Viable gravity-mediated supersymmetry breaking

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We present a complete, viable model of gravity-mediated supersymmetry breaking that is safe from all flavor constraints. The central new idea is to employ a supersymmetry breaking sector without singlets, but with D-terms comparable to F-terms, causing supersymmetry breaking to be dominantly communicated through $U(1)_R$-symmetric operators. We construct a visible sector that is an extension of the minimal supersymmetric standard model where an *accidental* $U(1)_R$-symmetry emerges naturally. Gauginos acquire Dirac masses from gravity-mediated D-terms and tiny Majorana masses from anomaly-mediated contributions. Contributions to soft breaking scalar (mass)^2 arise from flavor-arbitrary gravity-induced F-terms plus one-loop finite flavor-blind contributions from Dirac gaugino masses. Renormalization group evolution of the gluino causes it to naturally increase nearly an order of magnitude larger than the squark masses. This hierarchy, combined with an accidentally $U(1)_R$-symmetric visible sector, nearly eliminates all flavor violation constraints on the model. If we also freely tune couplings and phases within the modest range 0.1–1, while maintaining nearly flavor-anarchic Planck-suppression contributions, we find our model to be safe from Δm_K, ϵ_K, and $\mu \to e$ lepton flavor violation. Dangerous $U(1)_R$-violating Kähler operators in the Higgs sector are eliminated through a new gauged $U(1)_X$ symmetry that is spontaneously broken with electroweak symmetry breaking. Kinetic mixing between $U(1)_X$ and $U(1)_Y$ is present with loop-suppressed (but log-enhanced) size ϵ. The Z' associated with this $U(1)_X$ has very peculiar couplings—it has order one strength to Higgs doublets and approximately ϵ strength to hypercharge. The Z' could be remarkably light and yet have escaped direct and indirect detection.

I. INTRODUCTION

If supersymmetry plays a role in the physics near the electroweak scale, the most pressing question is how supersymmetry breaking is mediated to the "visible" supersymmetrized standard model. The nature of this mediation makes a qualitative impact on the superpartner spectrum and interactions, thereby being directly relevant to the physics probed by experiments including the Large Hadron Collider (LHC).

Gravity mediation [1–9] is far and away the simplest mediation mechanism. The hidden and visible sectors interact only through Planck-suppressed nonrenormalizable operators. Gaugino and scalar superpartners acquire comparable masses. Moreover, unlike low scale mediation, the μ and B_μ terms in the Higgsino and Higgs sector can easily acquire comparable sizes [10]. Also, with the gravitino comparable to other soft supersymmetry breaking masses, a neutral superpartner of the visible sector can be the lightest supersymmetric particle, serving as a dark matter candidate. Last, but not least, gauge coupling unification [11,12] is automatic.

Gravity mediation does, however, have a serious flavor problem. Since the mediation of supersymmetry breaking is at the Planck scale, this is necessarily above (or at) the scale of generating flavor. No symmetries enforce flavor universality in the operators mediating supersymmetry breaking. Thus, there is no mechanism to ensure the mass-squared matrices of squarks and sleptons to be almost exactly proportional to an identity matrix for consistency with constraints on flavor violation beyond the standard model [13–16]. Even though models with continuous flavor symmetries are proposed to avoid the flavor problem [17,18], it is unlikely that gravitational interactions respect any global symmetry [19]. Models in which squark and slepton squared-mass matrices are aligned with the corresponding Yukawa matrices [20,21] because of gauged discrete family symmetries [22,23] do solve the flavor problem but these implementations are rather complicated.

In this paper we present a complete model of supersymmetry breaking with gravity mediation and a visible sector that is safe from flavor constraints. We exploit the observation that flavor violation beyond the standard model can be greatly suppressed *even without flavor universality of the squark and slepton sectors* if the visible sector particle content respects an approximate global $U(1)_R$ symmetry [24]. For us, the origin of the approximate $U(1)_R$ symmetry is *accidental*. We do not enforce $U(1)_R$ symmetry *per se*, but arrange the hidden and visible sectors such that the only potentially dangerous $U(1)_R$-violation arises as a result of anomaly mediation (which, nevertheless, is small enough to maintain flavor safety). Getting this to work with gravity mediation—namely, sizeable D-terms, F-terms, a modest hierarchy between the Dirac gluino and squark

masses, not generating sizeable Majorana masses, but generating approximately R-symmetric Higgsino masses is the subject of the paper. Our philosophy is to exploit holomorphy of the superpotential (any terms we do not need or want can be naturally set to zero), but to strictly eliminate dangerous Kähler terms through symmetries. As a consequence, we are inevitably led to gauging an additional $U(1)_X$ symmetry in the visible sector. The phenomenology of this $U(1)_X$ is quite interesting: the Z' of this broken $U(1)_X$ couples with order one strength to the Higgs sector and order $10^{-4} \rightarrow 10^{-2}$ strength to standard model fields.

Our paper is organized as follows. In Sec. II, we outline the basic philosophy regarding the hidden sector, the visible sector, the mediation, the sizes of dimensionless constants, and so forth. In Sec. III we present the supersymmetry breaking sector—the "4-1" model—and its relevant characteristics to our framework. In Sec. IV we discuss the full set of operators that can mediate supersymmetry breaking up to dimension-6. This provides the basis to understanding the content and supersymmetry breaking induced into the visible sector, discussed in Sec. V. It is here where we show how the visible sector can remain approximately R-symmetric despite the spontaneous (and explicit) R-violation of the hidden sector. We then outline the main phenomenological implications of our model in Sec. VI. The gaugino and scalar spectrum are discussed in the context of a benchmark point, illustrating that all flavor constraints can be satisfied. Discussion of the Higgs sector, as well as the $U(1)_X$ sector phenomenology is begun. Finally, we present a brief discussion of unification in Sec. VII and conclude with a more general discussion in Sec. VIII.

II. PHILOSOPHY

The basic philosophy behind our model-building efforts can be summarized as follows:

   (i)   The supersymmetry breaking hidden sector is specified: we employ the $SU(4) \times U(1)$ (4-1) model [25]. This model has two virtues for our proposal: First, it contains no singlet fields. Second, it contains a $U(1)$ that acquires a D-term comparable to the F-terms.
   (ii) Supersymmetry breaking is only mediated by Planck-suppressed operators involving 4-1 fields and visible sector fields.
   (iii) "Chiral" Kähler potential[1] and all superpotential operators are present or absent as we deem appropriate for the model. We exploit the technical naturalness of these operators to eliminate some otherwise potentially dangerous mediation operators.

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[1] By chiral Kähler we refer to the operators that appear as $\int d^4\theta f$ in the global supersymmetric limit, where f are polynomials of only chiral or antichiral superfields.

