Supersymmetric electroweak corrections to top quark production at the Fermilab Tevatron

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We calculate the genuine supersymmetric electroweak corrections of order $\alpha m_t^2/m_W^2$, which arise from loops of the chargino, neutralino, and squark, to top quark production at the Fermilab Tevatron in the minimal supersymmetric model. The observable hadronic cross section can be enhanced by 20% for a top quark mass of 170 GeV and squark mass of 100 GeV. When the squark mass gets larger, the corrections decrease rapidly.

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I. INTRODUCTION

The evidence for top quark production, with a mass of $174 \pm 10^{+13}_{-12}$ GeV, has been reported by the Collider Detector at Fermilab (CDF) Collaboration at the Tevatron with an integrated luminosity of 19.3 pb^{-1} [1]. The cross section for top-quark pair production is found to be $13.9^{+6.1}_{-4.8}$ pb for $M_{\mathrm{top}} = 174$ GeV. This value is somewhat higher than the most recent theoretical estimate for the top quark production cross section with next-to-leading order QCD corrections which include gluon resummation [2]. In the standard model (SM), the one-loop electroweak corrections to the cross section were found to be only a few percent [3,4] and, thus, cannot significantly modify the result given by Ref. [2]. This difference between the SM prediction and the Tevatron experimental result for top production cross section indicates the presence of new physics. One way that new physics appears in top production is through radiative corrections. As presented in Ref. [4], in the two-Higgs-doublet model (2HDM) and in the minimal supersymmetric model (MSSM), the $O(\alpha m_t^2/m_W^2)$ Yukawa correction, which arises from the loops of Higgs bosons and the scalar components of virtual vector bosons, can be up to 15%. In the MSSM, in addition to this Yukawa correction, the genuine supersymmetric corrections, which arise from loops of superparticles, should also be taken into account. The dominant virtual effects of superparticles arise from the supersymmetric QCD correction of order α_s and supersymmetric electroweak correction of order $\alpha m_t^2/m_W^2$ due to a heavy top quark. The one-loop supersymmetric QCD correction has already been calculated [5] and was found to be significant in the favorable case. In this paper we present the genuine supersymmetric electroweak correction of order $\alpha m_t^2/m_W^2$, which arises from loops of the chargino, neutralino, and squarks, to top quark production at the Tevatron.

In Sec. II, we present the analytic results in terms of the well-known standard notation of one-loop Feynman integrals. In Sec. III, we present some numerical examples and discuss the implication of our results.

II. CALCULATIONS

At the Tevatron, the top quark is dominantly produced via quark-antiquark annihilation [6]. The genuine supersymmetric electroweak correction of order $\alpha m_t^2/m_W^2$ to the amplitude is contained in the correction to the vertex of top-quark color current. The relevant Feynman diagrams are shown in Fig. 1 and the Feynman rules can be found in Ref. [7]. In our calculation, we use dimensional regularization to regulate all the ultraviolet divergences in the virtual loop corrections and we adopt the on-mass-shell renormalization scheme [8]. The renormalized amplitude for $q\bar{q} \to t\bar{t}$ can be written as

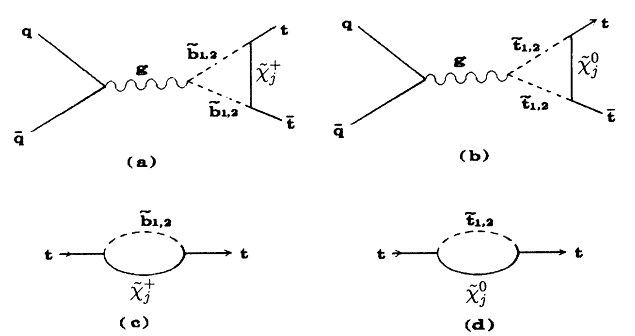

FIG. 1. Feynman diagrams: (a,b) vertex diagrams, (c,d) self-energy diagrams. Here $\tilde{t}_i, \tilde{b}_i, \tilde{\chi}_j^+$, and $\tilde{\chi}_j^0$ stand for scalar top quark, scalar bottom quark, charginos, and neutralinos, respectively.

