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Supersymmetry breaking and weakly vs. strongly coupled string theory *

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Abstract

In the context of the field theory limit of superstrings, we consider an almost realistic model of supersymmetry breaking by gaugino condensation which includes, through nonperturbative corrections to the Kähler potential, dilaton stabilization at a value compatible with a weak coupling regime. Invariance under modular transformations is ensured through a Green-Schwarz term and string threshold corrections, which lead to moduli stabilization at the self-dual point. We are thus in a position to discuss several issues of physical relevance: gravitino, dilaton and moduli masses, axion, soft supersymmetry breaking parameters and gauge coupling unification.

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1 Introduction

One of the thorniest problems encountered by supersymmetry breaking mechanisms in the context of superstring models is the stabilization of the dilaton field. Indeed, as long as supersymmetry is conserved, the dilaton corresponds to a flat direction of the scalar potential. Supersymmetry breaking lifts the corresponding degeneracy and should therefore account for a realistic ground state. This is a notoriously difficult problem. According to conventional wisdom:

- the problem of supersymmetry breaking is fundamentally coupled to the problem of dilaton stabilization, as was just explained.
- only nonperturbative string effects can account for dilaton stabilization; in other words only a precise knowledge of the strongly coupled string theory will allow a solution to this problem.

Once these two premises are accepted, it does not seem a far cry to conclude the following:

- supersymmetry breaking can be understood only if one knows how to deal with strongly coupled string theory.

Even in these days of duality and M-theory, this seems a remote possibility.

We argue in this article that, even though the former two statements are more than plausible, one should not readily jump to the latter conclusion. Our discussion will be based on an explicit model constructed in ref. [1] which agrees with the first two assertions but escapes the last one. This model includes supersymmetry broken at a realistic scale, a stabilized dilaton, moduli fields with couplings respecting modular invariance and a zero cosmological constant. We believe that it is sufficiently realistic to allow for a discussion of many issues associated with supersymmetry breaking and moduli physics, based on actual computations rather than educated guesses. Needless to say,

we have no miraculous solution for either dilaton stabilization or the vanishing of the cosmological constant. Although these are incorporated in the model by fixing some parameters (only the second constraint requires fine tuning), the model is still predictive enough in many respects to provide a counter-example to the grim prospects mentioned above.

Before coming to the explicit model, let us explain why one may evade the conclusion that supersymmetry is broken in a regime where string theory is strongly coupled. Dilaton stabilization deals with the behavior of the scalar potential at large and small values of the dilaton field. In what follows we describe the dilaton as the lowest component ℓ of a linear supermultiplet which also includes the antisymmetric tensor field: we believe that this is the natural way to proceed,[1] at least when one deals with the weakly coupled heterotic string. This is the formalism in which the Green-Schwarz mechanism for anomaly cancellation is most easily implemented. It is also the safest one when one starts to include nonperturbative effects –as we will– since the equivalence between the linear and chiral multiplet formulations may be blurred by these effects.

Moreover ℓ is the string coupling: the string perturbation expansion can be organized as a series in $\ell/(8\pi^2)$. As we will see shortly, this *string* coupling should not be confused with the coupling of the *effective field theory.* In the limit of $\ell \to 0$, the string is very weakly coupled and supersymmetry should be restored: one returns to the flat direction and the scalar potential vanishes[2]. If one is prepared to cope with a supersymmetric minimum at vanishing ℓ, the real problem lies in the limit of strong coupling: $\ell \to \infty$. It has been argued [3] that nonperturbative corrections to the Kähler potential indeed stabilize the dilaton for large ℓ. Our model is an explicit illustration of this mechanism [4, 5].

[1]This view is strongly reinforced by the recent results of [2].

[2]Presumably S-duality does not help in this case. If it has anything to say, it is that similarly in the limit of very strong coupling, the potential of the S-dual theory vanishes.

This is important to get a stable minimum but it does not imply that this minimum should lie in the large ℓ region. Indeed arguments based on the unification of couplings tend to indicate that, if there is a physical ground state, it lies in a region where at least the *effective theory* coupling is small. Let us take this opportunity to stress that this is one of the few predictions that one can infer from the low energy values of gauge couplings. It is often stated that the precise value of the gauge coupling unification scale (3×10^{16} GeV) is another one. We think that this is a misleading statement since most models constructed so far that hold a claim for being realistic include new forms of matter which perturb the evolution of the gauge couplings at some intermediate threshold.

