Dynamical supersymmetry breaking
in a superstring inspired model

Nikolaos Irges

Institute for Fundamental Theory,
Department of Physics, University of Florida
Gainesville FL 32611, USA

Abstract

We present a dilaton dominated scenario for supersymmetry breaking in a recently constructed realistic superstring inspired model with an anomalous $U(1)$ symmetry. Supersymmetry is broken via gaugino condensation due to a confining $SU(N_c)$ gauge group in the hidden sector. In particular, we find that by imposing on the model the phenomenological constraint of the absence of observed flavor changing neutral currents, there is a range of parameters related to the hidden sector and the Kähler potential for which we obtain a low energy spectrum consistent with present experimental bounds. As an illustrative example, we derive the low energy spectrum of a specific model. We find that the LSP is the lightest neutralino with a mass of 53 GeV and the lightest Higgs $h^0$ has a mass of 104 GeV.
1 Introduction

During the last few years there has been an increasing activity in trying to construct a complete and phenomenologically viable model with an anomalous $U(1)$ gauge symmetry-$X$, remnant of superstring compactification, with its anomalies canceled via the Green-Schwarz mechanism [1]. The Green-Schwarz anomaly cancelation mechanism occurs if the non zero anomaly coefficients $C_j$ and the corresponding Kac-Moody levels $k_j$, satisfy

$$\frac{C_j}{k_j} = 16\pi^2 \delta_{GS}, \quad \text{for all } j,$$

with

$$\delta_{GS} = \frac{C_g}{192\pi^2}, \quad C_g = Tr(X).$$

In these models, the vacuum expectation value of the dilaton generates a Fayet-Iliopoulos term that triggers the breaking, generating a scale $\xi$, slightly below the string scale [2]. In previous works, it has been demonstrated that models of this class can naturally explain the values of many low energy parameters [3]. Supersymmetry breaking at the same time still remains a mystery, so we can not say a lot about an important part of the phenomenological predictions that these models are capable of producing, the soft supersymmetry breaking parameters like squark and gaugino masses. These, contribute to processes on which there exist strict experimental constraints. For example, they contribute to fcnc effects, known to be very small. A viable model therefore, must have a supersymmetry breaking mechanism that yields squark masses compatible with low energy data on these processes. One class of models attempting to explain the suppression of the supersymmetric contribution to fcnc is decoupling of the third generation by making it much heavier than the first two -approximately degenerate in mass [4]. Another is quark-squark alignment, in which case even if supersymmetry breaking doesn’t yield degenerate or decoupled squarks, fcnc is automatically suppressed [5]. A third is a supersymmetry breaking mechanism that yields degenerate (or almost degenerate) squark masses. Some time ago, a class of models in which supersymmetry breaking is communicated to the low energy world by an anomalous $U(1)$ was proposed, in the context of global supersymmetry [6], [7]. It was soon realized that if the breaking is mainly due to nonzero vevs of $D$-terms the resulting soft masses are proportional to the anomalous $U(1)$ charges of the fields and the desired degeneracy is not achieved. This problem could be overcome if the anomalous $U(1)$ is family blind, but then the family structure of the quark and lepton sector would be trivial, unless there are additional family dependent, non-anomalous $U(1)$ factors ($Y^{(1)}$) in the gauge group. We will show below, that the vevs of the $D$-terms associated to the non-anomalous $U(1)$’s are proportional to that of the anomalous $U(1)$, so they are all contribute the same order of magnitude to flavor changing processes and the problem we had before, remains. The soft masses, in the context of these models, have two types of contributions:

$$m_{\phi_i}^2 = m_0^2 + m_i^2 = \frac{1}{12} K_2 |< F_S >|^2 - \sum_{a=X,Y^{(1)}} q_{\phi_i}^a < D_a >, \quad (1.3)$$

where $m_0$ is the (family independent) contribution from the dilaton $F$-term, $m_i$ the (generically family dependent) contribution from $D$-terms and $K_2$ is the second derivative of the Kähler potential $K$, with respect to the real part of the dilaton field $S$: $y = (1/2)(S + \bar{S})$. $K$ is an unknown function of $y$. We will try to answer the following question: Is it possible to construct a model in which the presence of the family dependent $U(1)$’s is not disastrous for flavor physics? We distinguish two possible phenomenologically viable scenarios.

- $m_0^2 << m_i^2 << m_{1/2}^2$.

The first is a “no scale” type of scenario, where $m_0$ is very close to zero and the non-universal $m_i$ are larger but still small enough, so that when extrapolated to the MSSM scale, they do not give dangerous contributions to fcnc because the running of the soft masses to low scale is dominated by (large) gaugino masses. Such a boundary condition is obtained when the dilaton is stabilized at a very small value of $K_2$. This can be achieved by assuming a weakly coupled form for $K$. An example of such a Kähler potential was proposed in [8]:

$$K = - \ln (2y) - \frac{s_0}{y} + \frac{(b + 4s_0^2)}{24y^2}, \quad (1.4)$$

We assume that the non-universal couplings of the dilaton are suppressed.
which as long as \( b > 0 \) and \( b \leq 1/m^2 \), has a minimum near \( y_0 = s_0 - (1/m) \) and the values of the derivatives at the minimum are given by:

\[
K_1 = -\frac{1}{3s_0}; \quad K_2 = \frac{1}{m^2 s_0}; \quad K_3 = -\frac{2}{ms_0}.
\] (1.5)

Indeed, \( |K_1| < 1 \) and \( K_2 << 1 \), since for reasonable values of the parameters, \( m \) is a number \( \sim 20 - 80 \) (see eq 1.8 and 2.28 below) and therefore \( s_0 \sim 3/2 \). This type of models however, as we will see, are plagued by charge/color breaking minima because of the absence of large contribution to \( m_0 \) from the dilaton, that tends to stabilize the vacuum. In order, therefore, to construct viable models of this type, we will have to assume the existence of additional family independent \( F \)-term contributions from other moduli that stabilize the low energy vacuum.

- \( m_i^2 << m_0^2 \sim m_{1/2}^2 \).

