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Naturally flavorful supersymmetry at the LHC

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The suppression of flavor and *CP* violation in supersymmetric theories may be due to the mechanism responsible for the structure of the Yukawa couplings. We study model independently the compatibility between low-energy flavor and *CP* constraints and observability of superparticles at the LHC, assuming a generic correlation between the Yukawa couplings and the supersymmetry breaking parameters. We find that the superpotential operators that generate scalar trilinear interactions are generically problematic. We discuss several ways in which this tension is naturally avoided. In particular, we focus on several frameworks in which the dangerous operators are naturally absent. These frameworks can be combined with many theories of flavor, including those with (flat or warped) extra dimensions, strong dynamics, or flavor symmetries. We show that the resulting theories can avoid all the low-energy constraints while keeping the superparticles light. The intergenerational mass splittings among the sfermions can reflect the structure of the underlying flavor theory, and can be large enough to be measurable at the LHC. Detailed observations of the superparticle spectrum may thus provide new handles on the origin of the flavor structure of the standard model.

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I. INTRODUCTION

The origin of the flavor structure is one of the deepest mysteries of the standard model (SM). In the absence of the Yukawa couplings (and neutrino masses), the standard model respects the following flavor symmetry[1]:

$$\begin{aligned} G_{\rm flavor}^{\rm SM} &\equiv U(3)^5 \\ &= U(3)_Q \times U(3)_U \times U(3)_D \times U(3)_L \times U(3)_E. \end{aligned} \tag{1}$$

Is this symmetry a mere artifact of the low-energy Lagrangian, or is it (or its subgroup) physically realized at high energies and spontaneously broken to produce the Yukawa couplings? If the latter, what is the fundamental flavor group $G_{\rm flavor}$ ($\subset G_{\rm flavor}^{\rm SM}$), and how is it broken? In the standard model, these questions can be explored only through the observed pattern of masses and mixings of the quarks and leptons, making it difficult to arrive at conclusive answers.

Theories beyond the standard model may provide additional clues to address the puzzle of flavor, since the new physics sector may contain new information on flavor. In supersymmetric theories, for example, supersymmetry breaking parameters for squarks and sleptons may carry such information. It is, however, not obvious how much new information one can expect. In supersymmetric theories, generic weak scale values for the supersymmetry (SUSY) breaking parameters

$$(m_\Phi^2)_{ij} \sim m_{\rm SUSY}^2, \qquad (a_f)_{ij} \sim m_{\rm SUSY}, \tag{2}$$

lead to flavor changing neutral currents far in excess of current experimental bounds. Here, $(m_\Phi^2)_{ij}$ ($\Phi = Q, U, D, L, E$) and $(a_f)_{ij}$ ($f = u, d, e$) represent scalar squared masses and trilinear interactions, respectively, $i, j = 1, 2, 3$ are generation indices, and $m_{\rm SUSY}$ is a parameter of order the weak scale. A common solution to this problem is to assume flavor universality

$$\begin{aligned} (m_\Phi^2)_{ij} &\propto \delta_{ij} \text{ or } \ll m_{\rm SUSY}^2, \\ (a_f)_{ij} &\propto (y_f)_{ij} \text{ or } \ll m_{\rm SUSY}, \end{aligned} \tag{3}$$

at a scale M where the supersymmetry breaking parameters are generated [1–5]. Here, $(y_f)_{ij}$ are the Yukawa matrices. This assumption, however, greatly reduces the flavor information encoded in low-energy supersymmetry. A nontrivial flavor structure can still be found in the low-energy squark and slepton masses due to renormalization group evolution below M. This structure, however, does not carry any information on flavor beyond Yukawa couplings Φ have, although it does allow us to explore some of these couplings that cannot be probed in the standard model [6]. In addition, the size of the relevant flavor nonuniversality is typically so small that most of the interesting parameters can be probed only indirectly through low-energy flavor and *CP* violating processes.

In this paper we study the question: is there a natural and generic framework for supersymmetry which is sufficiently "flavorful," i.e., which allows us to obtain more detailed information on flavor through measurements of superparticle masses and interactions at the LHC? This is not trivial because such a framework must satisfy stringent

[1]$G_{\rm flavor}^{\rm SM}$ contains hypercharge, baryon number, and lepton number. Out of the five $U(1)$ factors, only hypercharge and baryon minus lepton number are anomaly free with respect to the standard model gauge group.

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constraints from flavor and *CP* violation while the deviation from universality must be sufficiently large to be experimentally observable. In particular, in order for the flavor information to be extracted at the LHC, superparticles must be light enough to be produced at the LHC, making it more difficult to satisfy the bounds from the low-energy flavor and *CP* violating processes.

While it is not too difficult to consider an *ad hoc* deviation from universality that is measurable at the LHC and not excluded by the low-energy data, an important question is if there is a theoretically well-motivated setup which naturally produces measurable effects that are consistent with the low-energy experiments and encode information on the origin of flavor. In a previous paper with Papucci, we studied a simple setup in which flavor changing interactions in the supersymmetry breaking parameters are scaled by factors associated with the Yukawa couplings [7]. We showed that such a setup can avoid all the low-energy constraints, while giving interesting flavor signatures at the LHC. This clearly illustrates that there is an interesting, natural stage between Eqs. (2) and (3). In fact, there have been many models proposed to address the problem of flavor changing neutral currents, in which flavor violation in the supersymmetry breaking parameters is somehow related to the Yukawa couplings [8–21].[2] While many of these models require rather special structures or setups to avoid all the current experimental bounds, the analysis of Ref. [7] suggests that the minimal structure needed to obtain a consistent framework for flavor signatures at the LHC may, in fact, be much simpler.

In this paper we study the tension between LHC observability and constraints from low-energy flavor and *CP* violation in generic supersymmetric theories in which the structure of the supersymmetry breaking parameters is correlated with that of the Yukawa couplings. An interesting general point emphasized in Ref. [7] (see also [23]) is that among the operators giving supersymmetry breaking parameters, a class of operators in the superpotential

$$W \sim XQ_iU_jH_u, XQ_iD_jH_d, XL_iE_jH_d, \qquad (4)$$

generally leads to a strong tension. Here, X represents a chiral superfield whose F-term vacuum expectation value (VEV) is responsible for supersymmetry breaking. We first elucidate this point, and define what we call the superpotential flavor problem. We emphasize that the problem is general and does not depend on any particular theory of flavor. In fact, this problem is part of a more general problem associated with left-right propagation of sfermions, which is also discussed in detail.

We then discuss how the problems described above can be solved. We present several possibilities that can avoid the stringent constraints from left-right sfermion propagation without suppressing the operators of Eq. (4). We also present simple frameworks in which the operators of Eq. (4) are naturally absent. These frameworks can be combined with many theories of flavor, including theories with (flat or warped) extra dimensions, strong dynamics, or flavor symmetries. We perform detailed studies of the constraints from low-energy flavor and *CP* violation within these frameworks, and find that they can naturally avoid all the constraints while preserving the observability of superparticles at the LHC. We also find that the intergenerational mass splittings among sfermions can show a variety of patterns depending on the details of the underlying flavor theory, allowing us to gain additional handles on the origin of flavor at the LHC.

The organization of the paper is as follows. In Sec. II we discuss constraints from low-energy flavor and *CP* violation arising from left-right propagation of sfermions. We emphasize the model-independent nature of the problem associated with the operators of Eq. (4), but also discuss additional stringent model-dependent constraints. In Sec. III we discuss possibilities to avoid these constraints. In particular, we present simple frameworks in which the operators of Eq. (4) are naturally absent. In Sec. IV we study constraints from flavor and *CP* violation, including ones arising from left-left and right-right sfermion propagation, in these frameworks. We also analyze the size of the intergenerational mass splittings among sfermions, and find that they can differ significantly from the flavor universal case. Prospects for observing these features at the LHC are discussed in Sec. V. Finally, discussion and conclusions are given in Sec. VI.

II. THE SUPERSYMMETRIC LEFT-RIGHT FLAVOR PROBLEM

The flavor problem in supersymmetric models is typically phrased such that generic supersymmetry breaking parameters, Eq. (2), lead to excessive flavor and *CP* violation at low energies. This, however, neglects the possibility that the physics responsible for the observed Yukawa couplings also controls the pattern of the supersymmetry breaking parameters. Here we argue that there is a *generic* tension between weak scale supersymmetry and low-energy flavor and *CP* violation even if we take this possibility into account. Throughout the discussion, we assume that *CP* violating effects not associated with flavor, e.g., those arising from a nontrivial phase in the Higgs sector, are adequately suppressed. We also assume that the strong *CP* problem is solved.

A. Flavor (non)universality in the operator language

We begin our discussion by listing all the operators in the supersymmetric standard model (SSM). In the gauge sector, the relevant operators are

[2] The possibility of having large flavor violation in the supersymmetry breaking parameters which is not related to the Yukawa couplings was recently discussed in theories with an extended R symmetry [22].

$$\mathcal{O}_{g_A}\colon \int d^2\theta \frac{1}{4g_A^2} \mathcal{W}^{A\alpha}\mathcal{W}^A_\alpha + \text{H.c.},$$
$$\mathcal{O}_{\lambda_A}\colon \int d^2\theta \eta_A \frac{X}{M} \mathcal{W}^{A\alpha}\mathcal{W}^A_\alpha + \text{H.c.}, \tag{5}$$

where $A = 1, 2, 3$ represents the standard model gauge group, $U(1)_Y$, $SU(2)_L$, and $SU(3)_C$, and $\mathcal{O}_{g_A}$ and $\mathcal{O}_{\lambda_A}$ give the gauge kinetic terms and the gaugino masses, respectively.[3] Here, X is the supersymmetry breaking superfield, $\langle X\rangle = \theta^2 F_X$, and M characterizes a scale at which supersymmetry breaking effects are mediated to the SSM sector. Since F_X/M sets the scale for superparticle masses, we consider $F_X/M \approx O$ (TeV). In the minimal supersymmetric standard model, the Higgs sector operators are given by

$$\mathcal{O}_{Z_H}\colon \int d^4\theta Z_H H^\dagger H, \tag{6}$$

$$\mathcal{O}_{\kappa_H}\colon \int d^4\theta \kappa_H \frac{X^\dagger X}{M^2} H^\dagger H,$$
$$\mathcal{O}_{\eta_H}\colon \int d^4\theta \eta_H \frac{X}{M} H^\dagger H + \text{H.c.}, \tag{7}$$

$$\mathcal{O}_{\mu}\colon \int d^4\theta \eta_\mu \frac{X^\dagger}{M} H_u H_d + \text{H.c.},$$
$$\mathcal{O}_{b}\colon \int d^4\theta \kappa_b \frac{X^\dagger X}{M^2} H_u H_d + \text{H.c.}, \tag{8}$$

$$\mathcal{O}_{\text{SUGRA}}\colon \int d^4\theta \lambda_H H_u H_d + \text{H.c.}, \tag{9}$$

where $H = H_u, H_d$. The operators $\mathcal{O}_{Z_H}$ give the kinetic terms, while the rest provide the supersymmetric mass, μ, and the holomorphic and nonholomorphic supersymmetry breaking squared masses, b, $m^2_{H_u}$, and $m^2_{H_d}$, for the Higgs doublets.[4] The last operator is relevant only in the context of supergravity. In nonminimal models, e.g., in models with extra gauge groups and/or singlet fields, the set of operators in Eqs. (5)–(9) is extended.

The operators described above (or their extensions in nonminimal models) do not introduce flavor violation. Flavor violation may arise when we introduce matter fields. With matter fields, we can write operators

$$\mathcal{O}_{Z_\Phi}\colon \int d^4\theta (Z_\Phi)_{ij} \Phi_i^\dagger \Phi_j, \tag{10}$$

$$\mathcal{O}_{\kappa_\Phi}\colon \int d^4\theta (\kappa_\Phi)_{ij} \frac{X^\dagger X}{M^2} \Phi_i^\dagger \Phi_j,$$
$$\mathcal{O}_{\eta_\Phi}\colon \int d^4\theta (\eta_\Phi)_{ij} \frac{X}{M} \Phi_i^\dagger \Phi_j + \text{H.c.}, \tag{11}$$

$$\mathcal{O}_{\lambda_f}\colon \int d^2\theta (\lambda_f)_{ij} \Phi_{Li}\Phi_{Rj} H + \text{H.c.},$$
$$\mathcal{O}_{\zeta_f}\colon \int d^2\theta (\zeta_f)_{ij} \frac{X}{M} \Phi_{Li}\Phi_{Rj} H + \text{H.c.}, \tag{12}$$

where $\Phi = \Phi_L, \Phi_R$ with $\Phi_L = Q, L$ and $\Phi_R = U, D, E$ represents matter fields, $i, j = 1, 2, 3$ are generation indices, and $f = u, d, e$ corresponds to $\{\Phi_L, \Phi_R, H\} = \{Q, U, H_u\}, \{Q, D, H_d\}, \{L, E, H_d\}$. Here, Z_Φ and κ_Φ are 3×3 Hermitian matrices, while η_Φ, λ_f and ζ_f are general complex 3×3 matrices. The operators $\mathcal{O}_{Z_\Phi}$ give the kinetic terms, $\mathcal{O}_{\lambda_f}$ the Yukawa couplings, and $\mathcal{O}_{\kappa_\Phi}$, $\mathcal{O}_{\eta_\Phi}$ and $\mathcal{O}_{\zeta_f}$ the soft supersymmetry breaking parameters.

Flavor universality is the assumption that

$$(\kappa_\Phi)_{ij} \propto (Z_\Phi)_{ij}, \qquad (\eta_\Phi)_{ij} \propto (Z_\Phi)_{ij},$$
$$(\zeta_f)_{ij} \propto (\lambda_f)_{ij}, \tag{13}$$

for all $\Phi = Q, U, D, L, E$ and $f = u, d, e$, which leads to supersymmetry breaking parameters of the form

$$(m^2_\Phi)_{ij} \propto \delta_{ij}, \qquad (a_f)_{ij} \propto (y_f)_{ij}, \tag{14}$$

where $(m^2_\Phi)_{ij}$, $(a_f)_{ij}$, and $(y_f)_{ij}$ represent the scalar squared masses, scalar trilinear interactions, and the Yukawa couplings in the basis where the fields are canonically normalized. In fact, the three conditions of Eq. (13) could each be replaced by

$$|(\kappa_\Phi)_{ij}| \ll |\eta_A|^2, \qquad |(\eta_\Phi)_{ij}| \ll |\eta_A|,$$
$$|(\zeta_f)_{ij}| \ll |(\lambda_f)_{ij}\eta_A|, \tag{15}$$

since then the low-energy supersymmetry breaking parameters, generated by SSM renormalization group evolution, take approximately the form of Eq. (14).

Deviations from Eqs. (13) and (15) generically lead to flavor and CP violating effects. If the supersymmetry breaking parameters take the flavor universal form at some scale below the scale of flavor physics M_F, then small deviations from universality are caused only by SSM renormalization group evolution below that scale, which do not provide much insight into the origin of flavor at the LHC. On the other hand, if $M_F \lesssim M$ (or if the mediation mechanism of supersymmetry breaking somehow carries information on physics responsible for the Yukawa structure), then we expect that the supersymmetry breaking parameters have an intrinsic flavor nonuniversality, which contains information on the physics of flavor at M_F. Of course, this deviation from universality cannot be arbitrary. In order to satisfy all the constraints while keeping superparticles within the reach of the LHC, the devia-

[3]We define an operator $\mathcal{O}$ to be the entire term that appears in the Lagrangian, including the coefficient.

[4]Here we have neglected the tree-level superpotential operator $\int d^2\theta \mu_0 H_u H_d$ + H.c. In order to have weak scale values for μ and b, the coefficient of this operator must be suppressed: $\mu_0 \lesssim O$ (TeV).

tion must somehow be correlated with the Yukawa structure. This is, however, precisely what we expect if the supersymmetry breaking parameters feel the physics responsible for the flavor structure.

