Z' Gauge Models from Strings

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Potentially realistic string models often contain additional abelian gauge factors besides the standard model group. The consequences of such extended gauge structure are manifold both for theory and phenomenology as I show focussing in the simplest case of just one additional non-anomalous $U(1)'$. First, I discuss the possible symmetry breaking patterns according to the scale at which the $U(1)'$ symmetry gets broken: in the first case, that scale must be below $\sim 1$ TeV to avoid fine-tuning; in the second case, the breaking can take place along a flat direction at an intermediate scale between the string scale and the electroweak scale. In both cases, I present a number of the generic implications expected, e.g. for the $\mu$ problem, $Z'$ and Higgs physics, dark matter and fermion masses.

1. INTRODUCTION

Many string models predict a gauge group at the string scale which contains, besides the standard model group $G_{SM} = SU(3)_C \times SU(2)_L \times U(1)_Y$, extra $U(1)$ gauge factors. In this talk, I will examine the generic phenomenological and theoretical implications of this type of scenarios. For simplicity, I limit the analysis to one extra non-anomalous $U(1)'$. This is the simplest gauge extension of the SM one can consider and the phenomenology of particular $Z'$ models has been studied extensively. Here, no attempt is made at building a concrete realistic string model but rather at exploring the generic implications of string inspired $Z'$ models. As we shall see, such simple extended models can have manifold implications for, among others, the $\mu$-problem, Higgs physics, exotic matter, fermion masses, dark matter, the spectrum of superpartners and cosmology. I will describe some of these implications in two different scenarios which can be distinguished by the scale of $U(1)'$ breaking. For more details the reader is directed to refs. [1,2].

Let me be more specific about the string input in the analysis. I will consider a class of specific string vacua: orbifold models [4] with Wilson lines, in particular, models based on the free-fermionic construction (at special points of moduli space) [7-10]. I then assume:

- N=1 Supersymmetry
- The (visible sector) gauge group at $M_{string}(\simeq 5 \times 10^{17}$ GeV) is $G_{SM} \times U(1)'$, where the extra abelian factor has no anomaly.
- The matter content is that of the Minimal Supersymmetric Standard Model (MSSM) plus some number of SM singlets, $S_i$, which carry $U(1)'$ charges $Q_i$, plus some exotics (generally transforming non-trivially under $G_{SM}$). Both singlets and exotics appear generically in string models and can play an important role in breaking the $U(1)'$ symmetry.
- The gauge couplings $g_3, g_2, g_1, g_1'$ (for $SU(3)_C, SU(2)_L, U(1)_Y$ and $U(1)'$ respectively, with proper normalization for $g_1$ and $g_1'$) unify at $M_{string}$. This unification works only at the $10-15\%$ level in the simplest models (in which the couplings meet at a scale below $M_{string}$). The small numerical inconsistency that results does not affect the generic predictions of these models that I will present.

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• Supersymmetry breaking, the main uncertainty in extracting phenomenological consequences from string models, is parameterized by soft masses for scalars and gauginos and by trilinear masses associated to superpotential terms. Naturalness requires these masses (which in general are \( m_{soft} \)) to be of the order of the electroweak scale.

• No bilinear mass terms are present in the superpotential and the Yukawa couplings at \( M_{\text{string}} \) are either zero or related to the unified gauge coupling \( g_0 \) by \( \sqrt{2} \) factors (typically \( \sim g_0 \sqrt{2} \), \( g_0 \), or \( g_0 / \sqrt{2} \)) in the fermionic \( Z_2 \times Z_2 \) orbifold constructions).

For fixed values of the parameters at the string scale one can address, by a renormalization group analysis (with the chosen parameters as boundary conditions) what are the low-energy implications typical of these scenarios.

