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Soft RPV through the baryon portal

Gordan Krnjaic^a and Yuhsin Tsai^b

^a*Perimeter Institute for Theoretical Physics,
Waterloo, Ontario, Canada*

^b*Physics Department, University of California Davis,
Davis, California, U.S.A.*

E-mail: gkrnjaic@perimeterinstitute.ca, yhtsai@ucdavis.edu

ABSTRACT: Supersymmetric (SUSY) models with R -parity generically predict sparticle decays with invisible neutralinos, which yield distinctive missing energy events at colliders. Since most LHC searches are designed with this expectation, the putative bounds on sparticle masses become considerably weaker if R -parity is violated so that squarks and gluinos decay to jets with large QCD backgrounds. Here we introduce a scenario in which baryonic R -parity violation (RPV) arises effectively from soft SUSY breaking interactions, but leptonic RPV remains accidentally forbidden to evade constraints from proton decay and FCNCs. The model features a global R -symmetry that initially forbids RPV interactions, a hidden R -breaking sector, and a heavy mediator that communicates this breaking to the visible sector. After R -symmetry breaking, the mediator is integrated out and an effective RPV A -term arises at tree level; RPV couplings between quarks and squarks arise only at loop level and receive additional suppression. Although this mediator must be heavy compared to soft masses, the model introduces no new hierarchy since viable RPV can arise when the mediator mass is near the SUSY breaking scale. In generic regions of parameter space, a light thermally-produced gravitino is stable and can be a viable dark matter candidate.

KEYWORDS: Supersymmetry Phenomenology

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1 Introduction

Weak scale supersymmetry (SUSY) has long been the leading framework for addressing the hierarchy problem. However, after accumulating over 20 fb^{-1} of data, the LHC has yet to find any evidence of superpartners near the TeV scale and has already placed tight constraints on the most compelling regions of SUSY parameter space. As the lower bounds on stop and higgsino masses approach the TeV range, there is generic tension with naturalness; at least some fine tuning is required to stabilize the electroweak scale.

However, this interpretation of LHC results is model dependent since most SUSY searches assume R -parity conservation and, thus, require substantial MET in the final state. If this assumption is relaxed, sparticles can decay to standard model particles and the bounds become significantly weaker, thereby alleviating the tension with naturalness. Since none of SUSY's theoretically desirable features strictly requires R -parity, the current experimental situation motivates serious efforts to construct viable R -parity violating (RPV) alternatives.

Figure 1. The loop process that yields SUSY breaking fermion-scalar RPV interactions.

In the absence of R -parity, the MSSM allows dangerous baryon and lepton violating operators in the superpotential

$$W_{RPV} = \frac{\lambda_{ijk}}{2} L_i L_j \bar{E}_k + \lambda'_{ijk} Q_i L_j \bar{D}_k + \frac{\lambda''_{ijk}}{2} \bar{U}_i \bar{D}_j \bar{D}_k + \mu_{L_i} L_i H_u \quad , \quad (1.1)$$

and corresponding SUSY breaking terms in the soft Lagrangian

$$\mathcal{L}_{SUSY} \supset \frac{\mathcal{A}_{ijk}}{2} \tilde{L}_i \tilde{L}_j \tilde{\bar{E}}_k + \mathcal{A}'_{ijk} \tilde{Q}_i \tilde{L}_j \tilde{\bar{D}}_k + \frac{\mathcal{A}''_{ijk}}{2} \tilde{\bar{U}}_i \tilde{\bar{D}}_j \tilde{\bar{D}}_k + \mathcal{B}_i \tilde{L}_i H_u + h.c. \quad , \quad (1.2)$$

which induce rapid proton decay and unsuppressed FCNCs if the couplings in eqs. (1.1) and (1.2) are of natural size. Since proton decay typically requires both baryon and lepton number violation, the most stringent constraints can be evaded if leptonic RPV is strongly suppressed, but baryonic RPV via $\bar{U} \bar{D} \bar{D}$ is large enough to allow the lightest squarks to decay promptly without MET [1, 2].

Several models in the literature satisfy these criteria. Minimal Flavor Violating (MFV) SUSY [3, 4], for example, constrains all flavor violating processes with the appropriate Yukawa couplings, which also determine the size and scope of allowed RPV interactions. However, maintaining MFV structure in a UV complete scenario requires nontrivial model building [5–7]. Similarly, “Collective RPV” [8] only allows RPV in particular combinations of couplings, so their overall effect yields the requisite suppression. Other models with similar features are found in [9–14].

Here we propose a novel scenario in which baryonic RPV arises at tree level in the *soft terms*, but the scalar-fermion RPV interactions in figure 1 arise only at loop level with additional suppression. These loop suppressed couplings can still be dangerous if RPV $\mathcal{A}$ -terms are of order the weak scale. For instance, if the baryon number violating $\mathcal{A}$ -term ($\mathcal{A}''$) is comparable to a typical soft mass m_s ,

$$\lambda'' \simeq \frac{g_s^2}{16\pi^2} \frac{\mathcal{A}''}{m_s} \sim 10^{-2} \quad , \quad (1.3)$$

this effective scalar-fermion coupling is ruled out by precision flavor constraints, which require $\lambda'' \lesssim 10^{-7}$ for light flavors [15–17]. However, if these terms are generated effectively through a heavy mediator of mass M that ensures $\mathcal{A}'' \sim m_s^2/M$, then the amount of RPV is controlled dynamically. In this framework, viable soft RPV can arise when M is of order the SUSY breaking scale, so no additional hierarchy is required. Although some aspects of

soft RPV interactions have been studied from a phenomenological perspective in [18–22], to our knowledge, a realistic model has never been realized before.

Our model features a global R -symmetry that forbids RPV interactions in the superpotential. This symmetry is broken in a hidden sector and communicated to the MSSM through a heavy mediator that gets integrated out to induce effective RPV $\mathcal{A}$ -terms for squarks.¹ If gauge mediation communicates SUSY breaking to the visible sector, the spectrum will also feature a metastable gravitino LSP that can be a viable dark matter candidate if thermally produced in the early universe.

The outline of this paper is as follows: in section 2, we list the general criteria for soft RPV and present a concrete model based on gauge mediation; in section 3 we consider the experimental constraints and map out the allowed parameter space; and in section 4 we make some concluding remarks.

2 Model description

On general grounds, a viable model of soft RPV requires:

- Some symmetry G that forbids the usual RPV interactions in the visible sector.
- A hidden sector (generically distinct from the SUSY breaking sector) that interacts with visible fields through a heavy mediator.
- G -breaking triggered by soft terms in the hidden sector.

When the mediator is integrated out, the effective superpotential becomes

$$W_{eff} \supset \frac{X}{M} \mathcal{O}_{vis} + X F_{\mathcal{G}} , \quad (2.1)$$

where M is the heavy mediator mass, X is a hidden sector superfield, and $F_{\mathcal{G}}$ is a G breaking spurion. The F-term for X induces a G -breaking $\mathcal{A}$ -term $\sim F_{\mathcal{G}}/M$ for visible sector scalars, while RPV interactions involving only visible fermions are forbidden at tree level when $\langle \tilde{X} \rangle = 0$.

