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Flavorful Supersymmetry

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Abstract

Weak scale supersymmetry provides elegant solutions to many of the problems of the standard model, but it also generically gives rise to excessive flavor and CP violation. We show that if the mechanism that suppresses the Yukawa couplings also suppresses flavor changing interactions in the supersymmetry breaking parameters, essentially all the low energy flavor and CP constraints can be satisfied. The standard assumption of flavor universality in the supersymmetry breaking sector is not necessary. We study signatures of this framework at the LHC. The mass splitting among different generations of squarks and sleptons can be much larger than in conventional scenarios, and even the mass ordering can be changed. We find that there is a plausible scenario in which the NLSP is a long-lived right-handed selectron or smuon decaying into the LSP gravitino. This leads to the spectacular signature of monochromatic electrons or muons in a stopper detector, providing strong evidence for the framework.

1 Introduction

Despite many new alternatives, weak scale supersymmetry is still regarded as the leading candidate for physics beyond the standard model. It not only stabilizes the electroweak scale against potentially large radiative corrections, but also leads to successful gauge coupling unification at high energies and provides a candidate for dark matter. The fact that supersymmetry must be broken, however, leads to a severe flavor and CP problem. Including generic supersymmetry breaking parameters of order the weak scale causes flavor changing and CP violating processes with rates much greater than current experimental bounds. In fact, the problem has become more severe because of recent experimental progress, especially in B physics.

The most common approach to this problem is to assume that supersymmetry breaking and its mediation to the supersymmetric standard model sector preserve flavor. In other words, the mechanism for generating the Yukawa couplings for the quarks and leptons is separate from that which mediates supersymmetry breaking, so the fundamental supersymmetry breaking parameters do not contain any sources of flavor or CP violation. This is typically achieved in one of two ways. The first is to simply assume flavor universality and CP conservation for the supersymmetry breaking parameters at the scale where the low energy field theory arises [1], and the second is to impose a low energy mechanism which leads to flavor universality [2, 3].

A careful look at the problem, however, shows that the situation does not need to be as described above. We know that the Yukawa couplings for the first two generations of quarks and leptons are suppressed, implying that there is some mechanism responsible for this suppression. Suppose that this mechanism suppresses all non-gauge interactions associated with light generation fields, not just the Yukawa couplings. Then the supersymmetry breaking masses for the light generation squarks and sleptons are also suppressed at the scale where the mechanism is operative, suppressing flavor and CP violation associated with these masses. This scenario was considered before in Ref. [4] in the context of reducing fine-tuning in electroweak symmetry breaking. The necessary flavor universal contribution to the squark and slepton masses arises automatically at lower energies from the gaugino masses through renormalization group evolution. An additional contribution may also arise from a low energy mechanism leading to flavor universal supersymmetry breaking masses.

In this paper we study a scenario in which the physics responsible for the quark and lepton masses and mixings is also responsible for the structure of the supersymmetry breaking masses. We call this scenario flavorful supersymmetry in order to emphasize the direct connection between flavor physics and supersymmetry breaking. To preserve the success of gauge coupling unification in the most straightforward way, we assume that this physics lies at or above the unification scale $M_U \approx 10^{16}$ GeV. We find that, in contrast to naive expectations, a large portion of parameter space is not excluded by current experimental data. We study implications of this scenario on the low energy superparticle spectrum, which can be tested at future colliders. In particular, we point out distinct signatures at the LHC, arising in the plausible case where the gravitino is the lightest supersymmetric particle. Throughout the paper we assume that R parity is conserved, although the framework can be extended straightforwardly to the case of R parity violation.

The organization of the paper is as follows. In section 2 we describe our basic framework, and in section 3 we show that it satisfies experimental constraints from low energy flavor and CP violation. In section 4 we discuss implications on the weak scale superparticle spectrum, and we analyze signatures at the LHC in section 5. Finally, conclusions are given in section 6.

2 Framework

Suppose that the supersymmetric standard model, or supersymmetric grand unified theory, arises at a scale M_* ($\gtrsim M_U$) as an effective field theory of some more fundamental theory, which may or may not be a field theory. We consider that the physics generating the Yukawa couplings suppresses all non-gauge interactions associated with the quark, lepton and Higgs superfields Q_i, U_i, D_i, L_i, E_i, H_u and H_d (and N_i if we introduce the right-handed neutrinos), where $i = 1, 2, 3$ is the generation index. In particular, it suppresses the operators generating the supersymmetry breaking masses at the scale M_*:

$$\begin{aligned}
\mathcal{L} \;=\; & \left(\sum_{A=1,2,3} \int d^2\theta\, \eta_A \frac{X}{M_*} \mathcal{W}^{A\alpha} \mathcal{W}^A_\alpha + \text{h.c.} \right) \\
& + \int d^4\theta \Biggl[\kappa_{H_u} \frac{X^\dagger X}{M_*^2} H_u^\dagger H_u + \kappa_{H_d} \frac{X^\dagger X}{M_*^2} H_d^\dagger H_d \\
& \quad + \left(\kappa_\mu \frac{X^\dagger}{M_*} H_u H_d + \kappa_b \frac{X^\dagger X}{M_*^2} H_u H_d + \eta_{H_u} \frac{X}{M_*} H_u^\dagger H_u + \eta_{H_d} \frac{X}{M_*} H_d^\dagger H_d + \text{h.c.} \right) \\
& \quad + (\kappa_\Phi)_{ij} \frac{X^\dagger X}{M_*^2} \Phi_i^\dagger \Phi_j + \left((\eta_\Phi)_{ij} \frac{X}{M_*} \Phi_i^\dagger \Phi_j + \text{h.c.} \right) \Biggr] \\
& + \left[\int d^2\theta \left((\zeta_u)_{ij} \frac{X}{M_*} Q_i U_j H_u + (\zeta_d)_{ij} \frac{X}{M_*} Q_i D_j H_d + (\zeta_e)_{ij} \frac{X}{M_*} L_i E_j H_d \right) + \text{h.c.} \right],
\end{aligned} \qquad (1)$$

where $X = \theta^2 F_X$ is a chiral superfield whose F-term vacuum expectation value is responsible for supersymmetry breaking, $\mathcal{W}^A_\alpha$ ($A = 1, 2, 3$) are the field-strength superfields for $U(1)_Y$, $SU(2)_L$ and $SU(3)_C$, and $\Phi = Q, U, D, L$ and E. The κ_Φ are 3×3 Hermitian matrices, while η_Φ, ζ_u, ζ_d and ζ_e are general complex 3×3 matrices. (Here, we have omitted the operators involving N_i because in most cases they do not affect our analysis.)

Assuming that suppression factors ϵ_{Φ_i}, ϵ_{H_u} and ϵ_{H_d} appear associated with the fields Φ_i, H_u

and H_d, we obtain for the parameters in Eq. (1)

$$\kappa_{H_u} \approx \tilde{\kappa}_{H_u} \epsilon_{H_u}^2, \qquad \kappa_{H_d} \approx \tilde{\kappa}_{H_d} \epsilon_{H_d}^2, \tag{2}$$

$$\kappa_\mu \approx \tilde{\kappa}_\mu \, \epsilon_{H_u} \epsilon_{H_d}, \qquad \kappa_b \approx \tilde{\kappa}_b \, \epsilon_{H_u} \epsilon_{H_d}, \qquad \eta_{H_u} \approx \tilde{\eta}_{H_u} \epsilon_{H_u}^2, \qquad \eta_{H_d} \approx \tilde{\eta}_{H_d} \epsilon_{H_d}^2, \tag{3}$$

$$(\kappa_\Phi)_{ij} \approx \tilde{\kappa}_\Phi \, \epsilon_{\Phi_i} \epsilon_{\Phi_j}, \qquad (\eta_\Phi)_{ij} \approx \tilde{\eta}_\Phi \, \epsilon_{\Phi_i} \epsilon_{\Phi_j}, \tag{4}$$

$$(\zeta_u)_{ij} \approx \tilde{\zeta}_u \, \epsilon_{Q_i} \epsilon_{U_j} \epsilon_{H_u}, \qquad (\zeta_d)_{ij} \approx \tilde{\zeta}_d \, \epsilon_{Q_i} \epsilon_{D_j} \epsilon_{H_d}, \qquad (\zeta_e)_{ij} \approx \tilde{\zeta}_e \, \epsilon_{L_i} \epsilon_{E_j} \epsilon_{H_d}, \tag{5}$$

and for the Yukawa couplings

$$(y_u)_{ij} \approx \tilde{y}_u \, \epsilon_{Q_i} \epsilon_{U_j} \epsilon_{H_u}, \qquad (y_d)_{ij} \approx \tilde{y}_d \, \epsilon_{Q_i} \epsilon_{D_j} \epsilon_{H_d}, \qquad (y_e)_{ij} \approx \tilde{y}_e \, \epsilon_{L_i} \epsilon_{E_j} \epsilon_{H_d}, \tag{6}$$

where tilde parameters represent the "natural" size for the couplings without the suppression factors. For example, if the theory is strongly coupled at M_*, $\tilde{y}_u \sim \tilde{y}_d \sim \tilde{y}_e \sim O(4\pi)$, while if it is weakly coupled, we expect $\tilde{y}_u \sim \tilde{y}_d \sim \tilde{y}_e \sim O(1)$. Note that $O(1)$ coefficients are omitted in the expressions of Eqs. (2 – 6); for example, $(\kappa_\Phi)_{ij}$ is not proportional to $(\eta_\Phi)_{ij}$ because of an arbitrary $O(1)$ coefficient in each element.