2. RADIATIVE BREAKING OF \( U(1)' \)

To have an acceptable model, the extra \( U(1)' \) gauge factor has to be broken at low energies. Generically, this breaking is achieved radiatively (in a way which parallels electroweak radiative breaking in the MSSM [10]) when the soft mass squared \( m_S^2 \) of some singlet \( S \) [charged under \( U(1)' \)] is driven to negative values in the infrared (an alternate breaking mechanism is discussed in the next section). This causes an instability of the potential for \( \langle S \rangle = 0 \), and \( S \) acquires a non-zero vacuum expectation value (VEV) thus breaking the \( U(1)' \). The value of \( \langle S \rangle \), and thus the breaking scale, is fixed by quartic \( D \) terms [from \( U(1)' \)] in the potential for the field \( S \) (alternatively there may be also quartic \( F \)-terms),

\[
V(S) = m_S^2 S^2 + \frac{1}{2} g_1^2 Q_S^2 S^4,
\]

(1)
giving \( g' \langle S \rangle \sim |m_S| \).

If the two Higgs doublets \( H_1, H_2 \) are charged under \( U(1)' \) (the interesting case) electroweak breaking is influenced at tree-level by the breaking of \( U(1)' \). In order to have a natural electroweak breaking scale (or equivalently, to avoid fine tuning the soft masses), the \( U(1)' \) breaking scale cannot be much larger than \( O(1) \) TeV.

Thus, we see that the \( Z' \) models derived from strings naturally predict \( M_{Z'} \lesssim O(1) \) TeV.

One exception to this generic result [besides the trivial one of having the \( U(1)' \) decoupled from the electroweak Higgs sector] occurs if the \( U(1)' \) is broken radiatively along a flat direction. For example, if the model contains two SM singlets \( S_{1,2} \) with opposite \( U(1)' \) charges, the field direction \( S \equiv S_1 = S_2 \) is \( D \)-flat and (assuming also \( F \)-flatness) the tree-level potential \( V(S) = -m_S^2 S^2 \) has no quartic coupling and is unstable. Eventually, the potential gets stabilized at very large values of the field \( S \) due to radiative corrections [the most important of which are encoded in the running mass \( m_S^2(S) \)] or by non-renormalizable operators \( (NROs) \) which can no longer be neglected for such large field values. In both cases the \( U(1)' \) breaking scale is much larger than the electroweak scale (e.g. \( 10^{12} \) GeV). This is one example of an intermediate scale being generated in string models [4]. Still, these scenarios have also many implications for the low-energy effective theory. In the rest of the talk I will discuss separately the two general classes of scenarios just described: I) \( Z' \) below the TeV scale [2] and II) intermediate scale \( Z' \) [3]. For other related discussions in the literature see [3].

3. \( Z' \) BELOW THE TeV SCALE

Let us consider the simplest case, which requires the presence of just one singlet \( S \) to give a correct symmetry breaking pattern. We further assume that the \( U(1)' \) charge assignments of \( S \) and the two Higgs doublets \( H_{1,2} \) are such that the superpotential contains a term

\[
W = h_s S H_1 \cdot H_2,
\]

(2)

and thus \( Q_S + Q_1 + Q_2 = 0 \). After symmetry breaking, \( \langle S \rangle = s / \sqrt{2} \), \( \langle H_1^0 \rangle = v_1 / \sqrt{2} \) and \( \langle H_2^0 \rangle = v_2 / \sqrt{2} \), the gauge group gets broken down to \( U(1)_{em} \) with

\[
M_W^2 = \frac{1}{4} g_2^2 (v_1^2 + v_2^2),
\]

(3)

from which \( v^2 = v_1^2 + v_2^2 = (246 \text{ GeV})^2 \), and

\[
M^2_{Z-Z'} = \left( \begin{array}{cc}
\frac{M_Z^2}{\Delta^2} & \Delta^2 \\
\Delta^2 & M_{Z'}^2
\end{array} \right),
\]

(4)
where
\[ M_Z^2 = \frac{1}{4} G^2 (v_1^2 + v_2^2), \]  
\[ M_Z'^2 = g_1^2 (v_1^2 Q_1^2 + v_2^2 Q_2^2 + s^2 Q_3^2), \]  
\[ \Delta^2 = \frac{1}{2} g_1 G (v_1^2 Q_1 - v_2^2 Q_2). \]  

