A Light Scalar Top Quark in MSGUT

Tadashi Kon† and Toshihiko Nonaka ‡

Faculty of Engineering, Seikei University, Tokyo 180, Japan
†Department of Physics, Rikkyo University, Tokyo 171, Japan

Abstract

We reexamine a possibility for the existence of a light supersymmetric partner of the top quark (stop) with mass \(15\sim16\text{GeV}\) in the framework of the minimal supergravity GUT model (MSGUT). Such light stop could explain the slight excess of the high \(p_T\) cross section of the \(D^{*\pm}\)-meson production in the two-photon process at TRISTAN. We find two types of solution for the RGEs in the MSGUT allowing the existence of the light stop. The type I [type II] solution is characterized by the heavy [light] top quark, \(m_t \simeq 150\text{GeV} [100\text{GeV}]\), and the light [heavy] squarks, \(m_{\tilde{q}} \simeq 150\text{GeV} [300\text{GeV}]\). It is found that the type II solution is more favorable because these parameter sets seem to satisfy all constraints settled by the recent collider experiments as well as by the non-accelerator observations and the cosmological considerations. We point out that the existence of such a stop could change the dominant decay mode of sparticles, the top quark and the Higgs bosons. Consequently, the present experimental bounds on the supersymmetric parameters as well as on masses of the top and the Higgs could be weakened substantially. However, the allowed parameter region is rather restricted and in turn masses and mixing parameters of the other SUSY partners as well as masses of the Higgs and the top are severely constrained. For example, \(75\text{GeV} \lesssim m_{\tilde{g}} \lesssim 85\text{GeV}, m_{\tilde{W}_1} \lesssim 50\text{GeV}, m_{\tilde{\ell}} \simeq m_{\tilde{q}} \simeq 300\text{GeV}, \theta_t \simeq 0.9, m_h \lesssim 65\text{GeV} \) and \(90\text{GeV} \lesssim m_t \lesssim 100\text{GeV} \).

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† e-mail address : d34477@jpnac.bitnet, mtsk@jpnkektr.bitnet
‡ e-mail address : rra001c@jpnrky00.bitnet
1 Introduction

In spite of much efforts of many experimental colleagues the nature has not yet opened "the new physics windows" and the standard model (SM) of the electroweak interaction has been established precisely except for the lack of the top quark and the Higgs boson. The mass of top quark is now expected to be larger than about 100 GeV and the other quarks and the leptons are known to be much lighter than the top \cite{1}. One of the models beyond the SM, the supersymmetric (SUSY) standard models \cite{2}, draw much attention since they can solve the naturalness problem in the SM and can achieve naturally the unification of gauge interactions at the GUT scale \cite{3}. It is known that the naturalness criterion in the SUSY models inevitably predicts the existence of many SUSY partners (sparticles) with masses below the order of magnitude of one TeV. Therefore, the search for the sparticles is one of the most important purposes of the present and future collider experiments. Up to now, there is no clear evidence for the existence of the sparticles at the present collider experiments and the lower mass bounds for them are growing up day by day. Should we wait for LEP-II, LHC or some next linear colliders to open the SUSY window?

In the minimal SUSY standard model (MSSM) the slepton $\tilde{\ell}$, the lighter chargino $\tilde{W}_1$ or the scalar top quark (stop) $\tilde{t}_1$ is expected to be the lightest charged sparticle. While the model independent lower bounds on masses of the sleptons and the chargino have been settled by LEP, $m_{\tilde{\ell}}$, $m_{\tilde{W}_1} > m_Z/2$, there is a window allowing the existence of a very light stop, $m_{\tilde{t}_1} < m_Z/2$, even if we have no evidence at LEP \cite{4} and Tevatron \cite{5}. From a theoretical point of view, moreover, it is argued that one of stops $\tilde{t}_1$ is naturally lighter than the other squarks and the sleptons because of the large top mass \cite{6, 7}. Such a possibility is very interesting because this is the SUSY counterpart of the fact that the leptons and the other quarks are much lighter than the top. We should notice such a possibility seriously because the stop could be discovered first at the collider experiments even before the top.

Enomoto et al. in the TOPAZ group at TRISTAN have reported a slight excess of the high $p_T$ cross section of $D^{*\pm}$-meson production in a two-photon process \cite{8}. The disagreement between the measured value and the standard model prediction now becomes 3.5$\sigma$ level \cite{9}, which should be compared to 1.5$\sigma$ reported previously \cite{8}. Although there remains a possibility that such a excess could be explained by the large contribution of the gluon structure function of the photon \cite{10}, there is another exciting way to interpret this enhancement, i.e., it is the signature of the pair production of the stop with mass $\lesssim$ 20 GeV. Since such a light stop will decay into the charm-quark plus the lightest neutralino \cite{7}, the signature of the stop production will be the charmed meson production with large missing energies. This signature would resemble the charmed-hadron production in the two-photon process at $e^+e^-$ colliders. Enomoto et al. have pointed out that the stop with mass about 15 $\sim$ 16 GeV and the neutralino with mass about 13 $\sim$ 14 GeV could well explain the experimental data.

In previous papers \cite{11}, we have investigated the possibility for the existence of the stop with mass 15$\sim$16 GeV in the framework of the minimal supergravity GUT model (MSGUT) \cite{12}. We have pointed out that the existence of such a stop could change the dominant decay mode of some particles, especially that of the gluino, and could weaken the present experimental bound on the SUSY parameters. We have found the light stop solutions

2
for the renormalization group equations (RGEs) \cite{13} and such solutions inevitably set constraints on masses and mixing parameters of the other SUSY partners as well as masses of the Higgs and the top, for example, \( m_\tilde{q} \lesssim 85 \text{GeV} \), \( m_\tilde{W} \lesssim 55 \text{GeV} \), \( 110 \text{GeV} < m_\tilde{\ell} \lesssim 140 \text{GeV} \), \( 120 \text{GeV} < m_\tilde{\nu} \lesssim 160 \text{GeV} \), \( \theta_t \approx 0.9 \), \( m_H \lesssim 65 \text{GeV} \) and \( m_t \lesssim 135 \text{GeV} \).

In this paper we give a new set of solutions for the RGEs to allow the light stop \( m_\tilde{t} \lesssim 15 \text{GeV} \) called type II solution in addition to the solution obtained previously (type I), where we also present our calculational scheme in detail. The type II solution is characterized by the heavy squarks and sleptons, \( m_\tilde{\ell} \approx m_\tilde{\nu} \approx 300 \text{GeV} \), and the light top, \( m_t \lesssim 100 \text{GeV} \). We show this type of solution is more favorable if we consider the recent collider data as well as the nucleon decay and the dark matter constraints.

The paper is organized as follows. The theoretical bases in the MSSM for lightness of the stop are summarized in Sec.2. In Sec.3 we reconsider the present bounds on the stop mass and those on the gaugino parameters. The two types of solution for the RGEs in the MSGUT are given in Sec.4. In Sec.5 we choose favorable solution by a consideration of some experimental informations. Some phenomenological implications of the light stop are discussed in Sec.6. Sec.7 is devoted to summary and conclusion.

2 Light stop : its theoretical bases

In the framework of the MSSM \cite{2}, the stop mass matrix in the \((\tilde{t}_L, \tilde{t}_R)\) basis is expressed by

\[
\mathcal{M}_t^2 = \begin{pmatrix}
m^2_{\tilde{t}_L} & a_t m_t \\
a_t m_t & m^2_{\tilde{t}_R}
\end{pmatrix},
\]

where \( m_t \) is the top mass. The SUSY mass parameters \( m_{\tilde{t}_L,R} \) and \( a_t \) are parametrized in the following way \cite{14}:

\[
m^2_{\tilde{t}_L} = \tilde{m}^2_{Q_3} + m^2_Z \cos 2\beta \left( \frac{1}{2} - \frac{2}{3} \sin^2 \theta_W \right) + m_t^2,
\]
\[
m^2_{\tilde{t}_R} = \tilde{m}^2_{U_3} + \frac{2}{3} m^2_Z \cos 2\beta \sin^2 \theta_W + m_t^2,
\]
\[
a_t = A_t + \mu \cot \beta,
\]

where \( \tan \beta, \mu \) and \( A_t \) denote the ratio of two Higgs vacuum expectation values \((= v_2/v_1)\), the SUSY Higgs mass parameter and the trilinear coupling constant, respectively. The soft breaking masses of the third generation doublet \( \tilde{m}_{Q_3} \) and the up-type singlet \( \tilde{m}_{U_3} \) squarks are related to those of the first (and second) generation squarks as

\[
\tilde{m}^2_{Q_3} = \tilde{m}^2_{Q_1} - \tilde{I},
\]
\[
\tilde{m}^2_{U_3} = \tilde{m}^2_{U_1} - 2\tilde{I},
\]

where \( \tilde{I} \) is a function proportional to the top quark Yukawa coupling \( \alpha_t \) and is determined by the RGEs in the MSGUT (see Sec.4 and Appendices A and B). Throughout of this paper we adopt the notation of Ref.\cite{14}.
There are two origins for lightness of the stop compared to the other squarks and sleptons, i) smallness of the diagonal soft masses \( m^2_{\tilde{t}_L} \) and \( m^2_{\tilde{t}_R} \) and ii) the left-right stop mixing. Both effects are originated from the large Yukawa interaction of the top. The origin i) can be easily seen from Eqs.\((2)\)∼\((6)\). The diagonal mass parameters \( m^2_{\tilde{t}_L} \) and \( m^2_{\tilde{t}_R} \) in Eq.\((1)\) have possibly small values owing to the negative large contributions of \( \tilde{I} \) proportional to \( \alpha_t \) in Eqs.\((5)\) and \((6)\). It should be noted that this contribution is also important in the radiative SU(2)×U(1) breaking in the MSGUT. The Higgs mass squared has an expression similar to Eqs.\((5)\) and \((6)\);

\[
\tilde{m}_{H_2}^2 = \tilde{m}_{L_1}^2 - 3\tilde{I},
\]

where \( \tilde{m}_{L_1}^2 \) denotes the soft breaking mass of the first generation doublet slepton. The large contribution of \( \tilde{I} \) enables \( \tilde{m}_{H_2}^2 \) to become negative at an appropriate weak energy scale. In order to see another origin ii) we should diagonalize the mass matrix (1). The mass eigenvalues are obtained by

\[
m^2_{\tilde{t}_1} = m^2_{\tilde{t}_L} + m^2_{\tilde{t}_R} \mp (m^2_{\tilde{t}_L} - m^2_{\tilde{t}_R})^2 + (2a_t m_t)^2 \right)^{1/2},
\]

and the corresponding mass eigenstates are expressed by

\[
\begin{pmatrix}
\tilde{t}_1 \\
\tilde{t}_2
\end{pmatrix} = \begin{pmatrix}
\tilde{t}_L \cos \theta_t - \tilde{t}_R \sin \theta_t \\
\tilde{t}_L \sin \theta_t + \tilde{t}_R \cos \theta_t
\end{pmatrix},
\]

where \( \theta_t \) denotes the mixing angle of stops.

\[
\tan \theta_t = \frac{a_t m_t}{m^2_{\tilde{t}_L} - m^2_{\tilde{t}_R}}.
\]

We see that if the SUSY mass parameters and the top mass are of the same order of magnitude, small \( m^2_{\tilde{t}_1} \) is possible owing to the cancellation in the expression (3) \( \tilde{t}_1 \).

After the mass diagonalization we can obtain the interaction Lagrangian for the mass eigenstate \( \tilde{t}_1 \). We note, in particular, that the stop coupling to the Z-boson \( \tilde{t}_1 \tilde{t}_1^* Z \) depends sensitively on the mixing angle \( \theta_t \). More specifically, it is proportional to

\[
C_{\tilde{t}_1} \equiv \frac{2}{3} \sin^2 \theta_W - \frac{1}{2} \cos^2 \theta_t.
\]

Note that for the special value of \( \theta_t \sim 0.98 \), the Z-boson coupling completely vanishes [4].