In this section we present a concrete model in which G is an R -symmetry. To ensure predominantly baryonic RPV in the effective theory, we need lepton number to remain a good, accidental symmetry even after R -breaking. Fortunately this can be accomplished with an appropriate choice of hidden sector fields. However, SUSY breaking typically contributes an additional source of R -breaking, so we need to ensure that the mediation mechanism doesn't spoil the accidental lepton symmetry. Thus, we will use gauge mediation to communicate SUSY breaking to both visible and hidden sectors; perturbative gauge interactions preserve both lepton and baryon number, so leptonic RPV will not arise after R -breaking.

2.1 Soft RPV from a broken R -symmetry

Since R -symmetries are vital for generic SUSY breaking [24], we begin by imposing the following R -charge assignments for MSSM fields

$$R[Q, \bar{U}, \bar{D}] = 1, \quad R[L] = 4/3, \quad R[N, \bar{E}] = 2/3, \quad R[H_u, H_d] = 0 , \quad (2.2)$$

¹A global R -symmetry can also yield purely leptonic RPV operators [23] in the superpotential.

Figure 2. A schematic diagram of the relevant sectors. SUSY breaking is communicated to both the R -breaking and visible sectors through gauge mediation.

which forbid the RPV interactions in eq. (1.1) without imposing R -parity. Although this choice of R -charges is anomalous, heavy spectators can be added to cancel this anomaly without spoiling any of the model's features. The MSSM μ term is also forbidden at tree level, but one can arise if an additional singlet S with R -charge $+2$ gets a VEV to induce $\langle S \rangle H_u H_d$ in the superpotential. It is also possible to generate weak scale higgsino and (Dirac) gaugino masses with an *unbroken* R -symmetry, though additional electroweak doublets are required [25, 26]. Since the novel features of our model do not depend on the details of the Higgs sector, we leave this issue for future work.

The model contains three sectors depicted schematically in figure 2:

- *Visible sector:* contains the usual MSSM fields and interactions consistent with the R -symmetry, which forbids RPV.
- *SUSY breaking sector:* breaks both SUSY and the R -symmetry. SUSY breaking is mediated to the other sectors by gauge fields and decouples when all the gauge couplings vanish.
- *Soft R -breaking hidden sector:* features an additional $U(1)_H$ gauge symmetry so hidden scalars get soft masses from gauge mediation. These soft masses can explicitly break the R -symmetry or induce radiative symmetry breaking through renormalization group evolution. R -breaking in this sector is communicated to the visible fields by heavy mediators $\mathcal{D}$ and $\bar{\mathcal{D}}$.

Even though the R -symmetry is also generically broken in the SUSY breaking sector, perturbative gauge interactions preserve both lepton and baryon number, so R -parity is not violated by gauge mediation. Visible sector RPV can only arise if the mediator connecting the visible and R -breaking sectors carries either lepton or baryon number. In principle, the SUSY breaking and hidden sectors may be merged, but, for simplicity of exposition we ignore this possibility here.

	$SU(3)_c$	$U(1)_Y$	$U(1)_H$	R
$\bar{U}$	$\bar{3}$	$-2/3$	0	1
$\bar{D}$	$\bar{3}$	$1/3$	0	1
$\bar{\mathcal{D}}$	$\bar{3}$	$1/3$	0	0
$\mathcal{D}$	3	$-1/3$	0	2
X	1	0	0	-1
Σ	1	0	1	$3/2$
$\bar{\Sigma}$	1	0	-1	$3/2$

Figure 3. The charge assignments in our model. From top to bottom: the right-handed quarks in the visible sector, the heavy mediators $\mathcal{D}$ $\bar{\mathcal{D}}$, the singlet X connects the mediators to the Σ fields, which are charged under the gauged $U(1)_H$. The rightmost column lists R -charge assignments.

For the field content and charge assignments in table 3, the most general, renormalizable superpotential for the new states is

$$\kappa_{ij} \epsilon^{abc} \bar{U}_a^i \bar{D}_b^j \bar{\mathcal{D}}_c + \kappa'_i \bar{D}^i \mathcal{D} X + \eta \Sigma \bar{\Sigma} X + M_{\mathcal{D}} \bar{\mathcal{D}} \mathcal{D} \quad , \quad (2.3)$$

where a, b, c are color indices and i, j are flavor indices. For $M_{\mathcal{D}} \gg m_s$, the heavy mediators $\mathcal{D}$ and $\bar{\mathcal{D}}$ are integrated out and the effective superpotential becomes

$$- \frac{\kappa_{i[j} \kappa'_{k]}}{M_{\mathcal{D}}} \epsilon^{abc} \bar{U}_a^i \bar{D}_b^j \bar{D}_c^k X + \eta \Sigma \bar{\Sigma} X \quad , \quad (2.4)$$

where the j and k indices are antisymmetrized. If the scalar component of X gets a vacuum expectation value (VEV), there will be baryonic RPV in both the soft terms and in the effective superpotential. To emphasize the novel features of this model, we assume $\langle \tilde{X} \rangle = 0$ without essential loss of generality; we revisit this assumption in section 2.3. The effective scalar potential now contains

$$|F_X|^2 \supset - \frac{\kappa_{i[j} \kappa'_{k]} \eta^*}{M_{\mathcal{D}}} (\tilde{\Sigma} \tilde{\Sigma})^* \tilde{U}^i \tilde{D}^j \tilde{D}^k + c.c. \quad , \quad (2.5)$$

and baryonic RPV arises from a $\tilde{\Sigma}$ and $\tilde{\Sigma}$ loop with a $\mathcal{B}$ -term ($\mathcal{B}_{\Sigma}$) insertion in figure 4(a) or from Σ and $\bar{\Sigma}$ VEVs (v_{Σ}), which generate the diagram in figure 4(b). Note that the R -charges in eq. (2.2) are chosen to forbid the baryon and lepton number violating interaction $QL\bar{\mathcal{D}}$, which generates $QL\bar{D}$ when the mediator is integrated out. Lepton number remains a good accidental symmetry after R breaking, so the only B violation comes from $\bar{U}\bar{D}\bar{D}$ and higher order terms of the form $(\bar{U}\bar{D}\bar{D})^n$ with the same gauge-index contractions.² Since

²Since gravity violates all global and discrete symmetries, Planck suppressed operators — e.g. $\frac{1}{M_{pl}} QQQ L$ and $\frac{1}{M_{pl}} \bar{U}\bar{U}\bar{D}\bar{E}$ — can still be dangerous if their coefficients are not suppressed [7]. As in the R -parity conserving MSSM, we assume these to be negligible or absent in a full theory valid at the Planck scale.

Figure 4. Effective λ'' couplings from a nonzero $\mathcal{B}$ term (a) and from spontaneous R -breaking (b). Diagrams with electroweak gauginos in place of gluinos also give subdominant contributions to this process.