Here we use \( G^2 = g_2^2 + g_Y^2 = g_2^2 + 3g_1^2/5. \)

From Eqs. (4) and (7) we see that it is possible to accommodate a very small \( Z - Z' \) mixing angle (the experimental upper bound is \( O(10^{-3}) \) in typical models \([14]\)) due to a cancellation in the off-diagonal term \( \Delta^2 \) for a particular value of \( \tan \beta \equiv v_2/v_1 \approx \sqrt{|Q_1/Q_2|} \). The presence of the singlet VEV \( s \), on the other hand, helps in giving to \( Z' \) a sufficiently large mass to evade collider constraints \([19]\). In ref. \([2]\) two phenomenologically viable scenarios were identified:

- **(a) A-driven scenario.** Symmetry breaking is driven by a large value of the trilinear soft mass associated to the superpotential term \([2]\) [this is similar to \( SU(3)_C \) color breaking in the MSSM for large values of \( A_Q \), the trilinear coupling associated to the quark Yukawa couplings], that is, \( A \gg M_{1/2}, m_0 \), where \( M_{1/2} \) represents a typical gaugino mass and \( m_0 \) a scalar soft mass. In this scenario, the minimum of the potential has \( v_1 \approx v_2 \approx s \). This is shown in figure 1, where the adimensional ratio \( f_i = v_i/m_i \) is plotted as a function of \( c_A = A/m_0 \) in some particular example. For \( Q_1 = Q_2 \), values of \( \tan \beta \) close to 1 would give a small enough mixing angle.

- **(b) Large \( s \) scenario.** \( (s \gg v_1, v_2) \)

This case corresponds to a hierarchy in the masses \( M_Z' \gg M_Z \). The electroweak scale, as represented by \( M_Z \), is light due to accidental cancellations among soft masses, which are typically of the order \( M_{Z'} \). To avoid fine-tuning, \( M_{Z'} \) cannot be much larger than \( O(1 \text{ TeV}) \) as we already discussed before.

### 3.1. Soft Masses at \( M_{\text{string}} \)

One can then proceed to determine what pattern of soft supersymmetry breaking terms at the string scale will lead, via RG evolution, to the two low-energy acceptable scenarios described previously.

With the minimal particle content (MSSM + singlet) it is not possible to achieve the correct symmetry breaking if universal boundary conditions for the soft masses at \( M_{\text{string}} \) are used. Both scenarios, (a) and (b), can however be reached with non-universal boundary conditions. The allowed parameter space is somewhat constrained and typically corresponds to gaugino masses much smaller than the scalar masses.

It is easier to have an acceptable symmetry breaking pattern in the presence of exotics, in particular if the singlet \( S \) has couplings to them of the form

\[ W = h_E S E_1 E_2. \]  

The large value of the Yukawa coupling \( h_E \) (~...
The longitudinal components of symmetry breaking, 4 d.o.f. are eaten by the total of 10 degrees of freedom. After spontaneous doublets \( H_3 \). Higgs Physics for LEP II searches [20].

In the presence of the superpotential term \( h_s S H_1 \cdot H_2 \), the \( \mu \) term of the MSSM Higgs potential is generated dynamically by the VEV of the singlet \( S \) as

\[
\mu_{eff} = h_s \langle S \rangle.
\]

As discussed before, \( \langle S \rangle \) lies at the electroweak scale, so that \( \mu_{eff} \) is of the correct order of magnitude to give a proper electroweak breaking [the same mechanism generates a \( B \) term \( \sim h_s \langle S \rangle A \sim m_{soft} \)]. This solves the \( \mu \) problem \([15]\) of the MSSM (where \( \mu \) is a SUSY parameter which does not know about the electroweak scale and has to be set of that order by hand).

This solution to the \( \mu \) problem is similar in spirit to the mechanism in the Next to Minimal Supersymmetric Standard Model (NMSSM) \([16]\) in which the MSSM is extended by a chiral singlet but no additional \( U(1) \) gauge symmetry is present. The advantage of having a \( U(1)^{\prime} \) present is that there is no cosmological problem with the formation of domain walls after symmetry breaking: the discrete \( Z_3 \) symmetry of the NMSSM responsible for the domains is embedded in \( U(1)^{\prime} \).

