

Gauging the way to MFV

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ABSTRACT: We present a UV complete model with a gauged flavor symmetry which approximately realizes holomorphic Minimal Flavor Violation (MFV) in R -parity violating (RPV) supersymmetry. Previous work has shown that imposing MFV as an ansatz easily evades direct constraints and has interesting collider phenomenology. The model in this work spontaneously breaks the flavor symmetry and features the minimum “exotic” field content needed to cancel anomalies. The flavor gauge bosons exhibit an inverted hierarchy so that those associated with the third generation are the lightest. This allows low energy flavor constraints to be easily satisfied and leaves open the possibility of flavor gauge bosons accessible at the LHC. The usual MSSM RPV operators are all forbidden by the new gauge symmetry, but the model allows a purely exotic operator which violates both R -parity and baryon number. Since the exotic fields mix with MSSM-like right handed quarks, diagonalizing the full mass matrix after flavor-breaking transforms this operator into the trilinear baryon number violating operator $\bar{U}\bar{D}\bar{D}$ with flavor coefficients all suppressed by three powers of Yukawa couplings. There is a limit where this model realizes exact MFV; we compute corrections away from MFV, show that they are under theoretical control, and find that the model is viable in large regions of parameter space.

KEYWORDS: Beyond Standard Model, Supersymmetric Standard Model

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1 Introduction

Supersymmetry (SUSY) has long been considered an elegant solution to the hierarchy problem [1]. The LHC is now testing whether weak scale SUSY is found in nature, and as yet has found no evidence for superpartners. Searches for events with multiple jets and large missing energy — the generic signature of R -parity conserving SUSY — are currently in good agreement with Standard Model (SM) expectations up to the \sim TeV scale [2, 3] and place strong constraints on the parameter space that tames the hierarchy problem.

One possible response to this tension is to abandon R -parity, and, with it, the missing energy signatures at the LHC. Unfortunately, allowing R -parity violation (RPV) dramatically worsens the SUSY flavor problem — the parameters in the RPV superpotential all contain flavor indices, and generic values of these parameters are badly ruled out by low energy flavor tests [4]. Furthermore, abandoning R -parity introduces operators that violate baryon and lepton number, inducing rapid proton decay. One solution to this problem is to impose minimal flavor violation (MFV) [5–10] on the MSSM soft Lagrangian and the RPV superpotential [11].¹ MFV is rooted in the observation that the Standard Model (SM) contains a large flavor symmetry $U(3)_Q \times U(3)_U \times U(3)_D \times U(3)_L \times U(3)_E$ which is only broken by the Yukawa couplings. Any hypothetical extension of the SM whose flavor violation arises entirely from Yukawa couplings is said to be MFV. Applying this ansatz to SUSY [6, 7] forces the \mathcal{A} -terms to be aligned with the Yukawa matrices, \mathcal{Y} , and forces the soft masses to be proportional to the unit matrix up to corrections of the form $\mathcal{Y}^\dagger \mathcal{Y}$. Applying MFV to theories without R -parity [11] forces all the RPV couplings to be suppressed by powers of the Yukawa couplings.

In [13] it was shown that requiring the flavor breaking spurions to couple holomorphically, a feature that would be expected in any realistic SUSY model, drastically reduces the number of allowed RPV operators. In the absence of neutrino masses, only a single baryon number violating (BNV) operator is allowed at the renormalizable level:

$$W_{\text{RPV}} = (\mathcal{Y}_U \bar{U})(\mathcal{Y}_D \bar{D})(\mathcal{Y}_D \bar{D}), \quad (1.1)$$

and it requires three Yukawa insertions making even the largest element quite small. Since all other RPV operators are forbidden in the superpotential, the accidental lepton number symmetry of the SM is preserved up to Kahler corrections suppressed by a high power of the UV cutoff, so proton decay is extremely suppressed. Introducing neutrino masses via the seesaw mechanism does not alter the framework’s essential features as all additional lepton violating operators are suppressed by neutrino masses. The MFV structure also allows all flavor constraints to be easily satisfied, and this scenario has interesting phenomenology outlined in [13].

In the SM, MFV is automatically realized since the Yukawa couplings are the only allowed sources of flavor violation. However, in many models of physics beyond the SM, particularly those like the MSSM where many other flavor violating structures are allowed, MFV is merely an intriguing hypothesis. Complete models that implement MFV are

¹Ref. [12] also considers an MFV-like setup where a single lepton number violating RPV operator is the flavor breaking spurion in the lepton sector.

scarce. The most prominent is gauge mediated SUSY breaking [14–18], which is flavor universal at the scale where SUSY breaking is mediated, but RG running induces small MFV deviations away from flavor universality. There are few examples of UV complete models of MFV which deviate from flavor universality and also incorporate RPV.

In this paper, we build a model which has a gauged flavor symmetry

$$G_F = \text{SU}(3)_Q \times \text{SU}(3)_U \times \text{SU}(3)_D \times \text{SU}(3)_L \times \text{SU}(3)_E \times \text{SU}(3)_N \quad (1.2)$$

and approximately reproduces the RPV MFV framework of [13] in the IR. Here, the final factor in the gauge group comes from introducing a triplet of right handed neutrino fields \bar{N} . We introduce the minimal field content required to cancel the anomalies of the new gauge symmetry giving us a model which is a supersymmetrized version of [19]. Similar models have been presented in the context of Grand Unified Theories [20], and Left-Right models [21], and a supersymmetrized Left-Right model was briefly considered in [22]. This theory contains “flavon” superfields² Y which break the flavor symmetries, but, unlike many other models which Higgs a gauged flavor group [23–32], the Yukawa couplings are inversely proportional to the vacuum expectation value (VEV) of the flavon fields. Therefore the flavor gauge fields associated with the first two generations will be much heavier than that of the third generation. This is a realization of the seesaw mechanism in the quark sector which was first introduced in [33–35] because the heavier the gauge fields become, the lighter the SM quarks are. This is a valuable feature because the fields which can mediate flavor changing effects in the first two generations are much more constrained, but, in this setup, they will be very heavy. On the other hand, the third generation gauge bosons are poorly constrained and can be lighter, leaving open the possibility of third generation flavor gauge bosons being discovered at the LHC.

In addition to the MSSM field content, our model has flavon fields which break the flavor gauge symmetry, “exotics” which cancel anomalies and couple directly to flavons, and flavor gauge superfields. The model has the following general structure:

- SUSY breaking is communicated to the MSSM and exotic fields at a high scale, M_* .
- A gauged flavor group is broken by VEVs of flavons, $\langle Y \rangle$, at a scale below M_* but above the electroweak and soft scales.
- The flavon VEVs give mass to the exotic fields at tree level, but they do not directly couple to MSSM-like fields.
- Mass mixing between exotic and MSSM fields is diagonalized to generate Yukawa couplings with inverse dependence on the flavon VEV.
- The separation between the flavor and soft scales ensures approximate mass degeneracy among exotics and their superpartners.

²The flavon fields are denoted Y , while the SM Yukawa matrices are denoted \mathcal{Y} , and, as we note here, $\mathcal{Y} \propto Y^{-1}$.

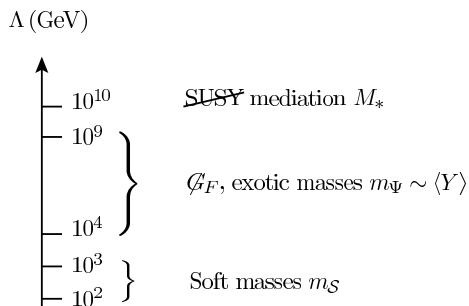


Figure 1. The hierarchies of scales in our model. The numerical values given are approximate and can be raised as long as their order remains the same.

- D -term masses from the flavon VEVs are small compared the hidden sector SUSY breaking contributions to the soft masses.
- Exotic interaction eigenstates allow baryon violation while the MSSM-like $\bar{U}\bar{D}\bar{D}$ is forbidden by flavor gauge invariance; the MSSM RPV operator arises when the mass matrix is diagonalized.

An approximate hierarchy of the scales in this model is given in figure 1, although the numerical values of the scales are all lower bounds. We note that many of the details of the model presented here are not necessary to achieve approximate MFV structure as long as the model has the above features. Particular details of our model such as the choice of the flavor group G_F are taken for concreteness, but are not strictly necessary.

In the limit of $M_* \gg \langle Y \rangle \gg m_S, v$, MFV is exact, but there are corrections which scale like the ratios of the scales. These corrections come from many different sectors of the model, but most of them are parametrically proportional to the SM Yukawa couplings. Namely, SUSY breaking parameters associated with the first and second generation are much smaller than those associated with the third, but the matrix structure of these parameters is such that they are *not* aligned with the SM Yukawa couplings. Therefore, at leading order, the model presented here is MFV, but there are corrections which realize the “flavorful SUSY” [36] paradigm.³ Finally, there are corrections coming from D -terms of the flavor gauge group which are anarchic in flavor space, but all corrections are under theoretical control and can be sufficiently small to evade all constraints.

The outline of this paper is as follows. In section 2 we describe the quark sector of our model in the SUSY preserving limit. In section 3 we describe how SUSY breaking is communicated and show that flavor violation in the soft Lagrangian inherits approximate MFV structure. We also describe sources of deviation from MFV. In section 4 we consider experimental constraints from direct production and various flavor violating processes, and we show that this model is viable for natural values of the flavor parameters. In section 5 we extend the model to include leptons and we show that MFV leads to a natural realization

³RPV in flavorful SUSY was previously considered in [37] in the context of partial compositeness within SUSY.

	SU(3) _Q	SU(3) _U	SU(3) _D	SU(3) _c	SU(2) _L	U(1) _Y
Q	3	1	1	3	2	+1/6
\bar{u}	1	3	1	$\bar{\mathbf{3}}$	1	-2/3
\bar{d}	1	1	3	$\bar{\mathbf{3}}$	1	+1/3
ψ_{uc}	$\bar{\mathbf{3}}$	1	1	$\bar{\mathbf{3}}$	1	-2/3
ψ_{dc}	$\bar{\mathbf{3}}$	1	1	$\bar{\mathbf{3}}$	1	+1/3
ψ_u	1	$\bar{\mathbf{3}}$	1	3	1	+2/3
ψ_d	1	1	$\bar{\mathbf{3}}$	3	1	-1/3
Y_u	3	3	1	1	1	0
Y_u^c	$\bar{\mathbf{3}}$	$\bar{\mathbf{3}}$	1	1	1	0
Y_d	3	1	3	1	1	0
Y_d^c	$\bar{\mathbf{3}}$	1	$\bar{\mathbf{3}}$	1	1	0

Table 1. The field content and charge assignments of the quark sector. The top section is MSSM-like fields, the middle section is exotics, and the bottom section is flavons.

of pure Dirac neutrino masses. We also briefly consider the more standard seesaw scenario. We give concluding remarks in section 6.