Figure 5. The hierarchies of scales in our model. Since $\mathcal{A}'' \sim \mathcal{B}_\Sigma/M_{\mathcal{D}} \sim v_\Sigma^2/M_{\mathcal{D}}$ and $M_{\mathcal{D}} \sim \sqrt{F}$, this setup introduces no energy scales beyond those already required in conventional SUSY models.

these higher order terms induce the same B violating phenomena as the leading-order $\bar{U}\bar{D}\bar{D}$ interaction, we will ignore them in our analysis of experimental constraints.

The R -charge assignments eq. (2.2) allow the $H_u LN$ operator, but forbid the Majorana term MNN , so Dirac neutrino masses can be generated without violating lepton number. If, instead, $L, \bar{E}$ and N were all assigned unit R -charges, an additional discrete or flavor symmetry would be necessary to forbid lepton violation in the IR. Examples of such constructions can be found in [4–6] and [27].

Since gauge mediation communicates SUSY breaking to both visible and hidden sectors, the essential features of this model are insensitive to the details of SUSY breaking and the field content of the messenger sector. These details will, however, determine the relative sizes of $\mathcal{B}_\Sigma$ and v_Σ , so for the remainder of this paper we will remain agnostic about which diagram in figure 4 dominates and consider only the limiting cases in which only $\mathcal{B}_\Sigma$ or v_Σ is nonzero. The general case with both contributions merely interpolates between these extremes, so our approach loses no essential generality. The ladder of scales in figure 5 summarizes the relative sizes of various inputs in our model for the convenience of the later discussion.

2.2 $\mathcal{B}$ -term R -breaking

As a warmup to see the essential features of the model, we first consider a toy situation in which all R -breaking arises from a nonzero $\Sigma\bar{\Sigma}$ $\mathcal{B}$ -term, but the $U(1)_H$ remains unbroken. The $\tilde{\Sigma}$ and $\bar{\tilde{\Sigma}}$ scalars get positive soft masses (m_Σ) from gauge mediation so $v_\Sigma = v_{\bar{\Sigma}} = 0$, visible sector RPV arises from the effective $\mathcal{A}$ term

$$\mathcal{A}_{ijk}'' \simeq \frac{\kappa_{i[j}\kappa'_{k]}\eta^*}{16\pi^2} \frac{\mathcal{B}_\Sigma}{M_\mathcal{O}} \log \frac{M_*^2}{m_\Sigma^2}, \quad (2.6)$$

where M_* is the messenger scale, so the diagram in figure 4(a) yields

$$\lambda_{ijk}'' \simeq \frac{\kappa_{i[j}\kappa'_{k]}\eta^* g_s^2 \mathcal{B}_\Sigma}{(16\pi^2)^2 M_\mathcal{O} M_{\tilde{g}}} \log \frac{M_*^2}{m_\Sigma^2}. \quad (2.7)$$

Although there are flavor-specific bounds on λ'' , in principle, the index structure in $\kappa_{i[j}\kappa'_{k]}$ can be flavor anarchic. In section 3, we will see that, for order one $\kappa_{ij}, \kappa'_{k'}$, and η with benchmark inputs $M_\mathcal{O} \sim 10^4$ TeV, $M_* \sim 10^9$ TeV, $\sqrt{\mathcal{B}_\Sigma} \sim M_{\tilde{g}} \sim m_\Sigma \sim 1$ TeV, the baryonic RPV coupling λ'' is naturally of order 10^{-7} and safe from experimental constraints; additional flavor structure in eq. (2.7) enlarges the viable parameter space, but is not an essential feature of this mechanism and lies beyond the scope of this work.

While fermion mass terms for $\Sigma, \bar{\Sigma}$ and X are forbidden at tree level, a Dirac mass μ_Σ arises from hidden gaugino (λ_H) interactions at one loop in figure 6,

$$\mu_\Sigma = \frac{g_H^2 \mathcal{B}_\Sigma M_{\lambda_H}}{16\pi^2 (M_{\lambda_H}^2 - m_\Sigma^2)^2} \left(M_{\lambda_H}^2 - m_\Sigma^2 + M_{\lambda_H}^2 \log \frac{m_\Sigma^2}{M_{\lambda_H}^2} \right) \sim \frac{g_H^2}{16\pi^2} \frac{\mathcal{B}_\Sigma}{M_{\lambda_H}}, \quad (2.8)$$

where M_{λ_H} is the hidden gaugino mass and the last expression is true in the $M_{\lambda_H} \sim m_\Sigma$ regime. Although the SUSY breaking parameter M_{λ_H} appears in the denominator of eq. (2.7), this diagram vanishes in the limit where SUSY is restored. An X fermion mass $\sim \mu_\Sigma/16\pi^2$ also arises with additional loop suppression from a similar diagram with $\Sigma, \bar{\Sigma} \rightarrow X$, $\lambda_H \rightarrow \Sigma$ and $M_{\lambda_H} \rightarrow \mu_\Sigma$. In this phase, the dark gauge symmetry is unbroken, so the stable Σ fermions annihilate to dark radiation in the early universe. The X fermions decay promptly through the $X\bar{U}\bar{D}\bar{D}$ operator so long as they are heavier than the proton. If they are lighter than the ~ 10 MeV gravitino dark matter candidate (see section 3.4), they can contribute to the dark matter abundance without overclosing the universe.

For $v_\Sigma = 0$ in this minimal setup, the X scalar is massless at tree level and acquires a tachyonic mass from loops of Σ and $\bar{\Sigma}$ fermions. The resulting VEV generates potentially large superpotential RPV via $\langle X \rangle \bar{U}\bar{D}\bar{D}$, so this toy scenario is unstable unless $\tilde{X}$ acquires mass by other means. Additional mass terms for X can arise either in the superpotential with additional R -charged fields or after SUSY breaking if the mediation mechanism gives gauge-singlets soft masses.

2.3 Spontaneous R -breaking

Now we present a more concrete scenario that generates soft RPV and solves this problem with nonzero VEVs $v_{\Sigma, \bar{\Sigma}}$ that break both $U(1)_H$ and the R -symmetry. For simplicity

Figure 6. Dirac mass μ_Σ for Σ and $\bar{\Sigma}$ from a nonzero $\mathcal{B}_\Sigma$ term.

we assume all hidden sector $\mathcal{A}$ and $\mathcal{B}$ terms vanish and set $m_\Sigma = m_{\bar{\Sigma}}$, so the scalar potential contains

$$\frac{g_H^2}{2} \left(|\tilde{\Sigma}|^2 - |\tilde{\bar{\Sigma}}|^2 \right)^2 + \eta^2 \left(|\tilde{\Sigma} \tilde{\bar{\Sigma}}|^2 + |\tilde{X} \tilde{\Sigma}|^2 + |\tilde{X} \tilde{\bar{\Sigma}}|^2 \right) - m_\Sigma^2 \left(|\tilde{\Sigma}|^2 + |\tilde{\bar{\Sigma}}|^2 \right) , \quad (2.9)$$

where the negative mass squared can arise through RG evolution if Σ and $\bar{\Sigma}$ couple to other fields with nonzero soft masses – see appendix A for a concrete example.

