

Aspects of Z' -mediated supersymmetry breaking

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In a recent paper, we proposed the possibility that supersymmetry breaking is communicated dominantly via a $U(1)'$ vector multiplet. We also required that the $U(1)'$ plays a crucial role in solving the μ problem. We discuss here in detail both the construction and the phenomenology of one class of such models. The low-energy spectrum generically contains heavy sfermions, Higgsinos and exotics ~ 10 – 100 TeV; an intermediate $M_{Z'} \sim 3$ – 30 TeV; light gauginos ~ 100 – 1000 GeV, of which the lightest can be W -ino-like; a light Higgs with a mass of ~ 140 GeV; and a singlino which can be very light. We present a set of possible consistent charge choices. Several benchmark models are used to demonstrate characteristic phenomenological features. Special attention is devoted to interesting LHC (Large Hadron Collider) signatures such as gluino decay and the decay patterns of the electroweak-inos. Implications for neutrino masses, exotic decays, R -parity, gauge unification, and the gravitino mass are briefly discussed.

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I. INTRODUCTION

Many supersymmetry (SUSY) breaking mediation mechanisms, such as gravity mediation [1–7], anomaly mediation [8,9], gauge mediation [10–20], and gaugino mediation [21,22], have been proposed (for a review, see [23]). In a recent paper [24], we proposed that supersymmetry breaking could instead be communicated naturally by some exotic gauge interactions. A typical example of such a mediator is an extra $U(1)'$.¹ The existence of low-energy supersymmetry would give indirect evidence that TeV scale new physics could be directly embedded into some high scale fundamental theory, such as string theory. Concrete semirealistic superstring constructions frequently lead to additional nonanomalous $U(1)'$ factors in the low-energy theory (see, e.g., [30–41]), and in some cases both the ordinary sector and hidden sector particles carry $U(1)'$ charges, allowing a $U(1)'$ -mediated communication between the two sectors. More recently [42], it was realized that there is a natural way of implementing such a mediation mechanism in a large class of D-brane constructions.

Motivated by the μ -problem of the minimal supersymmetric standard model (MSSM), we focused on one class of solutions, which invokes a spontaneously broken PQ symmetry (see, e.g., [43]). From the point of view of top-down constructions it is common that such a symmetry is promoted to a $U(1)'$ gauge symmetry [44,45]. It is natural to make this $U(1)'$ the mediator of SUSY breaking as well, since in this case μ (as well as μB) will be set by the scale of the other soft SUSY breaking parameters. Whether or not the electroweak symmetry breaking is finely tuned, μ and μB terms generated this way are of the right-size. We

would like to include this as a feature of the class of models we consider, though it is not absolutely essential.

In our setup, a supersymmetry breaking Z' -ino mass term, $M_{Z'}$, is generated due to $U(1)'$ coupling to the hidden sector. The observable sector fields feel the supersymmetry breaking through their couplings to $U(1)'$, implying interesting features of the sparticle spectrum. The sfermion masses are of the order of $m_{\tilde{f}}^2 \sim M_{Z'}^2/16\pi^2$. The $SU(3)_C \times SU(2)_L \times U(1)_Y$ gaugino masses are generated at higher loop order, $M_{1/2} \sim M_{Z'}/(16\pi^2)^2$, which is 2–3 orders of magnitudes lighter than the sfermions. LEP direct searches suggest electroweak-ino masses >100 GeV [46]. We therefore expect that the sfermions are heavy, typically about 100 TeV. In this sense, this scenario could be viewed as a mini-version of split supersymmetry [47,48]. One important difference is the μ -parameter, which is set by the scale of $U(1)'$ breaking. Although in principle a free parameter, we find it is naturally at the same order of magnitude as the sfermions. Similar to split supersymmetry, one fine-tuning is needed to maintain a low electroweak scale. The scenario does not have flavor or CP violation problems due to the decoupling of the sfermions. The flavor violation in the scenario will be further suppressed if we choose flavor universal $U(1)'$ charges for the standard model matter fields. Because of the same decoupling effect, we expect that the contribution to the muon anomalous magnetic moment is negligible in this scenario.

This paper is organized as follows. We first review the generic setup and resulting sparticle spectrum. Then, as an example, we construct a specific model with more assumptions about the consistency conditions and the existence of specific types of exotics. Finally, we comment on the phenomenology of this class of models, including the spectrum, the gluino life time, cold dark matter, possible ranges for the gravitino mass, exotic decays, possibilities for neutrino mass, R -parity, and gauge unification.

¹Scenarios involving an extra $U(1)'$ in supersymmetry mediation have been considered previously [25–29]. Here we assume that the Z' -mediation is the dominant source for both scalar and gaugino masses.

II. GENERIC FEATURES OF Z' -MEDIATED SUPERSYMMETRY BREAKING

The schematics of the $U(1)'$ mediation model are presented in Fig. 1. We extend the MSSM in the following ways. First, introduce an extra $U(1)'$ gauge symmetry. Second, promote the μ parameter into a dynamical field S which is charged under the $U(1)'$. Third, include other exotics with Yukawa couplings to S . The last assumption is included to drive the necessary radiative symmetry breaking and to cancel anomalies. Such exotics and couplings generically exist in string theory constructions. The superpotential is

$$W = y_u H_u Q u^c + y_d H_d Q d^c + y_e H_d L e^c + y_\nu H_u L \nu^c + \lambda S H_u H_d + \sum_{i \in \{\text{exotics}\}} y_i S X_i X_i^c, \quad (1)$$

where i labels the species of exotics.

A. Features of the spectrum

We begin by discussing the pattern of the soft supersymmetry breaking parameters, the masses of the Z' -ino and of the MSSM squarks and gauginos, which are the most robust predictions of this scenario. At the supersymmetry breaking scale, Λ_S , supersymmetry breaking in the hidden sector is assumed to generate a supersymmetry breaking mass for the fermionic component of the $U(1)'$ vector superfield. Given details of the hidden sector, its value could be evaluated via the standard technique of analytical continuation into superspace [49]. In particular, the gauge kinetic function of the field strength superfield \hat{Z}' at the supersymmetry breaking scale is

$$\mathcal{L}_{\hat{Z}'} = \int d^2\theta \left[\frac{1}{g_{Z'}^2(0)} + \beta_{Z'}^{\text{hid}} \log\left(\frac{\Lambda_S}{M}\right) + \beta_{Z'}^{\text{vis}} \log\left(\frac{\Lambda_S}{M_{Z'}}\right) \right] \hat{Z}' \hat{Z}', \quad (2)$$

where M is the messenger scale, which we have assumed to be around the supersymmetry breaking scale, $M \sim \Lambda_S$. $\beta_{Z'}^{\text{hid}}$ and $\beta_{Z'}^{\text{vis}}$ are β -functions induced by $U(1)'$ couplings to hidden and visible sector fields, respectively. Using analytical continuation, we replace M with $M + \theta^2 F$,

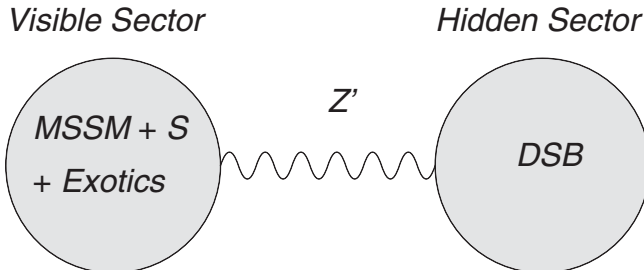


FIG. 1. Z' -mediated supersymmetry breaking.

where F is the supersymmetry breaking order parameter. We obtain the \tilde{Z}' mass as $M_{\tilde{Z}'} \sim g_{Z'}^2 \beta_{Z'}^{\text{hid}} F/M$. We assume that the $U(1)'$ gauge symmetry is not broken in the hidden sector.

We assume that all the chiral superfields in the visible sector are charged under $U(1)'$, so all the corresponding scalars receive soft mass terms at 1-loop of order,²

$$m_{\tilde{f}_i}^2 \sim \frac{g_{Z'}^2 Q_{f_i}^2}{16\pi^2} M_{\tilde{Z}'}^2 \log\left(\frac{\Lambda_S}{M_{\tilde{Z}'}}\right), \quad (3)$$

where $g_{Z'}$ is the $U(1)'$ gauge coupling and Q_{f_i} is the $U(1)'$ charge of f_i , which we take to be of order 1. [The exact expressions can be determined from the renormalization group equations (RGEs) given in Appendix B.]

The $SU(3)_C \times SU(2)_L \times U(1)_Y$ gaugino masses, however, can only be generated at 2-loop level since they do not directly couple to the $U(1)'$,

$$M_a \sim \frac{g_{Z'}^2 g_a^2}{(16\pi^2)^2} M_{\tilde{Z}'} \log\left(\frac{\Lambda_S}{M_{\tilde{Z}'}}\right), \quad (4)$$

where g_a is the gauge coupling for the gaugino $\tilde{\lambda}_a$, and the internal line is the sum over the chiral supermultiplets charged under the a^{th} gauge group. [We have suppressed the group and $U(1)'$ charge factors.] Since these gaugino masses are proportional to g_a^2 , we expect that the gluino will typically be significantly heavier than the others. However, that conclusion and the ordering of W -ino and b -ino masses depends on specific charge assignments and the exotic matter content.

From the discussion above, we see that the gauginos are considerably lighter than the sfermions. Taking $M_a \gtrsim 100$ GeV, we find

$$M_{\tilde{Z}'} \log\left(\frac{\Lambda_S}{M_{\tilde{Z}'}}\right) \sim 10^4 \text{ TeV} \quad (5)$$

and

$$m_{\tilde{f}_i} \sim \frac{(4\pi)^3}{g_{Z'} g_a^2} M_a \sim 100 \text{ TeV}, \quad (6)$$

where we have assumed that $g_{Z'}$ is of electroweak strength. Similarly, A terms associated with the Yukawa interactions in (1) are generated at one loop by \tilde{Z}' exchange, yielding

$$A \sim \frac{y g_{Z'}^2}{16\pi^2} M_{\tilde{Z}'} \log\left(\frac{\Lambda_S}{M_{\tilde{Z}'}}\right) \sim y \times 10 \text{ TeV}, \quad (7)$$

²Equation (3) cannot be the full story or we would not be able to drive the singlet scalar mass-square negative or keep the Higgs light. However, this contribution does serve to set the overall scale. To generate a much lighter mass scale requires fine-tuning.

where the Yukawa coupling y is absorbed into A . Again, the exact expressions, including the counting factors and dependence on the $U(1)'$ charges, can be found from the expressions in Appendix B.