2 Gauged supersymmetric model

We begin by describing the quark sector of the model which has a gauged flavor group $G_F \equiv \text{SU}(3)_Q \times \text{SU}(3)_U \times \text{SU}(3)_D$. The full gauge group is $G \equiv G_{\text{SM}} \times G_F$, with the usual $G_{\text{SM}} \equiv \text{SU}(3)_c \times \text{SU}(2)_L \times \text{U}(1)_Y$. In section 5 we will extend the model to include leptons, so the flavor group will be enlarged to contain SU(3) factors for the corresponding L, E and \bar{N} fields. The superfield content of the quark sector is given in table 1: in addition to the MSSM-like fields, Q, \bar{u}, \bar{d} , we add exotic fields $\psi_u, \psi_{uc}, \psi_d, \psi_{dc}$ which are also charged under the flavor group as well as SM color and hypercharge; there are no new fields with SU(2)_L charge. The $\psi_u (\psi_d)$ and $\bar{u} (\bar{d})$ fields now form a vectorlike pair under the full gauge group, so exotic-SM mass-mixing will be present.

Finally, the model also features the flavon fields Y_u and Y_u^c which break $\text{SU}(3)_Q \times \text{SU}(3)_U$, and Y_d and Y_d^c which break $\text{SU}(3)_Q \times \text{SU}(3)_D$. We now have a supersymmetrized setup of the model presented in [19], but here we are forced to add the Y_u^c and Y_d^c superfields to cancel the anomalies that come from the fermions in Y_u and Y_d . We will discuss the origin of flavor breaking in section 2.3, but we note here that, for the separation of scales we will require, the $Y_{u,d}$ scalar VEVs will be approximately aligned with those of their corresponding $Y_{u,d}^c$ conjugates.

The most general, renormalizable superpotential invariant under $G_{\text{SM}} \times G_F$ is

$$W = \lambda_u H_u Q \psi_{uc} + \lambda'_u Y_u \psi_u \psi_{uc} + M_u \psi_u \bar{u} + W_Y + (u \rightarrow d) + W_{\text{BNV}}, \quad (2.1)$$

where the flavon self interactions and baryon violating terms are given by

$$W_Y \equiv \lambda_{Y_u} Y_u Y_u Y_u + \lambda_{Y_u^c} Y_u^c Y_u^c Y_u^c + \mu_{Y_u} Y_u Y_u^c, \quad W_{\text{BNV}} \equiv \frac{1}{2} \eta \psi_{uc} \psi_{dc} \psi_{dc}. \quad (2.2)$$

For the remainder of this paper, whenever down-type interactions have the same structure as their up-type counterparts, they will be omitted for simplicity.

All dimensionless superpotential couplings $(\lambda, \lambda', \lambda_Y, \lambda_{Y^c}, \eta)$ are *flavor universal* coefficients and μ_Y is a free parameter. Note that W_Y contains all allowed $Y_{u,d}, Y_{u,d}^c$ interactions and W_{BNV} is the only source of R -parity violation. The familiar baryon violating trilinear coupling $\bar{u}\bar{d}\bar{d}$ is forbidden by the flavor symmetry, but will be induced upon spontaneous flavor symmetry breaking (see section 2.2).

2.1 Yukawa couplings

Following [19], the $Y_{u,d}$ fields completely break the flavor symmetry prior to electroweak symmetry breaking and give the exotic fields an additional source of mass mixing

$$\lambda'_u \langle Y_u \rangle^{\alpha\beta'} (\psi_{uc})_\alpha (\psi_u)_{\beta'} + M_u (\bar{u})^{\alpha'} (\psi_u)_{\alpha'} + (u \rightarrow d), \tag{2.3}$$

where unprimed and primed greek letters are $SU(3)_Q$ and $SU(3)_U$ indices respectively; in our convention, fundamental indices are raised and anti-fundamental indices are lowered. Following MSSM conventions for Higgs fields, we write the scalar component of $\langle Y_u \rangle$ without a tilde; all other scalars will carry tildes. After flavor breaking, \bar{u} and ψ_{uc} have the same quantum numbers under the remaining unbroken gauge symmetry and can therefore mix with each other. One triplet from the pair matches up with ψ_u and gets a Dirac mass that is $\mathcal{O}(\langle Y \rangle)$. We refer to these exotic mass eigenstates as Ψ . The three remaining fields from \bar{u} and ψ_{uc} are massless before electroweak symmetry and can be identified as the MSSM right handed up superfields $(\bar{U}, \bar{C}, \bar{T})$. More details of the mass diagonalization are given in the appendix.

Rewriting the Higgs coupling in the superpotential in terms of mass eigenstates gives rise to SM Yukawa couplings,

$$\lambda_u H_u Q^\alpha (\psi_{uc})_\alpha \longrightarrow \lambda_u H_u Q^\alpha (\mathcal{V}_{\alpha\beta'}^u \bar{U}^{\beta'}) \equiv (\mathcal{Y}_u)_{\alpha\beta'} H_u Q^\alpha \bar{U}^{\beta'}, \tag{2.4}$$

where $\mathcal{V}^{u,d}$ is the unitary matrix that transforms interaction states into mass eigenstates, and the SM Yukawa matrices are defined by $\mathcal{Y}_u \equiv \lambda_u \mathcal{V}_u$. In the $M/\langle Y \rangle \rightarrow 0$ limit, the \bar{u} are, themselves, the massless eigenstates \bar{U} and do not mix with any exotic fields. Thus, for the lighter generations $\mathcal{Y} \sim \mathcal{O}(M/\langle Y \rangle)$ while for the $\mathcal{O}(1)$ top Yukawa coupling, the relationship is more complicated since M and $\langle Y \rangle$ will be comparable. Note that the change of basis in eq. (2.4) also induces exotic $H_u Q \Psi$ couplings, which do not exhibit the Yukawa structure. This mass diagonalization is shown schematically in figure 2.

Since the Yukawas depend on the ratio $M/\langle Y \rangle$, we demand that $\langle Y \rangle \gg M$ for the lighter generations whose corresponding exotic fermions receive large $\mathcal{O}(\langle Y \rangle)$ masses and become exceedingly heavy; for the third generation, the Yukawa coupling is of order unity, so the corresponding exotic partners are naturally lighter. Note, however, that *the overall scale of M and Y is undetermined*. This feature will become instrumental in suppressing effects that deviate from MFV structure after SUSY is broken in section 3.

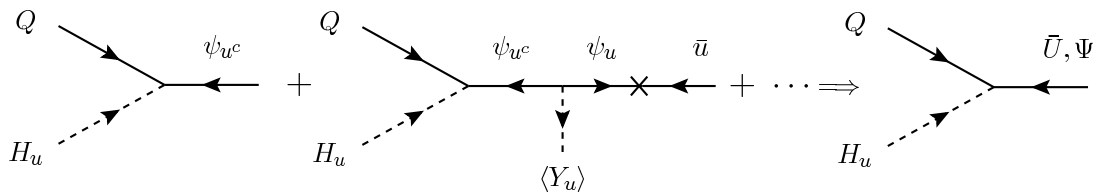


Figure 2. Diagrams that induce the effective Yukawa couplings in the flavor interaction eigenbasis (left) and the mass eigenbasis (right). Upon flavor symmetry breaking, the spurion $Y_u \rightarrow \langle Y_u \rangle$ gives Dirac masses to ψ_u and ψ_{uc} — the former of which also mixes with \bar{u} through the supersymmetric mass term $M_u \psi_u \bar{u}$. Diagonalizing the exotic masses gives rise to the mass-eigenstate \bar{U} which remains massless (prior to EWSB) and is identified with right-handed up-type quark of the MSSM.

2.2 Trilinear baryon violation

The usual BNV operator in the MSSM is

$$\lambda''_{\alpha\beta\gamma} \bar{U}^\alpha \bar{D}^\beta \bar{D}^\gamma, \tag{2.5}$$

where α, β, γ are flavor indices and we have suppressed color indices which are contracted by the epsilon tensor. In this model, that operator is forbidden by the flavor symmetry, and the only source of R -parity violation in is given by

$$W_{\text{BNV}} = \frac{1}{2} \eta \epsilon^{abc} \epsilon^{\alpha\beta\gamma} (\psi_{uc})_{a\alpha} (\psi_{dc})_{b\beta} (\psi_{dc})_{c\gamma}, \tag{2.6}$$

where a, b and c represent $SU(3)_c$ indices. We can rotate from the interaction into the mass basis using the \mathcal{V} matrices and get an operator with characteristic Yukawa suppression

$$W_{\text{BNV}} = \frac{1}{2} \eta \epsilon^{abc} \epsilon^{\alpha\beta\gamma} (\mathcal{V}_{\alpha\alpha'}^u \bar{U}_a^{\alpha'}) (\mathcal{V}_{\beta\beta'}^d \bar{D}_b^{\beta'}) (\mathcal{V}_{\gamma\gamma'}^d \bar{D}_c^{\gamma'}) + \mathcal{O}(\Psi). \tag{2.7}$$

The $\mathcal{O}(\Psi)$ terms are R -parity violating couplings between MSSM fields and at least one exotic. These do not play a large role in TeV scale phenomenology because they have both Yukawa suppression as well as suppression from the heavy exotic.

From eq. (2.7), we identify the BNV matrix between MSSM fields defined in eq. (2.5) as

$$\lambda''_{\alpha'\beta'\gamma'} = \left(\frac{\eta}{\lambda_u \lambda_d^2} \right) \epsilon^{\alpha\beta\gamma} (\mathcal{Y}_{\alpha\alpha'}^u) (\mathcal{Y}_{\beta\beta'}^d) (\mathcal{Y}_{\gamma\gamma'}^d), \tag{2.8}$$

where (α, β, γ) are $SU(3)_Q$ indices and used interchangeably⁴ with $(\alpha', \beta', \gamma')$, and the overall prefactor $\eta/\lambda_u \lambda_d^2$ is of order one. The way the exotic fields feed into the MSSM BNV operator is shown schematically in figure 3. Note that this is precisely the parametric dependence of the BNV operator in [13], which is the only source of R -parity violation in the absence of neutrino masses.

As shown in [13], the MFV structure of λ'' means that all the entries are very small, with the tsb entry being the largest with size $\sim 1 \times 10^{-5} (\tan \beta/10)^2$. Therefore, R -parity is a

⁴Technically, the primed indices that accompany \bar{U} and \bar{D} do not uniquely correspond to $SU(3)_{U,D}$ indices since the flavor group is now broken and these are no longer good quantum numbers.

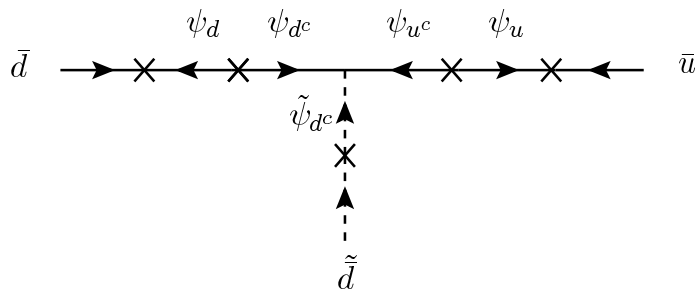


Figure 3. Example diagram that becomes the trilinear BNV operator $\bar{U}\bar{D}\bar{D}$ when the exotic interaction eigenstates are diagonalized. Note that the scalar mass insertion arises from the supersymmetric F -term $\lambda_d^* \langle Y_d^* \rangle M_d \tilde{\psi}_{dc}^* \tilde{d}$ in the scalar potential.

good approximate symmetry, and RPV will only be noticed in very sensitive processes such as the decay of the lightest supersymmetric particle (LSP) and baryon number violation. The constraints and phenomenology of these operators are detailed in [13], and we will review them in section 4.

