Radiative corrections to sfermion mass splittings

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Abstract

We study the one-loop radiative corrections to the SU(2) breaking mass splittings between sfermions in SU(2) doublets in the Minimal Supersymmetric Standard Model. At tree-level, the differences of mass squared $m^2_{\tilde{f}_1 L} - m^2_{\tilde{f}_2 L}$ of SU(2) doublet sfermions ($\tilde{f}_1$, $\tilde{f}_2)_L$ in the first two generations are determined by $\tan \beta$, and are universal to sleptons and squarks. The radiative correction, however, breaks this relation. The typical deviation from the universality between sleptons and squarks is within $\pm 0.05$ in terms of the “effective $\cos 2\beta$”. We also study the SUSY parameter dependence of the deviation. For very heavy sfermions, the relative corrections become large.
1 Introduction

In the Minimal Supersymmetric (SUSY) Standard Model (MSSM) \cite{1}, the tree-level masses of left-handed sfermions $\tilde{f}_L$ in SU(2) doublets $\tilde{F} = (\tilde{f}_1, \tilde{f}_2)_L$ in the first two generations are given in terms of SU(2) invariant masses $M_{\tilde{F}}$, the ratio of vacuum expectation values of two Higgs scalars $\tan \beta = v_u/v_d$, isospin $I_{3f_L}$, charge $Q_{f_L} = I_{3f_L} + Y_{f_L}$, $m_Z$ and $s_W \equiv \sin \theta_W$, ignoring masses of their fermionic superpartners $f$. The explicit form of their masses is

$$m^2_{\tilde{f}_L} = M^2_{\tilde{F}} + (I_{3f} - s^2_W Q_f)m^2_Z \cos 2\beta. \quad (1)$$

The difference of mass squared between $\tilde{f}_{1L}$ and $\tilde{f}_{2L}$ is then

$$m^2_{\tilde{f}_{1L}} - m^2_{\tilde{f}_{2L}} = m^2_W \cos 2\beta. \quad (2)$$

Eq. (2) is independent of flavors (mass, color, hypercharge) of sfermions and also of any unification conditions beyond the MSSM such as the unification of sfermion masses in minimal supergravity. As a special case, a mass sum rule between squarks and sleptons

$$m^2_{\tilde{\nu}_L} - m^2_{\tilde{\ell}_L} = m^2_{\tilde{u}_L} - m^2_{\tilde{d}_L} = m^2_W \cos 2\beta, \quad (3)$$

follows from eq. (2).

The reason for the relation (2) is that the mass splitting between $\tilde{f}_{1L}$ and $\tilde{f}_{2L}$ is generated by the D-term four-scalar interaction $\lambda(F^* \tau^a F)(H^* \tau^a H_i)$ ($i = (d, u)$, $\tau^a$: generators of SU(2)) which is related to the SU(2) gauge vector interaction by the SUSY Ward identity. The relations (2,3) then provide a detailed test of the MSSM \cite{2,3} and also a method for fixing $\tan \beta$ \cite{4,5}.

However, due to the SUSY violation, the relation (2) between sfermion masses, $m_W$ and $\tan \beta$ is modified by radiative corrections. In this paper, we present analytic and numerical results of one-loop corrections to mass splittings (2) of SU(2) doublet sfermions in the first two generations (i.e. SUSY partners of massless quarks and leptons). The breaking of the relation (2) was briefly discussed in ref. \cite{4} for $(\tilde{f}_1, \tilde{f}_2)_L = (\tilde{\nu}, \tilde{e}_L)$ case. Here we mainly discuss the breaking of the mass relation (3) between sleptons and squarks. It is shown that the deviations from the relation (3) are typically within ±0.05 in terms of “effective $\cos 2\beta$”, which is defined in section 3. The relative corrections to the tree-level results (3) become large for very heavy sfermions. The measurement of the violation of the relation (3) would therefore be important to know the nature of the MSSM beyond the tree-level.