For $g_H^2 > \eta^2/2$, the classical minimum is

$$v_\Sigma = v_{\bar{\Sigma}} = m_\Sigma/\eta \quad , \quad \langle \tilde{X} \rangle = 0 \quad , \quad (2.10)$$

but, quantum corrections still generate an $\tilde{X}$ VEV. However, unlike in section 2.2, $\tilde{X}$ now has a tree level mass of $m_X \sim v_\Sigma$, so minimizing the Coleman-Weinberg potential yields $\langle \tilde{X} \rangle \propto \mu_\Sigma^3/m_X^2$, where

$$\mu_\Sigma = \frac{g_H^2 \eta^2 v_\Sigma^2 M_{\lambda_H}}{16\pi^2 (M_{\lambda_H}^2 - m_\Sigma^2)^2} \left(M_{\lambda_H}^2 - m_\Sigma^2 + M_{\lambda_H}^2 \log \frac{m_\Sigma^2}{M_{\lambda_H}^2} \right) \sim \frac{g_H^2}{16\pi^2} \frac{\eta^2 v_\Sigma^2}{M_{\lambda_H}} , \quad (2.11)$$

is the $\Sigma\bar{\Sigma}$ Dirac mass that arises from the loop-diagram in figure 7. As in eq. (2.8), we work in the $M_{\lambda_H} \sim m_\Sigma$ limit and there is a SUSY breaking mass in the denominator. However, the diagram that generates μ_Σ vanishes when SUSY is restored since $M_\Lambda \rightarrow 0$ and $v_\Sigma \rightarrow 0$ as the latter requires $m_\Sigma > 0$ in eq. (2.9), which fails in this limit. Thus, the $\langle X \rangle \bar{U} \bar{D} \bar{D} / M_\mathcal{D}$ correction to fermionic RPV is subdominant to the soft contribution in figure 4(b) for which

$$\lambda''_{ijk} = \frac{\kappa_{i[j} \kappa'_{k]} \eta^* g_s^2 v_\Sigma^2}{32 \pi^2 M_\mathcal{D} M_{\tilde{g}}} . \quad (2.12)$$

Since the effective potential also contains

$$|F_X|^2 = \left| \eta \tilde{\Sigma} \tilde{\bar{\Sigma}} - \frac{\kappa \kappa'}{M_\mathcal{D}} \tilde{U} \tilde{D} \tilde{D} \right|^2 , \quad (2.13)$$

a nonzero v_Σ can, in principle, trigger color breaking, however, in appendix B we find that color remains unbroken so long as $v_\Sigma \lesssim M_\mathcal{D} (m_{\tilde{q}}/M_\mathcal{D})^{3/4}$, where $m_{\tilde{q}}$ is a typical squark mass of order the weak scale.

Figure 7. Loop level Dirac mass μ_{Σ} for Σ and $\bar{\Sigma}$ for $v_{\Sigma} \neq 0$.

The R -symmetry forbids superpotential mass terms for X , Σ and $\bar{\Sigma}$, so the hidden sector spectrum is entirely determined by SUSY breaking parameters. As in section 2.2, $\tilde{\Sigma}$ and $\tilde{\Sigma}$ get gauge mediated soft masses and $\tilde{X}$ gets a soft mass at one loop. After symmetry breaking, the $X, \Sigma, \bar{\Sigma}$, and λ_H fermions mix and the resulting mass eigenstates are of order the electroweak scale. For generic mixing angles, all hidden sector mass eigenstates will be linear combinations of all four interaction eigenstates, so they all decay promptly through the $X\bar{U}\bar{D}\bar{D}$ portal.

A spontaneously broken R -symmetry gives rise to a massless R -axion that can accelerate supernova cooling and cause cosmological problems [28]. Conventionally, R -breaking arises only in the SUSY breaking sector and the BPR mechanism [29] generates an R -axion mass from a constant term in the superpotential introduced to cancel the cosmological constant. In this scenario, our hidden sector also contributes to R -breaking, so the physical R -axion is now a linear combination of SUSY breaking and hidden sector states, but still acquires a BPR mass, so we will not consider it further. Although the BPR term explicitly breaks the R -symmetry, we assume its existence has no additional bearing on the symmetries of our superpotential; it serves merely as a placeholder for the cosmological constant problem, which is beyond the scope of this work.

3 Experimental bounds

In this section we consider the experimental constraints on our realization of soft RPV. For simplicity, we will follow the organization of section 2.1 and separately constrain the cases in which the B-term and $\Sigma, \bar{\Sigma}$ VEVs are solely responsible for R -breaking; the most general case interpolates between these extremes. The plots in figure 8 carve out the allowed parameter space in both $\mathcal{B}$ -term and spontaneously broken scenarios.

3.1 Direct production

Although the parameter space for RPV spectra with sparticles below a TeV has recently been reduced, the sensitivity of these bounds is driven primarily by lepton number violating processes. For purely baryonic RPV, the bounds are considerably weaker and can accommodate natural stops with ~ 100 GeV masses, provided they decay predominantly to dijets

Figure 8. The parameter space for our model in the $\mathcal{B}$ -term scenario (left) and in the spontaneously broken phase (right). In each case, the light (dark) green represents the allowed region where the stop decays with vertices smaller than 2 mm (10 cm). Here we assume the most conservative scenario with $|\kappa_{ij}| = |\kappa'_i| = |\eta| = 1$ for all coefficients. The rates that determine the red excluded regions are quadratically sensitive to these parameters, so if light flavors have smaller coefficients, the parameter space expands considerably.

via $\bar{U}\bar{D}\bar{D}$ [1]. For RPV gluinos decaying exclusively to $\tilde{g} \rightarrow t\bar{t}$, the strongest experimental bound is now $\gtrsim 670$ GeV [30–35], however, recasting R -parity conserving SUSY searches may place a stronger ~ 800 GeV bound on the gluino mass [36].

3.2 Baryon number violation

The $\bar{U}\bar{D}\bar{D}$ interaction explicitly violates baryon number, so our model faces constraints from the null results of several low energy searches. The strongest limits come from the bounds on the characteristic timescales for dinucleon decay ($pp \rightarrow K^+K^+$) [37] and neutron-antineutron oscillation ($n - \bar{n}$) [38]

$$\tau_{pp \rightarrow KK} \geq 1.7 \times 10^{32} \text{ yrs.}, \quad \tau_{n-\bar{n}} \geq 2.44 \times 10^8 \text{ sec.}, \quad (3.1)$$

and from proton decay via $p \rightarrow K^+\nu$, for which the bound is [39]

$$\tau_{p \rightarrow K^+\nu} \geq 2.3 \times 10^{33} \text{ yrs.} \quad (3.2)$$

Although our model doesn't violate lepton number, this bound conservatively constrains the $p \rightarrow K^+\tilde{G}$ decay, which has similar kinematics for a sufficiently light gravitino.