2.3 Flavor breaking

Like the non-supersymmetric model in [19], our framework does not explain the hierarchies of Yukawa couplings. Unlike [19], this model has the advantage of non-renormalization theorems [38, 39] for superpotential parameters. Therefore, we accept the tuning in the flavor sector of the SM and get the flavor hierarchy dynamically from superpotential interactions. This means we are trading the technically natural Yukawa couplings for technically natural superpotential parameters.

To generate nontrivial flavor textures, we can couple our flavons $Y_{u,d}$ and $Y_{u,d}^c$ to a singlet

$$W \supset \lambda_S S (Y_u Y_u^c - w^2) + \frac{1}{2} M_S S^2 + (u \rightarrow d), \tag{2.9}$$

where w sets the flavor breaking scale and M_S is chosen to ensure that S doesn't get a VEV. The vacuum of this theory is degenerate under all gauge rotations that preserve $\langle Y_u Y_u^c \rangle = w^2$, so minimizing the D -term potential

$$\frac{g_Q^2}{2} |Y_u^* T_Q^a Y_u - Y_u^c T_Q^a Y_u^{c*} + (u \rightarrow d) + \dots|^2 + \text{SU}(3)_Q \rightarrow \text{SU}(3)_{U,D}, \tag{2.10}$$

forces the supersymmetric relation $\langle Y_{u,d} \rangle = \langle Y_{u,d}^c \rangle$ up to small corrections that arise from the trilinear interactions in eq. (2.2); we will return to these corrections in section 3.5.

Although this mechanism will generate a VEV for Y and Y^c , obtaining the observed Yukawa and CKM matrices requires more structure since the potential derived from eq. (2.9) doesn't break the full gauge group. Following [19], a straightforward generalization involves multiple copies of identical Y, Y^c , and S fields

$$W \supset \lambda_{S_i} S^i (C_{ijk} Y_u^j (Y_u^c)^k - w_i^2) + (u \rightarrow d), \tag{2.11}$$

where C_{ijk} is a dimensionless matrix of coefficients whose entries are arbitrary for our present purposes. On general grounds, we are free to use N singlets S_i and perform M

gauge transformations to exhaust the $SU(3)_Q \times SU(3)_U$ flavor symmetry and rotate the Y_u^j and Y_u^{ck} VEVs in flavor space. For suitably chosen $N > M$, we will have used up all our gauge freedom, but will still have leftover minimization conditions for which Y_u^j and Y_u^{ck} can no longer be rotated into a particular direction of flavor space. Thus we can, in principle, break the full flavor group and give the flavons realistic VEVs that yield the observed Yukawa and CKM matrices.

Having multiple flavon fields does not change the essential structure of the model because Y_u in eq. (2.1) is just replaced with $c_i Y_u^i$, which becomes the effective flavon. Although the interaction superpotential in eq. (2.2) becomes more complicated with more copies of Y_u and Y_u^c — couplings acquire new indices: $\lambda_Y \rightarrow \lambda_{Y_{ijk}}$, $\mu_Y \rightarrow \mu_{Y_{ij}}$ — the essential requirements remain qualitatively similar. For the remainder of this paper, when we consider flavon self interactions we will ignore these details without loss of essential generality; it should be understood that any requirements we impose on these parameters in later sections should apply to the effective couplings that take into account contributions from all index combinations.

In this paper, we do not exhaust all possible flavor-breaking mechanisms. This discussion serves merely to demonstrate that the VEVs required for our scenario can, in principle, be engineered without spoiling any of the model’s essential features; the singlets introduced here play no other role and we will ignore them for the rest of the paper. For a discussion of alternative flavor breaking patterns and mechanisms, see [40–42].

2.4 Running of SM gauge couplings

Since the exotic states carry color and hypercharge, both beta functions are modified in the UV; for sufficiently light exotics, both groups also encounter Landau poles below the Planck scale. However, there is no upper bound on either $\langle Y \rangle$ or the mixing parameters M as long as they are below the scale of SUSY breaking mediation M_* (see section 3). Therefore, the scales $\langle Y \rangle$ and M can be chosen to decouple all exotic fields and avoid Landau poles while maintaining the appropriate ratios to generate the SM Yukawa matrices. Provided that all lower bounds on flavor VEVs are satisfied (see section 4), none of the model’s essential features depend on the absolute scale. The scales do, of course, affect the LHC discovery potential for exotic fermions and flavor gauge bosons.

We compute one loop running of the SM gauge couplings including the new fields given in table 1 as well as the fields in the lepton sector described in section 5 and given in table 2. In the lightest plausible realization (see section 4) with third generation exotic masses at 3 TeV and all other masses scaled up the appropriate inverse Yukawa couplings, the $SU(3)_c$ and $U(1)_Y$ gauge couplings diverge around 10^{18} and 10^{14} GeV, respectively. If naively interpreted as the scale associated with proton decay, the latter could pose marginal problems for our model in the lighter exotic regime, however our flavor symmetry forbids lepton violation (see section 5), so the high scale physics should also contain our flavor group and forbid the most dangerous operators even at this scale.

This framework unfortunately spoils the high accuracy of $SU(5)$ gauge coupling unification in the MSSM. It is straightforward to add more matter charged under $SU(2)_L$ and $SU(3)_c$ so that all three hit Landau poles around the same scale to accommodate the

possibility of unification. This model contains a larger gauge symmetry than the MSSM, so it is possible that there exist more natural embeddings in larger gauge groups which contain the full flavor group. It is also important to note that the computation of Landau poles is specific to the gauge group chosen in this paper; the mechanisms used to realize MFV are more generic and may work with other flavor groups. We leave Grand Unification possibilities as well as alternative choices of flavor gauge group to future work.

3 MFV SUSY breaking

Thus far, our discussion has assumed exact supersymmetry. When all F and D -terms vanish in the vacuum, the exotics and their superpartners have degenerate mass matrices (see the appendix) which are diagonalized by the \mathcal{V} matrices that define SM Yukawa couplings in eq. (2.4). In order for MFV to be preserved, the mass matrices for the exotic scalars and fermions need to remain approximately degenerate after SUSY breaking. This will allow soft terms to be diagonalized by the same matrices that yield the Yukawa couplings. For example, taking an \mathcal{A} -term

$$H_u \tilde{Q} \tilde{\psi}_{u^c} \rightarrow H_u \tilde{Q} \mathcal{V} \tilde{U} \propto \mathcal{Y}_u H_u \tilde{Q} \tilde{U}, \tag{3.1}$$

showing that we get the desired MFV structure. In this section we will describe all the contributions to SUSY breaking and explain how approximate degeneracy among the exotics is preserved.

In order for our model to be phenomenologically viable, the dominant source of SUSY breaking for the MSSM-like fields will be via a hidden sector mediated by higher dimensional operators. There are, of course, many alternative SUSY breaking mechanisms one could consider. A particularly interesting example is “Flavor Mediation” [43] in which the flavor gauge-interactions mediate soft masses with non-trivial flavor structure. However, flavor mediation with an inverted hierarchy of gauge boson masses yields light first and second generation squarks and a heavy third generation, which both spoils MFV and requires more fine tuning to cancel quadratic divergences, so we will not adopt this scenario here.

The dynamics introduced in section 2 induce two additional sources of SUSY breaking:

- Nonzero F_Y -terms that arise from flavon self interactions when Y acquires a VEV and $\lambda_Y, \mu_Y \neq 0$. This contribution will be highly suppressed in the MSSM sector because the Y s couple only to exotics at tree level.
- Nonzero D -terms for the flavor gauge groups. Those terms are schematically proportional to $g_F^2(\langle Y \rangle^2 - \langle Y^c \rangle^2)$ where g_F is the gauge coupling for the flavor group. The non-degeneracy between $\langle Y \rangle$ and $\langle Y^c \rangle$ is induced both by nonzero λ_Y, μ_Y and by the hidden sector SUSY breaking.

Below, we describe the requirements for these two contributions to be small perturbations, so that SUSY is still approximately preserved in the exotic sector and the scalars are roughly degenerate with the fermions.

3.1 Hidden sector SUSY breaking

In order to have sufficiently large SUSY breaking, we take an ansatz of a hidden sector field X with $\langle X \rangle = F\theta^2$, and we also assume that the hidden sector communicates to the visible sector through higher dimensional operators suppressed by a high scale M_* . We require $M_* \gg \langle Y \rangle$ so that the flavor symmetry is imposed on all allowed soft terms. As we will see in section 4.2.4, the bound on the gravitino mass requires M_* to be relatively large, $M_* \gtrsim 10^{10}$ GeV, which is completely consistent with having $M_* \gg \langle Y \rangle$.

Up to irrelevant $\mathcal{O}(1)$ variations in the coefficients, the soft Lagrangian contains the following \mathcal{A} , \mathcal{B} , and soft mass $m_{\mathcal{S}}$ terms for MSSM-charged fields at the mediation scale

$$\begin{aligned} \mathcal{A}_{\mathcal{S}}(H_u \tilde{Q} \tilde{\psi}_{u^c} + Y_u \tilde{\psi}_u \tilde{\psi}_{u^c} + \tilde{\psi}_{u^c} \tilde{\psi}_{d^c} \tilde{\psi}_{d^c} + \text{c.c.}) + \mathcal{B}_{\mathcal{S}}(H_u H_d + \tilde{\psi}_u \tilde{u} + \text{c.c.}) \\ + m_{\mathcal{S}}^2(\tilde{Q}^\dagger \tilde{Q} + \tilde{u}^\dagger \tilde{u} + \tilde{\psi}_u^\dagger \tilde{\psi}_u + \tilde{\psi}_{u^c}^\dagger \tilde{\psi}_{u^c}) + (u \rightarrow d), \end{aligned} \quad (3.2)$$

where, for simplicity and without loss of essential generality we assume that all allowed higher dimensional operators of the same order. This gives $\mathcal{A}_{\mathcal{S}} \simeq \sqrt{\mathcal{B}_{\mathcal{S}}} \simeq m_{\mathcal{S}} \equiv F/M_*$.

All the soft parameters are flavor blind since the SUSY spurion is a G_F singlet, but flavor violation will arise through the dynamics of exotic fields whose masses arise from flavon VEVs. In addition to requiring that the mediation scale is above the flavor scale, we also require that the soft scale is below the flavor scale, $\langle Y \rangle \gg m_{\mathcal{S}}$. This means that in the exotic sector, SUSY breaking masses will be small perturbations on the dominant masses that come from flavor breaking. Therefore, exotic boson diagonalization matrices will be modified by additive amounts proportional to the ratio of SUSY breaking to SUSY preserving contributions to the exotic mass squared matrix: $m_{\mathcal{S}}^2/\langle Y \rangle^2$. Because we know that $\langle Y \rangle \gg m_{\mathcal{S}}$, the deviations from MFV are under theoretical control.

