PROBING NEW PHYSICS: FROM CHARMM TO SUPERSTRINGS†

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Abstract

Effective theories based on experimental data provide powerful probes and tests of underlying theories in elementary particle physics. Examples within and beyond the Standard Model are discussed, including a specific model for supersymmetry breaking within the context of the weakly coupled heterotic string.

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1 Introduction

This talk begins with some historical remarks about work done during the period of the inception of the Standard Model, especially areas where my own work overlapped significantly with that of Arkady Vainshtein. The remainder of the talk will be devoted to more recent work, in which Arkady’s work on supersymmetry also had an influence. The common thread through the work of both periods is the interconnectedness of theory and experiment in unraveling the elementary structure of nature.

Experiment provides us with effective theories. That is, the data tell us how to write down effective Lagrangians with which we can calculate tree-level S-matrix elements that reproduce the data over some range of energy and distance scales with reasonable accuracy. When we try to take an effective theory seriously as a quantum field theory, we typically encounter difficulties (usually in the form of infinite amplitudes) that suggest a scale at which new physics must come into play. Given a hypothesis for the specifics of the new physics, detailed studies of the effective lower energy theory can test its validity. I will recall examples that contributed to the construction and study of the Standard Model and review approaches to the phenomenology of supergravity and superstring theory. A specific model for supersymmetry breaking within the context of the weakly coupled heterotic string will be discussed.

2 Effective Theories for the Standard Model

An early example of a successful effective Lagrangian is Fermi theory. A series of experiments led theorists such as Fermi, Gamov and Teller to postulate four fermion couplings to describe nuclear $\beta$-decay. Further data revealed the V-A nature of the couplings as well as a set of flavor selection rules for both strangeness-changing and strangeness-conserving (semi-)leptonic decays of hadrons. This led to the interpretation of the effective Fermi interaction as arising from the exchange of heavy, electrically charged vector bosons $W^{\pm}$, with the identification $G_F/\sqrt{2} = g_W^2/m_W^2$, coupled to bilinear quark currents that are Noether currents of the strong interaction. These studies provided early building blocks both of the GWS model of electroweak interactions and of the quark model and QCD.

The Fermi theory is nonrenormalizable. Attempts to treat it as a quantum theory revealed quadratic divergences which suggested that loop corrections must be effectively cut off by new physics at a scale $\Lambda_F \sim 300$ GeV. Similarly, analyses of high energy scattering amplitudes revealed
a breakdown of tree unitarity at cm energies above 600 GeV. These results were the motivation for the construction of the \( pp \) colliders that ultimately produced the \( W, Z \) bosons with masses \( \sim 100 \) GeV, providing the needed cut-off. As it happens, by the time of their discovery, the underlying electroweak theory had already been developed and tested, and the measurement of the weak boson masses provided spectacular confirmation of that theory. Once the underlying theory is known, there is no impediment to treating the low energy effective Fermi theory as a full quantum theory, provided loop integration is appropriately cut off at the physical threshold for new physics, for example

\[
G^2_\Delta \Lambda^2_F \sim G^2_F m_W^2 = g^4_W m_W^2.
\]  

Another example of an effective theory for the Standard Model is the low energy chiral Lagrangian for pions. The approximate \( SU(2)_R \otimes SU(2)_L \) invariance of the strong interactions that was uncovered in the study of pion couplings was a crucial building block for the theory of QCD with very light \( u,d \) quarks. The chiral Lagrangian is nonrenormalizable, but its study at the quantum level, with appropriate cut-offs related to the scale of confinement, has contributed important information on quark masses and other aspects of low energy QCD matrix elements.

In the following I discuss two applications of effective theories for the Standard Model that both Arkady and I were involved in.

### 2.1 The Charm Threshold

The effective Fermi theory constructed from experimental measurements of weak (semi-)leptonic decays contained both strangeness-changing (\( \bar{u}s \)) and strangeness-conserving (\( \bar{u}d \)) quark currents. At the quantum level this leads to a semi-leptonic (\( \bar{d}s \)\( \bar{\mu}\mu \)) four-quark Fermi coupling with effective Fermi constant

\[
G_{\Delta S} \sim \sin \theta_c G^2_F \Lambda^2_{\Delta S},
\]

and consistency with experiment requires \( \Lambda_{\Delta S} \sim 1 \) GeV. The new threshold was provided by the GIM mechanism [1] whereby the \( u \) loop contribution was canceled up to quark mass effects by the loop contribution from the postulated charm quark \( c \), and suggested a charm quark mass \( m_c \sim 1 \) GeV. The BIM mechanism [2], using the same new quark to cancel gauge anomalies in the context of the GWS electroweak theory, strengthened the case for the existence of charm.