In the model that we consider, nonperturbative contributions are included in the Kähler potential in order to stabilize the dilaton. They play an important role in providing a stable ground state. But it is not necessary to choose unnatural values of the parameters in order to ensure that this ground state lies in the phenomenologically preferred region of weak coupling.

Let us briefly describe the explicit model that we use[3]. Supersymmetry is broken through the condensation of gauginos associated with a hidden sector gauge group $\mathcal{G} = \prod_a \mathcal{G}_a$, subgroup of E_8. We use the linear multiplet formulation of gaugino condensation [6, 7] which introduces for each gaugino condensate a vector superfield V_a. The dilaton field is the lowest component of the vector superfield $V = \sum_a V_a$: $\ell = V|_{\theta=\bar{\theta}=0}$. The individual components $V_a|_{\theta=\bar{\theta}=0}$ do not appear in the effective theory component Lagrangian.

The gauge condensate superfields $U_a \simeq \mathrm{Tr}(\mathcal{W}^\alpha \mathcal{W}_\alpha)_a$, where $\mathcal{W}_a$ is the gauge chiral superfield for the group $\mathcal{G}_a$, appear as the chiral projections of V_a:

$$U_a = -(\mathcal{D}_{\dot{\alpha}}\mathcal{D}^{\dot{\alpha}} - 8R)V_a. \tag{1}$$

These condensates are taken to be static (nonpropagating) for reasons to be discussed below.

[3]For more details, we refer the reader to ref. [1].

The Lagrangian describing the gravitational sector including the dilaton simply reads:

$$\mathcal{L}_{KE} = \int \mathrm{d}^4\theta\, E\,[-2 + f(V)]\,, \quad k(V) = \ln V + g(V), \tag{2}$$

where $k(V)$ is the dilaton-dependent part of the Kähler potential and the functions $f(V)$, $g(V)$ parameterize nonperturbative string effects. One might wonder why one needs to introduce the function $f(V)$ besides the correction $g(V)$. It turns out that the two are related by the condition

$$V\frac{\mathrm{d}g(V)}{\mathrm{d}V} = -V\frac{\mathrm{d}f(V)}{\mathrm{d}V} + f, \tag{3}$$

which ensures that the Einstein term has canonical form.

Once these corrections are included, the *effective theory* coupling is $g^2/2 = \ell/[1 + f(\ell)]$ (which is nothing but $(s + \bar{s})^{-1}$ in the dual chiral formulation). Thus if $f(\ell)$ is large, the string theory may be strongly coupled when the effective field theory remains weakly coupled. It turns out in our example that this is not the case at the ground state that we consider (where we find $f(<\ell>)$ to be of order one).

The complete effective Lagrangian includes moduli fields and allows for the presence of matter condensates described by chiral superfields Π^α (taken to be nonpropagating). It includes the terms necessary to reproduce the modular anomaly [8]-[12] and two counterterms which allow the cancellation of this modular anomaly: the Green-Schwarz term [13] and the term induced by string threshold corrections [14][4]. A superpotential for the matter condensates which respects the symmetries of the underlying theory is also added.

The Kähler potential for the effective theory is taken to be:

$$K = k(V) + \sum_I g^I, \quad g^I = -\ln(T^I + \bar{T}^I), \tag{4}$$

[4]However we do not include the moduli-dependent contribution from $N = 2$ sectors that is not related to the Casimir operators of the gauge group [15].