B. Generic scalar trilinear interactions

We consider the case where $M_F \lesssim M$ and flavor nonuniversality in the operators of Eqs. (10)–(12) at M_F is controlled by the physics responsible for the structure of the Yukawa couplings. This provides a possibility of avoiding the low-energy constraints without imposing flavor universality, allowing us to probe the origin of flavor through the superparticle spectrum. In general, correlations between the structure of the Yukawa couplings and that of the nonuniversality in the operators of Eqs. (10)–(12) are model dependent. However, one class of operators, $\mathcal{O}_{\zeta_f}$ in Eq. (12), is expected to have a structure similar to the Yukawa couplings. This is relatively model-independent because the matter and Higgs fields appear in $\mathcal{O}_{\zeta_f}$ in precisely the same way as in $\mathcal{O}_{\lambda_f}$, which produces the Yukawa couplings.

Suppose that the Yukawa couplings

$$(y_f)_{ij} = (\lambda_f)_{kl}(Z_{\Phi_L}^{-1/2})_{ki}(Z_{\Phi_R}^{-1/2})_{lj}Z_H^{-1/2}, \qquad (16)$$

have a hierarchical structure as a result of some flavor physics, for example, physics associated with spontaneous breaking of a flavor symmetry or wave function profiles of matter fields in extra dimensions. We then expect that the scalar trilinear interactions generated by $\mathcal{O}_{\zeta_f}$,

$$(a_f)_{ij} = -(\zeta_f)_{kl}(Z_{\Phi_L}^{-1/2})_{ki}(Z_{\Phi_R}^{-1/2})_{lj}Z_H^{-1/2}\frac{F_X}{M}, \qquad (17)$$

also have a similar structure. Specifically, we can consider that the Yukawa couplings take the form

$$(y_u)_{ij} \approx \mathcal{E}^u_{ij}\tilde{y}, \qquad (y_d)_{ij} \approx \mathcal{E}^d_{ij}\tilde{y}, \qquad (y_e)_{ij} \approx \mathcal{E}^e_{ij}\tilde{y}, \quad (18)$$

and that the observed structure for the quark and lepton masses and mixings is generated by the "suppression factors" $\mathcal{E}^f_{ij}$. Here, $\tilde{y}$ represents the "natural" size of the couplings before taking into account the origin of the flavor structure; for example, we expect $\tilde{y} \approx O(1)$ if the relevant physics is weakly coupled, but it could be as large as of $O(4\pi)$ if strongly coupled. The scalar trilinear interactions are then expected to take the form

$$(a_u)_{ij} \approx \mathcal{E}^u_{ij}\frac{\tilde{\zeta}F_X}{M}, \qquad (a_d)_{ij} \approx \mathcal{E}^d_{ij}\frac{\tilde{\zeta}F_X}{M}, \qquad (a_e)_{ij} \approx \mathcal{E}^e_{ij}\frac{\tilde{\zeta}F_X}{M}, \qquad (19)$$

where $\tilde{\zeta}$ again represents the natural size of the coefficients. Note that $O(1)$ coefficients are omitted in the expressions of Eqs. (18) and (19); for example, we expect that $(a_f)_{ij}$ is in general not proportional to $(y_f)_{ij}$ because of an arbitrary $O(1)$ coefficient in each element.

The structure of Eqs. (18) and (19) is expected to appear in most theories of flavor. A special case is when $\mathcal{E}^f_{ij}$ factorize as $\mathcal{E}^u_{ij} = \epsilon_{Q_i}\epsilon_{U_j}$, $\mathcal{E}^d_{ij} = \epsilon_{Q_i}\epsilon_{D_j}$, and $\mathcal{E}^e_{ij} = \epsilon_{L_i}\epsilon_{E_j}$, so that each matter field carries its own suppression factor. This arises in many models of flavor, for example, in classes of models with flavor symmetries or strong dynamics. The important point is the similarity between the forms of $(y_f)_{ij}$ and $(a_f)_{ij}$. This comes from the fact that matter and Higgs fields appear identically in the two classes of operators in Eq. (12).

The gaugino masses, which arise from $\mathcal{O}_{\lambda_A}$ as $M_A = -2\eta_A g_A^2 F_X/M$, set the scale for the superparticle masses. Assuming that the mediation mechanism produces unsuppressed $\mathcal{O}_{\zeta_f}$, i.e., $\tilde{\zeta} \approx \eta_A$, we then obtain $(a_f)_{ij} \approx \mathcal{E}^f_{ij}M_A/g_A^2$ at the scale M_F.[5] Taking into account Eq. (18) and SSM renormalization group evolution below M, we can write the flavor nonuniversal part of the low-energy scalar trilinear interactions as

$$(a_u)_{ij} \approx (y_u)_{ij}\frac{a_{\rm C}}{\tilde{y}}, \qquad (a_d)_{ij} \approx (y_d)_{ij}\frac{a_{\rm C}}{\tilde{y}}, \qquad (a_e)_{ij} \approx (y_e)_{ij}\frac{a_{\rm N}}{\tilde{y}}, \qquad (20)$$

where $a_{\rm C}, a_{\rm N} \approx O(M_A)$ are characteristic mass scales for these interactions associated with colored and noncolored sfermions, and we expect $a_{\rm C} \gtrsim a_{\rm N}$ due to the structure of the SSM renormalization group equations. We note again that $O(1)$ coefficients are omitted for each element of Eq. (20), so that $(a_f)_{ij}$ is not proportional to $(y_f)_{ij}$ as a matrix.

C. The superpotential flavor problem

The flavor nonuniversality at the level of Eq. (20) can be problematic. Flavor and CP violation can in general be quantified by mass insertion parameters, which are obtained by dividing the off diagonal entry of the sfermion mass-squared matrix by the average diagonal entry in the super-Cabibbo-Kobayashi-Maskawa (CKM) basis [24,25]. The mass insertion parameters obtained from Eq. (20) are

$$(\delta^u_{LR})_{ij} \approx \frac{a_{\rm C}}{\tilde{y}m_{\rm C}^2}(M_u)_{ij}, \qquad (\delta^d_{LR})_{ij} \approx \frac{a_{\rm C}}{\tilde{y}m_{\rm C}^2}(M_d)_{ij}, \qquad (\delta^e_{LR})_{ij} \approx \frac{a_{\rm N}}{\tilde{y}m_{\rm N}^2}(M_e)_{ij}, \qquad (21)$$

for $i \neq j$, where $m_{\rm C}$ and $m_{\rm N}$ are characteristic masses for colored and noncolored superparticles, and we expect $m_{\rm C} \gtrsim m_{\rm N}$. Here, $(M_u)_{ij} = (y_u)_{ij}\langle H_u\rangle$, $(M_d)_{ij} = (y_d)_{ij}\langle H_d\rangle$, and $(M_e)_{ij} = (y_e)_{ij}\langle H_d\rangle$ are the quark and lepton mass matrices in the original (not super-CKM) basis. The diago-

[5] Note that in our notation, the grand unified relations for the gaugino masses correspond to $\eta_1 = \eta_2 = \eta_3$.

nal elements, $(\delta^f_{LR})_{ii}$, receive additional terms coming from the sfermion left-right masses proportional to $|\mu|$ and the flavor universal contribution to $(a_f)_{ij}$ generated by renormalization group evolution below M. However, assuming there is no intrinsic CP violation associated with supersymmetry breaking, these terms do not contribute to the imaginary parts of $(\delta^f_{LR})_{ii}$, which are relevant in the discussion below. We thus find

$$\mathrm{Im}\,(\delta^u_{LR})_{ii} \approx \frac{a_{\rm C}\sin\varphi_u}{\tilde{y}m^2_{\rm C}}(M_u)_{ii},$$
$$\mathrm{Im}(\delta^d_{LR})_{ii} \approx \frac{a_{\rm C}\sin\varphi_d}{\tilde{y}m^2_{\rm C}}(M_d)_{ii}, \qquad (22)$$
$$\mathrm{Im}(\delta^e_{LR})_{ii} \approx \frac{a_{\rm N}\sin\varphi_e}{\tilde{y}m^2_{\rm N}}(M_e)_{ii},$$

where φ_f are the phases of the contributions to $(\delta^f_{LR})_{ii}$ from the flavor nonuniversal part of $(a_f)_{ij}$. Note that $\varphi_f \approx O(1)$ is expected even if supersymmetry breaking does not introduce new CP violating phases because these complex phases arise generically from the Yukawa couplings when going into the super-CKM basis.[6]

The theoretical estimate of Eqs. (21) and (22) can be compared with experimental constraints from low-energy observables. The bound on the $\mu \to e\gamma$ process [28] gives

$$\frac{1}{\sqrt{2}}\sqrt{|(\delta^e_{LR})_{12}|^2 + |(\delta^e_{LR})_{21}|^2} \lesssim 4\times 10^{-6}\left(\frac{m_{\rm N}}{200\ {\rm GeV}}\right). \qquad (23)$$

The $\mu \to e$ conversion and $\mu \to eee$ processes also give comparable bounds. The limits on the electric dipole moments (EDMs) of the electron [29], neutron [30], and mercury atom [31] lead to

$$|\mathrm{Im}(\delta^e_{LR})_{11}| \lesssim 2\times 10^{-7}\left(\frac{m_{\rm N}}{200\ {\rm GeV}}\right), \qquad (24)$$

$$|\mathrm{Im}(\delta^u_{LR})_{11}| \lesssim 2\times 10^{-6}\left(\frac{m_{\rm C}}{600\ {\rm GeV}}\right),$$
$$|\mathrm{Im}(\delta^d_{LR})_{11}| \lesssim 1\times 10^{-6}\left(\frac{m_{\rm C}}{600\ {\rm GeV}}\right), \qquad (25)$$

$$|\mathrm{Im}(\delta^u_{LR})_{11}| \lesssim 4\times 10^{-7}\left(\frac{m_{\rm C}}{600\ {\rm GeV}}\right),$$
$$|\mathrm{Im}(\delta^d_{LR})_{11}| \lesssim 4\times 10^{-7}\left(\frac{m_{\rm C}}{600\ {\rm GeV}}\right), \qquad (26)$$

respectively. Here, the bounds of Eqs. (23)–(26) are obtained conservatively by scanning the ratios of the superparticle masses in a reasonable range (see, e.g., [25,32]).[7] The bounds from neutron and mercury EDMs are subject to large theoretical uncertainties [33], and we have used conservative estimates. The constraints from $(\epsilon'/\epsilon)_K$ and $b \to s\gamma$ also lead to bounds on $|\mathrm{Im}(\delta^d_{LR})_{12,21}|$ and $|(\delta^d_{LR})_{23,32}|$, but they are not as strong as the bounds above when the left-right mass insertion parameters scale naively with the quark masses.

We can obtain the bounds on the superparticle masses using Eqs. (21) and (22), with the approximation $a_{\rm C} \approx m_{\rm C}$ and $a_{\rm N} \approx m_{\rm N}$, which is sufficient for the level of analysis here. Taking $(M_u)_{11} \simeq m_u \simeq 2$ MeV and $(M_d)_{11} \simeq m_d \simeq 4$ MeV, the neutron EDM bound of Eq. (25) leads to the following bound on m_C:

$$m_{\rm C} \gtrsim \max\left\{800\ {\rm GeV}\left(\frac{\sin\varphi_u}{\tilde{y}}\right)^{1/2}, 1.5\ {\rm TeV}\left(\frac{\sin\varphi_d}{\tilde{y}}\right)^{1/2}\right\}, \qquad (27)$$

whereas the mercury EDM bound, Eq. (26), gives

$$m_{\rm C} \gtrsim \max\left\{1.3\ {\rm TeV}\left(\frac{\sin\varphi_u}{\tilde{y}}\right)^{1/2}, 1.9\ {\rm TeV}\left(\frac{\sin\varphi_d}{\tilde{y}}\right)^{1/2}\right\}. \qquad (28)$$

The bound on m_N depends on the assumption on the charged lepton mass matrix. If we conservatively take $(M_e)_{11} \simeq m_e \simeq 0.5$ MeV and $(M_e)_{12} \simeq (M_e)_{21} \simeq (m_e m_\mu)^{1/2} \simeq 7$ MeV, we obtain

$$m_{\rm N} \gtrsim \max\left\{600\ {\rm GeV}\frac{1}{\tilde{y}^{1/2}}, 700\ {\rm GeV}\left(\frac{\sin\varphi_e}{\tilde{y}}\right)^{1/2}\right\}. \qquad (29)$$

On the other hand, if the large neutrino mixing angle θ_{12} receives a significant contribution from the charged lepton Yukawa matrix, we expect $(M_e)_{12} \simeq m_\mu \tan\theta_{12} \simeq 70$ MeV, giving a much stronger bound

$$m_{\rm N} \gtrsim \max\left\{1.9\ {\rm TeV}\frac{1}{\tilde{y}^{1/2}}, 700\ {\rm GeV}\left(\frac{\sin\varphi_e}{\tilde{y}}\right)^{1/2}\right\}. \qquad (30)$$

In fact, this latter bound is expected to apply in the case where $\mathcal{E}^e_{ij}$ factorizes: $\mathcal{E}^e_{ij} = \epsilon_{L_i}\epsilon_{E_j}$, since then the large 1–2 neutrino mixing angle generically implies $\epsilon_{L_1} \approx \epsilon_{L_2}$, leading to a large 1–2 element of the charged lepton mass matrix of $O(m_\mu)$.

In addition to the uncertainties already described, the bounds on $m_{\rm C,N}$ derived above are subject to uncertainties coming from $O(1)$ coefficients in front of Eq. (20). Only the square root of these coefficients, however, appears in the bounds. For example, if we take the magnitude of these

[6]This may be avoided in certain models, e.g., models with Hermitian Yukawa and scalar trilinear interaction matrices [26] and those with spontaneous CP violation [27].

[7]Here we consider the ranges $m^2_{\tilde{g}}/m^2_{\tilde{q}} \lesssim 2$ and $m^2_{\chi}/m^2_{\tilde{l}} \lesssim 3$, where $m_{\tilde{g}}$, $m_{\tilde{q}}$, m_χ, and $m_{\tilde{l}}$ are the gluino, squark, weak gaugino, and slepton average masses. These ranges are motivated by renormalization group considerations with M_F well above the TeV scale, e.g., $M_F \gtrsim 10^{10}$ GeV. The case of smaller M_F will be discussed in Sec. III A.

coefficients to be between 0.5 and 2, the bounds receive unknown coefficients of $O(0.7\text{–}1.4)$, which do not significantly affect the results. These bounds also scale with the square root of the natural size of the scalar trilinear interactions at M, $(\tilde{\zeta}/\eta_A)^{1/2}$, which we have set unity. In addition, the bounds scale with the experimental limits on the $\mu \to e\gamma$ branching ratio, $\mathrm{Br}(\mu \to e\gamma)$, and the electron, neutron and mercury EDMs, d_e, d_n, and d_{Hg}, as

$$\left(\frac{\mathrm{Br}(\mu \to e\gamma)}{1.2 \times 10^{-11}}\right)^{-1/4}, \tag{31}$$

$$\left(\frac{d_e}{1.6 \times 10^{-27}\ e\,\mathrm{cm}}\right)^{-1/2}, \quad \left(\frac{d_n}{2.9 \times 10^{-26}\ e\,\mathrm{cm}}\right)^{-1/2}, \quad \left(\frac{d_{\mathrm{Hg}}}{2.1 \times 10^{-28}\ e\,\mathrm{cm}}\right)^{-1/2}. \tag{32}$$

Therefore, if future experiments such as ones in Refs. [34–41] improve the upper bounds on these (and other) quantities, the lower bounds on $m_{\mathrm{C,N}}$ increase accordingly.

The bounds of Eqs. (27)–(30) place lower limits on the superparticle masses, yielding a tension with the observability of supersymmetry at the LHC. In fact, the conservative bound of Eq. (29) already gives strong constraints on the superparticle spectrum for $\tilde{y} \approx 1$. In particular, in the case that colored and noncolored superparticles do not have a strong mass hierarchy at $M \gg$ TeV, we expect that $m_{\mathrm{C}} \approx (2\text{–}4)m_{\mathrm{N}}$ at low energies. This pushes up the masses of colored superparticles beyond 1 TeV, and, in many cases, beyond the reach of the LHC. The constraints are even stronger if the large neutrino mixing angle θ_{12} receives a sizable contribution from the charged lepton Yukawa matrix, as in Eq. (30). We call this generic tension between low-energy flavor and CP violation and the observability of supersymmetry at the LHC *the superpotential flavor problem*, since it is caused by the superpotential operators $\mathcal{O}_{\zeta_f}$. An important point, again, is that the problem is relatively model independent because the flavor structure of $\mathcal{O}_{\zeta_f}$ is expected to be correlated with that of $\mathcal{O}_{\lambda_f}$ in wide classes of flavor theories.