\section{Present bounds on stop mass and gaugino parameters}

\subsection{Stop mass bounds}

Before a discussion of experimental bounds on the stop mass \( m_{\tilde{t}_1} \), we examine the decay modes of the stop. In the MSSM, the stop lighter than the other squarks and gluino can decay into the various final states : \( \tilde{t}_1 \rightarrow t \tilde{Z}_1 \) (a), \( b\tilde{W}_1 \) (b), \( b\ell\tilde{\nu} \) (c), \( b\nu\tilde{\ell} \) (d), \( bW\tilde{Z}_1 \) (e), \( bf'\tilde{Z}_1 \) (f), \( c\tilde{Z}_1 \) (g), where \( \tilde{Z}_1, \tilde{W}_1, \tilde{\nu} \) and \( \tilde{\ell} \), respectively, denote the lightest neutralino, the
lighter chargino, the sneutrino and the charged slepton. If we consider the light stop with mass smaller than 20GeV, the first five decay modes (a) to (e) are kinematically forbidden due to the model independent lower mass bounds for respective particles: \( m_{\tilde{t}_1} \lesssim 60 \text{GeV} \), \( m_{\tilde{\nu}_1} \lesssim 45 \text{GeV} \), \( m_{\tilde{\chi}_1^0} \lesssim 45 \text{GeV} \) and \( m_{\tilde{\chi}_1^+} \lesssim 40 \text{GeV} \). So (f) and (g) left. Hikasa and Kobayashi \cite{4} have shown that the one-loop mode \( \tilde{t}_1 \to c\tilde{Z}_1 \) (g) dominates over the four-body mode \( \tilde{t}_1 \to bf'\tilde{Z}_1 \) (f). This is absolutely true in the case considered here, because the mode (f) is negligible by the kinematical suppression, \( m_{\tilde{t}_1} \sim m_{\tilde{Z}_1} + m_b \). So we can conclude that such a light stop will decay into the charm quark jet plus the missing momentum taken away by the neutralino with almost 100% branching ratio. Note that the width of stop in this case is very small, i.e., of the order of magnitude of eV.

Naively, it will be expected that Tevatron and/or LEP can set severe bounds on the stop mass through the processes: \( gg \to \tilde{t}_1\tilde{t}_1^* \to c\tilde{Z}_1\tilde{Z}_1 \) (Tevatron) and/or \( Z \to \tilde{t}_1\tilde{t}_1^* \) (LEP). However, the situation is not so obvious. Baer et al. \cite{5} have performed the analyses of the experimental data of 4pb\(^{-1}\) integrated luminosity Tevatron running, and have obtained the results that the stop could easily escape detection if \( m_{\tilde{Z}_1} \gtrsim 10 \text{GeV} \). Such a large neutralino mass could make the charm quark jets softer. Consequently the stop production cross section plotted against the missing transverse energy becomes smaller than the present upper bounds, where they impose cuts on the missing transverse energy \cite{2}. Moreover, we should point out that LEP cannot exclude a light stop \( m_{\tilde{t}_1} < m_{\tilde{Z}_1}/2 \) by measuring the extra width of the Z-boson. In fact we find that there is no bound on the stop mass if the mixing angle \( \theta_t \) is larger than about 0.7, where we take the experimental limit \( \Delta \Gamma_Z < 28 \text{MeV} \) (95% C.L.) \cite{14}. The origin of such a sensitivity of \( \Gamma(Z \to \tilde{t}_1\tilde{t}_1^*) \) is in the fact that the \( \tilde{t}_1\tilde{t}_1^*Z \) coupling is proportional to \( C_{\tilde{t}_1} \) \cite{14}.

Okada \cite{17} has investigated possible bounds on masses of the stop and the neutralino from the experimental data of the \( b \to s\gamma \) decay. He has shown that the light stop with mass \( m_{\tilde{t}_1} \lesssim 20 \text{GeV} \) has not been excluded by the data. It has been also pointed out by Fukugita et al. \cite{18} that the existence of the light stop does not conflict with the experimental bounds on \( \Delta \rho \) and \( K^0 \to \bar{K}^0 \) mixing.

Recently, the stringent limit comes from the direct searches for the stop at \( e^+e^- \) colliders. New analyses of the direct search by the VENUS group at TRISTAN show that the TOPAZ stop \( (m_{\tilde{t}_1} = 15 \sim 16 \text{GeV} \) and \( m_{\tilde{\chi}_1^0} = 13 \sim 14 \text{GeV} \)) just confronts the experimental bounds \cite{19}. In fact it seems that such a stop has been excluded for \( m_{\tilde{\chi}_1^0} - m_{\tilde{Z}_1} > 3 \text{GeV} \). Moreover, the OPAL group have reported the bounds on the \( (m_{\tilde{t}_1}, m_{\tilde{\chi}_1^0}, \theta_t) \) plane from the direct searches at LEP \cite{20}. They have shown the light stop \( m_{\tilde{t}_1} < m_{\tilde{Z}_1}/2 \) only survives if \( m_{\tilde{t}_1} - m_{\tilde{\chi}_1^0} < 2.2 \text{GeV} \) and \( 1.15 > \theta_t > 0.85 \). A similar bound has been reported by the DELPHI group \cite{21}.

We should say, after all, the TOPAZ stop has not been excluded only if \( m_{\tilde{t}_1} - m_{\tilde{\chi}_1^0} < 2.2 \text{GeV} \) and \( 1.15 > \theta_t > 0.85 \). In the following, we examine such a possibility in the framework of the MSGUT and investigate its phenomenological implications.

\footnote{This is the bound independent on the decay mode of the top. In fact, the CDF bound, \( m_t > 110 \text{GeV} \), will not be applicable in the case of the light stop, as we will discuss in Sec.6}
\footnote{Recent analyses has been reported in Ref.10}
3.2 bounds on gaugino parameters

In the MSSM, masses and mixing parameters of the gaugino-higgsino sector are determined by three parameters $\mu$, $\tan \beta$ and $M_2$, where $M_2$ denotes the soft breaking SU(2) gaugino mass. Some regions in the $(\mu, \tan \beta, M_2)$ parameter space have already excluded by the negative searches for the SUSY particles at some collider experiments. First, we consider the experimental data at LEP; lower bound on the mass of lighter chargino,

$$m_{\widetilde{W}_1} > 45 \text{GeV},$$

(12)

upper bound on the branching ratio of the visible neutralino modes of the $Z$,[21, 22]

$$\text{BR}(Z \to \text{vis.}) \equiv \sum_{i,j} \frac{\Gamma(Z \to \widetilde{Z}_i \widetilde{Z}_j)}{\Gamma_{\text{tot}}} < 5 \times 10^{-6},$$

(13)

and upper bound on the invisible width of the $Z$,[23]

$$\Gamma(Z \to \widetilde{Z}_1 \widetilde{Z}_1) < 16.2 \text{MeV},$$

(14)

Next we should discuss bounds on the gaugino parameters from the hadron collider experiments. If we assume the GUT relation,[24]

$$m_{\widetilde{g}} = M_3 = \frac{\alpha_s}{\alpha} \sin^2 \theta_W M_2$$

(15)

in the MSGUT, the gluino mass $m_{\widetilde{g}}$ bounds from the hadron colliders could be converted into the bounds on $M_2$.[24] Naively accepted gluino mass bound at CDF is

$$m_{\widetilde{g}} > 150 \text{GeV \ (90\% C.L.)},$$

(16)

which can be easily converted into the bound on $M_2$ by Eq.(15) as $M_2 > 44 \text{GeV}$. (Note that the GUT relation (15) depends sensitively on $\sin^2 \theta_W$ and $\alpha_s$. Here we take $\sin^2 \theta_W = 0.232$ and $\alpha_s = 0.113$.)

In Fig.1 we show the region excluded by the experimental data (12), (13), (14) and (16) in the $(\mu, M_2)$ plane for $\tan \beta = 2$, where the regions inside of each contour have been excluded. We also plot a contour of $m_{\widetilde{Z}_1} = 13 \sim 14 \text{GeV}$ which can explain the TRISTAN data as mentioned above. At first sight, the neutralino with mass $13 \sim 14 \text{GeV}$ has already been excluded by the bounds from LEP and Tevatron experiments. (Note that the contour of $m_{\widetilde{Z}_1} = 13 \sim 14 \text{GeV}$ lies also in the excluded region for $\mu > 0$.) We see that the stringest bound comes from the constraint on the visible width of the $Z$-boson at LEP[13] [B] or from the gluino search at CDF[10] [D]. Fortunately, however, these bounds are rather model dependent. As we will show in Sec.5, the existence of the light stop can weaken both bounds substantially. Here we should adopt the model independent bounds from the lower mass bound on the lighter chargino [2] [A] and upper bound on the invisible width of the $Z$-boson [14] [C]. Then we find the allowed region, $-170 \text{GeV} \lesssim \mu \lesssim -20 \text{GeV}$ for $\tan \beta = 2$. If we take larger (smaller) values of $\tan \beta$, the allowed region becomes narrower (wider). We find that the allowed region disappears for $\tan \beta \gtrsim 2.8$. Furthermore, we see that $m_{\widetilde{Z}_1} = 13 \sim 14 \text{GeV}$ corresponds to $M_2 = 22 \sim 24 \text{GeV}$ in the allowed region and we can find that this correspondence is independent of the values of $\tan \beta$. Consequently, we can take $M_2 = 22 \sim 24 \text{GeV}$ as an input value in the following calculation. The allowed region in the $(\mu, \tan \beta)$ plane fixed by $M_2 = 22 \text{GeV}$ or $24 \text{GeV}$ is shown in Fig.2. It is worth mentioning that the lightest neutralino $\widetilde{Z}_1$ is almost the photino $\widetilde{\gamma}$ in the allowed parameter range in Fig.2. In fact, the photino component of the neutralino is larger than 98% in the range.
4 MSGUT analysis

4.1 solutions for RGEs

Before presenting our results for analysis, we summarize briefly the calculational scheme in the MSGUT [13]. All physical parameters go from the GUT scale $M_X$ down to low energies governed by the RGEs [13]. All RGEs for the MSGUT parameters are presented in Appendix A. Here we neglect all Yukawa couplings except for the top. This is not a bad approximation as long as $\tan \beta$ is not too large ($\ll m_t/m_b$), which is the case we consider here, $\tan \beta \sim 2.8$. In order to solve the RGEs we must fix the boundary conditions for independent parameters. At $M_X$ we set the following conditions on the parameters; the SUSY Higgs mass parameter $\mu(M_X) = \mu_\infty$ and three soft breaking mass parameters: the common scalar mass $\tilde{m}^2_f(M_X) = \tilde{m}^2_{H_1}(M_X) = m^2_\infty$, the gaugino mass $f^{-1}_5(M_X) = f^{-1}_2(M_X) = M_\infty$ and the trilinear coupling $A_t(M_X) = A_b(M_X) = A_f(M_X) = \cdots = A_\infty$. As usual, we take the Higgs mixing parameter $B$ as $B(M_X) = A_\infty - m_\infty$. We present the analytical formulae for solutions of the RGEs in the general case $f_i \neq 1$ in Appendix B. Hereafter we take the GUT relation [15] i.e., $f_1 = f_2 = f_3 = 1$, and discuss briefly the case $f_3 \neq 1$ in Appendix C.

As for the evolution of the gauge couplings $\alpha_i(t)$ and the gaugino masses $M_i(t)$, we take the input values $\alpha^{-1}(m_Z) = 128.8$ and $\sin^2 \theta_W = 0.232$. Assuming that the SUSY scale is not too different from $m_Z$, we may use the SUSY beta function at all scales above $m_Z$ for simplification. Then one finds $M_X = 2.1 \times 10^{16}$GeV, $\alpha^{-1}_\infty = \alpha^{-1}_3(M_X) = \alpha^{-1}_2(M_X) = \alpha^{-1}_1(M_X) = 24.6$ and $\alpha_3(m_Z) = 0.113$.

