I. NEUTRALINOS AND CHARGINOS

The higgsinos and the electroweak gauginos mix with each other because of the effects of electroweak symmetry breaking. The neutral higgsinos ($\tilde{H}_u^0$ and $\tilde{H}_d^0$) and the neutral gauginos ($\tilde{B}$ and $\tilde{W}_0^0$) combine to form the mass eigenstates called neutralinos. The charged higgsinos ($\tilde{H}_u^+ + \tilde{H}_d^-$) and winos ($\tilde{W}_+^0$ and $\tilde{W}_-^0$) mix to form the mass eigenstates with charge $\pm 1$ called charginos. We denote the neutralinos mass eigenstates by $\tilde{N}_i$ with $i = 1, 2, 3, 4$. We denote the chargino mass eigenstates by $\tilde{C}_i^\pm$ with $i = 1, 2$. (Another common notation is $\tilde{\chi}_0^i$ and $\tilde{\chi}_i^\pm$.) By convention, these mass eigenstates are labeled in ascending mass order, so that $m_{\tilde{N}_1} < m_{\tilde{N}_2} < m_{\tilde{N}_3} < m_{\tilde{N}_4}$ and $m_{\tilde{C}_1} < m_{\tilde{C}_2}$. The lightest neutralino, $\tilde{N}_1$, is often assumed to be the LSP, unless there is a lighter gravitino, or unless $R$-parity is not conserved, because it is the only MSSM particle that can make a good dark matter candidate. In this section, we describe the mass spectrum and mixing of neutralinos and charginos in the MSSM.

In the interaction eigenstate basis, $\psi^0 = (\tilde{B}, \tilde{W}_0^0, \tilde{H}_d^0, \tilde{H}_u^0)$, the neutralino-mass part of the Lagrangian reads

$$L_{\tilde{N}} = -\frac{1}{2} (\psi^0)^T M_{\tilde{N}} \psi^0 + \text{h.c.},$$

where

$$M_{\tilde{N}} = \begin{pmatrix} M_1 & 0 & -g' v_d / \sqrt{2} & g' v_u / \sqrt{2} \\ 0 & M_2 & g v_d / \sqrt{2} & -g v_u / \sqrt{2} \\ -g' v_d / \sqrt{2} & g v_d / \sqrt{2} & 0 & -\mu \\ g' v_u / \sqrt{2} & -g v_u / \sqrt{2} & -\mu & 0 \end{pmatrix}. \tag{1}$$

The entries $M_1$ and $M_2$ come directly from the MSSM soft Lagrangian, while $\mu$ comes from the superpotential. Electroweak symmetry breaking generates the gaugino-higgsino mixing, when the Higgs scalar is replaced by its VEV in the Higgs-higgsino-gaugino coupling. Eq. (2) can be rewritten as

$$M_{\tilde{N}} = \begin{pmatrix} M_1 & 0 & -c_\beta s_W m_Z & s_\beta s_W m_Z \\ 0 & M_2 & c_\beta c_W m_Z & -s_\beta c_W m_Z \\ -c_\beta s_W m_Z & c_\beta c_W m_Z & 0 & -\mu \\ s_\beta s_W m_Z & -s_\beta c_W m_Z & -\mu & 0 \end{pmatrix}, \tag{3}$$

where $s_\beta = \sin \beta$, $c_\beta = \cos \beta$, $s_W = \sin \theta_W$, and $c_W = \cos \theta_W$. The mass matrix can be diagonalized by a unitary matrix $V_N$ to obtain mass eigenstates:

$$\tilde{N}_i = (V_N)_{ij} \psi^0_j, \tag{4}$$
so that

\[ V_N^* M_N V_N^{-1} = \text{diag}(m_{\tilde{N}_1}, m_{\tilde{N}_2}, m_{\tilde{N}_3}, m_{\tilde{N}_4}) \]

(5)

has real positive entries on the diagonal.

In general, the parameters \( M_1, M_2 \) and \( \mu \) can have arbitrary complex phases. A redefinition of the phases of \( \tilde{B} \) and \( \tilde{W} \) always allows us to choose a convention in which \( M_1 \) and \( M_2 \) are both real and positive. The phase of \( \mu \) within that convention is then really a physical parameter and cannot be rotated away. (We have already used the freedom to redefine the phases of the Higgs fields to make \( b, \langle H^0_u \rangle \) and \( \langle H^0_d \rangle \) real and positive.) However, if the phase of \( \mu \) is different from zero and, more precisely, of order one, there can be potentially disastrous predictions for CP violation, particularly unacceptably large EDMs for the electron and the neutron. (This situation can, in principle, be avoided by cancelations of contributions involves \( a \)-terms against those involving gaugino masses.) Therefore, it is often assumed that \( \mu \) is real in the same set of conventions where \( M_1, M_2, b, \langle H^0_u \rangle, \langle H^0_d \rangle \) are real. The sign of \( \mu \) is still undetermined by this constraint. We would like to stress that imposing the simultaneous vanishing (or smallness) of these CP violating phases is an ad-hoc requirement, meant the render the MSSM phenomenologically viable, not a generic prediction.

In models that satisfy

\[ \frac{M_1}{g_1^2} = \frac{M_2}{g_2^2} = \frac{M_3}{g_3^2} = \frac{m_1/2}{g_U^2}, \]

(6)

one has the prediction

\[ M_1 \approx \frac{5}{3} \tan^2 \theta_W M_2 \approx 0.5 M_2 \]

(7)

at the electroweak scale. If so, then the neutralino masses and mixing angles depend on only three unknown parameters. This assumption has been made in most phenomenological studies; nevertheless, it is an assumption, that can (and should) be experimentally tested.