Depending on the setup, some of the coefficients may be vanishing. For example, if the supersymmetry breaking sector does not contain an "elementary" gauge singlet at M_*, then $\eta_A = \tilde{\kappa}_\mu = \tilde{\eta}_{H_u} = \tilde{\eta}_{H_d} = \tilde{\eta}_\Phi = \tilde{\zeta}_u = \tilde{\zeta}_d = \tilde{\zeta}_e = 0$, and the gaugino masses must be generated by some low energy mechanism. (The supersymmetric Higgs mass, the μ term, must also be generated at low energies unless it exists at M_* in the superpotential.) The precise pattern for η_A and the tilde parameters affects low energy phenomenology, but our analysis of flavor and CP violation is independent of the detailed pattern.

In this paper we consider the case where $\epsilon_{H_u} \sim \epsilon_{H_d} \sim O(1)$, and assume for simplicity that the two Higgs doublets obey the same scaling, $\tilde{\kappa}_{H_u} \sim \tilde{\kappa}_{H_d} \sim \tilde{\kappa}_H$ and $\tilde{\eta}_{H_u} \sim \tilde{\eta}_{H_d} \sim \tilde{\eta}_H$, as do the matter fields, $\tilde{\kappa}_Q \sim \tilde{\kappa}_U \sim \tilde{\kappa}_D \sim \tilde{\kappa}_L \sim \tilde{\kappa}_E \sim \tilde{\kappa}_\Phi$ and $\tilde{\eta}_Q \sim \tilde{\eta}_U \sim \tilde{\eta}_D \sim \tilde{\eta}_L \sim \tilde{\eta}_E \sim \tilde{\eta}_\Phi$, leading to $\tilde{\zeta}_u \sim \tilde{\zeta}_d \sim \tilde{\zeta}_e \sim \tilde{\zeta}$ and $\tilde{y}_u \sim \tilde{y}_d \sim \tilde{y}_e \sim \tilde{y}$. An extension to more general cases is straightforward. The supersymmetry breaking (and μ) parameters are then obtained from Eqs. (1 – 5) as

$$M_A \approx \eta_A M_{\rm SUSY}, \qquad \mu \approx \tilde{\kappa}_\mu M_{\rm SUSY}^\dagger, \qquad b \approx (\tilde{\kappa}_b + \tilde{\kappa}_\mu \tilde{\eta}_H) |M_{\rm SUSY}|^2, \tag{7}$$

$$m_{H_u}^2 \approx m_{H_d}^2 \approx (\tilde{\kappa}_H + |\tilde{\eta}_H|^2) |M_{\rm SUSY}|^2, \qquad (m_\Phi^2)_{ij} \approx \{(\kappa_\Phi)_{ij} + (\eta_\Phi^\dagger \eta_\Phi)_{ij}\} |M_{\rm SUSY}|^2, \tag{8}$$

$$(a_u)_{ij} \approx \{(y_u)_{kj} (\eta_Q)_{ki} + (y_u)_{ik} (\eta_U)_{kj} + (y_u)_{ij} \tilde{\eta}_H\} M_{\rm SUSY} + \tilde{\zeta} \, \epsilon_{Q_i} \epsilon_{U_j} M_{\rm SUSY}, \tag{9}$$

$$(a_d)_{ij} \approx \{(y_d)_{kj} (\eta_Q)_{ki} + (y_d)_{ik} (\eta_D)_{kj} + (y_d)_{ij} \tilde{\eta}_H\} M_{\rm SUSY} + \tilde{\zeta} \, \epsilon_{Q_i} \epsilon_{D_j} M_{\rm SUSY}, \tag{10}$$

$$(a_e)_{ij} \approx \{(y_e)_{kj} (\eta_L)_{ki} + (y_e)_{ik} (\eta_E)_{kj} + (y_e)_{ij} \tilde{\eta}_H\} M_{\rm SUSY} + \tilde{\zeta} \, \epsilon_{L_i} \epsilon_{E_j} M_{\rm SUSY}, \tag{11}$$

where $M_{\rm SUSY} \equiv F_X/M_*$, and M_A are the gaugino masses, $m_{H_u}^2$, $m_{H_d}^2$ and m_Φ^2 are non-holomorphic supersymmetry breaking squared masses, b is the holomorphic supersymmetry breaking Higgs mass-squared, and $(a_u)_{ij}$, $(a_d)_{ij}$ and $(a_e)_{ij}$ are holomorphic supersymmetry breaking scalar trilinear interactions. We find that the pattern of the supersymmetry breaking parameters is correlated with that of the Yukawa couplings, which now read

$$(y_u)_{ij} \approx \tilde{y}\, \epsilon_{Q_i}\epsilon_{U_j}, \qquad (y_d)_{ij} \approx \tilde{y}\, \epsilon_{Q_i}\epsilon_{D_j}, \qquad (y_e)_{ij} \approx \tilde{y}\, \epsilon_{L_i}\epsilon_{E_j}. \tag{12}$$

In general, the correlation between Eqs. (7 – 11) and Eq. (12) significantly reduces the tension between supersymmetry breaking and flavor physics [4]. We note again that $O(1)$ coefficients are omitted in each term in Eqs. (7 – 12); for instance, the last terms of Eqs. (9 – 11) are not proportional to the corresponding Yukawa matrices, Eq. (12), because of these $O(1)$ coefficients.

Taking $\epsilon_{\Phi_1} \leq \epsilon_{\Phi_2} \leq \epsilon_{\Phi_3}$ without loss of generality, the Yukawa couplings of Eq. (12) lead to the following quark and lepton masses and mixings

$$\begin{array}{lcl} (m_t, m_c, m_u) & \approx & \tilde{y}\, \langle H_u \rangle\, (\epsilon_{Q_3}\epsilon_{U_3},\, \epsilon_{Q_2}\epsilon_{U_2},\, \epsilon_{Q_1}\epsilon_{U_1}), \\ (m_b, m_s, m_d) & \approx & \tilde{y}\, \langle H_d \rangle\, (\epsilon_{Q_3}\epsilon_{D_3},\, \epsilon_{Q_2}\epsilon_{D_2},\, \epsilon_{Q_1}\epsilon_{D_1}), \\ (m_\tau, m_\mu, m_e) & \approx & \tilde{y}\, \langle H_d \rangle\, (\epsilon_{L_3}\epsilon_{E_3},\, \epsilon_{L_2}\epsilon_{E_2},\, \epsilon_{L_1}\epsilon_{E_1}), \\ (m_{\nu_\tau}, m_{\nu_\mu}, m_{\nu_e}) & \approx & \frac{\tilde{y}^2 \langle H_u \rangle^2}{M_N}\, (\epsilon_{L_3}^2,\, \epsilon_{L_2}^2,\, \epsilon_{L_1}^2), \end{array} \tag{13}$$

and

$$V_{\rm CKM} \approx \begin{pmatrix} 1 & \epsilon_{Q_1}/\epsilon_{Q_2} & \epsilon_{Q_1}/\epsilon_{Q_3} \\ \epsilon_{Q_1}/\epsilon_{Q_2} & 1 & \epsilon_{Q_2}/\epsilon_{Q_3} \\ \epsilon_{Q_1}/\epsilon_{Q_3} & \epsilon_{Q_2}/\epsilon_{Q_3} & 1 \end{pmatrix}, \qquad V_{\rm MNS} \approx \begin{pmatrix} 1 & \epsilon_{L_1}/\epsilon_{L_2} & \epsilon_{L_1}/\epsilon_{L_3} \\ \epsilon_{L_1}/\epsilon_{L_2} & 1 & \epsilon_{L_2}/\epsilon_{L_3} \\ \epsilon_{L_1}/\epsilon_{L_3} & \epsilon_{L_2}/\epsilon_{L_3} & 1 \end{pmatrix}, \tag{14}$$

where we have included the neutrino masses through the seesaw mechanism with the right-handed neutrino Majorana masses $W \approx M_N \epsilon_{N_i}\epsilon_{N_j} N_i N_j$, and $V_{\rm CKM}$ and $V_{\rm MNS}$ are the quark and lepton mixing matrices, respectively. The values of the ϵ parameters are then constrained by the observed quark and lepton masses and mixings.

There are a variety of possibilities for the origin of the ϵ factors. They may arise, for example, from distributions of fields in higher dimensional spacetime or from strong conformal dynamics at or above the scale M_*. In a forthcoming paper we will discuss an explicit example of such models. In general, if the suppressions of the Yukawa couplings arise from wavefunction effects in a broad sense, as in the examples described above, we can obtain the correlation given in Eqs. (7 – 11) and Eq. (12). Another possibility is to introduce a non-Abelian flavor symmetry connecting all three generations. Flavor violating supersymmetry breaking parameters having a similar correlation to Eqs. (7 – 12) may then be generated through the breaking of that symmetry.[1] While this allows flavor universal contributions to the supersymmetry breaking parameters in addition to Eqs. (7 – 11), the essential features of the framework are not affected.