In 1973, Ben Lee and I noticed [3] that in the limit of \( u-c \) mass degeneracy, not only the amplitudes for \( K \to \mu\mu \) and \( K^0 \leftrightarrow \bar{K}^0 \) vanish, but so do those for processes like \( K \to \gamma\gamma \), that are observed to be unsuppressed. A similar observation was made by Ernest Ma [4]. A careful
analysis of these and other $K$-decays in the context of the GWS electroweak theory revealed that unsuppressed processes of the latter type have amplitudes $\propto \alpha G_F \ln(m_c/m_u)$, whereas suppressed processes of the former type have amplitudes $\propto \alpha G_F (m_c^2 - m_u^2)/m_W^2$. Consistency with data required $m_c^2 \ll m_c^2 \ll m_W^2$, further supporting very small $u,d$ masses. While the calculation of $K \to \mu\mu$ is theoretically very clean, it turns out that the leading contribution, enhanced by a factor $\ln(m_W/m_c)$, cancels between $W$ and $Z$ exchange diagrams; the result was a rather weak limit: $m_c \leq 9$ GeV. On the other hand, if one was brave enough to attempt to evaluate the matrix element for $K^0 \leftrightarrow \bar{K}^0$, a prediction on the order of a GeV for the charmed quark could be inferred and used to predict branching ratios for other rare kaon decays that had not yet been observed. Once we had done all this work, we learned from Bjorken that we had been scooped, at least in part, by our Russian colleagues Vainshtein and Kriplovich who had considered the same processes and made similar inferences about the charm quark mass.

Ben and I evaluated the $K^0 \leftrightarrow \bar{K}^0$ matrix element using a simple factorization ansatz, and, from the observed value of the $K_L-K_S$ mass difference, found $m_c \approx 1.5$ GeV, neglecting color (QCD was just beginning to emerge at that time as a candidate theory for the strong interactions). This low value worried us, since charm had not been seen in experiments (or so it was assumed; several hints were indeed in the experimental literature; see, e.g., [6]). So we “rounded it off” to 2 GeV, until a reader of the draft insisted that we really found 1.5. So we caved in, but decided to include a second evaluation using colored quarks which gave us back the seemingly safer value of 2. Of course 1.5 is the right answer even though QCD is correct. At the time we knew nothing of a third generation, and its appearance muddied the waters for a while, but since the top quark couplings to lighter quarks are very weak, its contribution to the neutral kaon mass difference is insignificant. Analyses based on a $1/N_c$ expansion suggest that one should neglect the color factor, and lattice QCD calculations support the naïve factorization hypothesis to a good approximation.