All physics at weak scale $m_Z$ are determined by the six parameters ($m_\infty$, $A_\infty$, $M_\infty$, $\mu$, $\tan \beta$, $m_t$). There are, moreover, two conditions imposed on the parameters to have the correct scale of SU(2)$\times$U(1) breaking. So we can reduce the number of the independent parameters to four out of the six. Here we take the four independent input parameters as ($M_\infty$, $\mu$, $\tan \beta$, $m_t$). As we have discussed in Sec.3.2, furthermore, we can fix one of the input values, $M_2 = 22 \sim 24$GeV, which corresponds to $M_\infty = 26.7 \sim 29.1$GeV for $\sin^2 \theta_W = 0.232$. After all, there remain the only three arbitrary parameters ($\mu$, $\tan \beta$, $m_t$).

4.2 two types of solutions

We seek numerical solutions for the RGEs to give a light stop with mass $m_{\tilde{t}_1} = 15 \sim 16$GeV varying the three parameters ($\mu$, $\tan \beta$, $m_t$). Then we find two types of solution, type I and type II. The type I solution has already been presented at the previous papers [11]. Common properties of both solutions are i) mass of lighter chargino $\tilde{W}_1$ is smaller than 55GeV, ii) the lighter CP even neutral Higgs boson is relatively light, $m_h \lesssim 65$GeV and iii) the stop mixing angle $\theta_t$ is severely limited as $\theta_t \simeq 0.9$. It is interesting that $\theta_t \simeq 0.9$ is not input but output of the MSGUT calculation. In fact this value can save the light stop from the recent bounds from the direct searches at LEP [20, 21] as discussed in Sec.3.1. The differences between the two types of solution are in masses of the top quark and the squarks (sleptons) in the first and the second generations. The relatively heavy [light] top and the light [heavy] squarks (sleptons) characterizes the type I [type II] solution. In the next section, we consider some experimental informations in order to choose favorable solution and to get further constraints on the parameters.
the solutions type I and type II. The branching ratio of the direct decay mode \( \tilde{g} \rightarrow \tilde{q}\tilde{Z}_1 \) has been reported as \( 95\text{GeV} \) and \( 60\text{GeV} \) for type I and type II solution.

R selection cuts \([25]\). (i) We coalesce partons within \( \Delta R = \sqrt{\Delta \eta^2 + (\Delta \phi)^2} < 0.7 \) into single jets. We also require that all jets satisfy \( |\eta_{ij}| < 3.5 \), and each jet must have \( E_T > 15\text{GeV} \). The highest \( E_T \) cluster is also required to be central \( (|\eta| < 1) \). (ii) We require that there be no jet with \( E_T > 5\text{GeV} \) within a 30° cone back to back in azimuth with the leading jet. (iii) We require \( E_T > 40\text{GeV} \). We can find that the lower mass bounds for gluino are respectively about 95GeV and 60GeV for type I and type II solution.

5 Choice of parameter sets and further constrants

5.1 gluino searches at CDF

It is known that the bound \([13]\) from the negative gluino search at CDF is not realistic. To get realistic bound we must include the cascade decays in the analyses \([23]\). The gluino mass bound at CDF taken into account of the cascade decays \( \tilde{g} \rightarrow q\bar{q}\tilde{Z}_{2,3,4} \) and \( \tilde{g} \rightarrow u\bar{d}W_{1,2} \) as well as the direct decay \( \tilde{g} \rightarrow \tilde{q}\tilde{Z}_1 \) has been reported as \( 95\text{GeV} \) (90\% C.L.) for \( m_\tilde{g} \)

\[ m_\tilde{g} > 95\text{GeV} \quad (90\% \text{C.L.}) \tag{17} \]

for \( \mu = -250\text{GeV} \) and \( \tan \beta = 2 \), for example. We should note, however, that the bound \([17]\) has been obtained based on the assumption that \( m_{\tilde{t}_1} > m_\tilde{g} \) and the gluino can not decay into the stop. This is not the case we consider here. In fact, the gluino can decay into the stop pairs, \( \tilde{g} \rightarrow \tilde{t}_1\tilde{t}_1^* \tilde{Z}_1 \), which becomes another seed for the cascade decay because the stop and neutralino could be light enough.

The Feynman diagrams for the gluino decay are depicted in Fig.3 and the analytical formula for the decay width is presented in Appendix D. In Fig.4 we show the \( m_\tilde{g} \) dependence of the branching ratio of gluino, where we include the mode \( \tilde{g} \rightarrow \tilde{t}_1\tilde{t}_1^* \tilde{Z}_1 \) and sum up quark flavors \( q, q' = u, d, c, s \). We take \( \tan \beta = 2.0, \mu = -150\text{GeV}, m_{\tilde{t}_1} = 15\text{GeV}, \theta_t = 0.9, \) and \( M_2 = 22\text{GeV} \), and take \( m_\tilde{g} \) as a free parameter. The squark masses and the top mass are taken as \( (m_\tilde{q}, m_t) = (2m_\tilde{g}, 130\text{GeV}) \) for (a) and \( (3.8m_\tilde{g}, 95\text{GeV}) \) for (b), where \( m_\tilde{g} \equiv m_{\tilde{g}_{L,R}} = m_{\tilde{Z}_{L,R}} = m_{\tilde{\phi}_{L,R}} = m_{\tilde{p}_{L,R}} \). Note that Fig.4 (a) and (b) respectively correspond to the solutions type I and type II. The branching ratio of the direct decay mode \( \tilde{g} \rightarrow \tilde{q}\tilde{Z}_1 \), which is important in the \( \tilde{g} \) search in terms of large \( E_T \) signature, is reduced substantially even for the light gluino with mass \( m_\tilde{g} > 60\text{GeV} \). Therefore, we should reconsider the UA2 bound \( m_\tilde{g} > 79\text{GeV} \) \([27]\) obtained under the assumption \( \text{BR}(\tilde{g} \rightarrow q\bar{q}\tilde{Z}_1) \sim 100\% \) as well as the CDF bound \([17]\). For the value \( m_\tilde{g} = 75\text{GeV} \) determined by the GUT relation, \( \text{BR}(\tilde{g} \rightarrow q\bar{q}\tilde{Z}_1) \sim 20\% (0.6\%) \) for the type I [type II] solution, which should be compared with \( \text{BR}(\tilde{g} \rightarrow q\bar{q}\tilde{Z}_1) \sim 70\% \) obtained when there is no stop mode. We can find that if we take larger values of \( m_\tilde{q} \) and/or smaller values of \( m_t \), \( \text{BR}(\tilde{g} \rightarrow q\bar{q}\tilde{Z}_1) \) is reduced rapidly. In this case the Tevatron bound \([17]\) would be diminished significantly. This is because the width of the stop mode \( \Gamma(\tilde{g} \rightarrow \tilde{t}_1\tilde{t}_1^* \tilde{Z}_1) \) becomes larger for smaller values of \( m_t \) and all the other widths become smaller for larger values of \( m_\tilde{q} \), as we can see from Figs.3 and 4.

We try to simulate the Monte-Carlo calculation in order to get the gluino mass bounds from the CDF gluino searches. In Fig.5 we show the expected number of events in \( 4.3pb^{-1} \) integrated luminosity Tevatron running. In this calculation, we take following kinematical selection cuts \([25]\). (i) We coalesce partons within \( \Delta R = \sqrt{\Delta \eta^2 + (\Delta \phi)^2} < 0.7 \) into single jets. We also require that all jets satisfy \( |\eta_{ij}| < 3.5 \), and each jet must have \( E_T > 15\text{GeV} \). The highest \( E_T \) cluster is also required to be central \( (|\eta| < 1) \). (ii) We require that there be no jet with \( E_T > 5\text{GeV} \) within a 30° cone back to back in azimuth with the leading jet. (iii) We require \( E_T > 40\text{GeV} \). We can find that the lower mass bounds for gluino are respectively about 95GeV and 60GeV for type I and type II solution.
5.2 second neutralino searches at LEP

In Sec.3 we have noted that the second neutralino search through measuring the visible width of the $Z$-boson could set severe constraints on the MSSM basic parameters ($M_2$, tan $\beta$, $\mu$). The experimental upper bound reported recently is $[13, 21, 22]$ and this bound can be converted to the excluded region in the ($M_2$, $\mu$) plane fixed tan $\beta = 2$ as shown in Fig.1. We should note that the leptons or quarks (jets) signatures have been used in serching for the second neutralino through the direct decay $\tilde{Z}_2 \rightarrow f\bar{f}\tilde{Z}_1$. The bound $[13]$ has been obtained under the assumption of BR($\tilde{Z}_2 \rightarrow f\bar{f}\tilde{Z}_1$) = 100%. Here again we should consider the appearance of the light stops in the final state of the second neutralino decay. In fact, if the mass of second neutralino $m_{\tilde{Z}_2}$ is larger than $2m_{\tilde{t}_1} + m_{\tilde{Z}_1} \approx 43$GeV, it can decay into the stops.

The Feynman diagrams for the second neutralino decay are depicted in Fig.6 and the analytical formula for the decay width is presented in Appendix D, where we neglect a contribution Fig.6(4) because $\tilde{Z}_1 \sim \tilde{\gamma}$ and the stop coupling to the $Z$-boson (11) is expected to be rather small in our case. In Fig.7 we show the $m_t$ dependence of the branching ratio of second neutralino. We take tan $\beta = 2.0$, $\mu = -150$GeV, $m_{\tilde{t}_1} = 15$GeV, $\theta_t = 0.9$, $m_h = 60$GeV, $\alpha = -0.6$ and $M_2 = 22$GeV, where $\alpha$ denotes the Higgs mixing angle $[28]$. The squark masses and the trilinear coupling $A_t$ are taken as $(m_{\tilde{q}}, A_t) = (2m_{\tilde{q}}, 300$GeV) in (a) and $(3.8m_{\tilde{q}}, 600$GeV) in (b), where we take $m_{\tilde{t}} = m_{\tilde{t}}$ for simplicity. It is worth mentioning that $\Gamma(\tilde{Z}_2 \rightarrow \tilde{t}_1\bar{t}_1\tilde{Z}_1)$ dominates over $\Gamma(\tilde{Z}_2 \rightarrow f\bar{f}\tilde{Z}_1)$ for large $m_t$ because the virtual Higgs exchange diagram Fig.6(3) gives large contribution to the stop width. The predicted values in the parameter sets for type I and type II can be read from Fig.7 (a) with $m_t \approx 130$GeV and (b) with $m_t \approx 95$GeV, respectively. The branching ratio of the direct decay mode $\tilde{Z}_2 \rightarrow f\bar{f}\tilde{Z}_1$, which is important in the $\tilde{Z}_2$ search in terms of the leptons or quarks (jets) signature, is reduced substantially as BR($\tilde{Z}_2 \rightarrow f\bar{f}\tilde{Z}_1$) $\approx 25\%$ [2.5\%] for type I [type II] solution. Therefore, we should reconsider the LEP bound $[13, 21, 22]$ obtained under the assumption BR($\tilde{Z}_2 \rightarrow f\bar{f}\tilde{Z}_1$) $\sim$ 100\%. Simply we can obtain the effective bounds on the observed branching ratio as

$$\text{BR}(Z \rightarrow \text{vis.})_{\text{obs}}^{\text{upper}} = \frac{\text{BR}(Z \rightarrow \text{vis.})_{\text{exp}}^{\text{upper}}}{\text{BR}(\tilde{Z}_2 \rightarrow f\bar{f}\tilde{Z}_1)}, \tag{18}$$

and we get its numerical value, $2 \times 10^{-5}$ [$2 \times 10^{-4}$] for the type I [type II] solution.