   (iv) All "nonchiral" operators in the Kähler potential allowed by symmetries are assumed to be present.
   (v)   Nonzero Planck-suppressed operators, and dimensionless couplings, have coefficients between about 0.1 to 1. We allow ourselves complete freedom in setting the sizes of different operators with coefficients in this range.

The 4-1 model is one of several examples of global supersymmetry breaking models that has a D-term generated with size of order an F-term [25–27]. For us, the crucial feature of this model is the absence of hidden sector singlets. Generically, at the local supersymmetry breaking minimum, the hidden sector fields acquire scalar and supersymmetry breaking expectation values that spontaneously break the $U(1)_R$ symmetry, in accord with Ref. [28]. Operators mediating supersymmetry breaking involving supersymmetry breaking spurions include one or more powers of the D-term, which automatically preserves $U(1)_R$, and F-terms, which must arise in gauge-invariant combinations. For us, the absence of singlets implies the lowest dimension gauge invariants are also $U(1)_R$ invariant (of the form $X^\dagger X$, where X is a hidden sector chiral superfield, as we will see). Hence, despite the sizeable spontaneous breaking of $U(1)_R$ in the hidden sector, the lowest dimension operators mediating supersymmetry breaking do so through $U(1)_R$-symmetric combinations of hidden sector fields. This observation was also exploited in a different context in Ref. [29].

Embedding a global supersymmetry breaking model into supergravity also requires adding a constant term to the superpotential, explicitly breaking $U(1)_R$, to fine-tune the cosmological constant to zero. The explicit breaking gives a mass to the would-be massless R-axion [30]. It also leads to the gravitino acquiring a Majorana mass, absorbing the would-be massless goldstino from supersymmetry breaking. The visible sector is thus not completely isolated from explicit $U(1)_R$ breaking, due to the anomaly-mediated contribution to visible sector fields proportional to the gravitino mass [31,32].

One might think that the approach should be to simply demand an approximate $U(1)_R$ symmetry on the visible sector by hand. But one could also hold the same viewpoint for flavor universality; if we demand the $U(3)^5$ flavor symmetry on the soft masses in isolation, flavor universality would be approximately maintained even though the Yukawa couplings violate $U(3)^5$. Therefore, if we wish to solve the flavor problem without appeal to an "initial condition" at the Planck scale, we must seek a mechanism where an approximate $U(1)_R$ symmetry in the visible sector emerges. Only then can we compare the $U(1)_R$-symmetric solution to the flavor problem with e.g. gauge mediation [25,33–40], gaugino mediation [41,42], or anomaly mediation [31,32], in which the $U(3)^5$-preserving soft masses emerge in the presence of the $U(3)^5$-violating Yukawa couplings.

III. SUPERSYMMETRY BREAKING

We employ the 4-1 model for our supersymmetry breaking sector. The 4-1 model is an SU(4) ⊗ U(1) supersymmetric gauge theory with four chiral superfields transforming under SU(4) ⊗ U(1) as

$$
\begin{array}{ccc}
\text{fields} & \text{SU(4)} & \text{U(1)} \\
X_1 & 4 & -3 \\
X_2 & \bar{4} & -1 \\
X_3 & 6 & 2 \\
X_4 & 1 & 4
\end{array}
\tag{1}
$$

At high energies, the theory is weakly coupled and the dimensions of all the chiral superfields are close to their classical dimensions. The field content allows a unique renormalizable superpotential $\lambda X_1 X_2 X_4$ permitted under the symmetries. At lower energies, SU(4) gets strong at a scale Λ_4. The dynamically generated nonperturbative superpotential consistent with all the symmetries is unique,

$$
W_{\mathrm{np}} = \frac{\Lambda_4^5}{\sqrt{X_1 X_2 X_3^2}}.
\tag{2}
$$

In the limit $g_4 \gg \lambda \gg g_1$, the minimum of the potential occurs in the D-flat direction of SU(4). By a suitable gauge rotation, the nonzero vevs have the form

$$
X_1 = X_2^{\mathrm{T}} = \begin{bmatrix} v_1 \\ 0 \\ 0 \\ 0 \end{bmatrix}, \quad
X_3 = \begin{bmatrix} v_2\sigma_2 & 0 \\ 0 & v_2\sigma_2 \end{bmatrix}, \quad
X_4 = v_3.
\tag{3}
$$

After rescaling all the fields, $\phi \to \frac{\Lambda_4}{\lambda^{1/4}} \phi$, the scalar potential is [25,26]

$$
V = V_F + V_D \quad \text{where} \quad V_F = \sum_i |F_{X_i}|^2 \quad \text{and} \quad V_D = \frac{1}{2} D_1^2,
$$

$$
V_F = \lambda^{6/5} \Lambda_4^4 \left[|v_1|^2 + \left| 2v_1 v_3 - \frac{1}{v_1^2 v_2} \right|^2 + \left| \frac{1}{v_1 v_2^2} \right|^2 \right],
$$

$$
V_D = \frac{1}{2} \left[4 g_1 \frac{\Lambda_4^2}{\lambda^{2/5}} (|v_1|^2 - |v_2|^2 - |v_3|^2) \right]^2.
\tag{4}
$$

Both V_F and V_D cannot be simultaneously zero for any value of v_1, v_2, v_3 and consequently supersymmetry is broken at the minimum. We can find the minimum of the potential numerically: For $g_1/\lambda \lesssim 0.1$, the potential is minimized at $\langle D_1 \rangle \sim 1.5 \, \mathrm{Max}[\langle F_{X_i} \rangle] \neq 0$.

Therefore, the 4-1 model provides both F-term and D-term supersymmetry breaking. Importantly, the absence of any gauge-singlet F-term vevs in the 4-1 model implies restrictions on the form of the operators that can mediate supersymmetry breaking. Amusingly, in the original context of gauge mediation, the 4-1 model was abandoned precisely because of the absence of singlets, since no Majorana gaugino masses are generated. This "bug" for

gauge mediation becomes an essential "feature" of our model that allows us to realize an accidental U(1)$_R$ symmetry.

IV. MEDIATING SUPERSYMMETRY BREAKING

Supersymmetry breaking is mediated through Planck-suppressed operators involving the visible sector fields. In this section we list and categorize all possible contact operators suppressed by up to two powers of M_{Pl}.