2.2 Penguins and the $\Delta I = \frac{1}{2}$ Rule

While the elementary couplings of the quarks and gluons of QCD become manifest at high momentum transfer $|q| \gg \Lambda_{QCD}$, at low energy’s we are forced to work with an effective theory of hadrons. Here I consider weak decays. Purely leptonic processes can of course trivially be dealt with using perturbation theory. Semi-leptonic processes are also tractable because the hadronic vertex can be factored out, and the assumption that the hadronic currents are the Noether currents of flavor/chiral SU(3) is highly predictive. In contrast, in nonleptonic decays, the underlying $\bar{q}q/W$ couplings are completely masked by gluon exchanges across all weak vertices. However we learned
from Ken Wilson how to find the correct effective quark Lagrangian at scales much lower than \( m_W \): the operator product expansion \([10]\). The dominant operator is the one of lowest dimension, in this case a four quark operator. Since gluon exchange conserves helicity, and only left-handed quarks participate in the charge-changing weak interaction, this is a V-A Fermi operator, but the I-spin structure is modified with respect to the original operator. This is because the eigenstates of the S-matrix are in fixed representations of color \( SU(3) \), and these are related by Fermi statistics to the final state ud I-spin in the scattering process \( us \rightarrow ud \). The scattering occurs in a \( J = L = S = 0 \) state, which is antisymmetric, so the color and I-spin states must have the same symmetry. These are 1) an anti-triplet of \( SU(3)_c \) with \( I_{ud} = 0 \), which is attractive, and 2) a sextet of \( SU(3)_c \) with \( I_{ud} = 1 \), which is repulsive. Since the initial state has \( I = \frac{1}{2} \), the former is a pure \( \Delta I = \frac{1}{2} \) transition and the second is a mixture of \( \Delta I = \frac{1}{2} \) and \( \Delta I = \frac{3}{2} \). The explicit calculation \([11]\) gives a mild enhancement of the former and suppression of the latter such that the ratio of the effective Fermi coupling constants is \( G_1^F / G_2^F \approx 5 \), whereas experiment suggests that a factor of about 20 is needed, if one evaluates the amplitudes using simple factorization. Subsequently, Arkady and his collaborators pointed out \([12]\) that we had forgotten a diagram, namely the penguin diagram, depicted in Figure 1. Although our Russian colleagues invented the diagram, its name was first introduced in a paper \([13]\) (where we also estimated the \( b\bar{b} \) photoproduction cross section using the SVZ QCD sum rules \([14]\)) on the properties of the newly discovered B states after a dart game that ended with John Ellis required to use the word “penguin” in his next paper. (More details on the dart game can be found in Misha Shifman’s contribution to a tribute \([15]\) to Arkady.) This diagram gives an effective local four-quark operator because the \( q^2 \) in the gluon propagator cancels the \( q^2 \) in the numerator associated with the color charge radius. It contributes only to \( \Delta I = \frac{1}{2} \) transitions because the penguin’s head has \( \Delta I = \frac{1}{2} \) \((s \rightarrow d)\) and its foot has \( \Delta I = 0 \) since gluon exchange conserves flavor. It further has a different spin structure from the operators analyzed in \([11]\) because, while the penguin’s head has only V-A couplings, its foot is a pure vector coupling because gluon exchange conserves parity. As a consequence, after a Fierz transformation so as to express the effective operator in terms of color singlet quark bilinears, one gets \( S, P \) interactions that can have enhanced matrix elements with respect to the usual V, A couplings.

One is still left with the thorny problem of evaluating hadronic matrix elements of the effective quark operators. This involves techniques such as chiral symmetry and lattice calculations; the latter have indicated an enhancement of penguin diagrams. The \( \Delta I = \frac{1}{2} \) rule in kaon decay was discussed extensively by Bijnens \([16]\). To the extent that the nonrelativistic quark model is a good approximation for baryons, the \( \Delta I = \frac{3}{2} \) operator discussed above cannot contribute because it has
no overlap with the baryon wave function \[17\]. It seems likely to this observer that the undeservedly accurate (\(\sim 2\%\)) \(\Delta I = \frac{1}{2}\) rule does not have a simple explanation, but results from the conjunction of a number of QCD effects, all of which point at least qualitatively in the right direction.

Aside from their contribution to the \(\Delta I = \frac{1}{2}\) rule, the SVZ penguin diagrams have played, and continue to play, an essential role in the analysis of heavy quark decays and CP violation.

3 Is the Standard Model an Effective Theory?

Over the last twenty years experiments have continued to verify the predictions of the Standard Model to great accuracy over a wide range of energy scales. Yet the majority of theorists believe that the Standard Model itself has a limited range of validity. There are a variety of reasons for this, one being the large number of arbitrary parameters. Here I concentrate on just two, very likely connected, difficulties: the gauge hierarchy problem and the failure of the Standard Model to incorporate gravity. The gauge hierarchy problem points to a scale of new physics in the TeV region, providing the motivation for the LHC and an upgraded Tevatron. The fact that gravity couples to everything else suggests new physics at a much higher scale, the Planck scale, that (at least in the “conservative” approach taken here) will never be probed directly by experiments; probing this physics must rely on the interpretation of indirect effects in the effective low energy theory. Incorporating gravity in a fully consistent, predictive and verifiable theory, often referred to as the “Theory of Everything” (ToE) is the holy grail of elementary particle physics.

3.1 Bottom Up Approach

This approach starts from experimental data with the aim of deciphering what it implies for an underlying, more fundamental theory. 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 are by now disfavored by experiment). 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 $m_{\text{scalars}} = m_0 > m_{\text{gauginos}} = m_\frac{1}{2} \sim m_{\text{gravitino}} = m_3$ at the weak scale.