After all, we obtain the correct excluded region in the ($M_2$, $\mu$) plane fixed tan $\beta = 2$ for type I and type II solutions depicted in Fig.8 (a) and (b), respectively. We realize that the neutralino with mass $13 \sim 14$GeV is still excluded for the type I solution. However, a window, $-170$GeV $\lesssim \mu \lesssim -70$GeV, which allows the existence of the light neutralino, opens up for the type II solution.

5.3 proton decay constraint

Besides the collider experiments we should consider the constraints on the model parameters from the non-accelerator obsevations and the cosmological considerations. In the

*Rigorously speaking, the contour of $m_{\tilde{g}}$ depends slightly on $\mu$. In Fig.8 we neglect such a sensitivity since our conclusion is not affected by it.*
MSGUT there are the dangerous dimension five operators which contribute to the fast nucleon decay \[29\]. The nucleon decay life time \(\tau_N\) in the framework of the SU(5) SUSY GUT model has been investigated in detail and the present upper bounds on \(\tau_N\) have been converted to the constraints on the masses of the chargino and the sfermions. The most stringent constraint comes from the non-observation for the decay \(N \rightarrow K \nu_{\mu}\). Hisano et al. have shown that the squarks and the sleptons with masses 200\(\text{GeV}\) have been excluded for \(m_{\tilde{W}_1} \lesssim 50\text{GeV}\) even if we take the most conservative parameters \[29\]. This result seems to indicate that the type II solution is more favorable.

5.4 dark matter candidate

It is expected that the LSP in the SUSY model can be a good candidate for the dark matter of the universe \[30\]. In our case the LSP is the lightest neutralino with mass \(13\text{GeV} \sim 14\text{GeV}\) and it is almost the photino. The relic abundance \(\Omega h^2\) of the light photino has been calculated and the upper bound on the sfermion masses can be obtained if we consider the universe will not be overclosed \(\Omega h^2 < 1\). Naively, the photino with mass \(\lesssim 20\text{GeV}\) can be a dark matter candidate only when \(m_{\tilde{q}}, m_{\tilde{\ell}} \lesssim 160\text{GeV}\). At first insight, the type I solution is more favorable and this fact seems to contradict with the result from the nucleon decay constraints discussed in the last subsection.

Recently, however, Fukugita et al. \[18\] have pointed out that the co-annihilation process \(\tilde{\gamma} \tilde{\gamma} \rightarrow gg\) can be occurred through the chain process \(\tilde{\gamma} + c \rightarrow \tilde{t} \) and \(\tilde{t} \tilde{t} \rightarrow gg\) in the case of the light stop. The cross section of such a process is enhanced when \(m_{\tilde{\gamma}}\) and \(m_{\tilde{\gamma}}\) are close to degenerate. They have shown that the cosmological limit \(\Omega h^2 > 1\) can be avoided even for \(m_{\tilde{q}}, m_{\tilde{\ell}} \simeq 300\text{GeV}\) in so far as \(m_{\tilde{t}} - m_{\tilde{\gamma}} \lesssim 5\text{GeV}\). Note that the type II solution gives \(m_{\tilde{q}}, m_{\tilde{\gamma}} \simeq 300\text{GeV}\) and the type II solution can survive. It has been pointed out that the proton decay favors a large value of \(\xi_0 \equiv m_{\infty}/M_{\infty}\) but the cosmology of the neutralino dark matter disfavors large value of \(\xi_0\) \[30\]. So the type II solution satisfies both constraints.

5.5 neutral Higgs searches at LEP

A result from the negative searches for the MSSM Higgs at LEP could set another constraint on the SUSY parameter space. The present limit on the MSSM (SM like) Higgs mass is

\[
    m_h \gtrsim 61.5 \text{GeV} \quad (95\% \text{C.L.})
\]

for \(\sin(\beta - \alpha) \simeq 1\) \[20\]. Note, however, that this bound has been obtained based upon the assumption that the Higgs boson does not decay into the stop. Here we must consider the fact that the neutral Higgs could have a dominant decay mode \(h \rightarrow \tilde{t}_1 \tilde{t}_1^*\) with almost 100\% branching ratio if the stop is light enough. In this case energies of visible jets from the Higgs production would become softer and it can be smaller than the detection lower cuts. Therefore, if we incorporate such a decay mode in data analyses, the present lower bounds are expected to be weakened.

We try to simulate the Monte-Carlo calculation in order to get the Higgs mass bounds from the LEP Higgs searches. In Fig.9 we show the expected number of events in 93.5 \(\text{pb}^{-1}\)
integrated luminosity LEP running. Here we assume the SM Higgs will decay into $b\bar{b}$ with 100% branching ratio for simplicity. It is expected that we can distinguish the Higgs signal from backgrounds by measuring the hadronic ($q\bar{q}$) invariant mass $M_{q\bar{q}}$ distribution of the candidate events, which should have a distinct peak for the SM Higgs events as we can see from Fig.9. However, if the Higgs decays into the stops $h \rightarrow \tilde{t}_1 \tilde{t}_1 \rightarrow c\bar{c}Z_1 \bar{Z}_1$, the peak in $M_{q\bar{q}}$ distribution will appear at $M_{q\bar{q}} \ll m_h$ because the large energies and momenta will be carried off by the neutralino $\tilde{Z}_1$. As a consequence, the signatures will be hidid the large backgrounds from the $\gamma\gamma$ processes and $Z \rightarrow q\bar{q}\ell^+\ell^-$ events for the neutrino channel $Z \rightarrow hZ^* \rightarrow h(\nu\bar{\nu})$ and the lepton channel $Z \rightarrow hZ^* \rightarrow h(\ell^+\ell^-)$, respectively.

Usually we can suppress the $\gamma\gamma$ background by taking a lower cut for the jet transverse momentum for the neutrino channel $Z \rightarrow hZ^* \rightarrow h(\nu\bar{\nu})$. Such a selection cut will not be useful when the Higgs decays into the stops since the jet transverse momentum will become very soft. So we suppose that the expected number of events of the neutrino channel reduced considerably. On the other hand, the selection cuts on the visible jet energies are not so essential in the lepton channel $Z \rightarrow hZ^* \rightarrow h(\ell^+\ell^-)$. The energy and momentum distribution of the scattered leptons from the virtual $Z$-boson are irrespective to the decay modes of the Higgs. In fact, the Higgs mass can be determined by the recoil mass

$$M^2_{\text{recoil}} = m_Z^2 - 2m_Z(E_{\ell^+} + E_{\ell^-}) + M^2_{\ell^+\ell^-},$$

instead of the direct measurement of $M_{q\bar{q}}$ in principle. We suppose, therefore, that the reduction rates of the events in the lepton channel are not so large and that the lower mass bounds on the Higgs from the lepton channel will be applicable even when the Higgs decays into the stop pairs. The Higgs mass bounds in terms of the lepton channel reported by the four groups at LEP are slightly different each other [21, 22, 31]. Here we assume conservative bound $m_h \gtrsim 55\text{GeV}$ (21) and adopt it as a further constraint on our parameter space $(\mu, \tan \beta, m_t)$.

### 5.6 typical parameter sets

From the discussions given above we find that the type II solution is more favorable. In this subsection we present the typical parameter values of the type II solution. Those for the type I solution can be found in Ref.[11]. Contours of $m_{\tilde{t}_1} = 15\text{GeV}$ in the $(\mu, \tan \beta)$ plane for $M_2 = 22\text{GeV}$ are shown in Fig.10. Each line corresponds to a contour of $m_{\tilde{t}_1} = 15\text{GeV}$ for a fixed $m_t$ value. We also plot the $m_h = 55\text{GeV}$ contour and $m_{\tilde{W}_1} = 45\text{GeV}$ contour.

In the region denoted by "excluded theoretically" the false vacuum is realized [32], i.e.,

$$A_t^2 > 3(m^2_{\tilde{t}_L} + m^2_{\tilde{t}_R} + m^2_\mu),$$

where $m^2_\mu = \bar{m}_{\mu\mu}^2 + \mu^2 - 3\bar{I}$. We see that there is a rather narrow (but finite) range allowing the existence of the light stop, if the top is slightly light too, $m_t \approx 95\text{GeV}$. Furthermore, we find that the light stop solution gives inevitably a light Higgs boson, $m_h \lesssim 60\text{GeV}$. While we have included the radiative correction in the calculation of the Higgs mass [33], deviations $\delta m_h$ from the tree level results are not so large, $|\delta m_h| \lesssim 2\text{GeV}$. The neutral Higgs is standard Higgs like, i.e., $\sin(\beta - \alpha) \simeq 1$.
Adopting the bound $m_h > 55 \text{GeV}$, we can choose three typical parameter sets (A), (B) and (C), denoted in Fig.10. Input and output values of the parameters of the sets (A), (B) and (C) are presented in Table I. Interested reader can compare those parameter sets of the type II solution with those of the type I solution presented in the previous papers [11]. The set (A) [(C)] has the largest [smallest] values of the scalar fermion masses, the neutral Higgs $h$ mass and the top mass, and has the smallest [largest] value of the lighter chargino mass. We find that masses and mixing parameters are severely constrained, for example, $m_{\tilde{g}} \simeq 75 \text{GeV}$, $m_{\tilde{W}_1} \simeq 50 \text{GeV}$, $m_{\tilde{\tau}} \simeq m_{\tilde{\nu}} \simeq 300 \text{GeV}$, $\theta_t \simeq 0.9$, $m_h \lesssim 60 \text{GeV}$ and $m_t \simeq 90 \text{GeV}$. These parameter sets are characterized by the rather light top quark and the heavy squarks and sleptons. To obtain those values we take $M_2 = 22 \text{GeV}$ and $\sin^2 \theta_W = 0.232$ ($\alpha_s = 0.113$). It should be noted that masses of the gluino, the Higgs and the top depend sensitively on those input parameters. For example, the allowed ranges are changed to $m_{\tilde{g}} \simeq 85 \text{GeV}$, $m_h \lesssim 65 \text{GeV}$ and $m_t \simeq 98 \text{GeV}$ for $M_2 = 24 \text{GeV}$ and $\sin^2 \theta_W = 0.230$ ($\alpha_s = 0.120$). Anyway, it seems that the expected top mass $m_t = 90 \sim 98 \text{GeV}$ is too small compared to the present experimental bounds from Tevatron. We will discuss the top mass bound taking into account of the stop modes of the top decay in the next section.

6 Phenomenological implications

Now we are in position to discuss some consequence of the light stop scenario in the MSGUT and give strategies to confirm or reject such a possibility in present and future experiments. Some numerical results are calculated with the typical parameter sets (A), (B) and (C) in Table I.

6.1 top decay

The existence of the light stop with mass $15 \sim 16 \text{GeV}$ will alter completely decay patterns of some ordinary and SUSY particles (sparticles). First we discuss the top decay [5, 35]. In our scenario, the top can decay into final states including the stop; $t \rightarrow \tilde{t}_1 \tilde{Z}_1$, $\tilde{t}_1 \tilde{Z}_2$ and $\tilde{t}_1 \tilde{g}$. Branching ratios of the top for the typical parameter sets are presented in Table II. We find that the stop modes $t \rightarrow \tilde{t}_1 \tilde{Z}_k$, $\tilde{t}_1 \tilde{g}$ have about 80% branching ratio and dominate over the standard mode $t \rightarrow bW^+ \simeq 20\%$. Strategies for the top search at Tevatron would be forced to change because the leptonic branching ratios of the top would be reduced by the dominance of the stop modes. The expected event rates for the ordinary $bW$ signatures will be reduced, i.e., the total cross section $\sigma(p\bar{p} \rightarrow t\bar{t}X \rightarrow b\bar{b}W^+W^-X)$ for $m_t \simeq 90 \text{GeV}$ is almost equivalent to that for $m_t \simeq 165 \text{GeV}$ in the standard model.