D. More general problem with left-right sfermion propagation

The superpotential flavor problem provides a strong, model-independent tension between weak scale supersymmetry and low-energy flavor and CP violating observables. This is, however, only one aspect of a more general problem associated with left-right propagation of the sfermions in flavor and CP violating amplitudes.

Suppose that the superpotential flavor problem is somehow solved, i.e., the operators $\mathcal{O}_{\zeta_f}$ are strongly suppressed. There will still be the contributions to lepton flavor violation and EDMs associated with left-right propagation of sfermions. First of all, there are flavor nonuniversal scalar trilinear interactions generated by $\mathcal{O}_{Z_\Phi,\eta_\Phi}$, yielding $(\delta^f_{LR})_{ij}$ $(i \neq j)$ and $\mathrm{Im}(\delta^f_{LR})_{ii}$. These contributions must be sufficiently suppressed. Moreover, even if they are small, lepton flavor violation and EDMs are induced by diagrams that use multiple mass insertion parameters $(\delta^f_{LR})_{ij}$, $(\delta^f_{LL})_{ij}$, and $(\delta^f_{RR})_{ij}$ (see Fig. 1) [42], instead of a single insertion of $(\delta^f_{LR})_{ij}$. Since the diagrams depend on parameters $(\delta^f_{LL})_{ij}$ and $(\delta^f_{RR})_{ij}$, whose correlations with the Yukawa couplings are model dependent, the tension caused by these diagrams is not as model independent as the superpotential flavor problem. Nevertheless, this provides strong constraints on supersymmetric models in which the structure of the supersymmetry breaking parameters is correlated with that of the Yukawa couplings.

One might naively think that because of the use of multiple mass insertion parameters, the diagrams of Fig. 1 are much smaller than those using a single $(\delta^f_{LR})_{ij}$. This is, however, not always the case for the following reasons:

(i) The left-right mass insertions used can be flavor universal, $\mathrm{Re}(\delta^f_{LR})_{ii}$, since the necessary flavor/CP violation can come from insertions of $(\delta^f_{LL})_{ij}$ and/or $(\delta^f_{RR})_{ij}$. This may enhance the contributions from multiple mass insertion diagrams relative to single insertion ones, especially for $f = d$, e, since the flavor universal part of $(\delta^{d,e}_{LR})_{ij}$ is enhanced by $\tan\beta \equiv \langle H_u \rangle / \langle H_d \rangle$. [For $f = e$, it is also enhanced by μ/m_{N}, which is typically of $O(m_{\mathrm{C}}/m_{\mathrm{N}})$.]

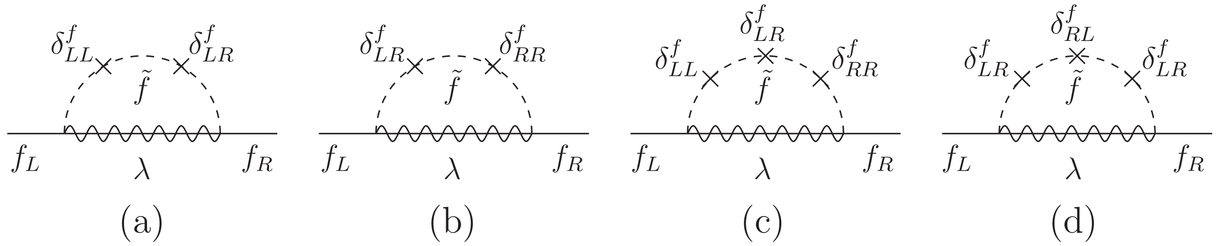

FIG. 1. Multiple mass insertion diagrams that lead to dangerous flavor and CP violating contributions. Here, $f_{L,R}$, $\tilde{f}$, and λ represent fermions, scalars, and gauginos, respectively.

(ii) The sfermions propagating between two mass insertions can be from a heavier generation. For diagrams with triple mass insertions, for example, the states propagating between mass insertions can be third generation states, minimizing extra suppressions arising from use of more mass insertion parameters.

In fact, these two ingredients can make the contributions from the diagrams of Fig. 1 comparable or even larger than those from the diagrams with a single insertion of $(\delta^f_{LR})_{ij}$ of Eqs. (21) and (22).

To illustrate this point, suppose that $\mathcal{E}^f_{ij}$ in Eq. (18) factorize, $\mathcal{E}^u_{ij} = \epsilon_{Q_i}\epsilon_{U_j}$, $\mathcal{E}^d_{ij} = \epsilon_{Q_i}\epsilon_{D_j}$, and $\mathcal{E}^e_{ij} = \epsilon_{L_i}\epsilon_{E_j}$, giving the Yukawa couplings $(y_u)_{ij} \approx \tilde{y}\epsilon_{Q_i}\epsilon_{U_j}$, $(y_d)_{ij} \approx \tilde{y}\epsilon_{Q_i}\epsilon_{D_j}$, and $(y_e)_{ij} \approx \tilde{y}\epsilon_{L_i}\epsilon_{E_j}$. Suppose also that the flavor nonuniversal part of the sfermion squared masses scale naively with the ϵ factors

$$(m^2_\Phi)_{ij} \approx \epsilon_{\Phi_i}\epsilon_{\Phi_j} m^2_{\rm S}, \tag{33}$$

where $m_{\rm S} \approx \eta_A F_X/M$ is the scale of supersymmetry breaking parameters at M, which we assume to be the same for colored and noncolored superparticles. The mass insertion parameters generated by Eq. (33) are then

$$\begin{gathered}(\delta^u_{LL})_{ij} \approx (\delta^d_{LL})_{ij} \approx \epsilon_{Q_i}\epsilon_{Q_j}\frac{m^2_{\rm S}}{m^2_{\rm C}},\\ (\delta^u_{RR})_{ij} \approx \epsilon_{U_i}\epsilon_{U_j}\frac{m^2_{\rm S}}{m^2_{\rm C}}, \qquad (\delta^d_{RR})_{ij} \approx \epsilon_{D_i}\epsilon_{D_j}\frac{m^2_{\rm S}}{m^2_{\rm C}},\end{gathered} \tag{34}$$

$$\begin{gathered}(\delta^e_{LL})_{ij} \approx (\delta^\nu_{LL})_{ij} \approx \epsilon_{L_i}\epsilon_{L_j}\frac{m^2_{\rm S}}{m^2_{\rm N}},\\ (\delta^e_{RR})_{ij} \approx \epsilon_{E_i}\epsilon_{E_j}\frac{m^2_{\rm S}}{m^2_{\rm N}},\end{gathered} \tag{35}$$

where $i \neq j$, and we have included the mass insertion parameters for the sneutrinos. On the other hand, the dominant contribution to the flavor universal part of the left-right mass insertion parameters are given by

$$\begin{gathered}{\rm Re}(\delta^u_{LR})_{ii} \approx \frac{1}{m_{\rm C}}(M_u)_{ii}, \qquad {\rm Re}(\delta^d_{LR})_{ii} \approx \frac{\mu\tan\beta}{m^2_{\rm C}}(M_d)_{ii},\\ {\rm Re}(\delta^e_{LR})_{ii} \approx \frac{\mu\tan\beta}{m^2_{\rm N}}(M_e)_{ii},\end{gathered} \tag{36}$$

where we have taken $\mu \approx m_{\rm C} \gtrsim m_{\rm N}$ and $\tan\beta \gtrsim 1$, and assumed that flavor universal scalar trilinear interactions $(a_{u,d})_{ii} \approx (y_{u,d})_{ii} m_{\rm C}$ and $(a_e)_{ii} \approx (y_e)_{ii} m_{\rm N}$ are generated by renormalization group evolution. [The expression of Eq. (36) also applies to the case where $\mathcal{E}^f_{ij}$ do not factorize.]

Consider, for example, the diagram of Fig. 1(c) with $f = u$. This leads to the contribution to the up quark EDM that scales with

$$(\delta^u_{LL})_{13}(\delta^u_{LR})_{33}(\delta^u_{RR})_{31} \approx \frac{(M_u)_{11}}{\tilde{y} m_{\rm C}}\left(\frac{m_{\rm S}}{m_{\rm C}}\right)^4 \frac{(y_u)^2_{33}}{\tilde{y}}, \tag{37}$$

which can be comparable to the dangerous contribution that scales with $(\delta^u_{LR})_{11} \approx (M_u)_{11}/\tilde{y} m_{\rm C}$ of Eq. (21) with $a_{\rm C} \approx m_{\rm C}$. The diagram of Fig. 1(a) with $f = e$ gives a contribution to the $\mu \to e\gamma$ process that scales with

$$(\delta^e_{LL})_{12}(\delta^e_{LR})_{22} \approx \frac{(M_e)_{12}}{\tilde{y} m_{\rm N}}\frac{m^2_{\rm S}\mu}{m^3_{\rm N}}\tilde{y}\epsilon^2_{L_2}\tan\beta, \tag{38}$$

which can also be dangerous because it could be comparable to the contribution from $(\delta^e_{LR})_{12} \approx (M_e)_{12}/\tilde{y} m_{\rm N}$ of Eq. (21) with $a_{\rm N} \approx m_{\rm N}$, especially for large $\tan\beta$. These examples show that the multiple mass insertion diagrams may lead to flavor and CP violation at a dangerous level even in the absence of flavor and CP violating $(\delta^f_{LR})_{ij}$.

In practice, the constraints from multiple mass insertion diagrams can be taken into account by considering the effective left-right mass insertion parameters

$$\begin{aligned}(\delta^f_{LR,{\rm eff}})_{ij} \equiv \max\{&c_d(\delta^f_{LL})_{ik}(\delta^f_{LR})_{kj},\\ &c_d(\delta^f_{LR})_{ik}(\delta^f_{RR})_{kj}, c_t(\delta^f_{LL})_{ik}(\delta^f_{LR})_{kk}(\delta^f_{RR})_{kj},\\ &c_t(\delta^f_{LR})_{ik}(\delta^f_{RL})_{kk}(\delta^f_{LR})_{kj}\},\end{aligned} \tag{39}$$

and requiring that $(\delta^f_{LR,{\rm eff}})_{ij}$ satisfy the bounds of Eqs. (23)–(26) with $(\delta^f_{LR})_{ij}$ replaced by $(\delta^f_{LR,{\rm eff}})_{ij}$. Here, $c_d \simeq (0.5\text{–}0.8)$ and $c_t \simeq (0.3\text{–}0.6)$ are numerical coefficients arising from the difference of momentum integral functions with various numbers of insertions. Once $(\delta^f_{LL})_{ij}$ and $(\delta^f_{RR})_{ij}$ are given, these constraints can be checked.

III. APPROACHES TO THE PROBLEM

In order to have a framework for weak scale supersymmetry in which the LHC can provide additional insight into the origin of the observed flavor structure, the supersymmetric left-right flavor problem must somehow be addressed. The bounds associated with left-left and right-right sfermion propagation must also be avoided, although they are, in general, less stringent. To address the issue of whether there are theories that naturally satisfy all these constraints, we start by identifying classes of theories that do not have the superpotential flavor problem, a robust part of the supersymmetric left-right flavor problem. While it is not automatic that these theories will be safe from low-energy constraints or even solve the supersymmetric left-right flavor problem, they provide frameworks with which to build more detailed theories that can avoid all low-energy constraints. The remaining constraints will be discussed in the next section.

A. General considerations

There are essentially two different directions to address the superpotential flavor problem. One is to assume that the operators $\mathcal{O}_{\zeta_f}$ exist with their natural size, but the bounds are somehow avoided. Barring accidental cancellations in the amplitudes for low-energy flavor and CP violating effects, this includes the following possibilities:

(i) The bounds are given by Eq. (29), and the superparticles are not too much heavier. If the superparticle masses satisfy $m_C \approx (2\text{–}4)m_N$, the viable parameter region is somewhat squeezed. The constraints are slightly relaxed if we allow the masses of colored and noncolored superparticles to be of similar size, $m_C \sim m_N$. Avoiding the bound of Eq. (30), however, is still not easy.

(ii) The intrinsic size of the Yukawa couplings is large, $\tilde{y} \gg O(1)$. In this case the bounds on m_C and m_N are not significant, especially in the case where the large neutrino mixing angle θ_{12} arises only from the neutrino mass matrix, Eqs. (27)–(29). If the bound is given by Eq. (30), the lower bound on m_N can be relaxed to about 500 GeV by taking the largest possible value of $\tilde{y} \approx 4\pi$. This constraint, however, is still significant.

(iii) The gauginos are significantly heavier than the sfermions. In this case the bounds of Eqs. (27) and (28) and Eqs. (29) and (30) are relaxed approximately by a factor of $(m_{\tilde{g}}/m_{\tilde{q}})/\sqrt{2}$ and $(m_\chi/m_{\tilde{l}})/\sqrt{3}$, respectively, with $m_{C,N}$ now interpreted as the masses of the sfermions. This situation can occur if M_F is close to the TeV scale, and the masses of the sfermions at M_F are suppressed by the dynamics generating the Yukawa hierarchy, as in flavor models of Ref. [16].[8] Note that we only need sfermions to be accessible at the LHC to probe the origin of the flavor structure.

These possibilities are certainly viable, especially given uncertainties in our estimates. The tension between flavor constraints and LHC observability, however, still exists. If one of these possibilities is realized, and the superparticles are within the LHC reach, we expect that $\mu \to e$ processes and/or atomic and nuclear EDMs will be discovered in the near future, for example, in the experiments of Refs. [34–41], which expect to improve present bounds by several orders of magnitude.

The other direction to address the superpotential flavor problem is to consider that the operators $\mathcal{O}_{\zeta_f}$ are somehow suppressed. This includes the following possibilities:

(iv) The coefficients of the operators $\mathcal{O}_{\zeta_f}$ (or at least those of the 1–2 and 2–1 elements of $\mathcal{O}_{\zeta_e}$) are accidentally suppressed. The required amount of suppression is not strong if we adopt Eq. (29). However, if we instead use Eq. (30), we need to have $\tilde{\zeta}/\eta_A \lesssim 0.07\tilde{y}(m_N/500\ \text{GeV})^2$, which provides a strong bound on $\tilde{\zeta}/\eta_A$ for $\tilde{y} \approx 1$.

(v) The scalar trilinear interactions are exactly proportional to the Yukawa matrices, $(a_f)_{ij} \propto (y_f)_{ij}$, leading to vanishing flavor and CP violating mass insertion parameters. This may be achieved, for example, if $(a_f)_{ij}$ and $(y_f)_{ij}$ arise from a single operator through the lowest and highest components VEVs of X, $\langle X\rangle = X_0 + \theta^2 F_X$, with $\arg(F_X/X_0) \approx \arg(M_A)$. The large top quark mass, however, requires that X_0 is close to the cutoff scale M_*, and the problem may be regenerated by higher order terms in X_0/M_*.

(vi) The operators $\mathcal{O}_{\zeta_f}$ are suppressed by some mechanism, $\tilde{\zeta} \ll \eta_A$. This mechanism may or may not operate in the regime where effective field theory is valid.

Note that (iv), (v), and (vi) above can also be combined with (i), (ii), and (iii) described before. For example, we can consider a setup where $\mathcal{O}_{\zeta_f}$ are suppressed by some mechanism, (vi), *and* the natural size of the Yukawa couplings, $\tilde{y}$, is large, (ii).

In the rest of this paper, we focus on the last possibility, (vi), and see how much the situation will be improved. As we discussed, we still need to address the more general left-right flavor problem and the constraints from left-left and right-right sfermion propagation, which we defer to the next section. Here we present simple classes of theories in which $\mathcal{O}_{\zeta_f}$ are naturally suppressed. In fact, this provides a platform for the analysis of the more model-dependent part of the flavor problem.

There are several possible ways that the operators $\mathcal{O}_{\zeta_f}$ can be suppressed. In fact, they may simply be absent at a scale where the SSM arises as an effective field theory, as a result of the dynamics of some more fundamental theory. This is not unreasonable because $\mathcal{O}_{\zeta_f}$ are the only superpotential operators associated with supersymmetry breaking, so if supersymmetry breaking is mediated to the SSM sector by loop processes then these operators may be absent. In the remainder of this section we discuss three simple classes of theories for suppressing $\mathcal{O}_{\zeta_f}$. We classify them according to the pattern of the supersymmetry breaking parameters obtained at M_F. The simple suppression of $\mathcal{O}_{\zeta_f}$ described above can effectively be classified into the first class, because it leads to the same pattern of the supersymmetry breaking masses at M_F. There are clearly many models within each class, and we explicitly discuss some of them.