Below, X represents any one of chiral superfields X_1, X_2, X_3, X_4 in the 4-1 model. As shown in Eq. (1), X is *not* a singlet and so all operators involving single power of X are forbidden. Majorana gaugino mass operators (namely, $\int d^2\theta X W_a W_a$), the A term operators ($\int d^2\theta X Q U H$), and the Giudice-Masiero μ term ($\int d^4\theta X^\dagger H_u H_d$) involve a single X and hence do not exist in our model. The absence of these operators will be essential for the emergence of an accidental U(1)$_R$. The gauge-invariant chiral combination $X_1 X_2 X_4$ can give rise to these operators but they are suppressed by additional powers of M_{Pl} and are inconsequential.

The leading gauge-invariant operator involving X is $X^\dagger X$, by which we refer to the operator $X_i^\dagger \exp(q_a^i g_a V_a) X_i$, where q_a^i is the charge of X_i under the "a"th gauge group. This combination is also U(1)$_R$ invariant.

If the visible sector is extended to include chiral superfields Σ_a in the adjoint representation of the gauge group a, gauginos can acquire U(1)$_R$-invariant Dirac masses through dimension-5 operators involving the hidden sector U(1) superfield strength W' [43–47],

$$
\int d^2\theta \frac{W'}{M_{\mathrm{Pl}}} W_a \Sigma_a.
\tag{5}
$$

For the rest of this section, we use Φ_i to refer to all visible sector chiral superfields (i.e. the index i runs over all matter, Higgs and Σ fields). All dimension-5 and -6 operators in the superpotential and Kähler potential involving fields from both the hidden and visible sectors can be written as

$$
\text{dim5:} \quad \int d^2\theta \frac{W'W'}{M_{\mathrm{Pl}}} \Phi_i,
\tag{6}
$$

$$
\int d^2\theta \frac{X_1 X_2 X_4}{M_{\mathrm{Pl}}} \Phi_i,
\tag{7}
$$

$$
\int d^4\theta \frac{X^\dagger X}{M_{\mathrm{Pl}}} \Phi_i,
\tag{8}
$$

$$
\text{dim6:} \quad \int d^2\theta \frac{W'W'}{M_{\mathrm{Pl}}^2} \Phi_i \Phi_j,
\tag{9}
$$

$$\int d^2\theta\, \frac{X_1 X_2 X_4}{M_{\mathrm{Pl}}^2} \Phi_i \Phi_j, \qquad (10)$$

$$\int d^4\theta\, \frac{X^\dagger X}{M_{\mathrm{Pl}}^2} \Phi_i \Phi_j, \qquad (11)$$

$$\int d^4\theta\, \frac{X^\dagger X}{M_{\mathrm{Pl}}^2} \Phi_i^\dagger \Phi_j. \qquad (12)$$

All the dimension-5 operators in Eqs. (6)–(8) involve only the gauge singlets of Φ (for example, Σ_1). Operators in Eqs. (9)–(11) involve either two chiral fields of opposite charge (for example, $H_u H_d$) or chiral fields in the real representation (Σ_a^2). In particular, the B_μ term (namely, $\int d^4\theta X^\dagger X H_u H_d$) is contained in Eq. (11). Finally, Eq. (12) involves every field of Φ and generates soft mass-squareds for all the scalars.

V. THE VISIBLE SECTOR

Having established that the visible sector is now shielded from $U(1)_R$ breaking in the hidden sector, if we construct the visible sector to be $U(1)_R$ invariant itself, then it will remain so after supersymmetry breaking.

We describe our $U(1)_R$-invariant visible sector in terms of its three subsectors: the gauge and gaugino sector, the matter (squark and slepton) sector, and the Higgs sector. In this section we first concentrate on the supersymmetry breaking contributions from gravity-mediated operators, and only later reintroduce the small corrections arising from anomaly mediation.

The contact operators suppressed by M_{Pl} must be renormalized down to the electroweak symmetry breaking (EWSB) scale before they are converted to soft mass terms and the flavor violating observables are calculated. In our model, we need to renormalize the Dirac gauginos mass operators in Eq. (5), scalar soft mass-squared operators in Eq. (12), and the B-type scalar mass operators in Eq. (11). As we will see, the renormalization of these operators leads to a substantial increase in the gluino Dirac mass relative to the squark masses even through $\langle D \rangle \sim \langle F \rangle$ at the Planck scale. Although these operators scale due to the interactions present in the visible sector as well as in the hidden sector [48], we do not consider renormalization from hidden sector interactions since the effect is small in the 4-1 model. Amusingly, however, the effect of hidden sector renormalization is to further lower the contributions to the soft mass-squared of the matter (and other scalar) sector relative to the Dirac gauginos, allowing for potentially an even larger mass ratio between gauginos and squarks/sleptons [49–54].

A. Gauge and gaugino sector

Since Majorana masses are not generated by gravity mediation from the 4-1 model, we introduce chiral adjoint superfields Σ_a for each gauge symmetry of the standard model [26,43–45,47,55–62]. This allows us to write the gravity-mediated operator Eq. (5), which generates Dirac gaugino masses (as well as scalar masses for the adjoint scalar).

The part of the Lagrangian involving visible sector gauge field-strength superfields W_a is then written as

$$\int d^2\theta \sum_a W^a W^a + \int d^2\theta \sum_a \sqrt{2}\, w_a \frac{W'}{M_{\mathrm{Pl}}} W_a \Sigma_a, \qquad (13)$$

where w_a are order one dimensionless coupling constants. These constants—leading to the Dirac masses—scale differently from operators that would lead to Majorana masses in models with gauge singlets in the hidden sector [47]. The evolution is given by

$$
w_a(\mu) = w_a(M_{\mathrm{Pl}}) \sqrt{\frac{Z_{\Sigma_a}(M_{\mathrm{Pl}})}{Z_{\Sigma_a}(\mu)}} \sqrt{\frac{\alpha_a(M_{\mathrm{Pl}})}{\alpha_a(\mu)}}
$$

$$
= w_a(M_{\mathrm{Pl}}) \begin{cases} \left(\dfrac{\mu}{\Lambda}\right)^{-(N_a/2\pi)\alpha_a} & \text{for } b_a = 0, \\[2ex] \left(\dfrac{g_a(\mu)}{g_a(\Lambda)}\right)^{1-(2Na/ba)} & \text{for } b_a \neq 0. \end{cases} \qquad (14)
$$