- Anomaly mediated SUSY-breaking [18, 19], in which $m_0 = m_\frac{1}{2} = 0$ classically; these models are characterized by $m_3 \gg m_0, m_\frac{1}{2}$, and typically $m_0 > m_\frac{1}{2}$. An exception is the Randall-Sundrum (RS) “separable potential”, constructed [18] 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 [20].

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

### 3.2 Top Down Approach

This approach starts from a ToE with the hope of deriving the Standard Model from it; the current favored candidate is superstring theory. The driving motivation is that this is at present the only known candidate for reconciling GR with quantum mechanics. Superstring theories are consistent in ten dimensions; in recent years it was discovered that all the consistent superstring theories are related to one another by dualities. These are, in my nomenclature: S-duality: $\alpha \to 1/\alpha$, and T-duality: $R \to 1/R$, where $\alpha$ is the fine structure constant of the gauge group(s) at the
Figure 2: M-theory according to John Schwarz.

Figure 3: M-theory according to Mike Green.
string scale, and $R$ is a radius of compactification from dimension $D$ to dimension $D - 1$. Figure 2 shows 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 2 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 under T-duality. We will be specifically concentrating on this theory, and T-duality will play an important role.

Another image of M-theory, the “puddle diagram” of Figure 3, indicates 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 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.

4 The $E_8 \otimes E_8$ Heterotic String

I first recall the reasons for the original appeal of the weakly coupled $E_8 \otimes E_8$ heterotic string theory compactified on a Calabi-Yau (CY) manifold (or a CY-like orbifold). The zero-slope (infinite string tension) limit of this superstring theory 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 unobservable in current experiments; here they are assumed to be compactified to a size 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 parameters 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: $E_8 \supset E_6 \otimes SU(3)$, 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 [29]. 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, i.e., the spectrum of states with masses very small compared with the scales of the string tension and of compactification. Massless particles 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 $A_m$, $m = M - 3 = 1, \ldots, 6$. Under the decomposition $E_8 \supset E_6 \otimes SU(3)$, the $E_8$ adjoint contains the adjoints of $E_6$ and $SU(3)$, and the representation $(27, 3) + (27, \overline{3})$. Thus the massless spectrum includes gauge fields in the adjoint representation of $E_6 \otimes G_o \otimes E_8$ with zero triality under each $SU(3)$, and scalar fields in $27 + \overline{27}$ of $E_6$, with triality $\pm 1$ under each $SU(3)$, together with their fermionic superpartners. The number of $27$ and $\overline{27}$ chiral supermultiplets that are massless depends on the detailed topology of the compact manifold. The important point for phenomenology is the decomposition under $E_6 \rightarrow SO(10) \rightarrow SU(5)$:

$$\begin{align*}
(27)_{E_6} &= (16 + 10 + 1)_{SO(10)} = (\{5 + 10 + 1\} + \{5 + \overline{5}\} + 1)_{SU(5)}.
\end{align*}$$

A $5 + 10 + 1$ contains one generation of quarks and leptons of the Standard Model, a right-handed neutrino and their scalar superpartners; a $5 + \overline{5}$ 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 $E_6$. In conventional models for grand unification, adjoints (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 in which gauge flux is trapped around “tubes” in the compact manifold, in a manner reminiscent of the Aharonov-Bohm effect. The vacuum value of the trapped flux $< \int d^m A_m >$ has the same effect as an adjoint Higgs, evading the difficulties of constructing a viable Higgs sector encountered in conventional GUTS. Wilson lines reduce the gauge group in four dimensions to

$$G_{\text{obs}} \otimes G_{\text{hid}}, \quad G_{\text{obs}} = G_{\text{SM}} \otimes G' \otimes G_o, \quad G_{\text{SM}} \otimes G' \in E_6, \quad G_o \in SU(3),$$

$$G_{\text{hid}} \in E_8, \quad G_{\text{SM}} = SU(3)_c \otimes SU(2)_L \otimes U(1)_w.$$ (4)

There are many other four dimensional string vacua with different features. The attractiveness of the above 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 includes a hidden sector that can provide a viable mechanism for spontaneous SUSY breaking. More specifically, if some subgroup $G_a$ of $G_{\text{hid}}$ is asymptotically free, with a $\beta$-function coefficient $b_a > b_{SU(3)}$, defined by the renormalization group equation (RGE)