The second is a "minimal sugra" type of scenario, where in order to suppress fcnc, we require that all the \( D \)-term contributions are very small. We will now argue that in the extreme case where these exactly vanish, we can make predictions for the soft parameters. Following \[\text{[7]}\], upon integrating out the heavy gauge field associated with \( X \) and taking the \( D \)-term part of its equation of motion at the minimum, we obtain a relation between the vacuum expectation values of the anomalous \( D \)-term \( D_X \) and that of the dilaton \( F \)-term \( F_S \):

\[
<D_X> = -\frac{1}{4} |<F_S>|^2 \left[ \frac{K_3}{K_1} - \left( \frac{K_2}{K_1} \right)^2 \right] \left( 1 - \frac{16\pi^2 \delta_{GS} K_2}{4 K_1} \right)^{-1},
\] (1.6)

In the presence of additional non-anomalous \( U(1) \) factors, we can similarly integrate out their heavy gauge fields and equation (1.6) still holds. The scale of the FI term can be evaluated from

\[
\frac{\xi}{M} = \frac{1}{2} \frac{1}{\sqrt{16\pi^2 \delta_{GS} K_1}}.
\] (1.7)

For a confining gauge group in the hidden sector, there is a non perturbative contribution to the superpotential that is of the form

\[
W^{(np)} = B e^{-mS},
\] (1.8)

where \( m \) is a model dependent (group theoretical) number and the prefactor \( B \) has units of mass cubed. The contribution of the dilaton to the scalar potential then becomes:

\[
V^S = \frac{1}{4} K_2 M^2 |F_S|^2 = \frac{1}{K_2 M^2} |W^{(np)}|^2 = 4m^2 B^2 e^{-2my} \frac{e^{-2my}}{K_2 M^2},
\] (1.9)

where we have used (1.8) and that

\[
F_S = -\frac{4}{K_2 M^2} \frac{\partial W^{(np)}*}{\partial S}.
\] (1.10)

In a dilaton dominated scenario (where \( V \sim V^S \)), the above term dominates the minimization condition \( V_1 = 0 \) and therefore at the minimum we get the condition

\[
\frac{K_3}{K_2} = -2m.
\] (1.11)

Substituting (1.11) into (1.4), we deduce that in order the \( D \)-term contribution to the soft masses to vanish, the following has to be satisfied at the minimum:

\[
\frac{K_2}{K_1} = -2m.
\] (1.12)

This implies that in order to have degenerate squarks after supersymmetry breaking, the form of the Kähler potential (at the minimum) has to be of the form:

\[
K = ce^{-2my_0}.
\] (1.13)

\(^3\text{We assume that there is no appreciable kinetic mixing between the } U(1) \text{’s.}\)
The constant $c$ can be fixed from (1.7). Doing so, we obtain for $K_2$ at the minimum:

$$K_2 = -\frac{8m_{16}}{16\pi^2 \delta G_S} \left( \frac{\xi}{M} \right)^2.$$  

(1.14)

Knowing $K_2$, allows us to compute the soft masses from (1.3), in a model with known superpotential.

We emphasize that the two types of limits are quite different. In the first “no scale” limit, we will use a perturbative form for the Kähler potential which stabilizes the dilaton at a very small value of $K_2$. In the “minimal sugra” limit, we do not assume a specific form for $K$, but instead we guess its value at the minimum, requiring that the $D$-term vanishes. The whole function $K$ is unknown in this case and it may or may not contain both perturbative and non-perturbative contributions. One has to be very careful with trying to guess the form of the whole function. In fact, by taking for $K$ the sum of the perturbative $-\log (2y)$ term and a typical non-perturbative term, makes the vanishing of the $D$-term rather difficult. This could mean that it is not correct to impose such a constraint or that the form of $K$ can not naively be guessed. The dilaton in this limit, is assumed to be stabilized to $y_0$ and if it happens, it happens at a much higher value of $K_2$ than in the “no scale” case.

Many models with a family dependent, anomalous $U(1)$ constructed in the past, have success explaining the mass ratios and mixings in the quark and charged lepton sectors. Recently however, it was argued that in order a model of this type to naturally relate vacuum stability with the seesaw mechanism and R-parity conservation, the anomalous $U(1)$ has to be family blind. Therefore, a model that can explain the mass hierarchies and mixings has to contain in addition to the anomalous, other, non-anomalous $U(1)$ factors that are family dependent. Interestingly, this is precisely what happens in most realistic superstring compactifications. Such a model was presented in [11], with one anomalous, family blind ($X$) and two non-anomalous family dependent ($Y^{(1)}$, $Y^{(2)}$) $U(1)$’s. The vacuum of this model was shown to be stable and free of flat directions associated with any invariant of the gauge group. It reproduced all quark and charged lepton masses and mixings and predicted neutrino masses and mixings compatible with the solar and atmospheric neutrino data. Proton decay was within the experimental bounds and R-parity was conserved, yielding a stable LSP. In addition, the model had a hidden sector capable of breaking supersymmetry via gaugino condensation. We will use the same observable sector as in [11] but assume a different, simpler hidden sector. In section 2 we show how supersymmetry is broken in a general model with many $U(1)$’s and give expressions for the the soft parameters. In section 3 we apply the general formalism to an explicit model. In section 4 we give our conclusions.