This solution is complementary to the Giudice-Masiero mechanism \([14]\), in the sense that the \( U(1)^{\prime} \) symmetry would forbid a \( H_1 H_2 \) term in the Kähler potential, which is in the basis of that mechanism.

The two scenarios (a) and (b) would correspond to a small and large \( \mu_{eff} \) respectively. In the first scenario (\( A \) dominated) light charginos and neutralinos are expected, with obvious implications for LEP II searches \([21]\).

3.3. Higgs Physics

The (CP conserving) Higgs sector contains two doublets \( H_1, H_2 \) and one singlet \( S \) making up a total of 10 degrees of freedom. After spontaneous symmetry breaking, 4 d.o.f. are eaten by the longitudinal components of \( Z, Z' \) and \( W^{\pm} \). The 6 d.o.f. left appear as physical Higgs bosons: three (CP even) scalars \( H_0^0 \), one (CP odd) pseudoscalar \( A^0 \) and one charged pair \( H^{\pm} \). That is, the model contains just one scalar more than in the MSSM.

The pattern of Higgs masses is very different in the two scenarios (a) and (b):

In the \( A \)-dominated scenario, all Higgs masses are controlled by one single mass scale \( v \) and three adimensional parameters: \( h_s, g_1^2 \) and \( Q_1 \). The spectrum is therefore light, although all Higgses can still evade detection at LEP II. One typical example, for \( h_s = 0.7, g_1^2 = (5/3)G^2 \sin^2 \theta_W \) and \( Q_1 = -1 \), has \( m_{H^{\pm}} = 146 \) GeV, \( m_{A^0} = 211 \) GeV, \( m_{H_3^0} = 230 \) GeV, \( m_{H_2^0} = 152 \) GeV and \( m_{H_1^0} = 122 \) GeV (tree-level masses).

In the large \( s \) scenario there are three mass scales in the Higgs spectrum. Typically \( m_{H_1^0} \sim M_{Z'} \) (with \( H_3^0 \) singlet dominated), \( m_{A^0} \gg M_Z \), with \( (H_2^0, A^0; H^\pm) \) forming a nearly degenerate doublet decoupled from electroweak symmetry breaking, and \( H_1^0 \) remains at the electroweak scale.

An upper bound on the tree-level mass of \( H_1^0 \) can be easily obtained as

\[
m_{H_1^0}^2 \leq M_Z^2 \cos^2 2\beta + \frac{1}{2} h_s^2 v^2 \sin^2 2\beta + g_1^2 \sqrt{Q_H} v^2 (10)
\]

with \( Q_H = Q_1 \cos^2 \beta + Q_2 \sin^2 \beta \). The first term in Eq. \((10)\), coming from \( SU(2)_L \times U(1)_Y \) D-terms, is the MSSM piece. The second, from F-terms, appears also in the NMSSM but the third piece, from \( U(1)^{\prime} \) D-terms, is a particular feature of this type of models \([21, 23]\). More refined bounds can be obtained under extra assumptions for the parameters of the model (see \([2]\)). This tree-level bound can be very large; 170 GeV is a typical example (it is usually forgotten that the bound \( \sim 150 \) GeV for any perturbative SUSY model \([23]\) applies only if the gauge sector is not extended) so that \( H_1^0 \) can easily evade detection at LEP II.

3.4. Superpartner Spectrum

The spectrum of superpartners is also typically light for the \( A \)-dominated scenario (with all particles expected to lie at the electroweak breaking scale) while in the large \( s \) scenario soft masses are generically larger, giving a decoupled SUSY mass
pattern.

There is, however, an extra effect associated with \(U(1)'\) D-term contributions to the masses. The squared mass of any scalar \(\phi_i\) (with \(U(1)'\) charge \(Q_i\)) is shifted by the amount

\[
\delta m_i^2 = \frac{1}{2} g_i^2 Q_i (Q_1 v_1^2 + Q_2 v_2^2 + Q_S s^2).
\]

(11)

This can be a sizable effect that can cause some superpartner to be lighter. In general, the mass sum-rules that can be derived in particular models will be violated by these contributions, leaving a characteristic imprint in the mass spectrum \[24\].