There is a not-unlikely limit in which the electroweak symmetry breaking effects can be viewed as a small perturbation on the neutralino mass matrix. If

\[ m_Z \ll |\mu + M_1|, |\mu + M_2|, \]

(8)

then the neutralino mass eigenstates are very nearly a “bino-like” \( \tilde{N}_1 \approx \tilde{B} \); a “wino-like” \( \tilde{N}_2 \approx \tilde{W}^0 \); and a “higgsino-like” \( \tilde{N}_3, \tilde{N}_4 \approx (\tilde{H}^0_u \pm \tilde{H}^0_d)/\sqrt{2} \), with mass eigenvalues

\[ m_{\tilde{N}_1} = M_1 - \frac{m_Z^2 s_W^2 (M_1 + \mu \sin 2\beta)}{\mu^2 - M_1^2} + \cdots, \]

(9)
\[ m_{\tilde{N}_2} = M_2 - \frac{m_W^2(M_2 + \mu \sin 2\beta)}{\mu^2 - M_2^2} + \cdots, \]  
\[ m_{\tilde{N}_{3,4}} = |\mu| + \frac{m_W^2(I \mp \sin 2\beta)(|\mu| \pm M_1 c_W \pm M_2 s_W)}{2(\mu \pm M_1)(\mu \pm M_2)} + \cdots, \]

where we take \( M_1, M_2 \) real and positive, and \( \mu \) real and with sign \( I = \pm 1 \). The subscripts of the mass eigenstates may need to be rearranged depending on the numerical values of the parameters. In particular, the above ordering assumes \( M_1 < M_2 \ll |\mu| \).

The chargino spectrum can be analyzed in a similar way. In the interaction eigenstate basis \( \psi^\pm = (\tilde{W}^+\tilde{H}_u^+, \tilde{W}^-\tilde{H}_d^-) \), the chargino mass terms in the Lagrangian are

\[ \mathcal{L}_{M_C} = -\frac{1}{2}(\psi^\pm)^T M_C \psi^\pm + \text{h.c.}, \]  

where, in a \( 2 \times 2 \) block form,

\[ M_C = \begin{pmatrix} 0 & X_C^T \\ X_C & 0 \end{pmatrix}, \]

with

\[ X_C = \begin{pmatrix} M_2 & gv_u \\ gv_d & \mu \end{pmatrix} = \begin{pmatrix} M_2 & \sqrt{2}s_\beta m_W \\ \sqrt{2}c_\beta m_W & \mu \end{pmatrix}. \]  

The mass eigenstates are related to the interaction eigenstates by two unitary \( 2 \times 2 \) matrices \( V_+ \) and \( V_- \) according to

\[ \begin{pmatrix} \tilde{C}_1^+ \\ \tilde{C}_2^+ \end{pmatrix} = V_+ \begin{pmatrix} \tilde{W}^+ \\ \tilde{H}_u^+ \end{pmatrix}, \quad \begin{pmatrix} \tilde{C}_1^- \\ \tilde{C}_2^- \end{pmatrix} = V_- \begin{pmatrix} \tilde{W}^- \\ \tilde{H}_d^- \end{pmatrix}. \]  

Note that the mixing matrix for the positively charged left-handed fermions is different from that for the negatively charged left-handed fermions. They are chosen so that

\[ V_-X_CV_+^{-1} = \text{diag}(m_{\tilde{C}_1}, m_{\tilde{C}_2}), \]  

with positive real entries \( m_{\tilde{C}_i} \). Because these are \( 2 \times 2 \) matrices, it is not difficult to solve for the masses explicitly:

\[ m_{\tilde{C}_{1,2}}^2 = \frac{1}{2} \left[ |M_2|^2 + |\mu|^2 + 2m_W^2 \mp \sqrt{(|M_2|^2 + |\mu|^2 + 2m_W^2)^2 - 4|\mu|M_2 - m_W^2 \sin 2\beta|^2} \right]. \]  

These are the (doubly degenerate) eigenvalues of the \( 4 \times 4 \) matrix \( M_C^T M_C \), or, equivalently, the eigenvalues of \( X_C^T X_C \). In the limit of Eq. (8), with real \( M_2 \) and \( \mu \), the chargino mass eigenstates consist of wino-like \( \tilde{C}_1^\pm \) and a higgsino-like \( \tilde{C}_2^\pm \), with masses

\[ m_{\tilde{C}_1} = M_2 - \frac{m_W^2(M_2 + \mu \sin 2\beta)}{\mu^2 - M_2^2} + \cdots, \]  
\[ m_{\tilde{C}_2} = |\mu| - \frac{m_W^2 I(\mu + M_2 \sin 2\beta)}{\mu^2 - M_2^2} + \cdots. \]
The labeling here, again, corresponds to $M_2 < |\mu|$, and $I = \text{sign}(\mu)$. Note that, in the limit that electroweak symmetry breaking effects are somewhat smaller than the electroweak conserving ones, the charged $\tilde{C}_1$ and the neutral $\tilde{N}_2$ winos are nearly degenerate, with mass of order $M_2$. Similarly, the higgsinos $\tilde{C}_3, \tilde{N}_3$ and $\tilde{N}_4$ are nearly degenerate, with mass of order $|\mu|$.

The Feynman rules involving neutralinos and charginos may be inferred in terms of $V_N, V_+, V_-$ from the MSSM Lagrangian.

II. GLUINOS

The gluino is a color octet fermion, so it cannot mix with any other particle in the SSM, even if $R$-parity is violated. In this regard, it is unique among all the SSM sparticles. In models with gaugino mass unification, or in models of gauge mediation, the gluino mass parameter $M_3$ is related to the bino and wino mass parameters $M_1$ and $M_2$ by Eq. (6), so

$$M_3 = \frac{\alpha_s}{\alpha} \sin^2 \theta_W M_2 = \frac{3 \alpha_s}{5 \alpha} \cos^2 \theta_W M_1$$

at any RG scale, up to small two-loop corrections. This implies a rough prediction,

$$M_3 : M_2 : M_1 \approx 6 : 2 : 1$$

near the TeV scale. It is therefore reasonable to suspect that the gluino is considerably heavier than the lighter neutralinos and charginos (even in many models where the gaugino mass unification condition is not imposed).