3.2.1 Dinucleon decay

Following Goity and Sher [16], the dinucleon decay rate for the dominant processes shown in figure 9 is

$$\Gamma_{pp \rightarrow KK} \sim \rho_N \frac{128 \pi \alpha_s^2 \Lambda^{10}}{m_p^2 m_{\tilde{u}}^8 M_{\tilde{g}}^2} (\lambda''_{uds})^2, \quad (3.3)$$

where $m_{\tilde{u}}$ is the lightest up-type squark mass, $\rho_N \sim 0.25/\text{fm}^3$ is the density of nuclear matter and Λ is the characteristic hadronic energy-scale. Here we assume $M_{\tilde{C}} > M_{\tilde{g}} \alpha/\alpha_s \gtrsim$

Figure 9. Dinucleon decay via baryonic RPV interactions. In the text we assume gluino exchange (left) dominates.

220 GeV, so the gluino exchange diagram in figure 9 dominates. Thus, satisfying the experimental bound $\tau_{pp \rightarrow KK} \geq 1.7 \times 10^{32}$ yrs. requires

$$\lambda''_{uds} \lesssim 2.5 \times 10^{-7} \left(\frac{150 \text{ MeV}}{\Lambda} \right)^{5/2} \left(\frac{M_{\tilde{g}}}{800 \text{ GeV}} \right)^{1/2} \left(\frac{m_{\tilde{u}}}{500 \text{ GeV}} \right)^2, \quad (3.4)$$

Translating this into a constraint on the $\mathcal{B}$ -term scenario ($v_\Sigma = 0$) in section 2.2, we have

$$\frac{\mathcal{B}_\Sigma}{M_\mathcal{Q}} \lesssim 81 \text{ MeV} \left(\frac{150 \text{ MeV}}{\Lambda} \right)^{5/2} \left(\frac{M_{\tilde{g}}}{800 \text{ GeV}} \right)^{1/2} \left(\frac{m_{\tilde{u}}}{500 \text{ GeV}} \right)^2 (\eta^* \kappa_{u[d\kappa'_s]})^{-1}, \quad (3.5)$$

where we have set $M_* = 10^9 \text{ GeV}$ and $m_\Sigma = 500 \text{ GeV}$ inside the log of eq. (2.7). Similarly, for the spontaneous R -breaking scenario ($\mathcal{B}_\Sigma = 0$) in section 2.3, the corresponding bound is extracted from eq. (2.12)

$$\frac{v_\Sigma^2}{M_\mathcal{Q}} \lesssim 42 \text{ MeV} \left(\frac{150 \text{ MeV}}{\Lambda} \right)^{5/2} \left(\frac{M_{\tilde{g}}}{800 \text{ GeV}} \right)^{1/2} \left(\frac{m_{\tilde{u}}}{500 \text{ GeV}} \right)^2 (\eta^* \kappa_{u[d\kappa'_s]})^{-1}. \quad (3.6)$$

Unlike similar processes in MFV SUSY [4] where the light quark couplings are Yukawa suppressed, our setup imposes no necessary hierarchies in the RPV couplings.

3.2.2 $n - \bar{n}$ oscillation

Unlike dinucleon decay, $n - \bar{n}$ oscillation also requires flavor violation from R -parity conserving vertices. However, aside from the baryon violating $\mathcal{A}$ -term, all visible sector soft masses arise directly from gauge mediation, so their flavor structure comes entirely from Yukawa couplings. Thus, up to an overall coefficient, our $n - \bar{n}$ oscillation amplitudes are identical to those computed in [4].

Chirality-preserving flavor-violating masses arise predominantly from MSSM F -terms after SUSY and electroweak symmetry breaking through

$$\tilde{Q}^\dagger \left(v_u^2 Y_u Y_u^\dagger + v_d^2 Y_d Y_d^\dagger \right) \tilde{Q}, \quad (3.7)$$

and similar terms for $\tilde{U}$ and $\tilde{D}$, where $Y_{u,d}$ are Yukawa matrices. For simplicity, we take the Higgs doublet VEVs $v_{u,d}$ to be at the soft scale $\sim m_\mathcal{S}$. In gauge mediation, chirality flipping

Figure 10. The dominant diagram for neutron anti-neutron oscillation. All R -parity conserving soft masses and $\mathcal{A}$ -terms are consistent with MFV.

$\mathcal{A}$ -terms arise only at higher order and suffer both Yukawa and loop suppression, so they are typically smaller than soft masses. However, different realizations of gauge mediation give rise to $\mathcal{A}$ terms with different degrees of suppression relative to the soft scale. Since we remain agnostic about the details of the messenger sector, we conservatively parametrize any possible suppression with the general ansatz $\mathcal{A} \equiv \epsilon m_s$.

Putting all the squarks at a common soft mass $m_{\tilde{q}} \sim m_s$, the amplitude for the dominant diagram shown in figure 10 is

$$\mathcal{M}_{n-\bar{n}} \sim g_s^2 \epsilon^2 \lambda^6 \Lambda \left(\frac{\Lambda}{m_{\tilde{q}}} \right)^4 \left(\frac{\Lambda}{M_{\tilde{g}}} \right) (\lambda''_{udb})^2, \quad (3.8)$$

where $\lambda \simeq 0.23$ comes from the approximate CKM matrix parametrization in [4]. The oscillation timescale is approximately $\tau_{n-\bar{n}} \sim \mathcal{M}^{-1}$, thus the experimental bound $\tau_{n-\bar{n}} \geq 2.44 \times 10^8$ sec. requires

$$\lambda''_{udb} \lesssim 1.7 \times 10^{-6} \epsilon^{-2} \left(\frac{m_{\tilde{q}}}{500 \text{ GeV}} \right)^4 \left(\frac{250 \text{ MeV}}{\Lambda} \right)^6 \left(\frac{M_{\tilde{g}}}{800 \text{ GeV}} \right), \quad (3.9)$$

which is weaker than the bound from dinucleon decay in eq. (3.4) even when ϵ is order one.