B. Framework I—Higgsphobic supersymmetry breaking

We here present the first class of theories in which $\mathcal{O}_{\zeta_f}$ are naturally suppressed. A unique feature of the operators $\mathcal{O}_{\zeta_f}$ is that among the operators relevant to flavor violation, Eqs. (10)–(12), these are the only operators that contain both Higgs fields, H, and the supersymmetry breaking field, X. Therefore, if the theory does not allow direct

[8] M_F needs to be low to prevent the scalar masses from becoming comparable to the gaugino masses through SSM renormalization group evolution.

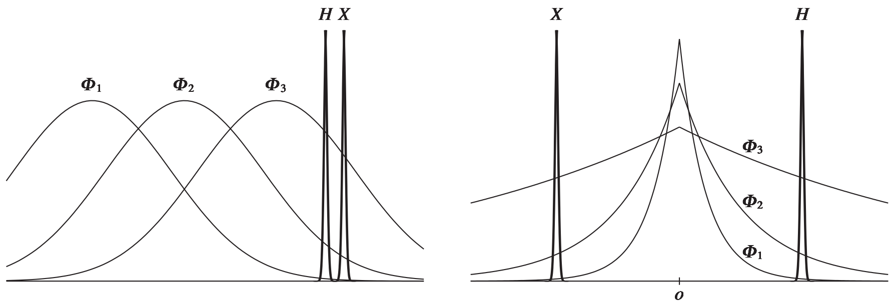

FIG. 2. A schematic depiction of possible configurations of the matter, Higgs and supersymmetry breaking fields. The Higgs and supersymmetry breaking fields are localized to separate but nearby points (left). In the case where all the matter wave functions are spherically symmetric and centered around the same point o, the Higgs and supersymmetry breaking fields can be localized (approximately) the same distance away from o (right). The gauge fields are assumed to propagate in the bulk.

coupling between H and X, a framework which we call *Higgsphobic supersymmetry breaking*, then $\mathcal{O}_{\zeta_f}$ are forbidden. The other operators $\mathcal{O}_{Z_\Phi,\kappa_\Phi,\eta_\Phi,\lambda_f}$ can exist as long as couplings between Φ and X and between Φ and H are allowed. A simple way of realizing this is to assume that H and X are localized in different points in extra dimensions, while Φ have broad wave functions overlapping with both H and X.[9]

Let us now focus on the extra dimensional way of suppressing $\mathcal{O}_{\zeta_f}$ described above. In theories with extra dimensions, wave function overlaps between Φ and H control the size of the 4D Yukawa couplings [18,43,44]. This motivates a configuration where heavier generation matter fields have larger wave function overlaps with the Higgs fields. For example, if the widths of all the matter wave functions are the same, then heavier generation fields have wave functions peaked closer to the Higgs fields. On the other hand, if all the matter wave functions are peaked at the same point (separated from where the Higgs fields reside), then heavier generation fields have wider wave functions. A schematic depiction of these possibilities is shown in Fig. 2.

The location of X cannot be arbitrary. If the X field is localized to a generic point in the region where matter wave functions are significant, the generated soft supersymmetry breaking parameters have a random structure, e.g., Eq. (2), leading to large flavor and CP violation. One way of avoiding this is to localize X far away from wave function peaks for all the matter fields, in which case flavor violation in the low-energy supersymmetry breaking masses arises only from loop effects across the bulk [18]. Another way is to localize X at a point close to where the Higgs fields are localized [20]. In this case the tree-level structure of the operators $\mathcal{O}_{\kappa_\Phi,\eta_\Phi}$ is correlated with that of the Yukawa couplings $\mathcal{O}_{\lambda_f}$, since they are both controlled by the wave function values of the matter fields in the region where the X and the Higgs fields reside. In the case where all the matter wave functions are spherically symmetric and peaked at the same point o, a similar correlation can be obtained by localizing X to a point (approximately) the same distance away from o as the Higgs fields. In all these cases, flavor violation also arises from loop effects across the bulk.

Some of the field configurations discussed above are depicted schematically in Fig. 2. Note that while the figures describe only three matter wave functions for illustrative purposes, all Q_i, U_i, D_i, L_i, and E_i fields can have distinct wave functions. The geometry of the extra dimensions can also be more general: the number of extra dimensions is arbitrary, and the spacetime need not be flat. While the scale of the extra dimensions, i.e., the scale of Kaluza-Klein resonances, is, in principle, arbitrary, it is simplest to consider it to be of order the unification scale, $1/R \approx M_{\rm unif} \approx 10^{16}$ GeV, to preserve the success of supersymmetric gauge coupling unification in the most straightforward manner. In the case where H and X are localized in the infrared region of warped spacetime, the scale can be lower, $10~{\rm TeV} \lesssim 1/R \lesssim M_{\rm unif}$ (in which case gauge coupling unification can occur through modified gauge coupling running above $1/R$, as in Ref. [45]). It is also possible to consider the framework in the context of grand unification in higher dimensions [46,47]. In the theories considered here, a natural scale for flavor physics is of order $1/R$, while a natural scale for supersymmetry breaking mediation is of order the (local) cutoff scale M_*, which we take to be somewhat above $1/R$.

One consequence of Higgsphobic supersymmetry breaking is that the Higgs sector operators $\mathcal{O}_{\kappa_H,\eta_H,\mu,b}$ in Eqs. (7) and (8) are forbidden in the minimal setup. There are, however, several ways to generate the desired μ and b parameters, which are discussed in Appendix A.

C. Framework II—Remote flavor-supersymmetry breaking

We now consider the second framework. An essential ingredient of this framework is a "separation" between

[9]Precisely speaking, in extra dimensional theories it is sufficient to assume that the Yukawa couplings are allowed only in places separated from the X field. The Higgs fields can be delocalized in that case.

supersymmetry breaking and flavor symmetry breaking. Consider a flavor symmetry $G_{\rm flavor}$ that prohibits the Yukawa operators $\mathcal{O}_{\lambda_f}$ in the unbroken limit. The SSM Yukawa couplings are then generated through breaking of $G_{\rm flavor}$, which we assume to be the origin of the observed Yukawa structure [48]. An important point is that among the operators relevant to flavor violation, $\mathcal{O}_{Z_\Phi,\kappa_\Phi,\eta_\Phi,\lambda_f,\zeta_f}$ in Eqs. (10)–(12), $\mathcal{O}_{\zeta_f}$ are the only operators that *require* both $G_{\rm flavor}$ breaking and supersymmetry breaking (see Table I). Therefore, if we assume that the theory possesses $G_{\rm flavor}$, and that $G_{\rm flavor}$ and supersymmetry are broken in different sectors of the theory that do not directly communicate with each other, a framework which we call *remote flavor-supersymmetry breaking*, then the operators $\mathcal{O}_{\zeta_f}$ are absent.

In the present framework, the Yukawa couplings are generated through breaking of $G_{\rm flavor}$. Assuming that the breaking is caused by the VEV of a chiral superfield ϕ, the relevant operators are written schematically as

$$\mathcal{L} \approx \int d^2\theta \sum_{i,j} \left(\frac{\phi}{M_*}\right)^{(n_f)_{ij}} \Phi_{Li}\Phi_{Rj} H + \text{H.c.}, \qquad (40)$$

where M_* is the (effective) cutoff scale and $(n_f)_{ij}$ are integers. In general, these operators could generate dangerous scalar trilinear interactions through supersymmetry breaking. Here we assume that they do not generate significant scalar trilinear interactions. The conditions under which this is indeed the case are discussed in Appendix B for the general case that $G_{\rm flavor}$ is broken by the VEVs of several fields ϕ_m ($m = 1, 2, \cdots$).

A symmetry group $G_{\rm flavor}$ needs to be chosen to avoid all the low-energy flavor and CP violating constraints. Suppose that $G_{\rm flavor}$ were a simple $U(1)$ symmetry, whose breaking controls the size of the Yukawa couplings. In this case the Cabibbo angle, θ_C, would be reproduced by the difference of the $U(1)$ charges of Q_1 and Q_2, q_{Q_1} and q_{Q_2}, as $\sin\theta_C \approx \epsilon^{q_{Q_1}-q_{Q_2}}$, where ϵ is the dimensionless $U(1)$ breaking parameter normalized to have a charge of -1, and $q_{Q_1} > q_{Q_2}$. This, however, would lead to too large flavor violation in $\mathcal{O}_{\kappa_Q}$, giving $(\delta^d_{LL})_{12} \approx (\delta^d_{LL})_{21} \approx O(\epsilon^{q_{Q_1}-q_{Q_2}}) \approx O(\sin\theta_C)$, which needs to be smaller than of order $10^{-2}(m_C/600\ \text{GeV})$ to avoid the bound from ϵ_K. (Here we have assumed $\eta_A^2 \approx \kappa_\Phi$ and a generic Yukawa structure.) Similar conflicts between the Yukawa structure and flavor violating processes also arise in other places. One possibility of avoiding these bounds is to consider more elaborate Abelian charge assignments, for example, under $G_{\rm flavor} = U(1) \times U(1)$ (see, e.g., [12]). Another, perhaps simpler, approach is to consider a non-Abelian $G_{\rm flavor}$ symmetry under which (at least) the first two generations of quarks and leptons having the same standard model gauge quantum numbers are in a single $G_{\rm flavor}$ multiplet. This makes the relevant coefficients of $\mathcal{O}_{\kappa_\Phi}$ proportional to the unit matrix, significantly reducing the problem.

TABLE I. Required symmetry breaking to write down operators in Eqs. (10)–(12). The operators $\mathcal{O}_{\lambda_f,\zeta_f}$ require $G_{\rm flavor}$ breaking, while $\mathcal{O}_{\kappa_\Phi,\eta_\Phi,\zeta_f}$ require supersymmetry breaking.

Operators	$\mathcal{O}_{Z_\Phi}$	$\mathcal{O}_{\kappa_\Phi}$	$\mathcal{O}_{\eta_\Phi}$	$\mathcal{O}_{\lambda_f}$	$\mathcal{O}_{\zeta_f}$
$G_{\rm flavor}$ breaking				√	√
Supersymmetry breaking		√	√		√

Note that flavor violation in this framework can come mainly from the operators $\mathcal{O}_{Z_\Phi}$ (in the $G_{\rm flavor}$ symmetric field basis). Consider that $G_{\rm flavor}$ is a sufficiently large subgroup of $SU(3)^5 = SU(3)_Q \times SU(3)_U \times SU(3)_D \times SU(3)_L \times SU(3)_E$ so that all the three generations are treated equally under $G_{\rm flavor}$. In this case the sector breaking supersymmetry generates the operators $\mathcal{O}_{\kappa_\Phi}$ and $\mathcal{O}_{\eta_\Phi}$, but they are completely flavor universal

$$(\kappa_\Phi)_{ij} \propto \delta_{ij}, \qquad (\eta_\Phi)_{ij} \propto \delta_{ij}. \qquad (41)$$

Flavor violation, however, still arises at M_F because the operators $\mathcal{O}_{Z_\Phi}$ receive flavor nonuniversal contributions from the sector breaking $G_{\rm flavor}$

$$(Z_\Phi)_{ij} \not\propto \delta_{ij}. \qquad (42)$$

This situation is depicted schematically in Fig. 3. The scalar squared masses and trilinear interactions in the basis where the fields are canonically normalized are, therefore, flavor nonuniversal at M_F.

A simple way of realizing the present framework is to consider higher dimensional theories in which the bulk flavor symmetry $G_{\rm flavor}$ and supersymmetry are broken on separate branes. Note that $G_{\rm flavor}$ can be broken on multiple branes, which could help address the issue of vacuum alignment, depending on $G_{\rm flavor}$ and its breaking pattern. If the relevant extra dimension is warped [49], $G_{\rm flavor}$ and supersymmetry can be broken at the ultraviolet and infrared branes, respectively. Through the AdS/CFT correspondence, these theories have a 4D interpretation that supersymmetry is dynamically broken by strong gauge dynamics that have an approximate flavor symmetry.

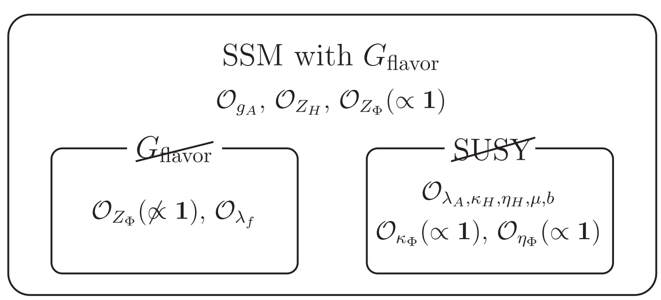

FIG. 3. The schematic picture of a remote flavor-supersymmetry breaking theory with $G_{\rm flavor}$ being a sufficiently large subgroup of $SU(3)^5$. Here, we have depicted only operators relevant for the analysis.

Exchanging the locations of supersymmetry and flavor breaking is also an interesting possibility, which corresponds to 4D theories in which nontrivial flavor structures arise dynamically at low energies.

D. Framework III—Charged supersymmetry breaking

The final framework we consider is one in which X carries a nontrivial charge of some symmetry, so that the operators $\mathcal{O}_{\zeta_f}$ are forbidden. (We assume that the Yukawa couplings, $\mathcal{O}_{\lambda_f}$, are allowed.) This symmetry should have anomalies with respect to the standard model gauge group so that the gaugino mass operators $\mathcal{O}_{\lambda_A}$ can be written. (For an example of this class of models, see [20].) An immediate consequence of this framework, which we call *charged supersymmetry breaking*, is that the operators $\mathcal{O}_{\eta_H}$ and $\mathcal{O}_{\eta_\Phi}$ are also forbidden. This class of theories, therefore, has vanishing scalar trilinear interactions at the scale M.[10]

The operator $\mathcal{O}_\mu$, which leads to the μ parameter, may or may not be forbidden, depending on the charge assignments of X and H. An interesting point is that once we choose the charge assignment such that $\mathcal{O}_\mu$ is allowed, $\mathcal{O}_b$ is always forbidden. Assuming that the gravitino mass is small, $m_{3/2} \ll m_{\rm C,N}$, this implies that $|b| \ll |\mu|^2$, solving the supersymmetric CP problem associated with the Higgs sector. In the case where $\mathcal{O}_\mu$ is not allowed, the μ and b parameters can be generated from $\mathcal{O}_{\rm SUGRA}$, as long as $m_{3/2}$ is of order the weak scale.

The framework of charged supersymmetry breaking is compatible with many theories of flavor, including theories with extra dimensions or flavor symmetries. For example, it can be combined with the framework described in the previous subsection. This will prohibit the operators $\mathcal{O}_{\eta_H,\eta_\Phi}$, which would otherwise be there. Another interesting way of obtaining the Yukawa hierarchy in this framework is to generate it at lower energies by some strong gauge dynamics [16]. This generates large anomalous dimensions for Φ, and, after canonically normalizing fields, the Yukawa and supersymmetry breaking parameters develop a hierarchy. Note that unlike Higgsphobic or remote flavor-supersymmetry breaking, charged supersymmetry breaking guarantees that the scalar trilinear interactions vanish at M, which is necessary to prevent reintroducing the superpotential flavor problem in models of the type given in Ref. [16] [although the problem may be avoided by making the gauginos much heavier than the scalars, as discussed in (iii) in Sec. III A]. In these theories, the scale for supersymmetry breaking mediation can be high to naturally preserve gauge coupling unification, e.g., $M \gtrsim M_{\rm unif}$, while the scale of flavor physics, M_F, can be much lower as long as the mechanism generating the Yukawa structure does not introduce large relative running between the standard model gauge couplings. In these models, the scale of flavor physics can be as low as 10–100 TeV.