[1] For earlier analyses on flavor violation in models with non-Abelian flavor symmetries, see e.g. [5].

3 Constraints from Low Energy Processes

The supersymmetry breaking parameters are subject to a number of constraints from low energy flavor and CP violating processes. Here we study these constraints for the parameters given in Eqs. (7 – 11). We assume that CP violating effects associated with the Higgs sector are sufficiently suppressed. This is achieved if either $b \ll |\mu|^2$ at M_* or the phases of μ and b are aligned in the basis where the M_A are real.

The values of low energy supersymmetry breaking parameters are obtained from Eqs. (7 – 11) by evolving them down to the weak scale using renormalization group equations. Contributions from other flavor universal sources, such as gauge mediation, may also be added. To parameterize these effects in a model-independent manner, we simply add universal squark and slepton squared masses, $m_{\tilde{q}}^2 \equiv \lambda_{\tilde{q}}^2 \, |M_{\rm SUSY}|^2$ and $m_{\tilde{l}}^2 \equiv \lambda_{\tilde{l}}^2 \, |M_{\rm SUSY}|^2$, to $(m_\Phi^2)_{ij}$:

$$(m_\Phi^2)_{ij} \rightarrow \left\{ \begin{array}{ll} (m_\Phi^2)_{ij} + \lambda_{\tilde{q}}^2 \, |M_{\rm SUSY}|^2 \delta_{ij} & \text{for } \Phi = Q, U, D \\ (m_\Phi^2)_{ij} + \lambda_{\tilde{l}}^2 \, |M_{\rm SUSY}|^2 \delta_{ij} & \text{for } \Phi = L, E \end{array} \right. . \tag{15}$$

We neglect the differences of the flavor universal contribution among various squarks and among various sleptons, but it is sufficient for our purposes here. Effects on the gaugino masses and the scalar trilinear interactions are absorbed into the redefinition of η_A and $\tilde{\eta}_H$, respectively. The $m_{H_u}^2$ and $m_{H_d}^2$ are also renormalized, but this effect is incorporated by treating $\tan\beta \equiv \langle H_u \rangle / \langle H_d \rangle$ as a free parameter.

We use the mass insertion method [6] to compare the expected amount of flavor violation in the present scenario to low energy data. In order to do a mass insertion analysis, we need to work in the super-CKM basis where the Yukawa matrices are diagonalized by supersymmetric rotations of Q_i, U_i, D_i, L_i and E_i. The mass insertion parameters, δ_{ij}, are then obtained by dividing the off-diagonal entry of the sfermion mass-squared matrix by the average diagonal entry. Using Eqs. (7 – 11, 15), we obtain

$$(\delta^u_{ij})_{LL} \approx \frac{1}{\lambda_{\tilde{q}}^2} \Big(\tilde{\kappa}_\Phi + |\tilde{\eta}_\Phi|^2 \epsilon_{Q_3}^2 \Big) \, \epsilon_{Q_i} \epsilon_{Q_j}, \qquad (\delta^u_{ij})_{RR} \approx \frac{1}{\lambda_{\tilde{q}}^2} \Big(\tilde{\kappa}_\Phi + |\tilde{\eta}_\Phi|^2 \epsilon_{U_3}^2 \Big) \, \epsilon_{U_i} \epsilon_{U_j}, \tag{16}$$

$$(\delta^u_{ij})_{LR} = (\delta^u_{ji})^*_{RL} \approx \frac{1}{\lambda_{\tilde{q}}^2} \Big\{ \tilde{y} \, \tilde{\eta}_\Phi (\epsilon_{Q_j}^2 + \epsilon_{U_i}^2) + \tilde{\zeta} \Big\} \, \epsilon_{Q_i} \epsilon_{U_j} \frac{v \sin\beta}{M_{\rm SUSY}}, \tag{17}$$

for the up-type squarks,

$$(\delta^d_{ij})_{LL} \approx \frac{1}{\lambda_{\tilde{q}}^2} \Big(\tilde{\kappa}_\Phi + |\tilde{\eta}_\Phi|^2 \epsilon_{Q_3}^2 \Big) \, \epsilon_{Q_i} \epsilon_{Q_j}, \qquad (\delta^d_{ij})_{RR} \approx \frac{1}{\lambda_{\tilde{q}}^2} \Big(\tilde{\kappa}_\Phi + |\tilde{\eta}_\Phi|^2 \epsilon_{D_3}^2 \Big) \, \epsilon_{D_i} \epsilon_{D_j}, \tag{18}$$

$$(\delta^d_{ij})_{LR} = (\delta^d_{ji})^*_{RL} \approx \frac{1}{\lambda_{\tilde{q}}^2} \Big\{ \tilde{y} \, \tilde{\eta}_\Phi (\epsilon_{Q_j}^2 + \epsilon_{D_i}^2) + \tilde{\zeta} \Big\} \, \epsilon_{Q_i} \epsilon_{D_j} \frac{v \cos\beta}{M_{\rm SUSY}}, \tag{19}$$

for the down-type squarks,

$$(\delta^e_{ij})_{LL} \approx \frac{1}{\lambda^2_{\tilde{l}}}\Big(\tilde{\kappa}_\Phi + |\tilde{\eta}_\Phi|^2\epsilon^2_{L_3}\Big)\,\epsilon_{L_i}\epsilon_{L_j}, \qquad (\delta^e_{ij})_{RR} \approx \frac{1}{\lambda^2_{\tilde{l}}}\Big(\tilde{\kappa}_\Phi + |\tilde{\eta}_\Phi|^2\epsilon^2_{E_3}\Big)\,\epsilon_{E_i}\epsilon_{E_j}, \tag{20}$$

$$(\delta^e_{ij})_{LR} = (\delta^e_{ji})^*_{RL} \approx \frac{1}{\lambda^2_{\tilde{l}}}\Big\{\tilde{y}\,\tilde{\eta}_\Phi(\epsilon^2_{L_j} + \epsilon^2_{E_i}) + \tilde{\zeta}\Big\}\,\epsilon_{L_i}\epsilon_{E_j}\frac{v\cos\beta}{M_{\rm SUSY}}, \tag{21}$$

for the charged sleptons, and

$$(\delta^\nu_{ij})_{LL} \approx \frac{1}{\lambda^2_{\tilde{l}}}\Big(\tilde{\kappa}_\Phi + |\tilde{\eta}_\Phi|^2\epsilon^2_{L_3}\Big)\,\epsilon_{L_i}\epsilon_{L_j}, \tag{22}$$

for the sneutrinos. Here, $v \equiv (\langle H_u\rangle^2 + \langle H_d\rangle^2)^{1/2} \simeq 174$ GeV and $\tan\beta = \langle H_u\rangle/\langle H_d\rangle$.

The values of the ϵ parameters are constrained to reproduce the observed quark and lepton masses and mixings through Eqs. (13, 14). They depend on $\tilde{y}$ as well as the value of $\tan\beta$. For illustrative purpose, we take the pattern

$$\begin{array}{lll} \epsilon_{Q_1} \approx \tilde{y}^{-\frac{1}{2}}\alpha_q\,\epsilon^2, & \epsilon_{U_1} \approx \tilde{y}^{-\frac{1}{2}}\alpha_q^{-1}\epsilon^2, & \epsilon_{D_1} \approx \tilde{y}^{-\frac{1}{2}}\alpha_q^{-1}\alpha_\beta\,\epsilon, \\ \epsilon_{Q_2} \approx \tilde{y}^{-\frac{1}{2}}\alpha_q\,\epsilon, & \epsilon_{U_2} \approx \tilde{y}^{-\frac{1}{2}}\alpha_q^{-1}\epsilon, & \epsilon_{D_2} \approx \tilde{y}^{-\frac{1}{2}}\alpha_q^{-1}\alpha_\beta\,\epsilon, \\ \epsilon_{Q_3} \approx \tilde{y}^{-\frac{1}{2}}\alpha_q, & \epsilon_{U_3} \approx \tilde{y}^{-\frac{1}{2}}\alpha_q^{-1}, & \epsilon_{D_3} \approx \tilde{y}^{-\frac{1}{2}}\alpha_q^{-1}\alpha_\beta\,\epsilon, \end{array} \tag{23}$$

$$\begin{array}{ll} \epsilon_{L_1} \approx \tilde{y}^{-\frac{1}{2}}\alpha_l\,\epsilon, & \epsilon_{E_1} \approx \tilde{y}^{-\frac{1}{2}}\alpha_l^{-1}\alpha_\beta\,\epsilon^2, \\ \epsilon_{L_2} \approx \tilde{y}^{-\frac{1}{2}}\alpha_l\,\epsilon, & \epsilon_{E_2} \approx \tilde{y}^{-\frac{1}{2}}\alpha_l^{-1}\alpha_\beta\,\epsilon, \\ \epsilon_{L_3} \approx \tilde{y}^{-\frac{1}{2}}\alpha_l\,\epsilon, & \epsilon_{E_3} \approx \tilde{y}^{-\frac{1}{2}}\alpha_l^{-1}\alpha_\beta, \end{array} \tag{24}$$

with

$$\tan\beta \approx \alpha_\beta\,\epsilon^{-1}, \tag{25}$$

where $\epsilon \sim O(0.1)$ and α_q, α_l and α_β are numbers parameterizing the freedoms unfixed by the data of the quark and lepton masses and mixings. Here, we have assumed that $\tan\beta$ is larger than ≈ 2, as suggested by the large top quark mass. The pattern of Eq. (23 – 25) leads to

$$\begin{array}{lcl} (m_t, m_c, m_u) & \approx & v\,(1, \epsilon^2, \epsilon^4), \\ (m_b, m_s, m_d) & \approx & v\,(\epsilon^2, \epsilon^3, \epsilon^4), \\ (m_\tau, m_\mu, m_e) & \approx & v\,(\epsilon^2, \epsilon^3, \epsilon^4), \\ (m_{\nu_\tau}, m_{\nu_\mu}, m_{\nu_e}) & \approx & \frac{v^2}{M_N}(1, 1, 1), \end{array} \tag{26}$$

and

$$V_{\rm CKM} \approx \begin{pmatrix} 1 & \epsilon & \epsilon^2 \\ \epsilon & 1 & \epsilon \\ \epsilon^2 & \epsilon & 1 \end{pmatrix}, \qquad V_{\rm MNS} \approx \begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix}, \tag{27}$$

which successfully reproduces the gross structure of the observed quark and lepton masses and mixings [7]. The mass insertion parameters are obtained by substituting Eqs. (23 – 25) into Eqs. (16 – 22).