The gravitino mass $m_{3/2} \sim F/M_P$ depends strongly on the size of supersymmetry breaking in the hidden sector. Requiring MSSM gaugino masses ≥ 100 GeV sets $F/M \geq 10^6/(g_z^2 g_a^2)$ GeV. Assuming \sqrt{F} , M , and Λ_S to be within an order of magnitude, we could have $\sqrt{F} \sim 10^7\text{--}10^{11}$ GeV. This gives a wide range of gravitino masses with very different phenomenologies, as will be discussed in Sec. IV B 2. This is very different from gauge mediated supersymmetry breaking, where a typically lower supersymmetry breaking scale ($\sim 10\text{--}1000$ TeV) implies a gravitino much lighter than the other superpartners. Without specifying a particular model of supersymmetry breaking and gravity mediation, we will treat it as a free parameter to begin with. The gravitino will be further constrained by cosmological data such as big bang nucleosynthesis and by the cold dark matter density Ω_{CDM} , which we will discuss in Sec. IV B 2. We will assume that \sqrt{F} is not very different, within a couple of orders of magnitude, from the supersymmetry breaking scale Λ_S . In the scenario under consideration, the supersymmetry breaking scale is constrained logarithmically by the requirement of radiative symmetry breaking. Since the relevant Yukawa contributions that ultimately fix the electroweak scale are proportional to $y_i^2 \log(\Lambda_S/M_{Z'})/16\pi^2$, the gravitino mass is exponentially sensitive to the choice of the parameters in the model, as $m_{3/2} \propto e^{1/y_i^2}$.

B. Kinetic mixing

Kinetic mixing between the Z' -ino and b -ino will be generated at one-loop level through interactions with the visible sector matter content. It is generically of the order $k \sim (g_z g_Y \text{tr}(Q_i Y_i)/16\pi^2) \log(\Lambda_S/M_{Z'})$ [where Q_i are the $U(1)'$ charges, Y_i are the hypercharges and the trace is taken over all charged matter]. The existence of such kinetic mixing implies that we must add a term $K \hat{B} \hat{Z}'$ to Eq. (2), where K is in general a holomorphic function whose lowest component acquires a vev k . If K involves some hidden sector field, the induced correction to the light b -ino mass will be on the order of $k^2 M_{Z'}$, which is at the same order as the contribution considered in the previous section. However, we have assumed that only visible sector fields, which do not participate in supersymmetry breaking, are charged under both hypercharge and $U(1)'$. Therefore, by construction, such a contribution is absent at the one-loop level in our scenario. It will enter at higher loop order, which is negligible in comparison with the two-loop contribution we have considered.

We now discuss the effect of the kinetic mixing. This will shift the mass of the b -ino, but such a shift is proportional to the square of the b -ino's mass.

The gaugino kinetic and mass terms are

$$\mathcal{L} = -i(\tilde{B}^\dagger \tilde{Z}'^\dagger) \begin{pmatrix} 1 & k \\ k & 1 \end{pmatrix} \bar{\sigma}^\mu D_\mu \begin{pmatrix} \tilde{B} \\ \tilde{Z}' \end{pmatrix} + (\tilde{B} \tilde{Z}') \begin{pmatrix} M_1 & 0 \\ 0 & M_{Z'} \end{pmatrix} \begin{pmatrix} \tilde{B} \\ \tilde{Z}' \end{pmatrix}, \quad (8)$$

where we have ignored nondiagonal terms which are of the same order as M_1 due to their negligible effect on the b -ino mass. After bringing the kinetic term to its canonical form, we find that the new mass eigenvalues are

$$M_1 \rightarrow M_1 \left(1 - \frac{M_1 k^2}{M_{Z'}}\right), \quad M_{Z'} \rightarrow M_{Z'}(1 + k^2), \quad (9)$$

where $M_1 \ll M_{Z'}$, $k < 1$, and we keep only the lowest order terms. This result can be understood in terms of chiral symmetry. In the limit of vanishing b -ino mass, the zero eigenvalue of the gaugino mass matrix is not changed by the congruence transformation that brings the kinetic term to its canonical form. As a result, once the b -ino mass is ‘‘turned on,’’ the shift must be proportional to it.

Similar results apply to kinetic mixing between the Z and Z' gauge bosons [50,51]. In particular when the Z' gauge boson becomes massive, the Z gauge boson remains massless. The only effect this mixing will have is to shift the Z' gauge coupling: $g_z Q_i \rightarrow g_z Q_i - k g_Y Y_i$.

C. Contribution from other mediation mechanisms

Since the soft parameters resulting from Z' mediation involve a large hierarchy, nondominant contributions from other mediation mechanisms could also be important. For example, there could be other gauge interactions between the hidden and observable sector. However, as long as they do not contain SM gauge interactions, we expect the main features of the soft breaking parameters to continue to hold, as most of our discussion above is independent of the gauge group [except the kinetic mixing between $U(1)$ factors]. The other obvious candidate is gravity mediation, which yields a contribution to the gaugino mass of order F/M_P . Since we have $M_{Z'} \sim (g_z^2/16\pi^2)F/M$, we expect $\sqrt{F} \sim 10^7\text{--}10^{11}$ GeV, without assuming a large hierarchy between Λ_S , \sqrt{F} , and M . Therefore, gravity mediation could give comparable contributions to the gaugino masses for higher values of F . On the other hand, its contribution to soft scalar mass-squares $\sim F^2/M_P^2$ is expected to be much smaller than the Z' -mediated contribution. This is very different from gauge mediation through the standard model gauge groups, where all the soft terms are of the same order. Therefore, while in principle the gravity mediation contribution to the gaugino masses could be comparable to the one from Z' -mediation, we expect the hierarchy between the sfermions and scalars to be a robust prediction of this scenario.

It is possible that the gravity mediation piece is sequestered [52–54]. In this case, the dominant supergravity contribution will come from anomaly mediation. Such contribution could be important in our case if $\Lambda_S \sim 10^{11} - 10^{12}$ GeV.

D. Symmetry breaking and fine-tuning

The $U(1)'$ gauge symmetry is broken by the vev $\langle S \rangle$. We assume that this symmetry breaking is triggered by radiative corrections to the soft mass m_S^2 , especially through Yukawa couplings to exotics.³

We are looking for parameters which result in solutions such that $\langle S \rangle \gg v$, where $v \equiv (|\langle H_u^0 \rangle|^2 + |\langle H_d^0 \rangle|^2)^{1/2} \sim 174$ GeV is the electroweak scale. It is therefore reasonable to first determine $\langle S \rangle$ ignoring the Higgs doublets, and then to consider the Higgs potential for the doublets regarding $\langle S \rangle$ as a fixed parameter. We have verified that the corrections from the shift in $\langle S \rangle$ due to the doublets is small. The scalar potential for S is

$$V(S) = m_S^2 |S|^2 + \frac{1}{2} g_{z'}^2 Q_S^2 |S|^4, \quad (10)$$

which is minimized for $\langle S \rangle^2 = -m_S^2 / g_{z'}^2 Q_S^2$ for $m_S^2 < 0$. The $U(1)'$ symmetry breaking is driven by the radiative corrections to m_S^2 .

$$16\pi^2 \frac{dm_S^2}{d \log \mu} = -8 g_{z'}^2 Q_S^2 M_{Z'}^2 + 4\lambda^2 (m_S^2 + m_{H_u}^2 + m_{H_d}^2) + 2 \sum_{\{\text{exotics}\}} y_i^2 (m_S^2 + m_{X_i}^2 + m_{X_i^c}^2). \quad (11)$$

The charges and Yukawa couplings have to be chosen so that radiative symmetry breaking actually occurs. The relative contribution to m_S^2 from the exotics goes as $-(y_i^2 / 16\pi^2) m_{X_i}^2 \log(\Lambda_S / M_{Z'})$.⁴ Therefore, successful radiative breaking of $U(1)'$ usually requires that the Yukawa couplings to the exotics are not small and that some hierarchy exists between Λ_S and $M_{Z'}$, which depends on the choice of the Yukawa couplings. We will illustrate such effects in the context of a specific model, for which we typically find $\langle S \rangle \sim 100$ TeV.

Meanwhile, to generate the electroweak scale we must fine-tune one linear combination of the two Higgs doublets to be much lighter than its natural scale. The Z' -ino mass, $M_{Z'}$, sets the overall scale in the visible sector, so the tuning must be between the dimensionless couplings in the model, namely, $g_{z'}$, λ , and the other Yukawas, as well as the ratio $\log(\Lambda_S / M_{Z'})$. While the restriction on the parameter space

³An alternative possibility, which we have not investigated, would be for the $U(1)'$ gauge symmetry to also be broken in the hidden sector. In that case, $M_{Z'}$, $M_{Z'}$, and Λ_S would all be free parameters.

⁴The Yukawa contribution to the running actually continues below $M_{Z'}$, but in most of the cases considered this is a small effect.

from $U(1)'$ breaking is model dependent, the need for fine-tuning to obtain the electroweak symmetry breaking is generic.

The full mass matrix for the two Higgs doublets is,

$$\mathcal{M}_H^2 = \begin{pmatrix} m_2^2 & -A_H \langle S \rangle \\ -A_H \langle S \rangle & m_1^2 \end{pmatrix}, \quad (12)$$

$$m_2^2 = m_{H_u}^2 + g_{z'}^2 Q_S Q_2 \langle S \rangle^2 + \lambda^2 \langle S \rangle^2,$$

$$m_1^2 = m_{H_d}^2 + g_{z'}^2 Q_S Q_1 \langle S \rangle^2 + \lambda^2 \langle S \rangle^2,$$

where $Q_2 \equiv Q_{H_u}$, $Q_1 \equiv Q_{H_d}$, and all the couplings and mass terms are evaluated at $M_{Z'}$. A_H is the radiatively generated soft trilinear coupling between the Higgs doublets and the singlet, $\mathcal{L} = -A_H H_u H_d S + \text{H.c.}$ Electroweak breaking requires one small eigenvalue $\mathcal{O}(v^2)$. Since A_H has unit mass dimension it is generally suppressed with respect to the scalar soft masses by about an order of magnitude ($\sim 4\pi$). Therefore, we will have to tune one of the diagonal terms to be small. The up-type Higgs soft mass can usually be driven negative owing to the large top Yukawa coupling. We can then tune it against the other contributions to the diagonal up-type entry. In particular, we will adopt the following scheme for the tuning. Since $\langle S \rangle^2$ has a different dependence on Λ_S than $m_{H_u}^2$ we will keep all the couplings and mass scales fixed and allow ourselves to vary only Λ_S . This suffices to generate a small eigenvalue and obtain the electroweak scale.