and the complete effective Lagrangian reads:

$$\begin{aligned}\mathcal{L}_{eff} &= \int \mathrm{d}^4\theta\, E\Big(-2 + f(V) + \sum_a V_a \Big\{ b'_a \ln(\bar{U}_a U_a / e^g V) \\ &+ \sum_\alpha b_a^\alpha \ln\left(\Pi_r^\alpha \bar{\Pi}_r^\alpha\right) - \sum_I \frac{b_a^I}{8\pi^2} \ln\left[\left(T^I + \bar{T}^I\right) |\eta^2(T^I)|^2\right] \Big\}\Big) \\ &+ \left(\frac{1}{2}\int \mathrm{d}^4\theta\, \frac{E}{R} e^{K/2} W(\Pi^\alpha, T^I) + \text{h.c.}\right) \end{aligned} \tag{5}$$

where $\Pi_r^\alpha = e^{\sum_I q_I^\alpha g^I/2}\Pi^\alpha$ is a modular invariant combination (q_I^α is the modular weight of the matter condensate Π^α). The coefficients b'_a and b_a^α are found by chiral and conformal anomaly matching. The one-loop beta function coefficient is simply $b_a = b'_a + \sum_\alpha b_a^\alpha$. One can extract the component Lagrangian from (5) using standard procedures. Its bosonic terms are given in ref. [1] and we will not write them explicitly here.

Since the gaugino condensates U_a are nonpropagating in this model, the equations of motion fix their scalar components $u_a \equiv \rho_a e^{i\omega_a}$ in terms of other fields. One might wonder whether some important physics is missed by doing so. We argue that this is the only correct procedure. The dynamical case has been studied in detail in ref. [16] using the simpler example of an E_8 gauge condensate. One may check for this example that both the condensate value ρ and its phase ω are fields of mass larger than the condensate scale. In order to be consistent, one should therefore integrate over them, in which case one recovers the theory with a *static* E_8 condensate [4][5]. We believe that this is general and that the only consistent effective theory below the condensation scale is the static theory. Hence the only dilaton-like scalar and axion-like pseudoscalar in this effective theory are the model independent dilaton and axion of the original string theory (barring some small mixing with the heavy static condensates and some high order mixing with the moduli).

[5] If we take the example of QCD, the same should be correct for the matter condensates, at least in the absence of Goldstone bosons associated with broken global chiral symmetry.

One finds:

$$\rho_a^2 = e^{-2\frac{b'_a}{b_a}} e^K e^{-\frac{(1+f)}{b_a \ell} - \frac{b}{b_a}\sum_I g^I} \prod_I |\eta(t^I)|^{\frac{4(b-b_a)}{b_a}} \prod_\alpha |b_a^\alpha / 4c_\alpha|^{-2\frac{b_a^\alpha}{b_a}}, \tag{6}$$

where $b = C_{E_8}/(8\pi^2)$, C_{E_8} is the quadratic Casimir operator in the adjoint of E_8, and c_α are couplings in the superpotential, presumably of order one.

This is an interesting formula in many respects. One recovers in particular a behavior in $e^{-2/b_a g^2}$ once one identifies correctly the squared gauge coupling of the effective theory as given by $2\ell/(1+f(\ell))$.

This behavior reveals the limitations of the dual formulation using the chiral supermultiplet S to describe the string dilaton. In such a formulation, the gaugino condensate is described by a chiral superfield H^3 which is found to depend on the dilaton as e^{-S/b_a}. This seems to be in agreement with the preceding behavior since a duality transformation yields *at lowest order*:

$$\frac{1+f(L)}{L} = S + \bar{S} \tag{7}$$

and it is often attributed to the "power of holomorphy".

However powerful, holomorphy should be taken with a grain of salt in this instance because higher order corrections modify the relation (7). In particular, the Green-Schwarz term contributes to the right-hand side in a nonholomorphic way. This is the standard observation that, in the chiral formulation, the expansion parameter which is given in terms of the dilaton field as $1/\mathrm{Re}\, S$ needs to be redefined at each order of perturbation theory.

One may wonder how the chiral superfield H^3 dependence in S may undergo these nonholomorphic redefinitions. The point is that it is incorrect to take H^3 as an unconstrained superfield because it describes $\mathrm{Tr}(\mathcal{W}^\alpha \mathcal{W}_\alpha)$ in the effective theory, and this field is subject to the usual Bianchi identities of the gauge sector. The correct treatment is through the use of a 3-form supermultiplet [17] and is naturally implemented in the linear multiplet formulation [7].

