Effects of Flavor Violation on Split Supersymmetry

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Abstract

We discuss consequences of flavor mixings in the scalar fermion sector of supersymmetric models endowed with ultra heavy scalars. We find that, under extreme fine-tunings different than but similar in size to that needed to obtain a light Higgs doublet, intergenerational mixings generically lead to light sfermion states which facilitate a number of phenomena ranging from rare processes to electric dipole moments. The number of scalar fermions to be discovered by collider searches is fewer than those in a complete supersymmetric model.
The models of physical phenomena are generically based on the principle of naturalness, that is, the mother Nature strongly disfavors fine-tunings. The present understanding of the four fundamental forces of Nature faces with two naturalness problems: (i) Higgs boson mass exhibits a quadratic sensitivity to the ultraviolet (UV) cut-off of the standard electroweak theory (SM), and (ii) experimental value of the cosmological constant (CC) turns out to be far below the theoretical estimates. Concerning the former, a resolution lies in embedding the SM into a UV-safe extension above Fermi energies. A solution for the CC problem, on the other hand, might eventually require modifications in the gravitational dynamics of the vacuum energy. If Nature indeed prefers a supersymmetric organizing principle above the Fermi scale then Higgs boson mass gets stabilized against unnatural quantum fluctuations. Furthermore, supersymmetric nature of interactions (i) predicts the unification of three gauge forces near the scale of strong gravitational interactions, (ii) realizes radiative electroweak breaking, and (iii) provides a viable dark matter candidate. In spite of these rather appealing aspects, however, supersymmetry does not offer a solution to the CC problem as it cannot be a good symmetry of Nature below the electroweak scale. In fact, a solution to the CC problem requires ‘new physics’ to show up at a scale around the neutrino mass. Given these features of the two naturalness problems, a simultaneous and unified resolution, as recently proposed by Arkani-Hamed and Dimopoulos [1], might be that Nature is inherently fine-tuned. In other words, Nature might admit small numbers $\epsilon_{CC}$ and $\epsilon_h$ such that

$$\Lambda \sim \epsilon_{CC} M_{SUSY}^4, \quad m_h^2 \sim \epsilon_h M_{SUSY}^2$$

where the scale of supersymmetry breaking, $M_{SUSY}$, can be well inside the Planckian territory or in some intermediate domain depending on what mechanism is responsible for the breaking. The smallness of $\epsilon_{CC,h}$ is a measure of the UV-sensitivity of the quantity under concern: for $M_{SUSY} \sim M_{Pl}$ one has $\epsilon_{CC} \sim 10^{-120}$ and $\epsilon_h \sim 10^{-32}$. In this picture it is the presence of fine-tunings, rather than the renormalization group flow of masses, that provides a SM-like Higgs doublet to condense for breaking the gauge symmetry. Highly appealing aspects of the supersymmetric models, unification of the gauge couplings and presence of a dark matter candidate, are still maintained thanks to the chiral protection of gaugino and Higgsino masses [2, 3]. In the minimal supersymmetric model (MSSM), for instance, the $\mu$ parameter, gaugino masses and triscalar couplings all explicitly break the continuous $R$ invariance of the model (because of which they are endowed with CP–odd phases).

In the following we will focus on the scalar fermion sector of the MSSM, endowed with mass hierarchies pertinent to split supersymmetry [1, 4], and show that part of the sfermions necessarily weigh at the weak scale if (i) the sfermion mass-squared matrices exhibit flavor mixings and (ii) extreme fine-tunings (different than but similar in size to that needed to induce a light Higgs doublet) are allowed. The motivations for such an analysis are twofold: (i) the mechanism that breaks the supersymmetry does not need to be flavor-blind, and (ii) there are no chiral symmetries that can protect any entry of a generic sfermion mass matrix. The detailed discussions below will give rise to the conclusion that, non-observation of light sfermion states at the LHC will guarantee approximate flavor-blindness of supersymmetry breaking.
It is useful to start our analysis with a brief review of the main points made in [1]. The superpotential of the MSSM

$$\hat{W} = \mu \hat{H}_u \cdot \hat{H}_d + \hat{Q} \cdot \hat{\bar{H}}_u Y_u \hat{U}^c + \hat{\bar{H}}_d \cdot \hat{Q} Y_d \hat{D}^c + \hat{\bar{H}}_d \cdot \bar{L} Y_e \hat{E}^c$$

