim 1$ as expected, and in most orbifolds they have the wrong sign to correct the discrepancy.

• Affine levels $k_a \neq 1$. The data require that the ratio of the $SU(3)_c$ and $SU(2)_L$ affine levels be close to unity. Since these are integers, they must be very large integers. The only known models have $k_a = 1, 2, 3$, (which is insufficient) and the dimensions of matter representations grows with k_a – taking us farther afield from the observed spectrum. Models with $k_2 = k_3 = 1$ can be made to fit the data provided that $k_1 < 1$, which has few [33] realizations in actual orbifold compactifications.

• Non-MSSM chiral matter. All known orbifold compactifications have additional chiral matter with respect to the MSSM. These transform according to real (reducible) representations of the SM gauge group: $(\mathbf{r}+\bar{\mathbf{r}})_{SM} \in \left(\mathbf{27}+\overline{\mathbf{27}}\right)_{E_6}$, and thus can acquire gauge invariant masses well above the scale of electroweak symmetry breaking. By including such states with masses below the string string scale, (11) can be brought into agreement with the data with weak coupling at the string scale $g(m_s) = O(1)$. Fits to the data require one or more additional states transforming as $\mathbf{3}+\overline{\mathbf{3}}$ under $SU(3)_c$ as well as pairs of $SU(2)_L$ doublets, and/or states transforming as $(\mathbf{3},\mathbf{2})+(\overline{\mathbf{3}},\mathbf{2})$ under $SU(3)_c \otimes SU(2)_L$; these fits can discriminate among models. This appears to be the most straightforward and natural source of threshold corrections to an MSSM fit to the prediction (11).

3 Gaugino condensation and the runaway dilaton

The superpotential (5) results in a potential for the dilaton of the form

$$V(s) \propto e^{-2\mathrm{Re}s/b_a}, \tag{14}$$

which has its minimum at vanishing vacuum energy with $<\mathrm{Re}s> \rightarrow \infty,\ g^2 \rightarrow 0$. This is the runaway dilaton problem. One possible way out is the introduction of a second source of SUSY breaking such that the vacuum energy vanishes, but the superpotential does not: $<W>=<W(S)+W'> \neq 0$. Then the gravitino acquires a mass, and local SUSY is broken. The only scenario of this type that has been realized explicitly is similar to the Hosotani mechanism: W' is a constant induced [14] by a *vev* of the form $<\int dv^{lmn} d_{[l}b_{mn]}>$, where dv^{lmn} is a volume element on the compact manifold, and $d_{[l}b_{mn]}$ is the curl of the anti-symmetric tensor of 10-d supergravity. The difficulty is that this *vev* satisfies a quantization condition [34], which means that if $<W> \neq 0$, the gravitino mass will be near the Planck scale and a large hierarchy for local SUSY breaking cannot be generated (although the generation of a large hierarchy for observable SUSY breaking is not *a priori* excluded [35]).

When the Green-Schwarz term is included, a second dilaton runaway direction is encountered. The potential is no longer positive definite. The small coupling ($\ell \to 0$) behavior is unaffected (with 2Res replaced by ℓ^{-1}), but the potential has a maximum at $b_a \ell = .5$, and is negative for $b_a \ell > .5(1+\sqrt{3}) \approx 1.37$. Since $b_a \leq b_{E_8} \approx .38$, V is negative for $\alpha = \ell/2\pi > 1.32/b_a\pi > .57$. This is the strong coupling regime, and nonperturbative string effects cannot be neglected; they are expected [36] to modify the Kähler potential for the dilaton, which in the perturbative limit is $k(\ell) = \ln \ell$. It has been shown [29, 37] that these contributions can indeed stabilize the dilaton.