This paper is organized as follows. In section 2, we present the analytic form of the one-loop corrected mass splitting for sfermions. The renormalization condition is also briefly discussed. In section 3, we show the numerical results of the radiative corrections to eqs. (4, 3) and their dependence on various SUSY parameters. Section 4 is devoted for conclusion.
2 One-loop corrections to sfermion masses

In this paper, we ignore the mixing of sfermions in different generations. Then the one-loop corrected mass of a sfermion \( \tilde{f}_L \) is expressed as

\[
m^2_{\tilde{f}_L} (\text{pole}) = \hat{m}^2_{\tilde{f}_L} - \text{Re} \Pi_{\tilde{f}_L} (q^2 = m^2_{\tilde{f}}) + 2m^2_Z (I_{3f} - s^2_W Q_f) (\cos^2 \beta \Delta v_d \over v_d - \sin^2 \beta \Delta v_u \over v_u),
\]

where \( \hat{m}_{\tilde{f}_L} \) denotes the tree-level mass \([\Pi]\) in terms of tree-level parameters \( (\hat{M}^2_F, \hat{m}_Z, \tilde{s}_W, \tan \beta) \):

\[
\hat{m}^2_{\tilde{f}_L} = \hat{M}^2_F + (I_{3f} - s^2_W Q_f) \hat{m}_Z \cos 2\hat{\beta}.
\]

\( \Pi_{\tilde{f}_L} (q^2) \) is the two-point \( \tilde{f}_L \tilde{f}_L \) function. \( \Delta v_{d,u} \) denotes the shift of vacuum expectation values from the tree-level values, namely sum of 1-loop tadpole contributions \( \Delta T^{(1)}_{d,u} \) and their counterterms \( \Delta T^{(1)CT}_{d,u} \). We comment on the renormalization of \( \Delta v_{d,u} \) later in this section.

We use the dimensional reduction \([3]\) for regularization. The contributions of the \( \epsilon \)-scalar masses \([7]\) are neglected here since they are common to \( m_{\tilde{f}_1} \) and \( m_{\tilde{f}_2} \) and cancel out in the mass difference \([7]\).

We calculate \( \Pi_{\tilde{f}_L} (q^2) \) in terms of the \('t\) Hooft-Veltman functions \([8]\) \( A, B_{0,1} \), using definitions given in Ref.\([9]\). We adopt the \('t\) Hooft-Feynman gauge for convenience. As for couplings of SUSY particles, we basically follow the notations in Ref.\([10]\). The explicit form of \( \Pi_{\tilde{f}_L} (q^2) \) is then (henceforth we omit the subscript \( L \) for \( \tilde{f}_1, \tilde{f}_{1,2} \)):

\[
\Pi_{\tilde{f}} (q^2) = \Pi_{\tilde{f}}^{g,\tilde{g}} (q^2) + \Pi_{\tilde{f}}^{Z,\tilde{X}} (q^2) + \Pi_{\tilde{f}}^{W,\tilde{X}^+} (q^2) + \Pi_{\tilde{f}}^{H} (q^2) + \Pi_{\tilde{f}}^{D},
\]

\[
\Pi_{\tilde{f}}^{g,\tilde{g}} (q^2) = \frac{C_f \alpha_3}{4\pi} \left[ -A(0) - (3q^2 + m^2_f)B_0(q^2, 0, \tilde{f}) - 2q^2B_1(q^2, 0, \tilde{f}) + A(\tilde{f}) 
- 4m^2_B_0(q^2, \tilde{g}, 0) - 4q^2B_1(q^2, \tilde{g}, 0) \right],
\]

\[
\Pi_{\tilde{f}}^{Z,\tilde{X}} (q^2) = \frac{Q_f s^2_W \alpha_2}{4\pi c_W} \left[ 3A(0) - (3q^2 + m^2_f)B_0(q^2, 0, \tilde{f}) - 2q^2B_1(q^2, 0, \tilde{f}) + A(\tilde{f}) 
+ \frac{\alpha_2}{4\pi c_W} (I_{3f} - s^2_W Q_f)^2 \left[ 3A(Z) - (3q^2 + m^2_f)B_0(q^2, Z, \tilde{f}) 
- 2q^2B_1(q^2, Z, \tilde{f}) + A(\tilde{f}) \right] 
- \frac{\alpha_2}{4\pi c_W} \sum_i |I_{3f}N_{i2}c_W + Y_fN_{i1}s_W|^2 \left[ A(0) + m^2_{\tilde{X}_i}B_0(q^2, \tilde{X}_i, 0) 
+ q^2B_1(q^2, \tilde{X}_i, 0) \right],
\]