Here we summarize the constraints from low energy flavor and CP violating processes, compiled from Ref. [8]. In the quark sector, the most stringent experimental constraints come from K-$\bar{K}$, D-$\bar{D}$ and B-$\bar{B}$ mixings, $\sin 2\beta$ and the $b \to s\gamma$ process. The model-independent constraints are obtained by turning on only one (or two) mass insertion parameter(s) and considering the gluino exchange diagrams. They are summarized as

$$\sqrt{|\mathrm{Re}(\delta^d_{12})^2_{LL/RR}|} \lesssim (10^{-2}\text{–}10^{-1}), \quad \sqrt{|\mathrm{Re}(\delta^d_{12})^2_{LR/RL}|} \lesssim (10^{-3}\text{–}10^{-2}), \quad \sqrt{|\mathrm{Re}(\delta^d_{12})_{LL}(\delta^d_{12})_{RR}|} \lesssim 10^{-3},$$

$$\sqrt{|\mathrm{Im}(\delta^d_{12})^2_{LL/RR}|} \lesssim (10^{-3}\text{–}10^{-2}), \quad \sqrt{|\mathrm{Im}(\delta^d_{12})^2_{LR/RL}|} \lesssim (10^{-4}\text{–}10^{-3}), \quad \sqrt{|\mathrm{Im}(\delta^d_{12})_{LL}(\delta^d_{12})_{RR}|} \lesssim 10^{-4},$$

$$|(\delta^u_{12})_{LL/RR}| \lesssim (10^{-2}\text{–}10^{-1}), \quad |(\delta^u_{12})_{LR/RL}| \lesssim 10^{-2}, \quad |(\delta^u_{12})_{LL}| = |(\delta^u_{12})_{RR}| \lesssim (10^{-3}\text{–}10^{-2}),$$

$$|(\delta^d_{13})_{LL/RR}| \lesssim (0.1\text{–}1), \quad |(\delta^d_{13})_{LR/RL}| \lesssim (10^{-2}\text{–}10^{-1}), \quad |(\delta^d_{13})_{LL}| = |(\delta^d_{13})_{RR}| \lesssim 10^{-2},$$

$$|(\delta^d_{23})_{LR/RL}| \lesssim 10^{-2},$$

where we have taken the gluino and squark masses to be the same order of magnitude and $m_{\tilde{q}} = 500$ GeV. For heavier superparticles, the bounds scale roughly linearly with $m_{\tilde{q}}$, i.e. all the bounds weaken for larger $m_{\tilde{q}}$ by a factor of $m_{\tilde{q}}/500$ GeV. In the lepton sector, the most stringent constraint comes from the $\mu \to e\gamma$ process, and is given by

$$|(\delta^e_{12})_{LL}| \lesssim (10^{-4}\text{–}10^{-3}), \qquad |(\delta^e_{12})_{LR/RL}| \lesssim (10^{-6}\text{–}10^{-5}),$$

where we have taken the weak gaugino and slepton masses to be the same order of magnitude and $m_{\tilde{l}} = 200$ GeV. For heavier superparticles, the bound on $|(\delta^e_{12})_{LL}|$ scales roughly quadratically with $m_{\tilde{l}}$, while that on $|(\delta^e_{12})_{LR/RL}|$ scales roughly linearly with $m_{\tilde{l}}$. Finally, the bounds from the neutron and electron electric dipole moments (EDMs) constrain the flavor conserving entry of the sfermion mass matrices:

$$|\mathrm{Im}(\delta^u_{11})_{LR}| \lesssim 10^{-6}, \qquad |\mathrm{Im}(\delta^d_{11})_{LR}| \lesssim 10^{-6}, \qquad |\mathrm{Im}(\delta^e_{11})_{LR}| \lesssim 10^{-7},$$

where we have again taken $m_{\tilde{q}} = 500$ GeV and $m_{\tilde{l}} = 200$ GeV, and the bounds become weaker linearly with increasing superparticle masses.

We now determine whether flavor and CP violation arising from Eqs. (16 – 25) is compatible with the experimental bounds given above. We take $\epsilon \approx (0.05 - 0.1)$ to reproduce the gross structure of the quark and lepton masses and mixings, and take $\tilde{\kappa}_\Phi \sim \tilde{\eta}_\Phi \sim O(1)$ for simplicity. For $\tilde{\zeta} \sim O(1)$, we find stringent constraints coming from the electron EDM and $\mu \to e\gamma$, which push the supersymmetry breaking scale up to $M_{\rm SUSY} \gtrsim 5$ TeV for $\tilde{y} \sim 1$ and $M_{\rm SUSY} \gtrsim 1.5$ TeV for $\tilde{y} \sim 4\pi$. Here, we have taken $M_{\rm SUSY} \approx m_{\tilde{q}} \approx (5/2) m_{\tilde{l}}$. This implies that the superpotential couplings in Eq. (1), $\tilde{\zeta}$, must somehow be suppressed, unless the superparticles are relatively heavy. This may naturally arise from physics above M_*, since the superpotential has the special property of not being renormalized at all orders in perturbation theory.

For $\tilde{\zeta} \ll 1$, a wide parameter region is open. For $\tilde{y} \sim 1$, we find that the region

$$0.2 \lesssim \alpha_q \lesssim 3, \qquad \frac{\alpha_q}{\alpha_\beta} \gtrsim 0.5, \qquad \alpha_l \lesssim 0.3, \qquad \frac{\alpha_l}{\alpha_\beta} \gtrsim 0.2, \tag{28}$$

satisfies all the constraints for $\epsilon = 0.05$, $M_{\rm SUSY} = m_{\tilde{q}} = 500$ GeV and $m_{\tilde{l}} = 200$ GeV. The region is somewhat smaller for $\epsilon = 0.1$. For $\tilde{y} \sim 4\pi$, we find the region

$$0.05 \lesssim \alpha_q \lesssim 10, \qquad \frac{\alpha_q}{\alpha_\beta} \gtrsim 0.1, \qquad \alpha_l \lesssim 1, \qquad \frac{\alpha_l}{\alpha_\beta} \gtrsim 0.04, \tag{29}$$

for $\epsilon = 0.05$, $M_{\rm SUSY} = m_{\tilde{q}} = 500$ GeV and $m_{\tilde{l}} = 200$ GeV, and somewhat smaller for $\epsilon = 0.1$.[2] This result agrees with that of Ref. [4], which analyzed the case of $\tilde{y} \sim 1$ without including the constraints from the EDM bounds.

We have used the particular parameterization of Eq. (23 – 25) in the analysis here, but we can adopt a more refined scheme for the values of the ϵ parameters to better accommodate the observed quark and lepton masses and mixings. For example, we can make ϵ_{L_1} somewhat smaller than Eq. (24) to explain the smallness of the $e3$ element of $V_{\rm MNS}$, which is experimentally smaller than about 0.2. A value of $\tan\beta$ somewhat larger than Eq. (25) also improves the top to bottom mass ratio. Our basic results above are not affected by these modifications.

We conclude that the current experimental constraints allow the existence of the general supersymmetry breaking parameters of Eq. (1) where the couplings are suppressed by the factors suggested by Yukawa couplings, as long as the superparticles are relatively heavy or the superpotential couplings, $\tilde{\zeta}$, are suppressed. For $\tilde{\zeta} \ll 1$, a wide parameter region is open even for light superparticles, $m_{\tilde{q}} \approx 500$ GeV and $m_{\tilde{l}} \approx 200$ GeV. In the next section, we discuss implications of this scenario, which we call flavorful supersymmetry, on the low energy spectrum.

4 Implications on the Superparticle Spectrum

Phenomenology of supersymmetric theories depends strongly on the spectrum of superparticles. In particular, the order of the superparticle masses controls decay chains, and thus affects collider signatures significantly. In this section we discuss the splitting and ordering of the superparticle masses among different generations and between different superparticle species.