An explicit model based on affine level one[5] orbifolds with three untwisted moduli T_I and a gauge group of the form (2) has been constructed. The dilaton is taken to be the $\theta = \bar{\theta} = 0$ component of a real superfield L that satisfies a modified linearity condition; *i.e.*, its chiral projection is

$$\left(\mathcal{D}_{\dot{\alpha}}\mathcal{D}^{\dot{\alpha}} - 8R\right) L = \sum_a W^\alpha_a W^a_\alpha + \sum_a U_a. \tag{15}$$

The right hand side is a chiral superfield of chiral weight 2. The first term is a sum over operators bilinear in the chiral superfields W^a_α of the unconfined Yang-Mills sectors, while the second sum is over condensate superfields of the confined Yang-Mills sectors (*i.e.*, strongly coupled at scales Λ^a_c). With this construction, the condensate superfields automatically satisfy a constraint [30, 38, 40] implied by the Bianchi identity, that is usually ignored in chiral supermultiplet formulations of gaugino condensation. The condensate self-couplings and their couplings to confined matter consist of the classical contribution obtained by the substitution $W^\alpha_a W^a_\alpha \to u_a$ in the standard Yang-Mills Lagrangian, a quantum field theory correction obtained by anomaly matching [39], the Green-Schwarz term [40] and string threshold

[5] This is a simplifying but not a necessary assumption.

corrections needed to restore modular invariance.

In this formalism it is convenient to introduce a function $f(\ell)$ that modifies the string scale coupling constant[6] g, and is related to $g(\ell) = k(\ell) - \ln \ell$ by a differential equation:

$$g^2 = \left\langle \frac{2\ell}{1 + f(\ell)} \right\rangle, \quad \ell g'(\ell) = -\ell f'(\ell) + f(\ell). \tag{16}$$

The results of [36] suggest a parameterization of the form

$$f(\ell) = \sum_{n=0} a_n \ell^{-n/2} e^{-c_n/\sqrt{\ell}}. \tag{17}$$

Retaining just the first one or two terms in the expansion (17), the potential can be made positive definite everywhere and the parameters can be chosen to fit two data points: the coupling constant $g^2 \sim 1$ and the cosmological constant $\Lambda = 0$ (or very nearly so). This is fine tuning, but reasonable values can be obtained for the parameters, *e.g.*, $c_0 = c_1 = 1$, $a_1/a_0 < 0$ with a_0, a_1 in the range 2–5 (positivity of the potential requires $a_0 > 2$).

It should be emphasized that only one condensate u_a is needed for dilaton stabilization. This picture is very different from previously studied "race-track" models [41] where dilaton stabilization is achieved through cancellations among different condensates with similar β-functions. If more than one condensate is present, nonperturbative corrections to the dilaton Kähler potential are still required to stabilize the dilaton.

If the gauge group for the dominant condensate (largest b_a) is not E_8, the moduli t_I are also stabilized through their couplings to twisted sector matter and/or moduli-dependent string threshold corrections. Their vacuum values [42] are at one of the two self dual points[7] in the fundamental domain(see, *e.g.* [7]): $t^I = 1, e^{i\pi/6}$; hence the 4-d string theory is weakly coupled ($\mathrm{Re}t \sim 1$), as well as the 4-d field theory ($g^2 \sim 1$). The moduli auxiliary

[6]If one performs a duality transformation *via* a Lagrange multiplier [27] $S + \bar{S}$, the equations of motion for L give $S + \bar{S} = [f(L) + 1]/L$, and $g^{-2} =< \mathrm{Re}s >$ in the chiral formulation of the classical effective field theory with no GS term.

[7]The coefficients c_n in (17) were assumed to be moduli-independent. If nonperturbative

fields vanish in the vacuum: $< F^I >= 0$, avoiding a potentially dangerous source of flavor changing neutral currents. The nonholomorphic constraint on the condensate superfield U_a implied by the chiral projection (15) is an essential ingredient in this last result [29].