Since the down-type mass term is much larger than all the other scales, $\tan \beta$ is well approximated by

$$\tan \beta = \frac{m_1^2}{A_H \langle S \rangle} \sim 10-100. \quad (13)$$

Next, we turn to discuss the part of the mass spectrum which will be determined by the $U(1)'$ symmetry breaking.

The effective μ term is $\mu = \lambda \langle S \rangle$. Assuming $\lambda = \mathcal{O}(0.1 - 1)$, we have $\mu \sim 10-100$ TeV. Similarly, the fermionic component of the exotic superfields X_i and X_i^c will acquire supersymmetric masses $y_i \langle S \rangle \sim 10-100$ TeV.

The Z' mass is

$$M_{Z'} = \sqrt{2} g_{z'} Q_S \langle S \rangle = \sqrt{2} |m_S|. \quad (14)$$

The singlino $\tilde{\nu}$ receives a mass through mixing with the Z' -ino. The mass matrix is given by

$$\mathcal{M}_{SZ} = \begin{pmatrix} 0 & \sqrt{2} g_{z'} Q_S \langle S \rangle \\ \sqrt{2} g_{z'} Q_S \langle S \rangle & M_{Z'} \end{pmatrix}, \quad (15)$$

where we ignore any possible phases. For $|m_S| \ll M_{Z'}$ the eigenvalues are given by the usual seesaw formula,

$$\mathcal{M}_{SZ}^{(1)} = -\frac{2|m_S|^2}{M_{Z'}} = -\frac{M_{Z'}^2}{M_{Z'}}, \quad \mathcal{M}_{SZ}^{(2)} = M_{Z'}. \quad (16)$$

The mass of both the Z' gauge boson and the singlino are governed by $|m_S|$ which is naively of the same order as the

other soft scalar masses ~ 100 TeV. However, there is an interesting limit with $g_{z'} \ll \lambda$ in which the fine-tuning required in the Higgs sector leads to smaller values for $|m_S|$. The singlet's vev, $\langle S \rangle$, contributes to the mass of the up-type Higgs as in Eq. (12). The necessary cancellation in Eq. (12) prevents the singlet's vev from becoming too large or it is impossible to tune the up-type Higgs mass to be of the order of the EW scale if λ is order unity. The typical value of m_{H_u} is $\sim g_{z'} M_{Z'}/4\pi$, a loop-factor below $M_{Z'}$, so we expect $\langle S \rangle \sim (g_{z'}/\lambda) M_{Z'}/4\pi$. But, this implies an even lower scale for m_S ,

$$|m_S| \sim \frac{g_{z'}^2/\lambda}{4\pi} M_{Z'} \sim (10^{-2}-10^{-3}) M_{Z'}. \quad (17)$$

We refer to this phenomenon, where $|m_S|$ is lighter than expected, as accidental tuning. It is accidental because it comes about as a result of the fine-tuning in the Higgs sector and the smallness of the gauge coupling, $g_{z'} \ll \mathcal{O}(1)$.

This accidental tuning leads to a Z' gauge boson and singlino much lighter than expected. The Z' gauge-boson mass $M_{Z'} = \sqrt{2}|m_S|$ can be light enough to be produced at the Large Hadron Collider (LHC). The singlino is even lighter with a mass $m_{\tilde{\chi}} = 2|m_S|^2/M_{Z'} \sim 10^{-3}-10^{-6} M_{Z'}$. It may even be the lightest super-partner (LSP) as we shall demonstrate below with explicit models.

At low energies there will be a single standard model-like Higgs scalar, while the other linear combination, as well as the charged Higgs and the pseudoscalar, is heavy at the 100 TeV scale. The Higgs mass is somewhat heavier than the typical prediction of the MSSM, due to the $U(1)'$ D term and the running of the effective quartic coupling from $M_{Z'}$ down to the electroweak scale.

III. MODEL BUILDING

A. Charge assignments

We first outline some general considerations for model building in this scenario and then present a particular model which satisfies all of these requirements. Variations on most of these assumptions are possible, but beyond the scope of this paper.

The free parameters are the $U(1)'$ charges of the particles, $g_{z'}$, λ , the exotic Yukawa couplings, $M_{Z'}$, and the supersymmetry breaking scale Λ_S .

We will consider scenarios in which $U(1)'$ is anomaly free under the visible sector fields. This, along with the need for radiative breaking, will require the introduction of exotic fields. In principle, since some of the hidden sector fields must carry $U(1)'$ charges they could also contribute to the anomaly cancellation. However, such hidden sector fields would have to be chiral. If they are, to have masses characteristic of the hidden sector dynamics the $U(1)'$ would have to be broken in the hidden sector. There would therefore be a tendency for the entire $U(1)'$ supermultiplet

to decouple at around the supersymmetry breaking scale, making it more difficult to mediate the supersymmetry breaking. We will therefore assume for simplicity that the hidden sector fields are nonchiral under $U(1)'$.

We will also assume that all of the visible sector fields carry $U(1)'$ charges, that there is a single standard model singlet S which not only breaks $U(1)'$ but also generates an effective μ parameter and exotic masses, and that all standard model Yukawa couplings are allowed. The latter will include the Dirac coupling for the right-handed neutrino, but we will comment on a variation in which this is forbidden.

With a large $\langle S \rangle$ there is a danger that the quark and/or slepton fields could become tachyonic due to the $U(1)'$ D terms, leading to charge and color breaking. We of course require that this does not occur.

Finally, the LSP in these models is usually one of the standard model gauginos. Because of the well-known difficulties with a b -ino LSP, we will choose the $U(1)'$ charges to ensure a W -ino LSP instead. We do not make any *a priori* requirements concerning gauge unification, exotic decays, kinetic mixing between the $U(1)_Y$ and $U(1)'$ gauge bosons or gauginos, or R -parity, but will comment on all of them below.

B. A model

There are many possible $U(1)'$ charge assignments for the ordinary and exotic fields in a supersymmetric theory [55,56]. The most commonly studied are based on the breaking of the E_6 group to $SU(5) \times U(1) \times U(1)$, which yields an anomaly-free model consistent with gauge unification [57–59]. However, it is rather complicated, involving three S -type fields, three pairs D and D^c of exotic charge $\mp 1/3$ quarks, as well as multiple $SU(2)$ doublets which can be interpreted as extra Higgs doublets or as exotic lepton doublets. The latter ensure a b -ino LSP when combined with the Z' mediation scenario.

We will therefore explore an alternative model, characterized by a single S field and family universal charges. To ensure a W -ino LSP we will not introduce any exotic $SU(2)$ doublets (i.e., no exotic leptons or extra Higgs pairs), but will allow n_D pairs D , D^c of exotic quarks with weak hypercharge $\pm y_D$, and n_E pairs E , E^c of exotic leptons with weak hypercharge $\pm y_E$. The exotics are nonchiral with respect to the standard model gauge group, but chiral with respect to $U(1)'$. Without loss of generality, we can assume family-diagonal exotic Yukawa couplings

$$W_{\text{exotic}} = S \left(\sum_{i=1}^{n_D} y_{D_i} D_i D_i^c + \sum_{j=1}^{n_E} y_{E_j} E_j E_j^c \right) \quad (18)$$

[cf. (1)]. In practice, we will take a common value y_D for each y_{D_i} and similarly for the y_{E_j} .

The anomaly conditions are analyzed in Appendix A. It is found that the simplest solution to the mixed anomaly

TABLE I. $U(1)'$ charges for a particular anomaly-free model. We assume $Q_{H_u} \equiv Q_2 \equiv x \neq -1$, and usually take $x = -1/4$ in our numerical examples.

H_d	1	L	$\frac{2}{3} - \frac{1}{3}x$
H_u	x	e^+	$-\frac{5}{3} + \frac{1}{3}x$
S	$-(1+x)$	N^c	$-\frac{2}{3}(1+x)$
Q	$-\frac{1}{3}$	D	$\frac{8}{9} + \frac{2}{9}x$
u^c	$\frac{1}{3} - x$	D^c	$\frac{1}{9} + \frac{7}{9}x$
d^c	$-\frac{2}{3}$	E	$\frac{5}{3} - \frac{1}{3}x$
		E^c	$-\frac{2}{3} + \frac{4}{3}x$

constraints requires $n_D = 3$ color triplet pairs with hypercharge (electric charge) $Y_D = \mp 1/3$, and $n_E = 2$ singlet pairs with $Y_E = \mp 1$. There are two 2-parameter solutions for the $U(1)'$ charges, for which the quark doublet and H_u charges Q_Q and Q_2 are free parameters (after making the normalization $Q_1 = 1$), and two especially simple 1-parameter special cases in which Q_Q is fixed. We will mainly but not exclusively consider one of these special cases, with charges listed in Table I. We reemphasize that this is only a particularly simple example of a large range of possibilities.

IV. PHENOMENOLOGY

The low-energy phenomenology depends on several free parameters, such as the charge assignments, the exotics' Yukawa couplings, the PQ-symmetry breaking coupling λ , and the $U(1)'$ gauge coupling $g_{z'}$. In this section we explore the parameter space spanned by these choices and arrive at a global picture of the low-energy phenomenology.

We begin with the charge assignment. After imposing the anomaly cancellation conditions and normalizing the down-type Higgs $U(1)'$ charge to unity, $Q_1 = 1$, one is left with two undetermined charges, namely, Q_2 and Q_Q (the up-type Higgs charge and the left-handed quarks' charge, respectively). The other charges are given by the equations in Appendix A. The up-type Higgs charge, Q_2 , has to be small or otherwise it is either impossible to turn m_S^2 negative or fine-tune $m_{H_u}^2$ to be small. In Fig. 2 we present a scan of the points in the (Q_Q, Q_2) space where a solution is possible, i.e., where it is possible to obtain the EW scale without driving any of the scalars tachyonic or have any other supersymmetry partner too light. The scan utilizes the “+” solution in (A10), but we have verified that the “-” solution is similar.