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$$M_{\rm ren} = M_0 + \delta M \ , \tag{1}$$

where M_0 is the amplitude at the tree level and δM is the correction to the amplitude, which are given by

$$M_0 = \bar{v}(p_2)(-ig_s T^A \gamma^\nu) u(p_1) \times \frac{-ig_{\mu\nu}}{\hat{s}} \bar{u}(p_3)(-ig_s T^A \gamma^\mu) v(p_4) \ , \tag{2}$$

$$\delta M = \bar{v}(p_2)(-ig_s T^A \gamma^\nu) u(p_1) \frac{-ig_{\mu\nu}}{\hat{s}} \bar{u}(p_3) \delta\Lambda^\mu v(p_4) \ . \tag{3}$$

Here, p_1, p_2 denote the momenta of the incoming partons, and p_3, p_4 are used for outgoing t quark and its antiparticle. $\delta\Lambda^\mu$ stands for the genuine supersymmetric electroweak correction to the vertex of top-quark color current, which is given by

$$\delta\Lambda^\mu = -ig_s T^A \frac{g^2 m_t^2}{32\pi^2 m_W^2 \sin^2\beta} [\gamma^\mu F_1 + \gamma^\mu \gamma_5 F_2 + k^\mu F_3 + k^\mu \gamma_5 F_4 + ik_\nu \sigma^{\mu\nu} F_5 + ik_\nu \sigma^{\mu\nu}\gamma_5 F_6] \ , \tag{4}$$

where $\sigma^{\mu\nu} = \frac{i}{2}[\gamma^\mu, \gamma^\nu]$ and F_i are form factors which are presented in the Appendix. From the Appendix we found that all the ultraviolet divergences have canceled in the renormalized amplitude, as they should.

The renormalized differential cross section for the parton level process $q\bar{q} \to t\bar{t}$ can be written as the sum of the tree-level part and the correction part

$$\frac{d\hat{\sigma}}{d\cos\theta}(\hat{s}, \cos\theta) = \frac{d\hat{\sigma}^0}{d\cos\theta} + \frac{d(\Delta\hat{\sigma})}{d\cos\theta} \ , \tag{5}$$

where the tree-level part is given by

$$\frac{d\hat{\sigma}^0}{d\cos\theta} = \frac{8\pi\alpha_s^2}{9\hat{s}^3} \beta_t [(p_1\cdot p_3)^2 + (p_2\cdot p_3)^2 + m_t^2 p_1\cdot p_2] \ , \tag{6}$$

and the correction part is given by

$$\frac{d(\Delta\hat{\sigma})}{d\cos\theta} = \frac{8\pi\alpha_s^2}{9\hat{s}^3} \beta_t \frac{g^2 m_t^2}{32\pi^2 m_W^2 \sin^2\beta} \mathrm{Re}\{2F_1[(p_1\cdot p_3)^2 + (p_2\cdot p_3)^2 + m_t^2 p_1\cdot p_2] + m_t \hat{s}^2 F_5\} \ . \tag{7}$$

In the above θ is the angle between the top quark and an incoming quark, and

$$\begin{aligned} \beta_t &= \sqrt{1 - 4m_t^2/\hat{s}} \ , \\ \hat{s} &= (p_3 + p_4)^2 \ , \\ p_1 \cdot p_2 &= \frac{1}{2}\hat{s} \ , \\ p_1 \cdot p_3 &= \frac{1}{4}\hat{s}(1 - \beta_t \cos\theta) \ , \\ p_2 \cdot p_3 &= \frac{1}{4}\hat{s}(1 + \beta_t \cos\theta) \ . \end{aligned} \tag{8}$$

Integrating the differential cross section over $\cos\theta$ we get the parton cross section

$$\hat{\sigma}(\hat{s}) = \hat{\sigma}^0 + \Delta\hat{\sigma} \ , \tag{9}$$

with

$$\hat{\sigma}^0 = \frac{8\pi\alpha_s^2}{27\hat{s}^2} \beta_t (\hat{s} + 2m_t^2) \ , \tag{10}$$

$$\Delta\hat{\sigma} = \frac{8\pi\alpha_s^2}{9\hat{s}^3} \beta_t \frac{g^2 m_t^2}{32\pi^2 m_W^2 \sin^2\beta} \times [\tfrac{2}{3} F_1 \hat{s}(\hat{s} + 2m_t^2) + 2F_5 m_t \hat{s}^2] \ . \tag{11}$$