## 2 Supersymmetry Breaking with $U(1)$’s

In this section we extend the supersymmetric breaking mechanism of [1] and [7] for the case of one anomalous and an arbitrary number of non-anomalous $U(1)$’s. The $U(1)$’s break slightly below the string scale by the vevs of a set of singlet fields that we call $\theta_k$. The number of these singlets is equal to the number of the additional $U(1)$’s, so that their charges form a nonsingular square matrix:

$$A = \begin{pmatrix}
  x_1^{(1)} & x_2^{(1)} & x_3^{(1)} & \ldots \\
  y_1^{(1)} & y_2^{(1)} & y_3^{(1)} & \ldots \\
  y_1^{(2)} & y_2^{(2)} & y_3^{(2)} & \ldots \\
  \vdots & \vdots & \vdots & \ddots \\
  \vdots & \vdots & \vdots & \ddots \\
 \end{pmatrix}, \tag{2.15}
$$

where the first row contains the charges of $\theta_k$ with respect to the anomalous symmetry $X$, the second row the charges with respect to the non-anomalous $Y^{(1)}$ and so forth. The supersymmetric vacuum is defined to be the solution of the equations

$$D_{X,Y^{(a)}} = 0. \tag{2.16}$$
Denoting the vevs $< \theta_k >$ by $v_k$, the $D$-flatness condition

$$
\mathcal{A} \begin{pmatrix} |\theta_1|^2 \\ |\theta_2|^2 \\ |\theta_3|^2 \\ \vdots \end{pmatrix} = \begin{pmatrix} \xi^2 \\ 0 \\ 0 \\ \vdots \end{pmatrix}
$$

(2.17)

and the gauge invariance condition for the mass term of a field $t$

$$
\frac{1}{2} M t^2 \left( \frac{\theta_1}{M} \right)^{p_1} \left( \frac{\theta_2}{M} \right)^{p_2} \left( \frac{\theta_3}{M} \right)^{p_3} \ldots
$$

(2.18)

in the superpotential:

$$
\mathcal{A} \begin{pmatrix} p_1 \\ p_2 \\ p_3 \\ \vdots \end{pmatrix} = \begin{pmatrix} n \\ 0 \\ 0 \\ \vdots \end{pmatrix},
$$

(2.19)

give the supersymmetric vacuum constraint

$$
\begin{pmatrix} v_1^2 \\ v_2^2 \\ v_3^2 \\ \vdots \end{pmatrix} = \frac{1}{\rho} \begin{pmatrix} p_1 \\ p_2 \\ p_3 \\ \vdots \end{pmatrix},
$$

(2.20)

where $\xi^2$ is the Fayet-Iliopoulos term generated by the breaking of the anomalous $U(1)$, and $\rho \equiv -n/\xi^2$, where $n$ is the $X$ charge of the field $t^2$. $M$ is the cut-off scale of our theory $\sim 10^{16-17} \text{GeV}$. From (2.20), we can see that the ratio $p_i/v_i^2 = \text{const.}$ We will be looking for a supersymmetry breaking vacuum in the vicinity of this vacuum.

In the following, we assume for simplicity that $G^h$ is a semi-simple, compact, non-Abelian gauge group and that there is only one type of hidden condensates. If there are other hidden fields besides those forming the condensates, they are singlets of $G^h$. We also assume that the number of hidden colors $N_c$, is greater than the number of hidden families $N_f$, in which case the non-perturbative superpotential is particularly simple. Gaugino condensation occurs at a scale where the hidden sector beta function blows up. This scale, is calculated from the renormalization group equation to be

$$
\Lambda = Me^{-8\pi^2 k_h (2S)/b_0} = Me^{-\frac{nN_c}{\omega_d} \rho},
$$

(2.21)

where $k_h$ is the Kac Moody level of the hidden group $G^h$ and $b_0$ is the one loop beta function of the hidden sector. Below this scale, condensates of the hidden “quark” fields $q_i$ will be formed:

$$
t_i = (2q_i\bar{q}_i)^{1/2},
$$

(2.22)

where the index $i$ counts the number of hidden families $N_f$. In the following we will always assume that it is possible to diagonalize the condensate’s mass matrix and in addition that all the condensates have the same mass. In this case $t$ becomes a diagonal matrix with equal entries along the diagonal so we can simplify the calculation by minimizing the scalar potential for a single $t$ and keeping in mind that it is multiplied by an $N_f \times N_f$ unit matrix.

We are ready now to write down the scalar potential to be minimized. It’s general form is

$$
V = V^0 + V^S,
$$

(2.23)

4For the sake of simplicity, we decided to use only one mass scale in our model, the scale at which the observable sector gauge couplings unify, even though in certain cases this might not be the most appropriate.
where

\[ V^0 = \sum_k \left| \frac{\partial W}{\partial \theta_k} \right|^2 + \left| \frac{\partial W}{\partial t} \right|^2 + \sum_{X,Y^{(1)},Y^{(2)},\ldots} \frac{1}{2g^2} D^2 ; \]  

(2.24)

the \(D\)-terms in the above are

\[ D_X = -g_X^2 [x_1 |\theta_1|^2 + x_2 |\theta_2|^2 + x_3 |\theta_3|^2 + \ldots + \frac{1}{2} n |t|^2 + \xi^2] + \ldots \]  

(2.25)

\[ D_{Y^{(a)}} = -g_a^2 [y_1^{(a)} |\theta_1|^2 + y_2^{(a)} |\theta_2|^2 + y_3^{(a)} |\theta_3|^2 + \ldots] + \ldots , \]  

(2.26)

where \(a\) runs over only the non-anomalous \(U(1)\)'s. The superpotential is given by

\[ W = W^{(p)} + W^{(np)} = \frac{1}{2} Mt^2 \left( \theta_1 \right)^{p_1} \left( \theta_2 \right)^{p_2} \left( \theta_3 \right)^{p_3} \ldots + \left( \frac{d_a}{2d_r} - N_f \right) \left( 2\Lambda_{\beta}^2 \right) \frac{1}{2t^2 - N_f} , \]  

(2.27)

where \(d_r\) is the Dynkin index of the representation \(r\) of the hidden gauge group \((r = a\) is the adjoint). Using (2.21) and (2.27), we can express the model dependent constant \(m\) in terms of group theoretical numbers:

\[ m = \frac{8\pi^2 k_{th}}{2d_r - N_f} . \]  

(2.28)

Consider now the minimization conditions

\[ \theta_1 \frac{\partial V}{\partial \theta_1} = (p_1 - 1)|F_{\theta_1}|^2 + p_1 |F_{\theta_2}|^2 + p_1 |F_{\theta_3}|^2 + p_1 \frac{2}{t} F^*_t W^{(p)} - |\theta_1|^2 (x_1 D_X + y_1^{(1)} D_{Y^{(1)}} + y_1^{(2)} D_{Y^{(2)}} + \ldots) = 0 \]  