\[
\Pi_{\tilde{f}}^{W,\tilde{X}^+} (q^2) = \frac{\alpha_2}{8\pi} \left[ 3A(W) - (3q^2 + m^2_f)B_0(q^2, W, \tilde{f}^\prime) - 2q^2B_1(q^2, W, \tilde{f}^\prime) + A(\tilde{f}^\prime) 
- \frac{\alpha_2}{2\pi} \sum_j \left\{ |V_{j1}|^2, I_{3f} = +\frac{1}{2} \right\} \left[ A(0) + m^2_{\tilde{X}_j}B_0(q^2, \tilde{X}_j^+, 0) 
- \frac{\alpha_2}{2\pi} \sum_j \left\{ |U_{j1}|^2, I_{3f} = -\frac{1}{2} \right\} \right].
\]
The cancellation of quadratic divergences. The last term of eq. (6), \( \Pi_{DA} \), retained all graphs corresponding to each term in eq. (6). The graphs are shown in Fig. 1. In eqs. (7–9), we have part of the scalar-loop contribution (Fig. 1c) which is determined by the SU(2) fermion mixing matrices [10]. The term \( \Sigma_W \) is then omitted here.

In Eq. (13), we eliminated the tree-level mass \( \hat{m}_W \) and the pole mass \( m_W \):

\[ m_W^2 = \hat{m}_W^2 - \text{Re}\Pi_{f}^{WW}(m_W^2) + 2m_W^2(\cos^2 \beta \frac{\Delta v_d}{v_d} + \sin^2 \beta \frac{\Delta v_u}{v_u}). \]  

Here \( C_q = 4/3 \), \( C_t = 0 \) and \( c_W \equiv \cos \gamma \). \( \tilde{f}' \) denotes the SU(2) partner of \( \tilde{f} \). The graphs corresponding to each term in eq. (6) are shown in Fig. 1. In eqs. (7–10), we have retained all \( A(0) \) terms, which vanish after the minimal subtraction, to check the cancellation of quadratic divergences. The last term of eq. (3), \( \Pi_f^P \), denotes a special part of the scalar-loop contribution (Fig. 1c) which is determined by the SU(2) \( \times U(1) \) quantum numbers of \( \tilde{f}' \). It takes a form

\[ \Pi_f^P = -\frac{\alpha_2}{4\pi} f_{\tilde{f}f} \Sigma_{W^3} - \frac{s_W^2 \alpha_2}{4\pi c_W^2} Y_f \Sigma_B, \]  

where \( \Sigma_{W^3} \) and \( \Sigma_B \) are contributions of scalar loops to auxiliary fields of \( W^3 \) and \( B \), respectively, and are independent of the flavors of \( \tilde{f} \). The explicit form of \( \Sigma_{W^3} \) is

\[ \Sigma_{W^3} = -\frac{3}{2}(A(\tilde{f}) - A(\tilde{f}^c)) - \frac{9}{2}(A(\tilde{u}_L) - A(\tilde{d}_L)) - \frac{3}{2}(\cos^2 \beta \cos \theta_t A(\tilde{t}_1) + \sin^2 \theta_t A(\tilde{t}_2) - A(\tilde{u}_L)) \]

\[ -\frac{3}{4}(2A(H^+) - 2A(W) - A(P^0) + A(Z)) - \frac{\cos 2\alpha}{4}(A(H^0) - A(H^0)). \]  

In eq. (12), we assumed the generation independence of sfermion masses other than top squarks. In contrast to other terms (7–10), \( \Sigma_{W^3} \) includes contributions of the Higgs pseudoscalar \( P^0 \) and top squarks with mass eigenstates \( \tilde{t}_1, \tilde{t}_2 \) and mixing angle \( \theta_t \) [10]. The term \( \Sigma_B \) does not contribute to the mass splitting [2] and is therefore omitted here.