The hadronic cross section is obtained by convoluting the subprocess cross section $\hat{\sigma}_{ij}$ of partons i, j with parton distribution functions $f_i^A(x_1, Q), f_j^B(x_2, Q)$, which is given by

$$\begin{aligned} \sigma(S) &= \sum_{i,j} dx_1 dx_2 [f_i^A(x_1, Q) f_j^B(x_2, Q) + (A \leftrightarrow B)] \hat{\sigma}_{ij}(\hat{s}, \alpha_s(\mu)) \\ &= \sum_{i,j} \int_{\tau_0}^1 \frac{d\tau}{\tau} \left(\frac{1}{S} \frac{dL_{ij}}{d\tau} \right) (\hat{s}\hat{\sigma}_{ij}) \ , \end{aligned} \tag{12}$$

with

$$\frac{dL_{ij}}{d\tau} = \int_\tau^1 \frac{dx_1}{x_1} [f_i^A(x_1, Q) f_j^B(\tau/x_1, Q) + (A \leftrightarrow B)] \ . \tag{13}$$

In the above the sum runs over all incoming partons carrying a fraction of the proton and antiproton momenta ($p_{1,2} = x_{1,2} P_{1,2}$), $\sqrt{S} = 1.8$ TeV is the center-of-mass energy of Tevatron, $\tau = x_1 x_2$ and $\tau_0 = 4m_t^2/S$. As in Ref. [3], we do not distinguish the factorization scale Q and the renormalization scale μ and take both as the top quark mass. In order to compare our results with the Yukawa corrections in Ref. [4], we use the same parton distribution function as in Ref. [4], i.e., the Morfin-Tung leading-order parton distribution function [9].

III. NUMERICAL EXAMPLES AND DISCUSSION

In the numerical examples presented in Figs. 2–5, the supersymmetric (SUSY) parameters M, μ are fixed to be 200 and -100 GeV, respectively. The parameter tanβ and squark mass will be retained as variables and the dependence on them will be shown in our numerical examples. The mass eigenstates of each flavor squark are obtained by the mixing of left- and right-handed squark eigenstates with mixing angle θ [7]. Since the off-diagonal elements of squark mass matrices are proportional to the corresponding quark mass [10], the mixing effects for scalar top squarks may be significant. In general the off-diagonal elements of the mass matrix for scalar top quarks can be parametrized as $-m_t(A_t + \mu\cot\beta)$ [11]. But in our numerical examples, for simplicity, we only consider the unmixing case corresponding to $A_t = -\mu\cot\beta$. Furthermore, we do not consider the mass splitting between squarks of different flavors and denote the squark mass by $\tilde{m}_q (\equiv \tilde{m}_{t_{1,2}} = \tilde{m}_{b_{1,2}})$.

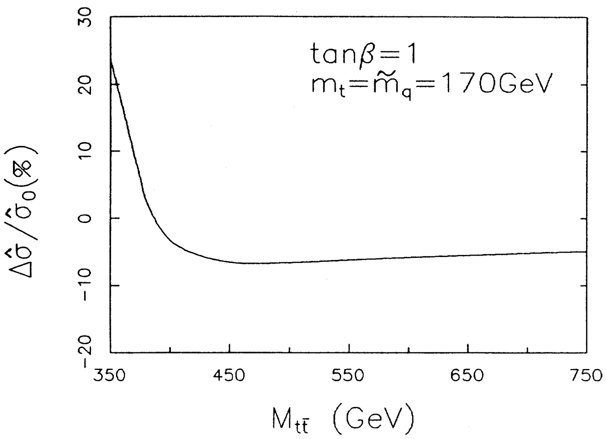

FIG. 2. Relative correction to parton $q\bar{q} \to t\bar{t}$ cross section versus $t\bar{t}$ invariant mass for $m_t = \tilde{m}_q = 170$ GeV and $\tan\beta = 1$.

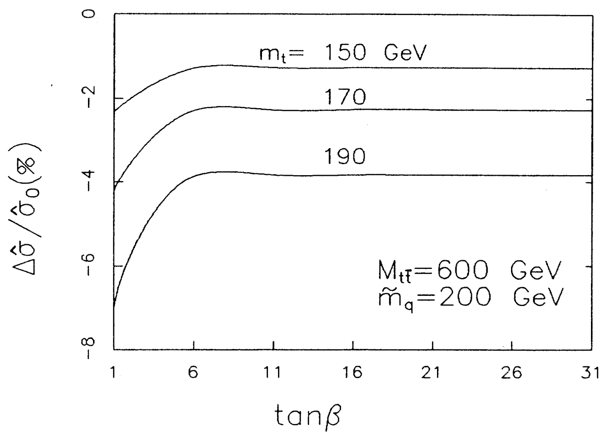

FIG. 4. Same as Fig. 2, but versus $\tan\beta$.