3.2.3 Proton decay

Since gauge-mediation typically features a light, sub-GeV gravitino, proton decay to $K^+ \tilde{G}$ through the diagram in figure 11 may be kinematically allowed. The rate for this process is

$$\Gamma_{p \rightarrow K^+ \tilde{G}} \sim \frac{m_p}{8\pi} \left(\frac{\Lambda}{m_{\tilde{u}}} \right)^4 \left(\frac{\Lambda^2}{\sqrt{3} m_{3/2} M_{pl}} \right)^2 (\lambda''_{uds})^2, \quad (3.10)$$

and the lifetime for this channel must be longer than 2.3×10^{33} yrs., so the gravitino mass bound is

$$m_{3/2} \geq 4.7 \text{ MeV} \left(\frac{\Lambda}{250 \text{ MeV}} \right)^4 \left(\frac{500 \text{ GeV}}{m_{\tilde{u}}} \right)^2 \left(\frac{\lambda''_{uds}}{10^{-7}} \right), \quad (3.11)$$

For $m_{3/2} \gtrsim 5 \text{ MeV}$, this implies a lower bound on the SUSY breaking scale

$$\sqrt{F} \gtrsim 3.2 \times 10^5 \text{ TeV} . \quad (3.12)$$

If minimal gauge mediation gives rise to soft masses, the messenger scale M_* must also satisfy

$$M_* \gtrsim 1.3 \times 10^9 \text{ TeV} \left(\frac{500 \text{ GeV}}{m_s} \right) . \quad (3.13)$$

3.3 Displaced vertices

To avoid MET searches at the LHC, sparticles must decay on collider timescales, so there is an upper bound on the lightest squark's lifetime. Although there are many LHC searches for displaced vertices [30, 40], hadronically-decaying long-lived particles are significantly harder to constrain [41]; viable decay lengths can even exceed $\sim 10 \text{ cm}$, so a dedicated search is necessary. Given these uncertainties, we consider the experimental bounds in two regimes: for prompt decays, we conservatively require decay lengths $\ell_{\tilde{q}} < 2 \text{ mm}$; for signatures with viable displaced vertices, we demand $\ell_{\tilde{q}} < 10 \text{ cm}$, so most sparticles decay inside the tracker before reaching the hadronic calorimeter (HCAL), but may still be found with a dedicated search.

The width for a hadronically decaying stop NLSP³ in its rest frame is

$$\Gamma_{\tilde{t} \rightarrow q\bar{q}} = \frac{m_{\tilde{t}}}{8\pi} \sin^2 \theta_t |\lambda''_{tq}|^2 , \quad (3.14)$$

where θ_t is the stop mixing angle. In the lab frame, the decay length is $\ell_t \simeq \gamma \Gamma_{\tilde{t} \rightarrow q\bar{q}}^{-1}$, where γ is the stop boost factor; for a 300 GeV stop and an 800 GeV gluino produced at rest, $\gamma \sim 2$. For the remainder of this section we assume, for simplicity, that $\gamma \sin^2 \theta_t = 1$.

Assuming the dominant stop decay is $\tilde{t} \rightarrow d\bar{s}$, the bound on λ''_{tds} is

$$\lambda''_{tds} > (0.26 - 1.8) \times 10^{-7} \left(\frac{300 \text{ GeV}}{m_{\tilde{t}}} \right)^{1/2} . \quad (3.15)$$

where the left and right numbers represent the bound assuming 10 cm and 2 mm displaced-vertex limits, respectively. For the $\mathcal{B}$ -term scenario ($v_\Sigma = 0, \mathcal{B}_\Sigma \neq 0$) in section 2.2, this implies

$$\frac{\mathcal{B}_\Sigma}{M_\mathcal{Q}} \gtrsim (8.3 - 58) \times \text{MeV} \left(\frac{m_{\tilde{g}}}{800 \text{ GeV}} \right) \left(\frac{300 \text{ GeV}}{m_{\tilde{t}}} \right)^{1/2} (\kappa_{t[d\kappa'_s]\eta^*})^{-1} , \quad (3.16)$$

with $m_\Sigma = 1 \text{ TeV}$ and $M_* \sim 10^9 \text{ TeV}$ inside the log in eq. (2.7). Similarly, for the spontaneously broken scenario ($v_\Sigma \neq 0, \mathcal{B}_\Sigma = 0$) in section 2.3, we have

$$\frac{v_\Sigma^2}{M_\mathcal{Q}} \gtrsim (4.3 - 31) \times \text{MeV} \left(\frac{m_{\tilde{g}}}{800 \text{ GeV}} \right) \left(\frac{300 \text{ GeV}}{m_{\tilde{t}}} \right)^{1/2} (\kappa_{t[d\kappa'_s]\eta^*})^{-1} . \quad (3.17)$$

³For typical SUSY breaking scales we consider, the gravitino is the LSP, though for extremely high SUSY breaking scales, this need not be the case.

Figure 11. Proton decay via $p \rightarrow K^+ \tilde{G}$.

These bounds assume the stop is the lightest squark and decays predominantly through RPV interactions. Thus, the only other kinematically allowed process $\tilde{t} \rightarrow t \tilde{G}$ must have a negligible branching ratio, which requires

$$\Gamma_{\tilde{t} \rightarrow q\bar{q}} \gg \Gamma_{t\tilde{G}} = \frac{m_{\tilde{t}}^5}{16\pi F^2} . \quad (3.18)$$

As long as the SUSY breaking scale satisfies $\sqrt{F} > 10^2 \text{ TeV}$, the RPV branching ratio exceeds 99%. This constraint is trivially satisfied by considerations from proton decay in section 3.2.3 above.

3.4 Gravitino dark matter

Since gauge mediation communicates SUSY breaking to the visible sector, the gravitino is the LSP with mass $m_{3/2} \sim F/M_{pl} \sim \mathcal{O}(10) \text{ MeV}$ for $\sqrt{F} \sim 10^8 \text{ GeV}$. In this mass range $m_{3/2} < m_p$, so the process $\tilde{G} \rightarrow qqq$ is kinematically forbidden and the gravitino is stable. Since sparticles rarely decay to gravitinos and their annihilation rate is suppressed by the SUSY breaking scale, the present day abundance is thermally generated [13]

$$\Omega_{3/2} h^2 \simeq 0.1 \left(\frac{T_R}{10^5 \text{ GeV}} \right) \left(\frac{m_{3/2}}{20 \text{ MeV}} \right)^{-1} \left(\frac{M_{\tilde{g}}}{800 \text{ GeV}} \right)^2 , \quad (3.19)$$

where T_R is the reheating temperature, so the RPV gravitino is a viable dark matter candidate. Although the scenario in [13] assumes R -parity conservation, the analysis is applicable so long as thermal gravitino production in the early universe is not upset by additional decays or annihilation processes.

Thermal gravitino dark matter can also arise in models of spontaneous bilinear RPV [42–44]. In that scenario, the gravitino is metastable on cosmological timescales, but can decay leptonically at late times and be observed through indirect detection. Our basic model features a stable gravitino and only generates baryonic RPV; lepton number remains a good accidental symmetry, so this potential signature is absent. However, in the presence of additional R -breaking states or an alternative SUSY breaking mediation mechanism, it is possible to generate small leptonic RPV to allow gravitino decay, but this is beyond the scope of the present work.

4 Conclusions

In this paper we have presented a new realization of weak scale SUSY with R -parity violation. Unlike conventional scenarios, suppressed baryonic RPV arises in the *soft terms* when an R -symmetry is broken in a hidden sector and a heavy mediator is integrated out; lepton number remains a good accidental symmetry. RPV interactions between quarks and squarks arise at one loop and receive additional suppression. The model features light ($\sim$ few 100 GeV) squarks that decay promptly to hadrons and evade LHC searches in viable regions of parameter space safe from flavor constraints. Without additional, colored field content, RG evolution from a high UV scale naturally yields comparable squark and gluino masses in the IR, so there is tension in maintaining light squarks to preserve naturalness while keeping gluinos near the TeV scale to evade stronger collider bounds. [45] Mild tuning or additional field content may be required for a fully natural spectrum, but a detailed analysis is beyond the scope of this work.