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

confinement and fermion condensation will occur at a scale $\Lambda_c \gg \Lambda_{QCD}$, and hidden sector gaugino condensation $< \bar{\lambda} \lambda > 
= 0$, may induce 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 $\phi$, the gravitino $\psi_M$ and the dilatino $\chi$. For the class of CY and orbifold compactifications described above, the zero-triality massless bosons in four dimensions are the 4-D metric $g_{\mu\nu}$, the antisymmetric tensor $b_{\mu\nu}$, the dilaton $\phi$, 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 $\phi, g_{mn}$, while the pseudoscalars $b_{mn}$ form axionic components of these supermultiplets.) The number of moduli is related to the number of particle generations ($\# \text{ of } 27\text{'s} - \# \text{ of } \overline{27}\text{'s}$). Typically, in a three generation orbifold model there are three moduli $t_I$; the vev’s $< \text{Re} t_I >$ 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 \(< Re u_I >\) determine the ratios of the two \textit{a priori} independent radii of each torus. These form chiral multiplets with fermions \( \chi^I, \chi^u_I \) obtained from components of \( \psi_m \). The 4-D dilatino \( \chi \) forms a chiral multiplet with with a complex scalar field \( s \) whose \( \text{vev} < s > = g^{-2} - \theta / 8\pi^2 \) determines the gauge coupling constant and the \( \theta \) parameter of the 4-D Yang-Mills theory. The “universal” axion \( \text{Im} \ s \) is obtained by a duality transformation \([32]\) from the antisymmetric tensor \( b_{\mu\nu} : \partial_{\mu} \text{Im} \ s \leftrightarrow \epsilon_{\mu\nu\rho\sigma} \partial_{\rho} b^{\sigma} \).

Because the dilaton couples to the (observable and hidden) Yang-Mills sector, gaugino condensation induces \([33]\) a superpotential for the dilaton superfield \( S \) (capital Greek and Roman letters denote chiral superfields, and the corresponding lower case letters denote their scalar components):

\[
W(S) \propto e^{-S/b_\alpha}. \tag{6}
\]

The vacuum value \(< W(S) >\) is governed by the condensation scale \(< e^{-S/b_\alpha} > = e^{-g^{-2}/b_\alpha} = \Lambda_c \) as determined by the RGE \([5]\). If it is nonzero, the gravitino acquires a mass \( m_3 \propto < W > \), and local supersymmetry is broken.

\section{A Runaway Dilaton?}

The superpotential \([3]\) results in a potential for the dilaton of the form \( V(s) \propto e^{-2 \text{Re} s/b_\alpha} \), which has its minimum at vanishing vacuum energy and vanishing gauge coupling: \(< \text{Re} s > \to \infty, g^2 \to 0 >\). This is the notorious runaway dilaton problem. The effective potential for \( s \) is determined by anomaly matching \([34]\): \( \delta L_{\text{eff}}(s, u) \leftrightarrow \delta L_{\text{hid}(\text{gauge})} \), where \( u, < u > = < \bar{\lambda} \lambda >_G \), 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 quantum anomalies 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 \([34]\) and a “Green-Schwarz” (GS) counter-term \([36]\), analogous to a similar anomaly canceling mechanism in 10-D SUGRA \([28]\). This introduces dilaton-moduli mixing, and the gauge coupling constant is now identified as \( g^2 = 2 < \ell >, \ell^{-1} = 2 \text{Re} s - b \sum_I \ln(2 \text{Re} t_I) \), where \( b \leq b_E = 30/8\pi^2 \) is the coefficient of the GS term. This term introduces a second runaway direction at strong coupling: \( V \to -\infty \) for \( g^2 \to \infty \). 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 \([37]\) to modify the Kähler
potential for the dilaton, and therefore the potential $V(\ell, u)$. It has been shown \cite{38, 39} that these contributions can indeed stabilize the dilaton.

