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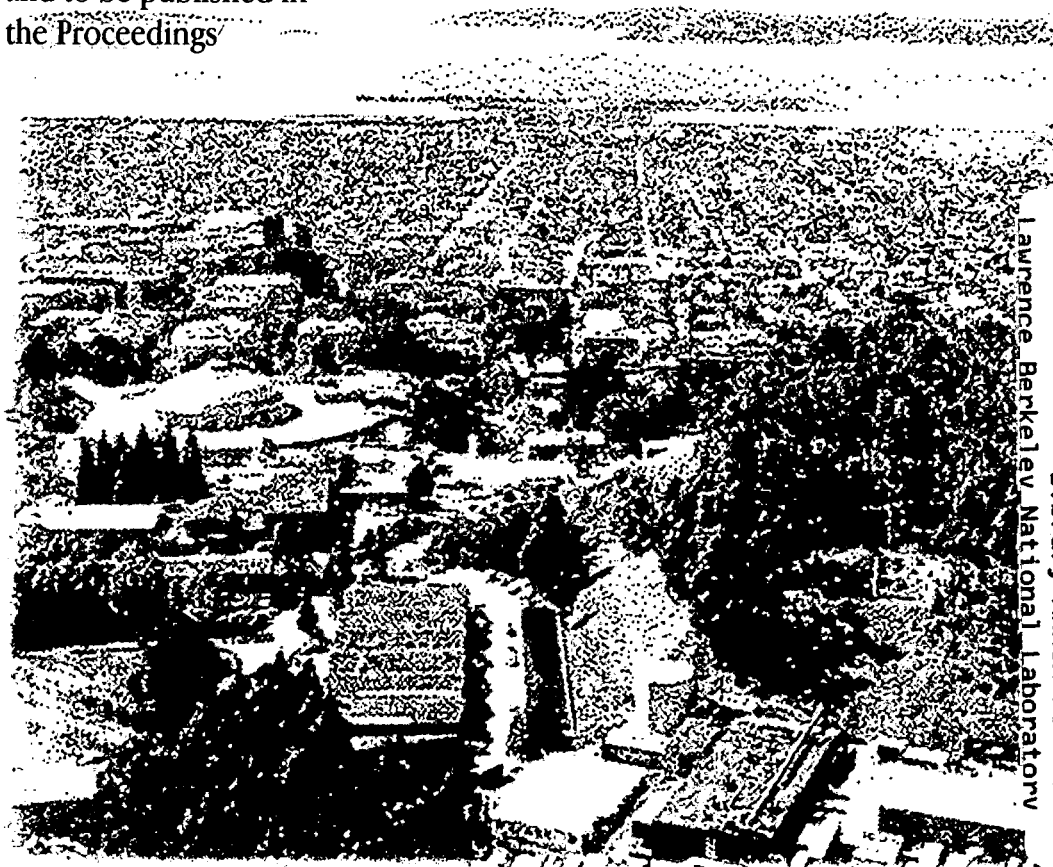
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July 2001

Presented at the
Spacetime Odyssey 2001,
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Progress in weakly coupled string phenomenology*†

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Abstract

The weakly coupled vacuum of $E_8 \otimes E_8$ heterotic string theory remains an attractive scenario for particle physics. The particle spectrum and the issue of dilaton stabilization are reviewed. A specific model for hidden sector condensation and supersymmetry breaking, that respects known constraints from string theory, is described, and its phenomenological and cosmological implications are discussed.

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1 Introduction: Approaches to the Theory of Everything

These days many theorists like to think that they are, in one way or another, working on the Theory of Everything (ToE). There are two basic approaches.