(2.29)

\[ t \frac{\partial V}{\partial t} = 2|F_\theta|^2 - |\theta_1|^2 + \frac{4}{t} F^*_t W^{(p)} + \frac{1}{\left( \frac{d_a}{2d_r} - N_f \right)^2} W^{(np)} - \frac{1}{2} nt^2 D_X + t |V^{S \ldots}| = 0 \]  

(2.30)

where \(|F_\theta|^2 = |F_{\theta_1}|^2 + |F_{\theta_2}|^2 + |F_{\theta_3}|^2 + \ldots\) and

\[ t V^{S} = \frac{-2}{2d_r - N_f} V^{S} \]  

where \(V^{S} = \frac{\partial V^{S}}{\partial t} . \)  

(2.31)

Defining \(N_c = \frac{2d_r}{d_s}\), (which for fields transforming in the fundamental of \(SU(N_c)\) is just the usual color \(N_c\), the \(F\)-terms entering the above equations are

\[ F^*_k = -\frac{\partial W}{\partial \theta_k} = -\frac{p_k}{\theta_k} W^{(p)} \]  

(2.32)

and

\[ F^*_t = -\frac{\partial W}{\partial t} = -\frac{2}{t} (W^{(p)} - \frac{1}{N_c - N_f} W^{(np)}) . \]  

(2.33)

We are looking for a minimum in the vicinity of the DSW [3] vacuum:

\[ < F_t > \approx 0 , \quad < t > \approx 0 \quad \text{and} \quad < \theta_k > \approx \xi . \]  

(2.34)

The first of the above conditions, implies that

\[ W^{(p)} = \frac{1}{N_c - N_f} W^{(np)} \]  

(2.35)

in the vacuum. Then, (2.30) becomes

\[ \frac{2}{t} F^*_t W^{(p)} = -f t |V^{S} - f|F_\theta|^2 , \]  

(2.36)

\[ ^5 \text{The dots in the two expressions stand for contributions from all other fields. These however should not be allowed to take vevs for obvious reason.} \]
where we have introduced $f = (N_c - N_f)/(N_c - N_f + 1)$. Substituting this into (2.20), we get

\[(p_1 - 1)|F_{\theta_1}|^2 + p_1(1 - f)(|F_{\theta_1}|^2 - |F_{\theta_1}|^2 - p_1f|F_{\theta_1}|^2 - p_1f|V^S|\) = \[\theta_1|^2(x_1D_X + y_1^{D_Y(1)} + y_1^{D_Y(2)} + ...)

which in the vacuum (where $<\theta_a> = v_a$ and everything is evaluated at the minimum), becomes

\[W(p)^\rho^2[(1 - f)(p_1 + p_2 + p_3 + ... - 1) - \rho f|V^S|\] = \[x_1 < D_X > + y_1^{D_Y(1)} + y_1^{D_Y(2)} + ...

where we have used (2.20). Notice that the left hand side of the above equation does not depend on the index \(a\), so all the minimization conditions with respect to all \(\theta_k\) can be obtained from this, by interchanging the subscripts of the right hand side by \(a\). This in turn implies that we can solve these equations for the vevs of the \(D\)-terms:

\[
\begin{pmatrix}
< D_X > \\
< D_Y(1) > \\
< D_Y(2) > \\
\vdots
\end{pmatrix} = \begin{pmatrix} 1 \\ 1 \\ 1 \\ \vdots \end{pmatrix} C(A^{-1})^T
\]

where

\[
C = W(p)^\rho^2 \left[ \frac{1}{N_c + 1 - N_f}(p_1 + p_2 + p_3 + ...) - 1 \right] - \rho \frac{N_c - N_f}{N_c + 1 - N_f}|V^S| - \frac{N_c - N_f}{N_c + 1 - N_f} \left[ \frac{8\pi^2 k_h}{(N_c - N_f + 1)nK_2} \left( \frac{\xi}{M} \right)^2 \right].
\]

The above relation implies that the values of the \(D\)-terms are proportional:

\[< D_Y(a) > = \frac{A_a}{A_X} < D_X >
\]

where \(A_X\) is the sum of the entries of the first column of \(A^{-1}\), \(A_1\) is the sum of the entries of the second column of \(A^{-1}\), etc. This shows that in general, the \(D\)-terms contribute to supersymmetry breaking, but also that for \(det A \neq 0\), if \(< D_X >\) vanishes, then all other \(< D_Y(a) >\) vanish as well.

Using (2.10), (2.21) and (2.27), we find that the vacuum value of the \(F\)-term associated with the dilaton is

\[< F_S > = \epsilon \tilde{m}(8\pi^2 k_h) \frac{4}{K_2} \left( \frac{\xi}{M} \right)^2
\]

and we have defined, as usual, the helpful variables

\[\tilde{m} = M \lambda_1^{p_1} \lambda_2^{p_2} \lambda_3^{p_3} ... \text{ with } \lambda_k \equiv < \frac{\theta_k}{M} >,\]

\[1 \gg \epsilon = \frac{t^2}{2\xi^2} = \left( \frac{\Lambda}{\xi} \right) \frac{\xi}{\tilde{m}} \left( \frac{\xi}{\tilde{m}} \right)^{1 - \frac{N_f}{N_c}} \text{, with } \Lambda = Me^{-8\pi^2 k_h (\frac{2\pi^2}{\alpha_0})}.\]

Here, \(y_0 = < y > = 1/g(M)^2\) where \(g(M)\) is the value of the gauge coupling at the unification scale \(M\) and we have assumed that the dilaton gets somehow stabilized to a reasonable value \(y_0\).\(^6\) The one loop beta function is given by

\[b_0 = 3d_{\alpha} - \sum d_{\tau} \]

\(^6\)We will see that by reasonable we mean a value $\sim 1.5$ (see Appendix)

\(^7\)In our normalization of the indices, for $SU(N_c)$ with $N_f$ families of “quarks” and “antiquarks”, the beta function is $b_0 = 2(3N_c - N_f)$.\(^7\)
We normalize the Dynkin indices so that
\[ Tr(T_a T_b) = d_r \delta_{ab} \] (2.46)
with \( T^r_a \) being the generators of \( G^h \) in the representation \( r \).