6.2 Higgs decay

Decay patterns of the Higgs particles will be changed too. The lighter CP-even neutral Higgs decays dominantly into the stop pair, $h \rightarrow \tilde{t}_1 \tilde{t}_1^*$, owing to the large Yukawa coupling of the top. In rough estimation, we obtain

$$\text{BR}(h \rightarrow \tilde{t}_1 \tilde{t}_1^*) \simeq \frac{1}{1 + \frac{3m_t^2m_h^2}{2m_t^4}} \simeq 1.$$ (23)
This fact would change the experimental methods of the Higgs searches at the present and future collider experiments. In particular, it could be expected that the number of events of the neutrino channel $Z \rightarrow hZ^* \rightarrow h(\nu \bar{\nu})$ will be smaller than that of the lepton channel $Z \rightarrow hZ^* \rightarrow h(\ell^+ \ell^-)$ in the Higgs $h$ searches at LEP. This is because the selection cuts on the visible jet energies in the neutrino channel are large enough to throw away the soft jets signatures from the Higgs decay $h \rightarrow t_1\bar{t}_1$. Consequently, when the Higgs decays into the stop pair, the searches for the lepton channel $Z \rightarrow hZ^* \rightarrow h(\ell^+ \ell^-)$ would be more important than the neutrino channel. More detail analyses of the charged [30] and neutral Higgs bosons [37] are presented separately.

6.3 sparticle decay

Now we discuss briefly the light stop impact on the sparticle decays. The lightest charged sparticle apart from the stop is the lighter chargino $\tilde{\chi}^+_1$. The two body stop mode $\tilde{W}_1 \rightarrow b\tilde{t}_1$ would dominate over the conventional three body mode $\tilde{W}_1 \rightarrow f\bar{f}Z_1$. As a consequence, it would be difficult to use the leptonic signature in the chargino search at $e^+e^-$ and hadron colliders. Since the chargino $\tilde{W}_1$, the second neutralino $\tilde{Z}_2$, the neutral Higgs $h$ and the gluino $\tilde{g}$, whose dominant decay modes are respectively $\tilde{W}_1 \rightarrow b\tilde{t}_1$, $\tilde{Z}_2 \rightarrow \tilde{t}_1\bar{t}_1 Z_1$, $h \rightarrow t_1\bar{t}_1$ and $\tilde{g} \rightarrow \tilde{t}_1\bar{t}_1 Z_1$, are copiously produced in the other sparticle decays, many stops would be expected in the final states of the sparticle production. For example, $\ell_L \rightarrow \nu\tilde{W}_1 \rightarrow \nu(b\tilde{t}_1)$, $\ell \rightarrow t\tilde{Z}_2 \rightarrow \ell(\tilde{t}_1\bar{t}_1 Z_1)$, $\tilde{q}_{L,R} \rightarrow q\tilde{g} \rightarrow q(\tilde{t}_1\bar{t}_1 Z_1)$, $\tilde{q}_L \rightarrow q'\tilde{W}_1 \rightarrow q'(b\tilde{t}_1)$ and $\tilde{Z}_{i(i=3,4)} \rightarrow \tilde{Z}_1 h \rightarrow \tilde{Z}_1(\tilde{t}_1\bar{t}_1)$.

6.4 stop and neutralino searches at LEP

We can see that the stop and its relatively light accompaniments, the gluino $\tilde{g}$, light neutralinos $\tilde{Z}_{1,2}$, and neutral Higgs $h$, should be visible at LEP, SLC, HERA and Tevatron. Especially, LEP could search the allowed region presented in Figs.8(b) and 10 in terms of the width of $Z$-boson and the direct stop search. We can see from Table I, the stop mixing angle $\theta_t$ is severely limited as $\theta_t \simeq 0.9$ in the allowed range. As the stop search in terms of $\Delta \Gamma_Z$ would be difficult in this case, the direct search for $e^+e^- \rightarrow \tilde{t}_1\bar{t}_1$ will be important [20, 21]. Second, the whole allowed region in Figs.8(b) and 10 can be explored by the precise measurement of BR($Z \rightarrow \text{vis}$). In fact, the smallest value of the neutralino contribution $\sum_{i \neq j} \Gamma(Z \rightarrow \tilde{Z}_i\tilde{Z}_j)/\Gamma_Z^{\text{obs}}$ to BR($Z \rightarrow \text{vis}$) is about $4 \times 10^{-7}$. Here we included the reduction of the observed branching ratio originated from the stop mode in the second neutralino decay. Values of branching ratios of the second neutralino for the typical parameter sets (A), (B) and (C) are tabulated in Table III. Clearly, the lighter chargino, $m_{\tilde{\chi}^+_1} \lesssim 50$GeV, would be visible at LEP-II. Moreover, LEP-II may be able to produce both the top and the stop through a process $e^+e^- \rightarrow \tilde{t}_1\tilde{Z}_1$.

6.5 gluino and stop searches at Tevatron

As mentioned before, Tevatron will play a crucial role in confirming or rejecting the light stop scenario in the MSGUT with the GUT relation. In this case the existence of relatively light gluino, $m_{\tilde{g}} = 75 \sim 85$GeV, with substantially large decay fraction $\tilde{g} \rightarrow \tilde{t}_1\bar{t}_1 Z_1$ is one
of definite prediction. Values of branching ratios of the gluino for the typical parameter sets (A), (B) and (C) are tabulated in Table IV. The branching ratio of the direct decay mode \( BR(\tilde{g} \to q\bar{q}Z_1) = 0.4 \sim 0.5\% \) is expected in the allowed range. In the gluino search at Tevatron the mixed signature, \( p\bar{p} \to \tilde{g}gX \to (\tilde{t}_1\tilde{t}_1^*Z_1)(q\bar{q}Z_1)X \), and in turn the two-jets events would be dominant signature. Of course, the main fraction of the events will be almost invisible because the resulting charm-jets from \( p\bar{p} \to \tilde{g}gX \to (\tilde{t}_1\tilde{t}_1^*Z_1)(\tilde{t}_1\tilde{t}_1^*Z_1)X \to c\bar{c}Z_1\bar{Z}_1Z_1\bar{Z}_1\bar{Z}_1 \) will be too soft to detect.

As has been pointed out by Baer et al. \([3, 10]\) that the softness of the charm jets makes direct searches for the light stop difficult at Tevatron. However, there is a possibility for the detection of the stop \( \gamma\gamma \) events originated from the decay of the \( s \)-wave stoponium \( \sigma_{\tilde{t}_1} = (\tilde{t}_1\tilde{t}_1^*) \) [34].

### 6.6 Stop searches at HERA

The \( ep \) collider HERA could search the light stop through its pair production process \( ep \to e\tilde{t}_1\tilde{t}_1^*X \) via boson-gluon fusion [38]. The total cross section of the process is larger than about 10\( pb \) for \( m_{\tilde{t}_1} \cepto 20\text{GeV} \), which is independent on the mixing angle \( \theta_t \). That is, \( \sim 10\text{pb} \) is expected for all parameters with \( m_{\tilde{t}_1} \cepto 20\text{GeV} \) in the allowed range in Fig.10. Detail analyses with Monte Carlo studies including possible dominant background process \( ep \to e\nu\pi X \) can be found in Ref.[39].

### 6.7 Stop searches at SLC

Polarized initial electron beams at SLC and at any linear \( e^+e^- \) colliders will be efficient to reveal the nature of left-right mixing in the stop sector, in other words, to measure the mixing angle of stop \( \theta_t \). In Fig.11 we show the \( \sqrt{s} \) dependence of the left-right asymmetry;

\[
A_{LR} \equiv \frac{\sigma(e_L) - \sigma(e_R)}{\sigma(e_L) + \sigma(e_R)} \tag{24}
\]

where \( \sigma(e_{L,R}) \equiv \sigma(e^+e_-_{L,R} \to \tilde{t}_1\tilde{t}_1^*) \), which are obtained by

\[
\sigma(e^+e_-_{L,R} \to \tilde{t}_1\tilde{t}_1^*) = \frac{\pi\alpha^2}{s} \beta_{\tilde{t}_1}^3 \left[ \frac{4}{9} + \frac{2}{3} C_{\tilde{t}_1} (v_e \pm a_e) \text{Re} \left( \frac{s}{D_Z} \right) + \frac{1}{4} C^2_{\tilde{t}_1} (v_e \pm a_e)^2 \left( \frac{s}{D_Z} \right)^2 \right] \left( 1 + \delta_{QCD} \right) \tag{25}
\]

where \( \beta_{\tilde{t}_1} = \sqrt{1 - 4m_{\tilde{t}_1}^2/s}, \; D_Z = s - m_Z^2 + im_Z \Gamma_Z, \; v_e \equiv (-\frac{1}{2} + 2 \sin^2 \theta_W)/(\sin^2 \theta_W \cos^2 \theta_W) \) and \( a_e \equiv -1/(2 \sin^2 \theta_W \cos^2 \theta_W) \). In the asymmetric combination in Eq.(24) the photon contribution is cancelled out and \( A_{LR} \) is proportional to \( C_{\tilde{t}_1} \) [11]. This is the reason for sensitive dependence on \( \theta_t \) of \( A_{LR} \) in Fig.11. Another important property is that \( A_{LR} \) is independent on the mass of stop \( m_{\tilde{t}_1} \) as well as on the QCD correction \( \delta_{QCD} \) since \( \beta_{\tilde{t}_1} \) and \( 1 + \delta_{QCD} \) disappeared in the fractional combination of \( \sigma \) in \( A_{LR} \). Therefore, this method for measuring \( \theta_t \) will be applicable for the stop with any mass satisfying \( m_{\tilde{t}_1} < \sqrt{s}/2 \).
7 Conclusion

We have investigated the possibility for the existence of the light stop \( m_{\tilde{t}_1} = 15 \sim 16 \text{GeV} \) and the neutralino \( m_{\tilde{\chi}^0_1} = 13 \sim 14 \text{GeV} \) in the MSGUT scenario taking into account of the present experimental bounds on the SUSY parameter space. We have found that the two types of solution for the RGEs in the MSGUT, type I and type II, allowing the existence of the light stop. It has been found that the type II solution is more favorable because these parameter sets seem to satisfy all constraints settled by the recent collider experiments as well as by the non-accelerator observations and the cosmological considerations. However, the allowed parameter region is rather restricted and in turn masses and mixing parameters of the other SUSY partners as well as masses of the Higgs and the top are severely constrained. For example, \( 75 \text{GeV} \lesssim m_{\tilde{g}} \lesssim 85 \text{GeV}, \ m_{\tilde{W}_1} \lesssim 50 \text{GeV}, \ m_{\tilde{\tau}} \sim m_{\tilde{q}} \sim 300 \text{GeV}, \ \theta_t \sim 0.9, \ m_h \lesssim 65 \text{GeV} \) and \( 90 \text{GeV} \lesssim m_t \lesssim 100 \text{GeV} \).

Our results can be stated in other words as follows. In the MSSM the very light stop and the light neutralino could be allowed only when the top is relatively light \( m_t \lesssim 100 \text{GeV} \), the sfermions are heavy \( m_{\tilde{f}} \gtrsim 300 \text{GeV} \) and the stop mixing angle is in the range \( 0.85 \lesssim \theta_t \lesssim 1.15 \). In this case, the existence of the light stop can compromise with the negative results of the searches for the top quark, the Higgs boson and the sparticles at LEP and Tevatron as well as with the non-observation of the nucleon decay and the cosmological dark matter constraints. Then we have found that such a peculiar solution can be realized in the framework of the MSGUT model. It is interesting, moreover, that the mass ratio of the top to the bottom quark is almost the same as the ratio of the sbottom to the stop, \( m_t/m_b \simeq m_{\tilde{b}}/m_{\tilde{t}} \), in our solution. This may have its origin in a specific model at the Plank scale such as superstrings.