1.1 Bottom Up

The first approach starts with experimental data with the aim of deciphering what it implies for an underlying, more fundamental theory. Its practitioners are usually called phenomenologists. One outstanding datum is the observed large gauge hierarchy, *i.e.*, the ratio of the Z mass, characteristic of the scale of electroweak symmetry breaking, to the reduced Planck scale m_P :

$$m_Z \approx 90\text{GeV} \ll m_P = \sqrt{\frac{8\pi}{G_N}} \approx 2 \times 10^{18}\text{GeV},$$

which can be technically resolved by supersymmetry (SUSY) (among other conjectures—that have been tightly constrained by experimental data, close to the point of being excluded as relevant to the gauge hierarchy). In addition, the conjunction of SUSY and general relativity (GR) inexorably implies supergravity (SUGRA). The absence of observed SUSY partners (sparticles) requires broken SUSY in the vacuum, and a more detailed analysis of the observed particle spectrum constrains the mechanism of SUSY-breaking in the observable sector: spontaneous SUSY-breaking is not viable, leaving soft SUSY-breaking as the only option that preserves the technical SUSY solution to the hierarchy problem. This means introducing SUSY-breaking operators of dimension three or less—such as gauge invariant masses—into the Lagrangian for the SUSY extension of the Standard Model (SM). The unattractiveness of these *ad hoc* soft terms strongly suggests that they arise from spontaneous SUSY breaking in a “hidden sector” of the underlying theory. Based on the above facts, a number of standard scenarios have emerged. These include:

- Gravity mediated SUSY-breaking, usually understood as “Minimal SUGRA” (MSUGRA), with masses of fixed spin particles set equal at the Planck scale; this scenario is typically characterized by

$m_{\text{scalars}} = m_0 > m_{\text{gauginos}} = m_{\frac{1}{2}} \sim m_{\text{gravitino}} = m_{\frac{3}{2}}$ at the weak scale.

- Anomaly mediated SUSY-breaking [1, 2], in which $m_0 = m_{\frac{1}{2}} = 0$ classically; these models are characterized by $m_{\frac{3}{2}} \gg m_0, m_{\frac{1}{2}}$, and typically $m_0 > m_{\frac{1}{2}}$. An exception is the Randall-Sundrum (RS) “separable potential”, constructed [1] to mimic SUSY-breaking on a brane spatially separated from our own in a fifth dimension; in this scenario $m_0^2 < 0$ and m_0 arises first at two loops. More

generally, the scalar masses at one loop depend on the details of Planck-scale physics [3].

- Gauge mediated SUSY uses a hidden sector that has renormalizable gauge interactions with the SM particles. These scenarios are typically characterized by small $m_{\frac{1}{2}}$.

1.2 Top Down

This approach starts from the ToE with the hope of deriving the Standard Model from it; these days most of its practitioners are known as string theorists. The driving motivation is that superstring theory is at present the only known candidate for reconciling GR with quantum mechanics. These theories are consistent in ten dimensions; over the last several years it was discovered that all the consistent superstring theories are related to one another by dualities. These are, in my nomenclature: S-duality: $\alpha \rightarrow 1/\alpha$, and T-duality: Radius $\rightarrow 1/\text{Radius}$, where α is the fine structure constant of the gauge group(s) at the string scale, and “Radius” is a radius of compactification from dimension D to dimension D – 1. Figure 1 shows [4] how these dualities relate the various 10-D superstring theories to one another, and to the currently presumed ToE, M-theory. Not a lot else is known about M-theory, except that it lives in 11 dimensions and involves membranes. In Figure 1 the small circles, line, torus and cylinder represent the relevant compact manifolds in reducing D by one or two. The two $O(32)$ theories are S-dual to one another, while the $E_8 \otimes E_8$ weakly coupled heterotic string theory (WCHS) is perturbatively invariant [5] under T-duality. We will be specifically concentrating on this theory, and T-duality will play an important role.

Another image of M-theory, shown in Figure 2, which I call the “puddle diagram”, indicates [6] that all the known superstring theories, as well as D = 11 SUGRA, are particular limits of M-theory. Currently, there is a lot of activity in type I and II theories, or more generally in theories with branes. Similarly the Hořava-Witten (HW) scenario [7] and its inspirations have received considerable attention. If one compactifies one dimension of the 11-D limit of M-theory, one gets the HW scenario with two 10-D branes, each having an E_8 gauge group. As the radius of this 11th dimension is shrunk to zero, the WCHS scenario is recovered. This is the scenario addressed here, in a marriage of the two approaches that may serve as an illustrative example of the diversity of possible SUSY breaking scenarios.