3.5. Dark Matter

The neutralino sector in these models is composed of \(\tilde{B}, \tilde{W}_3, \tilde{H}_1^0, \tilde{H}_2^0\) plus \(\tilde{B}'\) and \(\tilde{S}\). In most of the parameter space the LSP (which is the natural candidate for dark matter when \(R\)-parity is conserved) is the lightest neutralino with composition dominated by \(\tilde{B}\). However, in some regions of parameter space (with large \(s\) and large gaugino masses) the lightest neutralino is \(\tilde{S}\)-dominated making an atypical dark matter candidate \([25]\) with interesting rates for direct detection (\(\sim 10^{-2}\) events/kg/day in \(73\)Ge detectors).

4. INTERMEDIATE SCALE \(Z'\)

In this case, at least two SM singlets, \(S_1, S_2\), with \(Q_1 Q_2 < 0\) are required. Calling \(S\) the flat direction determined by the condition \(|Q_1| S_1^2 = |Q_2| S_2^2\) (we assume also \(F\)-flatness along \(S\)) the renormalization-group improved potential for \(S\) is simply

\[
V(S) = m_S^2 S S^2.
\]

(12)

If the couplings of \(S\) are such that \(m_S^2 < 0\) in the infrared (we assume \(m_S^2 > 0\) at the string scale) the shape of the potential is as depicted in figure 2 (solid line labeled RAD). The potential tracks the running of \(m_S^2\) (dashed line) and is stabilized when \(m_S^2\) turns positive at the scale \(\mu_{RAD}\). The minimum of the potential occurs at some intermediate scale \(M_I \lesssim \mu_{RAD} \gg M_Z\).

![Figure 2. Effective potential (solid lines) along the flat direction: 1. (marked RAD) in the case of stabilization by the running of \(m_S^2\) (dashed line) and 2. (marked NRO) in the case of stabilization by non-renormalizable operators below the scale \(\mu_{RAD}\).](image)

If \(\mu_{RAD}\) is very large, non-renormalizable operators like

\[
W_{NRO} \sim \alpha_K \left(\frac{S}{M}\right)^K S^3,
\]

(13)

which are generically present after integrating out physics at the string scale (\(M \approx M_{string}\)) cannot be neglected in the determination of the minimum. In \([13]\), we write the NRO term, originally given by powers of \(S_1\) and/or \(S_2\), projected along the flat direction \(S\). These terms contribute to the potential a positive-definite term

\[
\delta V(S) \sim \left(\frac{S^{2+K}}{M^K}\right)^2,
\]

(14)

that lifts the flat direction and can stabilize the potential below \(\mu_{RAD}\). The shape of the potential
4.2. Low-Energy Effective Theory

In this case is represented in figure 2 by the solid line labeled NRO.

We can distinguish then two different classes of models according to the mechanism for stabilization of the potential:

- **1.** Stabilization by pure running of $m_S^2$ (typical case for $\mu_{RAD} \lesssim 10^{12}$ GeV). The minimum is fixed at $\langle S \rangle \sim \mu_{RAD}$.

- **2.** Stabilization by NROs (typical case for $\mu_{RAD} \gtrsim 10^{12}$ GeV). The minimum is fixed at

$$\langle S \rangle \sim \mu_K \equiv (m_{soft} M^K)^{1/3} \lesssim \mu_{RAD}. \quad (15)$$

### 4.3. The $\mu$ parameter

In these models, the $\mu$ parameter can be generated dynamically by the VEV of $S$ through non-renormalizable operators in the superpotential that couple the MSSM Higgs doublets to the flat direction $\tilde{S}$:

$$W = hE S_i E_1 E_2. \quad (16)$$

In different cases (with different quantum numbers and multiplicities for the exotic fields $E_i$) were studied. It was found that a correct breaking can be obtained even in the case of universal soft masses at the string scale (provided the multiplicity of the exotics is large enough).