(2)

encodes the rigid parameters of the model: the Higgsino mass $\mu$ and the Yukawa matrices $Y_{u,d,e}$ of up quarks, down quarks and charged leptons, respectively. The $F$ terms of the Higgs doublets plus their soft-breaking terms induce the scalar potential

$$V_{\text{Higgs}} = m_H^2 H_u^H H_u + m_H^2 H_d^H H_d + m_{ud}^2 (H_u \cdot H_d + \text{h.c.}) + \text{D terms}$$

(3)

where the soft bilinear coupling $m_{ud}^2$, which can always be made real by a rephasing of $H_u,d$, mixes the two doublets. The mass splitting between the light and heavy physical Higgs doublets depends on the strength of this mixing. In fact, the light Higgs doublet, $h$, acquires a mass

$$2m_h^2 = m_{H_u}^2 + m_{H_d}^2 - \sqrt{(m_{H_u}^2 + m_{H_d}^2)^2 + 4(m_{ud}^4 - m_{H_u}^2 m_{H_d}^2)}$$

(4)

which can be forced to lie right at the electroweak scale under the admitted fine-tunings (1). Indeed, $h$ mimics the SM Higgs doublet [1, 4] if

$$m_{ud}^4 - m_{H_u}^2 m_{H_d}^2 \approx \text{TeV}^2 \left( m_{H_u}^2 + m_{H_d}^2 \right)$$

(5)

where TeV stands for the Fermi scale. Consequently, when the masses and mixings of the original Higgs doublets are fine-tuned with an accuracy $\epsilon_h \sim \text{TeV}^2/M_{\text{SUSY}}^2$ a light tachyonic Higgs doublet emerges automatically. The heavy Higgs doublet weighs $O(M_{\text{SUSY}})$ and its effect on the infrared (IR) dynamics is highly suppressed.

In the electroweak vacuum, $\langle h^0 \rangle \simeq m_{\text{top}}$, all quarks and charged leptons acquire masses:

$$m_u = \sin \beta \langle h^0 \rangle Y_u, \ m_{d,e} = \cos \beta \langle h^0 \rangle Y_{d,e}$$

where $\beta$ being the Higgs mixing angle [1, 4]. The unitary rotations of the superfields

$$\hat{Q} \rightarrow \left( \begin{array}{cc} V_{Q_u} & 0 \\ 0 & V_{Q_D} \end{array} \right) \hat{Q}, \ \hat{L} \rightarrow V_L \hat{L}, \ \hat{U}^c \rightarrow V_U \hat{U}^c, \ \hat{D}^c \rightarrow V_D \hat{D}^c, \ \hat{E}^c \rightarrow V_E \hat{E}^c$$

(6)

subject to the constraints

$$V_{Q_u}^{T} Y_u V_U = Y_u, \ V_{Q_D}^{T} Y_d V_D = Y_d, \ V_{L_E}^{T} Y_e V_E = Y_e$$