Having the expressions of the vevs of the \( D \) and \( F \)-terms, we can now calculate the soft masses. We can assign to each field -generically denoted by \( \phi_i \), with \( i \) being a family index- a set of numbers \( n_k \) such that the term
\[ \left( \frac{\theta_1}{M} \right)^{n_1} \left( \frac{\theta_2}{M} \right)^{n_2} \left( \frac{\theta_3}{M} \right)^{n_3} \cdots \phi_i \] (2.47)
is invariant under the \( U(1)'s \).

The soft masses then can be written as
\[ \tilde{m}_{\phi_i}^2 = m_{\bar{\phi}_i}^2 + m_i^2 = \left[ \frac{4}{\sqrt{3} K_2} (\epsilon \hat{m})(8\pi^2 k_h) \left( \frac{\xi}{M} \right)^2 \right]^2 + \left[ \sqrt{(n_1^1 + n_2^1 + n_3^1 + \cdots) C} \right]^2. \] (2.48)

The gaugino masses are
\[ m_{1/2} = \frac{< F_S >}{2y_0} = \frac{1}{y_0} \sqrt{\frac{3}{K_2}} m_0. \] (2.49)

The trilinear soft couplings are
\[ A_{ij}^{[u,d,e]} \sim < F_S > Y_{ij}^{[u,d,e]} = A_0 m_{1/2} Y_{ij}^{[u,d,e]} \equiv a_0 Y_{ij}^{[u,d,e]}, \] (2.50)
where \( A_0 \) is a constant of order of one and \( Y_{ij}^{[u,d,e]} \) is the corresponding Yukawa coupling in the superpotential. We now consider the two different boundary conditions.

- \( m_0^2 << m_i^2 << m_{1/2}^2 \)
  
  First, notice that in (2.40), the second term inside the brackets dominates over the first for reasonable values of the parameters, so the first term can safely be neglected. Then, the ratios that are expected to be small in this limit are
  \[ \frac{m_0^2}{m_i^2} = \frac{4}{3} \left( \frac{N_c - N_f + 1}{n} \right) \frac{(\xi/M)^2}{(n_1^1 + n_2^1 + n_3^1 + \cdots)} \] (2.51)
  and
  \[ \frac{m_i^2}{m_{1/2}^2} = \frac{y_0^2 K_2 m_0^2}{3 m_0^2}. \] (2.52)

In order both these ratios be simultaneously suppressed, \( K_2 \) has to be rather small, as we argued in the introduction.

- \( m_{1/2}^2 \simeq m_0^2 >> m_i^2 \)
  
  We saw that if the conditions (1.11) and (1.12) are satisfied then the only contribution to the soft masses comes from the dilaton \( F \)-term:
  \[ m_0 = \frac{4}{\sqrt{3} K_2} (\epsilon \hat{m})(8\pi^2 k_h) \left( \frac{\xi}{M} \right)^2 \] (2.53)
with \( K_2 \) given by (1.14) and it is manifestly flavor and family universal for all fermions. The common gaugino mass and the trilinear couplings are
  \[ a_0 \simeq m_{1/2} = \frac{1}{y_0} \sqrt{\frac{3}{K_2}} m_0. \] (2.54)

8The invariants of the whole gauge group are just gauge invariant polynomials of such terms.
3 The Model

We now apply the general formalism of the previous section to an explicit model. As we mentioned before, we will use the visible sector of \( \mathbb{H} \) and introduce a slightly different hidden sector. For a detailed discussion of the phenomenological consequences related purely to the visible sector we refer the reader there. For completeness, we give the Yukawa matrices of the visible sector in the Appendix.

The fields present in the visible sector are

1. Three singlets of the non-Abelian part of the gauge group that take vevs and uniquely break the \( U(1) \)'s (called \( \theta_i \)). This sector is necessarily anomalous.
2. Three chiral families in the \( 27(16 + 10 + 1) \) of \( E_6(SO(10)) \) except the singlet of \( SO(10) \).
3. One standard-model vector like pair of Higgs weak doublets. It turns out that this is a model with \( \tan \beta \) is of order of one (see Appendix).
4. Four singlets of the non Abelian part of the gauge group \( \Sigma_k \), that do not take vevs and they are introduced to cancel the anomalies of the \( \theta_i \) fields.

The hidden sector has three (\( N_f = 3 \)) families of vector-like fields \( q_i \) and \( \bar{q}_i \), transforming as \( N_c \) and \( \bar{N}_c \) of \( G_h \), with \( N_c \) being the fundamental representation of \( G_h \). There are, in addition, a set of \( G_h \) singlet hidden fields \( T_j \), \( j = 1, ..., N_T \).

The gauge group is

\[
SU(3)_c \times SU(2)_W \times U(1)_X \times U(1)_{Y(1)} \times U(1)_{Y(2)} \times D_1 \times G^h
\]

where \( D_1 \) is a discrete symmetry acting on the visible vector-like fields in the \( 10 \) of \( SO(10) \), as

\[
E_i \rightarrow -E_i, \quad \bar{E}_i \rightarrow -\bar{E}_i, \quad D_i \rightarrow -D_i, \quad \bar{D}_i \rightarrow -\bar{D}_i
\]

and

\[
G^h \equiv SU(N_c), \quad \text{with} \quad N_c > N_f.
\]

We denote by \( V \) and \( V' \) the non-anomalous \( U(1) \)'s in \( E_6 \) (besides the regular hypercharge \( Y \)), according to

\[
E_6 \subset SO(10) \times U(1)_{V'},
\]

with

\[
27 = 16_1 + 10_{-2} + 1_4
\]

where the subscript is the \( U(1) \) value. The other \( U(1) \) is

\[
SO(10) \subset SU(5) \times U(1)_V,
\]

with

\[
16 = 5_{-3} + 10_1 + 1_5; \quad \text{and} \quad 10 = 5_2 + 5_{-2}.
\]

The charges of the visible fields in the \( 16 + 10 \) of \( SO(10) \), \( (Q, \pi, \overline{q}, L, \bar{\pi}, \overline{N}, E, \overline{E}, D, \overline{D}) \), under the three \( U(1) \)'s \( (X, Y^{(1)}, Y^{(2)}) \) are

\[
X = (-1 - \frac{3}{20} V + \frac{1}{4} V') \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix},
\]

\[
Y^{(1)} = \frac{1}{5} (2Y + V) \begin{pmatrix} 2 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & -1 \end{pmatrix},
\]

\[
Y^{(2)} = \frac{1}{4} (V + 3V') \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -1 \end{pmatrix}
\]

and hypercharge is normalized so that the triangle anomaly coefficient with one anomalous gauge field and two hypercharge (and weak) gauge bosons obey the relation:

\[
(XYY) = \frac{5}{3} (XSU(2)SU(2)).
\]
In the following, we will call $(X SU(2) SU(2)) \equiv C_W$ for simplicity.