IV. SUPERPARTICLE SPECTRA AND LOW-ENERGY CONSTRAINTS

In the previous section we have presented three classes of theories in which $\mathcal{O}_{\zeta_f}$ are naturally suppressed. This, however, is not sufficient to avoid all the flavor and CP constraints while keeping superparticles light. The constraints from general left-right sfermion propagation discussed in Sec. II D, as well as those from left-left and right-right sfermion propagation, must still be addressed. The strongest constraints on left-left and right-right sfermion propagation arise from ϵ_K and the $\mu \to e\gamma$ process, giving

$$\sqrt{|{\rm Im}(\delta^d_{LL/RR})^2_{12}|} \lesssim 1\times 10^{-2}\left(\frac{m_{\rm C}}{600~{\rm GeV}}\right),\qquad \sqrt{|{\rm Im}(\delta^d_{LL})_{12}(\delta^d_{RR})_{12}|} \lesssim 2\times 10^{-4}\left(\frac{m_{\rm C}}{600~{\rm GeV}}\right), \tag{43}$$

$$|(\delta^e_{LL})_{12}| \lesssim 6\times 10^{-4}\frac{10}{\tan\beta}\left(\frac{m_{\rm N}}{200~{\rm GeV}}\right)^2,\qquad |(\delta^e_{RR})_{12}| \lesssim 3\times 10^{-3}\frac{10}{\tan\beta}\left(\frac{m_{\rm N}}{200~{\rm GeV}}\right)^2. \tag{44}$$

Here, the bounds are obtained conservatively by scanning the ratios of the superparticle masses in the same range as that leading to Eqs. (23)–(26) (see, e.g., [25]).[11] The bounds from ϵ_K are obtained from the conservative requirement that the supersymmetric contribution does not exceed the experimental value of $|\epsilon_{K,\rm exp}| \simeq 2.23\times 10^{-3}$ [50]. The $\mu \to e\gamma$ process also leads to bounds on $|(\delta^\nu_{LL})_{ij}|$ which provide similar constraints as the bound on $|(\delta^e_{LL})_{12}|$ in the theories considered below.

In this section we perform general analyses on flavor and CP constraints in the classes of theories discussed in Sec. III, assuming that the Yukawa structure coefficients $\mathcal{E}^f_{ij}$ factorize. This is possible because the supersymmetry

[10] It is possible that the mechanism generating $\mathcal{O}_{\lambda_A} \approx \int d^2\theta (\ln X)\mathcal{W}^{A\alpha}\mathcal{W}^A_\alpha$ + H.c. also generates other operators, e.g., $\mathcal{O}_{\eta_H}$ and $\mathcal{O}_{\eta_\Phi}$ of the form $\int d^4\theta \ln(X^\dagger X)H^\dagger H$ and $\int d^4\theta \ln(X^\dagger X)\Phi_i^\dagger\Phi_j$ + H.c., of similar size. In the minimal case such as the one in Ref. [20], $\mathcal{O}_{\lambda_A}$ are generated by gauge mediation and the other operators are suppressed (except for $\mathcal{O}_{\kappa_H}$ and $\mathcal{O}_{\kappa_\Phi}$ with $(\kappa_\Phi)_{ij} \propto (Z_\Phi)_{ij}$). The scales of the M_A and the other supersymmetry breaking masses are comparable if $\langle X\rangle \approx g_A^2 M/16\pi^2$. We here assume this structure: $\mathcal{O}_{\eta_H,\eta_\Phi}$ are absent, and the M_A are of the same order as the characteristic scale of the other supersymmetry breaking masses. If the symmetry is nonlinearly realized on X at M, the operators $\mathcal{O}_{\eta_\Phi} \approx \int d^4\theta (X + X^\dagger)\Phi_i^\dagger\Phi_j$ + H.c. are generically expected, leading to a spectrum similar to the case discussed in Sec. III B.

[11] Here, we have also scanned the region $0.3\alpha_1/\alpha_2 \lesssim M_1^2/M_2^2 \lesssim 3\alpha_1/\alpha_2$ and $1 \lesssim \mu^2/m_{\rm N}^2 \lesssim 16$, and required that 10% of the region evades the experimental constraints. If we change these conditions, the bounds would change by a factor of a few, but our conclusions would be unaffected.

breaking parameters follow a definite pattern in each class of theories, which, up to $O(1)$ coefficients, is described by a few free parameters.

A. Factorized flavor structure

In many of the theories discussed in Sec. III, the Yukawa structure coefficients $\mathcal{E}^f_{ij}$ take a factorized form: $\mathcal{E}^u_{ij} = \epsilon_{Q_i}\epsilon_{U_j}$, $\mathcal{E}^d_{ij} = \epsilon_{Q_i}\epsilon_{D_j}$, and $\mathcal{E}^e_{ij} = \epsilon_{L_i}\epsilon_{E_j}$, where $\epsilon_{\Phi_1} \leq \epsilon_{\Phi_2} \leq \epsilon_{\Phi_3}$ without loss of generality. The Yukawa couplings, Eq. (18), are then given by

$$(y_u)_{ij} \approx \tilde{y}\epsilon_{Q_i}\epsilon_{U_j}, \quad (y_d)_{ij} \approx \tilde{y}\epsilon_{Q_i}\epsilon_{D_j}, \quad (y_e)_{ij} \approx \tilde{y}\epsilon_{L_i}\epsilon_{E_j}. \tag{45}$$

In fact, this factorization generically appears in models with Abelian flavor symmetries and those with extra dimensions. Models with non-Abelian symmetries may also obey this, for example, if the $SU(3)_\Phi$ symmetry is broken by three spurions, $\approx (0, 0, \epsilon_{\Phi_3})$, $(0, \epsilon_{\Phi_2}, \epsilon_{\Phi_2})$, and $(\epsilon_{\Phi_1}, \epsilon_{\Phi_1}, \epsilon_{\Phi_1})$ for each $\Phi = Q, U, D, L, E$. The Yukawa couplings of Eq. (45) lead to the following quark and lepton masses and mixings

$$\begin{aligned}
(m_t, m_c, m_u) &\approx \tilde{y}\langle H_u\rangle(\epsilon_{Q_3}\epsilon_{U_3}, \epsilon_{Q_2}\epsilon_{U_2}, \epsilon_{Q_1}\epsilon_{U_1}),\\
(m_b, m_s, m_d) &\approx \tilde{y}\langle H_d\rangle(\epsilon_{Q_3}\epsilon_{D_3}, \epsilon_{Q_2}\epsilon_{D_2}, \epsilon_{Q_1}\epsilon_{D_1}),\\
(m_\tau, m_\mu, m_e) &\approx \tilde{y}\langle H_d\rangle(\epsilon_{L_3}\epsilon_{E_3}, \epsilon_{L_2}\epsilon_{E_2}, \epsilon_{L_1}\epsilon_{E_1}),\\
(m_{\nu_\tau}, m_{\nu_\mu}, m_{\nu_e}) &\approx \frac{\tilde{y}^2\langle H_u\rangle^2}{M_N}(\epsilon_{L_3}^2, \epsilon_{L_2}^2, \epsilon_{L_1}^2),
\end{aligned} \tag{46}$$

and

$$\begin{aligned}
V_{\rm CKM} &\approx \begin{pmatrix} 1 & \epsilon_{Q_1}/\epsilon_{Q_2} & \epsilon_{Q_1}/\epsilon_{Q_3}\\ \epsilon_{Q_1}/\epsilon_{Q_2} & 1 & \epsilon_{Q_2}/\epsilon_{Q_3}\\ \epsilon_{Q_1}/\epsilon_{Q_3} & \epsilon_{Q_2}/\epsilon_{Q_3} & 1 \end{pmatrix},\\
V_{\rm MNS} &\approx \begin{pmatrix} 1 & \epsilon_{L_1}/\epsilon_{L_2} & \epsilon_{L_1}/\epsilon_{L_3}\\ \epsilon_{L_1}/\epsilon_{L_2} & 1 & \epsilon_{L_2}/\epsilon_{L_3}\\ \epsilon_{L_1}/\epsilon_{L_3} & \epsilon_{L_2}/\epsilon_{L_3} & 1 \end{pmatrix},
\end{aligned} \tag{47}$$

where we have included the neutrino masses through the seesaw mechanism with the right-handed neutrino Majorana masses $W \approx M_N\epsilon_{N_i}\epsilon_{N_j}N_iN_j$, and $V_{\rm MNS}$ is the Maki-Nakagawa-Sakata matrix describing lepton mixing, while $V_{\rm CKM}$ is the usual quark mixing matrix.

The values of the ϵ parameters are constrained by the observed quark and lepton masses and mixings through Eqs. (46) and (47). They may also be constrained by possible grand unification. In the analysis of this section we use the following values for ϵ_{Φ_i}, inferred from the quark and lepton masses and mixing run up to the unification scale [51]:

$$\begin{aligned}
&\epsilon_{Q_1} \approx 0.003\tilde{y}^{-1/2}\alpha_q, && \epsilon_{U_1} \approx 0.001\tilde{y}^{-1/2}\alpha_q^{-1},\\
&\epsilon_{Q_2} \approx 0.03\tilde{y}^{-1/2}\alpha_q, && \epsilon_{U_2} \approx 0.04\tilde{y}^{-1/2}\alpha_q^{-1},\\
&\epsilon_{Q_3} \approx 0.7\tilde{y}^{-1/2}\alpha_q, && \epsilon_{U_3} \approx 0.7\tilde{y}^{-1/2}\alpha_q^{-1},\\
&\epsilon_{D_1} \approx 0.002\tilde{y}^{-1/2}\alpha_q^{-1}\tan\beta,\\
&\epsilon_{D_2} \approx 0.004\tilde{y}^{-1/2}\alpha_q^{-1}\tan\beta,\\
&\epsilon_{D_3} \approx 0.01\tilde{y}^{-1/2}\alpha_q^{-1}\tan\beta,
\end{aligned} \tag{48}$$

$$\begin{aligned}
&\epsilon_{L_1} \approx 0.002\tilde{y}^{-1/2}\alpha_l\tan\beta, && \epsilon_{E_1} \approx 0.001\tilde{y}^{-1/2}\alpha_l^{-1},\\
&\epsilon_{L_2} \approx 0.008\tilde{y}^{-1/2}\alpha_l\tan\beta, && \epsilon_{E_2} \approx 0.04\tilde{y}^{-1/2}\alpha_l^{-1},\\
&\epsilon_{L_3} \approx 0.01\tilde{y}^{-1/2}\alpha_l\tan\beta, && \epsilon_{E_3} \approx 0.7\tilde{y}^{-1/2}\alpha_l^{-1},
\end{aligned} \tag{49}$$

where $\alpha_{q,l}$ are numbers parametrizing the freedoms unfixed by the data. We have chosen $\alpha_{q,l}$ so that $SU(5)$ grand unified relations, $\epsilon_{Q_i} = \epsilon_{U_i} = \epsilon_{E_i}$ and $\epsilon_{D_i} = \epsilon_{L_i}$, are almost satisfied with $\alpha_q = \alpha_l = 1$. Note that the precise numbers in Eqs. (48) and (49) are not very important because of unknown $O(1)$ coefficients in the expressions of the quark and lepton masses and mixings as well as the bounds from low-energy flavor and CP violation. In addition, while we have used the data at $M_F \simeq M_{\rm unif}$, using a lower value of M_F would not qualitatively change the results, as it would only change the numbers in Eqs. (48) and (49) by additional $O(1)$ factors.

B. Higgsphobic supersymmetry breaking

Higgsphobic supersymmetry breaking theories discussed in section III B give the following pattern for the flavor-supersymmetry breaking parameters at M_F:

$$(m^2_\Phi)_{ij} \approx \{\epsilon_{\Phi_i}\epsilon_{\Phi_j} + (\eta^\dagger_\Phi\eta_\Phi)_{ij} + \Delta^\Phi_{ij}\}m^2_{\rm S}, \tag{50}$$

$$(a_f)_{ij} \approx \{(y_f)_{kj}(\eta_{\Phi_L})_{ki} + (y_f)_{ik}(\eta_{\Phi_R})_{kj}\}m_{\rm S}, \tag{51}$$

where we have suppressed flavor universal contributions as well as a possible difference between the colored and noncolored superparticle mass scales. Here, $m_{\rm S}$ is the characteristic scale for supersymmetry breaking parameters, $(\eta_\Phi)_{ij} \approx \epsilon_{\Phi_i}\epsilon_{\Phi_j}$ are complex 3×3 matrices, and Δ^Φ_{ij} parameterize flavor violating effects arising from bulk loops in higher dimensional theories discussed in section III B. In flat space models, we typically find $\Delta^\Phi_{ii} \lesssim g^2/16\pi^2 \approx O(10^{-2})$ because they arise from bulk gauge loops, where g represents the standard model gauge couplings. The off diagonal components are smaller, $\Delta^\Phi_{ij}(i \neq j) \ll \Delta^\Phi_{ii}$, since they arise through brane-localized terms which are volume suppressed. On the other hand, in warped space models where H and X are in the infrared region, we expect $\Delta^\Phi_{ij} \lesssim \epsilon_{\Phi_i}\epsilon_{\Phi_j}$, since in the dual 4D picture any couplings of a matter field to H and X, including flavor violating loops in higher dimensions, are controlled by the anomalous dimension of the corresponding strong dynamics operator, which provides a factor ϵ_{Φ_i} for each Φ_i.[12]

The pattern of Eqs. (50) and (51) is essentially the one discussed in Ref. [7] (for $\Delta^{\Phi}_{ij} \lesssim \epsilon_{\Phi_i}\epsilon_{\Phi_j}$). The left-left and right-right mass insertion parameters generated by Eq. (50) are

$$(\delta^u_{LL})_{ij} \approx \left\{(1+\epsilon^2_{Q_3})\epsilon_{Q_i}\epsilon_{Q_j} + (\Delta^Q_i - \Delta^Q_j)\frac{\epsilon_{Q_i}}{\epsilon_{Q_j}}\right\}\frac{m^2_{\rm S}}{m^2_{\rm C}}, \quad (\delta^u_{RR})_{ij} \approx (\delta^u_{LL})_{ij}|_{Q\to U}, \tag{52}$$

$$(\delta^d_{LL})_{ij} \approx \left\{(1+\epsilon^2_{Q_3})\epsilon_{Q_i}\epsilon_{Q_j} + (\Delta^Q_i - \Delta^Q_j)\frac{\epsilon_{Q_i}}{\epsilon_{Q_j}}\right\}\frac{m^2_{\rm S}}{m^2_{\rm C}}, \quad (\delta^d_{RR})_{ij} \approx (\delta^d_{LL})_{ij}|_{Q\to D}, \tag{53}$$

$$(\delta^e_{LL})_{ij} \approx \left\{(1+\epsilon^2_{L_3})\epsilon_{L_i}\epsilon_{L_j} + (\Delta^L_i - \Delta^L_j)\frac{\epsilon_{L_i}}{\epsilon_{L_j}}\right\}\frac{m^2_{\rm S}}{m^2_{\rm N}}, \quad (\delta^e_{RR})_{ij} \approx (\delta^e_{LL})_{ij}|_{L\to E}, \tag{54}$$

for $i<j$, and $(\delta^f_{LL})_{ij} \approx (\delta^f_{LL})_{ji}$ and $(\delta^f_{RR})_{ij} \approx (\delta^f_{RR})_{ji}$. Here, we have retained only the diagonal components of Δ^{Φ}_{ij}, $\Delta^{\Phi}_i \equiv \Delta^{\Phi}_{ii}$, which is justified in most models, as discussed above. The left-right mass insertion parameters given by Eq. (51) are

$$(\delta^u_{LR})_{ij} = (\delta^u_{RL})^*_{ji} \approx \epsilon_{Q_i}\epsilon_{U_j}(\epsilon^2_{Q_j}+\epsilon^2_{U_i})\frac{v\sin\beta}{m_{\rm C}}, \tag{55}$$

$$(\delta^d_{LR})_{ij} = (\delta^d_{RL})^*_{ji} \approx \epsilon_{Q_i}\epsilon_{D_j}(\epsilon^2_{Q_j}+\epsilon^2_{D_i})\frac{v\cos\beta}{m_{\rm C}}, \tag{56}$$

$$(\delta^e_{LR})_{ij} = (\delta^e_{RL})^*_{ji} \approx \epsilon_{L_i}\epsilon_{E_j}(\epsilon^2_{L_j}+\epsilon^2_{E_i})\frac{v\cos\beta}{m_{\rm N}}, \tag{57}$$

where we have assumed that the renormalization group effect makes $m_{\rm S} \to m_{\rm C}$ and $m_{\rm N}$ in the scalar trilinear interactions for colored and noncolored superparticles, respectively. Note that we have suppressed all the subleading terms as well as $O(1)$ coefficients in Eqs. (52)–(57).