As we will see, with the particle content of the model, the one-loop beta function coefficients are $b_3 = 0$, $b_2 = 4$, and $b_1 = 36/5$ [in the $SU(5)$ normalization], just as in Ref. [24]. In particular, assuming that there are no additional fields with standard model charges between the EWSB scale and the Planck scale, one finds

$$w_3(1 \text{ TeV}) \approx 5.4 \times w_3(M_{\mathrm{Pl}}), \qquad (15)$$

$$w_2(1 \text{ TeV}) = w_2(M_{\mathrm{Pl}}), \qquad (16)$$

$$w_1(1 \text{ TeV}) \approx 0.58 \times w_1(M_{\mathrm{Pl}}). \qquad (17)$$

Replacing Φ by Σ_a in Eqs. (6)–(12) one can find the gravity-mediated operators involving the Σ_a fields. The part of the Lagrangian in our model which contains these adjoint fields consists of the following operators:

$$\int d^4\theta \sum_a \left(1 + c_a \frac{X^\dagger X}{M_{\mathrm{Pl}}^2}\right) \mathrm{Tr}\big[\mathrm{e}^{-g_a V_a} \Sigma_a^\dagger \mathrm{e}^{g_a V_a} \Sigma_a\big]$$
$$+ \int d^4\theta \sum_a k_a \frac{X^\dagger X}{M_{\mathrm{Pl}}^2} \mathrm{Tr}\Sigma_a^2. \qquad (18)$$

Squared soft masses for the Σ fields are generated from operators with dimensionless coupling constants c_a, whereas the operators with couplings k_a give rise to B-type masses. In the absence of renormalizable superpotential interactions of Σ, the couplings c_a do not renormalize but k_a do:

$$c_a(\mu) = c_a(M_{\mathrm{Pl}}), \qquad (19)$$

$$k_a(\mu) = k_s(M_{\mathrm{Pl}})\left(\frac{Z_{\Sigma_a}(M_{\mathrm{Pl}})}{Z_{\Sigma_a}(\mu)}\right). \tag{20}$$

The Dirac gaugino mass operators in Eq. (13) also contribute to the squared mass of the Σ fields. When expanded in components, they give tree-level masses for the real parts of Σ_a. Once the Dirac gauginos are integrated out, the imaginary components of Σ_a also get masses due to one-loop finite contributions. In summary, the dominant contributions to the scalar mass-squareds can be summarized as

$$m^2_{\mathrm{Re}[\Sigma_a]} = 2|w_a|^2\frac{\langle D\rangle^2}{M_{\mathrm{Pl}}^2} + c_a\frac{|\langle F_X\rangle|^2}{M_{\mathrm{Pl}}^2}, \tag{21}$$

$$m^2_{\mathrm{Im}[\Sigma_a]} = \frac{N_a\alpha_a}{\pi}\log\left(\frac{m_{R_a}^2}{M_a^2}\right)|w_a|^2\frac{\langle D\rangle^2}{M_{\mathrm{Pl}}^2} + c_a\frac{|\langle F_X\rangle|^2}{M_{\mathrm{Pl}}^2}. \tag{22}$$

Here M_a, m_{R_a} are the *physical* masses of the gauginos and the real part of Σ_a, respectively.

B. Matter sector

The matter sector in our model comprises quark and lepton superfields in three generations, just like the minimal supersymmetric standard model (MSSM). We represent the field content by $\Phi_i^{(n)} \in \{Q_i, U_i, D_i, L_i, E_i\}$, where i is the generation index. Excluding the Yukawa terms, the Lagrangian in the matter sector may be written down as

$$\int d^4\theta \sum_n\left(\delta_{ij} + c_{ij}^n\frac{X^\dagger X}{M_{\mathrm{Pl}}^2}\right)\Phi_i^{(n)\dagger}\mathbf{e}^{(q_n^a g_a V_a)}\Phi_j^{(n)}, \tag{23}$$

where q_n^a designate charges of the superfield $\Phi^{(n)}$ under the ath gauge group and c_{ij}^n are dimensionless coupling constants. The second term in Eq. (23) belongs to the gravity-mediated operators in Eq. (12). No other operators in Eqs. (6)–(11) are allowed because of the charges of the matter fields.

The Planck-suppressed operator coefficient matrices c_{ij}^n are in general anarchic in flavor space, and thus give rise to flavor-arbitrary squark and slepton masses. Given the absence of Majorana gaugino masses, the scalar masses evolve with the renormalization group scale only because of the ordinary Yukawa couplings of the chiral superfields,

$$c_{ij}(\mu) = c_{kl}(M_{\mathrm{Pl}})\mathbf{P}\exp\left(-\int_0^t dt'\frac{\gamma_{ijkl}}{16\pi^2}\right)$$

where $\gamma_{ijkl} = \frac{1}{2}y_{ipq}^*y_{kpq}\delta_{jl} + \frac{1}{2}y_{lpq}^*y_{jpq}\delta_{ik} + 2y_{ipq}^*y_{jpq},$

$$\tag{24}$$

where $t = \log\mu$, and $\mathbf{P}\exp()$ refers to path ordered exponential to account for any noncommutativity of the matrices y^*y.

In addition to the renormalization group (RG) evolution of squark and slepton masses, the Dirac gauginos induce a one-loop finite contribution to their masses. This generates one-loop finite flavor-universal contributions:

$$\Delta m_{ij}^2 = \delta_{ij}c_i^0\frac{\langle D\rangle^2}{M_{\mathrm{Pl}}^2}, \tag{25}$$

where

$$c_i^0 = \sum_a C_2^a(\mathbf{R}_i)\frac{\alpha_a}{\pi}w_a^2\log\left(\frac{m_{R_a}^2}{M_a^2}\right). \tag{26}$$

Here M_a, m_{R_a} are the physical masses of the gauginos and the real part of the scalar gaugino, respectively, and $C_2^a(\mathbf{R}_i)$ is the quadratic Casimir of the representation $\mathbf{R}_i$. Therefore, the scalar soft masses at the EWSB scale are given as

$$m_{ij}^2 = \delta_{ij}c_i^0\frac{\langle D\rangle^2}{M_{\mathrm{Pl}}^2} + c_{ij}\frac{|\langle F_X\rangle|^2}{M_{\mathrm{Pl}}^2}. \tag{27}$$

C. Higgs sector

In the minimal R-symmetric supersymmetric standard model [24], the $U(1)_R$-extended Higgs sector does not have the usual Higgsino mass (namely, $\mu H_u H_d$). Instead, it contains an additional pair of electroweak doublets $R_{u,d}$ that do not couple to matter and have new mass terms $\mu_{u,d}H_{u,d}R_{u,d}$. Different R-charge assignments of the R and H fields ensure that they do not mix among themselves.

In contrast, we do not impose explicit $U(1)_R$ in our model, and so mixed terms such as $\int d^4\theta H^\dagger e^V R$, $\int d^4\theta X^\dagger X H^\dagger R$ can arise in the effective theory. These operators force R and H fields to be rotated among themselves that ultimately lead to the R-violating $\mu H_u H_d$ term and R-violating couplings of R fields with the matter sector fields. Therefore, we need a symmetry to distinguish H fields from R fields.