2 The $E_8 \otimes E_8$ Heterotic String

Let us recall the reasons for the original appeal of the weakly coupled $E_8 \otimes E_8$ heterotic string theory [8] compactified on a Calabi-Yau (CY) manifold [9] (or a CY-like orbifold [10]). The zero-

slope (infinite string tension) limit of superstring theory [11] 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. However 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: $E_8 \ni E_6 \otimes SU(3)$, resulting in $E_6 \otimes E_8$ as the gauge group in four dimensions. 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. 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 [12]. In this case the surviving gauge group is $E_6 \otimes \mathcal{G}_o \otimes E_8$, $\mathcal{G}_o \in SU(3)$. The low energy effective field theory is determined by the massless spectrum, *i.e.*, 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 , $M = 0, 1, \dots, 9$, appear in four dimensions as four-vectors A_μ , $\mu = M = 0, 1, \dots, 3$, and as scalars A_m , $m = M - 3 = 1, \dots, 6$. 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 $E_6 \otimes \mathcal{G}_o \otimes E_8$ 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 $\mathbf{27}$ 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 $E_6 \rightarrow SO(10) \rightarrow SU(5)$:

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

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 [13] in which gauge flux is trapped around “holes” or “tubes” in the compact manifold, in a manner reminiscent of the Aharonov-Bohm effect. The vacuum value of the trapped flux $\langle \int d\ell^m A_m \rangle$ 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 [14]. When this effect is included, the gauge group in four dimensions is

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

There are many other four dimensional string vacua in addition to the class of vacua described above. However the attractiveness of that picture is that the requirement of $N = 1$ 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. In addition, the $E_8 \otimes E_8$ string theory provides a hidden sector needed for a viable theory of spontaneously broken SUSY. 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)

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

confinement and fermion condensation will occur at a scale $\Lambda_c \gg \Lambda_{QCD}$, and hidden sector gaugino condensation $\langle \bar{\lambda}\lambda \rangle_{g_a} \neq 0$, may induce [15] 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. (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.) The number of moduli is related to the number of particle generations ($\#$ of $\mathbf{27}$'s $-$ $\#$ of $\overline{\mathbf{27}}$'s). Typically, in a three generation orbifold model there are three moduli t_I ; the *vev*'s $\langle \text{Re}t_I \rangle$ 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 $\langle \text{Re}u_I \rangle$ determine the ratios of the two *a priori* independent radii of each torus. These form chiral multiplets with fermions χ_I^t, χ_I^u obtained from components of ψ_m . The 4-D dilatino χ forms a chiral multiplet with with a complex scalar field s whose *vev* $\langle s \rangle = g^{-2} - i\theta/8\pi^2$ determines the gauge coupling constant and the θ parameter of the 4-D Yang-Mills theory. The "universal" axion $\text{Im}s$ is obtained by a duality transformation [16] from the antisymmetric tensor $b_{\mu\nu}$: $\partial_\mu \text{Im}s \leftrightarrow \epsilon_{\mu\nu\rho\sigma} \partial^\nu b^{\rho\sigma}$. Because the dilaton couples to the (observable and hidden) Yang-Mills sector, gaugino condensation induces [17] a superpotential for the dilaton superfield[†] S :

$$W(S) \propto e^{-S/b_a}. \quad (4)$$

The vacuum value $\langle W(S) \rangle \propto \langle e^{-S/b_a} \rangle = e^{-g^{-2}/b_a} = \Lambda_c$ is governed by the condensation scale Λ_c as determined by the RGE (3). If it is nonzero, the gravitino acquires a mass $m_{\frac{3}{2}} \propto \langle W \rangle$, and local supersymmetry is broken.