3.2 Flavon SUSY breaking

The exotic scalar mass matrix also receives SUSY breaking contributions from flavon self interaction operators after flavor breaking. In the absence of a symmetry to forbid the YY^c , YYY and $Y^c Y^c Y^c$ terms in the superpotential, the flavon VEVs break SUSY through nonzero F_Y and F_{Y^c} terms. This breaking induces exotic \mathcal{B} -terms of the form

$$\lambda'_u (\mu_Y^* \langle Y_u^c \rangle^* + 3\lambda_Y^* \langle Y_u \rangle^{*2}) \tilde{\psi}_u \tilde{\psi}_{u^c} + (u \rightarrow d) + \text{c.c.}, \quad (3.3)$$

which contribute to exotic scalar (but not fermion) masses and give rise to non-minimal flavor violation. To ensure that the corrections in eq. (3.3) are only perturbations away from mass degeneracy, we will need $\mu_Y \ll \langle Y \rangle$ and $\lambda_Y \ll 1$ when we consider constraints on non-MFV interactions in section 4.2.3.

3.3 Exotic scalar mass diagonalization

As discussed above and in more detail in the appendix, in the $m_{\mathcal{S}}, \lambda_Y, \mu_Y \rightarrow 0$ limit, the matrix which diagonalizes the exotic scalar mass matrix is the same as the one which diagonalizes the fermion exotic mass matrix, \mathcal{V} . The scalar mass matrix is modified when soft-masses are taken into account. After SUSY breaking, the mass eigenbases for exotic

states is no longer identical for fermions and bosons. These effects are proportional to $m_S^2/\langle Y \rangle^2 \ll 1$, which is small by construction since the flavor breaking scale is a free parameter in the model; only the ratio $M/\langle Y \rangle$ is constrained by the known Yukawa textures. For the lighter generations, there is also the hierarchy $\langle Y \rangle \gg M$ which ensures that (up to overall order-one coefficients) the diagonalization matrices have the parametric dependence

$$\mathcal{V}_{\text{fermion}} = \mathcal{V}_{\text{scalar}} \sim \frac{M}{\langle Y \rangle}, \tag{3.4}$$

where $\langle Y \rangle$ is a first or second generation VEV and both matrices are proportional to the Yukawa matrix \mathcal{Y} .

In a more typical regime, the flavon self couplings are nonzero and SUSY is also broken in the hidden sector, so the elements of the exotic scalar mass squared matrix become shifted by characteristic amounts $\langle Y \rangle^2 \rightarrow \langle Y \rangle^2 + m_S^2 + \dots$ where the additional terms all arise from SUSY breaking. This shift is reflected in the bosonic diagonalization matrix by small corrections of the form

$$\begin{aligned} \mathcal{V}_{\text{scalar}} &\sim \frac{M}{[\langle Y \rangle^2 + \mu_Y \langle Y \rangle + \lambda_Y \langle Y \rangle^2 + m_S^2 + \dots]^{1/2}} \\ &\approx \frac{M}{\langle Y \rangle} \left[1 + \mathcal{O}(\lambda_Y) + \mathcal{O}\left(\frac{\mu_Y}{\langle Y \rangle}\right) + \mathcal{O}\left(\frac{m_S^2}{\langle Y \rangle^2}\right) \right] \\ &= \mathcal{V}_{\text{fermion}} + \mathcal{O}(\mathcal{Y}\lambda_Y) + \mathcal{O}\left(\mathcal{Y}\frac{\mu_Y}{\langle Y \rangle}\right) + \mathcal{O}\left(\mathcal{Y}\frac{m_S^2}{\langle Y \rangle^2}\right), \end{aligned} \tag{3.5}$$

where the leading term is consistent with MFV since $\mathcal{V}_{\text{fermion}}$ is proportional to the full Yukawa matrix \mathcal{Y} , but the subleading terms inside the brackets have anarchic flavor structure. Nonetheless these terms always appear in combination with $\mathcal{V}_{\text{fermion}}$ so they are further Yukawa suppressed.

Note that in eq. (3.5) we have omitted order one coefficients and only keep track of the parametric dependence of SUSY breaking terms. Since this expansion is only valid for the first and second generation elements of \mathcal{V} , we have also omitted the indices on the flavon VEVs, which, on account of the inverted hierarchy, are predominantly the largest entries in $\langle Y \rangle$. In the appendix we give a more detailed treatment of this index dependence but the order-of-magnitude discussion in this section is sufficient to understand the leading corrections away from degeneracy. In section 4.2, these additional terms are subject to flavor constraints from their contributions to MSSM \mathcal{A} -terms and soft masses; we will see that our rough parameterization here will suffice to set order of magnitude bounds on SUSY breaking corrections.

3.4 Squark \mathcal{A} -terms

The \mathcal{A} -terms in the soft Lagrangian of eq. (3.2) can be transformed into the mass eigenbasis upon flavor symmetry-breaking giving the usual MSSM \mathcal{A} -terms,

$$\mathcal{L}_S \supset \mathcal{A}_S H_u \tilde{Q}^\alpha (\tilde{\psi}_{uc})_\alpha \longrightarrow \mathcal{A}_S H_u \tilde{Q}^\alpha (\mathcal{V}_{\alpha\beta'}^u \tilde{U}^{\beta'}) = \frac{\mathcal{A}_S}{\lambda_u} (\mathcal{Y}_u)_{\alpha\beta'} H_u \tilde{Q}^\alpha \tilde{U}^{\beta'}, \tag{3.6}$$

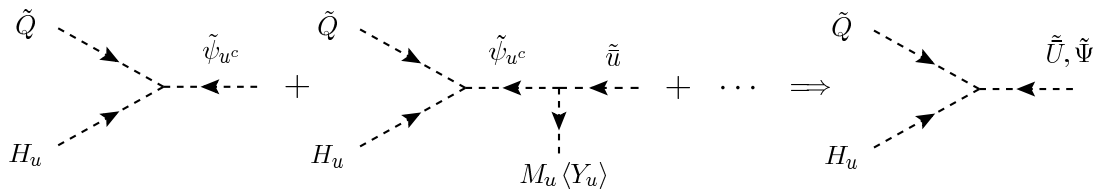


Figure 4. Interaction basis diagrams that yield \mathcal{A} -terms (left and middle) and mass basis diagrams that connect MSSM fields to the mass eigenstates \tilde{U} and $\tilde{\Psi}$ (right). The resulting $H_u \tilde{Q} \tilde{U}$ \mathcal{A} -term is consistent with MFV up to small corrections that arise from exotic scalar/fermion mass differences once SUSY is broken.

which are proportional to Yukawa *matrices* and preserve the features of MFV SUSY. This is shown schematically in figure 4, and we see that those diagrams have the same flavor violating couplings as those in figure 2. Since the mass-eigenvalues/eigenstates for exotic scalars and fermions are identical in the $\langle Y_{u,d} \rangle \gg m_S, \mathcal{A}_S, v$ limit (see the appendix), the \mathcal{A} -terms are exactly aligned with the Yukawa matrices from eq. (2.4). In this limit the scalars get nearly all their mass from SUSY preserving effects and the MFV structure arises in diagrams where the $\mathcal{A}_S H_u \tilde{Q} \tilde{\psi}_{uc}$ supply the only SUSY breaking vertices; the rest of the diagram involves supersymmetric $\tilde{\psi}_u, \tilde{\psi}_{uc}$ and \tilde{u} mass-mixing insertions.

The differences between figures 2 and 4 is that instead of exchanging exotic fermions, \mathcal{A} -terms are generated by exotic scalar exchange. So away from the pure MFV limit of $\langle Y_{u,d} \rangle \gg m_S, \mathcal{A}_S, v$, we must include the corrections in eq. (3.5) to the scalar mass matrix. Setting $\lambda_u = 1$ for simplicity, the full structure of a typical MSSM \mathcal{A} -term is now

$$\mathcal{A}_S \left[\mathcal{Y} + \mathcal{O}(\mathcal{Y} \lambda_Y) + \mathcal{O} \left(\mathcal{Y} \frac{\mu_Y}{\langle Y \rangle} \right) + \mathcal{O} \left(\mathcal{Y} \frac{m_S^2}{\langle Y \rangle^2} \right) \right] H_u \tilde{Q} \tilde{U}, \quad (3.7)$$

where only the first term $\propto \mathcal{Y}$ in the brackets is exactly MFV. To preserve the features of MFV SUSY, we will demand that these additional terms be small and in section 4 we will consider the relevant experimental bounds.

In this framework, as in the MSSM, only the third generation squarks can receive large \mathcal{A} -terms; all other such terms will be Yukawa suppressed. Thus, it is easy to accommodate light stop or sbottom LSPs with large mixing between the left and right handed states. Other sparticles will have very small mixing because the \mathcal{A} -terms are small.

3.5 Squark soft masses

In the interaction eigenbasis, the soft Lagrangian in eq. (3.2) contains the soft masses

$$\mathcal{L}_S \supset m_S^2 (\tilde{Q}^\dagger \tilde{Q} + \tilde{u}^\dagger \tilde{u} + \tilde{\psi}_u^\dagger \tilde{\psi}_u + \tilde{\psi}_{uc}^\dagger \tilde{\psi}_{uc}). \quad (3.8)$$

After diagonalizing the exotic states with unitary matrices, these terms are flavor universal at leading order

$$m_S^2 (\tilde{Q}^\dagger \tilde{Q} + \tilde{U}^\dagger \tilde{U} + \tilde{\Psi}^\dagger \tilde{\Psi}), \quad (3.9)$$

but receive corrections from several sources.

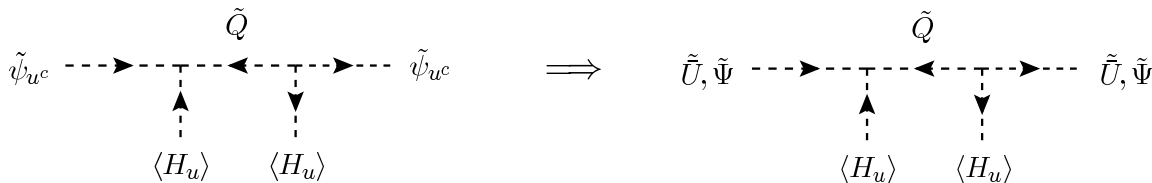


Figure 5. Interaction basis (left) correction to $m_S^2 \tilde{U}^\dagger \tilde{U}$ and $m_S^2 \tilde{\Psi}^\dagger \tilde{\Psi}$ operators; mass basis representation of the same diagram (right). Since the only SUSY breaking interactions in the diagram with $\tilde{U}^\dagger \tilde{U}$ external legs arise from $\mathcal{A}_S H_u \tilde{Q} \tilde{\psi}_{uc}$ vertices (in the interaction basis), the flavor violation in these processes will be proportional to Yukawa matrices and consistent with MFV.

The leading corrections are shown in figure 5 and obey the MFV structure; the only SUSY breaking vertices are due to the $\mathcal{A}_S H_u \tilde{Q} \tilde{\psi}_{uc}$ operator which carries Yukawa structure and is brought into the mass eigenbasis by the same matrices that define the Yukawa couplings. This correction to the soft mass exists in the MSSM as long as \mathcal{A} -terms are non-zero and it has the same structure. In both cases it is suppressed by v^2/m_S^2 because it requires two Higgs insertions and the exchange of a heavy scalar.