One simple choice of charges has $Q_2 = -1/4$ and $Q_Q = -1/3$ (see Table I). We normalize the coupling $g_{z'}$ to the hypercharge $U_Y(1)$ at the cutoff Λ_S ,

$$g_{z'}^2 = \mathcal{N}^2 g_Y^2 \frac{\text{Tr} Y_i^2}{\text{Tr} Q_i^2}. \quad (19)$$

We leave ourselves the freedom to choose a factor \mathcal{N}^2 of

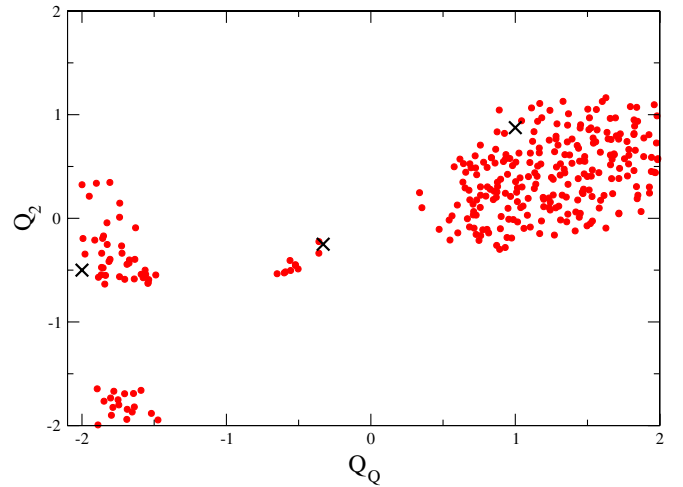


FIG. 2 (color online). The (red) dots in the $Q_Q - Q_2$ plane represent points for which a viable solution exists and where the electroweak scale is obtained, using the “+” solution in (A10). We fixed $\lambda = y_D = 0.5$ and $y_E = 0.1$. The regions of viable solutions will change as we vary these parameters, but we have verified that the overall structure remains unchanged. We then picked three representative points, one from each “island” (indicated by “x”), and investigated the resulting spectrum in detail. The point at $Q_Q = -1/3$ corresponds to the charges in Table I with $x = -1/4$.

order unity. With the above choice of charges, $g_{z'} \sim \frac{2}{3} g_Y \mathcal{N} \sim 0.23 \mathcal{N}$ at the SUSY breaking scale.

The other important parameters are the colored exotics' Yukawa, y_D , and λ . To gain a better insight into the range of possibilities in this class of models we performed a scan over both parameters and demanded that the electroweak scale is obtained by fine-tuning the SUSY breaking scale Λ_S . The details of this procedure are summarized in Appendix B.

As commented in Sec. II, we expect the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauginos to be light. In our particular model, where none of the exotics are charged under the SM $SU(2)$, the lightest gaugino is typically the W -ino. In addition, the singlino is usually also quite light as a consequence of the seesaw mass relations of Eq. (16). In Fig. 3, we plot the low-energy spectrum as a function of the colored exotics' Yukawa coupling, y_D .

It is important to consider the other islands of acceptable charge assignments shown in Fig. 2. As we already saw, there is not much freedom in the choice of Q_2 . Choose it too large and it becomes impossible to break $U(1)'$. However, we can try to choose a larger value for Q_Q . To pick a point on the left island we take $Q_Q = -2$ and $Q_2 = -1/2$.⁵ From the island on the right we take $Q_Q = 1$ and $Q_2 = 7/8$. We verified that the conclusions which

⁵If one requires rational charges one needs $Q_2 = -79/160$.

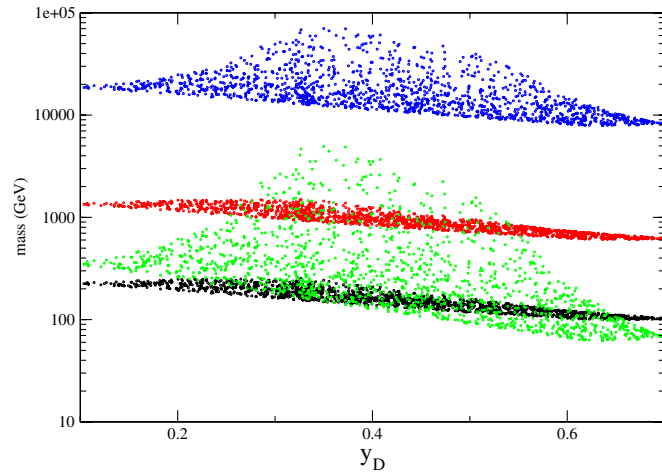


FIG. 3 (color online). A plot of the low-energy masses as a function of the colored exotic fields' Yukawa coupling, y_D . Z' gauge boson (blue, first from top right), gluino (red, second from top right), W -ino (black, third from top right), and singlino (green, fourth from top right). This spectrum corresponds to the charge assignment $Q_Q = -1/3$ and $Q_2 = -1/4$. The $U(1)'$ gauge coupling is set according to Eq. (19) with a factor of $\mathcal{N} = 0.5$. The b -ino mass is slightly lighter than the gluino mass and is not shown in order to reduce clutter. The spread corresponds to a variation in λ , the Higgs coupling to the singlet.

follow hold for other choices and are generic. The resulting spectra are shown in Figs. 4 and 5, respectively.

It is clear from the figures that the low-energy spectrum has a variety of patterns in the space of Z' -mediated supersymmetry breaking models. In particular, different ordering of the MSSM gauginos and the singlino could give rise to very different phenomenology, and the appearance of a light Z' gauge boson may prove to be the strongest indicator of the nature of the SUSY breaking

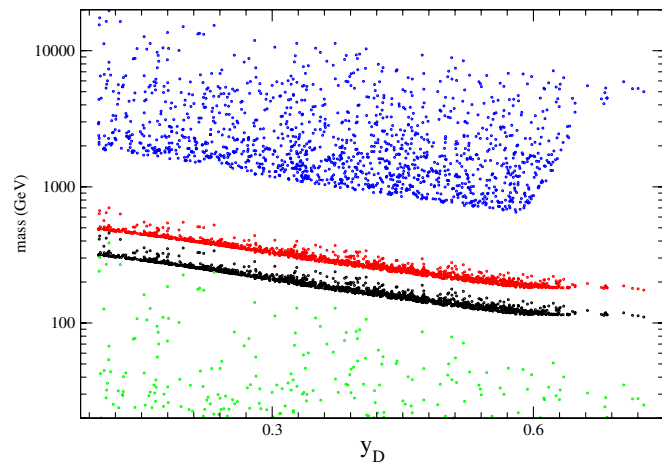


FIG. 4 (color online). Same as Fig. 3, except $Q_Q = -2$ and $Q_2 = -1/2$. The coloring and ordering from top right remains the same. The b -ino in this case is actually heavier than the gluino.

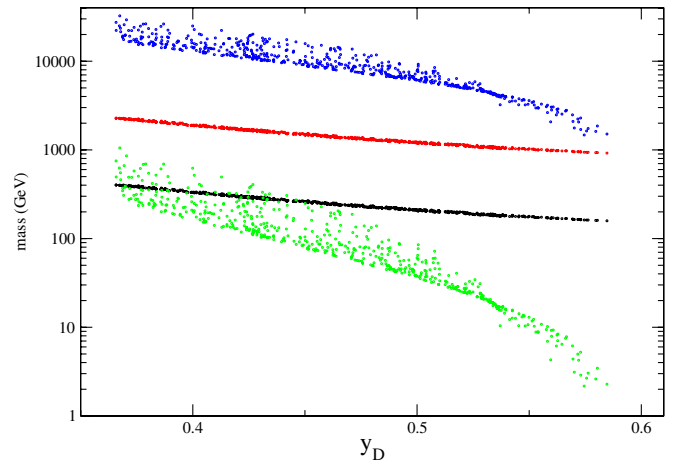


FIG. 5 (color online). Same as Fig. 3, except $Q_Q = 1$ and $Q_2 = 7/8$. The coloring and ordering from top right remains the same.

TABLE II. Model inputs and superpartner spectra of six representative models. The masses are in GeV. We fix $M_{Z'} = 10^6$ GeV. The masses of the first two generations of squarks and sfermions are typically larger than that of the third. The input parameters λ , $g_{z'}$, and $y_{D,E}$ are defined at Λ_S . The spectra are calculated using full renormalization group equations [60–64]. There is a theoretical uncertainty due to multiple RGE thresholds which mainly affects m_H , leading to a several GeV uncertainty. The gravitino mass is calculated by $m_{3/2} = \Lambda_S^2/M_P$, where M_P is the reduced Planck mass, assuming $\Lambda_S \sim \sqrt{F}$. There could be deviations from this relation in some SUSY breaking models which could lead to a gravitino mass that is different by up to a couple orders of magnitude (typically lower).

	1	2	3	4	5	6
Q_2	$-\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{4}$	$-\frac{1}{2}$	$-\frac{1}{2}$	$\frac{7}{8}$
Q_Q	$-\frac{1}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$	-2	-2	1
$g_{z'}$	0.45	0.23	0.23	0.06	0.04	0.055
λ	0.5	0.8	0.8	0.3	0.3	0.6
y_D	0.6	0.7	0.8	0.4	0.6	0.55
y_E	0.6	0.6	0.6	0.1	0.1	0.1
Λ_S	5×10^{10}	9×10^{10}	4×10^{10}	3×10^9	5×10^8	6×10^{11}
$\langle S \rangle$	2×10^5	7×10^4	6×10^4	2×10^5	8×10^4	2×10^4
$\tan\beta$	20	29	33	45	60	23
M_1	2700	735	650	760	270	185
M_2	710	195	180	340	123	178
M_3	4300	1200	1100	540	200	1040
m_H	140	140	140	140	140	140
$m_{\tilde{Q}_3}$	1×10^5	5×10^4	4×10^4	8×10^4	4×10^4	4×10^4
$m_{\tilde{L}_3}$	3×10^5	10^5	10^5	2×10^4	10^5	1.2×10^5
$m_{3/2}$	890	3600	810	3	0.1	10^5
$m_{\tilde{g}}$	4300	230	160	31	4	11
$M_{Z'}$	7×10^4	1.5×10^4	1.3×10^4	5600	2100	3400

mechanism. In Table II we give six benchmark points illustrating the possible variations in low-energy parameters for different charge choices and couplings. Point 6, which is the last column in the table, has to be interpreted

carefully. By itself this spectrum is inconsistent since such a large supersymmetry breaking scale $\Lambda_S = 6 \times 10^{11}$ GeV will induce gaugino masses much larger than the electroweak scale through gravity mediation. This conclusion may be evaded if some form of sequestering takes place, but we will not attempt such a construction here.