3 The Runaway Dilaton: A Brief Abridged History

The superpotential (4) results in a potential for the dilaton of the form $V(s) \propto e^{-2\text{Re}s/b_a}$, which has its minimum at vanishing vacuum energy and vanishing gauge coupling: $\langle \text{Re}s \rangle \rightarrow \infty, g^2 \rightarrow 0$. This is the notorious runaway dilaton problem. The effective potential for s is in fact determined

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

from anomaly matching [18]: $\delta\mathcal{L}_{eff}(s, u) \longleftrightarrow \delta\mathcal{L}_{hid}(\text{gauge})$, where u , $\langle u \rangle = \langle \bar{\lambda}\lambda \rangle_{G_a}$, is the lightest scalar bound state of the strongly interacting, confined gauge sector. Just as in QCD, the effective low energy theory of bound states must reflect both the symmetries and the anomalies—quantum induced breaking of classical symmetries—of the underlying Yang-Mills theory. It turns out that the effective quantum field theory (QFT) is anomalous under T-duality. Since this is an exact symmetry of heterotic string perturbation theory, it means that the effective QFT is incomplete. This is cured by including model dependent string-loop threshold corrections [19] as well as a “Green-Schwarz” (GS) counter-term [20], named in analogy to a similar anomaly canceling mechanism in 10-D SUGRA [11]. This introduces dilaton-moduli mixing, and the gauge coupling constant is now identified as $g^2 = 2\langle \ell \rangle$, $\ell^{-1} = 2\text{Re}s - b\sum_I \ln(2\text{Re}t^I)$, where $b \leq b_{E_8} = 30/8\pi^2$ is the coefficient of the GS term. It also introduces a second runaway direction, this time at strong coupling: $V \rightarrow -\infty$ for $g^2 \rightarrow \infty$. The small coupling behavior is unaffected, but the potential becomes negative for $\alpha = \ell/2\pi > .57$. This is the strong coupling regime, and nonperturbative string effects cannot be neglected; they are expected [21] to modify the Kähler potential for the dilaton, and therefore the potential $V(\ell, u)$. It has been shown [22, 23] that these contributions can indeed stabilize the dilaton.

The remainder of this paper describes an explicit model [22] based on affine level one[§] orbifolds with three untwisted moduli T_I and a gauge group of the form (2). Retaining just one or two terms of the suggested parameterizations [21] of the nonperturbative string corrections: $a_n \ell^{-n/2} e^{-c_n/\sqrt{\ell}}$ or $a_n \ell^{-n} e^{-c_n/\ell}$, the potential can be made positive definite everywhere and the parameters can be chosen to fit two data points: the coupling constant $g^2 \approx 1/2$ and the cosmological constant $\Lambda \simeq 0$. This is fine tuning, but it can be done with reasonable (order 1) values for the parameters c_n, a_n . If there are several condensates with different β -functions, the potential is dominated by the condensate with the largest β -function coefficient b_+ , and the result is essentially the same as in the single condensate case, except that a small mass is generated for the axion $\text{Im}s$.

4 Features of the Condensation Model

In this model, mass hierarchies arise from the presence of β -function coefficients; these have interesting implications for both cosmology and the spectrum of sparticles—the supersymmetric partners of the SM particles.[¶]

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

[¶]The soft SUSY breaking parameters were calculated in [22] for $\langle t_I \rangle = 1$; the results are similar if $\langle t_I \rangle = e^{i\pi/6}$.

4.1 Modular Cosmology

The masses of the dilaton $d = \text{Re}s$ and the complex t -moduli are related to the gravitino mass by [22]

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

Taking $b = b_{E_8} \approx .38 \approx 10b$, gives a hierarchy of order $m_{\frac{3}{2}} \sim 10^{-15} m_{Pl} \sim 10^3 \text{GeV}$ and $m_{t_I} \approx 20m_{\frac{3}{2}} \approx 20 \text{TeV}$, $m_d \sim 10^3 m_{\frac{3}{2}} \sim 10^6 \text{GeV}$, which is sufficient to evade the late moduli decay problem [24] in nucleosynthesis.