One finds that $(\delta^f_{LR})_{ij}$ in Eqs. (55)–(57) are typically much smaller than those in Eq. (21) with $a_{\rm C,N} \approx m_{\rm C,N}$ for $i, j = 1, 2$, so that the left-right flavor violation caused by this source is small. For the multiple mass insertion diagrams discussed in Sec. II D, we find that the effective mass insertion parameters of Eq. (39) for $i, j = 1, 2$, are given in terms of

$$d^f_{ij} \approx \max\left\{(\epsilon^2_{\Phi_{Lj}}+\epsilon^2_{\Phi_{R3}})\hat{\Delta}^{\Phi_L}_{i3}\frac{c_d m^2_{\rm S}}{m^2_{\rm C,N}}, (\epsilon^2_{\Phi_{L3}}+\epsilon^2_{\Phi_{Ri}})\hat{\Delta}^{\Phi_R}_{3j}\frac{c_d m^2_{\rm S}}{m^2_{\rm C,N}}, \hat{\Delta}^{\Phi_L}_{i3}\hat{\Delta}^{\Phi_R}_{3j}\frac{c_t\tilde{y}m_{\rm C}m^4_{\rm S}t}{m^5_{\rm C,N}}, (\epsilon^2_{\Phi_{L3}}+\epsilon^2_{\Phi_{Ri}})(\epsilon^2_{\Phi_{Lj}}+\epsilon^2_{\Phi_{R3}})\epsilon^2_{\Phi_{L3}}\epsilon^2_{\Phi_{R3}}\frac{c_t\tilde{y}v^2 m_{\rm C}}{m^3_{\rm C,N}t}\right\}, \tag{58}$$

as

$$(\delta^f_{LR,{\rm eff}})_{ij} \approx \epsilon_{\Phi_{Li}}\epsilon_{\Phi_{Rj}}\frac{v}{m_{\rm C,N}t} \times \max\left\{d^f_{ij}, \hat{\Delta}^{\Phi_L}_{ij}\frac{c_d\tilde{y}m_{\rm C}m^2_{\rm S}t}{m^3_{\rm C,N}}, \hat{\Delta}^{\Phi_R}_{ij}\frac{c_d\tilde{y}m_{\rm C}m^2_{\rm S}t}{m^3_{\rm C,N}}\right\}, \qquad (i \neq j), \tag{59}$$

and

$$(\delta^f_{LR,{\rm eff}})_{ii} \approx \epsilon_{\Phi_{Li}}\epsilon_{\Phi_{Ri}}\frac{v}{m_{\rm C,N}t}d^f_{ii}. \tag{60}$$

Here,

$$\hat{\Delta}^{\Phi_L}_{ij} \equiv (\Delta^{\Phi_L}_i - \Delta^{\Phi_L}_j)\theta_L + \epsilon^2_{\Phi_{Lj}}, \quad \theta_L = \begin{cases} 1 & \text{for } i<j \\ \epsilon^2_{\Phi_{Lj}}/\epsilon^2_{\Phi_{Li}} & \text{for } i>j \end{cases}, \tag{61}$$

$$\hat{\Delta}^{\Phi_R}_{ij} \equiv (\Delta^{\Phi_R}_i - \Delta^{\Phi_R}_j)\theta_R + \epsilon^2_{\Phi_{Ri}}, \quad \theta_R = \begin{cases} \epsilon^2_{\Phi_{Ri}}/\epsilon^2_{\Phi_{Rj}} & \text{for } i<j \\ 1 & \text{for } i>j \end{cases}, \tag{62}$$

$(\Phi_L, \Phi_R, m_{\rm C,N}, t) = (Q, U, m_{\rm C}, 1)$, $(Q, D, m_{\rm C}, \tan\beta)$, $(L, E, m_{\rm N}, \tan\beta)$ for $f = u, d, e$, and we have assumed $\epsilon_{\Phi_1} \lesssim \epsilon_{\Phi_2} \lesssim \epsilon_{\Phi_3} \lesssim O(1)$, $\tan\beta \gtrsim O(1)$, and $\mu \approx m_{\rm C}$. The values of $(\delta^f_{LR,{\rm eff}})_{ij}$ above should be compared with the "naive" $(\delta^f_{LR})_{ij}$ in Eq. (21) with $a_{\rm C,N} \approx m_{\rm C,N}$

$$(\delta^f_{LR,{\rm naive}})_{ij} \approx \epsilon_{\Phi_{Li}}\epsilon_{\Phi_{Rj}}\frac{v}{m_{\rm C,N}t}. \tag{63}$$

Using Eqs. (48) and (49) and $\Delta^{\Phi}_i \lesssim \max\{g^2/16\pi^2, \epsilon^2_{\Phi_i}\}$, and neglecting contributions sufficiently smaller than $(\delta^f_{LR,{\rm naive}})_{ij}$, the expressions of Eqs. (59) and (60) for moderate $\tan\beta \approx O(10)$ and $\tilde{y} \approx O(1)$ become

$$(\delta^f_{LR,{\rm eff}})_{ij} \approx (\delta^f_{LR,{\rm naive}})_{ij}\max\left\{(y_f)^2_{33}\frac{c_d m^2_{\rm S}}{\tilde{y}^2 m^2_{\rm C,N}}, (y_f)^2_{33}\frac{c_t m_{\rm C}m^4_{\rm S}t}{\tilde{y}m^5_{\rm C,N}}, (y_f)^4_{33}\frac{c_t v^2 m_{\rm C}}{\tilde{y}^3 m^3_{\rm C,N}t}\right\}. \tag{64}$$

[12] In these warped space models, the gaugino masses M_A are likely to be somewhat suppressed compared with $m_{\rm S}$, and a flavor universal contribution to the scalar squared masses at M_F, $\delta m^2_{\Phi}|_{\rm univ} \lesssim M^2_A$, is also expected. Using naive dimensional analysis in the dual 4D picture, we find $M_A \approx g^2_A(N/16\pi^2)^{3/4}m_{\rm S}/\tilde{y}^{1/2}$, where N is the size of the strongly coupled sector yielding H and X. This does not significantly affect the analysis below in some parameter region, although the top squarks may be somewhat heavy in these theories.

(For larger $\tilde{y}$, we have additional potentially relevant contributions of order $\Delta^{\Phi}_{ij} c_d \tilde{y} m_C m_S^2 t/m_{C,N}^3$ inside the curly brackets.) From this, we find that $(\delta^f_{LR,\mathrm{eff}})_{ij}$ can in fact be smaller than $(\delta^f_{LR,\mathrm{naive}})_{ij}$, so that the supersymmetric left-right flavor problem can be solved. Natural values of $(\delta^f_{LR,\mathrm{eff}})_{ij}$ inferred from Eq. (64), however, are not very much smaller than $(\delta^f_{LR,\mathrm{naive}})_{ij}$ (typically no more than an order of magnitude for $f = e$), so we can still expect positive signatures in future search on flavor and *CP* violation, e.g., in EDM experiments, in this class of theories. For larger $\tan\beta$, it becomes increasingly difficult to obtain $(\delta^f_{LR,\mathrm{eff}})_{ij} \ll (\delta^f_{LR,\mathrm{naive}})_{ij}$, so that very large $\tan\beta$, e.g., $\tan\beta \gtrsim 30$, is disfavored.

We now consider constraints from left-left and right-right sfermion propagation. Using Eqs. (48) and (49) in Eqs. (53) and (54), the bounds of Eqs. (43) and (44) give

$$(\Delta_1^Q - \Delta_2^Q) + 9\times 10^{-4}\frac{\alpha_q^2}{\tilde{y}} \lesssim 0.1\left(\frac{m_C}{600\ \mathrm{GeV}}\right)\frac{m_C^2}{m_S^2}, \quad (65)$$

$$(\Delta_1^D - \Delta_2^D) + 2\times 10^{-5}\frac{\tan^2\beta}{\tilde{y}\alpha_q^2} \lesssim 2\times 10^{-2}\left(\frac{m_C}{600\ \mathrm{GeV}}\right)\frac{m_C^2}{m_S^2}, \quad (66)$$

$$\left\{(\Delta_1^Q - \Delta_2^Q) + 9\times 10^{-4}\frac{\alpha_q^2}{\tilde{y}}\right\}^{1/2} \times \left\{(\Delta_1^D - \Delta_2^D) + 2\times 10^{-5}\frac{\tan^2\beta}{\tilde{y}\alpha_q^2}\right\}^{1/2} \lesssim 9\times 10^{-4}\left(\frac{m_C}{600\ \mathrm{GeV}}\right)\frac{m_C^2}{m_S^2}, \quad (67)$$

$$(\Delta_1^L - \Delta_2^L) + 6\times 10^{-5}\frac{\alpha_l^2\tan^2\beta}{\tilde{y}} \lesssim 2\times 10^{-3}\frac{10}{\tan\beta}\left(\frac{m_N}{200\ \mathrm{GeV}}\right)^2\frac{m_N^2}{m_S^2}, \quad (68)$$

$$(\Delta_1^E - \Delta_2^E) + 2\times 10^{-3}\frac{1}{\tilde{y}\alpha_l^2} \lesssim 0.1\frac{10}{\tan\beta}\left(\frac{m_N}{200\ \mathrm{GeV}}\right)^2\frac{m_N^2}{m_S^2}, \quad (69)$$

where we have assumed $\epsilon_{\Phi_3} \lesssim O(1)$. While there is an $O(1)$ coefficient omitted in front of each term, we still find some tension in Eqs. (66)–(68). For $m_S \approx m_N \approx m_C/(2\text{–}4)$, for example, these bounds require

$$\frac{1}{\tilde{y}\alpha_q^2} \lesssim 10^2\left(\frac{10}{\tan\beta}\right)^2\left(\frac{m_C}{600\ \mathrm{GeV}}\right), \quad (70)$$

$$(\Delta_1^Q - \Delta_2^Q)^{1/2}\left\{(\Delta_1^D - \Delta_2^D) + 2\times 10^{-5}\frac{\tan^2\beta}{\tilde{y}\alpha_q^2}\right\}^{1/2} \lesssim 10^{-2}\left(\frac{m_C}{600\ \mathrm{GeV}}\right), \quad (71)$$

$$(\Delta_1^L - \Delta_2^L) + 6\times 10^{-5}\frac{\alpha_l^2\tan^2\beta}{\tilde{y}} \lesssim 10^{-3}\frac{10}{\tan\beta}\left(\frac{m_N}{200\ \mathrm{GeV}}\right)^2, \quad (72)$$

at the order of magnitude level. The conditions of Eqs. (70) and (71) are satisfied in a wide parameter region, while the condition of Eq. (72) requires

$$(\Delta_1^L - \Delta_2^L) \lesssim 10^{-3}\frac{10}{\tan\beta}\left(\frac{m_N}{200\ \mathrm{GeV}}\right)^2, \qquad \frac{\alpha_l^2}{\tilde{y}} \lesssim 0.1\left(\frac{10}{\tan\beta}\right)^3\left(\frac{m_N}{200\ \mathrm{GeV}}\right)^2, \quad (73)$$

unless there is a strong cancellation. The first inequality leads to a tension in theories with $\Delta_i^{\Phi} \approx g^2/16\pi^2 \approx O(10^{-2})$, although it can be ameliorated by taking somewhat large m_N, e.g., $m_N \gtrsim 600\ \mathrm{GeV}(\tan\beta/10)^{1/2}$ or smaller $\tan\beta$. On the other hand, theories with $\Delta_i^{\Phi} \approx \epsilon^2_{\Phi_i}$ have little tension, and taking somewhat small α_l is enough to avoid the bounds. Note that very large $\tan\beta$ is, again, disfavored.

Finally, we discuss implications on the superparticle spectrum. The intergenerational mass splittings between the sfermions are controlled by Eq. (50), leading to

$$|m^2_{\Phi_i} - m^2_{\Phi_j}| \approx m_S^2 \max\{\Delta^{\Phi}_{i,j}, \epsilon^2_{\Phi_{i,j}}\}. \quad (74)$$

In particular, this gives

$$|m^2_{\tilde{\tau}_R} - m^2_{\tilde{\mu}_R}| \approx m_S^2 \max\{\Delta_2^E, \Delta_3^E, \epsilon^2_{E_3}\}, \quad (75)$$

$$|m^2_{\tilde{\tau}_R} - m^2_{\tilde{e}_R}| \approx m_S^2 \max\{\Delta_1^E, \Delta_3^E, \epsilon^2_{E_3}\}, \quad (76)$$

$$|m^2_{\tilde{\mu}_R} - m^2_{\tilde{e}_R}| \approx m_S^2 \max\{\Delta_1^E, \Delta_2^E, \epsilon^2_{E_2}\}, \quad (77)$$

which can lead to $O(1)$ fractional mass splittings between $\tilde{\tau}_R$ and $\tilde{e}_R$, $\tilde{\mu}_R$, and $O(10^{-3}\text{–}10^{-2})$ fractional mass splitting between $\tilde{\mu}_R$ and $\tilde{e}_R$. The signs of the splittings are arbitrary, so that, for example, $\tilde{\tau}_R$ can be heavier than $\tilde{e}_R$, $\tilde{\mu}_R$. Similar levels of intergenerational mass splittings are also possible for other sfermions, although for squarks, splittings will be somewhat diluted by large flavor universal renormalization effects by a factor of order m_S^2/m_C^2.

C. Remote flavor-supersymmetry breaking

The remote flavor-supersymmetry breaking scenario discussed in Sec. III C can give a variety of patterns for the supersymmetry breaking parameters, depending on G_{flavor} and its breaking. In general, G_{flavor} can be a product of a "three generation," "two generation," or "single generation" symmetry acting on each $\Phi = Q, U, D, L, E$. The first class corresponds to (a sufficiently large subgroup of) $SU(3)$ acting on three generations (Φ_1, Φ_2, Φ_3), the second to (a sufficiently large subgroup of) $SU(2)$ acting on the first two generations (Φ_1, Φ_2), and the third to products of $U(1)$ (or Z_N) symmetries.

In the limit of unbroken $G_{\rm flavor}$, the coefficients of the matter supersymmetry breaking operators are given by

$$(\kappa_\Phi)_{ij} \approx \delta_{ij}\kappa_i^\Phi, \qquad (\eta_\Phi)_{ij} \approx \delta_{ij}\eta_i^\Phi, \tag{78}$$

in the field basis where $(Z_\Phi)_{ij} = \delta_{ij}$. Here, depending on the component of $G_{\rm flavor}$ acting on Φ_i, the parameters κ_i^Φ and η_i^Φ exhibit the following pattern:

$$\kappa_1^\Phi = \kappa_2^\Phi = \kappa_3^\Phi, \qquad \eta_1^\Phi = \eta_2^\Phi = \eta_3^\Phi, \quad \text{for "three generation" symmetry,} \tag{79}$$

$$\kappa_1^\Phi = \kappa_2^\Phi \neq \kappa_3^\Phi, \qquad \eta_1^\Phi = \eta_2^\Phi \neq \eta_3^\Phi, \quad \text{for "two generation" symmetry,} \tag{80}$$

$$\kappa_1^\Phi \neq \kappa_2^\Phi \neq \kappa_3^\Phi, \qquad \eta_1^\Phi \neq \eta_2^\Phi \neq \eta_3^\Phi, \quad \text{for "single generation" symmetry.} \tag{81}$$

In any of these theories, our framework provides a solution to the superpotential problem. To see if all the constraints are avoided, we also need to study effects coming from $G_{\rm flavor}$ violation.