The charges of the rest of the visible sector singlet fields are

$$\theta_i : A = \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 1 \\ 1 & -1 & 0 \end{pmatrix} \text{ and } \Sigma_i : \begin{pmatrix} -1/2 & -1/2 & 0 & 0 \\ 0 & 0 & 1/2 & -1/2 \\ -9/4 & -7/4 & 9/4 & 7/4 \end{pmatrix}, \quad (3.66)$$

where the first row contains the charges under $X$ and the second (third) row contain the charges under $Y^{(1)}(Y^{(2)})$. The matrix $A$ implies that all three $U(1)$'s break at precisely the same scale. Then, the different expansion parameters $\lambda_k$, are all equal to the same $\lambda \equiv \xi/M$.

As mentioned before, there are in addition three families of vector like hidden fields transforming under the fundamental (anti-fundamental) representation $N_c (\overline{N}_c)$ of $SU(N_c)$ with charges under the $U(1)$'s as

| $X$ | $-3$ |
| $Y^{(1)}$ | $2/N_c$ |
| $Y^{(2)}$ | $1/2$ |

where $i = 1, 2, 3$ since $N_f = 3$. This implies that $p_1 = p_2 = p_3 = 6$.

The fields $T_j$ have no charges under the non-anomalous $U(1)$'s and their charge under $X$ is $-3$. This last set of fields is given for completeness, since their only purpose here is, to adjust the gravitational anomaly to be compatible with the Green-Schwarz mechanism. For a hidden sector with $k_h = 1$, we need

$$N_T = 45 - 6N_c. \quad (3.67)$$

As it stands, this model is anomaly free. The Green-Schwarz relations

$$\frac{C_j}{k_j} = \text{constant, \ for all } j$$

are all satisfied with the Kac-Moody levels of the non-Abelian factors all equal to 1, including $k_h$. $C_j$ in the above are the non-zero anomaly coefficients associated with the different gauge factors and $k_j$ the corresponding Kac-Moody levels. We distinguish again the two limits and present for each case examples in the context of this model.

- $m_0^2 \ll m_i^2 \ll m_{1/2}^2$

For this, “no scale” case, the universal contribution to the soft masses is very small (it vanishes for all practical purposes: $m_0 \sim 0$) and the family dependent contribution from the $D$-terms is

$$m_{i}^{(Q, \mathbf{n}, \mathbf{A}, L, \mathbf{E})} = \sqrt{C} \cdot \sqrt{n_i^{(Q, \mathbf{n}, \mathbf{A}, L, \mathbf{E})}}, \quad (3.69)$$

where $n_i^{\phi}$ is the sum of the exponents defined in (2.47) for the field $\phi$ and $i$ is its family index. They are easily calculable in our model (see Appendix). The gaugino masses and the trilinear couplings are

$$x \equiv m_{1/2} \simeq a_0. \quad (3.70)$$

In table 1, we show some typical values that our model gives for the parameters $\sqrt{C}$ and $x$ with $\xi/M$ as a free parameter. Unfortunately, all these models have problems associated with charge and/or color breaking minima of the scalar potential and therefore can not be considered as viable [12] as they stand. This is a generic feature of the no scale boundary conditions. There is however the possibility of other moduli $F$-terms contributing to $m_0$, with vevs large enough to protect the vacuum [1]. In such a case, we could have a viable model of the “no scale” type. An example is, if these contributions for the $N_c = 4$ and $\xi/M = 0.24$ case (see table 1.), were $\sim 200$ GeV. The resulting low energy spectrum would then be very similar to the “minimal sugra” spectrum that we present below and we show in table 2.

\footnote{I thank S. Martin for pointing this out.}
\[ m_{1/2} \simeq m_0^2 \gg m_i^2 \]

As an example, we take \( N_c = 5 \) and \( \xi/M = 0.28 \) \(^{10}\) and use as input \( y_0 = 1.48 \) (see Appendix). For these values, the small expansion parameter is \( \epsilon \simeq 0.86 \cdot 10^{-5} \) and the condensation scale becomes \( \Lambda \simeq 2 \cdot 10^{12} \text{ GeV} \). The dilaton is now stabilized at \( K_2 \simeq 1.3 \) (see I.14). Substituting these values into (2.53) and (2.49), we obtain

\[
m_0 \simeq 200 \text{ GeV} ; \quad m_{1/2} \simeq 200 \text{ GeV} ; \quad a_0 \simeq 200 \text{ GeV}.
\] (3.71)

Of course, these parameters, are predicted at the unification scale \( M \), so we have to extrapolate their values to \( M_Z \) to obtain the low energy spectrum. In table 2, we show the MSSM parameters corresponding to this particular model \(^{13}\). It is an example of a phenomenologically viable model, with no charge/color breaking minima, consistent with EWSB and fcnc. Experimental signatures that this type of models could imply is for example a trileptonic signal in \( pp \) collisions: \( pp \to \tilde{C}_{1}^{\pm} + \tilde{N}_2 + X \to l_1 \bar{\nu}_1 \tilde{N}_1 + l_2 \nu_2 \tilde{N}_2 + X \) (three leptons and missing energy) \(^{14}\), or at the LHC, the decay \( h^0 \to \gamma \gamma \) \(^{15}\). Also, \( e^+ e^- \) annihilation could pair produce the lowest mass charginos and sfermions. The exact degeneracy of the squark masses is a result of the simultaneous conditions (1.11) and (1.12). If these are not exactly obeyed, then the squark masses split with the splittings proportional to the vevs of the \( D \)-terms times the \( U(1) \) charges of the quark fields. Soft masses with non-universal contributions from \( D \)-terms may be a more realistic scenario but in that case the mass differences, especially between the first two families should be small, in order to avoid conflict with experimental bounds on fcnc.