If there is just one hidden sector condensate, the axion $a = \text{Im}s$ is massless up to QCD-induced effects: $m_a \sim (\Lambda_{QCD}/\Lambda_c)^{\frac{3}{2}} m_{\frac{3}{2}} \sim 10^{-9} \text{eV}$, 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 $b_+ \ell^2 \sqrt{6} \approx 1/50$ if $b_+ \approx .1b_{E_8}$, which may be sufficiently small to satisfy the (looser) constraints on f_a when moduli are present [25].

4.2 Sparticle Spectrum

In contrast to an enhancement of the dilaton and moduli masses, there is a suppression of gaugino masses: $m_{\frac{1}{2}} \approx b_+ m_{\frac{3}{2}}$, as evaluated at the scale Λ_a in the tree approximation. As a consequence quantum corrections can be important; for example there is an anomaly-like scenario in some regions of the (b_+, b_+^α) parameter space, where b_+^α is the hidden matter contribution to b_+ . If the gauge group for the dominant condensate (largest b_a) is not E_8 , the moduli t_I are stabilized through their couplings to twisted sector matter and/or moduli-dependent string threshold corrections at a self-dual point, and their auxiliary fields vanish in the vacuum. Thus SUSY-breaking is dilaton mediated, avoiding a potentially dangerous source of flavor changing neutral currents (FCNC). These results hold up to unknown couplings p_A of chiral matter ϕ^A to the GS term: at the scale Λ_a $m_{0A} = m_{\frac{3}{2}}$ if $p_A = 0$, while $m_{0A} = \frac{1}{2} m_{t_I} \approx 10m_{\frac{3}{2}}$ if the scalars couple with the same strength as the T-moduli: $p_A = b$. In addition, if $p_A = b$ for some gauge-charged chiral fields, there are enhanced loop corrections to gaugino masses [26]. Four sample scenarios were studied [27]: A) $p_A = 0$, B) $p_A = b$, C) $p_A = 0$ for the superpartners of the first two generations of SM particles and $p_A = b$ for the third, and D) $p_A = 0$ for the Higgs particles and $p_A = b$ otherwise. Imposing constraints from experiments and the correct electroweak symmetry-breaking vacuum rules out scenarios B and C. Scenario A is viable for $1.65 < \tan \beta < 4.5$, and scenario D is viable for all values of $\tan \beta$, which is

the ratio of Higgs v_{ev} 's in the supersymmetric extension of the SM that requires two Higgs chiral multiplets. The viable range of (b_+, b_+^α) parameter space is shown [28] in Figure 3 for $g^2 = \frac{1}{2}$. The dashed lines represent the possible dominant condensing hidden gauge groups $\mathcal{G}_+ \in E_8$ with chiral matter in the coset space E_8/\mathcal{G}_{hid} .

5 Other Issues in Cosmology

5.1 Flat Directions in the Early Universe

Many successful cosmological scenarios—such as an epoch of inflation—require flat directions in the potential. A promising scenario for baryogenesis suggested [29] by Affleck and Dine (AD) requires in particular flat directions during inflation in sparticle field space: $\langle \tilde{q} \rangle, \langle \tilde{\ell} \rangle \neq 0$, where \tilde{f} denotes the superpartner of the fermion f . While flat directions are common in SUSY theories, they are generally lifted [30] in the early universe by SUGRA couplings to the potential that drives inflation. This problem is evaded [31] in models with a “no-scale” structure, such as the classical potential for the untwisted sector of orbifold compactifications. Although the GS term breaks the no-scale property of the theory, quasi-flat directions can still be found. An explicit model [32] for inflation based on the effective theory described above allows dilaton stabilization within its domain of attraction with one or more moduli stabilized at the vacuum value $t_I = e^{i\pi/6}$. One of the moduli may be the inflaton. The moduli masses (1) are sufficiently large to evade the late moduli decay problem in nucleosynthesis, but unlike the dilaton, they are insufficient to avoid a large relic LSP density without violation [33] of R-parity (a quantum number that distinguishes SM particles from their superpartners). 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. 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 $\text{Im}t_I = 0$, the domain of attraction near $t_I = 1$ is rather limited: $0.6 < \text{Re}t_I < 1.6$, and the entropy produced by dilaton decay with an initial value in this range might be less than commonly assumed. The dilaton decay to its true ground state may provide [34] partial baryon number dilution, which is generally needed for a viable AD scenario.