In order to carry out the above program, one first needs to know the quantum corrections to the unconfined (i.e., at scales greater than $\Lambda_c$) Yang-Mills Lagrangian for the hidden sector, in order to do the correct anomaly matching. This is where we again encounter Arkady’s work. Just as we can correctly treat the effective Fermi Lagrangian as a quantum theory only if we incorporate the physical cut-off $m_W$, we must include a similar physical cut-off when encountering divergences in the effective SUGRA theory from superstrings. One way to do this is using Pauli-Villars regulators with couplings that respect SUSY \cite{40}. The required cut-offs \cite{40, 41} for regulating the coefficient of the Yang-Mills terms are (neglecting string nonperturbative corrections to the Kähler potential)

$$
\Lambda_A = \Lambda(\Lambda/g)^{2(1-3q_A)}, \quad \Lambda_g = \Lambda g^{\frac{2}{3}}, \quad (7)
$$

for matter and gauge loops, respectively, where $\Lambda = R^{-1}$ is the inverse radius of compactification, and $q_A$ is the average modular weight for the chiral supermultiplet $\Phi^A$. For “untwisted” matter $q_A = \frac{1}{3}$, giving the intuitive value $\Lambda_U = \Lambda = 1/R$, while the gauge-loop cut-off contains the “two-loop” factor $g^{\frac{2}{3}}$ which assures that the anomaly is a chiral superfield. Matching \cite{41} the field theory result with string loop calculations \cite{43} determines the GS term for $\Phi^A = 0$, and one obtains for the Wilson coefficient of the Yang-Mills operator

$$
(g^{-2}_a)_W = 2\ell^{-1} - \frac{1}{16\pi^2}(C_G^a - C_M^a) \ln \ell - \frac{1}{8\pi^2} \sum_A C_A^a \ln(1 - p_A \ell) + O(|\phi^A|^2) + \Delta_a(t), \quad (8)
$$

where $\Delta_a(t)$ is a string threshold correction that is present in some orbifold models, and $p_A$ is the coupling of $|\Phi^A|^2$ to the GS term. Neglecting $\Delta_a$, the RGE invariant \cite{8} can be identified with the general RGE invariant in SUSY Yang-Mills theories found by Shifman and Vainshtein \cite{12}

$$
g^{-2}_a(\mu) - \frac{1}{16\pi^2}(3C_G^a - C_M^a) \ln \mu^2 + C_G^a 8\pi^2 \ln g_a^2(\mu) + \frac{1}{8\pi^2} \sum_A C_A^a \ln Z_A^a(\mu), \quad (9)
$$

where $Z_A^a$ are the renormalization factors for the matter fields, provided we identify

$$
g_a(\mu_s) = 2\ell = (\mu_s/M_P)^2, \quad Z_A^a(\mu_s) = (1 - p_A \ell)^{-1}. \quad (10)
$$

The same boundary condition on the couplings was found in \cite{13} where different regularization procedures were used. (The inclusion of string nonperturbative effects can modify these results; their effects were found to be negligible in the model of \cite{35}).
6 A Condensation Model for SUSY Breaking

In this section I discuss features of an explicit model [38] based on affine level one orbifolds with three untwisted moduli $T_I$ and a gauge group of the form (4). Retaining just one or two terms of the suggested parameterizations [37] of the nonperturbative string corrections: $a_n\ell^{-n/2}e^{-c_n/\ell}$ 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 $\Lambda_{\text{cos}} \approx 0$. 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 $\beta$-function coefficients, the potential is dominated by the condensate with the largest $\beta$-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 $\text{Im}s$. In this model, mass hierarchies arise from the presence of $\beta$-function coefficients; these have interesting implications for both cosmology and the spectrum of sparticles.

6.1 Modular Cosmology

The masses of the dilaton $d = \text{Re}s$ and the complex $t$-moduli are related to the gravitino mass by [38]

$$m_d \sim \frac{1}{b_+^2} m_G, \quad m_{t_I} \approx \frac{2\pi}{3} \frac{(b - b_+)}{(1 + b < \ell >)} m_G.$$ (11)

Taking $b = b_{E8} \approx .38 \approx 10b$, gives a hierarchy of order $m_{\frac{3}{2}} \sim 10^{-15} m_{Pl} \sim 10^3$ GeV and $m_{t_I} \approx 20 m_{\frac{3}{2}} \sim 20$ TeV, $m_d \sim 10^3 m_{\frac{3}{2}} \sim 10^6$ GeV, which is sufficient to evade the late moduli decay problem [14] in nucleosynthesis.

If there is just one hidden sector condensate, the axion $a = \text{Im}s$ is massless up to QCD-induced effects: $m_a \sim (\Lambda_{QCD}/\Lambda_c)^{-\frac{1}{2}} m_{\frac{3}{2}} \sim 10^{-9}$ 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 $b_+ \approx b_{E8}$, which may be sufficiently small to satisfy the (looser) constraints on $f_a$ when moduli are present [15].