We now focus on the case where $G_{\rm flavor}$ contains a "three generation" symmetry for each $\Phi = Q, U, D, L, E$, e.g., $G_{\rm flavor} = SU(3)^5$. We also assume that $G_{\rm flavor}$ is broken by three spurions, $\approx (0, 0, \epsilon_{\Phi_3})$, $(0, \epsilon_{\Phi_2}, \epsilon_{\Phi_2})$, and $(\epsilon_{\Phi_1}, \epsilon_{\Phi_1}, \epsilon_{\Phi_1})$ for each $\Phi = Q, U, D, L, E$, which guarantees that $\mathcal{E}_{ij}^f$ take a factorized form. The operator coefficients at M_F in the field basis that naturally realizes $G_{\rm flavor}$ are then given by

$$(Z_\Phi)_{ij} \approx \delta_{ij} + \gamma\epsilon_{\Phi_i}\epsilon_{\Phi_j}, \qquad (\kappa_\Phi)_{ij} \approx \delta_{ij} + (D\text{-term}), \qquad (\eta_\Phi)_{ij} \approx \delta_{ij}, \tag{82}$$

where we have omitted $O(1)$ coefficients that appear in each (i, j) element of the second term of $(Z_\Phi)_{ij}$, in front of the first term of $(\kappa_\Phi)_{ij}$, and in front of the expression for $(\eta_\Phi)_{ij}$. Here, γ parametrizes the strength of the $G_{\rm flavor}$ breaking effect, which is suppressed by the volume of the extra dimensions in a higher dimensional realization of the scenario. (If the matter fields have nontrivial wave functions and/or $G_{\rm flavor}$ is broken on several different branes, γ can depend on Φ, i, j.) The second term of $(\kappa_\Phi)_{ij}$, denoted as D term, arises if $G_{\rm flavor}$ contains a continuous gauge symmetry component (for $M_F \lesssim M$), but it is absent if $G_{\rm flavor}$ is a global or discrete symmetry.

After canonically normalizing fields, $(Z_\Phi)_{ij} = \delta_{ij}$, Eq. (82) leads to the following operator coefficients:

$$(\kappa_\Phi)_{ij} \approx \delta_{ij}(1 + \gamma\epsilon_{\Phi_i}^2 + \Delta_i^\Phi), \qquad (\eta_\Phi)_{ij} \approx \delta_{ij}(1 + \gamma\epsilon_{\Phi_i}^2), \tag{83}$$

where we have diagonalized $(Z_\Phi)_{ij}$ of Eq. (82) and then rescaled fields so that the $(Z_\Phi)_{ij}$ become δ_{ij}. Here, the third term of $(\kappa_\Phi)_{ij}$ parametrizes the possible D-term contribution (which is always flavor diagonal in this basis). This contribution arises when a continuous gauged $G_{\rm flavor}$ symmetry is broken by the VEVs of fields ϕ_m $(m = 1, 2, \cdots)$ which have supersymmetry breaking masses [52]. For example, if $G_{\rm flavor}$ is broken by ϕ_1 and ϕ_2 whose transformation properties under $G_{\rm flavor}$ is opposite, a D-term contribution of order $m_{\phi_1}^2 - m_{\phi_2}^2$ generically arises, where $m_{\phi_m}^2$ represents the supersymmetry breaking mass squared of ϕ_m. While the D-term contribution is generically dangerous in theories with a gauged flavor symmetry, in the present framework the supersymmetry breaking masses of ϕ_m are suppressed because of the separation between $G_{\rm flavor}$ and supersymmetry breaking, so that the resulting D-term contribution Δ_i^Φ is also suppressed. (This was observed in Ref. [53] in theories with $G_{\rm flavor} = U(1)$.) The contribution is typically suppressed by a loop factor, as well as by powers of the ratio of the compactification scale, M_c, to the (effective) cutoff scale, M_*, in theories with extra dimensions. We therefore expect $\Delta_i^\Phi \approx (M_c/M_*)^n/16\pi^2 \lesssim O(10^{-2})$, where n is a model-dependent integer which can in general depend on Φ and i.[13]

Note that in the field basis leading to Eq. (83) the Yukawa couplings are still given by Eq. (45). The left-left and right-right mass insertion parameters are then given by

$$(\delta_{LL}^u)_{ij} \approx \left\{\gamma(1 + \gamma\epsilon_{Q_3}^2)\epsilon_{Q_i}\epsilon_{Q_j} + (\Delta_i^Q - \Delta_j^Q)\frac{\epsilon_{Q_i}}{\epsilon_{Q_j}}\right\}\frac{m_{\rm S}^2}{m_{\rm C}^2}, \qquad (\delta_{RR}^u)_{ij} \approx (\delta_{LL}^u)_{ij}|_{Q\to U}, \tag{84}$$

$$(\delta_{LL}^d)_{ij} \approx \left\{\gamma(1 + \gamma\epsilon_{Q_3}^2)\epsilon_{Q_i}\epsilon_{Q_j} + (\Delta_i^Q - \Delta_j^Q)\frac{\epsilon_{Q_i}}{\epsilon_{Q_j}}\right\}\frac{m_{\rm S}^2}{m_{\rm C}^2}, \qquad (\delta_{RR}^d)_{ij} \approx (\delta_{LL}^d)_{ij}|_{Q\to D}, \tag{85}$$

$$(\delta_{LL}^e)_{ij} \approx \left\{\gamma(1 + \gamma\epsilon_{L_3}^2)\epsilon_{L_i}\epsilon_{L_j} + (\Delta_i^L - \Delta_j^L)\frac{\epsilon_{L_i}}{\epsilon_{L_j}}\right\}\frac{m_{\rm S}^2}{m_{\rm N}^2}, \qquad (\delta_{RR}^e)_{ij} \approx (\delta_{LL}^e)_{ij}|_{L\to E}, \tag{86}$$

[13] If different ϕ_m have different renormalizable interactions of order λ, $\Delta_i^\Phi \approx (\lambda^2/16\pi^2)^2(m_{3/2}/m_{\rm S})^2$ can be generated through anomaly mediation. This contribution is typically of order 10^{-4} or smaller.

for $i < j$, and $(\delta^f_{LL})_{ij} \approx (\delta^f_{LL})_{ji}$ and $(\delta^f_{RR})_{ij} \approx (\delta^f_{RR})_{ji}$. The flavor and CP violating left-right mass insertion parameters are given by

$$(\delta^u_{LR})_{ij} = (\delta^u_{RL})^*_{ji} \approx \gamma\epsilon_{Q_i}\epsilon_{U_j}(\epsilon^2_{Q_j} + \epsilon^2_{U_i})\frac{v\sin\beta}{m_{\rm C}}, \quad (87)$$

$$(\delta^d_{LR})_{ij} = (\delta^d_{RL})^*_{ji} \approx \gamma\epsilon_{Q_i}\epsilon_{D_j}(\epsilon^2_{Q_j} + \epsilon^2_{D_i})\frac{v\cos\beta}{m_{\rm C}}, \quad (88)$$

$$(\delta^e_{LR})_{ij} = (\delta^e_{RL})^*_{ji} \approx \gamma\epsilon_{L_i}\epsilon_{E_j}(\epsilon^2_{L_j} + \epsilon^2_{E_i})\frac{v\cos\beta}{m_{\rm N}}. \quad (89)$$

We find that the mass insertion parameters of Eqs. (84)–(89) take the same form as those of Eqs. (52)–(57) for $\gamma = 1$, although the origins of Δ^Φ_i are different. Since γ is expected to be smaller than 1, the present class of theories is at least as safe as Higgsphobic theories with $\Delta^\Phi_i \approx O(10^{-2})$ from the flavor and CP violation point of view. Moreover, since we expect $\Delta^\Phi_i < O(10^{-2})$ in most cases due to power suppression of $(M_c/M_*)^n$, or even absent if $G_{\rm flavor}$ is a global or discrete symmetry, the low-energy constraints from flavor and CP violation generically give little tension with LHC observability in the present class of theories.

The intergenerational mass splittings between sfermions have a similar formula to the Higgsphobic case

$$|m^2_{\Phi_i} - m^2_{\Phi_j}| \approx m^2_{\rm S}\max\{\Delta^\Phi_{i,j}, \gamma\epsilon^2_{\Phi_{i,j}}\}. \quad (90)$$

In particular, the right-handed sleptons have the splittings

$$|m^2_{\tilde{\tau}_R} - m^2_{\tilde{\mu}_R}| \approx m^2_{\rm S}\max\{\Delta^E_2, \Delta^E_3, \gamma\epsilon^2_{E_3}\}, \quad (91)$$

$$|m^2_{\tilde{\tau}_R} - m^2_{\tilde{e}_R}| \approx m^2_{\rm S}\max\{\Delta^E_1, \Delta^E_3, \gamma\epsilon^2_{E_3}\}, \quad (92)$$

$$|m^2_{\tilde{\mu}_R} - m^2_{\tilde{e}_R}| \approx m^2_{\rm S}\max\{\Delta^E_1, \Delta^E_2, \gamma\epsilon^2_{E_2}\}. \quad (93)$$

The size of the splittings depends on the details of the theory, specifically the size of Δ^E_i and γ. We expect that the fractional mass splittings between $\tilde{\tau}_R$ and $\tilde{e}_R$, $\tilde{\mu}_R$ and between $\tilde{\mu}_R$ and $\tilde{e}_R$ are at most of $O(1)$ and $O(10^{-2})$, respectively, and typically smaller by at least a factor of a few because of the suppression by γ and $(M_c/M_*)^n$.

In the above analysis, we have focused on the case that $G_{\rm flavor}$ is the product of five ‘‘three generation’’ symmetries acting on Q, U, D, L, E. Similar analyses, however, can also be performed in the case where (some of) Q, U, D, L, E have only ‘‘two generation’’ or ‘‘single generation’’ symmetries. In particular, in the case of a ‘‘two generation’’ symmetry, we expect that the conclusion on flavor and CP violation does not change because the constraints from the processes involving the third generation particles are weak. The fractional mass splittings between the third and first two generation sfermions in this case are expected to be of $O(1)$, without a suppression from γ or $(M_c/M_*)^n$. In the case of a ‘‘single generation’’ symmetry, model by model analyses are needed. We expect, however, that the model can avoid the constraints if it involves ‘‘single generation’’ symmetries only for some Φ. For example, we can consider only E has a ‘‘single generation’’ symmetry while the rest have ‘‘three generation’’ symmetries, e.g., $G_{\rm flavor} = SU(3)_Q \times SU(3)_U \times SU(3)_D \times SU(3)_L \times U(1)_E$. In this case the fractional mass splittings between $\tilde{\tau}_R$ and $\tilde{\mu}_R$ and between $\tilde{\mu}_R$ and $\tilde{e}_R$ are of the same order, and presumably of $O(1)$.

D. Charged supersymmetry breaking

If the supersymmetry breaking field is charged under some symmetry, the operators $\mathcal{O}_{\eta_\Phi}$ are forbidden. In fact, this charged supersymmetry breaking framework can be combined with many flavor theories. To parametrize these wide classes of theories, we somewhat arbitrarily take

$$(\kappa_\Phi)_{ij} \approx \delta_{ij} + \gamma\epsilon_{\Phi_i}\epsilon_{\Phi_j} + \delta_{ij}\Delta^\Phi_i, \qquad (\eta_\Phi)_{ij} = 0, \quad (94)$$

at the scale M_F. This parametrization accommodates many theories of flavor, including ones based on extra dimensions, strong dynamics, and ‘‘three generation’’ flavor symmetries. [In some of these theories, the flavor universal part, i.e., the first term, of $(\kappa_\Phi)_{ij}$ is absent, but this does not affect the analysis of flavor and CP violation.] This parametrization needs to be modified appropriately for other types of theories, for example, those based on ‘‘two generation’’ or ‘‘single generation’’ flavor symmetries.

The left-left and right-right mass insertion parameters generated by Eq. (94) are

$$(\delta^u_{LL})_{ij} \approx \left\{\gamma\epsilon_{Q_i}\epsilon_{Q_j} + (\Delta^Q_i - \Delta^Q_j)\frac{\epsilon_{Q_i}}{\epsilon_{Q_j}}\right\}\frac{m^2_{\rm S}}{m^2_{\rm C}}, \qquad (\delta^u_{RR})_{ij} \approx (\delta^u_{LL})_{ij}|_{Q\to U}, \quad (95)$$

$$(\delta^d_{LL})_{ij} \approx \left\{\gamma\epsilon_{Q_i}\epsilon_{Q_j} + (\Delta^Q_i - \Delta^Q_j)\frac{\epsilon_{Q_i}}{\epsilon_{Q_j}}\right\}\frac{m^2_{\rm S}}{m^2_{\rm C}}, \qquad (\delta^d_{RR})_{ij} \approx (\delta^d_{LL})_{ij}|_{Q\to D}, \quad (96)$$

$$(\delta^e_{LL})_{ij} \approx \left\{\gamma\epsilon_{L_i}\epsilon_{L_j} + (\Delta^L_i - \Delta^L_j)\frac{\epsilon_{L_i}}{\epsilon_{L_j}}\right\}\frac{m^2_{\rm S}}{m^2_{\rm N}}, \qquad (\delta^e_{RR})_{ij} \approx (\delta^e_{LL})_{ij}|_{L\to E}, \quad (97)$$

for $i < j$, and $(\delta^f_{LL})_{ij} \approx (\delta^f_{LL})_{ji}$ and $(\delta^f_{RR})_{ij} \approx (\delta^f_{RR})_{ji}$, while the flavor and CP violating left-right mass insertion parameters are

$$(\delta^f_{LR})_{ij} = (\delta^f_{RL})^*_{ji} \approx 0. \quad (98)$$

The constraints from low-energy flavor and CP violation are obviously not stronger than in the classes of theories discussed in the previous two subsections for the same values of γ and Δ^Φ_i. Note that while the flavor and CP violating left-right mass insertion parameters are vanishing (up to the higher order effects from the Yukawa couplings), we still have flavor and *CP* violating effects generated by multiple mass insertion diagrams through the flavor universal part of $(\delta^f_{LR})_{ij}$ and through flavor and *CP* violating $(\delta^f_{LL,RR})_{ij}$. Nontrivial flavor and *CP* violation, therefore, can still be discovered in future experiments such as ones in Refs. [34–41].

The intergenerational mass splittings between sfermions are given by the formula in Eq. (90). The size of the splittings is controlled by the parameters γ and Δ_i^Φ, which are model dependent. In many models, $\gamma \lesssim O(1)$ and $\Delta_i^\Phi \lesssim O(10^{-2})$. We can, however, still expect that a variety of patterns for the intergenerational mass splittings can be obtained in this framework.

V. PROBING THE ORIGIN OF FLAVOR AT THE LHC

We have seen that supersymmetry with a flavorful spectrum is consistent with bounds from low-energy flavor and *CP* violating processes in a wide variety of models where the superpotential flavor problem is solved. Furthermore, the spectrum in these models can easily be light enough that it will have a substantial production cross section at the LHC. While it appears that the deviation from flavor universality is small, especially in the first two generations, the splittings are large enough that they can be significant at the LHC. Consider the right-handed sleptons, which from the structure of the SSM renormalization group equations are expected to be the lightest sfermions. If flavor nonuniversality comes only from renormalization group evolution, then the fractional mass splitting between $\tilde{e}_R$ and $\tilde{\mu}_R$ is controlled by the muon Yukawa coupling, and is expected to be of $O(10^{-4})$ for $\tan\beta \approx O(10)$. On the other hand, the analysis in Sec. IV shows that the fractional mass splitting from the contribution at M_F can be easily of $O(10^{-2})$. The splittings between $\tilde{\tau}_R$ and $\tilde{e}_R$, $\tilde{\mu}_R$ will also be larger than the renormalization group prediction in the moderate $\tan\beta$ regime. Furthermore, the contribution from renormalization group flow has a definite sign, with the mass ordering of the superparticles being anticorrelated with that of the standard model particles, i.e., $m_{\tilde{\tau}_R} < m_{\tilde{\mu}_R} < m_{\tilde{e}_R}$. On the other hand, the contributions considered in Sec. IV could be positive or negative, and could thus produce a spectrum which is unambiguously different from flavor universality.

An interesting possibility considered in Refs. [7,12] is the case where the lightest supersymmetric particle (LSP) is the gravitino, and the next-to-lightest supersymmetric particle (NLSP) is one of the right-handed sleptons. In the case where the supersymmetry breaking scale is large, i.e., M is large, the lifetime of the NLSP is long enough that it escapes the detector. A slepton NLSP will then appear as a heavy charged stable particle, something quite spectacular at the LHC. (Here we consider the case where R parity is conserved so that decay of the NLSP into only standard model particles is forbidden.) Furthermore, the LHC will usually produce squarks or gluinos, so NLSPs will generally be produced after a chain of decays. This cascade decay must produce a slepton in conjunction with a lepton, and because the NLSP is right-handed, the coupling to neutrinos is strongly suppressed. Therefore, NLSPs will be produced mostly with charged leptons that can be used to measure the flavor content of the NLSP.