A. LHC phenomenology

1. Gluino

Since the colored scalars are all very heavy, the LHC will predominantly produce gluino pairs. The gluinos will consequently decay either through a 3-body off-shell squark ($\tilde{g} \rightarrow q\bar{q}\tilde{\chi}_i$, where $\tilde{\chi}_i$ is one of the gauginos) or a 2-body loop induced process. The 3-body decay usually dominates, leading to a gluino life time

$$\tau_3 = 4 \times 10^{-16} \text{ sec} \left(\frac{m_{\tilde{Q}}}{10^2 \text{ TeV}} \right)^4 \left(\frac{1 \text{ TeV}}{M_3} \right)^5 \propto \frac{1}{g_{z'}^6}. \quad (20)$$

Interestingly enough, the decay induced by the loop with exotic matter, shown in Fig. 6, can be the leading effect. We will discuss here mainly the processes which result in a singlino, as it usually has the largest coupling to the exotics. Two-body decays into other MSSM gaugino states will be somewhat suppressed (although they could be important in certain cases) and the Higgsino does not couple to the exotic sector directly. The expression for the decay width of $\tilde{g} \rightarrow \tilde{S}g$ can be extracted from [65] with the appropriate changes. The gluino decay corresponds to the dimension-5 operator $\tilde{\sigma}\sigma^{\mu\nu}\gamma_5\tilde{g}^a G_{\mu\nu}^a$, where the presence of γ_5 is due to the Majorana nature of the gluino and singlino. As pointed out in [48], this operator is P (and C)-odd and therefore must vanish in the limit where the left- and right-handed heavy scalars are degenerate. Indeed, the decay width for this channel is given by,

$$\Gamma_{\tilde{g} \rightarrow \tilde{S}g} = \frac{1}{8\pi} \frac{2g_S^4}{(32\pi^2)^2} \left(\frac{M_3^2 - m_{\tilde{S}}^2}{M_3} \right)^3 \times n_D^2 m_D^2 y_D^2 (C_0^L - C_0^R)^2, \quad (21)$$

where $C_0^{L,R}$ are the Passarino-Veltman functions [66], in-

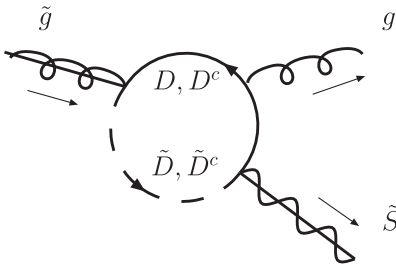


FIG. 6. The gluino can decay through colored exotic states into the singlino and a gluon. The other diagram in which the gluon is attached to the scalar propagator is suppressed. This decay channel can compete with the more standard decay of the gluino through off-shell squarks.

volving the left (right)-handed exotic scalars. They are given by (neglecting terms which are suppressed by ratios of the gaugino masses to the exotic matter mass),

$$C_0^{L,R} = \frac{m_{\tilde{D}_{L,R}}^2 - m_D^2 + m_{\tilde{D}_{L,R}}^2 \log(m_D^2/m_{\tilde{D}_{L,R}}^2)}{(m_{\tilde{D}_{L,R}}^2 - m_D^2)^2}, \quad (22)$$

where \tilde{D}_L (\tilde{D}_R) is the scalar component of D (D^c). Since these fields are chiral under $U(1)'$ they evolve differently under the RGE running and can differ significantly in mass. Parametrically, the 2-body channel leads to a life time (assuming no phase-space suppression),

$$\tau_2 \approx \frac{8}{n_D^2} 10^{-18} \text{ sec} \left(\frac{m_D}{10^2 \text{ TeV}} \right)^2 \left(\frac{1 \text{ TeV}}{M_3} \right)^3. \quad (23)$$

The exact value could be longer or shorter depending on the precise value of $C_0^{L,R}$. This analysis shows that it is potentially competitive with the standard 3-body mode and can lead to an interesting exotic decay of the gluino. In Table III we contrast the life time associated with the exotic 2-body mode versus the standard 3-body channel for the different benchmark points considered above. The relative branching ratio is very sensitive to the detailed model parameters. This is to be expected since the two-body width depends sensitively both on the mass splitting of left- and right-handed exotic scalars, as well as the mass of the exotic fermions. These quantities are in turn determined by charge assignments and exotic Yukawa couplings.

We remark here that the 2-body decay could give rise to very interesting collider signals if the singlino is not the LSP and decays subsequently (more on singlino decay in the next section).

It is interesting to compare the gluino decay signature in our case with that of the split-supersymmetry scenario. In split SUSY, the gluinos will also decay either through a 3-body off-shell squark ($\tilde{g} \rightarrow q\bar{q}\tilde{\chi}_i$, where $\tilde{\chi}_i$ is one of the gauginos) or a 2-body loop involving both the squarks and the standard model fermions [65,67]. A log enhancement of the 2-body channel associated with the third generation squark-quark loop, as well as the mixing between the LSP and the Higgsino, are important for the two-body decay to be comparable with the 3-body. Since in our case the Higgsinos are both very heavy, there is no such log enhancement and two-body decays are dominated by the exotic loop. Given that the two-body decay is induced by completely different virtual states, we expect the resulting branching ratio of $\tilde{g} \rightarrow g\tilde{S}$ will be quite different from that

TABLE III. The gluino life time (sec) for the 2-body channel versus the 3-body mode for the different benchmark points presented in Table II.

	1	2	3	4	5	6
τ_2	$9 \cdot 10^{-13}$	$8 \cdot 10^{-19}$	$6 \cdot 10^{-19}$	$6 \cdot 10^{-15}$	$5 \cdot 10^{-14}$	$3 \cdot 10^{-18}$
τ_3	$4 \cdot 10^{-19}$	$7 \cdot 10^{-18}$	$7 \cdot 10^{-18}$	10^{-16}	10^{-15}	$8 \cdot 10^{-18}$

of split supersymmetry. For example, for the squark masses in our scenario, the gluino life time is always too short to produce sizable displaced vertices. In the split-supersymmetry scenario considered in [65,67], the three body channel always dominates over the two-body one within the same range of squark masses, while the situation could be very different in our scenario.

2. The LSP and other inos

In general, the pattern of MSSM gaugino masses depends on both the charge assignments and the choice of the exotic sector. As a result of the absence of exotic doublets, which is a specific choice we made here, the W -ino is the lightest MSSM gaugino. The mass of the b -ino in our model is comparable to the gluino's and never serves as the LSP. The light inos include the W -ino, the singlino, and possibly the gravitino. As shown in the previous section and illustrated in Figs. 3–5, this model may admit different orderings of the light inos, and we discuss the different possibilities below.

The mass of the gravitino does not affect LHC phenomenology in this model. If the gravitino is not the LSP, it will not be produced at the LHC. At the same time, if it is the LSP, the range of gravitino mass implies that the next to lightest super-partner (NLSP) will decay outside the detector. Because of the decoupling of the scalars, the NLSP is neutral. Therefore, such decays will not be observable at the LHC.

The case of a singlino LSP with decoupled electroweak gauginos does not produce observable effects at the LHC either. In this case, the only way to produce the singlino is through the decay of the Z' . However, the decay mode will be dominated by $Z' \rightarrow \tilde{S} \tilde{S}$, which is again not observable.

There are several more interesting scenarios with either W -ino LSP or NLSP.

W-ino LSP only.—At tree level, the neutral and charged W -inos are degenerate, and the mass splitting induced by mixing with the Higgsinos is negligible for the large effective μ of this scenario. However, there is an important one-loop radiative correction which increases the charged W -ino mass by ~ 160 MeV with respect to the \tilde{W}^0 state [68–73]. This allows for the decay $\tilde{W}^+ \rightarrow \tilde{W}^0 + \pi^+$ with a life time around 1.4×10^{-10} sec, corresponding to a track length and displaced vertex around 4 cm from the production vertex in a detector, as has been studied extensively in connection with anomaly mediation [8,9].

W-ino NLSP and singlino LSP.—The W -ino can only decay to the singlino by mixing through the Higgsinos, leading to a suppression of the decay width. If there is no further phase-space suppression then the life time for $\tilde{W} \rightarrow h + \tilde{S}$ is approximately

$$\tau \sim \frac{4\pi}{g_W^2} \left(\frac{\langle S \rangle \tan\beta}{v} \right)^2 M_{\tilde{W}}^{-1} \sim 10^{-17} \text{ sec} \left(\frac{100 \text{ GeV}}{M_{\tilde{W}}} \right), \quad (24)$$

where the ratio of the singlet's vev to the electroweak scale

stems from the Higgsino-singlino mixing. Of course, the life time would be longer if there is phase-space suppression or the decay is via a virtual Higgs, and it is even possible in that case that there would be a displaced vertex.

Singlino NLSP and W-ino LSP.—The singlino decay into W -ino has a similar life time to the reversed process [with $M_{\tilde{W}}$ replaced by $m_{\tilde{S}}$ in Eq. (24)]. The singlino could be produced through Z' decay so this channel is potentially interesting and should be investigated further.

3. Z' production and decay

In some of the benchmarks we presented, the Z' is light enough to be produced at the LHC. This happens when $|Q_Q| \gtrsim 1$ and corresponds to the islands on the left and right in Fig. 2. In this case the normalization of $g_{Z'}$ becomes important since $\text{Tr} Q_i^2$ is larger, and $g_{Z'} \approx \frac{1}{10} g_Y$ and hence considerably smaller. This would normally be harmful, causing the W -ino to be too light. However, this is avoided here because the W -ino RGE has a term proportional to Q_Q^2 (and the other doublets' charges) and therefore the W -ino receives a large contribution as well. Together with the accidental tuning discussed above it is possible and even likely to have the Z' gauge boson in the observable spectrum as well as a very light singlino.

To have a light Z' gauge boson that is accessible at the LHC typically requires a smaller gauge coupling $g_{Z'}$. With a fixed spontaneous symmetry breaking scale, such a choice actually results in enhanced discovery potential at the LHC. Although the parton level total cross section is proportional to $g_{Z'}^2$, the parton distribution function depends inversely on a large power of $m_{Z'} \propto g_{Z'}$.

In Fig. 7 we plot the Z' production cross section times the leptonic branching ratio. If the Z' is not too heavy, $M_{Z'} < 4\text{--}5$ TeV it will likely be an easy task to observe

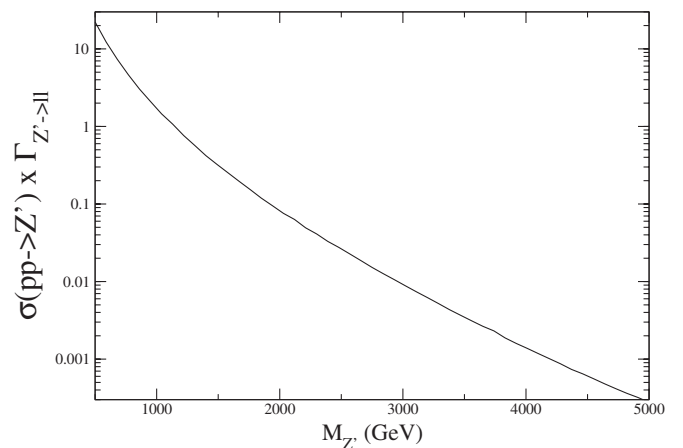


FIG. 7. A plot of the Z' gauge-boson production cross section times the leptonic branching ratio, which includes both $\mu^+ \mu^-$ and $e^+ e^-$ final states. The $U(1)'$ charge assignment used in generating this plot is $Q_2 = -1/2$ and $Q_Q = -2$ and the coupling was chosen nominally to be $g_{Z'} = 0.06$.

this resonance and determine its mass through its leptonic decay. Once its existence is established, it may be possible to uncover other and more difficult decay channels, such as $Z' \rightarrow \tilde{S} \tilde{S}$ etc. (for the possible utilization of a Z' in disentangling more difficult channels see Refs. [56,74]). A full discussion of the discovery reach and experimental challenges is beyond the scope of this paper, and we leave it for a future and more comprehensive study.