It should be emphasized that the signatures of the top and the neutral Higgs will be changed significantly. The branching ratio for the standard decay mode of the top \( t \to bW \) will be about 20% and then the expected event rates for the ordinary \( bW \) signatures will be reduced, i.e., the total cross section \( \sigma(pp \to t\bar{t}X \to b\bar{b}W^+W^-X) \) for \( m_t \simeq 90 \text{GeV} \) is almost equivalent to that for \( m_t \simeq 165 \text{GeV} \) in the standard model. For Higgs searches at LEP, it is expected that the Higgs with mass about 60GeV will be discovered more easily in searching for the lepton channel \( Z \to hZ^* \to h(\ell^+\ell^-) \) than in searching for the neutrino channel \( Z \to hZ^* \to h(\nu\bar{\nu}) \). Moreover, the light stop and its relatively light accompaniments, the gluino \( \tilde{g} \), the light neutralinos \( \tilde{\chi}^0_{1,2} \), and the lighter chargino \( \tilde{W}_1 \), should be visible near future at LEP, HERA and Tevatron.

We have exemplified that if we discover the light stop we will be able to constrain severely all the SUSY parameters at the unification scale. We can conclude that, therefore, the discovery of the stop will bring us a great physical impact. Not only will it be the first signature of the top flavor and supersymmetry but also it could shed a light on the physics at the unification scale.

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Appendix A

In this Appendix we list all renormalization group equations which are presented in Ref. [15]. In the following, dots (•) denote the derivative \( d/dt \) where \( t \equiv (4\pi)^{-1}\ln(M_X^2/q^2) \). The coefficients of the one loop \( \beta \) function in the MSGUT are given by \( b_1 = 33/5, b_2 = 1 \) and \( b_3 = -3 \). The canonical weak hypercharge coupling is defined by \( \alpha' = (3/5)\alpha_1 \).

Gauge couplings

\[
\dot{\alpha}_i = -b_i \alpha_i^2.
\]

Yukawa Couplings

\[
\dot{\alpha}_i = \left( \frac{16}{3} \alpha_3 + 3\alpha_2 + \frac{13}{9}\alpha' - 6\alpha_i \right)\alpha_i.
\]

Higgsino mass \( \mu \)

\[
\dot{\mu} = \left( \frac{3}{2}\alpha_2 + \frac{1}{2}\alpha' - \frac{3}{2}\alpha_t \right)\mu.
\]

Gaugino masses

\[
\dot{M}_i = -b_i \alpha_i M_i.
\]

Scalar masses

\[
\dot{\tilde{m}}_{Q_{1,2}}^2 = \frac{16}{3} \alpha_3 M_3^2 + 3\alpha_2 M_2^2 + \frac{1}{9} \alpha' M_1^2, \\
\dot{\tilde{m}}_{Q_3}^2 = \frac{16}{3} \alpha_3 M_3^2 + 3\alpha_2 M_2^2 + \frac{1}{9} \alpha' M_1^2 - \alpha_t (\tilde{m}_{Q_3}^2 + \tilde{m}_{U_3}^2 + \tilde{m}_{H_2}^2 + A_t^2), \\
\dot{\tilde{m}}_{U_{1,2}}^2 = \frac{16}{3} \alpha_3 M_3^2 + \frac{16}{9} \alpha' M_1^2, \\
\dot{\tilde{m}}_{D_{1,2,3}}^2 = \frac{16}{3} \alpha_3 M_3^2 + \frac{4}{9} \alpha' M_1^2, \\
\dot{\tilde{m}}_{U_3}^2 = \frac{16}{3} \alpha_3 M_3^2 + \frac{16}{9} \alpha' M_1^2 - 2\alpha_t (\tilde{m}_{Q_3}^2 + \tilde{m}_{U_3}^2 + \tilde{m}_{H_2}^2 + A_t^2), \\
\dot{\tilde{m}}_L^2 = 3\alpha_2 M_2^2 + \alpha' M_1^2, \\
\dot{\tilde{m}}_E^2 = 4\alpha' M_1^2, \\
\dot{\tilde{m}}_{H_1}^2 = 3\alpha_2 M_2^2 + \alpha' M_1^2, \\
\dot{\tilde{m}}_{H_2}^2 = 3\alpha_2 M_2^2 + \alpha' M_1^2 - 3\alpha_t (\tilde{m}_{Q_3}^2 + \tilde{m}_{U_3}^2 + \tilde{m}_{H_2}^2 + A_t^2).
\]

Trilinear parameters \( A \)

\[
\dot{A}_b = \frac{16}{3} \alpha_3 M_3 + 3\alpha_2 M_2 + \frac{7}{9} \alpha' M_1 - \alpha_t A_t, \\
\dot{A}_t = \frac{16}{3} \alpha_3 M_3 + 3\alpha_2 M_2 + \frac{13}{9} \alpha' M_1 - 6\alpha_t A_t.
\]
Higgs mixing parameter $B$

$$
\dot{B} = 3\alpha_2 M_2 + \alpha' M_1 - 3\alpha_i A_i.
$$

## Appendix B

Here we present the analytical solutions of all RGEs for the general case $f_i \neq 1$ in the Appendix A. To get following results we take the boundary conditions presented in the Sec. 4.

### Gauge couplings

$$
\alpha_i(t) = \frac{\alpha_\infty}{1 + b_i \alpha_\infty t}.
$$

### Gaugino masses

$$
M_i(t) = \alpha_i(t) f_i \frac{M_\infty}{\alpha_\infty} = \frac{M_\infty f_i}{1 + b_i \alpha_\infty t}.
$$

At weak scale, these mass parameters are related as

$$
M_3 = m_{\tilde{g}},
$$

$$
M_2 = \frac{\alpha}{\alpha_3 \sin^2 \theta_W} \frac{f_2}{f_3} M_3,
$$

$$
M_1 = \frac{5}{3} \tan^2 \theta_W \frac{f_1}{f_2} M_2.
$$

### Scalar masses (first and second generation sfermion)

$$
\tilde{m}^2(t) = m_\infty^2 + M_\infty \sum_{i=1}^{3} c_i \tilde{\xi}_i
$$

with

$$
\tilde{\xi}_i = \frac{1}{2b_i} \frac{f_i}{f^2} \left[ 1 - \frac{1}{(1 + b_i \alpha_\infty t)^2} \right],
$$

where $(c_1, c_2, c_3) = (1/15, 3, 16/3)$ for $\tilde{q}_L$, $(16/15, 0, 16/3)$ for $\tilde{u}_R$, $(4/15, 0, 16/3)$ for $\tilde{d}_R$, $(3/5, 3, 0)$ for $\tilde{\ell}_L$ and $(12/5, 0, 0)$ for $\tilde{e}_R$.

### Yukawa couplings

$$
\alpha_t(t) = \frac{\alpha_\infty \prod_{i=1}^{3} (1 + b_i \alpha_\infty t) c_i^{o_t}/b_i}{1 + 6\alpha_\infty F(t)}
$$

$$
= \frac{\alpha_\infty \left( \frac{\alpha_3(t)}{\alpha_\infty} \right)^{\frac{16}{13}} \left( \frac{\alpha_2(t)}{\alpha_\infty} \right)^{-3} \left( \frac{\alpha_1(t)}{\alpha_\infty} \right)^{-\frac{13}{13}}}{1 + 6\alpha_\infty F(t)},
$$

where

$$
F(t) = \int_0^t dt' \prod_{i=1}^{3} (1 + b_i \alpha_\infty t') c_i^{o_t}/b_i = \int_0^t dt' \left( \frac{\alpha_3(t')}{\alpha_\infty} \right)^{\frac{16}{13}} \left( \frac{\alpha_2(t')}{\alpha_\infty} \right)^{-3} \left( \frac{\alpha_1(t')}{\alpha_\infty} \right)^{-\frac{13}{13}}
$$

with $c_3^{o_t} = c_3^{o_b} = 16/6$, $c_2^{o_t} = c_2^{o_b} = 3$, $c_1^{o_t} = 13/15$ and $c_1^{o_b} = 7/15$. 

Scalar masses (Third Generation and Higgs)

\[ \mu(t) = \frac{\mu_\infty \left( \frac{\alpha_2(t)}{\alpha_\infty} \right)^{-\frac{3}{2}} \left( \frac{\alpha_1(t)}{\alpha_\infty} \right)^{-\frac{1}{2}}}{\left[ 1 + 6 \alpha_{t,\infty} F(t) \right]^{1/4}} \]

Trilinear parameters \( A \)

\[
A_e(t) = A_\infty + M_\infty t \left[ 3 f_2 \alpha_2(t) + \frac{9}{5} f_1 \alpha_1(t) \right],
\]

\[
A_d(t) = A_\infty + M_\infty t \left[ \frac{16}{3} f_3 \alpha_3(t) + 3 f_2 \alpha_2(t) + \frac{9}{5} f_1 \alpha_1(t) \right],
\]

\[
A_u(t) = A_\infty + M_\infty t \left[ \frac{16}{3} f_3 \alpha_3(t) + 3 f_2 \alpha_2(t) + \frac{13}{15} f_1 \alpha_1(t) \right],
\]

\[
A_f(t) = A_e(t),
\]

\[
A_t(t) = \frac{A_\infty}{ \left[ 1 + 6 \alpha_{t,\infty} F(t) \right] + M_\infty} \left[ t \sum_{i=1}^{3} f_i c_i^{\alpha_i} \alpha_i(t) - \frac{6 \alpha_{t,\infty}}{ \left[ 1 + 6 \alpha_{t,\infty} F(t) \right] } G(f_i; t) \right],
\]

\[
A_b(t) = A_\infty + M_\infty t \sum_{i=1}^{3} f_i c_i^{\alpha_i} \alpha_i(t) - \int_{0}^{t} dt' \alpha_i(t') A_i(t')
\]

Here we use the relation

\[
\int_{0}^{t} dt' \alpha_i(t') A_i(t') = A_\infty \frac{\alpha_{t,\infty} F(t)}{1 + 6 \alpha_{t,\infty} F(t)} + M_\infty \alpha_{t,\infty} \tilde{K}(t)
\]

with

\[
\tilde{K}(t) = \int_{0}^{t} dt' \frac{G'(f_i; t')}{1 + 6 \alpha_{t,\infty} F(t')} - \int_{0}^{t} dt' \frac{6 \alpha_{t,\infty} F'(t')}{ \left[ 1 + 6 \alpha_{t,\infty} F(t') \right]^2 } G(f_i; t')
\]

\[
= \frac{G(f_i; t)}{1 + 6 \alpha_{t,\infty} F(t)}.
\]

Higgs mixing parameter \( B \)

\[
B(t) = B_\infty + M_\infty t \left[ 3 f_2 \alpha_2(t) + \frac{3}{5} f_1 \alpha_1(t) \right] - 3 \int_{0}^{t} dt' \alpha_i(t') A_i(t')
\]

\[
= B_\infty - 3 A_\infty \frac{\alpha_{t,\infty} F(t)}{1 + 6 \alpha_{t,\infty} F(t)} + M_\infty \left\{ t \left[ 3 f_2 \alpha_2(t) + \frac{3}{5} f_1 \alpha_1(t) \right] - \frac{3 \alpha_{t,\infty} G(f_i; t)}{1 + 6 \alpha_{t,\infty} F(t)} \right\}.
\]

Scalar masses (Third Generation and Higgs)