If we could observe the decay of the NLSP, we could measure the flavor content of the NLSP more easily. In particular, we could precisely determine, by observation of flavor violating decays, if the flavor eigenbasis differs from the mass eigenbasis for the sleptons. Flavorful models including those discussed in Sec. III generically have this property, so this is an interesting test of intrinsic flavor nonuniversality. One way to measure decays of long-lived NLSP sleptons is to build a stopper detector outside one of the main LHC detectors, which would stop some of the NLSPs and measure their decays [54]. If the NLSP has a sufficiently long lifetime and a large number of decays are observed in a stopper detector, then the flavor composition of the NLSP can be measured very accurately.

While the scenario with a weak scale gravitino LSP has an NLSP with lifetime much longer than the flight time in the detector, the scenario can be extended to a much lighter gravitino LSP or another extremely weakly interacting particle, such as the axino. In fact, if the NLSP has a lifetime of order $c\tau \approx O(100\ \mu\text{m–}10\ \text{m})$, then there is a clean signal of a nonminimum ionizing track with a kink that turns into a minimum ionizing track [55]. Furthermore, an ATLAS study with the NLSP decaying to photons showed that a substantial number of decays can be measured if $c\tau \lesssim 100$ km [56], so it is conceivable that the decay of an NLSP can be observed in the detector for a large range of NLSP lifetimes. This kind of measurement would allow us to study the decays of the NLSP and gain knowledge on its flavor content.

We now focus on the case where the lightest neutralino is lighter than all of the right-handed sleptons, although the analysis also applies if all the superparticles promptly decay to a different particle which escapes the detector. In this case, the events are similar to well studied supersymmetric missing energy events, but many interesting studies of flavor violation can still be done. For example, if $m_{\tilde{\tau}_R} > m_{\chi_2^0} > m_{\tilde{e}_R,\tilde{\mu}_R}$, a natural spectrum in the models of Sec. III given the $O(1)$ mass splitting of the $\tilde{\tau}_R$ from the other sleptons, then the detectors can measure a fractional mass splitting between $\tilde{e}_R$ and $\tilde{\mu}_R$ as small as $O(10^{-4})$ [57]. This may help discriminate different classes of theories discussed in Sec. III. These and other studies could be performed not only in the right-handed slepton sector, but also with squarks and left-handed sleptons, which may determine whether the different SSM fields transform under different classes of flavor symmetries, as discussed in Sec. IV C. With a little luck, the LHC could discover a smoking gun for intrinsic flavor nonuniversal effects in the supersymmetry breaking sector, which may hint at the high energy theory that gives rise to the standard model flavor structure.

VI. DISCUSSION AND CONCLUSIONS

The problem of excessive flavor and CP violation arising from generic weak scale supersymmetry breaking parameters has been a guiding principle in searching for viable supersymmetric theories. In particular, this has been a strong motivation behind the search for flavor universal mediation mechanisms of supersymmetry breaking. On the other hand, the puzzle of flavor already exists in the standard model, and it is possible that the mechanism producing the observed Yukawa structure is also responsible for the suppression of possibly large flavor and CP violation in supersymmetric theories. How natural is this possibility? Is there any generic tension between the constraints from low-energy flavor and CP violation and the observability of superparticles at the LHC, even if we take into account the possibility of a correlation between the structures of the Yukawa couplings and supersymmetry breaking parameters?

In this paper we have studied, in a model-independent way, the question of compatibility between the low-energy flavor and CP constraints and the observability of a nontrivial flavor structure in superparticle spectra at the LHC. We have seen that there is a model-independent tension arising from the superpotential operators $\mathcal{O}_{\zeta_f}$ leading to scalar trilinear interactions. In particular, the constraint from the $\mu \to e\gamma$ process pushes the mass scale for noncolored superparticles $m_{\rm N}$ relatively high. Under the assumption of a factorized flavor structure, $m_{\rm N}$ should be of order a TeV or larger for a natural size of $\mathcal{O}_{\zeta_f}$. Assuming the usual hierarchy between colored and noncolored superparticle masses, this pushes up the masses of colored superparticles beyond several TeV, making supersymmetry unobservable at the LHC. Similar, though somewhat weaker constraints also arise from the bounds on the electron, neutron, and mercury EDMs. In fact, these observables also constrain flavor violation arising from other operators through multiple mass insertion diagrams.

We have discussed several ways in which these stringent constraints are naturally avoided. They include relaxing the mass hierarchy between colored and noncolored superparticles, making the fundamental strength of the Yukawa couplings strong, and making the gaugino masses larger than the scalar masses. We have also presented simple frameworks in which the dangerous operators $\mathcal{O}_{\zeta_f}$ are naturally absent. Since these operators are special, they can be absent in the low-energy effective theory. We have considered separating the Higgs and supersymmetry breaking fields (Higgsphobic), separating supersymmetry and flavor symmetry breaking (remote flavor-supersymmetry breaking), and assigning a nontrivial charge to the supersymmetry breaking field (charged supersymmetry breaking). These frameworks can be combined with a variety of flavor theories, including ones with (flat or warped) extra dimensions, strong dynamics, or flavor symmetries. In fact, we can consider many variations of flavor models using the basic setups discussed in this paper. The mediation scale of supersymmetry breaking, M, and the scale for flavor physics, M_F, can vary by many orders of magnitude in these theories.

We have performed detailed analyses on the constraints from low-energy flavor and CP violation in the frameworks described above. We have shown that the constraints, including ones arising from multiple mass insertion diagrams and left-left and right-right sfermion propagation, can be avoided in natural parameter regions while keeping the superparticles light. Expected sizes of flavor and CP violation, however, are not too much smaller than the current bounds, so signatures in future search on flavor and CP violation are not eliminated. The intergenerational mass splittings among sfermions in these theories can show a variety of patterns depending on the underlying mechanism responsible for the structure of the Yukawa couplings. For example, if SSM multiplets belonging to different representations of the standard model gauge group have different flavor symmetry structures, then it will show up in the spectrum of superparticles. In general, it is significant that the spectrum of superparticles contains information on left-handed and right-handed fields separately, while the Yukawa couplings contain only "products" of them.

While the size of the fractional mass splittings directly generated by the physics of flavor at M_F is not necessarily very large, e.g., of $O(10^{-2})$ or smaller for the first two generations, they are large enough to significantly affect the phenomenology at the LHC. This is because the intergenerational mass splittings generated by the standard SSM renormalization group evolution are typically very small, so that additional mass splittings can give large effects on the structure of the intergenerational superparticle spectrum. Moreover, because the signs of the intergenerational mass splittings caused by effects at M_F are arbitrary (at least from the low-energy effective field theory point of view), these additional splittings can change the mass ordering among different generation sfermions. In particular, this can make a third generation sfermion heavier than the corresponding first two generation sfermions, which can drastically affect the signatures at the LHC. We find it very possible that the intergenerational mass splittings of the size implied by the classes of theories discussed in this paper will be measured at the LHC.

The LHC will start running this year, and it is expected to give us meaningful data on TeV scale physics as early as next year. If supersymmetry is found, it will not only provide an explanation for the stability of the gauge hierarchy and a potential dark matter candidate, but it will also allow for a substantial number of new flavor measurements. While there are many viable supersymmetric models which are flavor universal, we have shown that there are also many nonuniversal models which avoid the stringent low-energy constraints. If supersymmetry is in fact flavorful, then its discovery at the LHC could shed new light on the longstanding mystery of the flavor pattern in the standard model.

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APPENDIX A: μ AND b IN HIGGSPHOBIC SUPERSYMMETRY BREAKING

In minimal Higgsphobic supersymmetry breaking, the operators $\mathcal{O}_{\kappa_H,\eta_H,\mu,b}$ in Eqs. (7) and (8) are forbidden. The μ and b parameters are then generated only by the operator $\mathcal{O}_{\mathrm{SUGRA}}$ through supergravity effects [58], giving

$$\mu = \frac{\lambda_H m_{3/2}^*}{(Z_{H_u} Z_{H_d})^{1/2}}, \qquad b = \frac{\lambda_H |m_{3/2}|^2}{(Z_{H_u} Z_{H_d})^{1/2}}, \tag{A1}$$

where $m_{3/2}$ is the gravitino mass. The nonholomorphic supersymmetry breaking masses are vanishing, $m_{H_u}^2 = m_{H_d}^2 = 0$, at the scale M. The expressions of Eq. (A1) imply that the gravitino mass should be of order the weak scale to use this contribution. The μ and b parameters of Eq. (A1) also lead to dangerous CP violation at low energies unless $\arg(m_{3/2}) \simeq \arg(M_A)$, providing an additional constraint on the setup. In the context of the higher dimensional models of Sec. III B, a weak scale gravitino mass is obtained if $M_* \approx M_{\mathrm{Pl}}$ or if there is an additional supersymmetry breaking field X' that does not couple to the SSM field and has $F_{X'} \approx F_X(M_{\mathrm{Pl}}/M_*)$.

It is possible to extend the minimal Higgsphobic setup by introducing fields B which directly couple with both H and X. In higher dimensional models, these B fields propagate in the bulk, and integrating them out can generate the operators $\mathcal{O}_{\kappa_H,\eta_H,\mu,b}$ in the low-energy effective field theory below $1/R$. The operators $\mathcal{O}_{\zeta_f}$ can still be absent by arranging the interactions of B appropriately, for example, by suppressing the couplings of B with matter fields. With these extensions, the generated Higgs sector parameters need not take the form in the minimal setup. In particular, the gravitino mass need not be of order the weak scale, and its phase need not be aligned with that of M_A. The requirement from suppressing CP violation, instead, constrains the representations and interactions of the B fields. For example, if the exchange of B generates μ but not b, and the contribution from Eq. (A1) is negligible, then the problem of CP violation disappears.

APPENDIX B: $(a_f)_{ij}$ IN REMOTE FLAVOR-SUPERSYMMETRY BREAKING

In remote flavor-supersymmetry breaking, the Yukawa couplings are generated through breaking of G_{flavor}. Suppose that the breaking is caused by the VEVs of several chiral superfields ϕ_m ($m = 1, 2, \cdots$). The Yukawa couplings are then generated from operators of the form

$$\mathcal{L} = \int d^2\theta \sum_{i,j} \sum_{\{(n_f)_{ij}^m\}} c_{\{(n_f)_{ij}^m\}} \frac{\prod_m \phi_m^{(n_f)_{ij}^m}}{(M_* C)^{(n_f)_{ij}}} \Phi_{Li} \Phi_{Rj} H + \mathrm{H.c.}, \tag{B1}$$

as

$$(y_f)_{ij} = \sum_{\{(n_f)_{ij}^m\}} c_{\{(n_f)_{ij}^m\}} \frac{\prod_m \phi_{m,0}^{(n_f)_{ij}^m}}{M_*^{(n_f)_{ij}}}, \tag{B2}$$

where Φ_{Li}, Φ_{Rj}, H, and ϕ_m are canonically normalized, M_* is the (effective) cutoff scale, $(n_f)_{ij}^m$ are integers with $(n_f)_{ij} \equiv \sum_m (n_f)_{ij}^m$, and $\phi_{m,0}$ is the lowest component VEV of ϕ_m. The sum $\sum_{\{(n_f)_{ij}^m\}}$ runs over all possible choices of integers $(n_f)_{ij}^m$ consistent with G_{flavor} invariance, and $c_{\{(n_f)_{ij}^m\}}$ are $O(1)$ coefficients in front of each term. Here, we have included the chiral compensator field $C = 1 + \theta^2 m_{3/2}$ which encodes supergravity effects [59].

The operators of Eq. (B1) may generate dangerous scalar trilinear interactions. These are given by

$$(a_f)_{ij} = \sum_{\{(n_f)_{ij}^m\}} c_{\{(n_f)_{ij}^m\}} \frac{\prod_m \phi_{m,0}^{(n_f)_{ij}^m}}{M_*^{(n_f)_{ij}}} \times \left\{ (n_f)_{ij} m_{3/2} - \sum_{m'} (n_f)_{ij}^{m'} \frac{F_{\phi_{m'}}}{\phi_{m',0}} \right\}, \tag{B3}$$

where F_{ϕ_m} is the F-term VEV of ϕ_m. This shows that even with $F_{\phi_m} = 0$, the scalar trilinear interactions are generated, which for $m_{3/2} \approx O(m_{\mathrm{C}}, m_{\mathrm{N}})$ lead to $(a_f)_{ij} \approx (y_f)_{ij} m_{\mathrm{C,N}}$, and are thus dangerous [60].[14] More generally, if some of the ϕ_m are stabilized using supersymmetry breaking effects (e.g., if these fields are flat directions lifted

[14]Our language here is different from that used in Ref. [60], in which the F-term VEV of a field is defined including a supergravity contribution so that the effect described here is viewed as arising from the F-term VEVs of the fields ϕ_m.

by higher dimension operators), we obtain $F_{\phi_m}/\phi_{m,0} \approx O(\max\{m_{3/2}, m_\phi\})$ with m_ϕ being the scale for the supersymmetry breaking masses of ϕ_m, and we obtain a contribution of order $(a_f)_{ij} \approx (y_f)_{ij} \max\{m_{3/2}, m_\phi\}$.

The contribution to $(a_f)_{ij}$ described above, however, is suppressed if one of the following conditions is satisfied[15]:

(a) The gravitino mass and the F-term VEVs for ϕ_m are all small, $m_a \equiv \max\{m_{3/2}, F_{\phi_m}/\phi_{m,0}\} \ll m_{\rm C,N}$. In this case, the effect of Eq. (B3) is suppressed by a factor of $m_a/m_{\rm C,N}$, giving $(\delta^{u,d}_{LR})_{ij} \approx (M_{u,d})_{ij} m_a/m^2_{\rm C}$ and $(\delta^e_{LR})_{ij} \approx (M_e)_{ij} m_a/m^2_{\rm N}$.

(b) The dimensions of the operators in Eq. (B1) are the same for all $i, j = 1, 2, 3$, i.e., $(n_f)_{ij} = n_f$, for $f = u, d, e$, and $F_{\phi_m}/\phi_{m,0} \ll m_{\rm C,N}$ with $\arg(m_{3/2}) = \arg(M_A)$ (or $(n_f)_{ij} = n_f$ and $F_{\phi_m}/\phi_{m,0}$ are nearly equal with their phases aligned with those of $m_{3/2}, M_A$). In this case, $(a_f)_{ij}$ is almost proportional to $(y_f)_{ij}$ as a matrix, giving a negligible contribution to $(\delta^f_{LR})_{ij}$.

(c) The VEVs of ϕ_m are stabilized in the supersymmetric limit. In this case, we obtain $\langle \phi_m \rangle = \phi_{m,0} C$, since any supersymmetric mass must be accompanied with C, leading to $F_{\phi_m}/\phi_{m,0} = m_{3/2}$. Equation (B3) then gives $(a_f)_{ij} = 0$, and the effect disappears.[16]

We find it simplest to stabilize ϕ_m supersymmetrically, (c), although we also leave the possibility open to (a) or (b). (In fact, the experimental bounds may be avoided with one of the above conditions satisfied only for $i, j = 1, 2$.) Note that the consideration here applies to any field that appears in front of $\Phi_{Li}\Phi_{Rj}H$ in the superpotential, and whose lowest component VEV gives a significant contribution to the Yukawa couplings.

The issue of scalar trilinear interactions generated by the F-term VEVs of C and the fields appearing in front of $\Phi_{Li}\Phi_{Rj}H$, in fact, exists in wider classes of theories. For example, in theories where the hierarchical Yukawa couplings are generated by wave function profiles of the matter and Higgs fields in extra dimensions, including ones discussed in Sec. III B, there are generally moduli fields appearing in front of $\Phi_{Li}\Phi_{Rj}H$. These moduli fields must satisfy condition (a) or (c) [option (b) is typically not available]. We assume that one of these conditions is satisfied when discussing the classes of theories in Sec. III.

[15]While completing this paper, Ref. [61] appeared, which discusses the issue considered in this Appendix. Their main solution corresponds to our (c) below. They also discuss the case (b).

[16]There is still an effect from the chiral compensator field at a loop level (anomaly mediation) [62]. This effect, however, does not lead to flavor or CP violation at a dangerous level. There could also be higher order corrections suppressed by powers of $m_\phi/\phi_{m,0}$, where m_ϕ represents generic supersymmetry breaking masses in the $G_{\rm flavor}$ breaking sector. These corrections are also negligible unless the scale of flavor physics is close to the TeV scale.

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