[2] If we require the absence of cancellation among diagrams for $\mu \to e\gamma$, the bound $|(\delta^e_{12})_{RR}| \lesssim (10^{-3}\text{–}10^{-2})$ arises. This, however, changes the regions of Eqs. (28, 29) only slightly. The lower bound on α_l/α_β becomes 0.5 in Eq. (28), and 0.1 in Eq. (29).

4.1 Mass splitting and ordering among generations

Among the various sfermions, the lightest species are most likely the right-handed sleptons: $\tilde{e}_R$, $\tilde{\mu}_R$ and $\tilde{\tau}_R$. This is because the sfermion squared masses receive positive contributions from the gaugino masses through renormalization group evolution at one loop, and these contributions are proportional to the square of the relevant gauge couplings. Since the right-handed sleptons are charged under only $U(1)_Y$, they receive contributions from just the hypercharge gaugino and are expected to be lighter than the other sfermions. A possible low energy gauge mediated contribution will not change the situation because it gives positive contributions to the sfermion squared masses proportional to the fourth power of the relevant gauge couplings (at least in the simplest case). Thus we focus on the right-handed sleptons and analyze the mass splitting among the three generations. A similar analysis, however, can also be performed for the other sfermion species.[3]

We consider the field basis in which the lepton Yukawa couplings, $(y_e)_{ij}$, are real and diagonal. If there is no intrinsic flavor violation in the sfermion masses, the 3×3 mass-squared matrix for the right-handed sleptons, m_E^2, receives a flavor universal contribution, $m_{\tilde{e}}^2\,\mathrm{diag}(1,1,1)$, and flavor dependent contributions through renormalization group evolution. This leads to

$$m_E^2 = \begin{pmatrix} m_{\tilde{e}}^2 - I_e & 0 & 0 \\ 0 & m_{\tilde{e}}^2 - I_\mu & 0 \\ 0 & 0 & m_{\tilde{e}}^2 - I_\tau \end{pmatrix}, \qquad (30)$$

at the weak scale, where I_e, I_μ and I_τ parameterize the effect of renormalization group evolution from the Yukawa and scalar trilinear couplings, and $I_e : I_\mu : I_\tau \approx (y_e)_{11}^2 : (y_e)_{22}^2 : (y_e)_{33}^2 \approx m_e^2 : m_\mu^2 : m_\tau^2$. (The effects from the neutrino Yukawa couplings that may exist above the scale of right-handed neutrino masses, M_N, are neglected here since they are expected to be small.) The expression of Eq. (30) tells us that, in the absence of a flavor violating contribution, (i) the interaction and mass eigenstates of the right-handed sleptons coincide, and (ii) the mass of a slepton corresponding to a heavier lepton is always lighter, since $I_\tau > I_\mu > I_e > 0$ due to the form of the renormalization group equations of the supersymmetric standard model when $(m_E^2)_{ii}, (m_L^2)_{ii}, m_{H_d}^2 > 0$. Inclusion of flavor universal left-right mixing does not change these conclusions.

The situation is very different if there exists intrinsic flavor violation in supersymmetry breaking. The supersymmetry breaking parameters at M_* in our scenario are given parametrically by

[3] The mass splitting and ordering for heavier species may also provide important tests for flavorful supersymmetry. Moreover, if there exists a $U(1)_Y$ D-term contribution, i.e. $m_{H_u}^2 - m_{H_d}^2 + \mathrm{Tr}[m_Q^2 - 2m_U^2 + m_D^2 - m_L^2 + m_E^2] \neq 0$, then the left-handed sleptons and sneutrinos may be lighter than the right-handed sleptons. It is also possible to consider the case in which a squark is the lightest sfermion if M_3 is significantly smaller than $M_{1,2}$ at M_*.

Eqs. (7 – 11, 4) even in the basis where $(y_e)_{ij}$ is diagonal. In particular,

$$m_E^2(M_*) \approx \begin{pmatrix} \epsilon_{E_1}^2 & \epsilon_{E_1}\epsilon_{E_2} & \epsilon_{E_1}\epsilon_{E_3} \\ \epsilon_{E_1}\epsilon_{E_2} & \epsilon_{E_2}^2 & \epsilon_{E_2}\epsilon_{E_3} \\ \epsilon_{E_1}\epsilon_{E_3} & \epsilon_{E_2}\epsilon_{E_3} & \epsilon_{E_3}^2 \end{pmatrix} |M_{\rm SUSY}|^2, \tag{31}$$

for $\tilde{\kappa}_\Phi \sim \tilde{\eta}_\Phi \sim O(1)$. In addition, m_E^2 receives universal contributions from the $U(1)_Y$ gaugino mass through renormalization group evolution, as well as possibly from other sources such as low energy gauge mediation. It also receives flavor violating contributions from the Yukawa and scalar trilinear couplings through renormalization group evolution. We find that the evolution effect on the off-diagonal elements is not significant; the changes of the coefficients are at most of order unity. The diagonal elements receive flavor universal contributions, which we denote as $m_{\tilde{e}}^2 \equiv \xi_{\tilde{e}}^2 |M_{\rm SUSY}|^2$, as well as flavor dependent contributions. Defining the flavor dependent part as $\hat{m}_{E_i}^2 \equiv (m_E^2)_{ii} - (m_E^2)_{ii}|_{y_e=a_e=0}$, the renormalization group equation for $\hat{m}_{E_i}^2$ is given by

$$\frac{d}{d\ln\mu_R}\hat{m}_{E_i}^2 = \frac{1}{4\pi^2}\left[(y_e)_{ii}^2\left\{(m_E^2)_{ii} + (m_L^2)_{ii} + m_{H_d}^2\right\} + \sum_k |(a_e)_{ki}|^2\right], \tag{32}$$

where i in the right-hand-side is not summed. This leads to m_E^2 at the weak scale of the form

$$m_E^2 \approx \begin{pmatrix} m_{\tilde{e}}^2 - K_e & 0 & 0 \\ 0 & m_{\tilde{e}}^2 - K_\mu & 0 \\ 0 & 0 & m_{\tilde{e}}^2 - K_\tau \end{pmatrix} + \begin{pmatrix} \epsilon_{E_1}^2 & \epsilon_{E_1}\epsilon_{E_2} & \epsilon_{E_1}\epsilon_{E_3} \\ \epsilon_{E_1}\epsilon_{E_2} & \epsilon_{E_2}^2 & \epsilon_{E_2}\epsilon_{E_3} \\ \epsilon_{E_1}\epsilon_{E_3} & \epsilon_{E_2}\epsilon_{E_3} & \epsilon_{E_3}^2 \end{pmatrix} |M_{\rm SUSY}|^2, \tag{33}$$

where $O(1)$ coefficients are omitted in each element in the second term, but not in the first term. The quantities K_e, K_μ and K_τ are defined by $K_\tau \equiv \hat{m}_{E_3}^2(M_*) - \hat{m}_{E_3}^2(M_{\rm SUSY})$ and $\{\tau, 3\} \rightarrow \{e, 1\}, \{\mu, 2\}$, and are given by solving Eq. (32). They are always positive for $(m_E^2)_{ii}, (m_L^2)_{ii}, m_{H_d}^2 > 0$, and $K_e : K_\mu : K_\tau \approx (y_e)_{11}^2 : (y_e)_{22}^2 : (y_e)_{33}^2$ for $(a_e)_{ij} \propto (y_e)_{ij}$.

The contributions K_e, K_μ and K_τ compete in general with the second term in Eq. (33). For $\tilde{\kappa}_\Phi \sim \tilde{\eta}_\Phi \sim \tilde{\eta}_H \sim O(1)$ and $\tilde{\zeta} \lesssim \tilde{y}$, for example, Eq. (32) scales as

$$\frac{d}{d\ln\mu_R}\hat{m}_{E_i}^2 \approx \frac{1}{4\pi^2}(y_e)_{ii}^2\left(\xi_{\tilde{e}}^2 + 2\xi_{\tilde{l}}^2 + \sum_k \frac{\epsilon_{L_k}^2}{\epsilon_{L_i}^2}\right)|M_{\rm SUSY}|^2, \tag{34}$$

where we have set $(m_L^2)_{ii} \approx m_{H_d}^2 \equiv \xi_{\tilde{l}}^2 |M_{\rm SUSY}|^2$. With the choice of Eqs. (23 – 25), this gives

$$K_\tau \approx \frac{1}{4\pi^2}(y_e)_{33}^2\left(\xi_{\tilde{e}}^2 + 2\xi_{\tilde{l}}^2 + O(1)\right)|M_{\rm SUSY}|^2 \ln\frac{M_*}{M_{\rm SUSY}} \sim \tilde{y}^2\epsilon_{L_3}^2\epsilon_{E_3}^2|M_{\rm SUSY}|^2, \tag{35}$$

and $\{\tau, 3\} \rightarrow \{e, 1\}, \{\mu, 2\}$, leading to

$$m_E^2 \approx \begin{pmatrix} \xi_{\tilde{e}}^2 - \tilde{y}^2\epsilon_{L_1}^2\epsilon_{E_1}^2 + \epsilon_{E_1}^2 & \epsilon_{E_1}\epsilon_{E_2} & \epsilon_{E_1}\epsilon_{E_3} \\ \epsilon_{E_1}\epsilon_{E_2} & \xi_{\tilde{e}}^2 - \tilde{y}^2\epsilon_{L_2}^2\epsilon_{E_2}^2 + \epsilon_{E_2}^2 & \epsilon_{E_2}\epsilon_{E_3} \\ \epsilon_{E_1}\epsilon_{E_3} & \epsilon_{E_2}\epsilon_{E_3} & \xi_{\tilde{e}}^2 - \tilde{y}^2\epsilon_{L_3}^2\epsilon_{E_3}^2 + \epsilon_{E_3}^2 \end{pmatrix} |M_{\rm SUSY}|^2. \tag{36}$$

Note that the signs of the $\tilde{y}^2\epsilon_{L_i}^2\epsilon_{E_i}^2$ terms are all negative, while each $\epsilon_{E_i}\epsilon_{E_j}$ term has an $O(1)$ coefficient whose sign can be either positive or negative.