Other input parameters are [12] $m_Z = 91.176$ GeV, $\alpha_{\rm em} = 1/128.8$, and $G_F = 1.166\,372 \times 10^{-5}$ (GeV)$^{-2}$. m_W is determined through [13]

$$m_W^2\left(1 - \frac{m_W^2}{M_Z^2}\right) = \frac{\pi\alpha}{\sqrt{2}G_F}\frac{1}{1-\Delta r}\ , \qquad (14)$$

where, to order $O(\alpha m_t^2/m_W^2)$, Δr is given by [14]

$$\Delta r \sim -\frac{\alpha N_C c_W^2 m_t^2}{16\pi^2 s_W^4 m_W^2}\ . \qquad (15)$$

Figure 2 is an example of the relative correction to the parton cross section as a function of $t\bar{t}$ invariant mass $M_{t\bar{t}} \equiv \sqrt{\hat{s}}$. In a small region ($M_{t\bar{t}} < 450$ GeV), the corrections depend strongly on $M_{t\bar{t}}$ and can enhance the Born cross section significantly when close to the $t\bar{t}$ production threshold. But in the remaining large region ($M_{t\bar{t}} > 450$ GeV), the correction is negative and insensitive to $M_{t\bar{t}}$. The dependence of the parton cross section on squark mass is shown in Fig. 3 for $M_{t\bar{t}} = 400$ GeV. In a region of $\tilde{m}_q < 350$ GeV, the correction is very sensitive to squark mass and can be positively large for $\tilde{m}_q < 150$ GeV. The correction gets its negative maximum size at $\tilde{m}_q = 350$ GeV, then decreases rapidly with squark mass and vanishes for $\tilde{m}_q > 500$ GeV. Figure 4 is an example showing the dependence on $\tan\beta$. Apart from a small region ($\tan\beta < 6$), the correction is not sensitive to $\tan\beta$.

The relative correction to the hadronic cross section as a function of squark mass is presented in Fig. 5 for different top quark masses. The correction increases with top quark mass. For $m_t = 120$ GeV the correction is negligibly small, but when m_t increases to 170 GeV the correction can be quite significant for favorable squark masses. For $m_t = 170$ GeV the correction can reach $+20\%$ for $\tilde{m}_q = 100$ GeV and -10% for $\tilde{m}_q = 350$ GeV. For $\tilde{m}_q < 350$ GeV, the correction is very sensitive to squark mass. In the range of $\tilde{m}_q > 350$ GeV, the correction drops rapidly with the increase of squark mass and finally tends to zero, showing the decoupling effects.

In Ref. [4], the Yukawa correction from the SUSY Higgs sector to the hadronic cross section was found to be very small for $\tan\beta = 1$, on the order of a percent. So the

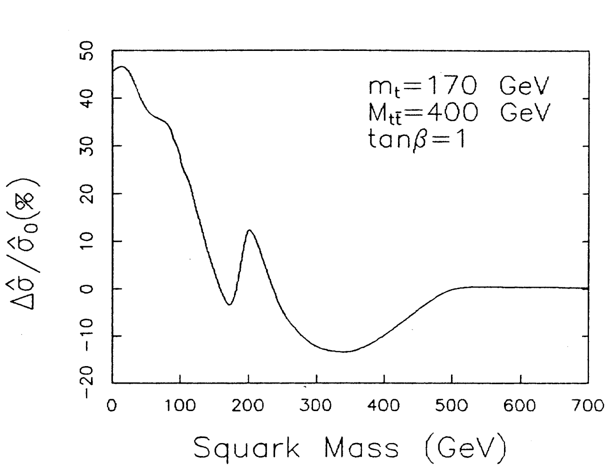

FIG. 3. Same as Fig. 2, but versus squark mass.