(7)

project all fermions into their physical bases via strictly diagonal $Y_{u,d,e}$. Under these transformations, the neutral current vertices remain flavor-diagonal as in the gauge basis whereas charged current vertices of quarks shuffle different flavors via the CKM matrix $V_{\text{CKM}} = V_{Q_u}^{T} V_{Q_D}$. This very basis, the super-CKM basis, is highly useful for analyzing the supersymmetric flavor mixings as additional sources with respect to the standard flavor violation. Indeed, under (6), the soft mass-squareds $M^2_{Q,d,e}$ of $\hat{Q}, \ldots, \hat{E}^c$ as well as their triscalar couplings $Y^A_{u,d,e}$ get dressed by the associated unitary matrices. These dressings modify the
gauge-basis flavor structures significantly. For instance, the left-chirality squarks \( \tilde{Q}_u \equiv \tilde{U}_L \) and \( \tilde{Q}_{d\downarrow} \equiv \tilde{D}_L \) acquire distinct masses \( M_{U_L}^2 = V_{Q_u}^T M_Q^2 V_{Q_u}^* \) and \( M_{D_L}^2 = V_{Q_d}^T M_Q^2 V_{Q_d}^* \) even if \( M_Q^2 \) is strictly diagonal. Furthermore, couplings of neutral and charged components of Higgs fields differ \( e.g., H_u^0 (H_d^+) \) couples to \( \tilde{U}_L \tilde{U}^c (\tilde{D}_L \tilde{D}^c) \) with \( V_{Q_u}^T Y_u^A V_U \equiv Y_u^A \left( V_{Q_d}^T Y_u^A V_U \equiv \tilde{Y}_u^A \right) \). Consequently, corresponding to superpotential (2), in super-CKM basis, the most general holomorphic and \( R \) parity conserving operator structures parameterizing soft supersymmetry breaking in the scalar fermion sector are given by

\[
\begin{align*}
\tilde{U}_L^T M_{U_L}^2 \tilde{U}_L^* + \tilde{D}_L^T M_{D_L}^2 \tilde{D}_L^* + \tilde{U}^c M_{UR}^2 \tilde{U}^c + \tilde{D}^c M_{DR}^2 \tilde{D}^c + \tilde{L}^T M_{LL}^2 \tilde{L}^* + \tilde{E}^c M_{LR}^2 \tilde{E}^c + [H_u^0 \tilde{U}_L Y_u^A \tilde{U}^c + H_d^0 \tilde{D}_L Y_d^A \tilde{D}^c + H_\mu \tilde{L}_L Y_e^A \tilde{E}^c + h.c.] \end{align*}
\]

(8)

where only the neutral Higgs couplings are displayed. The soft-breaking masses \( M_{U_L}^2, \ldots, L_R^2 \) and \( Y_{u,d,e}^A \) are \( 3 \times 3 \) matrices in flavor space. The mass matrices are obviously hermitian. They conserve parity and contribute to both CP and flavor violations via their off-diagonal entries. On the other hand, the triscalar couplings \( Y_{u,d,e}^A \) like the Yukawas themselves, are general parity-violating non-hermitian matrices. These flavor structures are generic; neither their textures nor their correlations are known \textit{a priori}. They are completely independent of each other; one’s texture is independent of another’s. The gauge-basis soft mass-squareds \( M_Q^2, \ldots, E \) break no symmetry but supersymmetry; therefore, there is no smallness argument for any of their entries due to the absence of chiral protection. On the other hand, \( Y_{u,d,e}^A \) explicitly break the continuous \( R \) invariance of the theory hence their chiral protection.

Given that the mechanism which breaks the supersymmetry is not known, use of the experimental bounds is the only way for reconstructing the moduli and phases of \( M_{U_L}^2, \ldots, L_R^2 \) and \( Y_{u,d,e}^A \). Indeed, if supersymmetry survives down to the Fermi scale then these flavor matrices induce sizeable modifications in rare processes (scaling as \( 1/M_{\text{SUSY}} \) to appropriate power [5]) as well as in Higgs-fermion couplings (scaling as \( |\mu||\text{triscalars}|/M_{\text{SUSY}}^2 \) or \( |\mu||\text{gaugino}|/M_{\text{SUSY}}^2 \) [6]). Consequently, if supersymmetry is not a weak-scale invariance of Nature then constraints from the former fade away whereas those from the latter remain intact. In the framework of [1], however, Higgs-fermion couplings, too, turn out to be completely insensitive to supersymmetric flavor structures due to the chiral protection of dimension-3 soft-breaking terms. Irrespective of the scale of supersymmetry breaking, vacuum stability arguments [7] impose rather strong constraints on the off-diagonal entries of triscalar couplings; however, such bounds are automatically satisfied in the framework of [1]. In light of these observations one concludes that flavor mixings in the sfermion sector are rather generic and remain unconstrained by phenomenological bounds.

