Many extensions of the standard model, especially grand unified theories and superstring models, predict the existence of additional $Z'$ bosons and associated exotic chiral supermultiplets. It has recently been argued that for classes of string motivated models with supergravity mediated supersymmetry breaking there are two scenarios for the additional $Z'$s: either the mass is in the accessible range $< O(1 \, \text{TeV})$, providing a natural solution to the $\mu$ problem and implications for the Higgs and sparticle masses and for the LSP; or, when the breaking is associated with a $D$-flat direction, at an intermediate scale, providing a possible explanation for the hierarchies of quark and charged lepton masses and new possibilities for neutrino masses. Related work, examining the detailed structure of specific perturbative string vacua for $D$ and $F$-flat directions, surviving $U(1)'s$ and exotics, and effective couplings, is briefly described.
Z' PHYSICS FROM STRINGS

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Many extensions of the standard model, especially grand unified theories and superstring models, predict the existence of additional Z' bosons and associated exotic chiral supermultiplets. It has recently been argued that for classes of string motivated models with supergravity mediated supersymmetry breaking there are two scenarios for the additional Z's: either the mass is in the accessible range $< O(1 \text{ TeV})$, providing a natural solution to the $\mu$ problem and implications for the Higgs and sparticle masses and for the LSP; or, when the breaking is associated with a $D$-flat direction, at an intermediate scale, providing a possible explanation for the hierarchies of quark and charged lepton masses and new possibilities for neutrino masses. Related work, examining the detailed structure of specific perturbative string vacua for $D$ and $F$-flat directions, surviving $U(1)'s$ and exotics, and effective couplings, is briefly described.

1 Z' Phenomenology

If the standard model (SM) gauge group is extended by an additional $U(1)$, then the mass eigenstates $Z_{1,2}$ will be mixtures of the SM $Z$ and new $Z'$ with mixing angle $\theta$. There are stringent limits on $M_{Z_2}$ and $\theta$ from precision $Z$ pole and neutral current experiments, because: (i) $M_{Z_1}$ is shifted from the SM prediction by mixing; (ii) the $Z_1$ couplings are changed by the mixing; (iii) $Z_2$ exchange may be important in neutral current amplitudes. There are also Tevatron limits on $M_{Z_2}$ from the non-observation of $Z_2$ decays into $e^+e^-$ or $\mu^+\mu^-$. The limits are model dependent, depending not only on the chiral couplings to $e, \nu, u$, and $d$, but (in the case of the direct production limits) on the number of open decay channels into exotics, superpartners, etc. Typically, for $Z'$ properties motivated by grand unification (GUTs) one finds $M_{Z_2} > 600 - 1000 \text{ GeV}$ and $|\theta| < \text{few} \times 10^{-3}$. For $M_{Z_2} \gg M_{Z_1}$ one expects $\theta \sim C g_1' M_{Z_2}^2 / GM_{Z_2}^2$, where $G = \sqrt{g^2 + g_Y^2}$ and $g_Y$ are respectively the ordinary and new $U(1)$ gauge couplings ($g_Y$ is the weak hypercharge coupling), and $C$ depends on the Higgs charges under the $U(1)'$ and their VEVs. The most stringent limits on $M_{Z_2}$, which occur in those specific models in which $C$ is fixed, actually come from $\theta$. For models with suppressed couplings to ordinary fermions, such as leptophobic models, much smaller $M_{Z_2}$ is allowed (e.g., 150

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Work in collaboration with J. Cleaver, M. Cvetič, D. Demir, J. R. Espinosa, L. Everett, and J. Wang.
GeV; one could even have \( M_{Z_2} < M_{Z_1} \), where \( Z_1 \) is the boson that is mainly
the SM one, as is larger \(|\theta| < \text{few} \times 10^{-2} \).

It should be possible to extend the direct limits on \( Z' \) with GUT-type
couplings to around a TeV at the Tevatron. At the LHC (with 100 fb\(^{-1}\)), one
should be able to discover a \( Z' \) via its leptonic decays up to around 4 TeV,[1]
well above the range \( M_{Z'} < 1 \) TeV expected in superstring theories.[5] At an
NLC (500 GeV, 50 fb\(^{-1}\)) one has \( e^+e^- \rightarrow \gamma , Z , Z' \rightarrow e^+e^- , \mu^+\mu^- , q\bar{q} , c\bar{c} , b\bar{b} \).
Observations of cross sections, forward-backward and polarization asymme-
tries, etc., should allow a sensitivity to a (virtual) \( Z' \) up to \( \sim 1-3 \) TeV,[4]
increasing rapidly with energy[5]. Once a \( Z' \) is observed, one will want to determine
not only its mass and mixing, but its chiral couplings to identify its origin. At
the LHC, a combination of forward-backward asymmetries (as a function of
rapidity), rapidity distributions, rare decays (\( Z' \rightarrow W\ell\nu \)), and associated pro-
duction of \( Z'Z, Z'W, Z'\gamma \) should provide significant diagnostic ability up to 1 -
2 TeV[4], with the information provided by the LHC and NLC complementary.