The one-loop corrected mass squared splitting of the sfermion doublet \( (\tilde{f}_1, \tilde{f}_2) \) is then

\[ (m^2_{\tilde{f}_1} - m^2_{\tilde{f}_2})(\text{pole}) = m_W^2 \cos 2\beta - \text{Re}\Pi_{f}^{f_1}(m^2_{f_1}) + \text{Re}\Pi_{f}^{f_2}(m^2_{f_2}) \]

\[ +\text{Re}\Pi_{f}^{WW}(m_W^2) \cos 2\beta + m_W^2 \sin^2 \beta \frac{\Delta v_d}{v_d} - m^2_W \sin^2 \beta \frac{\Delta v_u}{v_u}. \]  

In Eq. (13), we eliminated the tree-level mass \( \hat{m}_W \) by using the relation between \( \hat{m}_W \) and the pole mass \( m_W \):

\[ m_W^2 = \hat{m}_W^2 - \text{Re}\Pi_{f}^{WW}(m_W^2) + 2m_W^2(\cos^2 \beta \frac{\Delta v_d}{v_d} + \sin^2 \beta \frac{\Delta v_u}{v_u}). \]  

In Eq. (14), we used the relation between the mass splittings and the mixing parameters.
The explicit form of the transverse $WW$ two-point function $\Pi_T^{WW}(q^2)$ in the MSSM is given in Refs. [11, 12, 13].

We have to specify the renormalization condition for $\Delta v_{d,u}$ and $\tan \beta$ in order to study the correction to the relation (2). In this paper, we require that $\Delta T_{d,u}^{(1)CT}$ are generated from the renormalization of the Higgs scalar sector and adjusted so that the condition

$$\Delta v_d/v_d = \Delta v_u/v_u,$$

is satisfied. We then define the renormalized $\tan \beta$ by the modified minimal subtraction. $\tan \beta$ is then a DR running parameter which satisfies the renormalization group equation

$$\frac{d}{d \log Q} \tan \beta = -\frac{3h_t^2}{16\pi^2} \tan \beta, \quad h_t = \frac{gm_t}{\sqrt{2m_W} \sin \beta},$$

for $\tan \beta \ll m_t/m_b$. This definition of $\tan \beta$ corresponds to ones adopted in refs. [4, 5, 14] but differs from ones in refs. [12, 13]. Note, however, that the dependence on the renormalization condition for $\Delta v_{d,u}$ and $\tan \beta$ is common to both squarks and sleptons, as is seen in eq.(13). Therefore, its choice, as well as contributions of $\Sigma_W^3$ and $\Pi_T^{WW}(m_W^2)$, does not affect the violation of the sum rule (3) between squarks and sleptons.

We have checked that the overall dependence of the right hand side of eq.(13) on the renormalization scale $Q$ vanishes. We have also checked that the results (13) is independent of gauge fixing parameters in general $R_\xi$ gauge [15], as long as eq.(15) is satisfied.

3 Numerical Results and Discussion

For numerical presentation of our results, we define an “effective $\cos 2\beta$” for a doublet $\tilde{F} = (\tilde{f}_1, \tilde{f}_2)_L$ as

$$\cos 2\beta|_{eff}^{\tilde{F}} \equiv (m_{\tilde{f}_1}^2 - m_{\tilde{f}_2}^2)/m_W^2,$$

where all masses on the right hand side are pole masses. At tree-level, (16) is independent of the flavor of $\tilde{F}$ and identical to $\cos 2\beta = (v_d^2 - v_u^2)/(v_d^2 + v_u^2)$. The “effective $\tan \beta$” is not useful for large $\tan \beta$ cases since, in such cases, eq.(17) is insensitive to $\tan \beta$ and also $\cos 2\beta|_{eff}^{\tilde{F}} < -1$ is possible, as we will see in this section.