The expression of Eq. (36) shows that in flavorful supersymmetry (i) the interaction and mass eigenstates of the right-handed sleptons do not in general coincide, and (ii) the mass ordering of the sleptons is not necessarily anticorrelated with that of the leptons. In particular, we find that the lightest sfermion can easily be $\tilde{e}_R$ or $\tilde{\mu}_R$ (with slight mixtures from other flavors), in contrast to the usual supersymmetry breaking scenarios in which $\tilde{\tau}_R$ is the lightest because $I_\tau > I_\mu > I_e > 0$ in Eq. (30). In our case, $\tilde{\tau}_R$ is heavier than $\tilde{e}_R$ and $\tilde{\mu}_R$ if, for example, $\tilde{y} \sim 1$ and the $\epsilon_{E_3}^2$ term in the 3-3 entry of Eq. (36) has a positive coefficient. As we will see in section 5, this can lead to distinct signatures at the LHC which provide strong evidence for the present scenario. Note that even when the mass ordering is not flipped, the amount of mass splitting between the generations differs from the conventional scenarios, which may provide a direct test of this scenario at future colliders. In particular, with our choice of Eqs. (23 – 25), the flavor dependent contribution to the 3-3 entry of Eq. (36), $\epsilon_{E_3}^2|M_{\rm SUSY}|^2$, can be of the same order as the flavor universal contributions. This implies that the $\tilde{\tau}_R$ mass may be significantly split from those of $\tilde{e}_R$ and $\tilde{\mu}_R$, giving a window into the effect of intrinsic flavor violation in the supersymmetry breaking sector. The mass splitting between $\tilde{e}_R$ and $\tilde{\mu}_R$ is of order $\epsilon_{E_2}^2|M_{\rm SUSY}|^2$, which can also be much larger than the conventional scenarios and may be measurable.

4.2 The lightest and next-to-lightest supersymmetric particles

Phenomenology at colliders depends strongly on the species of the lightest superparticle (LSP) and the next-to-lightest superparticle (NLSP). As we have seen, it is natural to expect that (any) one of the right-handed sleptons is the lightest sfermion. For the gauginos, we expect that the bino, $\tilde{B}$, is naturally the lightest because of the renormalization group property of the gaugino masses, $M_A(\mu_R) = (g_A^2(\mu_R)/g_A^2(M_*))M_A(M_*)$, where g_A $(A = 1, 2, 3)$ are the $U(1)_Y$, $SU(2)_L$ and $SU(3)_C$ gauge couplings. This implies that the LSP and NLSP are determined by the competition between the right-handed sleptons, bino, and gravitino, which may also be lighter than the other superparticles.

The mass ordering between the right-handed sleptons, bino, and gravitino depends on the mechanism generating the gaugino masses and the universal contributions to the sfermion masses. Here we consider three representative cases. The first is the simplest case that all the operators of Eq. (1) exist with all η_A and tilde parameters of order unity, except that $\tilde{\zeta}$ is somewhat smaller (to suppress dangerous low energy processes). The second is that the theory does not contain an elementary singlet at M_* $(> M_U)$, so that $\eta_A = \tilde{\kappa}_\mu = \tilde{\eta}_H = \tilde{\eta}_\Phi = \tilde{\zeta} = 0$, and the gaugino and scalar masses are generated by gauge mediation with the messenger scale of order the unification scale, M_U. An interesting aspect of this theory is that the μ term can be generated from the interaction $\mathcal{L} \approx \int d^4\theta\, (H_u H_d + \text{h.c.})$ via supergravity effects, which are comparable to the gaugino and scalar masses: $\mu \approx F_X/M_{\rm Pl} \approx (g_A^2/16\pi^2) F_X/M_U \approx m_{\lambda,\tilde{q},\tilde{l}}$, where $M_{\rm Pl} \approx 10^{18}$ GeV is the reduced Planck scale, and $m_{\lambda,\tilde{q},\tilde{l}}$ represents the gaugino, squark and slepton masses ($\tilde{\kappa}_b$ must be suppressed for M_* smaller than $M_{\rm Pl}$). The third is a class of theories considered in Ref. [9], where $M_* \approx M_U$, and the gaugino and universal scalar masses arise from low energy gauge mediation.

The right-handed slepton mass-squared, m_E^2, and the bino mass, M_1, at the weak scale are given in terms of their values, $m_{E,H}^2$ and $M_{1,H}$, at some high energy scale M_H by

$$m_E^2 \simeq m_{E,H}^2 + \frac{2}{11}\left(1 - \frac{g_1^4}{g_{1,H}^4}\right)|M_{1,H}|^2, \tag{37}$$

$$M_1 \simeq \frac{g_1^2}{g_{1,H}^2}|M_{1,H}|, \tag{38}$$

where g_1 and $g_{1,H}$ are the $U(1)_Y$ gauge couplings at the weak scale and M_H, respectively. In the first case described above, we take $M_H \approx M_*$ and $m_{E,H}^2 \approx 0$ for $\tilde{e}_R$ and $\tilde{\mu}_R$. Neglecting model-dependent effects above M_U, we can set $M_H \approx M_U$, and we find that $m_E^2 < M_1^2$ at the weak scale for these particles, i.e. $\tilde{e}_R$ and $\tilde{\mu}_R$ are lighter than $\tilde{B}$. The mass of $\tilde{\tau}_R$ depends on the sign and size of $m_{E,H}^2 \approx \epsilon_{E_3}^2 |M_{\rm SUSY}|^2$, and may be lighter or heavier than $\tilde{e}_R$, $\tilde{\mu}_R$. In the case where the origin of the gaugino and sfermion masses is gauge mediation, as in the second and third cases above, we should take M_H to be the messenger scale, $M_{\rm mess}$. We find that $\tilde{B}$ is lighter than $\tilde{e}_R$ and $\tilde{\mu}_R$ for $M_{\rm mess} \approx M_U$, but the opposite is possible for lower $M_{\rm mess}$, depending on the number of messenger fields. The mass of $\tilde{\tau}_R$, again, depends on $m_{E,H}^2$.

The gravitino mass is given by

$$m_{3/2} \simeq \frac{F_X}{\sqrt{3} M_{\rm Pl}}, \tag{39}$$

which should be compared to the gaugino and sfermion masses. In the case that all the operators of Eq. (1) exist (except for the superpotential ones) with order one η_A and tilde parameters, the gaugino and sfermion masses are given by

$$m_{\lambda,\tilde{q},\tilde{l}} \approx \frac{F_X}{M_*}. \tag{40}$$

We consider that M_* is at least as large as M_U to preserve successful gauge coupling unification and at most of order $M_{\rm Pl}$ to stay in the field theory regime with weakly coupled gravity. This then leads to

$$\frac{M_U}{M_{\rm Pl}} m_{\lambda,\tilde{q},\tilde{l}} \lesssim m_{3/2} \lesssim m_{\lambda,\tilde{q},\tilde{l}}, \tag{41}$$

where $M_U/M_{\rm Pl} \approx 10^{-2}$. Note that order one coefficients are omitted in Eq. (41), so that the gravitino can be heavier than some (or all) of the superparticles in the supersymmetric standard model sector. Nevertheless, a natural range for the gravitino mass is below the typical superparticle mass by up to two orders of magnitude.

The gravitino mass in the other two cases also falls in the range of Eq. (41). In our second example, the superparticle masses are given approximately by $(g_A^2/16\pi^2)F_X/M_U \approx F_X/M_{\rm Pl}$, leading to $m_{3/2} \approx m_{\lambda,\tilde{q},\tilde{l}}$. The third example has superparticle masses of order $(g_A^2/16\pi^2)F_X/(M_U^2/M_{\rm Pl}) \approx F_X/M_U$, leading to $m_{3/2} \approx (M_U/M_{\rm Pl})\, m_{\lambda,\tilde{q},\tilde{l}} \approx 10^{-2} m_{\lambda,\tilde{q},\tilde{l}}$.

We conclude that the mass ordering between the right-handed sleptons, bino, and gravitino is model dependent. We find, however, that a natural range for the gravitino mass is given by Eq. (41).[4] Thus, it is plausible that the LSP is the gravitino with mass smaller than the typical superparticle mass by a factor of a few to a few hundred.

5 Signatures at the LHC

Signatures of flavorful supersymmetry at the LHC depend strongly on the mass ordering between the right-handed sleptons, $\tilde{l}_R = \tilde{e}_R, \tilde{\mu}_R, \tilde{\tau}_R$, the bino, $\tilde{B}$, and the gravitino, $\tilde{G}$. Based on signatures at the LHC, the six possible orderings can be classified into three cases.