4 Modular cosmology

The soft SUSY breaking parameters were calculated in [29] for $< t_I >= 1$; the results are similar if $< t_I >= e^{i\pi/6}$. If there is only one condensate, the universal axion[8] is massless [44], and the masses of the dilaton and the complex moduli are related to the gravitino mass by

$$m_d \sim \frac{1}{b_a^2} m_{\tilde{G}}, \quad m_{t_I} \approx \frac{2\pi}{3} \frac{(b_{E_8} - b_a)}{(1 + b_{E_8} < \ell >)} m_{\tilde{G}}. \tag{18}$$

In order to generate a hierarchy of order $m_{\tilde{G}} \sim 10^{-15} m_{Pl} \sim 10^3 GeV$ we require [29] $b_{E_8}/b_a \approx 10$ in which case $m_{t_I} \approx 20 m_{\tilde{G}}$, $m_d \sim 10^3 m_{\tilde{G}}$, which may be sufficient [29, 46] to solve the so-called cosmological moduli problem [47, 2, 48]. Since $m_d \sim 10^6 GeV$, its decay does not contribute to the moduli problem. The moduli masses are about $20 TeV$, which is sufficient to evade the late moduli decay problem [47], but requires R-parity violation [48] to avoid a large relic LSP density. If R-parity is conserved, this problem can be evaded if the moduli are stabilized at or near their vacuum values – or for a modulus that is itself the inflaton.

string contributions turn out to be moduli-dependent, and hence not modular invariant, as found [43] in a different orbifold from the class considered here, the moduli *vevs* would be slightly shifted from their self-dual points.

[8] In [29] the condensate superfields U_a are introduced as nonpropagating fields that are determined by the their equations of motion as functions of the other fields. A dynamical condensate has been studied [45] for the case $\mathcal{G}_{hid} = E_8$, and it was shown that the mass of the condensate superfield is larger than the condensate scale Λ_c. When this field is integrated out, the static E_8 model of [29] is recovered. In particular, the massless axion is essentially the universal one, up to $O(m_{\tilde{G}}^2/\Lambda_c^2)$ mixing effects.

In [42], an explicit model for inflation, based on this effective theory, was constructed in which the dilaton is stabilized within its domain of attraction, one or more moduli are stabilized at the vacuum value value $t_I = e^{i\pi/6}$, and one of the moduli may be the inflaton. It is possible that the requirement that the remaining moduli be in the domain of attraction is sufficient to avoid the problem altogether. For example, if $\mathrm{Im}t_I = 0$, the domain of attraction near $t_I = 1$ is rather limited: $0.6 < \mathrm{Re}t_I < 1.6$, and the entropy produced by dilaton decay with an initial value in this range might be less than commonly assumed.

If there are several condensates with different β-functions, the potential and the masses (18) are dominated by the condensate with the largest β-function coefficient b_a, and the result is essentially the same as in the single condensate case, except that a small mass is generated for the dynamical axion. If there is just one hidden sector condensate, the axion is massless[9] up to QCD-induced effects: $m_a \sim (\Lambda_{QCD}/\Lambda_c)^{\frac{3}{2}} m_{\tilde{G}}$, and it is the natural candidate for the Peccei-Quinn axion. Because of string nonperturbative corrections to its gauge kinetic term, the decay constant f_a of the canonically normalized axion is reduced with respect to the standard result by a factor $b_a \ell^2 \sqrt{6} \approx 1/50$ if $b_a \approx .1 b_{E_8}$, which may be sufficiently small to satisfy the (looser) constraints on f_a when moduli are present [2].