4. Higgs mass

At low energies there remains one light Higgs in the spectrum. Its mass is given as usual by $m_H^2 = 2\lambda_H v^2$, where $v = 174$ GeV and λ_H is the quartic coupling. The value of λ_H at low energies is determined by matching it to the supersymmetric contribution at $M_{\tilde{Z}'}$ and running it down to the electroweak scale,

$$\begin{aligned} 16\pi^2 \frac{d\lambda_H}{dt} &= 12(\lambda_H^2 + \lambda_H y_i^2 - y_i^4) \lambda_H (\mu \approx M_{\tilde{Z}'}) \\ &= \frac{1}{4}(g_2^2 + g_Y^2) + g_{Z'}^2 Q_2^2 + \frac{1}{2} \lambda^2 \sin^2 2\beta. \end{aligned} \quad (25)$$

The F -term contribution to the quartic, $\lambda^2 \sin^2 2\beta$, is negligible since $\tan\beta \gg 1$. The D -term contribution from the $U(1)'$ vector multiplet, $g_{Z'}^2 Q_2^2$, is usually smaller than the $SU(2) \times U(1)_Y$ D -term because both $g_{Z'}$ and Q_2 are not very large.

This leads to a prediction of the Higgs mass which is insensitive to the precise details of the high-energy parameters. It is predominantly affected by the running from $M_{\tilde{Z}'}$ down to the electroweak scale and yields,

$$m_H = 140 \text{ GeV} \quad (26)$$

with an uncertainty of a few percent coming from the precise matching and the value of $M_{\tilde{Z}'}$ (which we fixed at $M_{\tilde{Z}'} = 1000$ TeV for concreteness).

B. Cosmology

1. The W -ino

We have deliberately chosen the $U(1)'$ charges and exotics in our example construction to avoid a b -ino LSP. This is because the b -ino lacks any efficient annihilation or coannihilation mechanism for the large scalar masses and effective μ parameter favored in the scenario, leading to too much cold dark matter (CDM). (For a recent discussion, see, e.g., [75].) On the other hand, a W -ino LSP and its nearly degenerate charged partner, which have been studied extensively, especially in connection with anomaly mediated models [8,9], can annihilate efficiently into gauge bosons. In fact, for pure thermal production the CDM density is too low for the several hundred GeV mass range we have assumed, yielding [75]

$$\Omega h^2 \sim 0.021 \left(\frac{M_2}{1 \text{ TeV}} \right)^2, \quad (27)$$

compared to the observed value 0.111 ± 0.006 from WMAP (Wilkinson Microwave Anisotropy Probe) and galaxy surveys [76,77]. However, the CDM density can be considerably larger for nonstandard cosmological scenarios [68,78–84].

2. The gravitino

Another particle of interest to low-energy phenomena is the gravitino, with a mass given by

$$m_{3/2} \sim \frac{F}{k\sqrt{3}M_P} \sim \frac{2.4\text{eV}}{k} \left(\frac{\sqrt{F}}{100 \text{ TeV}} \right)^2, \quad (28)$$

where $M_P = 2.4 \times 10^{18}$ GeV is the reduced Planck mass and k depends on the details of the supersymmetry breaking mechanism in the hidden sector, but is typically ≤ 1 . We will take $\sqrt{F} = \Lambda_S$ and $k = 1$. The gravitino then depends very strongly on the SUSY breaking scale. The value of Λ_S does affect the other masses, because it determines the overall scale separation (recall that we tune Λ_S to obtain the EW scale while keeping $M_{\tilde{Z}'}$ fixed), but the dependence is only logarithmic. The symmetry breaking pattern in our model depends only logarithmically on the supersymmetry breaking scale. Therefore, the gravitino mass is exponentially sensitive to the choice of the charges and couplings, as shown in Fig. 8, and it may provide a sharp discriminator in the model space.

In this section, we will focus on cosmological implications and constraints on the gravitino mass (a good summary is given in [20]). However, one should keep in mind that these constraints are fairly indirect and can be overcome as mentioned below.

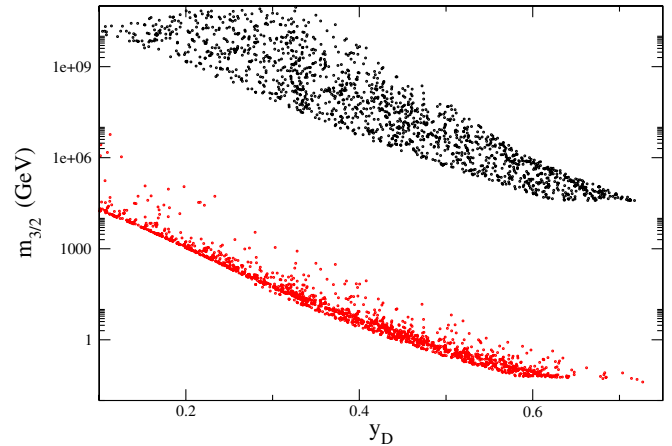


FIG. 8 (color online). A plot of the gravitino mass, $m_{3/2}$ as a function of the colored exotics' Yukawa coupling. The black (upper) points correspond to the charge assignment with $Q_Q = -1/3$ and $Q_2 = -1/4$, whereas the red (lower) points correspond to the assignment $Q_Q = -2$ and $Q_2 = -1/2$. The high value of the black points to the left are because a large Λ_S is needed to compensate the small Yukawa for those charges.

A stable (LSP) gravitino could overclose the universe unless it is lighter than a few keV (as in normal gauge mediation). However, this difficulty could be evaded if the reheating temperature T_R after inflation is rather low (i.e., $T_R \lesssim 10^8 m_{3/2}$).

The strongest constraint on gravitino mass comes from its interactions with other superpartners present in the early universe. Exact constraints on the parameter space are quite sensitive to the particle spectrum and interactions. A detailed study based on the spectrum, which is quite unique, is beyond the scope of this paper. In the following, we will very briefly summarize the results from early studies, see, for example, [85,86], comment on the relevance to our scenario, and point out cases where more careful studies need to be done.

Decay processes involving the gravitino are typically constrained by big bang nucleosynthesis (BBN), due to its long life time. If the decay products involve hadrons, any such decay with a life time longer than 1 second is strongly constrained by deuterium and helium abundances. On the other hand, if the decay process only induces electromagnetic showers then only a life time longer than 10^4 seconds is strongly constrained.

We begin with the case in which the gravitino is not the LSP. Such gravitinos will be produced during reheating. The gravitino decay into the LSP may lead to unacceptable modifications of BBN if there is any significant component of hadrons in the decay. Such effects have been studied carefully, e.g., in [86]. One way around this is for the gravitino life time to be shorter than ~ 1 s.⁶ This typically requires $m_{3/2} \gtrsim 10$ TeV.⁷ Alternatively, the BBN constraints can be satisfied for a relatively low reheating temperature, $T_R < 10^6 - 10^7$ GeV, suppressing the gravitino production. The only difference in our case from those well-studied scenarios is the decoupled sfermions. It is expected to affect more significantly the case where the gravitino mass is heavier than the gluino mass. The enhanced branching ratio of this channel in the absence of sfermions makes the constraint on the reheating temperature slightly stronger [86].

Alternatively, the gravitino could be the LSP. In this case the analysis becomes more complicated since the constraints depend on the identity of the NLSP. A scenario with a W -ino as the NLSP is similar to one with a b -ino

⁶If the decay products only contain photons, BBN constraints could be easily satisfied if the life time is less than 10^4 sec, which corresponds to a gravitino mass of about 1 TeV. However, in our examples, we always have a light W -ino. Therefore, $\tilde{G} \rightarrow W/Z + \tilde{W}$ will usually lead to hadrons.

⁷This requires that the gravity mediation effects are sequestered. In this case one would again have contributions from the anomaly mediation [8,9] to the standard model gauginos comparable to those of the two-loop Z' mediation, while the anomaly mediation would be irrelevant for the other soft parameters. This hybrid scenario could also use the mechanism of [78] to increase the CDM density due to the gravitino decay into the LSP W -ino.

although the numerics are different because of the smaller branching ratio into photons and the larger annihilation cross section during freeze out [85,87]. (For the b -ino case one requires the gravitino to be lighter than about 100 MeV.) A singlino as the NLSP is even less favorable because its decay into the gravitino must involve mixing with the Higgsino states, which leads to a suppression of the decay width. Furthermore, the life time is very sensitive to the precise value of \sqrt{F} , which we do not have a precise prediction for. All together the life time is generically much longer than a second regardless of the precise decay mode. It seems that having a singlino as the NLSP with appreciable density is pretty much ruled out. The singlino is expected to be produced in the thermal soup since it couples to the Z' . Therefore, it is hard to see how to make this case viable, without resorting to more exotic cosmologies with large late-time entropy production, such as thermal inflation [88–90].

V. COMMENTS AND ALTERNATIVES

A. Other possibilities of Z' -mediation

In this paper, we have focused on a particular scenario of Z' -mediation. Motivated by solving the μ -problem, we have considered a singlet-extended MSSM with a PQ-like Z' . More generally, there are of course many other possibilities of Z' which can play the role of the mediator of supersymmetry breaking, such as $B - L$, or any other well-studied or yet unknown exotic $U(1)$. As we have demonstrated in the examples presented in this paper, the detailed spectrum from Z' -mediation depends quite sensitively on the choice of model. However, we would like to emphasize that the sizable hierarchy between the scalars and the electroweak-inos will be a very generic feature of the Z' mediation.

It is of course possible to combine other mediation mechanisms with the Z' mediation. In those scenarios, we generically expect that the Z' mediation contribution to the electroweak-ino masses will be negligible, while the contribution to the scalar masses will be significant. The challenge of such scenarios is to give plausible reasons to why some other mediation mechanism will give comparable contributions as the Z' -mediation. Recently, one scenario of such a combination with anomaly mediation and a hypercharge mediation has been studied [91], and a combination with D -term mediation in [92]. Further studies on other possibilities for combining Z' -mediation with other mechanisms are certainly interesting and worth pursuing.