The scalar potential reads:

$$16\ell^2 V = \left(1+\ell\frac{dg}{d\ell}\right)\left|\sum_a (1+b_a\ell)\,u_a\right|^2 - 3\ell^2\left|\sum_a b_a u_a\right|^2 + 4\ell^2\,(1+b\ell)\sum_I\left|\frac{F^I}{\mathrm{Re}t^I}\right|^2. \tag{8}$$

where

$$F^I = \sum_a \frac{u_a(b-b_a)}{4(1+b\ell)^2}\left(1+\frac{4}{\eta(t^I)}\frac{\partial\eta}{\partial t}(t^I)\mathrm{Re}t^I\right) \tag{9}$$

is the T^I auxiliary field.

The minimum with respect to the modulus t^I is obtained at the self-dual point $t^I = 1$ where $F^I = 0$. This has some important phenomenological consequences as we will see below.

As for the remaining terms, the potential appears to be dominated by the condensate with the largest one-loop β-function coefficient, so the general case is qualitatively very similar to the single condensate case[6], and it appears that positivity of the potential can always be imposed. One thus does not need to appeal to another source of supersymmetry breaking to cancel the cosmological constant.

The gravitino mass is found to be

$$m_{\tilde{G}} = \frac{1}{3}\langle|M|\rangle = \frac{1}{4}\,\langle|\sum_a b_a u_a|\rangle \tag{10}$$

where M is the supergravity auxiliary field. The mass is thus also governed by the *vev* of the condensate with the largest beta-function coefficient. The scale of supersymmetry breaking is therefore found to be naturally low as long as the gauge group is smaller than E_8. There is in fact a further reduction due to the dependence of the condensate (6) on the moduli: at the self-dual point the reduction factor is approximately $e^{-\pi(b-b_a)/2b_a}$.

Of course, what plays an important role in stabilizing the dilaton as well as fine-tuning the cosmological constant to zero is the exact form of the

[6]This is quite different from racetrack models [18] where one has to play one condensate against another.

functions $f(\ell)$ and $g(\ell)$. They are constrained by (3) and subject to the weak coupling boundary conditions:

$$f(\ell = 0) = 0, \quad g(\ell = 0) = 0. \tag{11}$$

Thus $g(\ell)$ is completely determined once we have chosen $f(\ell)$. Also, the requirement of boundedness of the potential (8) in the strong coupling limit gives:

$$\left(\ell \frac{dg}{d\ell}\right)(\ell \to \infty) \geq 2. \tag{12}$$

Two possible choices for the function f are [3] $f = Ae^{-B/\ell}$ and [19] $f = A_p(\sqrt{\ell})^{-p}e^{-B/\sqrt{\ell}}$. Since all the constraints on the functions f and g are invariant under a rescaling of ℓ, the coefficient B can be fixed to obtain the right value of $< \ell >$ as measured by the gauge coupling at unification. And the parameter A (or A_p) can be used to fine-tune the cosmological constant to zero.

One may look more closely at the second case which is a genuine stringy nonperturbative effect[7]. Taking for illustrative purposes $f = [A_0 + A_1\ell^{-1/2}]\, e^{-B\ell^{-1/2}}$, where the condition (12) requires A_0 to be larger than 2, one finds a realistic minimum for values of the parameters of order one: $B < \ell >^{-1/2} \sim 1.1$ to 1.3, $A_0 \sim 2.7$ to 5.3 and $A_1 \sim -3.1$ to -4.6.

One particularly interesting aspect of the model is axion physics. Pseudoscalar fields are the phases ω_a of the condensates and the so-called model-independent axion which is dual to the fundamental antisymmetric tensor field. The latter couples in a universal way to the $F^{a\mu\nu}\tilde{F}_{a\mu\nu}$ term of each gauge group factor. If again we look at the dynamical model with one E_8 condensate [16] we find that out of the two possible pseudoscalars, the condensate phase is very heavy whereas the model-independent axion remains massless. This is obviously the supersymmetric counterpart of what happens

[7] We do not consider here the case where the coefficient B in the exponent is moduli-dependent [20]. If such nonperturbative stringy contributions turn out not to be modular invariant, this would perturb the moduli ground state away from the self-dual point.