2 String Motivated Models

It is well known that that electroweak (EW) breaking in the MSSM with
supergravity mediated supersymmetry breaking (SUGRA) can be radiative;
i.e., a positive Higgs mass square from SUSY breaking at the Planck scale can be
driven negative at low energy due to the large Yukawa coupling associated
with the \( t \) quark. In perturbative string models there are often extra non-
anomalous \( U(1)'s \) which are not broken at the string scale. These can be
broken radiatively[5] either at the electroweak scale (i.e., less than 1 TeV)[5],
or, when the breaking is associated with a \( D \)-flat direction, at an intermediate
scale[5].

2.1 Electroweak Breaking

In the ordinary MSSM the potential for the two Higgs doublets is
\( V(H_1, H_2) = V_F + V_D + V_S \), where

\[
V_F = \mu^2 (|H_1|^2 + |H_2|^2) \\
V_D = \frac{G^2}{8} (|H_1|^2 - |H_2|^2)^2 \\
V_S = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 - (B\mu H_1 \cdot H_2 + h.c.) .
\]

(A term in \( V_D \) involving charged fields has been omitted.) The \( F \) term \( V_F \) is
derived from the superpotential

\[ W = \mu \hat{H}_1 \cdot \hat{H}_2 + h_Q \hat{Q}_3 \hat{Q}_3 \cdot \hat{H}_2. \]
Unlike the ordinary SM, in which the quartic coefficient $\lambda$ in the Higgs potential is arbitrary, the coefficient $G^2 = g^2 + g_Y^2$ of the quartic $D$ term $V_D$ is associated with gauge couplings, leading to the the upper bound on the lightest Higgs scalar $m_{H_0} < M_Z$ (tree level) or $\lesssim 130$ GeV (including loops). The EW scale is $v^2 = v_1^2 + v_2^2 = (246 \text{ GeV})^2$, where $v_i = \sqrt{2} (H_i^0)$, and $M_Z = G v / 2$. The scale of $v$ is set not only by the soft SUSY breaking parameters $m_t^2$ and $B$ in $V_S$, but also by the supersymmetry preserving parameter $\mu$.

In SUGRA models one assumes that SUSY is broken in a hidden sector at some intermediate scale $M_I$ and then transmitted to the observable sector by supergravity. The soft breaking parameters are then all of the same order of magnitude $m_{3/2} \sim 1$ TeV (e.g., if $m_{3/2} \sim m_3^2 / M_{pl}$, where $M_{pl}$ is the Planck scale, then $M_I \sim 10^{11}$ GeV). In particular, all scalars in the theory (Higgs, squarks, sleptons) typically acquire positive mass squares of $O(m_{3/2}^2)$ at $M_{pl}$.

(Universal soft breaking, which we do not assume, is the stronger assumption that the scalar mass squares are all equal at $M_{pl}$.) This is the wrong sign for EW breaking. However, for sufficiently large $h_Q = O(1)$ the Yukawa interactions can drive $m_z^2$ negative (and of $O(-m_{3/2}^2)$) at low energies. Hence, radiative breaking requires a large $m_t$, consistent with the experimental value $\sim 175$ GeV.

Thus, the SUGRA mechanism can yield the needed soft parameters. However, one also requires $\mu = O(m_{3/2})$. Since $\mu$ is a supersymmetric parameter, this requires fine-tuning in the context of the MSSM, the famous $\mu$ problem. If one has some mechanism to force $\mu = 0$, then one can generate an effective $\mu_{\text{eff}} = O(m_{3/2})$ by several mechanisms, including: (1) The Giudice-Masiero mechanism in which $\mu_{\text{eff}}$ is transmitted to the observable sector by SUGRA along with the soft breaking terms. (2) The NMSSM, in which one introduces a SM singlet field $S$, with superpotential terms $W_S = h_S S \hat{H}_1 \cdot \hat{H}_2 + \kappa \hat{S}^3$, so that $\mu_{\text{eff}} = h_S \langle S \rangle$. However, the cubic term, needed to avoid an axion, allows a discrete symmetry and undesirable cosmological domain walls. (3) An extra gauge $U(1)'$ symmetry broken by the VEV of a SM singlet $S$ with $W_S = h_S S \hat{H}_1 \cdot \hat{H}_2$ can force $\mu = 0$ with $\mu_{\text{eff}} = h_S \langle S \rangle$. Unlike the NMSSM there is no domain wall problem.