For more precise estimates, one must take into account the fact that $M_3$ is really a running mass parameter with an implicit dependence on the RG scale $Q$. Because the gluino is a strongly interacting particle, $M_3$ runs rather quickly with $Q$,

$$\beta_{M_3} = \frac{d}{dt} M_3 = \frac{1}{8 \pi^2} b_3 g_3^2 M_3, \quad b_3 = 33/5.$$  \hspace{1cm} (22)

A more useful quantity physically is the RG scale independent mass $m_\tilde{g}$ at which the renormalized gluino propagator has a pole. Including one-loop corrections to the gluino propagator due to gluon exchange and quark-squark loops, one finds that the pole mass is given in terms of the running mass in the $\overline{\text{DR}}$ scheme by

$$m_\tilde{g} = M_3(Q) \left(1 + \frac{\alpha_s}{4 \pi} [15 + 6 \ln(Q/M_3) + \sum A_{\tilde{q}}]\right),$$  \hspace{1cm} (23)
where
\[ A_q = \int_0^1 dx \, x \ln[x m_{\tilde{q}}^2/M_3^2 + (1 - x)m_{\tilde{q}}^2/M_3^2 - x(1 - x) - i\epsilon]. \] (24)

The sum in Eq. (23) is over all 12 quark-squark multiplets, and we neglect small effects due to squark mixing. [As a check, requiring \( m_{\tilde{q}} \) to be independent of \( Q \) in Eq. (23) reproduces the one-loop RG equation in Eq. (22).] The correction terms proportional to \( \alpha_s \) in Eq. (23) can be quite significant, because the gluino is strongly interacting, with a large group theory factor [the 15 in Eq. (23)] due to its color octet nature, and because it couples to all the quark-squark pairs. The leading two-loop corrections to the gluino pole mass typically increase the prediction by another 1–2%.

### III. SQUARKS AND SLEPTONS

In principle, any scalars with the same electric charge, \( R \)-parity, and color quantum numbers can mix with each other. This means that with completely arbitrary soft terms, the mass eigenstates of the squarks and sleptons of the MSSM should be obtained by diagonalizing the three \( 6 \times 6 \) mass-squared matrices for up-type squarks (\( \tilde{u}_L, \tilde{c}_L, \tilde{t}_L, \tilde{u}_R, \tilde{c}_R, \tilde{t}_R \)), down-type squarks (\( \tilde{d}_L, \tilde{s}_L, \tilde{b}_L, \tilde{d}_R, \tilde{s}_R, \tilde{b}_R \)), and charged sleptons (\( \tilde{e}_L, \tilde{\mu}_L, \tilde{\tau}_L, \tilde{e}_R, \tilde{\mu}_R, \tilde{\tau}_R \)), and one \( 3 \times 3 \) mass-squared matrix for sneutrinos (\( \tilde{\nu}_e, \tilde{\nu}_\mu, \tilde{\nu}_\tau \)). The general form of the \( 6 \times 6 \) matrix is
\[ M_{\tilde{f}}^2 = \left( m_{\tilde{f}}^2 + (T_{3f} - Q_f s_W^2) c_{2\beta} m_Z^2 + M_f M_f^\dagger \right) a_f v_f - \mu^* Y_f v_f \]
\[ a_f v_f - \mu Y_f^c v_f \]
\[ m_{\tilde{f}}^2 - Q_f s_W^2 c_{2\beta} m_Z^2 + M_f M_f^\dagger \].
(25)

The situation is considerably simplified if two special conditions are imposed on the flavor structure of the soft supersymmetry breaking terms:

1. Squark and slepton mass-squared matrices are universal:
\[ M_Q^2 = m_Q^2 \mathbf{1}, \quad M_{\tilde{u}}^2 = m_{\tilde{u}}^2 \mathbf{1}, \quad M_{\tilde{c}}^2 = m_{\tilde{c}}^2 \mathbf{1}, \quad M_{\tilde{d}}^2 = m_{\tilde{d}}^2 \mathbf{1}, \quad M_{\tilde{L}}^2 = m_{\tilde{L}}^2 \mathbf{1}, \quad M_{\tilde{e}}^2 = m_{\tilde{e}}^2 \mathbf{1}. \] (26)

2. The \( A \)-matrices of trilinear scalar couplings are proportional to the corresponding Yukawa matrices:
\[ a_u = A_u Y_u, \quad a_d = A_d Y_d, \quad a_e = A_e Y_e. \] (27)

We will later come back to a detailed discussion whether this structure is theoretically motivated and phenomenologically unavoidable. For now, we analyze the spectrum that
would follow if (26) and (27) hold. Then, most of the mixing angles are very small. The third-generation squarks and sleptons can have masses very different from those first- and second-generation counterparts, because of the effects of large Yukawa \((y_t, y_b, y_\tau)\) and soft \((a_t, a_b, a_\tau)\) couplings in the RG equations:

\[
16\pi^2 \frac{d}{dt} m_{Q_3}^2 = X_t + X_b - \frac{32}{3} g_3^2 |M_3|^2 - 6g_2^2 |M_2|^2 - \frac{2}{15} g_1^2 |M_1|^2 + \frac{1}{5} g_1^2 S , \tag{28}
\]

\[
16\pi^2 \frac{d}{dt} m_{u_3}^2 = 2X_t - \frac{32}{3} g_3^2 |M_3|^2 - \frac{32}{15} g_1^2 |M_1|^2 - \frac{4}{5} g_1^2 S , \tag{29}
\]

\[
16\pi^2 \frac{d}{dt} m_{d_3}^2 = 2X_b - \frac{32}{3} g_3^2 |M_3|^2 - \frac{8}{15} g_1^2 |M_1|^2 + \frac{2}{5} g_1^2 S , \tag{30}
\]

\[
16\pi^2 \frac{d}{dt} m_{L_3}^2 = X_\tau + X_b - 6g_2^2 |M_2|^2 - \frac{6}{5} g_1^2 |M_1|^2 - \frac{3}{5} g_1^2 S , \tag{31}
\]

\[
16\pi^2 \frac{d}{dt} m_{e_3}^2 = 2X_\tau - \frac{24}{5} g_1^2 |M_1|^2 + \frac{6}{5} g_1^2 S , \tag{32}
\]

where

\[
S \equiv \text{Tr}[Y f m_{\phi^2}] = m_{H_u}^2 - m_{H_d}^2 + \text{Tr}[M_Q^2 - M_L^2 - 2 M_{\nu_{\tau}}^2 + M_{\nu_{e}}^2 + m_{\nu_{e}}^2],
\]

\[
X_t = 2|y_t|^2 (m_{H_u}^2 + m_{Q_3}^2 + m_{u_3}^2) + 2|a_t|^2 ,
\]

\[
X_b = 2|y_b|^2 (m_{H_d}^2 + m_{Q_3}^2 + m_{d_3}^2) + 2|a_b|^2 ,
\]

\[
X_\tau = 2|y_{\tau}|^2 (m_{H_d}^2 + m_{L_3}^2 + m_{e_3}^2) + 2|a_\tau|^2 . \tag{33}
\]

Furthermore, the third generation squarks and charged sleptons can have substantial mixing in pairs \((\tilde{t}_L, \tilde{t}_R), (\tilde{b}_L, \tilde{b}_R), (\tilde{\tau}_L, \tilde{\tau}_R)\). In contrast, the first and second generation squarks and sleptons have negligible Yukawa couplings (and, consequently, under our assumption (27), also negligible \(a\)-couplings), so they end up in seven very nearly degenerate, unmixed pairs \((\tilde{e}_R, \tilde{\mu}_R), (\tilde{\nu}_e, \tilde{\nu}_\mu), (\tilde{\tau}_L, \tilde{\tau}_R), (\tilde{d}_R, \tilde{s}_R), (\tilde{u}_L, \tilde{c}_L), (\tilde{d}_L, \tilde{s}_L)\).