Our proposal extends the gauge symmetry of the visible sector by an additional $U(1)_X$, under which R transforms but H does not.[2] Fields in the Higgs sector of our model and their quantum numbers are summarized in the following table:

	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_X$	
$H_{u,d}$	**1**	**2**	$\pm 1/2$	0	
$R_{u,d}$	**1**	**2**	$\mp 1/2$	± 1	(28)
$S_{u,d}$	**1**	**1**	0	∓ 1	
$T_{u,d}$	**1**	**1**	0	± 2	

We utilize the following set of marginal interactions to ensure masses for all matter and Higgs fields once electroweak symmetry is broken:

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[2]A global or discrete symmetry can also be invoked to distinguish H and R. However, global symmetries are not compatible with quantum gravity and typically, a discrete symmetry would be broken spontaneously at the EWSB scale leading to formation of domain walls.

$$\int d^2\theta \left[y_u H_u Q U^c + y_d H_d Q D^c + y_\ell H_d L E^c + \alpha_u S_u R_u H_u \right.$$

$$\left. + \alpha_d S_d R_d H_d + \frac{1}{2}\beta_u T_u S_u^2 + \frac{1}{2}\beta_d T_d S_d^2 \right] + \text{c.c.} \quad (29)$$

The Kähler potential terms in our model involving Higgs fields can conveniently be written using the field $\Phi_{u,d}^{(n)} \in \{H_{u,d}, R_{u,d}, S_{u,d}, T_{u,d}\}$:

$$\int d^4\theta \sum_n \sum_{i \in \{u,d\}} \left(1 + c_n \frac{X^\dagger X}{M_{\text{Pl}}^2}\right) \Phi_i^{(n)\dagger} \mathbf{e}^{(q_{ni}^a V_a)} \Phi_i^{(n)}$$

$$+ \int d^4\theta \sum_n k_n \frac{X^\dagger X}{M_{\text{Pl}}^2} \Phi_u^{(n)} \Phi_d^{(n)}. \quad (30)$$

The operator proportional to dimensionless coupling constants c_n give rise to scalar soft masses of the Higgs fields, whereas the ones proportional to k_n generate the B-type masses. The RG evolution of c_n and k_n are identical to the ones listed in Eqs. (24) and (20), respectively.

D. Emergence of an accidental $U(1)_R$

All marginal operators written down in this section are exactly invariant under the following R-charge assignment:

$$\mathbf{R}[Q, U, D, L, E] = 1, \qquad \mathbf{R}[Wa] = 1,$$
$$\mathbf{R}[\Sigma a] = 0, \qquad \mathbf{R}[H, S] = 0, \qquad \mathbf{R}[R, T] = 2. \quad (31)$$

This assignment is an extended version of the $U(1)_R$ symmetry that the authors in Ref. [24] imposed explicitly at the EWSB scale in order to reduce the contribution of flavor-arbitrary scalar soft masses to the flavor violating observables. They exploited the fact that the dominant (dimension-5) operators that are responsible for large loop contribution to these observables also violate $U(1)_R$ by two units.

It is striking to realize that all but two contact terms in our gravity-mediated potential from the previous subsection preserve this $U(1)_R$, even though we have not imposed it anywhere in our construction. Only the b-type operators $\int d^4\theta R_u R_d$ and $\int d^4\theta T_u T_d$ violate $U(1)_R$. However, note that each of these terms violate $U(1)_R$ by four units and since these are the only terms that violate $U(1)_R$, all effective operators that are generated at the EWSB scale either preserve $U(1)_R$ or violate $U(1)_R$ by units of 4, 8, 12,…. The dimension-5 operators that put severe constraints on arbitrary c_{ij} would not be generated at all in our model.

E. CP violation in the Higgs sector

In our Higgs sector, new phases can appear in the superpotential couplings $\alpha_{u,d}$ and $\beta_{u,d}$ and in the B-type soft masses $B_{h,r,s,t}$. However, there are only two combinations that are invariant under arbitrary phase redefinitions of these fields:

$$I_1 = \alpha_u \alpha_d (B_h B_r B_s)^*, \qquad I_2 = \beta_u \beta_d (B_s^2 B_t)^*. \quad (32)$$

In addition to I_1, using $\alpha_{u,d}$ and $B_{h,r,s}$, we could also form combinations like $\alpha_u \alpha_d / (B_h B_r B_s)$ or $B_h B_r B_s / (\alpha_u \alpha_d)$ which would be invariant under phase redefinition. However, these are not relevant for CP violation for the following reasons. First, physical observables must not involve a negative power of any of the B parameters, because no field becomes massless in the limit $B \rightarrow 0$ (as long as we keep nonzero soft mass-squareds for the scalars) so we should be able to expand any amplitude in positive powers of B. Second, $1/\alpha_{u,d}$ are not allowed because if one of $\alpha_{u,d}$ happens to be zero, we can completely remove all phases from $\alpha_{u,d}$ and $B_{h,r,s}$, so any CP violation from this set of parameters should vanish as $\alpha_{u,d} \rightarrow 0$. Therefore, all of $\alpha_{u,d}$ and $B_{h,r,s}$ must appear with positive powers, hence I_1 is the unique invariant. A similar argument can be made to single out I_2 for the β-B_s-B_t sector. Combined together, any CP violating physical observable must involve positive powers of $I_{1,2}$.

Furthermore, recall that with our R-charge assignment, B_r and B_t each break R-charge by 4 while all other parameters are R-neutral, so I_1 and I_2 each carry R-charge 4. Then, since all CP-violating observables in the standard model are R-neutral, they must appear in the combination $I_1 I_2^*$. This is a dimension-12 object which involves all four of $\alpha_{u,d}$ and $\beta_{u,d}$, so CP is broken not only really softly but also at a very high loop order, which therefore should be inconsequential.

VI. PHENOMENOLOGY

A comprehensive study of the parameter space in our model is beyond the scope of this paper. We rather find one point that is allowed under the flavor and CP violation constraints and satisfies other experimental bounds.

There are several novel phenomenological features of this point. At the weak scale, the point contains large flavor violation and Dirac gauginos that have been studied before in the context of the minimal R-symmetric supersymmetric standard model (see for example, [63–65]). In addition to these, our benchmark point also contains a light Z' with nontrivial couplings to the Higgs sector.