5 Implications for phenomenology and open questions

The string nonperturbative corrections necessary to stabilize the dilaton modify the boundary condition (11) for gauge coupling unification. Including

[9] Higher dimension operators might give additional contributions [44] to the axion mass.

the functions $f(\ell)$ and $g(\ell)$ we obtain, with g given in (16):

$$\begin{aligned} g_a^{-2}(m_s) &= g^{-2} + \frac{C_a}{8\pi^2}\ln(\lambda e) - \frac{1}{16\pi^2}\sum_I b_a^I \ln(t_I + \bar{t}_I)|\eta^2(t_I)|^2, \\ m_s^2 &= \lambda g^2 m_{Pl}^2, \quad \lambda = \frac{1}{2}\left\langle e^{g(\ell)-1}[f(\ell)+1]\right\rangle . \end{aligned} \tag{19}$$

In the perturbative case $\lambda = 1/(2e) \approx .18$, while a specific fit [29] with $\alpha = 1/25$ gives a negligible correction: $\lambda = e^{-1.65} \approx .19$. Another fit [42] with $\alpha \approx .17$, used to stabilize the dilaton during inflation, gives $\lambda \approx .15$.

The gaugino masses, as determined at the condensate scale Λ_c, are

$$m_{\lambda_b}(\Lambda_c) \approx -\frac{3g_b^2(\Lambda_c)b_a}{2\,(1 + b_a < \ell >)} m_{\tilde{G}}. \tag{20}$$

The one-loop RGE's predict the same ratios among gaugino masses as is conventionally assumed, but with absolute values that are below experimental bounds if $m_{\tilde{G}} \sim$TeV. However, because the masses (20) are negative (in the phase convention of *e.g.* [49]), they can be driven to much larger values by two-loop corrections [49] if there are sufficiently massive gauge-charged scalars with large Yukawa couplings.

The soft terms in the scalar potential are sensitive to the – as yet unknown – details of matter-dependent contributions to string threshold corrections and to the Green-Schwarz term. Neglecting the former,[10] the Green-Schwarz term is

$$V_{GS} = b\sum_I g^I + \sum_A p_A e^{\sum_I q_I^A g^I}|\Phi^A|^2 + O(|\Phi^A|^4), \quad g^I = -\ln(t^I + \bar{t}^I), \tag{21}$$

and the full Kähler potential reads

$$K = \ln(L) + g(L) + \sum_I g^I + \sum_A e^{\sum_I q_I^A g^I}|\Phi^A|^2 + O(|\Phi^A|^4). \tag{22}$$

[10]If the threshold corrections are determined by a holomorphic function, they cannot contribute to scalar masses.

The cubic "A-terms" and scalar masses are given, respectively, by

$$\begin{aligned} V_A(\phi) &\approx e^{K/2}\sum_A m_{\tilde G}\left[\sum_A \frac{p_A-b_a}{1+p_A<\ell>}\phi^A W_A(\phi) + \frac{3b_a}{1+b_a<\ell>}W(\phi)\right] + \text{h.c.}, \\ m_A^2 &\approx m_{\tilde G}^2\frac{(p_A-b_a)^2}{(1+p_A<\ell>)^2}, \end{aligned} \tag{23}$$

where $W(\Phi)$ is the cubic superpotential for chiral matter. The scalar squared masses are positive and independent of their modular weights by virtue of the fact that $< F^I >$ vanishes in the vacuum. They are universal – and unwanted flavor-changing neutral currents are thereby suppressed – if their couplings to the Green-Schwarz term are universal, in which case the A-terms reduce to

$$V_A(\phi) \approx 3m_{\tilde G}e^{K/2}W(\phi)\frac{p_A\,(1+2b_a<\ell>) - b_a^2<\ell>}{(1+p_A<\ell>)(1+b_a<\ell>)} + \text{h.c.}. \tag{24}$$

If the Green-Schwarz term is independent of the matter fields Φ^A, $p_A = 0$ and we have $m_A = m_{\tilde G},\ A \approx 2m_\lambda$. A plausible alternative is that the Green-Schwarz term depends only on the radii R_I of the three compact tori that determine the untwisted sector part of the Kähler potential (17):

$$K = \ln(L) + g(L) - \sum_I \ln(2R_I^2) + O(|\Phi^A_{\text{twisted}}|^2),$$

where $2R_I^2 = T^I + \bar T^I - \sum_A|\Phi_I^A|^2$ in string units. In this case $p_A = b$ for the untwisted chiral multiplets Φ_I^A and the untwisted scalars have masses comparable to the moduli masses:[11] $m_A = m_t/2 \approx A/3$. If there is a sector with $p_A = b$ and a Yukawa coupling of order one involving $SU(3)$ triplets (*e.g.* $\bar D DN$, where N is a standard model singlet), its two-loop contribution to gaugino masses can generate gluino masses that are well within experimental bounds if $m_{\tilde G} \sim$ TeV. Such a coupling could also generate a *vev* for N, thus