Let us first consider the spectrum of the first- and second-generation squarks and sleptons. In models with universal boundary conditions, the running masses-squared can be conveniently summarized as follows:

\[
m_{Q_1}^2 = m_{Q_2}^2 = m_{Q_3}^2 + K_3 + K_2 + \frac{1}{36} K_1 ,
\]

\[
m_{u_1}^2 = m_{u_2}^2 = m_{Q_3}^2 + \frac{4}{9} K_1 ,
\]

\[
m_{d_1}^2 = m_{d_2}^2 = m_{Q_3}^2 + \frac{1}{9} K_1 ,
\]

\[
m_{L_1}^2 = m_{L_2}^2 = m_{Q_3}^2 + K_2 + \frac{1}{4} K_1 ,
\]

\[
m_{e_1}^2 = m_{e_2}^2 = m_{Q_3}^2 + K_1 . \tag{34}
\]
A key point is that the same $K_3$, $K_2$ and $K_1$ appear in these equations, since all the chiral supermultiplets couple to the same gauginos with the same gauge couplings. The different couplings in front of $K_1$ just correspond to the various values of the weak hypercharge. The $K_i$’s arise from the RG running proportional to the gaugino masses. Explicitly,

$$K_a(Q) = \begin{cases} 
3/5 \
3/4 \
4/3 
\end{cases} \times \frac{1}{2\pi^2} \int_{m}^{Q_0} dt g_a^2(t) |M_a(t)|^2 \quad (a = 1, 2, 3). \quad (35)$$

Here $Q_0$ is the input RG scale at which the universal boundary conditions apply, and $Q$ should be taken to be evaluated near the squark and slepton mass under consideration, presumably less than 1 TeV. The running parameters $g_a(Q)$ and $M_a(Q)$ obey

$$\beta_{g_a} \equiv \frac{d}{dt} g_a = \frac{1}{16\pi^2} b_ag_a^3, \quad (b_1, b_2, b_3) = (33/5, 1, -3),$$

$$M_1/g_1^2 = M_2/g_2^2 = M_3/g_3^2 = m_{1/2}/g_U^2. \quad (36)$$

If the input scale is taken to be the scale of coupling unification, $M_U \approx 2 \times 10^{16}$ GeV, one finds that numerically

$$K_1 \approx 0.15 m_{1/2}^2, \quad K_2 \approx 0.5 m_{1/2}^2, \quad K_3 \approx (4.5 - 6.5) m_{1/2}^2 \quad (37)$$

for $Q \sim m_Z$. Note that $K_3 \gg K_2 \gg K_1$; this is a consequence of the relative sizes of the gauge couplings. The large uncertainty in $K_3$ is due, in part, to the experimental uncertainty in the QCD coupling constant, and in part due to the uncertainty in where to choose $Q$, since $K_3$ runs rather quickly below TeV.

Quite generally (for example, in gauge mediation, where there is a different $m_0^2$ for each type of gauge representation), one expects that the squarks are heavier than the sleptons.

Regardless of the type of model, there is also a “hyperfine” splitting in the squark and slepton mass spectrum produced by electroweak symmetry breaking. Each squark and slepton $\phi$ gets a contribution $\Delta_\phi$ to its mass-squared, coming from the $SU(2)_L$ and $U(1)_Y$ $D$-term quartic interactions of the form (squark)$^2$(Higgs)$^2$ and (slepton)$^2$(Higgs)$^2$, where the neutral Higgs scalars $H_u^0$ and $H_d^0$ get VEVs. They are model independent for a given value of $\tan \beta$:

$$\Delta_\phi = (T_{3\phi} g^2 - Y_{\phi} g^2) (v_d^2 - v_u^2) = (T_{3\phi} - Q_{\phi} \sin^2 \theta_W) \cos 2\beta \ m_Z^2, \quad (38)$$

where $T_{3\phi}$, $Y_{\phi}$ and $Q_{\phi}$ are the third component of weak isospin, the weak hypercharge, and the electric charge of the left-handed chiral supermultiplet to which $\phi$ belongs. For
example, $\Delta_{u_L} = \left( \frac{1}{2} - \frac{2}{3}\sin^2 \theta_W \right) \cos 2\beta \ m_2^2$, $\Delta_{d_L} = \left( -\frac{1}{2} + \frac{1}{3}\sin^2 \theta_W \right) \cos 2\beta \ m_2^2$, and $\Delta_{\tilde u_R} = \frac{2}{3}\sin^2 \theta_W \cos 2\beta \ m_2^2$. These $D$-term contributions are typically smaller than the $m_0^2$ and $K_{1,2,3}$ contributions, but should not be neglected. They split apart the components of $SU(2)_L$ doublets. The mass splittings between the members of the doublet within each of the first two generations is governed by the model independent sum rules

$$m_{\tilde e_L}^2 - m_{\tilde\nu_e}^2 = m_{\tilde d_L}^2 - m_{\tilde u_L}^2 = g^2(v_u^2 - v_d^2)/2 = -\cos 2\beta m_W^2. \quad (39)$$

In the allowed range $\tan \beta > 1$. It follows that $m_{\tilde e_L} > m_{\tilde\nu_e}$ and $m_{\tilde u_L} > m_{\tilde d_L}$, with the magnitude of the splitting constrained by electroweak symmetry breaking.