We finally emphasize that using at certain places \( M_{\text{string}} \) instead of \( M \equiv M_{\text{GUT}} \), can alter the numerology quite significantly but our purpose here is just to show that one can build a complete model based on \( U(1) \)'s, which is consistent with present experimental bounds.

### 4 Conclusion

We extended the dilaton dominated, (global) supersymmetry breaking mechanism by gaugino condensation in the hidden sector, communicated to the visible sector by an anomalous \( U(1) \), to the case when the gauge group contains additional non-anomalous \( U(1) \)'s. We saw that our model was capable of producing phenomenology in two different limits, depending on the choice of the hidden sector and the Kähler potential. One was a “no scale” type limit, where \( m_0 \sim 0, m_i \) small but nonzero and \( m_{1/2} = a_0 \) large. The dilaton was stabilized at a very low value of \( K_2 \) due to a weak coupling choice for \( K \). We could not find a viable model in this regime without some vacuum stabilization mechanism. The other, was a “minimal sugra” type limit. In this scenario, we did not stabilize the dilaton, but by requiring the \( D \)-term contributions to the soft masses be zero, we were able to obtain the value of the Kähler function and its derivatives in the vacuum, assuming that the dilaton gets stabilized to a reasonable value. For a particular choice of parameters, we obtained at unification scale degenerate soft masses and gaugino masses both \( \sim 200 \text{ GeV} \). We extrapolated the values of the soft parameters to \( M_Z \) and made low energy numerical predictions for this example.

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### 5 Appendix

We review some crucial facts concerning the visible sector of the model of ref. \(^{10}\).

\(^{10}\) There is nothing particularly deep in the choice of these parameters. The choice of \( N_c = 5 \) is motivated by the fact that it is probably the only value that results in reasonable squark masses and the choice \( \xi/M = 0.28 \) is the value of the expansion parameter \( \lambda_c \) that we found in ref. \(^{14}\), a number close to the Cabbibo angle (see Appendix). If we use a different value for \( \xi/M, m_0 \) and \( m_{1/2} \) will change.
The powers appearing in (3.69), are

\[ W_{Yukawa} = Y_{ij}^{[u]} Q_i \pi_j H_u + Y_{ij}^{[d]} Q_i \bar{\pi}_j H_d + Y_{ij}^{[e]} L_i \pi_j H_d + Y_{ij}^{[\nu]} L_i \bar{\pi}_j H_u + Y_{ij}^{[0]} M \pi_i \bar{\pi}_j + \cdots. \]  (5.72)

The Yukawa matrices that give the suppression to the operators are result of the breaking of the \(U(1)\)'s. They are parametrized by \(\lambda_c\), the Cabbibo angle. In the model under consideration they turn out to be in the quark sector:

\[ Y^{[u]} = \left( \begin{array}{ccc} \lambda_5 & \lambda_5 & \lambda_3 \\ 0 & \lambda_4 & \lambda_2 \\ \lambda_5 & \lambda_5 & 1 \end{array} \right), \quad Y^{[d]} = \left( \begin{array}{ccc} \lambda_5 & \lambda_3 & \lambda_3 \\ \lambda_5 & \lambda_2 & \lambda_2 \\ 1 & 1 & 1 \end{array} \right) \]  (5.73)

and in the lepton sector:

\[ Y^{[e]} = \omega \left( \begin{array}{ccc} \lambda_4 & \lambda_5 & \lambda_3 \\ 0 & \lambda_4 & \lambda_2 \\ \omega & \omega & 1 \end{array} \right), \quad Y^{[\nu]} = \omega \left( \begin{array}{ccc} \lambda_4 & \lambda_3 & \lambda_3 \\ \lambda_5 & \lambda_2 & \lambda_2 \\ \lambda_4 & \lambda_4 & 1 \end{array} \right), \quad Y^{[0]} = \omega \left( \begin{array}{ccc} \lambda_5 & \lambda_5 & \lambda_5 \\ \lambda_5 & \lambda_5 & \lambda_5 \\ 1 & 1 & 1 \end{array} \right). \]  (5.74)

One consequence of these matrices is that \(\tan \beta\) is of order of one, as one can read off from the (33) elements of \(Y^{[u]}\) and \(Y^{[d]}\).

- The vector-like matter in the 10 of \(SO(10)\) enters \(W\) with mass terms:

\[ W_{VL} = M Y_{ij}^{[E]} \bar{\pi}_i E_j + M Y_{ij}^{[D]} \bar{\pi}_i D_j + \cdots, \]  (5.75)

where

\[ Y^{[E]} = \left( \begin{array}{ccc} \lambda_2 & \lambda_7 & \lambda_{11} \\ \lambda_5 & \lambda_5 & \lambda_5 \\ \lambda_2 & \lambda_{11} & \lambda_{15} \end{array} \right) \] and \(Y^{[D]} = \left( \begin{array}{ccc} \lambda_5 & \lambda_7 & \lambda_7 \\ \lambda_5 & \lambda_5 & \lambda_5 \\ \lambda_2 & \lambda_{13} & \lambda_{15} \end{array} \right). \]  (5.76)

The mass eigenvalues of both fields are

\(\{\lambda_2^2 M, \lambda_5^2 M, \lambda_6^2 M\}. \)  (5.77)