We discuss the difference between $\cos 2\beta|_{eff}^{\tilde{Q}}$ for $\tilde{Q} = (\tilde{u}_L, \tilde{d}_L)$, $\cos 2\beta|_{eff}^{\tilde{L}}$ for $\tilde{L} = (\tilde{\nu}, \tilde{e}_L)$ and the DR running parameter $\cos 2\beta$. The Standard Model parameters are set as $m_W = 80.3$ GeV, $\alpha = 1/129$, $s_W^2 = 0.233$ and $\alpha_3 = 0.11$. The DR renormalization scale is set at $m_W$. As for SUSY parameters, we impose the one-loop grand unification of gaugino masses $m_{\tilde{g}}/\alpha_3 = M_2/\alpha_2 = 3M_1/(5\alpha_2 \tan^2 \theta_W)$ and generation independence of $M_{\tilde{F}}$. The difference between $\cos 2\beta|_{eff}^{\tilde{Q}}$ and $\cos 2\beta|_{eff}^{\tilde{F}}$
is then determined by unknown parameters \((M_L, M_{\tilde{Q}}, \tan \beta, M_2, \mu, m_P)\). The additional parameters \((m_t, m_{t_1}, \theta_t)\), which are necessary to calculate the difference between \(\cos 2\beta |_{\text{eff}}^L\) and \(\cos 2\beta |_{\text{eff}}^Q\) as functions of \(\tan \beta\), for a set of typical values of SUSY parameters. The differences between \(\cos 2\beta |_{\text{eff}}^L\), \(\cos 2\beta |_{\text{eff}}^Q\) and \(\cos 2\beta\) are within \(\pm 0.05\). For the parameter choice in this figure, these differences correspond to about \(\pm 0.5\) GeV deviations of \(m_{\tilde{L}} - m_{\tilde{Q}}\) from the tree-level results. It is also shown that for \(\tan \beta > 9\), \(\cos 2\beta |_{\text{eff}}^L\) is below \(-1\). In such region, obviously, we cannot define “on-shell \(\tan \beta\)” in terms of \(m_f\) and \(m_W\) by using eq.(16) and \(\tan \beta |_{\text{eff}}^F = [(1 - \cos 2\beta |_{\text{eff}}^F)/(1 + \cos 2\beta |_{\text{eff}}^F)]^{1/2}\).

In Fig.2, we show the values of \(\cos 2\beta\), \(\cos 2\beta |_{\text{eff}}^L\) and \(\cos 2\beta |_{\text{eff}}^Q\) as functions of \(\tan \beta\), for several values of \((M_L, M_{\tilde{Q}})\). The difference \(\delta c_{2\beta}\) tends to move to negative direction as \(m_{\tilde{L}, \tilde{Q}}\) increases, and move to positive direction as \(\tan \beta\) increases. We can see that the mass difference between sleptons and squarks is not a main cause for the difference between \(\cos 2\beta |_{\text{eff}}^L\). For example, its absolute value is larger for \((M_L = 300\) GeV, \(M_{\tilde{Q}} = 300\) GeV) than for \((M_L = 100\) GeV, \(M_{\tilde{Q}} = 300\) GeV). We find that the QCD \(\mathcal{O}(\alpha_3)\) contribution (8), which is shown in Fig.3 for \(m_{\tilde{Q}} = 300\) GeV, and the electroweak one are of the same order.

In Fig.4a-4c, we show \(\delta c_{2\beta}\) as a function of \((M_2, \mu)\). The behavior of \(\delta c_{2\beta}\) strongly depends on \(\tan \beta\). For example, in the limit of \((M_2, |\mu|) \to \infty\) with fixed \(\mu/M_2\), we find that \(\delta c_{2\beta}\) logarithmically decreases for \(\tan \beta < 1.2\) but increases for \(\tan \beta > 1.2\). In Fig.4d, it is shown that the main part of the dependence of \(\delta c_{2\beta}\) on \((M_2, \mu)\) comes from \(\cos 2\beta |_{\text{eff}}^Q\). In addition, we checked that the dependence of \(\delta c_{2\beta}\) on \(m_P\), which is not shown here, is smaller than \(0.005\) for \(100\) GeV < \(m_P < 1\) TeV.