When the SUGRA MSSM is considered in the context of a class of perturbative string models, one obtains in addition: (1) $\mu = 0$ by string selection rules. (2) There are typically additional non-anomalous $U(1)$'s as well as exotic chiral supermultiplets (which can play a role in radiative breaking). (3) The Yukawa couplings at the string scale are either zero or $O(g) \sim 1$, as needed for radiative breaking.
2.2 Symmetry Breaking with an Extra $U(1)'$

An additional non-anomalous $U(1)'$ gauge symmetry can be broken by the VEV of a SM singlet $S$ with nonzero $U(1)'$ charge $Q_S$. The addition of the $U(1)'$ and $S$ to the ordinary SM results in new arbitrary parameters in the scalar potential, so there is in general no prediction for the $Z'$ mass scale. However, things are much more constrained in the $U(1)'$ extension of the MSSM. Let us assume that $Q_1 + Q_2 \neq 0$, where $Q_{1,2}$ are the $U(1)'$ charges of $H_{1,2}$, so that $U(1)'$ invariance forces $\mu = 0$. If $Q_1 + Q_2 + Q_S = 0$ one can have

$$W = h_S S H_1 \cdot H_2 + h_Q Q_3 \cdot H_2 + [h_D D_1 D_2],$$

where the last term is an optional coupling of $S$ to new exotic multiplets $D_{1,2}$. The analogue of (1) becomes

$$D = \frac{h^2}{\sqrt{2}} (|H_1|^2 |H_2|^2 + |S|^2 |H_1|^2 + |S|^2 |H_2|^2)$$

$$V_D = \frac{G^2}{8} (|H_1|^2 - |H_2|^2)^2 + \frac{g_1^2}{2} (Q_1 |H_1|^2 + Q_2 |H_2|^2 + Q_S |S|^2)^2$$

$$V_S = m_S^2 |H_1|^2 + m_S^2 |H_2|^2 + m_S^2 |S|^2 - (A_S h_S S H_1 \cdot H_2 + \text{h.c.}),$$

where $g_1$ is the $U(1)'$ gauge coupling. Thus, if $S$ acquires a VEV, one has an effective parameter $\mu_{\text{eff}} = h_S(S)$, and the corresponding $(B\mu)_{\text{eff}} = A_S h_S(S)$. Acceptable EW breaking can occur if $\langle S \rangle$ and $A_S$ are of $O(\text{TeV})$. If all of the soft SUSY breaking parameters are of $O(m_{3/2})$, then one expects not only $\mu_{\text{eff}}$ and $(B\mu)_{\text{eff}}$ but also $M_Z$ and $M_{Z'}$ to be of $O(m_{3/2})$. Only some limiting (or somewhat tuned) cases will yield allowed $\theta$ and $M_{Z'}$. The $Z - Z'$ mixing angle $\theta$ is given by

$$\theta = \frac{1}{2} \arctan \left( \frac{2\Delta^2}{M_{Z'}^2 - M_Z^2} \right),$$

where

$$M_Z^2 = \frac{1}{4} G^2 (v_1^2 + v_2^2),$$

$$M_{Z'}^2 = g_1^2 (v_1^2 Q_1^2 + v_2^2 Q_2^2 + s^2 Q_3^2),$$

$$\Delta^2 = \frac{1}{2} g_1^2 G (v_1^2 Q_1 - v_2^2 Q_2)$$

are respectively the $Z$ and $Z'$ mass squares in the absence of mixing and the mixing mass squared, and $s \equiv \sqrt{2} \langle S \rangle$. Small mixing requires small $\Delta$ and/or $M_Z \ll M_{Z'}$.