Having summarized the status of flavor violation in low-scale supersymmetry, we now discuss implications of flavor mixings in \( M_{U_L}^2, \ldots, L_R^2 \) for sparticle spectrum when fine-tunings \( \mathcal{O}(\epsilon_h) \) are built-in properties of Nature. It is convenient to analyze first a rather plain flavor
structure by specializing to one of the mass-squareds e.g. $M^2_{UL}$:

$$
M^2_{UL} = \begin{pmatrix}
m_1^2 & 0 & 0 \\
0 & m_2^2 & m_{23}^2 \\
0 & m_{23}^2 & m_3^2 
\end{pmatrix}
$$

(9)

with two texture zeroes and $m_{1,2,3}^2 \sim |m_{23}^2| \sim M_{SUSY}^2$. That the second and third generations exhibit such a strong mixing implies that the mass

$$
2m_{u_L}^2 = m_2^2 + m_3^2 - \sqrt{(m_2^2 + m_3^2)^2 + 4 \left(|m_{23}^2|^2 - m_2^2 m_3^2\right)}
$$

(10)

of the lightest flavor eigenstate, with $\gamma_{u_L} = \text{Arg}[m_{23}^2]$, 

$$
\tilde{u}_L \equiv \cos \theta_{u_L} \tilde{U}_{L2} + \sin \theta_{u_L} e^{-i\gamma_{u_L}} \tilde{U}_{L3}
$$

(11)

can be fine-tuned to lie right at the weak scale if

$$
m_3^2 m_3^2 - |m_{23}^2|^2 \simeq \text{TeV}^2 \left(m_2^2 + m_3^2\right)
$$

(12)

holds. The amount of fine-tuning involved here is of the same size as the one needed for generating a light Higgs doublet. The weak-scale effective theory is obtained by the replacements

$$
\tilde{U}_{L2} \rightarrow \cos \theta_{u_L} \tilde{u}_L , \quad \tilde{U}_{L3} \rightarrow \sin \theta_{u_L} e^{i\gamma_{u_L}} \tilde{u}_L
$$

(13)

in the lagrangian. Clearly, dressing of the interaction vertices by $\cos \theta_{u_L}$ or $\sin \theta_{u_L}$ signals the decoupling of heavy squark up squark from the light spectrum.

In general, $M^2_{UL}, \ldots, M^2_{LR}$ are independent matrices, and thus, none, part or all of them can exhibit a structure similar to (9). In case each of them accommodates a strong mixing between any two flavors then the weak-scale effective theory consists of seven distinct sfermion states within the reach of LHC. Of course, all these light sfermions are only flavor eigenstates; whether they are physical fields or not depend on if triscalar couplings and $F$ terms feed further mixings. To see this point in detail, suppose that intergenerational mixings are all small except that $M^2_{UL}$ and $M^2_{UR}$ possess sizeable $(2,3)$ and $(1,2)$ entries, respectively. Then, the light sparticle spectrum consists of two up squarks $\tilde{u}_L$ and $\tilde{u}_R$ of opposite chirality and a single Higgs doublet $h$ in excess of gauginos and Higgsinos. An interesting aspect of this spectrum is that squarks develop left-right mixings via the triscalar couplings in (8) and $F$ term contributions. Indeed, in $(\tilde{u}_L, \tilde{u}_R)$ basis the mass-squared matrix of up squarks takes the form

$$
M^2_u = \begin{pmatrix}
m_{u_L}^2 + \Delta_L \langle h^0 \rangle^2 & \sin \beta \langle h^0 \rangle m_{LR} \\
\sin \beta \langle h^0 \rangle m_{LR}^* & m_{u_R}^2 + \Delta_R \langle h^0 \rangle^2
\end{pmatrix}
$$