Finally, we consider a special case where \(\tilde{L}, \tilde{Q}\) are much heavier than all other particles in the MSSM. This limit is theoretically interesting for the following reason: Suppose a possibility that the mass splitting (2) and its deviation from the tree-level result can be described by effective couplings \(\lambda_{\tilde{f} \cdot H^* \cdot H}(Q = M_{\tilde{f}})\) in the effective theory where SUSY particles heavier than \(M_{\tilde{f}}\) are integrated out. If this is the case, we can expect that the violation of the sum rule (8) is very small in this limit, since, as stated in section 1, the difference between \(\cos 2\beta |_{\text{eff}}^L\) and \(\cos 2\beta |_{\text{eff}}^Q\) is generated by the violation of the SUSY Ward identity. However, this is not the case. Fig.5 shows the dependence on \(M_{\tilde{Q}} = M_{\tilde{L}} = m_{\tilde{f}}\) for heavy sfermions. We can see that in this case the difference \(\delta c_{2\beta}\) rapidly increases as \(m_{\tilde{f}}\) increase, contrary to naive expectation above. In fact, the asymptotical form of the difference \(\delta c_{2\beta}\) in this limit is

\[
\delta c_{2\beta} = -\frac{4\alpha m_{\tilde{f}}}{3c_W m_W} + O(\alpha \log \frac{m_{\tilde{f}}}{m_W}) + \text{(finite term)}.
\]

The first term of eq.(18), proportional to \(m_{\tilde{f}}\), emerges from \(f - Z\) loop contribution in (8) due to the singularity of the two-point functions near the threshold \(\sqrt{q^2} \sim \ldots\)
$m_{\tilde{f}} + m_Z$. It is always larger than the second “leading logarithmic” term and dominant for $m_{\tilde{F}} > m_{\tilde{g}}$. Unfortunately, the large contribution to $m_{\tilde{f}_1}^2 - m_{\tilde{f}_2}^2$ does not necessarily imply large $m_{\tilde{f}_1} - m_{\tilde{f}_2}$ which is experimentally observable, due to the relation

$$m_{\tilde{f}_1} - m_{\tilde{f}_2} = \frac{m_{\tilde{f}_1}^2 - m_{\tilde{f}_2}^2}{m_{\tilde{f}_1} + m_{\tilde{f}_2}} \sim \frac{m_W^2}{2 m_{\tilde{F}}} \cos 2\beta|_{\tilde{F} \text{ eff}}.$$  \hspace{1cm} (19)

Therefore, in the limit $m_{\tilde{F}} \gg m_{\text{others}}$, the mass splitting $m_{\tilde{f}_1} - m_{\tilde{f}_2}$ itself remains very small, about $-0.5$ GeV. Nevertheless, the large deviation from tree-level sum rule is theoretically very interesting.

4 Conclusion

We have studied the one-loop radiative corrections to the SU(2) breaking mass splittings $m_{\tilde{f}_1}^2 - m_{\tilde{f}_2}^2$ between SU(2) doublet sfermions ($\tilde{f}_1$, $\tilde{f}_2$) in the MSSM. The analytic and numerical results of the radiative corrections to the mass splittings of sleptons and squarks in the first two generations have been presented. The corrections to the tree-level universal relation (3) between sleptons and squarks is shown to be within ±0.05 in terms of the effective $\cos 2\beta|_{\tilde{F} \text{ eff}}$ \hspace{0.5cm} (17) for typical values of the SUSY parameters. We have also studied the dependence on SUSY parameters. The difference between $\cos 2\beta|_{\tilde{L} \text{ eff}}$ and $\cos 2\beta|_{\tilde{Q} \text{ eff}}$ becomes large for very heavy sfermions. The measurement of the violation of the relation (3) would therefore help us to understand the nature of the MSSM beyond the tree-level.