If one neglects the Yukawa couplings except for the top (valid for \( \tan \beta \ll m_t/m_b \)),
masses for $\tilde{\tau}_L$, $\tau_R$, $\nu_\tau$ and $\tilde{\nu}_R$ are the same as their lower generation counterparts. Furthermore, the soft-breaking mass term for $H_1$, $\tilde{m}_{H_1}^2$, is equal to $\tilde{m}_L^2$. The solutions for the others are found to be

\[
\begin{align*}
\tilde{m}_{Q_3}^2(t) &= \tilde{m}_{Q_1}^2(t) - \bar{I}(t), \\
\tilde{m}_{U_3}^2(t) &= \tilde{m}_{U_1}^2(t) - 2\bar{I}(t), \\
\tilde{m}_{H_2}^2(t) &= \tilde{m}_{H_1}^2(t) - 3\bar{I}(t).
\end{align*}
\]

where

\[
\bar{I}(t) = \frac{\alpha_{1\infty}}{1 + 6\alpha_{1\infty}F(t)} \left\{ 3m_\infty^2F(t) + M_\infty^2 \sum_{i=1}^{3} c_i^\Sigma f_i^2 \int_0^t dt' F'(t') \left[ \frac{1 - (\frac{\alpha_i(t')}{\alpha_\infty})^2}{1 + 6\alpha_{1\infty}F(t')} \right] \right\}
\]

with $c_3^\Sigma = \frac{32}{3}, c_2^\Sigma = 6$ and $c_1^\Sigma = \frac{26}{15}$. We can write the last integral as follows:

\[
\int_0^t dt' F'(t')[A_i(t')]^2 = A_\infty^2 \frac{F(t)}{1 + 6\alpha_{1\infty}F(t)} + 2A_{1\infty}M_\infty \tilde{K}(t) + M_\infty^2 \tilde{K}_2(t),
\]

where

\[
\tilde{K}_2(t) = \int_0^t dt' F'(t') \left[ t' \sum_{i=1}^{3} f_i c_i^\alpha \alpha_i(t') - \frac{6\alpha_{1\infty}G(f_i; t')}{1 + 6\alpha_{1\infty}F(t')} \right]^2.
\]

$\bar{I}(t)$ is a quadratic function of $m_\infty, M_\infty$, and $A_\infty$:

\[
\bar{I}(t) = \bar{I}_{SS}m_\infty^2 + \bar{I}_{GG}M_\infty^2 + \bar{I}_{GA}M_\infty A_\infty + \bar{I}_{AA}A_\infty^2,
\]

where

\[
\begin{align*}
\bar{I}_{SS} &= \frac{3\alpha_{1\infty}F(t)}{1 + 6\alpha_{1\infty}F(t)}, \\
\bar{I}_{GG} &= \frac{\alpha_{1\infty}}{1 + 6\alpha_{1\infty}F(t)} \int_0^t dt' F'(t') \left\{ \sum_{i=1}^{3} c_i^\Sigma f_i^2 \left[ \frac{1 - (\frac{\alpha_i(t')}{\alpha_\infty})^2}{1 + 6\alpha_{1\infty}F(t')} \right] \right\}^2 \\
&+ \left[ t' \sum_{i=1}^{3} f_i c_i^\alpha \alpha_i(t') - \frac{6\alpha_{1\infty}G(f_i; t')}{1 + 6\alpha_{1\infty}F(t')} \right]^2 - 6 \left[ \frac{\alpha_{1\infty}G(f_i; t')}{1 + 6\alpha_{1\infty}F(t')} \right]^2, \\
\bar{I}_{GA} &= \frac{2\alpha_{1\infty}}{[1 + 6\alpha_{1\infty}F(t)]^2} G(f_i; t), \\
\bar{I}_{AA} &= \frac{\alpha_{1\infty}F(t)}{[1 + 6\alpha_{1\infty}F(t)]^2}.
\end{align*}
\]

**Appendix C**

Here we present the results of the MSGUT analyses without the GUT relation ($f_i \neq 1$). The calculational scheme is not different from that in Sec.4 except for $f_3 \neq 1$. Now we seek
solutions to give the light stop with its mass $m^2_{\tilde{t}_1} = 15\text{GeV}$ varying the three parameters $(\mu, \tan \beta, m_t)$ for $f_1 = f_2 = 1$ and $f_3 \neq 1$. In this case we can take $m^2_{\tilde{g}}$ values freely because the GUT relation is modified as

$$m^2_{\tilde{g}} = M_3 = f_3 \frac{\alpha_s}{\alpha} \sin^2 \theta_W M_2.$$  

The typical results are shown in Table V.

**Appendix D**

In this Appendix we present the formulae for the widths of the second neutralino decay $\tilde{Z}_2 \to \tilde{t}_1 \tilde{t}_1^* \tilde{Z}_1$ and the gluino decay $\tilde{g} \to \tilde{t}_1 \tilde{t}_1^* \tilde{Z}_1$. Formulae for the other modes can be found in Ref. [40]. The total decay width of the three-body decay $\tilde{X}(k) \to \tilde{t}_1(p_2)\tilde{t}_1^*(p_3)\tilde{Z}_1(p_1)$ can be written as

$$\Gamma = \frac{1}{(2\pi)^3} \frac{1}{32 m^2_{\tilde{X}}} \int \frac{(m^2_{\tilde{X}} - m^2_{\tilde{t}_1})^2}{(m^2_{\tilde{Z}_1} + m^2_{\tilde{t}_1})^2} dm^2_{12} \int \frac{(m^2_{13})_{\text{max}}}{(m^2_{13})_{\text{min}}} dm^2_{13} \sum_{\text{spin}} |M|^2,$$

where

$$m^2_{13} = (E^*_1 + E^*_3)^2 - \left( \sqrt{E^*_1^2 - m^2_{Z_1}} \pm \sqrt{E^*_3^2 - m^2_{\tilde{t}_1}} \right)^2,$$

and $m^2_{ij} \equiv (p_i + p_j)^2$.

The matrix elements $M = \sum_{\alpha} M_{\alpha}$ corresponding to the Feynman diagrams Figs.3 and 6 are given by

$$M_1 = \frac{2i e g_s T^a}{m^2_{12} - m^2_{\tilde{t}}} \left[ \bar{u}(p_1) \left( F^L P_R + F^R P_L \right) (\phi_1 + \phi_2 + m_t) (\cos \theta_t P_R + \sin \theta_t P_L) u(k) \right],$$

$$M_2 = \frac{2i e g_s T^a}{m^2_{13} - m^2_{\tilde{t}}} \left[ \bar{u}(p_1) \left( F^L P_R + F^R P_L \right) (\phi_1 + \phi_3 + m_t) (\cos \theta_t P_R + \sin \theta_t P_L) u(k) \right]$$

for $\tilde{g} \to \tilde{t}_1 \tilde{t}_1^* \tilde{Z}_1$ and

$$M_1 = \frac{2i e^2}{m^2_{12} - m^2_{\tilde{t}}} \left[ \bar{u}(p_1) \left( F^L P_R + F^R P_L \right) (\phi_1 + \phi_2 + m_t) \left( F^L P_L + F^R P_R \right) u(k) \right],$$

$$M_2 = \frac{2i e^2}{m^2_{13} - m^2_{\tilde{t}}} \left[ \bar{u}(p_1) \left( F^L P_R + F^R P_L \right) (\phi_1 + \phi_3 + m_t) \left( F^L P_L + F^R P_R \right) u(k) \right],$$

$$M_3 = \frac{2i e^2}{m^2_{2a} - m^2_{\tilde{h}}} X G_{12} \left[ \bar{u}(p_1) u(k) \right]$$

for $\tilde{Z}_2 \to \tilde{t}_1 \tilde{t}_1^* \tilde{Z}_1$. The coupling parameters are defined by

$$F^L_k = \frac{m_t N^L_{k1}}{2 m_W \sin \theta_W \sin \beta} \cos \theta_t + (e_a N^L_{k1} + A^R_{k1} N^L_{k2}) \sin \theta_t$$
\[ F_k^R = \left( e_u N'_k + A^L_k N'_2 \right) \cos \theta_t \frac{m_t N'_k}{2m_W \sin \theta_W \sin \beta} \sin \theta_t \]

\[ G_{kl} = \frac{1}{2} \left( \sin \alpha N_{l3} + \cos \alpha N_{k1} \right) \left( \frac{N_{l2}}{\sin \theta_W} - \frac{N_{l1}}{\cos \theta_W} \right) + \frac{1}{2} \left( \sin \alpha N_{l3} + \cos \alpha N_{l4} \right) \left( \frac{N_{k2}}{\sin \theta_W} - \frac{N_{k1}}{\cos \theta_W} \right), \]

\[ X = \frac{1}{4} \left[ m_Z \sin (\alpha + \beta) (\cos^2 \theta_t A^L_k + \sin^2 \theta_t A^R_k) - \frac{m_t^2}{m_W \sin \theta_W \sin \beta} \cos \alpha \sin \theta_t \cos \theta_t (\mu \sin \alpha - A_t \cos \alpha) \right], \]

where \( A_u^L = (1/2 - e_u \sin^2 \theta_W)/(\cos \theta_W \sin \theta_W) \) and \( A_u^R = e_u \tan \theta_W \). After calculating the matrix element squared, we get

\[ \Gamma(\bar{Z}_2 \to \bar{t}_1 t_1 Z_1) = \frac{\alpha^2}{8\pi m_Z^2} \int dm_{12}^2 dm_{13}^2 \left[ \frac{T_{11}}{(m_{12}^2 - m_t^2)^2} + \frac{T_{22}}{(m_{13}^2 - m_t^2)^2} \right. \]

\[ \left. + \frac{T_{12}}{(m_{12}^2 - m_t^2)(m_{13}^2 - m_t^2)} + \frac{U_{33}}{(m_{23}^2 - m_h^2)^2} \right. \]

\[ \left. + \frac{U_{13}}{(m_{13}^2 - m_t^2)(m_{23}^2 - m_h^2)} + \frac{U_{23}}{(m_{12}^2 - m_t^2)(m_{23}^2 - m_h^2)} \right] \]

where

\[ T_{ij} = [F^L_t F^L_j]^2 + (F^R_t F^R_j)^2] T^A_{ij} + 2([F^L_t]^2 + (F^R_t)^2] F^L_t F^R_j T^B_{ij} \]

\[ + 2[(F^L_t)^2 + (F^R_j)^2] F^L_t F^R_t T^C_{ij} + [(F^L_t F^R_j)^2 + (F^R_t F^L_j)^2] T^D_{ij} \]

\[ + 4F^L_t F^R_t F^L_j F^R_j T^E_{ij}, \]

\[ T^A_{11} = m_t^2 (m_{Z_2}^2 + m_{Z_1}^2 - m_{t_1}^2) + (m_{Z_2}^2 - m_{t_1}^2)(m_{Z_1}^2 - m_{t_1}^2) \]

\[ T^B_{11} = m_t m_{Z_2}^2 (m_{Z_2}^2 + m_{Z_1}^2 - m_{t_1}^2) \]

\[ T^C_{11} = m_t m_{Z_1}^2 (m_{Z_2}^2 + m_{Z_1}^2 - m_{t_1}^2) \]

\[ T^D_{11} = m_t^2 (m_{Z_2}^2 + m_{Z_1}^2 - 2m_{t_1}^2) \]

\[ T^E_{11} = m_{Z_1} m_{Z_2} (m_{Z_1}^2 + m_{Z_2}^2) \]

\[ T^A_{12} = 2(m_{12}^2 + m_{13}^2 + m_{Z_2}^2 - m_{t_1}^2) \]

\[ T^B_{12} = m_t m_{Z_2}^2 (m_{12}^2 + m_{13}^2 + 2m_{Z_2}^2 - m_{t_1}^2) \]

\[ T^C_{12} = m_t m_{Z_1}^2 (m_{12}^2 + m_{13}^2 + 2m_{Z_1}^2 - m_{t_1}^2) \]

\[ T^D_{12} = 2m_t^2 (m_{12}^2 + m_{13}^2 - 2m_{t_1}^2) \]

\[ T^E_{12} = m_{Z_1} m_{Z_2} (m_{Z_1}^2 + m_{Z_2}^2 - 2m_{t_1}^2 + 2m_t^2) \]

\[ U_{33} = 2X^2 G^L_{12} (m_{12}^2 + m_{13}^2 - 2m_{t_1}^2 + 2m_{Z_1} m_{Z_2}) \]

\[ U_{13} = 2X G_{12} [(F^L_t F^L_t + F^R_t F^L_j) m_t (m_{12}^2 + m_{13}^2 + 2m_{Z_1} m_{Z_2} - 2m_{t_1}^2) \]

\[ + (F^L_t F^L_t + F^R_t F^R_j) (m_{Z_1}^2 + m_{Z_2}^2)(m_{12}^2 - m_{t_1}^2 + m_{Z_1} m_{Z_2})] \]

and

\[ T^A_{22} = T^A_{11} (m_{12}^2 \leftrightarrow m_{13}^2) \]

\[ U_{23} = U_{13} (m_{12}^2 \leftrightarrow m_{13}^2). \]
The formula for $\Gamma(g \rightarrow \tilde{t}_1 \tilde{t}_1 \tilde{Z}_1)$ can be obtained by following replacements in $\Gamma(\tilde{Z}_2 \rightarrow \tilde{t}_1 \tilde{t}_1 \tilde{Z}_1)$,

\begin{align*}
m_{Z_2} &\Rightarrow m_{\tilde{g}} \\
\alpha^2 &\Rightarrow \frac{1}{2}\alpha\alpha_s \\
U_{ij} &= 0 \\
F_2^L &\Rightarrow \cos \theta_t \\
F_2^R &\Rightarrow \sin \theta_t.
\end{align*}
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Figure Captions

Figure 1: Excluded regions in ($\mu$, $M_2$) plane by LEP and Tevatron for $\tan \beta = 2$. Regions denoted by A, B, C and D are excluded by the chargino mass, the visible branching ratio of the $Z$, the invisible width of the $Z$ and the gluino mass, respectively. Contours of $m_{\tilde{Z}_1} = 13\text{GeV}$ and $14\text{GeV}$ are also depicted.