5.2 Relic Density of the Lightest SUSY Particle (LSP)

Two pertinent questions for SUSY cosmology are:

- Does the LSP overclose the Universe?

- Can the LSP be dark matter?

As discussed by Joe Silk [35], the window for LSP dark matter in the much-studied MSUGRA scenario [36], has become ever more tiny as the Higgs mass limit has increased; in fact there is not much parameter space in which the LSP does not overclose the universe. The ratios of electroweak sparticle masses at the Planck scale determine the composition of the LSP (which must be neutral) in terms of the Bino (superpartner of the SM $U(1)$ gauge boson), the Wino (superpartner of the SM $SU(2)$ gauge boson), and the higgsino (superpartner of the Higgs boson). The MSUGRA assumption of equal gaugino masses at the Planck scale leads to a Bino LSP with rather weak couplings, resulting in little annihilation and hence the tendency to overclose the universe, except in a narrow range of parameter space where the LSP is nearly degenerate with the next to lightest sparticle (in this case a stau $\tilde{\tau}$), allowing significant coannihilation. Relaxing this assumption [28] it was found that a predominantly Bino LSP with a small admixture of Wino can provide the observed amount Ω_d of dark matter. In the condensation model, this occurs in the region indicated by fine points in Figure 3. In this model the deviation from the MSUGRA scenario is due to the importance of loop corrections to small tree-level gaugino masses; in addition to a small Wino component in the LSP, its near degeneracy in mass with the lightest charged gaugino enhances coannihilation. For larger b_+ the LSP becomes pure Bino as in MSUGRA, and for smaller values it becomes Wino-dominated as in anomaly-mediated models which are cosmologically safe, but do not provide LSP dark matter, because Wino annihilation is too fast.

6 Issues: Realistic Orbifold Models

Orbifold compactifications with the Wilson line/Hosotani mechanism needed to break E_6 to the SM gauge group generally have $b_+ \leq b \leq b_{E_8}$. An example is a model [12] with hidden gauge group $O(10)$ and $b_+ = b = b_{O(10)}$. It is clear from (1) that this would lead to disastrous modular cosmology, since the t -moduli are massless. Moreover, in typical orbifold compactifications, the gauge group $\mathcal{G}_{obs} \otimes \mathcal{G}_{hid}$ obtained at the string scale has no asymptotically free subgroup that could condense to trigger SUSY-breaking. However in many compactifications with realistic particle spectra [12], the effective field theory has an anomalous $U(1)$ gauge subgroup, which is not anomalous at the string theory level. The anomaly is canceled [37] 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 masses of order Λ_D to some chiral multiplets, so that the β -function of

some of the surviving gauge subgroups may be negative below the scale Λ_D , typically an order of magnitude below the string scale. The presence of such a D-term was explicitly invoked in the above-mentioned inflationary model [32]. Its incorporation into the effective condensation potential is under study.

There is a large vacuum degeneracy associated with the D-term induced breaking of the anomalous $U(1)$, resulting in many massless “D-moduli” that have the potential for a yet more disastrous modular cosmology [38]. However preliminary results indicate that the D-moduli couplings to matter condensates lift the degeneracy to give cosmologically safe D-moduli masses. Although the D-term modifies the potential for the dilaton, one still obtains moduli stabilized at self-dual points giving FCNC-free dilaton dominated SUSY-breaking, an enhanced dilaton mass m_d and a suppressed axion coupling f_d . An enhancement of the ratio $m_d/m_{\frac{3}{2}}$ can result from couplings to condensates of $U(1)$ -charged D-moduli, that also carry T-modular weights.

7 Conclusions

The lessons of this talk are three-fold:

- Quantitative studies with predictions for observable phenomena are possible within the context of the WCHS.
- Experiments can place restrictions on the underlying theory, such as the hidden gauge sector physics through restriction on the allowed (b_+, b_+^{α}) parameter space, and the couplings and modular weights of D-moduli when an anomalous $U(1)$ is present. Experiments can also inform us about Planck scale physics, such as matter couplings to the GS term. The one-loop corrections to the soft scalar potential are also sensitive to the details of Planck scale physics.
- Searches for sparticles should avoid restrictive assumptions!