(14)
where

\[
\begin{align*}
\Delta_L &= \cos^2 \beta \left| (Y_u)_{22} \cos \theta_{uL} + (Y_u)_{33} \sin \theta_{uL} e^{i\gamma_{uL}} \right|^2 + \frac{1}{4} \cos 2\beta \left( g_2^2 - \frac{1}{3} g_1^2 \right) \\
\Delta_R &= \cos^2 \beta \left| (Y_u)_{11} \cos \theta_{uR} + (Y_u)_{22} \sin \theta_{uR} e^{i\gamma_{uR}} \right|^2 + \frac{1}{3} \cos 2\beta g_1^2
\end{align*}
\]

are generated by sfermion \( F \) and \( D \) terms quadratic in \( \bar{u}_{L,R} \), and

\[
m_{LR} = \left( Y^\Lambda_u \right)_{31} \sin \theta_{uL} \cos \theta_{uR} e^{i\gamma_{uL}} + \left( Y^\Lambda_u \right)_{32} \sin \theta_{uL} \sin \theta_{uR} e^{i(\gamma_{uL} + \gamma_{uR})}
+ \left( Y^\Lambda_u \right)_{21} \cos \theta_{uL} \cos \theta_{uR} + \left( Y^\Lambda_u - \mu^* \cot \beta \gamma_u \right)_{22} \cos \theta_{uL} \sin \theta_{uR} e^{i\gamma_{uR}}
\]

is induced by the \( F \) terms of Higgs fields and triscalar couplings in (8). This quantity receives contributions from all relevant entries of \( Y^\Lambda_u \) not just from its \((2,2)\) corner, and it comprises phases of both large and small scales of the model. The physical up squark states are achieved by diagonalizing (14); its mixing matrix will involve \( \text{Arg}[m_{LR}] \) to contribute to CP violation in processes proceeding with \( \bar{u}_{L,R} \) mediation.

Going to flavor-diagonal basis directly modifies the supergauge vertices even if squarks do not exhibit any LR mixing. Indeed, the light squarks \( \bar{u}_{L,R} \) above interact with gluino and quarks via

\[
\cos \theta_{uL} \bar{u}_L g_R \lambda^a c_L + \sin \theta_{uL} e^{-i\gamma_{uL}} \bar{u}_L g_R \lambda^a t_L - \cos \theta_{uR} \bar{u}_R g_L \lambda^a u_R - \sin \theta_{uR} e^{-i\gamma_{uR}} \bar{u}_R g_L \lambda^a e_R
\]

in units of \(-g_s/\sqrt{2}\). It is highly remarkable that a given light squark couples to distinct quark fields: a crucial property for flavor-changing neutral current (FCNC) transitions. In general, gaugino-fermion-sfermion couplings of this form generically exist for down quark and lepton fields: a crucial property for flavor-changing neutral current (FCNC) transitions. In general, gaugino-fermion-sfermion couplings of this form generically exist for down quark and lepton fields, too, provided that their mass-squared matrices exhibit two texture zeroes like (9).

The results derived above are not special to the mixing pattern in (9). Indeed, mass-squared matrices having a single texture zero

\[
\left( \begin{array}{ccc} \# & 0 & \# \\ 0 & \# & \# \\ \# & \# & \# \end{array} \right), \quad \left( \begin{array}{ccc} \# & \# & 0 \\ \# & \# & \# \\ 0 & \# & \# \end{array} \right), \quad \left( \begin{array}{ccc} \# & \# & \# \\ \# & \# & \# \\ \# & \# & \# \end{array} \right)
\]