Another correction to the scalar soft masses comes from the corrections to the matrix that diagonalizes the exotic scalars $\mathcal{V}_{\text{scalar}} = \mathcal{V}_{\text{fermion}} + \epsilon$, where ϵ is given in eq. (3.5) and $\mathcal{V}_{\text{fermion}} \propto \mathcal{Y}$. This correction depends linearly on the Yukawa couplings and arises because diagrams that induce corrections to chirality-preserving soft-masses require at least two insertions of interaction-eigenstate vertices (see figure 5). Thus, transforming to the mass basis involves two diagonalization matrices in the combination

$$(\mathcal{V}_{\text{scalar}})^\dagger (\mathcal{V}_{\text{scalar}}) = (\mathcal{V}_{\text{fermion}} + \epsilon)^\dagger (\mathcal{V}_{\text{fermion}} + \epsilon) = \mathcal{Y}^\dagger \mathcal{Y} + \mathcal{Y}^\dagger \epsilon + \epsilon^\dagger \mathcal{Y} + \mathcal{O}(\epsilon^2), \quad (3.10)$$

and the cross terms give us the $\mathcal{O}(\mathcal{Y})$ corrections.

The scalar soft masses are also corrected by non-zero D -term VEVs for the flavor gauge bosons. In section 2.3 we saw that minimizing the D -term potential forced $\langle Y \rangle = \langle Y^c \rangle$ in the limit where flavon self-interactions are absent and SUSY is preserved. More generically, nonzero λ_Y , μ_Y and m_S induce small splittings, so the VEVs are no longer aligned and the model acquires nonzero D -terms

$$\frac{g_F^2}{2} |\tilde{Q}^* T_Q^a \tilde{Q} - \tilde{\psi}_{uc} T_Q^a \tilde{\psi}_{uc}^* + Y_u^* T_Q^a Y_u - Y_u^c T_Q^a Y_u^{c*} + (u \rightarrow d)|^2 + \text{SU}(3)_Q \rightarrow \text{SU}(3)_{U,D}. \quad (3.11)$$

There are now cross-terms that yield squark masses squared of order $\sim g_F^2 (\langle Y_u^* \rangle^2 - \langle Y_u^{c*} \rangle^2)$. To estimate the effect of flavon self interactions, we note that (for $m_S = 0$) the flavon VEVs depend inversely on the quartic couplings

$$\langle Y \rangle \simeq \frac{\lambda_S w}{(\lambda_S^2 + \lambda_Y^2)^{1/2}} \approx w \left(1 - \frac{\lambda_Y^2}{2 \lambda_S^2} \right), \quad (3.12)$$

where λ_S is the Yukawa coupling to the singlet defined in eq. (2.9), and λ_Y is the flavon trilinear coupling defined in eq. (2.2). We note that λ_S is identical for Y and Y^c potentials, but λ_Y and λ_{Y^c} typically differ by small amounts, so we have done an expansion in small

λ_Y . Thus, the D -term squark mass squared is of order

$$g_F^2(\langle Y_u^* \rangle^2 - \langle Y_u^{c*} \rangle^2) \sim g_F^2 \frac{\lambda_Y^4}{2\lambda_S^4} \langle Y \rangle^2. \tag{3.13}$$

When hidden sector soft masses are included in the flavon potential, they also upset the $\langle Y \rangle, \langle Y^c \rangle$ cancellation by shifting the VEVs by small, uncorrelated amounts of order the soft-mass scale. The difference in eq. (3.13) now becomes

$$g_F^2(\langle Y_u^* \rangle^2 - \langle Y_u^{c*} \rangle^2) \rightarrow g_F^2 \left(m_S^2 + \frac{\lambda_Y^4}{2\lambda_S^4} \langle Y \rangle^2 \right), \tag{3.14}$$

where we have taken the typical flavon soft mass to be of order m_S .

Summarizing, the full squark soft mass squared for \tilde{Q}, \tilde{U} and \tilde{D} has the form

$$m_S^2 \left\{ 1 + \frac{v^2}{m_S^2} \mathcal{Y}^\dagger \mathcal{Y} + \mathcal{Y} \left[\mathcal{O}(\mathcal{Y} \lambda_Y) + \mathcal{O} \left(\mathcal{Y} \frac{\mu_Y}{\langle Y \rangle} \right) + \mathcal{O} \left(\mathcal{Y} \frac{m_S^2}{\langle Y \rangle^2} \right) \right] + g_F^2 \left[\mathcal{O}(1) + \mathcal{O} \left(\frac{\lambda_Y^4 \langle Y \rangle^2}{\lambda_S^4 m_S^2} \right) \right] \right\}, \tag{3.15}$$

where only the first two terms are properly MFV, and we will estimate the constraints on the remaining terms in section 4. Again, for our parametric estimates, we have set $\mathcal{A}_S = m_S$ for simplicity. Note that for the $\tilde{Q}^* \tilde{Q}$ soft mass, there are also corrections of order $v^2 / \langle Y \rangle^2$ from the tree level exchange of Ψ with two Higgs insertions, but these are always suppressed relative to the $m_S^2 / \langle Y \rangle^2$ terms in eq. (3.15) so we will not consider them further.

4 Experimental constraints

In this section we consider the experimental bounds on this model coming from both direct and indirect constraints. Many of the constraints, especially the indirect ones, are unchanged from the analyses in [13] and [19] and apply similarly to the model presented here. A more complete analysis of the low energy flavor parameters arising from the flavor gauge bosons and the exotics is given in [44].

4.1 Direct constraints

4.1.1 RPV limits

Recent LHC searches have dramatically reduced the parameter space for natural RPV spectra with sparticle masses below the \sim TeV range. However, most of this sensitivity is driven by signals from more dangerous lepton violating operators $QL\bar{D}, LL\bar{E}$, and LH_u . Scenarios with light, “natural” stops are still viable for masses above ~ 100 GeV provided they decay primarily to dijets through the trilinear RPV operator $\bar{U}\bar{D}\bar{D}$ [45]. LHC trijet searches have also recently excluded gluinos below ~ 650 GeV [46] by simple jet-counting searches, but the general features of our scenario are experimentally viable for natural, sub TeV choices of sparticle masses.

4.1.2 Flavor gauge bosons

The inverted hierarchy mechanism ensures that most of the gauge bosons are extremely heavy; their masses are proportional to inverse Yukawa couplings. Gauge bosons that mediate third generation interactions can be considerably lighter, but are constrained by limits on dijet production. Adding SUSY to the model in [19] does not affect the relevant signals, so the same considerations apply in our case, but the experimental limits have improved since 2010. Current LHC limits [47, 48] exclude Z' and W' states which decay to dijets below 1.5 and 2.1 TeV, respectively assuming SM gauge couplings. In section 4.2.3 we show that constraints on D -term flavor violation conservatively require the gauge couplings to be $g \lesssim 0.03$, so naively the bound is weaker in our case, but each of our gauge groups can supply a light gauge boson, so the effective gauge coupling is roughly of standard model size. This sets the strongest bound on the lightest gauge boson masses.

Supersymmetrization means that there will also be gauginos. These are approximately degenerate with the gauge bosons because $m_S/\langle Y \rangle \ll 1$. Since R -parity is approximately conserved up to small baryon number violation, gauginos must be pair produced at the LHC and will be much more difficult to observe than the gauge bosons.

4.1.3 Exotic Ψ states

The lightest exotic fermions will most often decay to Wq , Zq , or hq , however, the flavor violating gauge boson couplings are mixing-angle suppressed, so as long as the hq mode is kinematically accessible, it will be the dominant decay channel. Signals from $\Psi \rightarrow hq \rightarrow b\bar{b}q$ decays are observable in trijet resonance searches [46], which set lower bounds on RPV gluino masses of order ~ 650 GeV, however, the total rate for triplet fermions is reduced by a smaller color factor. Although a dedicated search is necessary to set limits on this scenario, we expect the current bound to be of order a few hundred GeV, which is automatically satisfied if we demand $\lambda'_{u,d} \sim \mathcal{O}(1)$ from eq. (2.1) and $\langle Y_{u,d} \rangle \gtrsim \mathcal{O}(10 \text{ TeV})$ to be large enough for the flavor gauge bosons to satisfy direct production bounds. For $\lambda'_{u,d} \sim 10^{-3}$ – 10^{-2} , these states may be observable at the LHC and can give rise to striking $pp \rightarrow \Psi\Psi \rightarrow (hb)(h\bar{b}), (ht)(h\bar{t}) \rightarrow (bb\bar{b})(b\bar{b}\bar{b}), (b\bar{b}t)(b\bar{b}\bar{t})$ signatures. The scalar partners $\tilde{\Psi}$ will still be roughly degenerate, but will be more difficult to produce because scalar cross sections are smaller than fermions. The scalars will decay down to MSSM superpartners and eventually the LSP which will decay via the RPV operator.

4.2 Indirect constraints

4.2.1 Gauge boson exchange

Unlike MFV proper, the gauged flavor model allows tree-level FCNCs through flavor gauge-boson exchange even in the absence of Yukawa couplings. In the $M_{u,d} \rightarrow 0$ limit, the \bar{u} fields no longer mix with the exotic fermions through the $M_u\psi_u\bar{u}$ coupling and every element of $\mathcal{V}_{\alpha\beta'}^{u,d}$ in eq. (2.4) becomes zero, but the four-Fermi operators that induce FCNCs persist:

$$\frac{g_F^2}{m_V^2} (\bar{f}\gamma^\mu f)(\bar{f}\gamma_\mu f) \sim \frac{1}{\langle Y \rangle^2} (\bar{f}\gamma^\mu f)(\bar{f}\gamma_\mu f), \quad (4.1)$$

where the gauge boson masses are given by $m_V \sim g_F \langle Y \rangle$, g_F is a typical flavor gauge coupling, and f is a SM fermion field whose flavor indices have been omitted. Since our model inherits the same gauge interactions as [19], the experimental constraints on these operators are identical.

Note that the gauge coupling cancels in eq. (4.1) and all the suppression comes from the scale of the appropriate flavon VEV. For transitions involving light generation quarks, the strongest bounds require $\langle Y \rangle_{11,12} \gtrsim 10^3\text{--}10^4$ TeV [49]. The corresponding bounds on third generation VEVs are much weaker, but these will be strongly constrained in section 4.2.3 where limits on g_F are combined with the direct production bound on the lightest gauge boson $m_V \sim g_F \langle Y \rangle_{33}$ from section 4.1.2.

4.2.2 Mixing of exotic fermions

Since the exotic states ψ mix with MSSM quark fields, there are new currents that couple to SM gauge bosons and modify the theoretical prediction of $R_b \equiv \Gamma(Z \rightarrow b\bar{b})/\Gamma(Z \rightarrow \text{all})$, the branching ratio of the Z to b -quarks. Since our modifications to the model do not affect these bounds relative to the non-supersymmetric case, we can read off the relevant constraint from figure 1 of [19]. For $\mathcal{O}(1)$ values of λ_d and $M_d \gtrsim \text{few } 100 \text{ GeV}$ the bounds are easily satisfied when the lightest down-type VEV is $\langle Y_d \rangle \gtrsim \text{TeV}$, which is automatic given the constraints on $\langle Y_d \rangle$ from Z' and W' searches — see section 4.1.2.