with the scalars. Again, it justifies our approach which treats the condensate degrees of freedom as static in the effective theory. If we allow for more than one condensate, the model-independent axion acquires a very small mass (typically exponentially suppressed relative to the gravitino mass by a factor of order $\sqrt{\rho_2/\rho_1}$ in the two-condensate case [1]). Thus we are left with a very light pseudoscalar which has the right couplings to be the QCD axion. This was actually noted by Banks and Dine [3] who used an argument based on the breaking of a continuous R-symmetry. Our model provides an explicit realization of this phenomenon. Because of corrections to its gauge kinetic term, the model-independent axion must be normalized properly, which gives a reduction factor for the axion decay constant f_a equal to $\ell\sqrt{2[1+\ell dg/d\ell]}$ in reduced Planck mass units ($m_{Pl}=1$). This factor is approximately equal to $b_a\ell^2\sqrt{6}$ at the vacuum for the single gauge condensate case studied in [1]. This gives a suppression factor of about 1/50 if the gravitino mass is found around 10^3 GeV. Higher-dimension operators might give extra contributions to the mass of this axion field [3].

One may easily extract from the scalar potential the masses of the dilaton and of the moduli, which are in particular relevant for cosmology. One finds for the moduli a mass

$$m_t \approx \left\langle \frac{\pi\rho_+}{6}\frac{(b-b_+)}{(1+b\ell)} \right\rangle . \tag{13}$$

where ρ_+ is the hidden sector condensate with the largest β-function coefficient b_+, and for the dilaton

$$m_d \sim \frac{1}{b_+^2} m_{\tilde{G}}. \tag{14}$$

In order to generate a hierarchy of order $m_{\tilde{G}} \sim 10^{-15} m_{Pl} \sim 10^3 GeV$ we require [1] $b/b_+ \approx 10$ in which case $m_t \approx 20 m_{\tilde{G}}$, $m_d \sim 10^3 m_{\tilde{G}}$, which may be sufficient to solve the so-called cosmological moduli problem.

It is also straightforward to determine the soft supersymmetry breaking terms, that are generated at the condensation scale $\mu_{cond} = \langle \rho_+^{\frac{1}{3}} \rangle$, in our

model [1]. The gaugino masses are

$$m_{\lambda_b} = -\left\langle \frac{g_b^2(\mu_{cond})}{8\ell^2}\left(\ell\frac{\partial g}{\partial \ell}+1\right)\sum_a \bar{u}_a\left(1+b_a\ell\right)\right\rangle \approx -\frac{3}{8}\frac{g_b^2(\mu_{cond})b_+^2}{1+b_+\langle\ell\rangle}\langle\bar{u}_+\rangle\,. \tag{15}$$

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. We neglect the former,[8] and write the Green-Schwarz term as

$$V_{GS} = b\sum_I g^I + \sum_A p_A e^{\sum_I q_I^A g^I}|\Phi^A|^2 + O(|\Phi^A|^4), \tag{16}$$

where the Φ^A are gauge nonsinglet chiral superfields, the q_A^I are their modular weights, and the full Kähler potential reads

$$K = k(V) + \sum_I g^I + \sum_A e^{\sum_I q_I^A g^I}|\Phi^A|^2 + O(|\Phi^A|^4). \tag{17}$$

With these assumptions the scalar masses and cubic "A-terms" are given, respectively, by

$$\begin{aligned} m_A^2 &= \frac{1}{16}\left\langle\left|\sum_a u_a\frac{(p_A-b_a)}{(1+p_A\ell)}\right|^2\right\rangle \approx \frac{1}{16}\left\langle\rho_+^2\frac{(p_A-b_+)^2}{(1+p_A\ell)^2}\right\rangle, \\ V_A(\phi) &= \frac{1}{4}e^{K/2}\sum_{a,A}\bar{u}_a\phi^A W_A(\phi)\left[\frac{p_A-b_a}{1+p_A\ell}+b_a-\left(\ell\frac{\partial g}{\partial\ell}+1\right)\frac{1+b_a\ell}{3\ell}\right]+\text{h.c.} \\ &\approx \frac{1}{4}e^{K/2}\bar{u}_+\left[\sum_A\frac{p_A-b_+}{1+p_A\ell}\phi^A W_A(\phi)+\frac{3b_+}{1+b_+\ell}W(\phi)\right]+\text{h.c.}, \end{aligned} \tag{18}$$

where $\phi = \Phi_{\theta=\bar{\theta}=0}$ and $W(\Phi)$ 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