By varying the parameters of the model the scale $\mu_{RAD}$ can be moved in the wide range $10^5$ GeV to $10^{16}$ GeV.

### 4.4. Low-Energy Effective Theory

In both cases, 1 and 2, the $U(1)'$ breaking takes place at an intermediate scale (e.g. $10^{12}$ GeV) much larger than the electroweak scale. The $Z'$ boson (and the exotics coupled to the flat direction) is thus very heavy (this can have interesting effects for gauge coupling unification [20]). The effective theory at the TeV scale is just the MSSM, but with an extra chiral multiplet $\hat{S}$, remnant of $U(1)'$ breaking. However, the interactions of $\hat{S}$ with MSSM fields are suppressed by powers of the intermediate scale. Nevertheless, these scenarios are interesting because many parameters of the MSSM can be determined by intermediate scale physics.

For example, the mass spectrum of superpartners is affected by the intermediate scale $U(1)'$ breaking. At the renormalizable level, the only interaction between the MSSM fields and the intermediate scale fields arises from the $U(1)'$ $D$-terms in the scalar potential. The resulting effect after integrating out the fields which have heavy intermediate scale masses [24] is a shift of the soft masses of MSSM fields charged under the extra $U(1)'$:

$$\delta m_i^2 = -Q_i \frac{m_1^2 - m_2^2}{Q_1 - Q_2}. \quad (18)$$

### 4.5. Radiative Breaking

Both scenarios, 1 and 2, are based on radiative breaking of the $U(1)'$ symmetry along the flat direction $S$. The precise condition that has to be satisfied at the electroweak scale to ensure the breaking (the $m_S^2 < 0$ condition of the previous subsection) is

$$\frac{m_1^2}{|Q_1|} + \frac{m_2^2}{|Q_2|} < 0, \quad (19)$$

where $m_1^2, m_2^2$ are the (electroweak scale) soft masses of $S_1$ and $S_2$ respectively. From Eq. (19) we see that $m_1^2 < 0$ and/or $m_2^2 < 0$ are required (and a hierarchy in the charges can make [10] easier to fulfill).

This breaking can be easily achieved when $S_1$ (and/or $S_2$) couples to exotics in a superpotential term

$$W = hE S_i E_1 E_2. \quad (20)$$

In different cases (with different quantum numbers and multiplicities for the exotic fields $E_i$) were studied. It was found that a correct breaking can be obtained even in the case of universal soft masses at the string scale (provided the multiplicity of the exotics is large enough).

By varying the parameters of the model the scale $\mu_{RAD}$ can be moved in the wide range $10^5$ GeV to $10^{16}$ GeV.
2. Stabilization by NROs. Eq. (21) is exactly of the form expected for $\mu_K$ [see Eq. (15)], the scale of the minimum in these scenarios. In other words, for $K$ in Eq. (13) equal to $P$ in Eq. (19), one gets automatically $\mu_{\text{eff}} \sim m_{\text{soft}}$. The same applies to the $B$-term in the MSSM Higgs potential, which can be generated by $\langle S \rangle$ in a similar way.

4.4. Non-Renormalizable Operators

As we have seen, NROs play a crucial role in these intermediate scale models. Interestingly, NROs are calculable in a large class of string models [27,28]. Writing

$$W = \alpha_K \left( \frac{S}{MP'} \right)^K \equiv \left( \frac{S}{M_K} \right)^K,$$ (22)

the coefficient $\alpha_K$ is basically given by a complicated world-sheet integral and some Clebsch-Gordan coefficients. Reabsorbing $\alpha_K$ in the mass $M_K$, the general tendency is that $M_K$ increases with $K$ so that the weight of higher order terms is further suppressed [27,28,3].

It is also important to determine which NROs will be present, that is, what values of the powers $P'$ (for the $\mu$ parameter) and $K$ (to fix the minimum) will be allowed. The values of these powers are constrained by

- $U(1)'$ gauge invariance. For example, if $Q_1 = 4/5$ and $Q_2 = -1/5$,

$$W \sim (S_1 S_2)^{P'} \sim S^{5n},$$ (23)

so that $K = 5n$.