Next we consider the masses of the top squarks, for which there are several non-negligible contributions. First, there are the mass-squared terms for $\tilde t_L^1\tilde t_L$ and $\tilde t_R^1\tilde t_R$, that are just equal to $m_t^2$, similar as for the first two generations. Second, there are contributions equal to $m_t^2$ for each of $\tilde t_L^1\tilde t_L$ and $\tilde t_R^1\tilde t_R$. These come from the $F$-terms in the scalar potential of the form $y_t^2 H_u^0 H_d^0 (\tilde t_L^1\tilde t_L + \tilde t_R^1\tilde t_R)$, with the Higgs fields replaced by their VEVs. (Of course, there are similar contributions to all other squarks and sleptons, but, for most purposes, they are negligibly small.) Third, there are contributions to the scalar potential from $F$-terms of the form $-\mu^* y_t^* \tilde t_L^1\tilde t_R^1 H_d^* + h.c.$ These contribute to the stop mass terms as $-\mu^* v y_t \cos \beta \tilde t_L^1\tilde t_R^1 + h.c.$ Finally, there are contributions to the scalar potential from the soft (scalar)$^3$ couplings $a_t \tilde t_R^1 \tilde Q_3 H_u + h.c.$ Putting all of these contributions together, we have a mass-squared $2 \times 2$ matrix for the top squarks, which in the interaction basis ($\tilde t_L$, $\tilde t_R$) is given by

$$\mathcal{L}_{\tilde t} = -(\tilde t_L^* \tilde t_R) M^2_{\tilde t} \left( \begin{array}{c} \tilde t_L \\ \tilde t_R \end{array} \right), \quad (40)$$

where

$$M^2_{\tilde t} = \begin{pmatrix} m_{Q_3}^2 + m_t^2 + \Delta_{\tilde u_L} & v(a_t^* \sin \beta - \mu y_t \cos \beta) \\ v(a_t \sin \beta - \mu^* y_t \cos \beta) & m_{\tilde u_3}^2 + m_t^2 + \Delta_{\tilde u_R} \end{pmatrix}. \quad (41)$$

This hermitian mass matrix can be diagonalized by a unitary matrix to give the mass eigenstates:

$$\begin{pmatrix} \tilde t_1 \\ \tilde t_2 \end{pmatrix} = \begin{pmatrix} c_\theta & -s_\theta^* \\ s_\theta & c_\theta^* \end{pmatrix} \begin{pmatrix} \tilde t_L \\ \tilde t_R \end{pmatrix}, \quad (42)$$

of masses $m_{\tilde t_1}^2 < m_{\tilde t_2}^2$. Here $|c_\theta|^2 + |s_\theta|^2 = 1$. If the off-diagonal terms of Eq. (41) are real, then $c_\theta$ and $s_\theta$ are the cosine and sine of a stop mixing angle $\theta_\tilde t$, which can be chosen in the range $0 \leq \theta_\tilde t < \pi$. Because of the large RG effects proportional to $X_t$ in Eqs. (28) and (29), at the electroweak scale one often finds that $m_{\tilde u_3}^2 < m_{Q_3}^2$, and both of these quantities are
usually significantly smaller than the squark masses-squared for the first two generations. The diagonal terms in (41) tend to somewhat mitigate this effect. The off-diagonal terms typically induce, however, significant mixing, which always reduces the lighter top-squark mass-squared eigenvalue. Therefore, models often predict that $\tilde{t}_1$ is the lightest squark of all, and that it is predominantly $\tilde{t}_R$.

A very similar analysis can be performed for the bottom squarks and for the charged tau sleptons, which in the interaction bases $(\tilde{b}_L, \tilde{b}_R)$ and $(\tilde{\tau}_L, \tilde{\tau}_R)$ have the following mass-squared matrices:

$$M_{\tilde{b}}^2 = \begin{pmatrix} m_{Q_3}^2 + m_t^2 + \Delta_{d_L} & v(a_t^* \cos \beta - \mu y_b \sin \beta) \\ v(a_b \cos \beta - \mu^* y_b \sin \beta) & m_{d_3}^2 + m_t^2 + \Delta_{d_R} \end{pmatrix},$$

(43)

$$M_{\tilde{\tau}}^2 = \begin{pmatrix} m_{\tilde{e}_3}^2 + m_t^2 + \Delta_{\tilde{e}_L} & v(a_\tau^* \cos \beta - \mu y_\tau \sin \beta) \\ v(a_\tau \cos \beta - \mu^* y_\tau \sin \beta) & m_{\tilde{e}_3}^2 + m_t^2 + \Delta_{\tilde{e}_R} \end{pmatrix}.$$  

(44)

These can be diagonalized to give the mass eigenstates $\tilde{b}_1, \tilde{b}_2$ and $\tilde{\tau}_1, \tilde{\tau}_2$, in exact analogy to Eq. (42).

The magnitude and importance of mixing in the sbottom and stau sectors depend on how big $\tan \beta$ is. If $\tan \beta$ is not very large (usually this means $\tan \beta \lesssim 10$), the sbottoms and staus are not significantly affected from the mixing terms and the RG effects from $X_b$ and $X_\tau$, because $y_b, y_\tau \ll y_t$. In that case, the mass eigenstates are very nearly the interaction eigenstates, $\tilde{b}_L, \tilde{b}_R, \tilde{\tau}_L$ and $\tilde{\tau}_R$. The latter three, and $\tilde{\nu}_\tau$, are then nearly degenerate with their first- and second-generation counterparts. However, even in the case of small $\tan \beta$, $\tilde{b}_L$ feels the effects of the large top Yukawa coupling. In particular, from Eq. (28) we see that $X_t$ acts to decrease $m_{Q_3}^2$ as it is RG-evolved from the input scale to the electroweak scale. Consequently, the mass of $\tilde{b}_L$ can be significantly lighter than those of $\tilde{d}_L$ and $\tilde{s}_L$, even if the boundary conditions are universal.

For larger values of $\tan \beta$, the mixings in Eqs. (43) and (44) can be quite significant, because $y_b, y_\tau$ and $a_b, a_\tau$ are non-negligible. Just as in the case of top squarks, the lighter sbottom ($\tilde{b}_1$) and stau ($\tilde{\tau}_1$) can be significantly lighter than the first and second generation counterparts, even for universal boundary conditions. Similarly, $\tilde{\nu}_\tau$ can be significantly lighter than $\tilde{\nu}_\mu, \tilde{\nu}_e$.

The requirement that the third generation squarks and sleptons have positive mass-squared implies limits on the magnitudes of $a_t^* \sin \beta - \mu y_b \cos \beta$, $a_b^* \cos \beta - \mu y_b \sin \beta$, and $a_\tau^* \cos \beta - \mu y_\tau \sin \beta$. If they are too large, then the smaller eigenvalue of Eqs. (41), (43) or
TABLE I: Interaction and mass eigenstates of the MSSM.