B. An alternative model of neutrino masses

$U(1)'$ models usually do not allow the large Majorana masses necessary for the canonical seesaw model [93]. The specific model constructed in Sec. III allows Dirac masses by assumption, which would have to be made small by fine-tuning. However, in a simple variant,⁸ the $U(1)'$ sym-

metry forbids Dirac Yukawa couplings $y_\nu H_u L \nu^c$ at the renormalizable level, but allows them to be generated by a higher-dimensional operator,

$$W_\nu = c_\nu \frac{S}{M_P} H_u L \nu^c. \quad (29)$$

This naturally yields small Dirac neutrino masses of order $(0.01 c_\nu)$ eV for $S = 100$ TeV, in accordance with observation. (This mechanism has been studied previously in a more general context [94].) One cannot say more about the hierarchy of neutrino masses or mixings without additional assumptions.

C. Exotics and R -parity

Exotic particles are necessary for anomaly cancellation in most $U(1)'$ models. These are usually nonchiral under the standard model gauge group, but chiral under $U(1)'$. As discussed in Sec. II our scenario typically involves exotic chiral supermultiplets with supersymmetric masses in the 10–100 TeV range, such as the $D + D^c$ quark pairs or $E + E^c$ lepton pairs in the model of Sec. III. Our focus is not on the specific model, but rather on the general Z' -mediation scenario, so we will mainly comment on the more general case.

There are several possibilities for the lightest exotic scalar or fermion of a given type⁹: (a) One is that it is absolutely stable. This possibility is severely constrained by cosmology and by direct searches for heavy stable particles. However, it would be viable if the reheating temperature after inflation was sufficiently low [82], i.e. $T_R < 20\text{--}200$ GeV for an exotic mass in the 10–100 TeV range [96]. (b) The most commonly studied case, especially for nonsupersymmetric models, is that the exotic decays by mixing with ordinary quarks and leptons, allowing decays such as $D \rightarrow (dZ, uW, dH)$ [97,98]. However, such mixings are often forbidden in supersymmetric $U(1)'$ models, at least at the renormalizable level, by $U(1)'$ and/or R -parity conservation. For example, in the specific models in Sec. III there are no allowed renormalizable-level operators that could lead to $D - d$ mixing. However, $E - e$ mixing could be induced by a nonholomorphic soft operator $LH_u^* E^c$ or a bilinear Ee^c , if present, for the $Q_Q = -Q_1/3$ model, or by LLE^c or $Ee^c \nu^c$ operators for $Q_Q = (Q_2 - Q_1)/6$. The latter case would require spontaneous R -parity violation via the vevs of a scalar ν or ν^c . (c) Another possibility is the existence of renormalizable-level couplings allowing the direct decay of an exotic into ordinary particles, such as the leptoquark (diquark) cou-

plings $Du^c e^c$ ($D^c u^c d^c$) [95]. One or the other could be present without inducing proton decay, and they would still allow a stable LSP (the exotic scalar would be the normal particle). No such D couplings are allowed in the models in Sec. III, but analogous couplings for the E or E^c (listed above in connection with mixing) could allow the rapid decays of E and E^c .¹⁰ (d) Finally, exotics could decay by higher-dimensional operators, analogous to (29), which could induce highly suppressed mixing with the ordinary particles or lead directly to the decays. They would therefore be stable on collider time scales, leading to exiting tracks or delayed decays in the detector [95]. Only dimension-5 operators would decay fast enough to satisfy constraints from big bang nucleosynthesis [95,99]. The only example in the models considered here is $LH_d E^c S/M_P$, occurring in the $Q_Q = -Q_1/3$ case.

Thus, the lightest D fermion or scalar would be stable in the specific models of Secs. III or VB, which is unacceptable unless T_R is very low. However, such operators can always be allowed for both the D - and E -type exotics by extending the particle content to include nonchiral exotics which obtain vevs. We emphasize, however, that these models are only examples of a general scenario.

Finally, we comment briefly on R -parity, which is frequently guaranteed by $U(1)'$ invariance [55]. In the present case, the operators $S^n LH_u$, $S^n LLe^c$, $S^n u^c d^c d^c$, and $S^n QLd^c$, $n \geq 0$ are forbidden for the specific models considered in Sec. III by the $U(1)'$ symmetry, so there is an automatic R -parity in the Lagrangian, even after $U(1)'$ breaking. The alternative model in Sec. VB with $Q_Q = -Q_1 - 2Q_2/3$ would allow the R -parity violating operator $Su^c d^c d^c/M_P$. This operator would lead to LSP decay, but with a life time much larger than (comparable to) the age of the universe for a W -ino (b -ino) LSP.

D. Gauge unification

We comment briefly on gauge unification for the standard model couplings. The successful unification in the MSSM is modified in the specific model considered here by the large Higgsino scale and (especially) by the exotics. (This would be less of a problem in the E_6 motivated models, which, however, lead to a b -ino LSP.) The gauge unification could easily be restored by additional nonchiral exotics, which could also lead to $\text{Tr}(QY) = 0$ at a high scale, and possibly by a noncanonical normalization of the $U(1)_Y$ coupling [100], which occurs frequently in string constructions. As an example, approximate gauge unification at around 3×10^{15} GeV would be achieved by the addition of four pairs of $SU(2)$ doublets with $Y = 0$ at around 2×10^{10} GeV. (These fractional charged states could be confined at that scale.) We reemphasize that these issues are very dependent on the specific model.

⁸The anomaly conditions in this variant require 4 singlets S and the ν^c charge inferred from (29). The other conditions are unchanged except for the form of the discriminant in (A10), which we will not display. The variant discriminant vanishes for $Q_Q = Q_2/3$ or for $Q_Q = -Q_1 - 2Q_2/3$.

⁹For a recent general discussion, see [95].

¹⁰The alternative models for a small neutrino mass do not allow either possibility (b) or (c).

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APPENDIX A: CHARGES AND ANOMALY CANCELLATION

In order to generate masses for all of the scalars, we assume that all of the visible sector chiral superfields, including the singlet, are charged under $U(1)'$. We also assume that there is only one singlet field S , and that there are no exotic $SU(2)$ doublets or additional Higgs pairs. The charge assignments are constrained by the requirements of family universality, anomaly cancellation, and that the superpotential terms in (1) and (18) are allowed. The superpotential condition implies

$$\left. \begin{aligned} Q_2 + Q_Q + Q_{u^c} &= 0 \\ Q_1 + Q_Q + Q_{d^c} &= 0 \\ Q_2 + Q_L + Q_{\nu^c} &= 0 \\ Q_1 + Q_L + Q_{e^c} &= 0 \end{aligned} \right\} \text{Yukawa couplings,} \quad (\text{A1})$$

$$\left. \begin{aligned} Q_S + Q_D + Q_{D^c} &= 0 \\ Q_S + Q_E + Q_{E^c} &= 0 \end{aligned} \right\} \text{exotics couplings,} \quad (\text{A2})$$

$$Q_1 + Q_2 + Q_S = 0 \} \text{ singlet coupling.} \quad (\text{A3})$$

Based on the choice of exotics in the model in Sec. III, the anomaly cancellation conditions lead to the following constraints:

$SU_C(3)^2 \times U(1)'$ anomaly cancellation:

$$n_D = 3. \quad (\text{A4})$$

$SU_L(2)^2 \times U(1)'$ anomaly cancellation:

$$Q_L = -3Q_Q - \frac{1}{3}(Q_1 + Q_2). \quad (\text{A5})$$

$U(1)'$ gravitational anomaly cancellation:

$$n_E = 2. \quad (\text{A6})$$

$U(1)_Y^2 \times U(1)'$ anomaly cancellation:

$$9Y_D^2 + 2Y_E^2 = 3, \quad (\text{A7})$$

where $Y_D = -Y_{D^c}$ and $Y_E = -Y_{E^c}$ are the hypercharges of D and E . We will choose the hypercharges in analogy with the SM, $Y_D = -1/3$ and $Y_E = -1$.

$U(1)_Y \times U(1)'^2$ anomaly cancellation:

$$Q_E = -3Q_Q - \frac{3}{2}Q_D + 2Q_1. \quad (\text{A8})$$

$U(1)^3$ anomaly cancellation:

$$\begin{aligned} 81Q_D^2 - 36Q_D(3Q_1 + Q_2 - 3Q_Q) \\ + 4(7Q_1^2 + 8Q_1Q_2 + Q_2^2 - 36Q_1Q_Q - 27Q_Q^2) &= 0. \end{aligned} \quad (\text{A9})$$

There are two possible choices for Q_D as solutions to the quadratic equation,

$$\begin{aligned} Q_D &= \frac{2}{9}(3Q_1 + Q_2 - 3Q_Q) \\ &\pm \sqrt{2(Q_1 + 3Q_Q)(Q_1 - Q_2 + 6Q_Q)}. \end{aligned} \quad (\text{A10})$$

These correspond to two 2-parameter solutions in terms of Q_2/Q_1 and Q_Q/Q_1 , with the other charges obtained from the previous constraints. Two simplified 1-parameter solutions are obtained by requiring the discriminant to vanish. We will mainly consider the case

$$Q_Q = -\frac{1}{3}Q_1, \quad (\text{A11})$$

and normalize $Q_1 = 1$, so the other charges are all determined by Q_2 , as listed in Table I.

We note that

$$\text{Tr}(QY) = 14Q_2 - 8Q_1 + 36Q_Q \quad (\text{A12})$$

does not vanish in general, and not for the special 1-parameter solutions. However, the vanishing can be restored by the addition of nonchiral states. These do not affect the anomaly conditions and can also restore gauge unification.

APPENDIX B: RENORMALIZATION GROUP EQUATIONS

In calculating the various masses we distinguish between two regions: $M_{Z'} < \mu < \Lambda_S$ and $\mu < M_{Z'}$. We use $t = \log(\mu/\Lambda_S)$.

1. $M_{Z'} < \mu < \Lambda_S$

For this region we use the RGEs given in [60–62]. To calculate the spectrum we need the one-loop RGEs for the gauge and Yukawa couplings, \tilde{Z}' , soft scalar masses, and the A terms, as well as the two-loop RGEs for the gaugino masses.