Two viable scenarios were described in (i) In the Large $A_S$ Scenario the EW and $U(1)'$ breaking is driven by a large $A_S h_S$ in the last term in (1).
This leads to $v_1 \sim v_2 \sim s$ (a generalization of $\tan \beta \sim 1$). In the special case $Q_1 = Q_2 = -Q_S/2$ one finds $\theta \sim 0$ and the prediction $M_{Z'}^2 / M_Z^2 \approx 12 g_Q^2 Q_1 / G^2$. For example, a concretes model with the couplings of the $E_6$ model yields $M_{Z'} \sim 84$ GeV. This model is not leptophobic and is therefore excluded, but it illustrates a scenario that may be viable in string-derived models with suppressed couplings to ordinary fermions.

(ii) In the Large $S$ Scenario one assumes that all of the soft parameters are of $O(1 \text{ TeV})$, with $m_{22}^2 < 0$. Then $s^2 \sim -2 m_{22}^2 / g^2 Q_S^2$ and $M_{Z'}^2 \sim -2 m_{22}^2$. One can have a smaller EW scale $v_{1,2} \ll s$ by accidental cancellations (because of $V_D$ this often only involves one constraint). To avoid excessive tuning this implies $M_{Z'} \lesssim O(1 \text{ TeV})$. Then $\theta \sim (\Delta^2 / M_{Z'}^2)(M_Z^2 / M_{Z'}^2)$ is small due to $M_{Z'}^2 \ll M_Z^2$, and can be further suppressed for small $\Delta^2$.

Both scenarios have a number of interesting consequences. These include:

(i) A solution to the $\mu$ problem, with $\mu_{\text{eff}}$ naturally of $O(M_Z)$ (large $A_S$) or $O(\text{TeV})$ (large $S$). (ii) A $Z'$ and associated exotics with masses $\lesssim O(\text{TeV})$. (iii) a predictive pattern of Higgs masses (large $A_S$) or weakened upper limit on the lightest Higgs (large $S$). (iv) Characteristic shifts in the scalar masses due to the $U(1)'$ D term. (v) New dark matter possibilities (e.g., $\tilde{S}$).

The weak scale parameters needed for both scenarios can be generated by radiative breaking. As motivated by SUGRA, we assume that at $M_{pl}$ all of the scalar mass squares $(m_{1,2}^2, m_{22}^2, m_1^2, m_2^2)$ are positive and of $O(m_{3/2}^2)$, but not necessarily universal. We also assume that the gaugino masses $M_i$ and the $A$ parameters are of $O(m_{3/2})$. The coupled one-loop RGE equations for the running gauge and Yukawa couplings and the soft parameters $m^2, M, A$ were studied for various toy models and models with $E_6$ couplings to relate the initial parameters at $M_{pl}$ to the EW scale parameters. It was found that the large $|A_S|$ scenario was possible though somewhat fine-tuned (it is necessary to ensure moderate $|A_Q|$, the $A$ term associated with $h_Q$, to avoid dangerous charge-color breaking minima). The large $S$ scenario, which requires $m_S^2 < 0$ at the EW scale is most easily obtained if there is a Yukawa coupling of $S$ to exotic multiplets (the optional $h_D$ term in (3)), but can be obtained without such couplings for some (non-universal) initial conditions.

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5This model, which has the matter content and couplings of three 27-plets as well as two Higgs-like doublets from $27 + 27^*$, is anomaly free and consistent with gauge unification. It is string-motivated, i.e., the Yukawa couplings are of $O(g)$ or zero, and the $B$ and $L$ violating GUT Yukawa relations are not respected, so that $S$ can be light.

6An alternative $E_6$ model involving matter from an extra 78 has much larger kinetic mixing and can lead to leptophobic couplings.
2.3 Intermediate Scale Breaking

Another possibility is $U(1)'$ breaking associated with a $D$-flat direction at an intermediate scale\[7\], which is expected to occur in many string models and which may be associated with fermion mass hierarchies\[8\]. This can occur, for example, if there are two SM singlets $S_{1,2}$ with $Q_{S_i}Q_{S_j} < 0$. If the model is also $F$-flat at the renormalizable level (i.e., there are no terms $\hat{S}_i\hat{S}_j\hat{S}_k$ in $W$), the low energy potential for $S_{1,2}$ is

$$V(S_1, S_2) = m_1^2|S_1^2| + m_2^2|S_2^2| + \frac{g_d'^2}{2}(Q_{S_1}|S_1^2| + Q_{S_2}|S_2^2|)^2,$$  \hspace{1cm} (9)

where the quartic term vanishes for $|S_2^2|/|S_1^2| = -Q_{S_1}/Q_{S_2}$.