<table>
<thead>
<tr>
<th>Name</th>
<th>Spin</th>
<th>R-parity</th>
<th>Interaction e.s.</th>
<th>Mass e.s.</th>
</tr>
</thead>
<tbody>
<tr>
<td>Higgs bosons</td>
<td>0</td>
<td>+1</td>
<td>$H_u, H_d$</td>
<td>$h^0, H^0, A^0, H^\pm$</td>
</tr>
<tr>
<td>Squarks</td>
<td>0</td>
<td>−1</td>
<td>$\tilde{Q}_{1,2,3}, \tilde{c}, \tilde{t}, \tilde{d}$</td>
<td>$\tilde{u}<em>{1,2,3,4,5,6}$, $\tilde{d}</em>{1,2,3,4,5,6}$</td>
</tr>
<tr>
<td>Sleptons</td>
<td>0</td>
<td>−1</td>
<td>$\tilde{L}_{1,2,3}, \tilde{\ell}$</td>
<td>$\tilde{\nu}_{1,2,3}$</td>
</tr>
<tr>
<td>Neutralinos</td>
<td>$\frac{1}{2}$</td>
<td>−1</td>
<td>$\tilde{B}^0, \tilde{W}^0, \tilde{H}^0_u, \tilde{H}^0_d$</td>
<td>$\tilde{N}_{1,2,3,4}$</td>
</tr>
<tr>
<td>Charginos</td>
<td>$\frac{1}{2}$</td>
<td>−1</td>
<td>$\tilde{W}^\pm, \tilde{H}^+_u, \tilde{H}^-_d$</td>
<td>$\tilde{C}^\pm_{1,2}$</td>
</tr>
<tr>
<td>Gluinos</td>
<td>$\frac{1}{2}$</td>
<td>−1</td>
<td>$\tilde{g}$</td>
<td>$\tilde{g}$</td>
</tr>
<tr>
<td>Gravitino</td>
<td>$\frac{1}{2}, \frac{3}{2}$</td>
<td>−1</td>
<td>$\tilde{G}$</td>
<td>$\tilde{G}$</td>
</tr>
</tbody>
</table>

(44) is driven negative, implying that a squark or a charged slepton gets a VEV, breaking $SU(3)_C$ and/or $U(1)_{EM}$.

IV. SUMMARY

The MSSM interaction and mass states are presented in Table I.

V. THE SUPERSYMMETRIC FLAVOR PROBLEM

A. The new physics flavor puzzle

We now derive bounds on the scale of new physics and/or its flavor structure. We do so in a model-independent language, and will later apply it to supersymmetry. Consider, for example, the following dimension-six, four-fermion, flavor changing operators:

\[
\mathcal{L}_{\Delta F=2} = \frac{z_{sd}}{\Lambda_{NP}^2} \langle \overline{d}_L \gamma_\mu s_L \rangle^2 + \frac{z_{cu}}{\Lambda_{NP}^2} \langle \overline{c}_L \gamma_\mu u_L \rangle^2 + \frac{z_{bd}}{\Lambda_{NP}^2} \langle \overline{d}_L \gamma_\mu b_L \rangle^2 + \frac{z_{bs}}{\Lambda_{NP}^2} \langle \overline{s}_L \gamma_\mu b_L \rangle^2.
\]  

Each of these terms contributes to the mass splitting between the corresponding two neutral mesons. For example, the term $\mathcal{L}_{\Delta B=2} \propto \langle \overline{d}_L \gamma_\mu b_L \rangle^2$ contributes to $\Delta m_B$, the mass difference between the two neutral $B$-mesons. We use $M^B_{12} = \frac{1}{2m_B} \langle B^0 | \mathcal{L}_{\Delta F=2} | B^0 \rangle$ and

\[
\langle B^0 | (\overline{d}_L \gamma_\mu b_L) (\overline{d}_L \gamma_\mu b_L) | B^0 \rangle = -\frac{1}{3} m_B^2 f_B^2 B_B. \tag{46}
\]
(Analogous expressions hold for the other neutral mesons.) This leads to \( \Delta m_B/m_B = 2|M_{121}^B|/m_B \sim (z_{bd}/3)(f_B/\Lambda_{NP})^2 \). Experiments give:

\[
\begin{align*}
\epsilon_K &\sim 2.3 \times 10^{-3}, \\
\Delta m_K/m_K &\sim 7.0 \times 10^{-15}, \\
\Delta m_D/m_D &\lesssim 2 \times 10^{-14}, \\
\Delta m_B/m_B &\sim 6.3 \times 10^{-14}, \\
\Delta m_{B_s}/m_{B_s} &\sim 2.1 \times 10^{-12}.
\end{align*}
\]

These measurements give then the following constraints (the bound on \( \Im(z_{sd}) \) is stronger by a factor of \((2\sqrt{\epsilon_K})^{-1}\) than the bound on \(|z_{sd}|\):

\[
\Lambda_{NP} \gtrsim \begin{cases} \\
\sqrt{\Im(z_{sd})} \times 10^4 \text{ TeV} & \epsilon_K \\
\sqrt{z_{sd}} \times 10^3 \text{ TeV} & \Delta m_K \\
\sqrt{z_{cu}} \times 10^2 \text{ TeV} & \Delta m_D \\
\sqrt{z_{bd}} \times 10^1 \text{ TeV} & \Delta m_B \\
\sqrt{z_{bs}} \times 10^1 \text{ TeV} & \Delta m_{B_s}
\end{cases}
\]

If the new physics has a generic flavor structure, that is \( z_{ij} = \mathcal{O}(1) \), then its scale must be above \( 10^3 - 10^4 \text{ TeV} \) (or, if the leading contributions involve electroweak loops, above \( 10^2 - 10^3 \text{ TeV} \)). If indeed \( \Lambda_{NP} \gg \text{TeV} \), it means that we have misinterpreted the hints from the fine-tuning problem and the dark matter puzzle. There is, however, another way to look at these constraints:

\[
\begin{align*}
\Im(z_{sd}) &\lesssim 6 \times 10^{-9} (\Lambda_{NP}/\text{TeV})^2, \\
z_{sd} &\lesssim 8 \times 10^{-7} (\Lambda_{NP}/\text{TeV})^2, \\
z_{cu} &\lesssim 1 \times 10^{-6} (\Lambda_{NP}/\text{TeV})^2, \\
z_{bd} &\lesssim 6 \times 10^{-6} (\Lambda_{NP}/\text{TeV})^2, \\
z_{bs} &\lesssim 2 \times 10^{-4} (\Lambda_{NP}/\text{TeV})^2.
\end{align*}
\]

It could be that the scale of new physics is of order TeV, but its flavor structure is far from generic.