6.2 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 $\Lambda_c$ 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 \((b_+, b_+^\alpha)\) parameter space, where \(b_+^\alpha\) is the hidden matter contribution to \(b_+\). If the gauge group for the dominant condensate (largest \(b_\alpha\)) is not \(E_8\), the moduli \(t_I\) are stabilized at self-dual points through their couplings to twisted sector matter and/or moduli-dependent string threshold corrections, 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 the unknown couplings \(p_A\) in Eq. (8): at the scale \(\Lambda_c\) \(m_{0A} = m_3\) if \(p_A = 0\), while \(m_{0A} = \frac{1}{2} m_{t_I} \approx 10 m_3\) 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 \[46\]. Four sample scenarios were studied \[47\]: A) \(p_A = 0\), B) \(p_A = b\) for the superpartners of the first two generations of SM particles and \(p_A = 0\) for the third, C) \(p_A = b\), 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 vev’s in the supersymmetric extension of the SM. The viable range of \((b_+, b_+^\alpha)\) parameter space is shown \[48\] in Figure 4 for \(g^2 = \frac{1}{2}\). The dashed lines represent the possible dominant condensing hidden gauge groups \(G_+ \in E_8\) with chiral matter in the coset space \(E_8/G_{hid}\).

### 6.3 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 \[13\] by Affleck and Dine (AD) requires in particular flat directions during inflation in sparticle field space: \(< \tilde{q} >, < \tilde{\ell} > \neq 0\), where \(\tilde{f}\) denotes the superpartner of the fermion \(f\). While flat directions are common in SUSY theories, they are generally lifted \[50\] in the early universe by SUGRA couplings to the potential that drives inflation. This problem is evaded \[51\] 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 \[52\] 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 \[11\] 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 \[53\] 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...
Figure 4: 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 $0.1 < m_3/\text{TeV}, \lambda_c < 10$, where $\lambda_c$ is a condensate superpotential coupling constant. The fine points correspond to $0.1 \leq \Omega_d h^2 \leq 0.3$, and the course points to $0.3 < \Omega_d h^2 \leq 1$. 

15
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 partial baryon number dilution, which is generally needed for a viable AD scenario.

6.4 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?

The window for LSP dark matter in the much-studied MSUGRA scenario [55], 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 [48] it was found that a predominantly Bino LSP with a small admixture of Wino can provide the observed density fraction $\Omega_d$ of dark matter. In the model of [38], this occurs in the region indicated by fine points in Figure 4. 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.
6.5 Realistic Orbifold Models?

Orbifold compactifications with the Wilson line/Hosotani mechanism needed to break $E_6$ to the SM gauge group generally have $b_+ \leq b \leq b_{E_6}$. An example is a model \cite{29} with hidden gauge group $O(10)$ and $b_+ = b = b_{O(10)}$. It is clear from (11) that this would lead to disastrous modular cosmology, since the $t$-moduli are massless. Moreover, in typical orbifold compactifications, the gauge group $G_{\text{obs}} \otimes G_{\text{hid}}$ 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 \cite{29}, the effective field theory has an anomalous $U(1)$ gauge subgroup, which is not anomalous at the string theory level. The anomaly is canceled \cite{56} 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 be lifted, inducing scalar vev's that further break the gauge symmetry and give masses of order $\Lambda_D$ to some chiral multiplets, so that the $\beta$-function of some of the surviving gauge subgroups may be negative below the scale $\Lambda_D$, 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 \cite{52}. Its incorporation into the effective condensation potential is under study \cite{57}.

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 \cite{58}. 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 $m_{\ell_1}/m_3$ can result from couplings to condensates of $U(1)$-charged D-moduli, that also carry T-modular weights. However, the requirement of a viable scalar/gaugino mass ratio may impose severe restrictions on the details of the effective theory.

7 Lessons

History has taught us that high energy physics can be successfully studied using lower energy effective theories. Hopefully this will be the case for string theory. In particular I have argued that

- 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 sector spectrum and the couplings and modular weights of D-moduli when an anomalous $U(1)$ is present. Data can also inform us about Plank scale physics through matter couplings to the GS term and one-loop corrections to the soft SUSY-breaking scalar potential.

Finally, the SUSY-breaking scenario presented here illustrates the need for sparticle searches to avoid restrictive assumptions based on “standard scenarios” that may be misleading in the absence of concrete models.

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