(a) $m_{\tilde{G}} < m_{\tilde{l}_R} < m_{\tilde{B}}$:

One of the right-handed sleptons is the NLSP, which decays into the LSP gravitino. The lifetime is given by

$$\tau_{\tilde{l}_R} \simeq \frac{48\pi\, m_{\tilde{G}}^2 M_{\rm Pl}^2}{m_{\tilde{l}_R}^5} \left(1 - \frac{m_{\tilde{G}}^2}{m_{\tilde{l}_R}^2} \right)^{-4}, \tag{42}$$

which is longer than ~ 100 sec for $m_{3/2}$ in the range of Eq. (41). Signatures are therefore stable charged tracks inside the main detectors, as well as the late decay of the lightest slepton in a stopper which could be placed outside the main detector.

(b) $m_{\tilde{l}_R} < m_{\tilde{B}}, m_{\tilde{G}}$:

One of the right-handed sleptons is the LSP, leaving charged tracks inside the detector. This case, however, has the cosmological problem of charged stable relics.

(c) $m_{\tilde{B}}, m_{\tilde{G}} < m_{\tilde{l}_R}$ or $m_{\tilde{B}} < m_{\tilde{l}_R} < m_{\tilde{G}}$:

A slepton decays into a bino and a lepton inside the detector, so that characteristic signatures are conventional missing energy events. Intrinsic flavor violation in the supersymmetry breaking masses, however, may still be measured by looking at various distributions of kinematic variables.

[4] The gravitino mass can be outside this range. A smaller gravitino mass could arise, for example, if the physics of flavor and supersymmetry breaking occurs below M_U consistently with gauge coupling unification. A larger gravitino mass is also possible if the couplings between X and the matter and Higgs fields are somehow suppressed. For example, if the X field carries a suppression factor ϵ_X then the gravitino mass is enhanced by ϵ_X^{-1}.

5.1 Long-lived slepton

We begin our discussion with case (a) above, in which (one of) the right-handed sleptons is the NLSP decaying into the LSP gravitino. The lifetime of the decay, however, is longer than ~ 100 sec, so that the NLSP is stable for collider analyses.

In the LHC, a stable charged particle interacts with the detector in much the same way as a muon. Therefore its momentum can be measured in both the inner tracker and the muon system. Because of its large mass, however, it will generally move slower than a muon. If its speed is in the range $0.6 \lesssim \beta \lesssim 0.8$, then its mass will be measured to a precision of order a few percent [10, 11]. While not all NLSP's produced have velocity in this range, it is reasonable to expect that a substantial fraction will. Even though they are produced from decays of much heavier strongly interacting superparticles, there will usually be several decay branches, each of which will divide the energy of the event. This reasoning is confirmed by more detailed study [12]. With a measurement of the NLSP mass, we can do full reconstruction of decay chains which will reduce the uncertainty in the NLSP mass to of order 0.1% [13],[5] and can measure more parameters of the low energy theory.

To determine the relationship between supersymmetry breaking and flavor physics, a critical measurement is the flavor content of the leptonic NLSP. Because flavor mixing is generically suppressed by ϵ factors, the NLSP will be mostly of a single flavor. In addition, the NLSP is right handed, so the coupling to the charginos will be small, except possibly the Higgsino to $\tilde{\tau}_R$. The coupling to the neutralino with mostly $\tilde{B}$ content, however, will be large, so the NLSP will usually be produced with a charged lepton of the same flavor. Therefore, we can look for events with only two isolated leptons and two NLSP's, and a (large) number of jets. Most such events will have leptons of the same flavor as the NLSP. The high effective mass of the event should significantly reduce the standard model backgrounds (mostly coming from fakes in events with heavy flavors plus jets or electroweak gauge bosons plus jets). Further background rejection, including supersymmetric and combinatoric, will be possible by reconstructing the masses of the intermediate particles. This could be complicated if the NLSP is mostly $\tilde{\tau}_R$, because we cannot fully reconstruct τ's, but the invariant mass can still be reconstructed and the flavor of the NLSP can be identified.

We now analyze the possibility of probing the flavor properties of the heavier sleptons. In particular, we focus on the situation where $\tilde{l}_1$ and $\tilde{l}_2$ are mostly $\tilde{e}_R$ and $\tilde{\mu}_R$. We consider, for definitiveness, the case where $\tilde{e}_R$ is the NLSP, although the same analysis applies if $\tilde{\mu}_R$ is the NLSP. As shown in section 4.1, it is likely that $\tilde{\mu}_R$ is only $\epsilon_{E_2}^2 M_{\rm SUSY} \approx$ a few GeV heavier than the NLSP, so the decay of $\tilde{B}$ will produce $\tilde{e}_R$ about half the time and $\tilde{\mu}_R$ just as often. When a

[5] This precision can be achieved if the systematic uncertainties are ~ 100 MeV and the squarks and gluinos are not too heavy.

$\tilde{\mu}_R$ is produced, it will decay into a $\tilde{e}_R$ and two leptons. The leptons produced in this decay will be soft in the $\tilde{\mu}_R$ rest frame, having energy of order only a few GeV, but in general the system will be boosted. This possibly poses a problem: for a fast $\tilde{\mu}_R$ the leptons will be harder but highly collimated with the NLSP track, $\theta \lesssim 0.1$, while for a slow $\tilde{\mu}_R$ the opening angle will be larger but the leptons will have low p_T. One expects that in some intermediate kinematic regime a reconstruction may be feasible, but a detailed study of this issue is beyond the scope of this paper. If this reconstruction turns out to be possible, one can look for events where one $\tilde{\mu}_R$ is produced, decaying to $\tilde{e}_R$. These events will have two hard leptons, two soft leptons, and two NLSP's. This event topology should make it possible to measure the mass difference between the two lightest sleptons, as well as to provide information on the flavor content of the (N)NLSP by looking at the flavor of the four leptons.

In the region of parameter space where $\alpha_\beta/\alpha_l \ll 1$, the flavor non-universal contribution will be very small and the sleptons will be degenerate. In this co-NLSP region the decay of one slepton into another is suppressed because the decay into charged sleptons is not kinematically allowed and the right-handed sleptons do not couple to neutrinos. In this region, all three right-handed sleptons are long lived, and extracting information on intrinsic flavor violation in the supersymmetry breaking parameters requires careful analyses. Since this is a small region of parameter space, we do not focus on it here.

The above analysis was for case (a) where the slepton was the NLSP and the gravitino the LSP, but it also applies to case (b) where the slepton is the LSP. While this scenario is disfavored cosmologically by limits on charged relics, the situation could be ameliorated by, for example, slight R parity violation in the lepton sector, along with a solution to the dark matter problem independent of supersymmetry.[6]

5.2 Late decay of the long-lived slepton

In order to determine the lifetime of the NLSP slepton, we would like to observe its decays. The NLSP's produced with $\beta \lesssim 0.4$ will be stopped within the detector. One can then detect NLSP decays by looking for particles which do not point back to the interaction area. Another possibility is that the NLSP will be stopped in the rock just outside the detector, and then some of the decay products will re-enter the detector. Unfortunately, very few NLSP's will be produced with low enough velocity, and one has to deal with cosmic neutrino background. A further possibility is to use the tracker to determine where in the surrounding rock the NLSP is stopped. If the lifetime is longer than a few weeks, we could then extract pieces of the rock where the NLSP is stopped

[6] An alternative possibility is that the slepton decays into the axino, the fermionic superpartner of the axion, with a lifetime (much) longer than the collider time scale. The phenomenology of this scenario is similar to the case with a gravitino LSP.

and study the decay in a more quiet environment [11]. This will also not have very many events, but it will allow very precise measurement of the mass and decay properties of the NLSP.

In addition, a large stopper detector can be built outside the main detector to trap the NLSP's and measure their decay products [12, 14]. Conventional scenarios only consider a $\tilde{\tau}_R$ NLSP, but in flavorful supersymmetry the NLSP could be one of the other sleptons, which would decay to a monochromatic electron or muon. This would make it very easy to (i) measure the mass of the gravitino given the mass of the NLSP measured in the collider, (ii) measure the lifetime of the NLSP by counting the number of decays as a function of time, and (iii) test supergravity relations such as Eq. (42) [15], and make sure that the gravitino is indeed the LSP. The stopper detector proposed in Ref. [12] did not include a magnetic field, so it could not measure the energy of muons, only of electrons and taus. Perhaps this design can be modified to include a magnetic field to measure the momentum of the muons.

A stopper detector can very precisely measure the flavor content of the NLSP. If a sufficient number of NLSP's are trapped and there is flavor mixing, then a few of the NLSP's will decay to a lepton with different flavor. This occurs in a very clean environment so there should be almost no fakes once the accelerator is turned off. A stopper detector can very efficiently separate electrons from muons, and it can use the monochromatic spectrum of the first two generation slepton decays to distinguish τ decay products. Mixing angles as small as about 10^{-2} can be measured [16]. The main background comes from cosmic neutrino events, but those should all have much lower energy than the NLSP decays.