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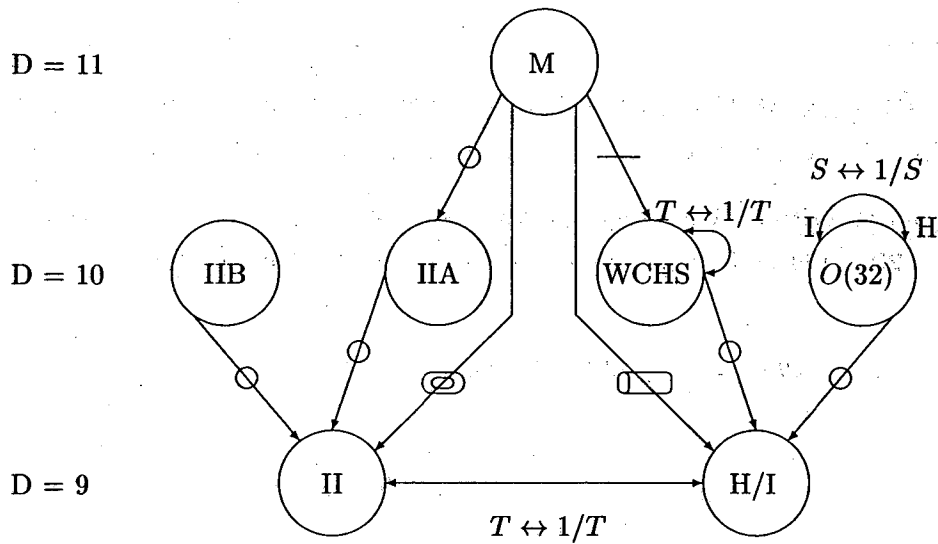


Figure 1: M-theory according to John Schwarz.

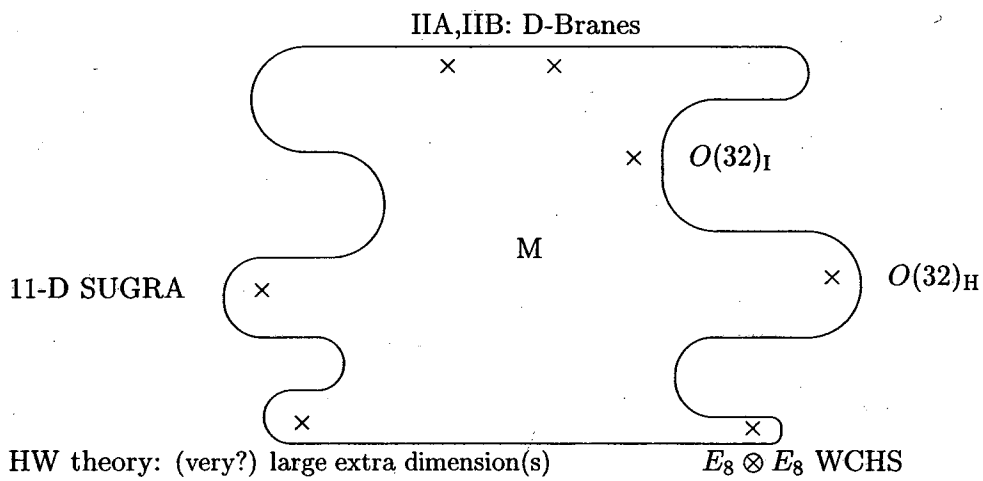


Figure 2: M-theory according to Mike Green.