One can use that language of effective operators also for the SM, integrating out all particles significantly heavier than the neutral mesons (that is, the top, the Higgs and the
weak gauge bosons). Thus, the scale is $\Lambda_{SM} \sim m_W$. Since the leading contributions to neutral meson mixings come from box diagrams, the $z_{ij}$ coefficients are suppressed by $\alpha_\text{q}^2$. To identify the relevant flavor suppression factor, one can employ the spurion formalism. For example, the flavor transition that is relevant to $B^0 - \bar{B}^0$ mixing involves $d_L b_L$ which transforms as $(8, 1, 1)_{SU(3)^c}$. The leading contribution must then be proportional to $(Y_u^T Y_u)^{13} \propto y_t^2 V_{tb} V_{td}^\ast$. Indeed, an explicit calculation (using VIA for the matrix element and neglecting QCD corrections) gives

$$2M_{12}^B/m_B \approx -\frac{\alpha_\text{q}^2 f_B^2}{12 m_W^2} S_0(x_i)(V_{ib} V_{td}^\ast)^2,$$

where $x_i = m_i^2/m_W^2$ and

$$S_0(x) = \frac{x}{(1-x)^2} \left[ 1 - \frac{11x}{4} + \frac{x^2}{4} - \frac{3x^2 \ln x}{2(1-x)} \right].$$

(50)

Similar spurion analyses, or explicit calculations, allow us to extract the weak and flavor suppression factors that apply in the SM:

$$\Im(z_{SM}^{sd}) \sim \alpha_\text{q}^2 y_t^2 |V_{td} V_{ts}|^2 \sim 1 \times 10^{-10},$$
$$z_{SM}^{sd} \sim \alpha_\text{q}^2 y_t^2 |V_{td} V_{ts}|^2 \sim 5 \times 10^{-9},$$
$$z_{SM}^{bd} \sim \alpha_\text{q}^2 y_t^2 |V_{td} V_{tb}|^2 \sim 7 \times 10^{-8},$$
$$z_{SM}^{bs} \sim \alpha_\text{q}^2 y_t^2 |V_{ts} V_{tb}|^2 \sim 2 \times 10^{-6}.$$  

(52)

(We did not include $z_{SM}^{cu}$ in the list because it requires a more detailed consideration. The naively leading short distance contribution is $\propto \alpha_\text{q}^2 (y_s^4/\alpha_\text{q}^2) |V_{cs} V_{us}|^2 \sim 5 \times 10^{-13}$. However, higher dimension terms can replace a $y_s^2$ factor with $(\Lambda/m_D)^2$. Moreover, long distance contributions are expected to dominate. In particular, peculiar phase space effects have been identified which are expected to enhance $\Delta m_D$ to within an order of magnitude of the present upper bound.)

It is clear than contributions from new physics at $\Lambda_{NP} \sim 1 \text{ TeV}$ should be suppressed by factors that are comparable or smaller than the SM ones. Why does that happen? This is the new physics flavor puzzle.

The fact that the flavor structure of new physics at the TeV scale must be non-generic means that flavor measurements are a good probe of the new physics. Perhaps the best-studied example is that of supersymmetry. Here, the spectrum of the superpartners and the structure of their couplings to the SM fermions will allow us to probe the mechanism of dynamical supersymmetry breaking.
B. The supersymmetric flavor puzzle

We consider the contributions from the box diagrams involving the squark doublets of the first two generations, $\tilde{Q}_{L1,2}$, to the $D^0 - \bar{D}^0$ and $K^0 - \bar{K}^0$ mixing amplitudes. The contributions that are relevant to the neutral $D$ system are proportional to $K_u^u K_1^u K_{2j}^u K_{1j}^u$, where $K_u$ is the mixing matrix of the gluino couplings to a left-handed up quark and their supersymmetric squark partners. (In the language of the mass insertion approximation, we calculate here the contribution that is $\propto (\delta_{LL})_{12}^2$.) The contributions that are relevant to the neutral $K$ system are proportional to $K_d^d K_1^d K_{2j}^d K_{1j}^d$, where $K_d$ is the mixing matrix of the gluino couplings to a left-handed down quark and their supersymmetric squark partners ($\propto (\delta_{LL})_{12}^2$ in the mass insertion approximation). We work in the mass basis for both quarks and squarks. A detailed derivation can be found in [2]. It gives:

$$M^{D}_{12} = \frac{\alpha_s^2 m_D f_D^2 B_{D\eta_{\text{QCD}}}}{160 m_a^2} [11 \tilde{f}_6(x_u) + 4 x_u f_6(x_u)] \left[ \frac{(\Delta m_q^2)}{m_a^4} \right] (K_u^u K_{1u}^u)^2,$$

$$M^{K}_{12} = \frac{\alpha_s^2 m_K f_K^2 B_{K\eta_{\text{QCD}}}}{160 m_d^2} [11 \tilde{f}_6(x_d) + 4 x_d f_6(x_d)] \left[ \frac{(\Delta m_q^2)}{m_d^4} \right] (K_d^d K_{1d}^d)^2.$$

(53)

(54)

Here $m_{\tilde{u},\tilde{d}}$ is the average mass of the corresponding two squark generations, $\Delta m_{\tilde{u},\tilde{d}}^2$ is the mass-squared difference, and $x_{u,d} = m_{\tilde{q}}^2/m_{\tilde{u},\tilde{d}}^2$.

One can immediately identify three generic ways in which supersymmetric contributions to neutral meson mixing can be suppressed:

1. Heaviness: $m_{\tilde{q}} \gg 1 \text{ TeV}$;

2. Degeneracy: $\Delta m_{\tilde{q}}^2 \ll m_{\tilde{q}}^2$;

3. Alignment: $K_{21}^{d,u} \ll 1$.

When heaviness is the only suppression mechanism, as in split supersymmetry [3], the squarks are very heavy and supersymmetry no longer solves the fine tuning problem.\(^1\) If we want to maintain supersymmetry as a solution to the fine tuning problem, either degeneracy or alignment or a combination of both is needed. This means that the flavor structure of supersymmetry is not generic, as argued in the previous section.