For weak-scale R -breaking, the heavy mediator masses can be near the SUSY breaking scale $\sqrt{F} \sim 10^8$ GeV to generate RPV couplings with the requisite suppression, so the model requires no new scales beyond those already present in conventional SUSY models. If gauge mediation communicates SUSY breaking, the model also features a light $\sim 1 - 100$ MeV gravitino with a thermal abundance. For a reheating temperature of order 10^5 GeV and a weak scale gluino, a gravitino in this mass range is a viable dark matter candidate.

Like MFV SUSY [4], our soft-RPV scenario predicts hadronically decaying squarks and gluinos with high multiplicity multijet final states for all colored LHC-accessible states. Both frameworks predict predominantly baryonic RPV and can accommodate light squarks, but the flavor structure in our case is not fixed by the Yukawa and CKM matrices, so the branching ratios to different quark flavors can differ by order one amounts. Furthermore, our model is also easily falsifiable: if SUSY stabilizes the hierarchy problem and features baryonic soft-RPV, then squark or gluino resonances will be discovered at the LHC, while leptonic RPV will be absent.

Although small baryonic RPV ameliorates the tension between weak-scale sparticle masses and LHC direct production constraints, this feature, by itself, does not yield a ~ 126 GeV Higgs. This issue is well known and various solutions have been proposed in the literature [46–55]. Our scenario is a proof of concept to show that baryonic RPV arising from soft terms can evade production bounds; generating a viable Higgs mass in this framework is compatible with many of the known strategies. The benchmark example in this paper uses gauge mediation to generate soft masses, so stop mixing is small, but other tools including non-decoupling D-terms [50, 53], induced electroweak symmetry breaking [54], or heavy superpartners [55] can lift the Higgs mass without spoiling any of the model’s essential features. For example, following [50], we can extend the gauged $U(1)_H$ symmetry to the SM sector and include the Higgs in the D-term potential

$$\frac{g_H^2}{2} \left(q_H |H_u|^2 - q_H |H_d|^2 + |\tilde{\Sigma}|^2 - |\tilde{\bar{\Sigma}}|^2 \right)^2. \quad (4.1)$$

Since the Σ VEVs and masses are at the TeV scale, we can integrate them out after $U(1)_H$ is broken in eq. (2.11) to induce an extra contribution to the Higgs quartic coupling

$$\frac{g_H^2 q_H^2}{2} (|H_u|^2 - |H_d|^2)^2 \times \left(1 + \frac{g_H^2 v_\Sigma^2}{4 m_\Sigma^2}\right)^{-1}. \quad (4.2)$$

With these corrections, the Higgs having an MSSM potential can now obtain the required mass with an order one g_H used in appendix A. The few TeV scale $g_H v_\Sigma$ studied there also satisfy the bounds from the electroweak precision measurement and the Z' searches.⁴ The bound constraints the mixing between the SM and $U(1)_H$ gauge bosons. Although a detailed analysis is required, this can be a reasonable way to generate the observed Higgs mass even without a large stop A-term in a gauge mediated scenario. However, we emphasize that, for our purposes, gauge mediation is merely a convenient mechanism to generate soft masses without violating lepton or baryon number; any alternative for which this holds true is similarly viable and would be interesting to study.

If R -breaking arises from a $\mathcal{B}$ -term for Σ and $\bar{\Sigma}$ as in section 2.2, the model requires either non-minimal gauge mediation to generate sizable $\mathcal{B}$ -terms, or another mediation mechanism that preserves the accidental lepton symmetry. We leave these model building details for future work. For the more-concrete spontaneous R -breaking scenario in section 2.3, the model requires either additional fields to drive radiative symmetry breaking for Σ and $\bar{\Sigma}$ or an alternative to gauge mediation that results in tachyonic soft masses in the hidden sector. In appendix A we show that radiative symmetry breaking is feasible, but leave other alternatives for future work.

Grand unification with RPV is challenging because both lepton and baryon number violating RPV interactions generally arise from the same interaction term. In $SU(5)$, for instance, $\bar{U}\bar{D}\bar{D}, QL\bar{D}$ and $LL\bar{E}$ all live in the same $10\bar{5}\bar{5}$ UV operator, so generating predominantly baryonic RPV at low energies requires additional model building gymnastics [11]. In our case, the R -charge assignments differ for quark and lepton superfields, so it is not clear whether grand unification is possible.

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Addendum (January 2014). Since this paper was originally submitted for publication, the LHC constraints on SUSY scenarios have improved [56]. Nearly model independent bounds on gluino masses are now in the $\sim$ TeV range and squark mass-bounds have also

⁴See, for example, the figure 4 in [10].

improved considerably. In this light, the LHC motivation for baryonic RPV is somewhat diminished, though the essential features of the soft RPV scenario remain viable even at higher particle mass scales.

A Hidden sector VEVs

Throughout the paper, we have assumed that the Σ and $\bar{\Sigma}$ scalars acquire negative mass-squared parameters to induce spontaneous symmetry breaking. Since the minimal superpotential only allows the $\Sigma X \bar{\Sigma}$ interaction and gauge mediation gives rise to positive soft masses, the setup requires either a nonminimal messenger sector to generate negative soft masses or substantial RG evolution. Since we are agnostic about the details of gauge mediation, here we present a concrete example of radiative R -breaking in the hidden sector as a proof of concept.

If the Σ scalars also couple to chiral fields Y and $\bar{Y}$ with identical R -charges of $1/4$ and $U(1)_H$ charges of $\mp 1/2$, the superpotential also contains

$$W \supset \eta \Sigma X \bar{\Sigma} + \lambda_Y \Sigma Y^2 + \lambda_{\bar{Y}} \bar{\Sigma} \bar{Y}^2, \quad (\text{A.1})$$

where η , λ_Y , and $\lambda_{\bar{Y}}$ are order one parameters. Including $U(1)_H$ gauge interactions, the full set of RGEs is

$$\frac{dg_H}{dt} = \frac{5g_H^3}{32\pi^2} \quad (\text{A.2})$$

$$\frac{d\lambda_{Y,\bar{Y}}}{dt} = \frac{\lambda_{Y,\bar{Y}}}{16\pi^2} \left(\frac{5}{2} \lambda_{Y,\bar{Y}}^2 - 3g_H^2 \right) \quad (\text{A.3})$$

$$\frac{d\eta}{dt} = \frac{\eta}{16\pi^2} (3\eta^2 - 4g_H^2) \quad (\text{A.4})$$