(18)

with all \#’s being \( \mathcal{O}(M_{\text{SUSY}}) \), also possess one small eigenvalue \( \mathcal{O}(\text{TeV}) \) and two large ones \( \mathcal{O}(M_{\text{SUSY}}) \) provided that fine-tunings \( \mathcal{O}(\epsilon_h) \) are admitted. For such flavor structures, all three generations of sfermions of a given chirality possess a light component, and that light sfermion interacts with gauginos and all three generations of fermions in a way generalizing (17). Furthermore, if at least one of \( (M^2_{uL}, M^2_{uR}), (M^2_{D_L}, M^2_{D_R}), (M^2_{L_L}, M^2_{L_R}) \) exhibit flavor mixings as in (18) then the resulting opposite-chirality light sfermions develop left-right mixings such that \( m_{LR} \) now involves all entries of triscalar couplings.

In addition to (9) and (18), the sfermion mass-squared matrices can assume a democratic texture as well

\[
\left( \begin{array}{ccc} \# & \# & \# \\ \# & \# & \# \\ \# & \# & \# \end{array} \right)
\]

(19)
which accommodates a heavy state weighing $\mathcal{O}(M_{\text{SUSY}})$ and two light states with masses $\mathcal{O}(\text{TeV})$. If all of $M^2_{U_L}, \ldots, M^2_{U_R}$ exhibit democratic flavor mixings then light spectrum contains a total of fourteen sfermion states. In case, for instance, $M^2_{U_L}$ and $M^2_{U_R}$ are of the form (19) then there will be four light up squark states and they will experience both intra- and inter-generational mixings in the left-right block in a way involving all entries of $\mathbf{Y}_u^A$. Furthermore, each of these light up squarks will interact with all three up quarks via supergauge vertices.

Having derived the light sparticle spectrum arising from $\mathcal{O}(M_{\text{SUSY}})$ flavor mixings, we now discuss their phenomenological implications:

- Even if there exists a single light sfermion flavor, fermion-sfermion-gaugino vertices necessarily lead to FCNC transitions. This is clear from (17) in which a light sfermion develops interactions with fermions of varying flavor. Consider, for instance, the last two terms of (17) which involve couplings of right-handed $u$ and $c$ quarks to a single up type squark. Obviously, these vertices generate supersymmetric contributions to rare processes involving $u_R-c_R$ transition. For instance, the effective Hamiltonian for $D^0-\overline{D}^0$ mixing receives the contribution

$$\Delta H_{\text{eff}} = \alpha_s^2 \cos^2 \theta_{u_R} \sin^2 \theta_{u_R} e^{2\gamma_{u_R}} \frac{1}{|M_\tilde{g}|^2} \left( \frac{x + 9}{18(1-x)^2} + \frac{3x + 2}{9(1-x)^3} \ln x \right) (\overline{c_R} \gamma_{\mu} u_R)^2$$

which contributes to both the lifetime and CP impurity of the neutral $D$ meson. This operator can lead to significant effects depending on the sizes of $x = m^2_{\tilde{u}_R}/|M_\tilde{g}|^2$ and the gluino mass so that the present bounds can put stringent limits on $\sin 2\theta_{u_R}$ [5]. Obviously, in case the associated mass-squared matrices admit light states the structure above repeats for other meson mixings as well. For example, if $M^2_{D_L}$ exhibits a texture like (9) then $B^0_s-\overline{B}^0_s$ mixing receives a contribution similar to (20) with the replacements $\theta_{u_R} \rightarrow \theta_{d_L}$, $\gamma_{u_R} \rightarrow \gamma_{d_L}$, $m_{\tilde{u}_R} \rightarrow m_{\tilde{d}_L}$ and $\overline{c_R} \gamma_{\mu} u_R \rightarrow \overline{b_L} \gamma_{\mu} s_L$. In addition to mixings of mesons, their decays are also influenced by light sparticle spectrum since, for instance, the first two terms of (17) give rise to the rare decays $t_L \rightarrow c_L (\gamma, g)$.