The largest mixing is among the top-like quarks and exotics, and there is an additional constraint from deviations from the CKM matrix element V_{tb} . The effective CKM matrix depends on the mixing angles that relate the mass/interaction eigenstates. Unitarity of the CKM is, therefore, only valid up to corrections that depend on the masses of the heavy, new particles. Like the R_b constraints, this observable is unaffected by supersymmetrization, so the bounds are identical to those summarized in figure 2 of [19] which shows a large, viable region of parameter space in the λ_u, M_u plane. In particular, for $\mathcal{O}(1)$ choices of λ_u and $M_u \gtrsim \text{few } 100 \text{ GeV}$, the bounds are easily satisfied as long as the smallest VEVs $\langle Y_{u,d} \rangle$ are in our range of interest $\gtrsim \text{few TeV}$.

4.2.3 Flavor violation in the MSSM sector

Since the exotic fermions and bosons have non-degenerate masses due to SUSY breaking and flavon self interactions, this splitting gives rise to soft-mass corrections that deviate from MFV structure; in full generality, the model realizes the “flavorful supersymmetry” scenario [36]. However in the weakly gauged, high flavor-scale limit, the deviations away from MFV are under theoretical control.

In order to estimate the bounds on low energy observables we use the mass insertion method [50], where we define δ as the ratio of the off diagonal to the diagonal entries in the scalar mass squared matrix in the basis where the fermion masses are diagonal. For pure MFV, $\delta_{\alpha\beta} = 0$ for $\alpha \neq \beta$. We then use the bounds on the mass insertion parameters from [51, 52] compiled in [36, 53].

The most stringent constraints in flavorful supersymmetry come from chirality changing operators such as \mathcal{A} -terms, shown in figure 4. We can parametrize the non-MFV

contributions to \mathcal{A} -terms from eq. (3.7) by defining

$$\Delta_{\mathcal{A}} \equiv \mathcal{A}_{\mathcal{S}} \left[\mathcal{O}(\mathcal{Y}\lambda_Y) + \mathcal{O}\left(\mathcal{Y}\frac{\mu_Y}{\langle Y \rangle}\right) + \mathcal{O}\left(\mathcal{Y}\frac{m_{\mathcal{S}}^2}{\langle Y \rangle^2}\right) \right], \quad (4.2)$$

where we take the overall coefficient $\mathcal{A}_{\mathcal{S}}$ be of order the soft scale $m_{\mathcal{S}}$. The $\mathcal{O}(1)$ numbers in front of these contributions are irrelevant for setting order-of-magnitude bounds, so we ignore them for simplicity and clarity.

From [36], the strongest bounds on the *non*-MFV contributions to \mathcal{A} -terms come from various EDM measurements. The left/right mass insertion parameter in the up sector is given by

$$(\delta_{\alpha\beta}^u)_{LR} \equiv \frac{v(\Delta_{\mathcal{A}})_{\alpha\beta}}{m_{\mathcal{S}}^2} = \frac{(\Delta_{\mathcal{A}})_{\alpha\beta}(m_u)_{\alpha\beta}}{m_{\mathcal{S}}^2 (\mathcal{Y}_u)_{\alpha\beta}}, \quad (4.3)$$

where, v is SM Higgs VEV, $(m_u)_{\alpha\beta} \equiv \sqrt{(m_u)_{\alpha}(m_u)_{\beta}}$ comes from the fermion mass matrix, $(\Delta_{\mathcal{A}})_{\alpha\beta}$ is the non-MFV correction to the $\alpha\beta$ \mathcal{A} -term from eq. (4.2), and $m_{\mathcal{S}}$ is the typical size of a squark mass. Note that the terms in $\Delta_{\mathcal{A}}$ are linear in Yukawa-couplings (but not the full matrices), so there is a cancellation between this dependence and the \mathcal{Y} in the denominator of eq. (4.3) when δ_{LR}^u is written in terms of the quark mass matrix. Using the parametric dependencies of $\Delta_{\mathcal{A}}$, we can compare our order-of-magnitude predictions for δ_{LR}^u against the strongest representative bound on Higgs-squark \mathcal{A} -terms coming from the non-observation of a neutron EDM:

$$|\text{Im}(\delta_{11}^u)_{LR}| \lesssim 4 \times 10^{-7} \left(\frac{m_{\mathcal{S}}}{600 \text{ GeV}} \right), \quad (4.4)$$

which translates into bounds on the first generation flavon VEVs, $\langle Y \rangle_{11}$, as well as on λ_Y , and $\mu_Y/\langle Y \rangle_{11}$. The index structure of $\langle Y \rangle$ that appears in eq. (4.2) is discussed further in the appendix.

Assuming $\mathcal{O}(1)$ complex phases and using $(m_u)_{11} \approx 3 \text{ MeV}$, we find that the $\mathcal{O}(\mathcal{Y}m_{\mathcal{S}}^2/\langle Y \rangle^2)$ correction from $\Delta_{\mathcal{A}}$ yields the constraint $\langle Y \rangle_{11} \gtrsim 2 \text{ TeV}$, which is irrelevant since the scale of flavor VEVs associated with light generations must already exceed $\sim 10^3\text{--}10^4 \text{ TeV}$ from limits on four-Fermi operators in section 4.2.1. The $\mathcal{O}(\mathcal{Y}\lambda_Y)$ and $\mathcal{O}(\mathcal{Y}\mu_Y/\langle Y \rangle)$ in $\Delta_{\mathcal{A}}$ are also constrained by the bounds on $(\delta_{11}^u)_{LR}$. Setting $m_{\mathcal{S}} = 600 \text{ GeV}$, we require λ_Y and $\mu_Y/\langle Y \rangle_{11} \lesssim 10^{-3}$.

We next compute the bounds from chirality conserving soft masses by defining

$$(\delta_{\alpha\beta}^q)_{LL,RR} \equiv \frac{(\Delta m_{\tilde{q}}^2)_{\alpha\beta}}{m_{\mathcal{S}}^2}, \quad (4.5)$$

where $\Delta m_{\tilde{q}}^2$ is any flavor violating (non MFV) squark mass associated with any of the $\tilde{Q}^*\tilde{Q}$, $\tilde{U}^*\tilde{U}$, or $\tilde{D}^*\tilde{D}$ operators. From eq. (3.15) we can read off the parametric dependence of our chirality conserving corrections

$$\Delta m_{\tilde{q}}^2 = m_{\mathcal{S}}^2 \left\{ \mathcal{Y} \left[\mathcal{O}(\mathcal{Y}\lambda_Y) + \mathcal{O}\left(\mathcal{Y}\frac{\mu_Y}{\langle Y \rangle}\right) + \mathcal{O}\left(\mathcal{Y}\frac{m_{\mathcal{S}}^2}{\langle Y \rangle^2}\right) \right] + g_F^2 \left[\mathcal{O}(1) + \mathcal{O}\left(\frac{\lambda_Y^4}{\lambda_S^4} \frac{\langle Y \rangle^2}{m_{\mathcal{S}}^2}\right) \right] \right\}, \quad (4.6)$$

where the flavor indices are suppressed. Although there are many terms to consider, chirality-violating (left-right) constraints set much stronger bounds on λ_Y and $\mu_Y/\langle Y \rangle$,

so we need not evaluate them again. Furthermore, the overall scale of $\langle Y \rangle$ already has to be large to satisfy FCNC constraints from gauge-boson exchange, so the only new constraints from chirality preserving interactions applies to the $\propto g_F^2$ terms in eq. (4.6).

Unlike the left/right transitions, flavor violation arising from left/left and right/right processes is far less constrained. The strongest bounds on this ratio arise from $K\bar{K}$ mixing which requires [36]

$$|(\delta_{12}^u)_{LL}| = |(\delta_{12}^u)_{RR}| \lesssim 10^{-3} - 10^{-2}, \tag{4.7}$$

$$|(\delta_{12}^d)_{LL}| = |(\delta_{12}^d)_{RR}| \lesssim 10^{-2} - 10^{-1}, \tag{4.8}$$

$$|(\delta_{12}^d)_{LL}(\delta_{12}^d)_{RR}| \lesssim 10^{-3}. \tag{4.9}$$

Translating these into bounds on our model, we find

$$|(\delta_{12}^{u,d})_{LL}| \sim |(\delta_{12}^{u,d})_{RR}| \simeq g_F^2 \left(1 + \frac{\lambda_Y^4 \langle Y \rangle^2}{2\lambda_S^4 m_S^2} \right) \lesssim 10^{-3}, \tag{4.10}$$

which, for the first term implies $g_F \lesssim 0.03$ and thereby places the strongest lower bound on the smallest flavor VEV since the lightest gauge boson masses satisfy $m_V \sim g_F \langle Y \rangle_{33} \gtrsim (1-2)$ TeV from the considerations in section 4.1.2. We, thus, require the lightest flavor VEVs to satisfy $\langle Y \rangle_{33} \gtrsim 30$ TeV. To estimate the bound on the second term in eq. (4.10), we set $m_S = 600$ GeV and conservatively choose the largest flavor VEV⁵ to be $\max\{\langle Y \rangle\} = \langle Y \rangle_{11} = 10^9$ GeV, which imposes $\lambda_Y/\lambda_S \lesssim 10^{-4} - 10^{-3}$.

In this section we have extracted bounds on up-type parameters because these are the most tightly constrained in [36, 53], but slightly weaker bounds exist for the down-type sector as well. Since the SUSY breaking corrections considered here are parametrically identical in the down-sector, the constraints are similar.

4.2.4 Baryon number violation

The primary constraints on the baryon number violating operators in eq. (2.7) are the $\Delta B = 2$ processes neutron-antineutron oscillations and dinucleon decay. In [13], it was shown that, under reasonable assumptions about $\tan\beta$ and hadronic matrix elements, holomorphic MFV sufficiently suppresses these processes to allow light (\sim few 100 GeV) scalars at the weak scale. Since the relevant diagrams involve two insertions of flavor-violating $\tilde{b}^* \tilde{d}$ soft masses and two trilinear BNV vertices all exhibiting MFV structure, our non-MFV effects will appear as small corrections to squark soft masses, which shift the predicted oscillation timescale $\tau_{\bar{n}-n}$ by

$$\tau_{\bar{n}-n} \sim (\tau_{\bar{n}-n})_{\text{MFV}} [1 + \mathcal{O}(\delta_{LL/RR})], \tag{4.11}$$

which arises from non-minimal flavor transitions among chirality preserving squark masses. The additional correction term is constrained to be $\lesssim 10^{-3}$ by eqs. (4.7)–(4.9). Similar

⁵As long as the flavor symmetry is intact when hidden sector SUSY breaking is communicated to the MSSM — i.e. $M_* > \max\{\langle Y \rangle\}$ — there is no upper bound on the largest VEV. We arrive at the upper value ($\sim 10^9$ GeV) by saturating the lower bound on the *smallest* VEV $\langle Y \rangle_{33} \gtrsim 30$ TeV and scaling this value by the smallest (inverse) Yukawa coupling.

considerations apply to dinucleon decay, so the non-MFV effects do not upset the relevant constraints on BNV processes; satisfying the bounds described previously in this section is sufficient to ensure that the desirable features of MFV SUSY are preserved.