$$\frac{dm_\Sigma^2}{dt} = \frac{1}{16\pi^2} \left[\eta^2 (2m_\Sigma^2 + m_{\bar{\Sigma}}^2 + m_X^2) + 4\lambda_Y^2 (m_\Sigma^2 + m_Y^2) + \frac{g_H^2}{2} (-2m_\Sigma^2 + m_Y^2 - m_{\bar{Y}}^2) \right] \quad (\text{A.5})$$

$$\frac{dm_{\bar{\Sigma}}^2}{dt} = \frac{1}{16\pi^2} \left[\eta^2 (2m_\Sigma^2 + m_{\bar{\Sigma}}^2 + m_X^2) + 4\lambda_{\bar{Y}}^2 (\bar{m}_\Sigma^2 + \bar{m}_Y^2) + \frac{g_H^2}{2} (-2m_\Sigma^2 + m_Y^2 - m_{\bar{Y}}^2) \right] \quad (\text{A.6})$$

$$\frac{dm_Y^2}{dt} = \frac{1}{16\pi^2} \left[\lambda_Y^2 (4m_\Sigma^2 + 6m_Y^2) + \frac{g_H^2}{2} \left(-m_\Sigma^2 + m_{\bar{\Sigma}}^2 - \frac{1}{2}m_Y^2 \right) \right] \quad (\text{A.7})$$

$$\frac{dm_{\bar{Y}}^2}{dt} = \frac{1}{16\pi^2} \left[\lambda_{\bar{Y}}^2 (4m_{\bar{\Sigma}}^2 + 6m_{\bar{Y}}^2) + \frac{g_H^2}{2} \left(-m_{\bar{\Sigma}}^2 + m_\Sigma^2 - \frac{1}{2}m_{\bar{Y}}^2 \right) \right] \quad (\text{A.8})$$

$$\frac{dm_X^2}{dt} = \frac{1}{16\pi^2} \left[\eta^2 (m_\Sigma^2 + m_{\bar{\Sigma}}^2 + 2m_X^2) \right] \quad (\text{A.9})$$

Note that, without the interactions in eq. (A.1), the $m_{\Sigma,\bar{\Sigma}}^2$ equations can be rewritten in terms of $x \equiv m_\Sigma^2 + m_{\bar{\Sigma}}^2$ so that both become $dx/dt \propto x$ whose solution never runs negative.

In figure 12, we plot contours of radiatively generated λ'' from eq. (2.12) in the $M_{\mathcal{Q}}, \lambda_Y$ plane. For each contour, minimal gauge mediation defines the UV boundary condition $m_\Sigma = \frac{g_H^2}{16\pi^2} \frac{F}{M_*}$ where F saturates the bound in eq. (3.12). The allowed region assumes all couplings $\eta, \lambda_{Y,\bar{Y}}, g_H$ are all unity and we choose $\eta = 0.1$ at the EW scale to generate a larger v_Σ and satisfy the bounds on λ'' from eqs. (3.4) and (3.15).

Figure 12. The allowed parameters space for $M_{\mathcal{D}}$ and λ_Y with contours of λ'' derived from eq. (2.7). The white region is excluded by the dinucleon decay and displaced vertex bounds in section 3. For the left plot, we assume prompt stop decays with lengths < 2 mm; for the right plot we assume displaced stop decays inside the tracker (< 10 cm). The VEVs are computed after RG evolution with a UV boundary condition at the messenger scale, $M_* = 10^9$ TeV, and IR boundary at the soft mass scale $m_s = 1$ TeV. We also assume flavor universal couplings $|\kappa|$ and $|\kappa'|$ and soft masses dictated by gauge mediation. Note that the range of $M_{\mathcal{D}}$ is of order the benchmark SUSY breaking scale $\sqrt{F} \sim 10^5$ TeV.

For this field content, radiative symmetry breaking requires $Y \bar{Y}$ to have larger soft masses than Σ and $\bar{\Sigma}$ at the mediation scale, which is not realized in the minimal gauge mediation; Σ and $\bar{\Sigma}$ have larger gauge charges. However this can be accommodated if the Y and $\bar{Y}$ carry additional gauge charges to give them larger soft masses at the mediation scale. Our example here assumes $m_{Y,\bar{Y}}(M_*) = 2m_{\Sigma,\bar{\Sigma}}(M_*)$ and suffices to demonstrate that radiative R -breaking is possible.

B Color breaking?

After R -breaking, up to order-one coefficients, the scalar potential derived from eq. (2.4) contains the terms

$$V \supset \left| \frac{\tilde{U} \tilde{D}_i \tilde{D}_j}{M_{\mathcal{D}}} + v_{\Sigma}^2 \right|^2 + m_u^2 |\tilde{U}|^2 + m_{\tilde{d}_i}^2 |\tilde{D}_i|^2 + m_{\tilde{d}_j}^2 |\tilde{D}_j|^2. \quad (\text{B.1})$$

which can break color if squark masses are too small. For simplicity, assuming identical squark soft masses and positive superpotential couplings, we can rewrite the potential in terms of dimensionless variables

$$\hat{V} \equiv \frac{m_{\tilde{d}_i}^2 m_{\tilde{d}_j}^2}{m_u^4 M_{\mathcal{D}}^4} V \supset \left| x y z + s^2 \right|^2 + \hat{m}^2 (x^2 + y^2 + z^2), \quad (\text{B.2})$$

where

$$x \equiv \frac{\langle \tilde{U} \rangle}{M_{\mathcal{D}}}, \quad y \equiv \frac{m_{\tilde{d}_i} \langle \tilde{D}_i \rangle}{m_u M_{\mathcal{D}}}, \quad z \equiv \frac{m_{\tilde{d}_j} \langle \tilde{D}_j \rangle}{m_u M_{\mathcal{D}}}, \quad s \equiv \frac{\sqrt{m_{\tilde{d}_i} m_{\tilde{d}_j}}}{m_u M_{\mathcal{D}}} v_{\Sigma}, \quad \hat{m} \equiv \frac{m_{\tilde{d}_i} m_{\tilde{d}_j}}{m_u M_{\mathcal{D}}}. \quad (\text{B.3})$$

At their extremal values, $x = y = z$, we demand

$$(x^3 + s^2)^2 + 3\hat{m}^2 x^2 \geq s^4, \quad (\text{B.4})$$

to avoid color breaking at the global minimum. This conditions implies, $s \lesssim \hat{m}^{3/4}$, so we need

$$v_\Sigma \lesssim \left(\frac{m_{\tilde{u}} m_{\tilde{d}_i} m_{\tilde{d}_j}}{M_{\mathcal{D}}^3} \right)^{1/4} M_{\mathcal{D}}. \quad (\text{B.5})$$

For the model's relevant parameter space, $m_{\tilde{d}} \gtrsim 500 \text{ GeV}$ and $M \simeq 10^5 \text{ TeV}$, this constraint becomes $v_\Sigma^2/M_{\mathcal{D}} \lesssim 10^{-4} \text{ TeV}$, which is an order of magnitude weaker than the dinucleon decay bound in eq. (3.6), so color remains unbroken for the viable parameter space we consider.

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