Depending on how crowded the light spectrum is, rare processes could be mediated by various sfermions. For instance, when there are two squarks experiencing LR mixings as in (14) then meson mixings receive contributions from $(V-A) \otimes (V-A)$, $(V-A) \otimes (V+A)$ and $(V-A) \otimes (V+A)$ type $\Delta F = 2$ operators. Furthermore, the flavor textures (18) and (19) enable light sfermions to couple to all three fermions of given electric charge and thus they contribute to mixings and decays of more than one meson such that the sfermion line can flip both chirality and flavor [8].

- The very existence of light sfermion states necessitates radiative corrections to Higgs boson mass beyond those in the SM. This statement is valid even if there exists a single sfermion field at the TeV scale. For exemplifying this point, consider a light $\tilde{u}_L$
springing from (9). At one loop this state shifts the Higgs boson mass by an amount

\[ \delta m_h^2 = \frac{3}{8\pi^2} \Delta_L \left\{ \left( m_{u_i}^2 + 2\Delta_L \langle h^0 \rangle^2 \right) \ln \frac{m_{\tilde{u}_i}^2}{Q^2} - m_{\tilde{u}_i}^2 \right\} \] (21)

at the renormalization scale \( Q \), in the \( \overline{\text{MS}} \) scheme, with physical squark mass \( m_{\tilde{u}_i}^2 = m_{u_i}^2 + \Delta_L \langle h^0 \rangle^2 \). It is straightforward to generalize (21) to cases with several sfermions of varying flavor and chirality. Clearly, unlike the conventional MSSM [9], the Higgs sector does not violate CP invariance.

- The electric dipole moments (EDM) are no exception; they are induced already at one-loop level in the presence of light sfermions. Indeed, the first and last terms of (17), for instance, induce a finite EDM for charm quark via gluino-squark exchange:

\[ \frac{D_c}{e} = -\frac{4\alpha_s}{9\pi} \frac{|m_{LR}| \langle h^0 \rangle}{M_g^3} \cos \theta_{u_L} \sin \theta_{u_R} \sin \beta \sin \left( \gamma_{LR} + \gamma_{u_R} - \gamma_{\tilde{g}} \right) \]

\[ \times \frac{x_1 x_2}{x_1 - x_2} \left( f(x_1) - f(x_2) \right) \] (22)

where \( m_{u_i}^2 \), with \( i = 1, 2 \), are the eigenvalues of (14), \( x_i = m_{u_i}^2 / |M_g|^2 \) and \( f(x) = (x/2(1 - x)^3) \left( 1 - x^2 + 2x \ln x \right) \). The EDMs, though genuinely flavor-diagonal, gain a direct sensitivity to the flavor mixings in triscalar couplings via its dependence on the modulus and phase of \( m_{LR} \). Clearly, as the light sfermions populate in flavor and chirality, the EDMs of fermions are contributed by a variety of diagrams [10].

- The question of if and how gauge couplings unify depends on the light spectrum and on the scale at which heavy spectrum lies. For the model under concern, the light spectrum is composed of SM fermions (which come in complete SU(5) irreps), the gauginos and Higgsinos, a light Higgs doublet and a set of light sfermions. Therefore, the one-loop beta function coefficients for SU(3)_c, SU(2)_L and U(1)_Y are given, respectively, by \( b_3 = -5 + \Delta_3/3 \), \( b_2 = -7/6 + \Delta_2/3 \) and \( b_1 = 9/2 + \Delta_1/3 \) where \( \Delta_i \) represent the contributions of light sfermions (they all vanish in the ‘minimal’ split supersymmetry [1, 4]). A straightforward calculation gives \( (\Delta_3, \Delta_2, \Delta_1) = (1/2, 1/2, 1/20), (1/2, 0, 4/5), (1/2, 0, 1/5), (0, 1/2, 3/20), (0, 0, 3/5) \) for a light \( \tilde{u}_L, \tilde{u}_R, \tilde{d}_R, \tilde{e}_L \) and \( \tilde{e}_R \), respectively. Consequently, each gauge coupling runs with a different slope depending on the type of the light sfermion field. For recovering complete gauge coupling unification, the MSSM unification scale must lie above the scale of heavy scalars between which the running is precisely that of the MSSM [1, 4].