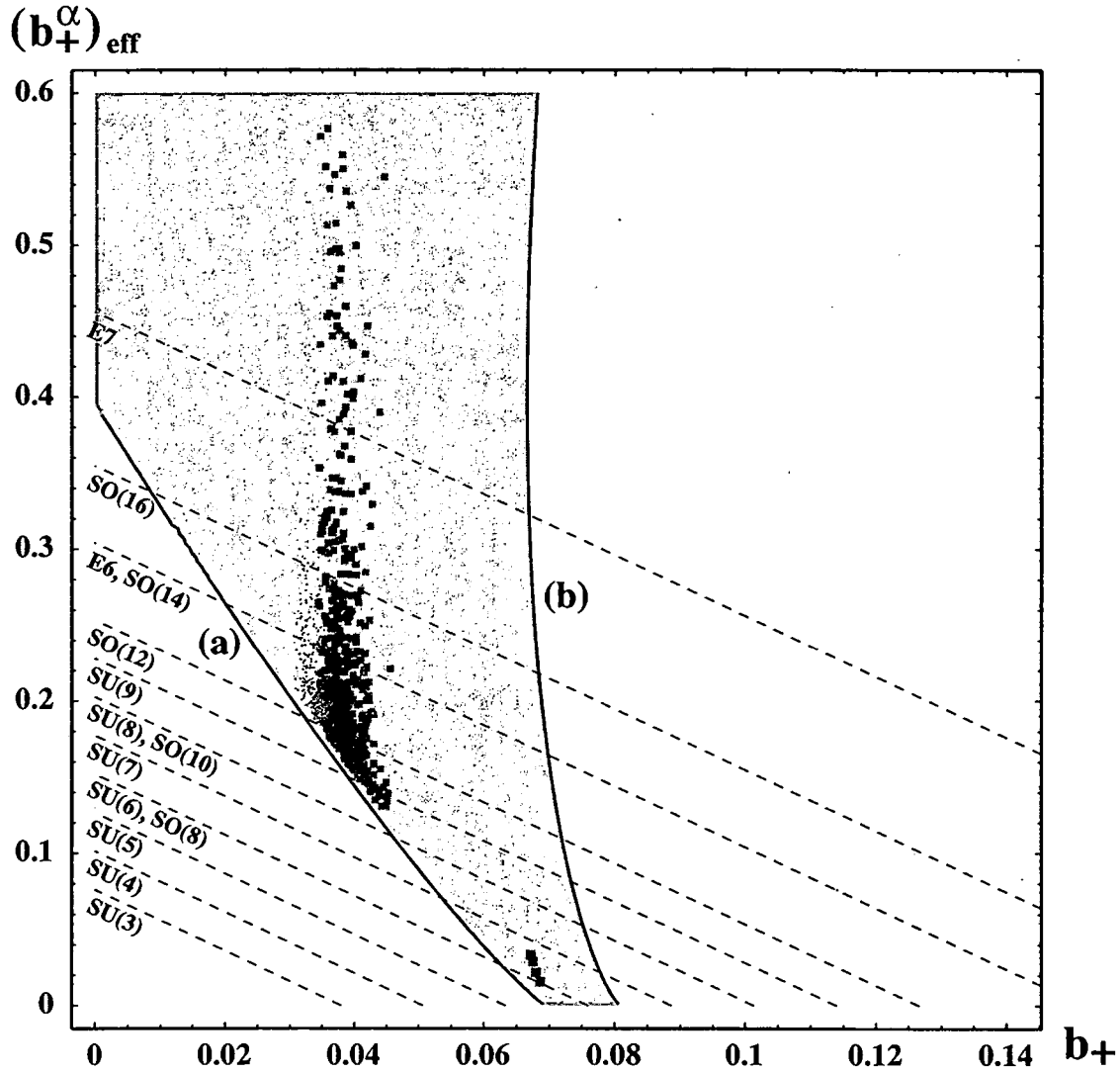


Figure 3: Viable hidden sector gauge groups for scenario A of the condensation model. The swath bounded by lines (a) and (b) is the region defined by $.1 < m_{\frac{3}{2}}/\text{TeV}, \lambda_c < 10$, where λ_c is a condensate superpotential coupling constant. The fine points correspond to $.1 \leq \Omega_d h^2 \leq .3$, and the course points to $.3 < \Omega_d h^2 \leq 1$.

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