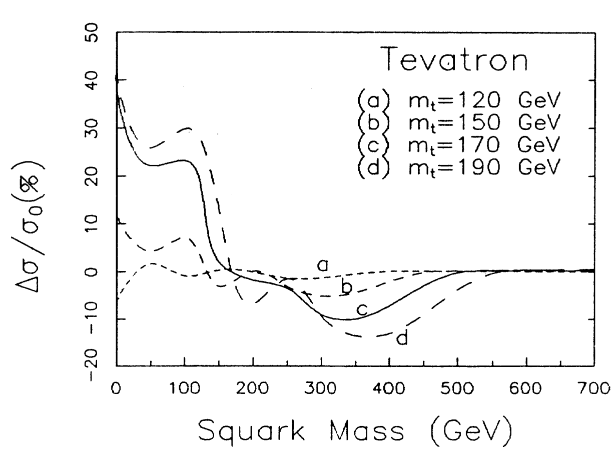

FIG. 5. Relative correction to hadronic cross section at Tevatron versus squark mass for $\tan\beta = 1$.

genuine supersymmetric electroweak correction is much larger than the Yukawa correction for favorable squark masses. This phenomenon also occurs in the corrections to the $Zb\bar{b}$ vertex [15] and top pair production at the e^+e^- collider [16]. It should be noted that more free parameters are involved in the genuine supersymmetric electroweak correction than in the Yukawa correction from the SUSY Higgs sector. The parameters in the genuine supersymmetric electroweak correction are the squark mass matrices, $\tan\beta, M$, and μ. In general, the correction can be quite large or negligibly small, depending on the values of all these parameters.

Charginos (neutralinos) are the superpartners of charged (neutral) Higgs bosons and vector bosons. The genuine supersymmetric electroweak correction arises from the loops of charginos and neutralinos. The Yukawa correction calculated in Ref. [4] arises from loops of charged and neutral Higgs bosons and vector bosons. An old theorem [17] says the anomalous magnetic moment for a spin-$\frac{1}{2}$ fermion vanishes in the SUSY limit. Away from the SUSY limit the cancellations have somewhat less effect. From Fig. 5 and the results of Ref. [4] we found that if the squark and charged Higgs boson have about the same mass, the genuine supersymmetric electroweak correction and the Yukawa correction have the opposite sign. But their effects will cancel only to a small extent since the genuine supersymmetric electroweak correction is much larger than the Yukawa correction. Thus the combined effects are much larger than the SM electroweak correction which is a few percent, and comparable to the supersymmetric QCD correction in the favorable case.

The correction can be even larger for smaller values of $\tan\beta$ and lighter squarks. For example, the correction can reach 30% for $m_t = 170$ GeV. If squarks are lighter than 10 GeV. But the experiment at the Tevatron has pushed the lower bound of squark mass to rather high values. The lower limit on squark mass from the D0 Collaboration is about 150 GeV [18]. The CDF limit [19] on squark mass without cascade decays is $\tilde{m}_q > 126$ GeV. The lower limit of $\tan\beta$ is 0.6 from perturbative bounds [20]. Reference [4] argues for lower values of $\tan\beta$ from perturbative unitarity than Ref. [20].

In conclusion, we presented the genuine supersymmetric electroweak corrections of order $\alpha m_t^2/m_W^2$ to top quark pair production at the Tevatron in the minimal supersymmetric model. The corrections can enhance the hadronic cross section by 20% for top quark mass of 170 GeV and squark mass of 100 GeV. But when squark mass gets larger, the corrections decrease rapidly. Just as the Yukawa correction presented in Ref. [4], if we assume a theoretical uncertainty in the cross section of $\pm 20\%$ from QCD [22], these supersymmetric corrections in the favorable case are potentially observable and could be used to place restrictions on MSSM.

ACKNOWLEDGMENTS

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APPENDIX

The form factors F_i are obtained by

$$F_i = F_i^c + F_i^n , \tag{A1}$$

where F_i^c are given as

$$F_1^c = \sum_{j=1,2} V_{j2}V_{j2}^* \left[c_{24} + m_t^2(c_{11} + c_{12}) + \frac{1}{2}B_1(m_t, \tilde{M}_j, \tilde{m}_b) + m_t^2 \frac{\partial B_1(p, \tilde{M}_j, \tilde{m}_b)}{\partial p^2}\bigg|_{p^2=m_t^2} \right] , \tag{A2}$$

$$F_2^c = \sum_{j=1,2} V_{j2}V_{j2}^* \left[c_{24} + \frac{1}{2}B_1(m_t, \tilde{M}_j, \tilde{m}_b) \right] , \tag{A3}$$