Using $SU(5)$ normalization ($g_1^2 = 5g_Y^2/3$), the one loop $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge-couplings RGEs are given by

$$\frac{dg_a}{dt} = \frac{g_a^3}{16\pi^2} b_a \quad (\text{B1})$$

where $b_a = (51/5, 1, 0)$ for $a = 1, 2, 3$. The $U(1)'$ gauge-coupling RGE is given by

$$\frac{dg_{z'}}{dt} = \frac{g_{z'}^3}{16\pi^2} \text{Tr} Q_i^2. \quad (\text{B2})$$

Keeping only the dominant terms proportional to $M_{Z'}$, the two loop $SU(3)_C \times SU(2)_L \times U(1)_Y$ gaugino RGEs are

$$\frac{d\tilde{M}_a}{dt} = \frac{4g_a^2 c_a}{(16\pi^2)^2} g_{z'}^2 M_{Z'}, \quad (\text{B3})$$

with

$$c_1 = \frac{6}{5} \sum_{\text{all scalars}} Q_i^2 Y_i^2, \quad c_2 = 9Q_Q^2 + 3Q_L^2 + Q_1^2 + Q_2^2, \\ c_3 = 3(2Q_Q^2 + Q_{u^c}^2 + Q_{d^c}^2) + n_D(Q_D^2 + Q_{D^c}^2). \quad (\text{B4})$$

The $U(1)'$ gaugino RGE is at one loop,

$$\frac{dM_{Z'}}{dt} = \frac{g_{z'}^2}{8\pi^2} M_{Z'} \text{Tr} Q_i^2. \quad (\text{B5})$$

Within these approximations it is easy to solve analytically the gaugino RGEs.

With the obvious definitions of the A terms [see below (12)] their RGEs are

$$16\pi^2 \frac{dA_D}{dt} = 4g_{z'}^2 y_D (Q_D^2 + Q_{D^c}^2 + Q_S^2) M_{Z'}, \\ 16\pi^2 \frac{dA_E}{dt} = 4g_{z'}^2 y_E (Q_E^2 + Q_{E^c}^2 + Q_S^2) M_{Z'}, \quad (\text{B6}) \\ 16\pi^2 \frac{dA_H}{dt} = 4g_{z'}^2 \lambda (Q_{H_u}^2 + Q_{H_d}^2 + Q_S^2) M_{Z'},$$

where we have neglected all the terms on the right-hand side that are not proportional to $M_{Z'}$.

The RGEs for the soft masses are

$$16\pi^2 \frac{dm_S^2}{dt} = -8g_{z'}^2 Q_S^2 M_{Z'}^2 + 4\lambda^2 (m_S^2 + m_{H_u}^2 + m_{H_d}^2) + 6n_D y_D^2 (m_S^2 + m_D^2 + m_{D^c}^2) + 2n_E y_E^2 (m_S^2 + m_E^2 + m_{E^c}^2), \\ 16\pi^2 \frac{dm_D^2}{dt} = -8g_{z'}^2 Q_D^2 M_{Z'}^2 + 2y_D^2 (m_S^2 + m_D^2 + m_{D^c}^2), \\ 16\pi^2 \frac{dm_{D^c}^2}{dt} = -8g_{z'}^2 Q_{D^c}^2 M_{Z'}^2 + 2y_D^2 (m_S^2 + m_D^2 + m_{D^c}^2), \\ 16\pi^2 \frac{dm_E^2}{dt} = -8g_{z'}^2 Q_E^2 M_{Z'}^2 + 2y_E^2 (m_S^2 + m_E^2 + m_{E^c}^2), \\ 16\pi^2 \frac{dm_{E^c}^2}{dt} = -8g_{z'}^2 Q_{E^c}^2 M_{Z'}^2 + 2y_{E^c}^2 (m_S^2 + m_E^2 + m_{E^c}^2), \\ 16\pi^2 \frac{dm_{H_u}^2}{dt} = -8g_{z'}^2 Q_2^2 M_{Z'}^2 + 2\lambda^2 (m_S^2 + m_{H_u}^2 + m_{H_d}^2) + 6Y_u^2 (m_{H_u}^2 + m_{Q_3}^2 + m_{u_3^c}^2), \quad (\text{B7}) \\ 16\pi^2 \frac{dm_{H_d}^2}{dt} = -8g_{z'}^2 Q_1^2 M_{Z'}^2 + 2\lambda^2 (m_S^2 + m_{H_u}^2 + m_{H_d}^2) + 6y_d^2 (m_{H_d}^2 + m_{Q_3}^2 + m_{d_3^c}^2), \\ 16\pi^2 \frac{dm_{Q_3}^2}{dt} = -8g_{z'}^2 Q_Q^2 M_{Z'}^2 + 2y_u^2 (m_{H_u}^2 + m_{Q_3}^2 + m_{u_3^c}^2) + 2y_d^2 (m_{H_d}^2 + m_{Q_3}^2 + m_{d_3^c}^2), \\ 16\pi^2 \frac{dm_{u_3^c}^2}{dt} = -8g_{z'}^2 Q_{u^c}^2 M_{Z'}^2 + 4y_u^2 (m_{H_u}^2 + m_{Q_3}^2 + m_{u_3^c}^2), \\ 16\pi^2 \frac{dm_{d_3^c}^2}{dt} = -8g_{z'}^2 Q_{d^c}^2 M_{Z'}^2 + 4y_d^2 (m_{H_d}^2 + m_{Q_3}^2 + m_{d_3^c}^2),$$

where we have ignored the (small) A term contributions on the right-hand side.¹¹

The one-loop RGEs for the superpotential couplings are

¹¹In general, there are also $U(1)_Y$ and $U(1)'$ D -term contributions to the scalar RGEs, which are of the form $\text{Tr}(Y m_i^2)$ and $\text{Tr}(Q_i m_i^2)$. To order $\mathcal{O}(g_{z'}^2)$ the contributions from the visible sector fields vanish in our scenario. Being the sum of scalar masses they vanish at the boundary $\mu = \Lambda_S$ like all the scalar masses. The RGE for this sum of masses is easily shown to be proportional to the sum itself by making use of the anomaly cancellation conditions on the charges. Being a homogeneous equation with vanishing boundary condition, the solution must vanish everywhere. Nonchiral hidden sector fields Ψ and Ψ^c could in principle yield nonvanishing $U(1)'$ D -term contributions, but only if their soft mass-squares are unequal. Such effects would be of the same order as the Z' contributions to the scalar masses.

$$\begin{aligned}
16\pi^2 \frac{d\lambda}{dt} &= \lambda \left[4\lambda^2 + 3n_D y_D^2 + n_E y_E^2 + 3y_u^2 + 3y_d^2 \right. \\
&\quad \left. - 3g_2^2 - \frac{3}{5}g_1^2 - 2g_{z'}^2(Q_S^2 + Q_1^2 + Q_2^2) \right], \\
16\pi^2 \frac{dy_D}{dt} &= y_D \left[2\lambda^2 + (3n_D + 2)y_D^2 + n_E y_E^2 - \frac{16}{3}g_3^2 \right. \\
&\quad \left. - \frac{6}{5}g_1^2(Y_D^2 + Y_{D^c}^2) - 2g_{z'}^2(Q_S^2 + Q_D^2 + Q_{D^c}^2) \right], \\
16\pi^2 \frac{dy_E}{dt} &= y_E \left[2\lambda^2 + 3n_D y_D^2 + (n_E + 2)y_E^2 - \frac{6}{5}g_1^2(Y_E^2 \right. \\
&\quad \left. + Y_{E^c}^2) - 2g_{z'}^2(Q_S^2 + Q_E^2 + Q_{E^c}^2) \right], \quad (\text{B8})
\end{aligned}$$

with similar expressions for y_u , y_d , and y_e . In practice, we ignored the relatively small running effects of y_E and y_e .

To obtain the Higgs potential at the electroweak scale one must run down below $M_{Z'}$. In doing so, one encounters several heavy thresholds. First, one must integrate out the Z' -ino and then the squarks, sleptons, Higgsinos, and exotics one by one. For simplicity and since these masses are not greatly separated from $M_{Z'}$ we will ignore the running between these scales.¹²

2. $\mu < M_{Z'}$

Below the mass scale of the scalars the Higgs mass and quartic's RGEs are those of the standard model. (Unlike [47], there are no contributions from the Higgs-gaugino-Higgsino couplings in the low-energy theory. Using the standard form of the Higgs potential, $m_H^2 \phi^\dagger \phi + \lambda_H (\phi^\dagger \phi)^2/2$, we have [63],

¹²This and other approximations we have made throughout, while small compared to the terms retained, are not negligible compared with the electroweak scale. However, the fine-tuning needed to obtain the electroweak scale is not restricted to a very small range of parameter space, so the approximations can be compensated by small changes in the values of the parameters such as λ or exotic Yukawa couplings

$$\begin{aligned}
16\pi^2 \frac{d\lambda_H}{dt} &= 12(\lambda_H^2 + \lambda_H y_t^2 - y_t^4), \\
16\pi^2 \frac{dm_H^2}{dt} &= 6m_H^2(\lambda_H + y_t^2), \quad (\text{B9})
\end{aligned}$$

where y_t is the top Yukawa and we have neglected other smaller contributions.

The RGE for y_t is [64]

$$16\pi^2 \frac{dy_t}{dt} = \frac{9}{2}y_t^3 - y_t \left(\frac{17}{20}g_1^2 + \frac{9}{4}g_2^2 + 8g_3^2 \right). \quad (\text{B10})$$

We must also specify the matching conditions in passing from the high-energy effective theory containing the Z' -ino and scalars to the low-energy theory with only SM fields and gauginos. The Higgs mass receives a quadratically divergent threshold correction from integrating out the squarks. The quartic coupling receives a contribution from the different D -terms as well as a contribution from an F -term,

$$\begin{aligned}
\lambda_H(\mu \approx m_{\varphi_i}) &= \frac{1}{4}(g_2^2 + g_{Y'}^2) + g_{z'}^2 Q_2^2 + \frac{1}{2}\lambda^2 \sin^2 2\beta, \\
m_H^2(\mu \approx m_{\varphi_i}) &= \min(\mathcal{M}_H^2) - \frac{3y_t^2}{16\pi^2} m_{\varphi_i}^2. \quad (\text{B11})
\end{aligned}$$

Notice that the F -term contribution is small for large $\tan\beta$.

The gauge-coupling RGEs in this region are

$$\frac{dg_a}{dt} = \frac{g_a^3}{16\pi^2} b_a, \quad (\text{B12})$$

where $b_a = (41/10, -11/6, -5)$. We do not run $g_{z'}$.

The one-loop gauginos RGEs in this region are

$$\frac{d\tilde{M}_a}{dt} = \frac{g_a^2}{16\pi^2} \tilde{M}_a c_a, \quad (\text{B13})$$

where $c_a = (0, -12, -18)$.

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