While lepton number is a very good approximate symmetry and the usual proton decay process cannot occur, proton decay via $p \rightarrow K^+ \tilde{G}$ can be mediated by just baryon number violating operators if the gravitino is sufficiently light. The diagrams for this process require one λ'' insertion and one off-diagonal soft mass insertion. For pure MFV SUSY, this imposes the rough constraint [13]

$$m_{3/2} \gtrsim (300 \text{ keV}) \left(\frac{300 \text{ GeV}}{m_S} \right)^2 \left(\frac{\tan \beta}{10} \right)^4. \quad (4.12)$$

Since $m_{3/2} \sim F/M_{\text{Pl}}$ and we can relate the SUSY breaking scale to the soft-mass and mediation scales $F \sim m_S M_*$, we place a lower bound on the mediation scale

$$M_* \gtrsim (10^{10} \text{ GeV}) \left(\frac{300 \text{ GeV}}{m_S} \right)^3 \left(\frac{\tan \beta}{10} \right)^4. \quad (4.13)$$

Therefore, considering SUSY breaking on the effective theory without flavor breaking is consistent because the mediation scale for SUSY breaking is higher than the scales of flavor breaking.

5 Lepton sector

In the Standard Model, neutrinos are exactly massless, so fitting MFV to the known fact that neutrinos do have mass requires physics beyond the SM. Therefore, there is no unique prescription to apply MFV to the lepton sector [10, 12]. Here, we take the attitude that the neutrino sector is analogous to the quark sector and imagine that the neutrinos get Dirac masses with a triplet of right handed neutrinos denoted \bar{N} , and we extend the model to include the gauged lepton group $SU(3)_L \times SU(3)_E \times SU(3)_N$ for the right and left-handed leptons. The field content in this sector is given in table 2 and gives rise to the same dynamics as the quark sector; Yukawa couplings arise when flavor is broken by VEVs of flavon superfields and heavy exotic fields are integrated out.

5.1 Dirac neutrinos

The simplest lepton sector completely mimics the organization of section 2 where the flavon VEVs are the only sources of flavor violation and the Yukawa couplings arise when $SU(3)_L \times SU(3)_E \times SU(3)_N$ is completely broken. In this sector, we use $Y_{\nu,e}$ to break the flavor group and $Y_{\nu,e}^c$ to cancel anomalies introduced by the flavon fermionic partners. This differs from the usual seesaw scenario because there is no Majorana mass for the right handed neutrinos. While in the standard seesaw scenario the right handed neutrino is a gauge singlet which can get a Majorana mass, in our case the \bar{N} field is charged under a flavor symmetry so it is natural to forbid its mass. In section 5.2 we will add another flavon which allows \bar{N} to get a mass and explore the traditional seesaw scenario.

	$SU(3)_L$	$SU(3)_E$	$SU(3)_N$	$U(1)_Y$
L	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$	$-1/2$
\bar{e}	$\mathbf{1}$	$\mathbf{3}$	$\mathbf{1}$	$+1$
\bar{N}	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{3}$	0
ψ_{ec}	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\mathbf{1}$	$+1$
ψ_N	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\mathbf{1}$	0
ψ_e	$\mathbf{1}$	$\bar{\mathbf{3}}$	$\mathbf{1}$	-1
ψ_ν	$\mathbf{1}$	$\mathbf{1}$	$\bar{\mathbf{3}}$	0
Y_ν	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{3}$	0
Y_ν^c	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\bar{\mathbf{3}}$	0
Y_e	$\mathbf{3}$	$\mathbf{3}$	$\mathbf{1}$	0
Y_e^c	$\bar{\mathbf{3}}$	$\bar{\mathbf{3}}$	$\mathbf{1}$	0

Table 2. The charge assignments for the lepton sector.

For the pure Dirac case, the superpotential in this sector is analogous to eq. (2.1)

$$W_L = \lambda_e H_d L \psi_{ec} + \lambda'_e Y_e \psi_e \psi_{ec} + M_e \psi_e \bar{e} + W_{Y_L} + (d \rightarrow u, e \rightarrow \nu, \bar{e} \rightarrow \bar{N}) \quad (5.1)$$

where the flavon self couplings are

$$Y_L = \lambda_{Y_L} Y_e Y_e Y_e + \lambda_{Y_L^c} Y_e^c Y_e^c Y_e^c + \mu_{Y_L} Y_e Y_e^c + (e \rightarrow \nu). \quad (5.2)$$

Because the Majorana mass term for the right handed neutrinos \bar{N} is forbidden, the neutrinos are purely Dirac particles whose masses are set by Yukawa couplings $\mathcal{Y}_\nu \sim \lambda_\nu M_\nu / \lambda'_\nu \langle Y_\nu \rangle$.

The effective field theory description of MFV SUSY in [13] allows lepton number violation in the Kähler potential even in the absence of neutrino masses. With the appropriate Yukawa matrix insertions, the operator LH_d^* is allowed by the $\mathcal{Z}_3 \subset SU(3)_L \times SU(3)_E$, but only arises at dimension eight and is strongly suppressed by the cutoff scale. Since our model is UV complete, it inherits the accidental lepton symmetry of the SM, so these operators are suppressed by either the planck scale or some potential unification scale. In either case, these are irrelevant from the standpoint of IR flavor constraints.

5.1.1 Indirect constraints on leptons

It is straightforward to generalize the discussion of SUSY breaking from section 3. Just like in the quark sector, the gauge symmetry gives rise to flavor-diagonal soft masses $\sim \mathcal{O}(m_S)$ with soft flavor violation arising primarily from induced Yukawa couplings that appear in soft masses and \mathcal{A} -terms. Additional flavor violation will arise in the soft terms after SUSY breaking from both the hidden sector described in section 3.1 and flavon self-interactions in section 3.2, but the deviations from MFV structure will again be under theoretical control and vanish in certain limits.

As in section 4.2.3, we adopt the mass insertion method to parameterize the size of FCNCs by defining the chirality violating mass insertion parameter

$$(\delta_{\alpha\beta}^e)_{LR} \equiv \frac{(\Delta_{\mathcal{A}})_{\alpha\beta}(m_e)_{\alpha\beta}}{m_{\mathcal{S}}^2 (\mathcal{Y}_e)_{\alpha\beta}}, \quad (5.3)$$

where α, β are generation indices, $\Delta_{\mathcal{A}}$ is the leading *non* MFV contribution to the $\alpha\beta$ \mathcal{A} -term, and $(m_e)_{\alpha\beta} \equiv \sqrt{(m_e)_{\alpha}(m_e)_{\beta}}$ is the charged lepton mass matrix. From [53], the strongest bounds in the lepton sector come from $\mu \rightarrow e\gamma$ and $\mu \rightarrow eee$ rare decays, and $\mu \rightarrow e$ conversion in nuclei, which constrain

$$\frac{1}{\sqrt{2}} \sqrt{|(\delta_{12}^e)_{LR}|^2 + |(\delta_{21}^e)_{LR}|^2} \lesssim 4 \times 10^{-6} \left(\frac{m_{\mathcal{S}}}{200 \text{ GeV}} \right), \quad (5.4)$$

and the electron EDM, which requires

$$|\text{Im}(\delta_{11}^e)_{LR}| \lesssim 2 \times 10^{-7} \left(\frac{m_{\mathcal{S}}}{200 \text{ GeV}} \right). \quad (5.5)$$

The non-MFV corrections to lepton \mathcal{A} -terms are analogous to those from the quark sector; the deviation $\Delta_{\mathcal{A}}$ is parametrically identical to that of eq. (4.2) with quarks exchanged for leptons. Assuming order-one complex phases, we can constrain the lepton VEVs $\langle Y_e \rangle_{12} \gtrsim 600 \text{ GeV}$ and $\langle Y_e \rangle_{11} \gtrsim 700 \text{ GeV}$, which are slightly weaker than the corresponding bound for the quark-flavor scales given in section 4.2.3. For $m_{\mathcal{S}} = 200 \text{ GeV}$, the bounds on flavon self-interaction parameters λ_{Y_e} and $\mu_{Y_e}/\langle Y_e \rangle_{11} \lesssim 10^{-1}$ are also slightly weaker than those in the quark sector.

Turning to chirality-preserving flavor violation, we can define

$$(\delta_{\alpha\beta}^e)_{LL,RR} \equiv \frac{(\Delta m_{\tilde{e}}^2)_{\alpha\beta}}{m_{\mathcal{S}}^2}, \quad (5.6)$$

in analogy with eq. (4.5), where we swap $\Delta m_{\tilde{q}}^2$ with $\Delta m_{\tilde{\ell}}^2$, which represents all non-MFV soft-mass contributions to the $\tilde{L}^* \tilde{L}$ or $\tilde{E}^* \tilde{E}$ operators. For slepton soft masses at $m_{\mathcal{S}} = 200 \text{ GeV}$, bounds on $\mu \rightarrow e\gamma$ require [36]

$$|(\delta_{12}^e)_{LL}| \lesssim 10^{-4} - 10^{-3}. \quad (5.7)$$

Since the $\Delta m_{\tilde{e}}^2$ corrections are identical to those found in eq. (4.6) with lepton-sector parameters (e.g. $\lambda_Y \rightarrow \lambda_{Y_e}$), we find comparable bounds on lepton sector gauge-couplings $g_F \lesssim 10^{-2}$. The flavon self interaction parameters $\lambda_{Y_e}/\lambda_{S_e}$ also contribute corrections of order $g_F^2 (\lambda_{Y_e}^4/\lambda_{S_e}^4) \langle Y_e \rangle^2$. Following our approach in the quark sector, we constrain the latter ratio by choosing the largest flavor VEV $\langle Y_e \rangle \rightarrow \max\{\langle Y_e \rangle\} \simeq \langle Y_e \rangle_{11} (\mathcal{Y}_{\tau}/\mathcal{Y}_e) \sim 10^6 \text{ GeV}$. We find that, for $m_{\mathcal{S}} = 200 \text{ GeV}$, the ratio is bound by $\lambda_{Y_e}/\lambda_{S_e} \lesssim 10^{-2}$.

Even though neutrino masses are purely Dirac in this case, the origin of the tiny neutrino Yukawa couplings is reminiscent of the seesaw mechanism as the couplings to H_u depend on a ratio of dimensionful scales, $M_{\nu}/\langle Y_{\nu} \rangle$. Since these Yukawa couplings are of order $\sim 10^{-12}$, it may be necessary for the $\lambda'_{\nu} M_{\nu}$ combination to be smaller than

its corresponding size in the quark sector to keep the flavon VEVs below the benchmark mediation scale $M_* \sim 10^{10}$ GeV, though we note that our benchmark is merely a lower bound and the mediation scale can be raised.