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

suppressed – if their couplings to the Green-Schwarz term are universal, in which case the A-terms reduce to

$$V_A(\phi) \approx \frac{3}{4} e^{K/2} W(\phi) \bar{u}_+ \frac{p_A\,(1+2b_+\ell) - b_a^2 \ell}{(1+p_A\ell)(1+b_+\ell)} + \text{h.c.} \equiv A e^{K/2} W(\phi) + \text{h.c.}. \quad (19)$$

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):

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

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: $m_A = m_t/2 \approx A/3$.[9] Finally, we note that if $b_+ \approx b/10 \approx 1/30$, gaugino masses are suppressed relative to the gravitino mass at the condensation scale $\mu_{cond} \sim 10^{-4} m_{Pl}$: $m_\lambda \sim m_{scalar}/40$. If there is a sector with $p_A = b$ and a Yukawa coupling of order one involving $SU(3)$ (anti-) triplets (*e.g.* $\bar{D}DN$, where N is a standard model singlet), its two-loop contribution to gaugino masses [22] can be more important than the standard one-loop contribution, generating a physical mass for gluinos that is well within experimental bounds for $m_{\tilde{G}} \sim$ TeV. Such a coupling could also generate a *vev* for N, thus 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 [23]. Clearly, a better understanding of the Φ-dependence of the string scale gauge coupling functions is required to make precise predictions for soft supersymmetry breaking. Neverthless our model suggests soft supersymmetry breaking patterns that may differ significantly from those generally assumed in the

[9] Scenarios in which the sparticles of the first two generations have masses as high as 20 TeV have in fact been proposed [21].

context of the MSSM. Phenomenological constraints such as current limits on sparticle masses, gauge coupling unification and a charge and color invariant vacuum 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.

String non-perturbative corrections necessary to stabilize the dilaton could make significant corrections to the unification of gauge couplings. The functions $f(\ell)$ and $g(\ell)$ introduced above and the threshold corrections whose form is dictated by T-duality invariance contribute as follows to the value of couplings at unification:

$$g_a^{-2}(\mu_s) = g_s^{-2} + \frac{C_a}{8\pi^2}\ln(\lambda e) - \frac{1}{16\pi^2}\sum_I b_a^I \ln(t^I + \bar{t}^I)|\eta^2(t^I)|^2, \quad (20)$$

$$g_s^{-2} = \frac{f+1}{2\ell}, \quad \mu_s^2 = \lambda g_s^2 m_{Pl}^2, \quad (21)$$

with

$$\lambda = \frac{1}{2}e^{g-1}(f+1) \quad (22)$$

Let us note however that this parameter is worth $1/(2e) \sim .18$ in the perturbative case and $e^{-1.65} \sim .19$ in the one condensate model.

We stress that the dependence on the radii moduli T^I does not allow an interpretation of the unification scale as the inverse radius of compactification. While the result (20) has been derived only for orbifold compactifications, its large T^I limit is consistent with the behavior found in the large T^I limit of Calabi-Yau compactification. (Note that our moduli are fixed at the self-dual point, therefore far from this limit.)

To conclude, we would like to stress that the model presented above is certainly not final and some of the results obtained, especially on the low energy sector of the theory, may be modified. Possible sources of modification are the presence of an anomalous $U(1)$ symmetry [24] or a constant term in the superpotential that breaks the modular invariance [25, 26]. Our model should be understood as an existence proof of the fact that, even when non-perturbative effects play an important role –especially in the stabilisation

of the dilaton– one may compute in a reliable way quantities of relevance in the low energy world. This opens the way to a discussion of a certain number of issues not addressed here, in particular the cosmology of the dilaton, the axion and the moduli. Moreover, we have shown on our example that predictions may be somewhat different from what appears to be the standard lore. Finally, if one assumes that the strongly interacting string is described along the lines of M-theory [27], one may obtain information about non-perturbative contributions [26] but we expect that the general picture for low energy physics will not be modified drastically.

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