Figure 2: Allowed region in ($\mu$, $\tan \beta$) plane for $M_2 = 22\text{GeV}$ and $24\text{GeV}$.

Figure 3: Feynman diagrams for the gluino decay $\tilde{g} \rightarrow \tilde{t}_1 \tilde{t}_1^* \tilde{Z}_1$.

Figure 4: $m_{\tilde{g}}$ dependence of branching ratios of gluino. Sum over quark flavors $q, q' = u, d, c, s$ are taken. Input parameters are $\tan \beta = 2.0$, $\mu = -150\text{GeV}$, $m_{\tilde{t}_1} = 15\text{GeV}$, $\theta_t = 0.9$ and $M_2 = 22\text{GeV}$. We take $(m_{\tilde{q}^-}, m_t) = (2m_{\tilde{g}^-}, 130\text{GeV})$ for (a) and $(3.8m_{\tilde{g}^-}, 95\text{GeV})$ for (b).

Figure 5: Expected number of events from $p\bar{p} \rightarrow \tilde{g}\tilde{g}X$ at CDF. Input parameters are $\tan \beta = 2.0$, $\mu = -150\text{GeV}$, $m_{\tilde{t}_1} = 15\text{GeV}$, $\theta_t = 0.9$ and $M_2 = 22\text{GeV}$. We take $(m_{\tilde{q}^-}, m_t) = (2m_{\tilde{g}^-}, 130\text{GeV})$ for type I and $(3.8m_{\tilde{g}^-}, 95\text{GeV})$ for type II. Straight line corresponds to upper bound settled by CDF gluino search.

Figure 6: Feynman diagrams for the second neutralino decay $\tilde{Z}_2 \rightarrow \tilde{t}_1 \tilde{t}_1^* \tilde{Z}_1$.

Figure 7: $m_t$ dependence of branching ratios of second neutralino. Input parameters are $\tan \beta = 2.0$, $\mu = -150\text{GeV}$, $m_{\tilde{t}_1} = 15\text{GeV}$, $\theta_t = 0.9$, $M_2 = 22\text{GeV}$, $m_h = 60\text{GeV}$, $\alpha = -0.6$ and $m_{\tilde{t}} = m_{\tilde{q}^-}$. We take $(m_{\tilde{q}^-}, A_t) = (2m_{\tilde{g}^-}, 300\text{GeV})$ for (a) and $(3.8m_{\tilde{g}^-}, 600\text{GeV})$ for (b).

Figure 8: Excluded regions in ($\mu$, $M_2$) plane ($\tan \beta = 2$) by LEP and Tevatron for type I (a) and type II (b) solution. Definitions of regions are the same as those given in Fig.1.

Figure 9: Expected number of events from $Z \rightarrow hZ^*$ at LEP for neutrino channel (a) and lepton (muon) channel (b). Input parameters are $\tan \beta = 2$, $m_{\tilde{t}_1} = 15\text{GeV}$, $m_{\tilde{Z}_1} = 13\text{GeV}$, $\alpha = -0.5$ and $L = 93.5\text{pb}^{-1}$. Here kinematical cuts $|\cos \theta_q| < 0.9$ for (a) and $M_{\ell^+\ell^-} > 20\text{ GeV}$ for (b) are adopted.

Figure 10: Stop mass contours in ($\mu$, $\tan \beta$) plane ($M_2 = 22\text{GeV}$) for fixed $m_t$. Each line corresponds to contour of $m_{\tilde{t}_1} = 15\text{GeV}$ for the fixed $m_t$ value. Points denoted by A, B and C are correspond to typical parameter sets in the text.

Figure 11: Total energy $\sqrt{s}$ dependence of left-right asymmetry for the stop production at $e^+e^-$ colliders. For comparison we also plot $A_{LR}$ for the up-type quark production.
**Table I**  Typical parameter sets

| masses in GeV | A     | B     | C     |
|--------------|-------|-------|-------|
| $M_2$        | 22    | 22    | 22    |
| tan $\beta$  | 2.17  | 2.0   | 2.02  |
| $\mu$        | -141  | -159  | -146.4|
| $m_t$        | 92.5  | 92.0  | 90    |
| $M_\infty$   | 26.7  | 26.7  | 26.7  |
| $m_\infty$   | 287.3 | 276.6 | 276.2 |
| $A_\infty$   | 692.6 | 667.1 | 672.3 |
| $\mu_\infty$ | -108.9 | -122.8 | -112.7 |
| $m_{\tilde{t}_1}$ | 15.0  | 15.0  | 15.0  |
| $m_{\tilde{t}_2}$ | 322.5 | 310.7 | 312.8 |
| $\theta_t$   | 0.919 | 0.923 | 0.915 |
| $m_{\tilde{b}_1}$ | 247.4 | 237.8 | 239.1 |
| $m_{\tilde{b}_2}$ | 295.8 | 285.3 | 285.0 |
| $m_{\tilde{\nu}_L}$ | 292.1 | 281.9 | 281.0 |
| $m_{\tilde{\nu}_R}$ | 293.4 | 283.0 | 282.6 |
| $m_{\tilde{d}_L}$ | 299.1 | 288.6 | 288.2 |
| $m_{\tilde{d}_R}$ | 295.4 | 285.0 | 284.6 |
| $m_{\tilde{\ell}_L}$ | 290.5 | 279.7 | 279.3 |
| $m_{\tilde{\ell}_R}$ | 289.7 | 278.9 | 278.5 |
| $m_{\tilde{\nu}}$ | 283.3 | 272.8 | 272.3 |
| $m_h$        | 58.6  | 55.6  | 54.8  |
| $m_A$        | 351.7 | 356.1 | 348.0 |
| $m_H$        | 357.6 | 362.6 | 354.6 |
| $m_{H^+}$    | 357.5 | 361.8 | 354.1 |
| $\alpha$     | -0.46 | -0.49 | -0.49 |
| $m_{\tilde{Z}_1}$ | 13.1  | 13.1  | 13.1  |
| $m_{\tilde{Z}_2}$ | 48.0  | 48.6  | 49.4  |
| $m_{\tilde{Z}_3}$ | 150.5 | 165.5 | 154.1 |
| $m_{\tilde{Z}_4}$ | 178.6 | 194.2 | 183.6 |
| $m_{\tilde{W}_1}$ | 45.1  | 45.1  | 46.0  |
| $m_{\tilde{W}_2}$ | 176.4 | 190.8 | 180.5 |
| $m_{\tilde{q}}$ | 74.4  | 74.4  | 74.4  |
### Table II  Branching ratios of top

|       | A    | B    | C    |
|-------|------|------|------|
| \(t \to t_1 Z_1\) | 0.297 | 0.318 | 0.400 |
| \(t \to \tilde{t}_1 \tilde{Z}_2\) | 0.231 | 0.214 | 0.286 |
| \(t \to \tilde{t}_1 \tilde{g}\) | 0.261 | 0.259 | 0.128 |
| \(t \to bW^+\) | 0.211 | 0.209 | 0.186 |

### Table III  Branching ratios of second neutralino

|       | A    | B    | C    |
|-------|------|------|------|
| \(Z_2 \to ff Z_1\) | 0.032 | 0.024 | 0.019 |
| \(\tilde{Z}_2 \to \tilde{t}_1 \tilde{t}_1 \tilde{Z}_1\) | 0.968 | 0.976 | 0.981 |

### Table IV  Branching ratios of gluino

|       | A      | B      | C      |
|-------|--------|--------|--------|
| \(\tilde{g} \to q\tilde{q} Z_1\) | \(4.4 \times 10^{-3}\) | \(5.0 \times 10^{-3}\) | \(4.6 \times 10^{-3}\) |
| \(\tilde{g} \to q\tilde{q} Z_2\) | \(4 \times 10^{-4}\) | \(4 \times 10^{-4}\) | \(3 \times 10^{-4}\) |
| \(\tilde{g} \to q\tilde{q} \tilde{W}_1\) | \(1.4 \times 10^{-3}\) | \(1.7 \times 10^{-3}\) | \(1.3 \times 10^{-3}\) |
| \(\tilde{g} \to \tilde{t}_1 \tilde{t}_1 \tilde{Z}_1\) | 0.994 | 0.993 | 0.994 |
Table V  Typical parameter sets for $f_3 \neq 1$

| masses in GeV | $f_3 = 0.8$ | $f_3 = 2.0$ |
|--------------|--------------|--------------|
| $M_2$        | 22           | 22           |
| $\tan \beta$| 2.25         | 2.0          |
| $\mu$        | -133.4       | -156.5       |
| $m_t$        | 93           | 95           |
| $M_\infty$   | 26.7         | 26.7         |
| $m_\infty$   | 294.7        | 220.1        |
| $A_\infty$   | 715.1        | 511.7        |
| $\mu_\infty$| -103.0       | -121.7       |
| $m_{t_1}$    | 15.0         | 15.0         |
| $m_{\tau_2}$| 326.0        | 297.4        |
| $\theta_t$   | 0.921        | 0.892        |
| $m_{\nu_1}$  | 250.9        | 220.1        |
| $m_{\nu_2}$  | 300.5        | 257.3        |
| $m_{\nu_L}$  | 296.7        | 253.4        |
| $m_{\nu_R}$  | 298.0        | 254.7        |
| $m_{d_L}$    | 303.8        | 260.9        |
| $m_{d_R}$    | 300.1        | 256.9        |
| $m_{\ell_L}$ | 297.9        | 215.1        |
| $m_{\ell_R}$ | 297.1        | 223.0        |
| $m_\nu$      | 290.6        | 215.1        |
| $m_\tilde{\nu}$ | 60.2    | 58.1         |
| $m_A$        | 353.1        | 300.3        |
| $m_H$        | 358.8        | 308.2        |
| $m_{H^+}$    | 358.9        | 307.4        |
| $\alpha$     | -0.45        | -0.50        |
| $m_{Z_1}$    | 13.2         | 13.1         |
| $m_{Z_2}$    | 47.6         | 48.6         |
| $m_{Z_3}$    | 144.3        | 163.5        |
| $m_{Z_i}$    | 171.9        | 192.2        |
| $m_{\tilde{\nu}_1}$ | 45.0 | 45.1 |
| $m_{\tilde{\nu}_2}$ | 170.3 | 189.0 |
| $m_{\tilde{\nu}}$ | 59.5 | 148.9 |
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