A. Benchmark point

First, let us choose gaugino masses. To reduce the number of parameters, let us make a simplifying assumption that the $SU(3) \otimes SU(2) \otimes U(1)_Y$ gauginos have an equal mass at the Planck scale, i.e.,

$$w_{3,2,1}|_{M_{\text{Pl}}} = 1 \quad \text{and} \quad \frac{\langle D \rangle}{M_{\text{Pl}}}\bigg|_{M_{\text{Pl}}} = 1 \text{ TeV}, \quad (33)$$

where $w_{1,2,3}$ are defined in Eq. (5) with the "SU(5) normalization" for $U(1)_Y$, although it is not our intention here to imply or require grand unification. Then, after evolving

the gaugino masses with renormalization group running Eq. (14) down to $\mu = 1$ TeV, we find

$$M_3 = 5.4 \text{ TeV},$$
$$M_2 = 1.0 \text{ TeV}, \qquad (34)$$
$$M_1 = 580 \text{ GeV}.$$

Instead of treating w_X at M_{Pl} as input parameter and then renormalizing it down to the EWSB scale, we rather choose the mass parameter M_X at 1 TeV to be

$$M_X = 500 \text{ GeV}. \qquad (35)$$

B. Flavor/*CP* violation

Let us analyze the most stringent constraints, the flavor and *CP* violation in K^0-$\bar{K}^0$ mixing. When generated at the Planck scale, the coefficient matrix c_{ij} of the Kähler term $X^\dagger X Q_i^\dagger Q_j$ (and similarly for U^c and D^c) should be "anarchic"; to be more precise, we take each element of c_{ij} to have a magnitude and a phase in the range from 0.1 to 1.

As in Eq. (27), the low-energy squark mass matrix also receives contributions from integrating out the heavy Dirac gauginos (dominantly the gluino). For the benchmark gluino mass, we obtain

$$m_{ij}^2 = \delta_{ij}(1.3 \text{ TeV})^2 + c_{ij} \frac{|\langle F_X \rangle|^2}{M_{\text{Pl}}^2}$$
$$= (1.3 \text{ TeV})^2 \left(\delta_{ij} + c_{ij} \frac{|\langle F_X \rangle|^2}{\langle D \rangle^2} \right). \qquad (36)$$

(Here, the running of $X^\dagger X Q^\dagger Q$ operator was neglected since it depends only on small Yukawa couplings at one loop when gauginos are Dirac. To the extent that there is RG evolution, its effect is to further lower the squark mass with respect to the gluino mass, only making our discussion stronger.) Thus, the flavor violating parameter $\delta_{ij} \equiv \Delta m_{ij}^2 / m_{\text{avg}}^2$ is given by

$$\delta_{ij} \approx \frac{\langle F_X \rangle^2}{\langle D \rangle^2} \frac{c_{ij}}{c^0} \quad \text{for } i \neq j. \qquad (37)$$

The strongest constraint from $\Delta F = 2$ observables follow from K^0-$\bar{K}^0$ mixing ($\delta_L = \delta_R \leq 0.08$) for gluino and squark masses to be order of 5 TeV and 1 TeV, respectively [66]). Assuming that $\langle D \rangle \sim 1.5 |\langle F_X \rangle|$, we find that

$$c_{ij}^q \leq 0.3 \times \begin{pmatrix} \sim 1 & \sim 1 & \sim 1 \\ \sim 1 & \sim 1 & \sim 1 \\ \sim 1 & \sim 1 & \sim 1 \end{pmatrix}. \qquad (38)$$

An even stronger bound arises due to *CP* violation in K^0-$\bar{K}^0$ mixing. Roughly, the *CP* violating quantity $\text{Im}(\delta_L^* \delta_R) \lesssim 10^{-3}$, which can be accomplished by taking $c_{ij}^q \leq 0.1$ (multiplying an anarchic matrix) as well as the relative phase to be smaller than 0.1.

Therefore, we conclude that it is possible to satisfy the stringent bound from the *CP* violation in K^0-$\bar{K}^0$ even with quite "flavorful" sfermion mass-squareds at the Planck scale, given the ratio $m_{\text{squark}}/m_{\text{gluino}} \sim \mathcal{O}(0.1)$, which arises naturally from RG evolution in our model. As discussed in Refs. [24,66], all other squark flavor bounds are satisfied once this ratio is assumed.

Lepton flavor violation also provides constraints on the slepton mass matrices [24,67]. Maximal mixing is strongly constrained, particularly in the right-handed slepton sector, due to constraints from both $\mu \rightarrow e$ conversion as well as $\mu \rightarrow e\gamma$ [67]. For a Dirac bino of mass 600 GeV, a relatively mild restriction on the mixing angle $\sin 2\theta_l \lesssim 0.3 \rightarrow 0.5$ for nondegenerate slepton mass eigenstates is necessary. For right-handed sleptons, since the dominant contribution to their masses is from the flavor-arbitrary Kähler potential operators, this implies some modest degeneracy needed, roughly $c_{12}^l / c_{ii}^l \lesssim 0.15 \rightarrow 0.3$. This is well within our stated goal of coupling constants remaining within the range $0.1 \rightarrow 1$.

C. Electroweak and U(1)$_X$ breaking

We seek a minimum of the scalar potential where $\langle H_{u,d} \rangle \neq 0$ and $\langle S_{u,d} \rangle \neq 0$ and $\langle R_{u,d} \rangle = \langle T_{u,d} \rangle = 0$, so that the electroweak symmetry is broken down to electromagnetism [SU(2)$_L \times$ U(1)$_Y \rightarrow$ U(1)$_{\text{EM}}$] and U(1)$_X$ is broken completely, while preserving the "accidental" U(1)$_R$ of the visible sector. The potential in our model depends on soft masses of all the multiplets in the Higgs sector (i.e. $H_{u,d}$, $R_{u,d}$, $S_{u,d}$, and $T_{u,d}$), soft masses of the adjoints Σ_a and the Yukawa couplings $\alpha_{u,d}$ and $\beta_{u,d}$.

Let us first consider the effect of the scalar adjoints. As shown in Eq. (21), the real parts of the scalar adjoints receive large masses from Dirac gaugino mass operators. For our benchmark scenario, their physical masses are at or above 1 TeV, while we expect the mass parameters in the Higgs sector are typically of the order of the weak scale. Hence, we integrate out the scalar adjoints to analyze the scalar potential in the Higgs sector.

In this limit, as is well known, the quartic terms in the Higgs-U(1)$_X$ sector vanish in the absence of soft masses $m_{\tilde\Sigma_a}^2$ and $b_{\tilde\Sigma_a^2}$. Once soft masses for these scalars are generated, the quartic is partially restored. In particular the *D*-term due to each gauge group shifts according to

$$\frac{g_a^2}{8} \rightarrow \eta_a \frac{g_a^2}{8} = \left(\frac{m_{\tilde\Sigma_a}^2 - b_{\Sigma_a}}{4M_a^2 + m_{\tilde\Sigma_a}^2 + b_{\Sigma_a}} \right) \frac{g_a^2}{8}. \qquad (39)$$

With our gaugino masses and using $m_{\tilde\Sigma_a}^2 \sim -b_{\Sigma_a} \sim$ (670 GeV)^2 (it corresponds to $c_{\Sigma_a} \sim -k_{\Sigma_a} \sim 1$), we find that the MSSM *D*-term effectively scales approximately as $(g_1^2 + g_2^2)/8 \rightarrow (1/4)(g_1^2 + g_2^2)/8$. This is not meant to be the best or even indicative of the quartic suppression—it simply illustrates the actual suppression for the particular

benchmark point that we have chosen to study. Finally, for the $U(1)_X$ sector we use $\eta_X = 1/3$.