- The light spectrum will leave observable signatures at the LHC; however, unlike the predictions of flavor-conserving split supersymmetry [1], gluino does not need to be a long-lived fermion; moreover, LHC does not need to be a gluino factory; such statements are in variant with the mass spectra of inos and sfermions. The LHC will
produce sfermions as well though they differ from the conventional MSSM spectrum in population and couplings. It is interesting to compare predictions of flavor-blind and flavor-violating split-supersymmetric models in connection with [11]. Indeed, in this work it was found that long-lived gluinos, produced in distant galactic nuclei, induce cosmic ray showers with rates detectable in upcoming Pierre Auger Observatory. Therefore, in case such signals are experimentally confirmed, it will be possible to conclude that the mechanism that breaks supersymmetry is flavor-blind. In the opposite case, we will have to live with light sparticles which can facilitate a number of processes listed above.

By explicit examples ranging from rare processes to collider expectations, we have briefly discussed certain phenomenological consequences of supersymmetric models endowed with ultra heavy scalars and nontrivial flavor mixings when Nature admits fine-tunings in (1). Our findings show that, flavor mixings give rise to light sfermion states whose contributions to various observables can be bounded even with the present level of experimental exploration. The TeV scale effective theory has interesting differences from the conventional MSSM in terms of the couplings, masses and CP properties of the particles. It is after a global analysis of the existing bounds on FCNCs, Higgs mass, EDMs, . . . that one can determine the likelihood ranges of various model parameters.

In conclusion, if the supersymmetric spectrum is to be completely split, that is, if gauginos, Higgsinos and a Higgs doublet are to be the only sparticles to weigh at the Fermi scale then one must prevent flavor mixings among the sfermions at $M_{SUSY}$ in addition to implementing chiral symmetries for protecting dimension-3 soft terms. Recently, the work of Arkani-Hamed and Dimopoulos has been rectified by constructing explicit models [12] in the spirit of [3]. It has there been shown that, spurion fields of the form $\hat{X} = 1 + \theta^4 M_{SUSY}^2$ break supersymmetry but preserve the $R$ invariance in accord with the requirements of splitting the sparticle spectrum. At the renormalizable level, such spurion fields couple to the visible sector fields either as $\int d^4\theta \hat{Q}_i^\dagger \hat{X}_{ij} \hat{Q}_j$ or as $\int d^4\theta H^u \hat{X}_{ud} H_d$ where $\hat{X}_{ij,ud}$ are dimension-less superfields. These interactions generate the soft-mass squareds $(M^2_Q)_{ij}$ and $m^2_{ud}$, all being naturally $O(M_{SUSY})^2$. The textures of the sfermion mass-squareds is determined by $D$ components of $\hat{X}_{ij}$; in case spurions with $i \neq j$ are sufficiently suppressed compared to those with $i = j$ then weak-scale effective theory is precisely the one put forward in [1]. However, in the absence of a symmetry principle that can impose such a hierarchy it will not be possible to clean up the Fermi scale from sfermions. A candidate symmetry would be the invariance of entire sfermion sector (excluding the triscalar couplings) under unitary rotations of the form $\tilde{f} \to U \tilde{f}$ with $UU^\dagger = 1$. Indeed, such an invariance would be operative if and only if the gauge-basis soft mass-squareds are proportional to the identity matrix. Putting differently, the sfermion mass-squareds must remain unchanged as one switches from gauge basis to super-CKM basis. Clearly, this symmetry does not need to be exact; it can be an approximate symmetry provided that off-diagonal entries of the sfermion mass-squared matrices are not comparable with those at the diagonal.
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