- The powers appearing in (3.69), are

\[ n_i^Q = \left[ \frac{9}{10} + \frac{19}{10} + \frac{37}{30}, \frac{9}{10} + \frac{9}{10} + \frac{37}{30}, \frac{9}{10} - \frac{1}{10} + \frac{7}{30} \right], \]  (5.78)

\[ n_i^{\pi} = \left[ \frac{9}{10} + \frac{19}{10} + \frac{77}{30}, \frac{9}{10} + \frac{9}{10} + \frac{17}{30}, \frac{9}{10} - \frac{1}{10} - \frac{13}{30} \right], \]  (5.79)

\[ n_i^{\bar{\pi}} = \left[ \frac{3}{10} + \frac{3}{10} + \frac{29}{30}, \frac{3}{10} + \frac{3}{10} - \frac{1}{30}, \frac{3}{10} + \frac{3}{10} + \frac{1}{30} \right], \]  (5.80)

\[ n_i^{L} = \left[ \frac{3}{10} + \frac{3}{10} + \frac{23}{30}, \frac{3}{10} + \frac{3}{10} - \frac{1}{30}, \frac{3}{10} + \frac{3}{10} - \frac{7}{30} \right], \]  (5.81)

\[ n_i^{\bar{\nu}} = \left[ \frac{9}{10} + \frac{19}{10} - \frac{1}{10}, \frac{9}{10} + \frac{9}{10} + \frac{19}{10}, \frac{9}{10} - \frac{1}{10} + \frac{9}{10} \right], \]  (5.82)

\[ n_i^{H_u} = \left[ \frac{-9}{5} + \frac{1}{5} + \frac{1}{5} \right], \]  (5.83)

\[ n_i^{H_d} = \left[ \frac{9}{5} - \frac{1}{5} - \frac{1}{5} \right]. \]  (5.84)

- To calculate the unification scale and the value of the gauge coupling at unification, we have to take into account the three new thresholds due to vector-like matter. By doing so, we find

\[ \frac{1}{g(M)^2} = 1.48 \]  (5.85)

and \(M = 3.4 \cdot 10^{16} GeV\). This is the \(M\) that is used consistently as our only mass parameter, even though we are aware that this might not be the most appropriate in some cases.
The value of the parameter $\xi/M$ is undetermined, as long as the Kähler potential is unknown. Since however we use it as the expansion parameter in the mass matrices, it is implicitly assumed to be a number close to the Cabbibo angle ($\sim 0.22$). Assuming, for example, for $K$ its usual tree level form $K = -\log (2y)$ and using $M$ and $g(M)$ from above, we obtain $\xi/M = 0.28$.

The $\mu$ term has zero charge under all $U(1)$’s. It could therefore appear in $W$ and in the Kähler potential, but since one does not get pure mass terms in a string spectrum, we assume its presence only in the Kähler potential. Therefore, after supersymmetry breaking, at low energy, a $\mu$ term of the correct order (few hundred GeV) will be generated by the Guidice-Masiero mechanism.

There is no kinetic mixing between the $U(1)$’s.

R-parity is exactly conserved.

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Table 1: Values for $x$ and $\sqrt{C}$ at the high scale $M$, for different choices of $N_c$ and $\xi/M$, in the “no scale” regime.

| $N_c$ = 5 | $\xi/M$ | $x$(GeV) | $\sqrt{C}$(GeV) | $N_c$ = 4 | $\xi/M$ | $x$(GeV) | $\sqrt{C}$(GeV) |
|-----------|--------|----------|-----------------|-----------|--------|----------|-----------------|
| 0.10      | 60     | 4        |                 | 0.22      | 60     | 1        |                 |
| 0.12      | 435    | 22       |                 | 0.24      | 190    | 3        |                 |
| 0.125     | 680    | 33       |                 | 0.26      | 565    | 8        |                 |
| 0.13      | 1035   | 48       |                 | 0.28      | 1533   | 20       |                 |

Table 2: The first column contains the name of the parameter. The second column contains the low energy value of the parameter for a “minimal sugra” type model corresponding to $N_c = 5$, $\xi/M = 0.28$ and $y_0 = 1.48$. The input to the RGE’s at $M$, is $m_0 = m_{1/2} = a_0 = 200$ GeV, $\tan\beta(M_Z) = 4$, $sgn(\mu) = +1$ and $m_t = 175$ GeV. The values of $(\tilde{B}, \tilde{W}, \tilde{g})$ are their lowest order pole masses.

| Parameter | Value at $M_Z$ (GeV if a mass) |
|-----------|--------------------------------|
| $(M_Z, v_{Higgs})$ | (90.4, 174.1) |
| $\tan\beta$ | 4 |
| $(\tilde{B}, \tilde{W}, \tilde{g})$ | (62, 122, 367) |
| $(\alpha_3, \alpha_2, \alpha_1, \sin\theta_W)$ | (0.116, 0.033, 0.0165, 0.232) |
| $(\tilde{Y}^{[u]}_{33}, \tilde{Y}^{[d]}_{33}, \tilde{Y}^{[e]}_{33})$ | (1, 0.08, 0.042) |
| $(B, \mu)$ | (125, 233) |
| $(\tilde{u}_L, \tilde{u}_R, \tilde{d}_L, \tilde{d}_R, \tilde{c}_L, \tilde{c}_R, \tilde{\nu}_e, L)$ | (407, 399, 414, 394, 428, 216, 226) |
| $(\tilde{c}_L, \tilde{c}_R, \tilde{s}_L, \tilde{s}_R, \tilde{\mu}_L, \tilde{\mu}_R, \tilde{\nu}_\mu, L)$ | (407, 399, 414, 394, 428, 216, 226) |
| $(\tilde{t}_L, \tilde{t}_R, \tilde{b}_L, \tilde{b}_R, \tilde{\tau}_L, \tilde{\tau}_R, \tilde{\nu}_\tau, L)$ | (435, 283, 395, 364, 238, 215, 226) |
| $(h^0, H^0, A^0, H^\pm)$ | (104, 346, 343, 352) |
| $(\tilde{C}_1, \tilde{C}_2)$ | (95, 270) |
| $(\tilde{N}_1, \tilde{N}_2, \tilde{N}_3, \tilde{N}_4)$ | (53, 99, −240, 272) |
| $LSP \rightarrow (\tilde{N}_1)$ | 53 |