As an example, suppose $Q_{S_1} = -Q_{S_2}$, and further that at low energies $m_{S_1}^2 < 0$ and $m_{S_2}^2 > 0$, as would typically occur by the radiative mechanism if $W$ contains a term $h_{0}\hat{S}_1\hat{D}_1\hat{D}_2$. If $m^2 \equiv m_{S_1}^2 + m_{S_2}^2 > 0$ the minimum will occur at $\langle S_1 \rangle \neq 0$, $\langle S_2 \rangle = 0$. Then, $\langle S_1 \rangle$ and $M_{Z'}$ will be at the EW scale ($\lesssim 1$ TeV), just as in the case of a single $S$. On the other hand, for $m^2 < 0$, the potential along the $F$ and $D$ flat direction $S_1 = S_2 \equiv S$ is

$$V(S) = m^2 S^2,$$  \hspace{1cm} (10)

which appears to be unbounded from below. However, $V(S)$ can be stabilized by either of two mechanisms: a) The leading loop corrections to the effective (RGE-improved) potential result in $m^2 \rightarrow m^2(S)$ in [10]. Since $m^2$ runs from a positive value at $M_{pl}$ to a negative value at low energies, the RGE-improved potential will have a minimum close to but slightly below the scale $\mu_{RAD}$ at which $m^2$ goes through zero. It was shown in [8] that $\mu_{RAD}$ can occur anywhere in the range $10^3 - 10^{17}$ GeV, depending on the soft breaking parameters and the exotic ($\hat{D}_i$) quantum numbers. b) Another possibility is that the $F$-flatness is lifted by higher-dimensional nonrenormalizable operators (NRO) in $W$, as are expected in string models, such as $W = (\hat{S}_1\hat{S}_2)^2/M$, where $M \sim 10^{17} - 10^{18}$ GeV is of the order of the string scale. For example, if $W$ contains

$$W_S = \frac{\hat{S}^{K+3}}{M^K},$$ \hspace{1cm} (11)

when evaluated along the flat direction $\hat{S}$, then the potential will be minimized at the scale

$$\mu_{NRO} \sim \left[m_{3/2}^{1/2}M^K\right]^{1/3},$$ \hspace{1cm} (12)

\[a\] A similar mechanism could occur for a total gauge singlet field.
which is around $10^{10}$ GeV for $K = 1$.

In general, both the radiative and NRO stabilization mechanisms can occur, and $\langle S \rangle$ will be of the order of the smaller of $\mu_{\text{RAD}}$ and $\mu_{\text{NRO}}$. In both cases, one expects $M_{Z'}, M_D \sim \langle S \rangle$. Also, $V''$ is of $O(m_3^2/2)$ at the minimum, leading to an electroweak scale invisible scalar. There are also characteristic $D$-induced shifts in the effective soft masses. A effective $\mu$ parameter can be generated by the superpotential term

$$W_\mu = \hat{S}\hat{H}_1 \cdot \hat{H}_2 \left( \frac{\hat{S}}{M} \right)^{P_\mu} \Rightarrow \mu_{\text{eff}} \sim \langle S \rangle \left( \frac{\langle S \rangle}{M} \right)^{P_\mu}.$$  (13)

(The special case $P_\mu = 0$ is needed for the EW scale breaking scenario.) For radiative stabilization, one obtains the needed $\mu_{\text{eff}} \sim 1$ TeV for, e.g., $P_\mu = 1$ and $\mu_{\text{RAD}} \sim 10^{10}$ GeV. For NRO stabilization,

$$\mu_{\text{eff}} \sim m_{3/2} \left( \frac{m_{3/2}}{M} \right)^{P_\mu} \frac{\langle S \rangle}{M}.$$  (14)

which is of the order of the soft breaking (and EW) scale $m_{3/2}$ for $P_\mu = K$.