Four nontrivial conditions to attain the right vacuum structure follow from the minimization conditions, $\partial V/\partial H^0_{u,d} = \partial V/\partial S_{u,d} = 0$, which we use to eliminate four input parameters, $m^2_{H_{u,d}}$, $m^2_{S_{u,d}}$ in terms of other soft and supersymmetry breaking parameters. In particular, we choose the following point (all parameters in the EWSB scale):

$$\langle S_u \rangle^2 + \langle S_d \rangle^2 = v^2_s = m^2_X/(2g^2_X) = (600 \text{ GeV})^2,$$

$$\frac{\langle S_u \rangle}{\langle S_d \rangle} = \tan\beta_X = 0.2, \qquad \frac{\langle H^0_u \rangle}{\langle H^0_d \rangle} = \tan\beta = 10,$$

$$m_X = 150 \text{ GeV} \qquad B_{h,r,t} = (300 \text{ GeV})^2,$$

$$B_s = (500 \text{ GeV})^2, \qquad \alpha_u = 1,$$

$$\alpha_d = 0.3, \qquad \beta_u = 0.9, \qquad \beta_d = 0.7,$$

$$m^2_{R_{u,d}} = m^2_{T_{u,d}} = (700 \text{ GeV})^2. \tag{40}$$

(We have checked that these couplings do not hit Landau poles below M_{Pl}.) The gauge coupling constant g_X can be inferred from the gauge boson mass m_X and total $U(1)_X$ breaking vacuum expectation value (vev) v_s. With our parameters in Eq. (40) we find $g_X = 0.18$. These values as well as the gaugino masses Eqs. (34) and (35) for all gauge groups define our benchmark point.

Now, at this benchmark point, the chargino and neutralino mass matrices are completely determined at tree level from which we find that the lightest two neutralinos have masses of 54, 176 GeV, and the lightest chargino has a mass of 117 GeV. These are safely above the LEP limits.

The CP-even neutral Higgs mass matrix is more complicated, where the scalar components of H^0_u, H^0_d, S_u, S_d mix among themselves. Naively, it might appear challenging to get Higgs masses above the LEP limit, given the η-suppressions of the quartic couplings, Eq. (39). For example, the lightest mass eigenstate in the CP-even sector would attain only $\simeq 20$ GeV mass at tree level at our benchmark point. Therefore, as in the MSSM, the one-loop quantum corrections to the quartic coupling from top/ stop loops are very important. At our benchmark point, there are additional important one-loop contributions from the superpotential term proportional to the α_u coupling. The radiative corrections to the Higgs quartic can be roughly estimated as

$$\frac{3y^4_t}{32\pi^2}\log\left(\frac{m_{\tilde{t}_1}m_{\tilde{t}_2}}{m^2_t}\right) + \frac{\alpha^4_u}{32\pi^2}\log\left(\frac{m_{R_u}m_{S_u}}{m_{\tilde{R}}m_{\tilde{S}}}\right). \tag{41}$$

The first term in Eq. (41) is due to the familiar top-stop mass splitting in the MSSM and the second contribution is due to the splittings among the neutralinos and the neutral scalars in R_u and S_u. Strictly speaking, α_u in Eq. (41) should have been replaced by the Higgs couplings of

various neutralino/neutral scalar mass eigenstates. We can, however, give a rough estimate of Higgs mass. For our benchmark point, stop masses are ~ 1.3 TeV and the top-stop splitting generates ~ 96 GeV for Higgs mass. Approximating $m_{\tilde{R},\tilde{S}} \sim 100$ GeV we see that $m_{R_u,S_u} \sim 1$ TeV is sufficient to obtain a lightest Higgs mass above the LEP bound.

We should emphasize that the difficulty to get the Higgs mass above the LEP bound is for our benchmark point, and does not in general apply to our model. There are at least two interesting ways to raise the Higgs mass beyond what we considered above: One is to raise the scalar masses $m_{\tilde{R},\tilde{T}}$, to increase the one-loop contribution. The second is to increase the masses of the adjoint scalars $m_{\tilde{\Sigma}_{1,2}}$, to increase the tree-level quartic coupling. Nevertheless, the lightest Higss boson does generically tend be rather light, close to the LEP bound, given the scales in the model.

D. $U(1)_R$ violation from anomaly mediation

Since the gravitino mass necessarily breaks $U(1)_R$, there are small but irreducible contributions to Majorana gaugino masses, slightly violating the accidental $U(1)_R$ of the visible sector. Fortunately, they do not upset our flavor protection mechanism. This issue was already studied in Ref. [24], which we briefly recast here.

The anomaly-mediated contributions to the Majorana-type gaugino masses are [31,32], in our model,

$$\tilde{M}_a = b_a \frac{\alpha_a}{4\pi} m_{3/2} \simeq \begin{cases} 0.01m_{3/2} & a = 1 \\ 0.01m_{3/2} & a = 2 \\ 0 & a = 3, \end{cases} \tag{42}$$

where we have evaluated the β-functions for our model content with couplings at the weak scale. Interestingly, the anomaly-mediation-induced Majorana mass for the gluino is zero at one loop because the SU(3) β-function vanishes for our field content. Indeed, any visible sector model minimally extended to become R-symmetric has this property. In practice, two-loop contributions to the Majorana gluino mass are present, as well as one-loop threshold contributions resulting from differences between the masses of the Dirac gluino versus the color octet scalars. These are expected to lead to a contribution smaller than about $10^{-3}m_{3/2}$.

The size of the gravitino mass $m_{3/2}$ relative to the visible sector is, in general, sensitive to the hidden sector [48,68–70]. In our case, however, this effect is small. Noting that $m_{3/2}$ is related to the cosmological constant [71] which is, in turn, approximately equal to the expectation value of the total potential, we find that for our benchmark point the gravitino mass is given as

$$m_{3/2} \approx 730 \text{ GeV}. \tag{43}$$

This implies the Majorana masses for all of the gauginos are below about 10 GeV. The small splittings in the

gaugino sector do not induce excessive R-violation back into flavor violating processes. They are, however, quite relevant to determining the identity of the lightest supersymmetric particle, its relic abundance, and the associated collider phenomenology.