- World-sheet selection rules, e.g. in free-fermionic models, can forbid terms otherwise consistent with all gauge symmetries (see, for example, [28,29]). So, one should keep in mind that not all the NROs satisfying all gauge symmetry constraints will be necessarily present in the superpotential. This is a typical stringy feature which is in contrast with the field theory lore.

4.5. Fermion Masses

The large value of string Yukawa couplings offers a natural explanation for the large value of the top mass [18]. The light fermions would have no Yukawa couplings in the superpotential and would therefore be massless, which is a good first order approximation. Models with an intermediate scale provide a mechanism to give small non-zero masses to these fermions using NROs like

$$W_{u_i} \sim H_2 \cdot Q_i U_i^c \left( \frac{S}{M} \right)^{P'_u},$$

$$W_{d_i} \sim H_1 \cdot Q_i D_i^c \left( \frac{S}{M} \right)^{P'_d},$$

$$W_{e_i} \sim H_1 \cdot L_i E_i^c \left( \frac{S}{M} \right)^{P'_e},$$ (24)

where $i$ is a generation index and I use standard notation for the MSSM chiral supermultiplets. Small Yukawa couplings are generated after $U(1)'$ symmetry breaking:

$$Y \sim \left( \frac{\langle S \rangle}{M} \right)^{P''}.$$ (25)

Limiting the discussion to scenarios with NRO stabilization, which give more definite predictions, it was shown in [3] how the family pattern of quark and lepton masses [30] can be qualitatively reproduced for particular choices of the powers $P'$ in Eq. (23) and $K$ in Eq. (13). In particular, as shown in Table 1, $K = 5, 6$ in combination with $P' = 1$ for the second family and $P' = 2$ for the first are satisfactorily close to the observed pattern. The mass of the light fermions of the third family do not fit quite as well and call for some other mechanism. However, given the roughness of the estimates and the simplicity of the model, the overall pattern of masses is quite encouraging.

There are also interesting implications for neutrino masses [31]. A very light (non-seesaw) doublet neutrino Majorana mass is possible from the superpotential

$$W \sim \frac{H_2 \cdot L_i}{M} \left( \frac{S}{M} \right)^{P''_{L_i}}.$$ (26)
can give a Majorana mass to the singlet neutrinos $\nu^c$. Experiments and dark matter.

The superpotential $W$ of the standard seesaw mechanism, with mass $m_{\nu^c}$ gives masses heavier than those coming from Eq. (26). The superpotential $W$ is

$$W_D \sim H_2 \cdot L_i \nu^c \left( \frac{S}{M} \right) P_{L_i \nu^c}^\prime,$$

$$W_M \sim \nu^c \nu^c S \left( \frac{S}{M} \right) P_{\nu^c \nu^c}^\prime.$$

The superpotential $W_D$ can yield Dirac neutrino masses heavier than those coming from Eq. (26). For example, taking $K = 5$ and $P_{L_i \nu^c}^\prime = 4$ or $P_{L_i \nu^c}^\prime = 5$ gives masses $m_{\nu^c \nu^c} = 0.9 \text{ eV}$ or $10^{-2} \text{ eV}$, respectively, which are in the interesting range for solar and atmospheric neutrinos, oscillation experiments and dark matter.

The superpotential $W_M$, on the other hand, can give a Majorana mass to the singlet neutrinos $\nu^c$ which can be very large or small, depending on the sign of $P_{\nu^c \nu^c}^\prime - K$. Combining the effects of $W_M$ and $W_D$, light neutrinos can be produced by the standard seesaw mechanism, with mass

$$m_{\text{seesaw}}^\text{light} \sim m_{\nu^c \nu^c}^2 / m_{\nu^c \nu^c}^\prime$$

$$\sim m_{\text{soft}} \left( \frac{m_{\text{soft}}}{M} \right)^{2 P_{L_i \nu^c}^\prime + K - P_{\nu^c \nu^c}^\prime}.$$  