Since three generations of Dirac neutrinos contain six light states that participate in $SU(3)_L \times SU(3)_N$ gauge interactions, they can potentially be in thermal equilibrium around the time of Big Bang Nucleosynthesis (BBN) and double the successful SM prediction of neutrino species in this epoch: $N_\nu(T_{\text{BBN}}) \simeq 3$ [54]. To maintain the success of the standard scenario in a gauged flavor model with right handed Dirac neutrinos, the gauge bosons that couple to the right handed neutrinos must be sufficiently weakly coupled. For the case of gauge bosons with couplings similar to that of SM $SU(2)_L$, this translate to requiring the gauge bosons be heavier than 3 TeV [21]. In our case, the gauge coupling is smaller by an order of magnitude, so the bound is negligible.

5.2 Seesaw scenario

To include Majorana masses for right handed neutrinos, we have to add additional flavor-charged fields M_N and M_N^c which transform under $SU(3)_N$ as $\bar{6}$ and 6, respectively. In addition to the Y_ν and Y_ν^c bifundamentals, M_N and M_N^c must also acquire holomorphic expectation values in order to properly execute the seesaw mechanism. Modifying the superpotential from eq. (5.1), we now include

$$\Delta W_L = \frac{1}{2} M_N \bar{N} \bar{N} + \mu_N M_N M_N^c, \tag{5.8}$$

which yield Majorana masses for the right handed neutrino \bar{N} when M_N acquires a holomorphic VEV in the UV.

Although this scenario requires additional spurions, the seesaw mechanism loses some of its appeal since the same physics determines the M_N scale and the Yukawa coupling scale. In this framework, making the neutrino-Higgs couplings comparable to those of other SM fields implies $\langle Y_\nu \rangle \sim \langle Y_e \rangle$, however, the same physics that sets the Higgs-Yukawa couplings through $\langle Y \rangle$ also determines the scale of $\langle M_N \rangle$. Thus, it's not clear that the seesaw mechanism in this context improves the naturalness of neutrino masses relative to the purely Dirac case. Nonetheless, the added structure does not significantly modify the lepton story from [13], so we will not consider it further.

6 Conclusions

Since the LHC has already put strong limits on models that predict multijet signals with large missing energy, R -parity violating SUSY is, perhaps, the most compelling viable solution to the hierarchy problem. In particular, baryon number violating RPV features multijet final states with little missing energy, so the dominant signals are difficult to distinguish from the large QCD background at the LHC. These models are, therefore, quite poorly constrained by current data. Unfortunately, abandoning R -parity makes the already-difficult SUSY flavor problem even worse. In the RPV scenario, both superpotential and SUSY breaking sectors introduce dangerous flavor violating operators whose

coefficients must be extremely small for the theory to be phenomenologically viable. Without additional structure to suppress flavor violation, this scenario requires extreme fine tuning.

In this paper we have presented a model which approximately realizes the MFV framework described in [13]. This scenario, while only a hypothesis in [13], provides a solution to SUSY flavor problem by forcing all flavor breaking parameters to be aligned with the Standard Model Yukawa matrices. Requiring the flavor breaking spurions to couple holomorphically ensures that the dominant RPV operators are baryon (not lepton) number violating, so the lightest (pair produced) superpartner decays predominantly to dijets, which are much harder to detect above the large QCD background. In our model, relatively light superpartners are still allowed and the LSP can be any particle in the MSSM. Since the SUSY breaking \mathcal{A} -terms are proportional to Yukawa matrices, it is straightforward to engineer an IR spectrum whose LSP is either a stop or sbottom squark.

The model presented here features a gauged $SU(3)_Q \times SU(3)_U \times SU(3)_D \times SU(3)_L \times SU(3)_E \times SU(3)_N$ flavor group and the minimal additional field content required for anomaly cancellation. On one hand, it is a supersymmetrized version of the model presented in [19], which realizes “flavorful SUSY” [36] whose flavor violating parameters scale with individual Yukawa couplings, but are not necessarily aligned with the full Yukawa matrices. The only terms that are not parametrically of Yukawa size arise from flavor D -terms, which induce anarchic, non-minimal flavor violation. However, in this framework, all deviations away from MFV can be made small enough to avoid all experimental bounds.

On the other hand, this model can be thought of as exactly MFV in a particular limit. This limit is as follows:

- The largest scale in the theory is the scale at which SUSY breaking is communicated to the visible sector, M_* .
- The flavor symmetry is Higgsed by flavon VEVs $\langle Y \rangle$ at a scale parametrically below M_* .
- The self interaction parameters of the flavons defined in eq. (2.2) are small, namely $\mu_Y \ll \langle Y \rangle$ and $\lambda_Y \ll 1$.
- The gauge coupling for the flavor gauge groups are small, $g_F \ll 1$.
- The size of soft SUSY breaking, $m_{\mathcal{S}}$, and the electroweak scale v are parametrically below the flavor breaking scale $\langle Y \rangle$.

While this is a rather complicated limit, we view this not as a corner of parameter space that the model must realize, but instead as a tool for computation. We have elucidated the corrections away from MFV in more general regions of parameter space and find there is much room for them to be small. All low energy and collider constraints have been detailed in section 4 and we find that this model can be viable with weak scale MSSM superpartners.

In this work, we have focused mostly on the quark sector, but we consider the lepton sector in section 5. Introducing a triplet of right handed neutrinos \bar{N} which transform

under their own $SU(3)_N$ leads to a natural realization of pure Dirac neutrinos. While the \bar{N} 's are normally gauge singlets which can get a Majorana mass, the flavor gauge symmetry forbids this Majorana mass, making pure Dirac neutrinos a more natural possibility than the standard seesaw scenario. The neutrino Yukawa couplings must be very small, but they remain technically natural because they arise from superpotential parameters. In order to implement the seesaw scenario, one must add additional flavons which transform as the symmetric 6 and $\bar{6}$ representations of $SU(3)_N$. We briefly consider this scenario in section 5.2.

In addition to the full MSSM spectrum, this model predicts many of the same features as the low-scale gauged flavor model with some exotic states that could be discovered at the LHC. Most of the bounds derived in [19] apply to this scenario, and we have updated the experimental limits to reflect the latest constraints. Current LHC Z' and W' limits require the lightest flavor gauge boson masses to be in the low TeV range, while fourth generation searches require the exotic fermion masses to be above a few hundred GeV. While this model only contains the minimum number of states consistent with supersymmetry and the new gauge symmetry, it has a large number of new states which could be accessible to the LHC in the near future.

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A Exotic mass diagonalization

In the MFV limit the flavor scale satisfies $\langle Y \rangle \gg v, m_S, \mu_Y$ and the flavon self-interaction couplings are negligible $\lambda_Y \ll 1$, the exotics acquire masses almost exclusively from the superpotential in eq. (2.1), and SUSY is an approximate symmetry among the exotic states. In this appendix, we will diagonalize the exotic mass matrix in this limit. For simplicity of exposition, we set the $\mathcal{O}(1)$ coefficients $\lambda_{u,d}, \lambda'_{u,d}$ to unity throughout.

The mass terms in the scalar potential are

$$\sum_{\alpha} \left| \frac{\partial W}{\partial \psi_{\bar{u}\alpha}} \right|^2 = |\langle Y_u \rangle^{\alpha\beta'} \tilde{\psi}_{u\beta'}|^2 = \langle Y_u \rangle_{\alpha\gamma'}^* \langle Y_u \rangle^{\alpha\beta'} \tilde{\psi}_u^{*\gamma'} \tilde{\psi}_{u\beta'} \quad (\text{A.1})$$

$$\sum_{\beta'} \left| \frac{\partial W}{\partial \bar{u}^{\beta'}} \right|^2 = |M_u \tilde{\psi}_{u\beta'}|^2 = M_u^2 \tilde{\psi}_u^{*\beta'} \tilde{\psi}_{u\beta'} \quad (\text{A.2})$$

$$\begin{aligned} \sum_{\beta'} \left| \frac{\partial W}{\partial \psi_{u\beta'}} \right|^2 &= |\langle Y_u \rangle^{\alpha\beta'} \tilde{\psi}_{u^c\alpha} + M_u \tilde{u}^{\beta'}|^2 \\ &= \langle Y_u \rangle_{\gamma\beta'}^* \langle Y_u \rangle^{\alpha\beta'} \tilde{\psi}_{u^c}^{*\gamma} \tilde{\psi}_{u^c\alpha} + M_u^* \langle Y_u \rangle^{\alpha\beta'} \tilde{u}_{\beta'}^* \tilde{\psi}_{u^c\alpha} \\ &\quad + \langle Y_u \rangle_{\gamma\beta'}^* M_u \tilde{\psi}_{u^c}^{*\gamma} \tilde{u}^{\beta'} + M_u^2 \tilde{u}_{\beta'}^* \tilde{u}^{\beta'} \end{aligned} \quad (\text{A.3})$$

which populate the scalar mass-squared matrix $m_{\tilde{\psi}}^2$. In the same basis (prior to EWSB) the fermion mass terms are

$$\langle Y_u \rangle^{\alpha\beta'} \psi_{u^c\alpha} \psi_{u\beta'} + M_u \bar{u}^{\beta'} \psi_{u\beta'} + \text{h.c.}, \quad (\text{A.4})$$

which give rise to the mass matrix $M_{\tilde{\psi}}$. Note that $M_{\tilde{\psi}}^\dagger M_{\tilde{\psi}} = m_{\tilde{\psi}}^2$, so exotic particles are mass degenerate with their superpartners; both matrices $M_{\tilde{\psi}}$ and $m_{\tilde{\psi}}^2$ are diagonalized by the same unitary transformations. In the EWSB preserving limit, this system has three zero eigenvalues whose eigenvectors become the MSSM \bar{U} fields.

For the purpose of illustration, we can perform gauge transformations to make $\langle Y_u \rangle$ diagonal and real so that $\langle Y_u \rangle^{\alpha\beta'} = \langle Y_u^* \rangle_{\alpha\beta'}$, which moves the CKM matrix over to the down-type VEVs. After diagonalizing the mass matrix, the massless MSSM eigenstates $\bar{U}^\alpha \equiv (\bar{U}^1, \bar{U}^2, \bar{U}^3)$ can be written in terms of the corresponding flavor eigenstates $\psi_{u^c}^\alpha$ and \bar{u}^α

$$\bar{U}^\alpha = \frac{1}{\sqrt{1 + \frac{M_u^2}{\langle Y_u \rangle_{\alpha\alpha}^2}}} \left(\bar{u}^\alpha - \frac{M_u}{\langle Y_u \rangle_{\alpha\alpha}} \psi_{u^c}^\alpha \right). \quad (\text{A.5})$$

Since the elements of the diagonalization matrix \mathcal{V} are defined by the ψ_{u^c} coefficients, and $\mathcal{V} \propto \mathcal{Y}$, then (up to order-one factors) the scale of Yukawa couplings is set by the appropriate $\sim M_u/\langle Y_u \rangle$ ratio. This justifies our approach in section 4.2.3 where the bounds on $(\delta_{11}^u)_{LR}$ translate into a lower bound on the $\langle Y_u \rangle_{11}$ element.

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