$$F_3^c = \frac{1}{2}m_t \sum_{j=1,2} V_{j2}V_{j2}^*(c_{21} - 2c_{23}) , \tag{A4}$$

$$F_4^c = \frac{1}{2}m_t \sum_{j=1,2} V_{j2}V_{j2}^*(c_{21} + 4c_{22} - 4c_{23}) , \tag{A5}$$

$$F_5^c = -\frac{1}{2}m_t \sum_{j=1,2} V_{j2}V_{j2}^*(c_{11} + c_{21}) , \tag{A6}$$

$$F_6^c = -\frac{1}{2}m_t \sum_{j=1,2} V_{j2}V_{j2}^*(c_{11} + c_{21} - 2c_{12} - 2c_{23}) , \tag{A7}$$

and F_i^n are given as

$$F_1^n = \sum_{j=1}^{4} N_{j4}N_{j4}^* \left[2c_{24} + 2m_t^2(c_{11} + c_{12}) + B_1(m_t, \tilde{M}_{0j}, \tilde{m}_t) + 2m_t^2 \frac{\partial B_1(p, \tilde{M}_{0j}, \tilde{m}_t)}{\partial p^2}\bigg|_{p^2=m_t^2} \right] , \tag{A8}$$

$$F_3^n = m_t \sum_{j=1}^{4} N_{j4}N_{j4}^*(c_{21} - 2c_{23}) , \tag{A9}$$

$$F_5^n = -m_t \sum_{j=1}^{4} N_{j4}N_{j4}^*(c_{11} + c_{21}) , \tag{A10}$$

$$F_2^n = F_4^n = F_6^n = 0 . \tag{A11}$$

The functions $c_{ij}(-p_3, p_3+p_4, \tilde{M}_j, \tilde{m}_b, \tilde{m}_b)$ in Eqs. (A2)–(A7), $c_{ij}(-p_3, p_3+p_4, \tilde{M}_{0j}, \tilde{m}_t, \tilde{m}_t)$ in Eqs. (A8)–(A10), and B_1 in (A2) and (A8) are the three-point and two-point Feynman integrals, the definition and expression for which can be found in Ref. [21].

In the above, V_{ij} are the elements of 2×2 matrix V which can be found in Ref. [7]. The chargino masses are given by [7]

$$\tilde{M}_{1,2} = \frac{1}{2}\{M^2 + \mu^2 + 2m_W^2 \pm[(M^2-\mu^2)^2 + 4m_W^4\cos^2 2\beta + 4m_W^2(M^2+\mu^2+2M\mu\sin 2\beta)]^{1/2}\} . \quad \text{(A12)}$$

N_{ij} are the elements of 4×4 matrix N and the neutralino masses $\tilde{M}_{0i}$ are the elements of the diagonal mass matrix N_D satisfying

$$N^* Y N^{-1} = N_D , \quad \text{(A13)}$$

where N_D is a diagonal mass matrix with non-negative entries and Y is given by

$$Y = \begin{pmatrix} M' & 0 & -m_Z s_W\cos\beta & m_Z s_W \sin\beta \\ 0 & M & m_Z c_W\cos\beta & -m_Z c_W\sin\beta \\ -m_Z s_W\cos\beta & m_Z c_W\cos\beta & 0 & -\mu \\ m_Z s_W\sin\beta & -m_Z c_W\sin\beta & -\mu & 0 \end{pmatrix} . \quad \text{(A14)}$$

Giving the values for the parameters $M, M', \mu, \tan\beta$, the matrices N and N_D can be obtained numerically. Here, the parameters M, M' are the masses of gauginos corresponding to SU(2) and U(1), respectively. μ is the coefficient of the H_1, H_2 mixing term in the superpotential and $\tan\beta = v_2/v_1$ is the ratio of vacuum expectation values of the two Higgs doublets H_1, H_2. With the grand unification assumption, i.e., SU(2)$\times$U(1) is embedded in a grand unified theory, we have the relation $M' = \frac{5}{3}(g'^2/g^2)M$. So the chargino masses and neutralino masses only depend on $M, \mu, \tan\beta$. For $M = 200$ GeV, $\mu = -100$ GeV, and $\tan\beta = 1$, the chargino masses are (120,220) GeV and the neutralino masses are (128.8, 84.6, 220.8, 100) GeV.

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