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Figure Captions

Fig. 1 Feynman graphs for one-loop two-point $\tilde{f}_L\tilde{f}_L$ functions in eqs. (8, 12): (a) $\Pi_{\tilde{f}}^{g\tilde{g}}, \Pi_{\tilde{f}}^{\gamma Z, \tilde{f}}, \Pi_{\tilde{f}}^{W, \tilde{f}+}$; (b) $\Pi_{\tilde{f}}^{H}$; (c) $\Pi_{\tilde{f}}^D$. Double lines denote auxiliary fields $D$ of corresponding gauge supermultiplets.

Fig. 2 $\cos 2\beta (\text{DR})$ and $\cos 2\beta |_{\text{eff}}^{\tilde{L}, \tilde{Q}}$ as functions of $\tan \beta$ for $M_{\tilde{L}} = M_{\tilde{Q}} = 300$ GeV, $M_2 = 100$ GeV, $\mu = -400$ GeV and $m_p = 300$ GeV.

Fig. 3 $\delta_{c2\beta} \equiv \cos 2\beta |_{\text{eff}}^{\tilde{L}} - \cos 2\beta |_{\text{eff}}^{\tilde{Q}}$ as a function of $\tan \beta$ for $(M_{\tilde{L}}, M_{\tilde{Q}})$ (GeV) = (100, 200), (200, 200), (100, 300) and (300, 300). Other SUSY parameters are set as $M_2 = 100$ GeV, $\mu = -400$ GeV and $m_p = 300$ GeV. The QCD contribution is also shown with a dashed line for $M_{\tilde{Q}} = 300$ GeV.

Fig. 4 $\delta_{c2\beta}$ in the $(M_2, \mu)$ plane for $\tan \beta = 1.1$ (a), 2 (b), 10 (c) and (cos $2\beta |_{\text{eff}}^{\tilde{L}}$ (thin dashed lines), cos $2\beta |_{\text{eff}}^{\tilde{Q}}$ (solid lines)) for $\tan \beta = 2$ (d). Other SUSY parameters are set as $M_{\tilde{L}} = M_{\tilde{Q}} = m_p = 300$ GeV. The regions below the thick dashed lines, where either $2m_{\tilde{x}_1} < m_Z$ or $m_{\tilde{x}_1} + m_{\tilde{x}_2} < m_Z$ holds, are excluded by LEP-I constraints.

Fig. 5 $\cos 2\beta |_{\text{eff}}^{\tilde{L}, \tilde{Q}}$ as functions of $m_F \equiv M_{\tilde{L}} = M_{\tilde{Q}}$. Other SUSY parameters are set as $\tan \beta = 2$ (cos $2\beta = -0.6$), $M_2 = 100$ GeV, $\mu = -400$ GeV and $m_p = 300$ GeV.
Fig. 1

(a) $g, \gamma, Z, W^\pm$

(b) $\tilde{\gamma}, \tilde{\chi}_i^0, \tilde{\chi}_j^\pm$

(c) $h^0, H^0, H^\pm, \phi_W^\pm$

D$_{g,W^\pm,W^3,B}$

(all scalars)
Fig. 2

\[ \cos^2 \beta(DR_{\text{bar}}) \]

\[ \cos^2 \beta(\tilde{L}) \]

\[ \cos^2 \beta(\tilde{Q}) \]
Fig. 3

\begin{align*}
\delta_\c_{2\beta} (M_L, M_Q) & \approx (300, 300) \\
\delta_\c_{2\beta} (100,200) \\
\delta_\c_{2\beta} (200,200) \\
\delta_\c_{2\beta} (100,300) \\
\delta_\c_{2\beta} (M_L, M_Q) & \approx (300, 300)
\end{align*}
Fig. 4a
Fig. 4b

$M_2 \, (\text{GeV})$

$\mu \, (\text{GeV})$

Values shown in the plot include 0.015, 0.00, and various other values indicated by dashed lines.
Fig. 4c
Fig. 5

\[ \cos^2 \beta_{(\text{eff})} \]

\[ m_{\tilde{F}} (\text{GeV}) \]

\[ \tilde{Q} \]
\[ \tilde{L} \]