[11] Scenarios in which the sparticles of the first two generations have masses as high as 20 TeV have in fact been proposed [50].

breaking possible additional $U(1)$'s at a scale ~ 10 TeV. The phenomenologically required μ-term of the MSSM may also be generated by the *vev* of a Standard Model gauge singlet or by one of the other mechanisms that have been proposed in the literature [51]. Finally, a flat direction af the classical scalar potential with a mass of $\sim 10TeV$ would attenuate [52] the baryon dilution problem in the Afflick-Dine mechanism for baryogenesis.

More complete information on the Φ-dependence of the string scale gauge coupling functions is required to make precise predictions for soft supersymmetry breaking. Nevertheless this model suggests soft supersymmetry breaking patterns that may differ significantly from those generally assumed in the context of the MSSM. Phenomenological constraints such as current limits on sparticle masses, gauge coupling unification and a charge and color invariant vacuum [53] can be used to restrict the allowed values of the p_A as well as the low energy spectrum of the string effective field theory. A numerical analysis of these issues is in progress [54].

The soft symmetry breaking parameters given above were calculated for the CP invariant vacuum $< t_I >= 1$. If some $< t_I >= e^{i\pi/6}$ in the true vacuum, the results are expected to be similar, except for the possible presence of CP violating phases. It remains to be determined whether these effects can provide the source of the observed CP violation.

In typical orbifold compactifications, the gauge group $\mathcal{G}_{obs} \otimes \mathcal{G}_{hid}$ obtained at the string scale has no asymptotically free subgroup. However in many compactifications with realistic particle spectra [7], the effective field theory has an anomalous $U(1)$ gauge subgroup, which is not anomalous at the string theory level. The anomaly is cancelled [55] by a GS counterterm, similar to the GS term introduced above to cancel the modular anomaly. This results in a D-term that forces some otherwise flat direction in scalar field space to acquire a vacuum expectation value, further breaking the gauge symmetry, and giving large masses to some chiral multiplets, so that the β-function of some of the surviving gauge subgroups may be negative at lower energy. It has been observed [56] that D-term may play a significant role in supersymmetry breaking. Its presence was explicitly invoked in the above-mentioned inflationary model. [42] Its incorporation into the effective condensation potential is under study [57].

6 Conclusions

There have been exciting developments [17] in string dualities that allow the study of a strongly coupled theory by relating it *via* a duality transformation to a different, weakly coupled theory, where perturbative methods can be used. The dilaton runaway problem ($g^2 \to 0$) for the weakly coupled theory then implies, however, that the strongly coupled theory is dual to a different weakly coupled theory with the same problem. This has led to the suggestion [2] that the true vacuum is at strong – but not too strong – coupling, meaning that neither the theory nor its dual is weakly coupled. Then instead of perturbation theory, symmetry arguments, reminiscent of chiral and flavor symmetry arguments used to study low energy QCD, must be used to make predictions for effective field theories from strings.

The results presented here show that, to the extent that one can do reliable calculations (in practice in the context of orbifold compactification) in effective field theories that satisfy the known constraints of string theory, weakly coupled string theory is compatible with phenomenology, provided string nonperturbative effects are taken into account. These effects can stabilize the dilaton at a weakly coupled vacuum value.

The developments in string dualites have led to the unification of all known string theories as different vacua of a single theory: M-theory [17]. Among these vacua the weakly coupled heterotic string remains a viable and attractive possiblity for the description of nature.

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