\(^1\) When the first two squark generations are mildly heavy and the third generation is light, as in effective supersymmetry [4], the fine tuning problem is still solved, but additional suppression mechanisms are needed.
C. Non-degenerate squarks at the LHC?

The $2 \times 2$ mass-squared matrices for the relevant squarks have the following form:

$$
\tilde{M}_{U_L}^2 = \tilde{m}_{Q_L}^2 + \left(\frac{1}{2} - \frac{2}{3} s_w^2\right) m_Z^2 \cos 2\beta + M_u M_u^\dagger, \\
\tilde{M}_{D_L}^2 = \tilde{m}_{Q_L}^2 - \left(\frac{1}{2} - \frac{1}{3} s_w^2\right) m_Z^2 \cos 2\beta + M_d M_d^\dagger.
$$

(55)

We note the following features of the various terms:

- $\tilde{m}_{Q_L}^2$ is a $2 \times 2$ hermitian matrix of soft supersymmetry breaking terms. It does not break $SU(2)_L$ and consequently it is common to $\tilde{M}_{U_L}^2$ and $\tilde{M}_{D_L}^2$. On the other hand, it breaks in general the $SU(2)_Q$ flavor symmetry.

- The terms proportional to $m_Z^2$ are the D-terms. They break supersymmetry (since they involve $D_T \neq 0$ and for $D_Y \neq 0$) and $SU(2)_L$ but conserve $SU(2)_Q$.

- The terms proportional to $M_q^2$ come from the $F_{U_R}$- and $F_{D_R}$-terms. They break the gauge $SU(2)_L$ and the global $SU(2)_Q$ but, since $F_{U_R} = F_{D_R} = 0$, conserve supersymmetry.

Given that we are interested in squark masses close to the TeV scale (and the experimental lower bounds are of order 300 GeV), the scale of the eigenvalues of $\tilde{m}_{Q_L}^2$ is much higher than $m_Z^2$ which, in turn, is much higher than $m_c^2$, the largest eigenvalue in $M_q M_q^\dagger$. We can draw the following conclusions:

1. $m_u^2 = m_d^2 = m_{\tilde{q}}^2$ up to effects of order $m_Z^2$, namely to an accuracy of $\mathcal{O}(10^{-2})$.

2. $\Delta m_u^2 = \Delta m_d^2 = \Delta m_{\tilde{q}}^2$ up to effects of order $m_c^2$, namely to an accuracy of $\mathcal{O}(10^{-5})$.

3. Since $K_u \simeq V_{uL} \tilde{V}_L^\dagger$ and $K_d \simeq V_{dL} \tilde{V}_L^\dagger$ (the matrices $V_{qL}$ diagonalize the quark mass-squared matrices $M_q M_q^\dagger$ while $\tilde{V}_L$ diagonalizes $\tilde{m}_{Q_L}^2$), the mixing matrices $K_u$ and $K_d$ are different from each other, but the following relation to the CKM matrix holds to an accuracy of $\mathcal{O}(10^{-5})$:

$$
K_u K_d^\dagger = V.
$$

(56)

Eqs. (53) and (54) can be translated into our generic language:

$$
\Lambda_{NP} = m_{\tilde{q}}.
$$

(57)
\[ z_{cu} = z_{12} \sin^2 \theta_u, \]
\[ z_{sd} = z_{12} \sin^2 \theta_d, \]
\[ z_{12} = \frac{11f_6(x) + 4xf_6(x)}{18} \alpha_s^2 \left( \frac{\Delta m_q^2}{m_q^2} \right)^2, \]  
(58)

with Eq. (56) giving
\[ \sin \theta_u - \sin \theta_d \approx \sin \theta_c = 0.23. \]  
(59)

We now ask the following question: Is it possible that the first two generation squarks, \( \tilde{Q}_{L,1,2} \), are accessible to the LHC \( (m_{\tilde{q}} \lesssim 1 \text{TeV}) \), and are not degenerate \( (\Delta m_q^2/m_q^2 = \mathcal{O}(1)) \)?

To answer this question, we use Eqs. (49). For \( \Lambda_{\text{NP}} \lesssim 1 \text{TeV} \), we have \( z_{cu} \lesssim 1 \times 10^{-6} \) and, for a phase that is \( \ll 0.1 \), \( z_{sd} \lesssim 6 \times 10^{-8} \). On the other hand, for non-degenerate squarks, and, for example, \( 11f_6(1) + 4f_6(1) = 1/6 \), we have \( z_{12} = 8 \times 10^{-5} \). Then we need, simultaneously, \( \sin \theta_u \lesssim 0.11 \) and \( \sin \theta_d \lesssim 0.03 \), but this is inconsistent with Eq. (59).

There are three ways out of this situation:

1. The first two generation squarks are quasi-degenerate. The minimal level of degeneracy is \( (\tilde{m}_2 - \tilde{m}_1)/(\tilde{m}_2 + \tilde{m}_1) \lesssim 0.12 \). It could be the result of RGE [5].

2. The first two generation squarks are heavy. Putting \( \sin \theta_u = 0.23 \) and \( \sin \theta_d \approx 0 \), as in models of alignment [6, 7], Eq. (48) leads to
\[ m_{\tilde{q}} \gtrsim 2 \text{TeV}. \]  
(60)

3. The ratio \( x = \tilde{m}_g^2/\tilde{m}_q^2 \) is in a fine-tuned region of parameter space where there are accidental cancellations in \( 11f_6(x) + 4xf_6(x) \). For example, for \( x = 2.33 \), this combination is \( \sim 0.003 \) and the bound (60) is relaxed by a factor of 7.

Barring such accidental cancellations, the *model independent* conclusion is that, if the first two generations of squark doublets are within the reach of the LHC, they must be quasi-degenerate [8, 9].

**Exercise:** Does \( K_{31}^d \sim |V_{ub}| \) suffice to satisfy the \( \Delta m_B \) constraint with neither degeneracy nor heaviness? (Use the two generation approximation and ignore the second generation.)

Is there a natural way to make the squarks degenerate? Examining Eqs. (55) we learn that degeneracy requires \( \tilde{m}_{Q_L}^2 \simeq \tilde{m}_q^2 1 \). We have mentioned already that flavor universality
is a generic feature of gauge interactions. Thus, the requirement of degeneracy is perhaps a hint that supersymmetry breaking is *gauge mediated* to the MSSM fields.