5.3 Neutralino (N)LSP

Finally we consider case (c) where the neutralino is lighter than the sleptons. With this spectrum, all sleptons will decay promptly, and measuring flavor violation is more difficult. Because the neutralino will escape the detector without interacting, every event has missing energy, making event reconstruction much more difficult. For direct slepton production one is forced to use kinematic variables such as M_{T2} [17], but they require very high statistics. The low Drell-Yan production cross section quickly prevents this strategy as the slepton mass is increased. The $\tilde{\tau}_R$ is expected to be very split from the other two generations, but looking for τ's means even more particles contributing to missing energy.

We are then driven to study lepton flavor violation in cascade decays by looking at multiple edges in flavor-tagged dilepton invariant mass distributions, along the lines of Refs. [18, 19]. This method requires sizable flavor violating couplings and will probe both those and any modifications to the slepton spectrum. However, in order to perform this study with right-handed sleptons, they must be produced by the second lightest neutralino, χ_2^0, which will be mostly wino, so it has a small branching fraction to right-handed sleptons, typically of order 1%. On the other hand, the χ_2^0 and the left-handed sleptons are expected to be in the same mass range. So if the spectrum is such that the left-handed sleptons are lighter than χ_2^0, then the branching ratio of χ_2^0 to $\tilde{l}_L$ will be large. One can then repeat the analysis of section 4.1 in the left-handed slepton sector and use the methods of Ref. [18] to probe flavor violation.

6 Conclusions

Weak scale supersymmetry provides elegant solutions to many of the problems of the standard model, but it also generically gives rise to excessive flavor and CP violation. While most existing models assume that the mechanism of mediating supersymmetry breaking to the supersymmetric standard model sector is flavor universal, we have shown that this is not necessary to satisfy all low energy flavor and CP constraints. We have considered a scenario, flavorful supersymmetry, in which the mechanism that suppresses the Yukawa couplings also suppresses flavor changing interactions in the supersymmetry breaking parameters. We find that a broad region of parameter space is allowed, as long as the superpotential couplings generating scalar trilinear interactions are suppressed or the superparticles have masses of at least a TeV.

The flavorful supersymmetry framework can lead to mass splitting among different generations of squarks and sleptons much larger than in conventional scenarios. This has interesting implications on collider physics. In particular, the mass ordering and splitting among the three right-handed sleptons, which are expected to be the lightest sfermion species, can easily differ significantly from the conventional scenarios. Signatures at colliders depend strongly on the species of the LSP and the NLSP, and we have argued that these are likely to be one of the right-handed sleptons, the bino, or the gravitino. The gravitino mass is typically in the range $10^{-2} m_{\lambda,\tilde{q},\tilde{l}} \lesssim m_{3/2} \lesssim m_{\lambda,\tilde{q},\tilde{l}}$, where $m_{\lambda,\tilde{q},\tilde{l}}$ is a characteristic superparticle mass, so that it is plausible to expect that the LSP is the gravitino with mass smaller than the typical superparticle mass by a factor of a few to a few hundred.

In the case that the lightest right-handed slepton is lighter than the bino, we expect to see the dramatic presence of long-lived charged particles at the LHC. This allows us to do full reconstruction of decay chains and reduce the uncertainty in the NLSP mass determination. Moreover, it is natural to expect that the lightest right-handed slepton is, in fact, the NLSP decaying into the gravitino with the lifetime longer than ~ 100 sec. Because of intrinsic flavor violation in the superparticle masses in flavorful supersymmetry, the NLSP is not necessarily $\tilde{\tau}_R$ but can be $\tilde{e}_R$ or $\tilde{\mu}_R$, leading to the spectacular signature of monochromatic electrons or muons in a stopper detector. This provides a simple method to measure the gravitino mass, as well as the lifetime and the flavor content of the NLSP, and will be a smoking gun signature for flavorful supersymmetry. In general, flavorful supersymmetry predicts flavor violation in both production and decay of sleptons. Precision measurements of these processes will also test the flavor content of the sleptons. While these precision measurements are difficult at the LHC, they can be done by using certain event topologies, regardless of the LSP species. Further study of flavor violation can also be done at a future linear collider.

The origin of the flavor structure is a deep mystery both in the context of the standard model and beyond the standard model. The framework of flavorful supersymmetry allows the LHC to probe this physics which may lie at a scale close to the Planck scale. Precise study of processes such as the ones discussed in this paper will be crucial to uncover the mechanism that leads to the distinct flavor pattern we see in nature.

Note Added:
While completing this paper, we received Ref. [20] which discusses similar ideas.

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References

[1] A. H. Chamseddine, R. Arnowitt and P. Nath, Phys. Rev. Lett. **49**, 970 (1982); R. Barbieri, S. Ferrara and C. A. Savoy, Phys. Lett. B **119**, 343 (1982); L. J. Hall, J. Lykken and S. Weinberg, Phys. Rev. D **27**, 2359 (1983).

[2] M. Dine and W. Fischler, Phys. Lett. B **110**, 227 (1982); Nucl. Phys. B **204**, 346 (1982); L. Alvarez-Gaumé, M. Claudson and M. B. Wise, Nucl. Phys. B **207**, 96 (1982); S. Dimopoulos and S. Raby, Nucl. Phys. B **219**, 479 (1983).

[3] M. Dine, A. E. Nelson and Y. Shirman, Phys. Rev. D **51**, 1362 (1995) [arXiv:hep-ph/9408384]; M. Dine, A. E. Nelson, Y. Nir and Y. Shirman, Phys. Rev. D **53**, 2658 (1996) [arXiv:hep-ph/9507378].

[4] R. Kitano and Y. Nomura, arXiv:hep-ph/0606134.

[5] A. Pomarol and D. Tommasini, Nucl. Phys. B **466**, 3 (1996) [arXiv:hep-ph/9507462]; R. Barbieri, L. J. Hall and A. Romanino, Phys. Lett. B **401**, 47 (1997) [arXiv:hep-ph/9702315]; G. G. Ross, L. Velasco-Sevilla and O. Vives, Nucl. Phys. B **692**, 50 (2004) [arXiv:hep-ph/0401064].

[6] L. J. Hall, V. A. Kostelecky and S. Raby, Nucl. Phys. B **267**, 415 (1986).

[7] See, e.g., L. J. Hall, H. Murayama and N. Weiner, Phys. Rev. Lett. **84**, 2572 (2000) [arXiv:hep-ph/9911341]; T. Yanagida and J. Sato, Nucl. Phys. Proc. Suppl. **77**, 293 (1999) [arXiv:hep-ph/9809307]; P. Ramond, Nucl. Phys. Proc. Suppl. **77**, 3 (1999) [arXiv:hep-ph/9809401].

[8] F. Gabbiani, E. Gabrielli, A. Masiero and L. Silvestrini, Nucl. Phys. B **477**, 321 (1996) [arXiv:hep-ph/9604387]; A. Masiero, S. K. Vempati and O. Vives, arXiv:0711.2903 [hep-ph], and references therein.

[9] M. Ibe and R. Kitano, JHEP **0708**, 016 (2007) [arXiv:0705.3686 [hep-ph]]; Y. Nomura and M. Papucci, arXiv:0709.4060 [hep-ph].

[10] ATLAS Detector and Physics Performance Technical Design Report, `http://atlas.web.cern.ch/Atlas/GROUPS/PHYSICS/TDR/access.html`

[11] A. De Roeck, J. R. Ellis, F. Gianotti, F. Moortgat, K. A. Olive and L. Pape, Eur. Phys. J. C **49**, 1041 (2007) [arXiv:hep-ph/0508198].

[12] K. Hamaguchi, M. M. Nojiri and A. de Roeck, JHEP **0703**, 046 (2007) [arXiv:hep-ph/0612060].

[13] I. Hinchliffe and F. E. Paige, Phys. Rev. D **60**, 095002 (1999) [arXiv:hep-ph/9812233]; S. Ambrosanio, B. Mele, S. Petrarca, G. Polesello and A. Rimoldi, JHEP **0101**, 014 (2001) [arXiv:hep-ph/0010081].

[14] K. Hamaguchi, Y. Kuno, T. Nakaya and M. M. Nojiri, Phys. Rev. D **70**, 115007 (2004) [arXiv:hep-ph/0409248]; J. L. Feng and B. T. Smith, Phys. Rev. D **71**, 015004 (2005) [Erratum-ibid. D **71**, 0109904 (2005)] [arXiv:hep-ph/0409278].

[15] W. Buchmüller, K. Hamaguchi, M. Ratz and T. Yanagida, Phys. Lett. B **588**, 90 (2004) [arXiv:hep-ph/0402179].

[16] K. Hamaguchi and A. Ibarra, JHEP **0502**, 028 (2005) [arXiv:hep-ph/0412229].

[17] C. G. Lester and D. J. Summers, Phys. Lett. B **463**, 99 (1999) [arXiv:hep-ph/9906349].

[18] A. Bartl, K. Hidaka, K. Hohenwarter-Sodek, T. Kernreiter, W. Majerotto and W. Porod, Eur. Phys. J. C **46**, 783 (2006) [arXiv:hep-ph/0510074].

[19] G. L. Bayatian *et al.* [CMS Collaboration], J. Phys. G **34**, 995 (2007).

[20] J. L. Feng, C. G. Lester, Y. Nir and Y. Shadmi, arXiv:0712.0674 [hep-ph].