Intermediate scale breaking scenarios have interesting implications for quark, charged lepton, and neutrino masses. For example, a $u$-type quark mass may be generated by the term

$$W_u = h_u \bar{u}^c \hat{Q} \cdot \hat{H}_2 \left( \frac{\hat{S}}{M} \right)^{P_u},$$  (15)

where in string models the nonzero coefficients $h_u$ are of $O(g) \sim 1$ for $P_u = 0$, and can be absorbed into $M$ for $P_u > 0$. (15) leads to an effective Yukawa coupling and fermion mass (in the case of NRO stabilization)

$$y_u \sim \left( \frac{\langle S \rangle}{M} \right)^{P_u} \Rightarrow m_u \sim \left( \frac{m_{3/2}}{M} \right)^{P_u} \langle H_2 \rangle.$$  (16)

Presumably, the $t$ mass is associated with $P_t = 0$, while the $u$ and $c$ masses, and any inter-generational masses associated with family mixing, could be due to operators of higher dimension. Similar hierarchies of dimensions of operators could lead to small $d$ and $e$ type masses and mixings, especially for the first two families, as well as naturally tiny Dirac neutrino masses, without the need for invoking a seesaw. Which terms actually have non-zero coefficients is determined not only by gauge invariance in the four dimensional effective field theory, but by string selection rules as well. This mechanism of small effective
Yukawas suppressed by intermediate scale VEV’s is somewhat analogous to the attempts\textsuperscript{22,23} to generate Yukawas suppressed by powers of $\langle S_A \rangle / M \sim 1/10$, where $S_A$ is a field which breaks the anomalous $U(1)'$ present in many free fermionic models.\textsuperscript{24} However, the lower intermediate scale considered here allows for the use of lower dimension operators\textsuperscript{1}. It is also possible to generate Majorana masses $m_M$ for sterile ($SU(2)$-singlet) neutrinos $N_L^c$ by the operators

$$W_M \sim \bar{N}_L^c \bar{S} \left( \frac{\hat{S}}{M} \right)^{P_M},$$

implying

$$m_M \sim \langle S \rangle \left( \frac{\langle S \rangle}{M} \right)^{P_M} \sim m_{3/2} \left( \frac{m_{3/2}}{M} \right)^{P_M - K},$$

which can be large (leading to a seesaw) or small, depending on the sign of $P_M - K$. From $\langle H_{1,2} \rangle \sim m_{3/2}$ for radiative breaking, one finds that neutrino Dirac and Majorana masses can be naturally small and comparable in the special case $P_D = P_M - K$, where $P_D$ is the power analogous to $P_u$ in (15) for a Dirac neutrino mass term.\textsuperscript{19} This can lead to significant mixing between ordinary neutrinos and light sterile neutrinos, as is suggested phenomenologically by the experimental hints of neutrino mass.\textsuperscript{25}

\section{Perturbative String Vacua}

The work discussed in Section 2 was motivated by certain general features of perturbative string models, especially (a) the existence of additional $U(1)'$s and exotics, (b) that the Yukawa couplings at the string scale are either zero or of $O(g) \sim 1$, (c) that world-sheet selection rules often forbid terms in the superpotential $W$ that would be allowed by the gauge symmetries of the effective four-dimensional field theory, and (d) that there are no elementary bilinear (mass) terms in $W$. A more ambitious project is to try to derive the consequences of specific string vacua. There are many possible string vacua, and none that have been studied are fully realistic. However, there are models based on the free fermionic construction\textsuperscript{26,24} that are quasi-realistic, containing the ingredients of the MSSM (gauge group, and candidates for three ordinary families and two Higgs doublets) and some form of gauge unification. They typically also contain a (partially) hidden sector non-abelian group, an anomalous $U(1)'$, a number of extra non-anomalous $U(1)'$s, and many additional exotic chiral

\footnote{Most of the studies\textsuperscript{23} have assumed that the non-zero coefficients could be classified according to the anomalous $U(1)_A$ symmetry. However, this is not the case in free fermionic models.\textsuperscript{24}}
supermultiplets. The latter include non-chiral exotic multiplets, fractionally charged states, and mixed states transforming non-trivially under both the ordinary and hidden sector groups.

A first step in studying the low energy consequences of such models is to determine which fields acquire VEVs at or near the string scale, in a way that breaks the anomalous $U(1)'$ but maintains $D$ and $F$ flatness. Techniques have recently been developed to compute classes of $D$-flat directions that can be proved $F$-flat to all orders. A number of models were considered, and it was found that those which have such flat directions composed on non-abelian singlet fields typically leave one or more non-anomalous $U(1)'$’s unbroken at the string scale. A next step, currently in progress, is to study the effective superpotential of the resulting model in these flat directions, after replacing the scalar fields which appear in the flat directions by their VEVs. In particular, it will be possible to study the $U(1)'$ breaking patterns and low energy consequences, after making appropriate ansätze for the soft supersymmetry breaking terms and the structure of the Kähler potential. It is unlikely that any realistic models will be found, but it is hoped that the analysis will give useful insights into the type of physics consequences that may derive from perturbative string theories.

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