With the choice $K = 5$ and $P_{L_i \nu^c}^\prime = 2, 1$ for $i = 1, 2$, respectively and with either $P_{L_3 \nu^c}^\prime = 1$ or $P_{L_3 \nu^c}^\prime = P_{u_3}^\prime = 0$ (involving a renormalizable Dirac neutrino term) the light eigenvalues of the three generations fall into the hierarchy of $3 \times 10^{-5} \text{ eV}, 1 \times 10^{-2} \text{ eV}$, and either $1 \times 10^{-2} \text{ eV}$ or $5 \text{ eV}$ for $P_{\nu^c \nu^c}^\prime = P_{L_i \nu^c}^\prime + 1$. This range is again of interest for laboratory and non-accelerator experiments.

5. SUMMARY AND CONCLUSIONS

A large class of potentially realistic string models predict a gauge group of the observable sector with additional abelian factors besides the standard group. In this talk I have discussed some phenomenological and theoretical implications of this type of models concentrating for simplicity in the simplest case, with just one additional (non-anomalous) $U(1)^\prime$. This modest extension already has a plethora of interesting consequences, some of which I have described. The goal has been to explore the generic features of these models rather than to construct a realistic string model. In this respect one can divide the scenarios in two different classes according to the scale at which the $U(1)^\prime$ symmetry gets broken.

Generically, $U(1)^\prime$ breaking takes place below the TeV scale much in the same way that electroweak breaking occurs. Actually, when the two MSSM Higgs doublets that break $SU(2)_L \times U(1)_Y$
are charged under $U(1)'$ (which is the case of interest) the breaking of $U(1)'$ is linked to electroweak breaking. To avoid fine-tuning the parameters of the model, the scale of $U(1)'$ breaking cannot be much larger than the Fermi scale. Thus, the new $Z'$ boson is expected to be on the reach of the next generation of colliders. An acceptable symmetry breaking pattern requires the presence of at least one standard model singlet $S$ charged under the $U(1)'$ and taking a VEV. If the coupling $SH_1 \cdot H_2$ is present in the superpotential, the $\mu$ parameter is generated dynamically and with the correct order of magnitude.

In this setting, two main possibilities exist for a phenomenologically viable scenario: a) symmetry breaking driven by the trilinear soft term associated to $SH_1 \cdot H_2$ and b) radiative breaking with a singlet VEV large compared to the electroweak scale (by say, one order of magnitude). Both cases can give an acceptably small scale (by say, one order of magnitude). Both cases are not only well motivated from the string perspective but have a remarkably large number of potentially interesting implications, both phenomenologically and theoretically.

An exception to the generic case described above can occur if the $U(1)'$ is broken along a flat direction. Such flat directions exists if at least two SM singlets have $U(1)'$ charges of opposite signs ($F$-flatness is also required). In that case, the field along the flat direction, $S$, can take a very large VEV. This is what happens if the potential $V(S)$ gets destabilized at the origin by a negative soft mass squared $m_S^2$ for $S$ in the infrared. Two main classes of scenarios are possible depending on the mechanism that stabilizes the potential at large $S$ and fixes the VEV [and thus the scale of $U(1)'$ breaking]: in one, the potential is stabilized by $m_S^2$ running to positive values in the ultraviolet; in the second, stabilization is due to nonrenormalizable operators. This latter case is more predictive as the scale involved is fixed mainly by the (discrete) order of the NROs. In both cases the low-energy effective theory is simply the MSSM but many of its parameters are influenced by the intermediate scale breaking of the $U(1)'$. There are, for example, effects on the soft masses for superpartners, which receive new $U(1)'$ contributions. In addition, NROs coupling MSSM fields to the flat direction can offer an explanation for the origin of the $\mu$ parameter, as well as for the hierarchy of the observed fermion masses (and there are also interesting implications for neutrino masses).

In conclusion, it is worth pursuing the detailed study of such extensions of the minimal supersymmetric standard model. These simple models are not only well motivated from the string perspective but have a remarkably large number of potentially interesting implications, both phenomenologically and theoretically.

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