E. $U(1)_X$ phenomenology

The $U(1)_X$ gauge boson Z' can be potentially very interesting given its rather light mass [in Eq. (40) at the benchmark point]. How can we produce Z' at colliders? At tree level, $U(1)_X$ only couples to $S_{u,d}$ and $T_{u,d}$, and the $S_{u,d}$ scalars can mix with $H_{u,d}$ with large mixing angles. Therefore, one of the more important production channels is $gg \rightarrow h^* \rightarrow Z' + Z'$.

At one-loop level, kinetic mixing of $U(1)_X$ and $U(1)_Y$ is inevitably induced from the loops of $R_{u,d}$, which carry both charges. Therefore, through this kinetic mixing, the Z' emitted off the Higgs boson will in turn decay to a pair of standard model fermions. We leave a detailed analysis of this interesting signal and other potentially useful channels for probing $U(1)_X$ to future work. Here, we check that such a light Z' is allowed by existing constraints.

First, note that the $U(1)_X$-$U(1)_Y$ kinetic mixing is much smaller compared to the analogous mixing for the well-studied extension involving $U(1)_{B-L}$. This is because the only fields running in the loop in our case are $R_{u,d}$, compared to the entire matter content of the standard model in the $B - L$ case. More explicitly, the mixing coefficient can be estimated as

$$\epsilon = \sum_i q_Y^i q_X^i \frac{g_Y g_X}{8\pi^2} \log\frac{\Lambda}{m_i} \approx 1.6 \times 10^{-3} \left(\frac{g_X}{0.18}\right) \log\frac{\Lambda}{1 \text{ TeV}},$$

(44)

where q_Y^i and q_X^i are charges of field i under $U(1)_Y$ and $U(1)_X$, respectively, and m_i is the mass. The scale Λ is some high mass scale above which there is no mixing (e.g. embedding $U(1)_X$ into a non-Abelian group).

The kinetic mixing implies the Z' couples to all standard model matter particles and also contributes to precision electroweak observables. The strength of the Z' coupling to standard model fields is given by $q_Y^i \epsilon g_Y$, where q_Y^i are the hypercharges of corresponding standard model fields. The bound on the Z' mass can be directly obtained from Ref. [72]: $m_X \gtrsim \epsilon g_Y \times 6.7$ TeV.

Considering only the Z' couplings proportional to ϵ, an interesting generalized bound was recently analyzed in Ref. [73]. For $\epsilon \lesssim 3 \times 10^{-2}$, m_X is allowed to be in the entire range between 10–1000 GeV. For our benchmark point, we find that ϵ would exceed 3×10^{-2} only when Λ exceeds 10^{11} GeV. The precise bounds in our model requires further study, however. This is because the Z'-stralung process is undoubtedly constrained for small Z' masses. We will return to this very interesting phenomenological issue in future work.

VII. UNIFICATION, UNIFONS, AND SINGLETS

As pointed out in Ref. [24], gauge coupling unification of the standard model gauge couplings is not automatic. The SU(3) color and SU(2) weak interactions do continue to unify to a perturbative value near 10^{16} GeV as in the MSSM, since the shift in the β-functions from the enlarged field content (Σ_3, Σ_2, R_u, R_d) is equivalent to three sets of fundamentals and antifundamentals for SU(3) and SU(2). The $U(1)_Y$ coupling β-function, however, is only shifted by one unit. Hence, two pairs of vectorlike unifons with hypercharge $q_Y = \pm 1$ transforming under $U(1)_Y$ are sufficient to have the standard model couplings unify near 10^{16} GeV.

However, this is not the end of the story. There are also dangerous operators in supersymmetric models involving singlets [74]. Dirac gaugino masses for $U(1)_Y$ and $U(1)_X$ require gauge singlets in the visible sector. A completely general Kähler potential would allow operators such as

$$\int d^4\theta \frac{X^\dagger X}{M_{\text{Pl}}} \Sigma_{Y,X}$$

(45)

that lead to large tadpoles for the singlet fields $\Sigma_{Y,X}$.

There are two possibilities for dealing with this. One is to assume that the $U(1)$s are embedded into one or more non-Abelian (unified) groups. So long as this unification is not too close to the Planck scale, higher dimensional operators involving breaking fields of the unified sector would keep the coefficients of the tadpole operators small enough.

Another possibility is to not have singlets at all in the visible sector. This would mean both $U(1)_Y$ and $U(1)_X$ would acquire purely Majorana-type masses. While this may be a viable solution for $U(1)_X$, a pure Majorana bino would reintroduce large lepton flavor violation. The most dangerous lepton flavor violation would arise in operators involving the right-handed sleptons, since their masses are (in general) entirely anarchic. The naive bounds on δ_{RR} suggest tuning at level of 10^{-2} would be needed.

VIII. DISCUSSION

We have presented a complete model of gravity-mediated supersymmetry breaking with a visible sector that is safe from flavor constraints. We exploited the observation that flavor violation beyond the standard model can be greatly suppressed *even without flavor universality of the squark and slepton sectors* if the visible sector particle content respects an approximate global $U(1)_R$ symmetry [24]. The key difference between previous approaches to R-symmetry and the one pursued here is that the origin of the approximate $U(1)_R$ symmetry is *accidental*. Since we employ gravity mediation, all operators must be RG evolved to the weak scale, and this causes a large upward renormalization of the gluino mass. In this way, the "little hierarchy" in which the Dirac gluino mass was

needed to be higher than the squark masses is automatic in our framework.

We emphasize that all flavor problems are solved with our framework, including Δm_K, ϵ_K, and lepton flavor violation constraints including $\mu \to e$ conversion and $\mu \to e\gamma$. The solution can be mostly attributed to R-symmetry and the large renormalization of the gluino mass. This in itself is not quite enough [66], and so we must also allow ourselves to tune $c_{ij}^{\tilde{q}}$ coefficients to order of 0.1 in magnitude and phase to fully comply with the constraint from ϵ_K. Somewhat more relaxed restrictions on the c_{ij}^l maintain safely from lepton flavor violation constraints.

Our philosophy exploited holomorphy of the superpotential while we strictly eliminated dangerous nonchiral Kähler terms through an additional gauged $U(1)_X$ symmetry in the Higgs sector. The phenomenology of this $U(1)_X$ is quite interesting: the Z' of this broken $U(1)_X$ couples with order one strength to Higgs fields and order $10^{-4} \to 10^{-2}$ strength to particles carrying hypercharge. More detailed phenomenology is left to future work, but clearly the prospects for detection at LHC and Tevatron are bright.

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