Supersymmetric Particle Production at HERA

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Abstract

In the framework of the minimal supersymmetric standard model and the $R$-parity breaking model, we investigate various production processes of the supersymmetric partner at HERA energies. Our emphasis is paid upon the scalar top quark, the partner of top quark, characterized by its lighter mass than the top quark and other scalar quarks in a model. We propose experimentally feasible approaches to search for clean signals of the stop from either its production or decay processes.
1 Introduction

It is very amazing that the Standard Model (SM) has been so successful [1]. Discoveries of $W^\pm$ and $Z^0$ and high level precision electroweak tests by LEP have been vindicating the SM over and over again. The discovery of the top quark at TEVATRON [2] seems to be its culmination. The theoretical prediction agrees very well with experiments. However, the problems of $R_b$ and $R_c$ may shed a light on the fate of the SM [3]. As far as the first approximation is concerned there is no need to go beyond the SM. However, the SM cannot be a fundamental theory in any sense. Actually, there are a large number of free parameters, the arbitrariness of particle masses and mixing angles, and the lack of any explanation for the replication of generations and so on. Moreover, it is known that the gauge hierarchy problem exists in the SM. Theorists have so far attempted to find a way to go beyond the SM via various standpoints: unification [4], technicolor [5], supersymmetry (SUSY) [6] or superstring [7].

Among them the SUSY seems to be the most reasonable candidate to our expectation. A number of theoretical and phenomenological reasons make low energy SUSY attractive with respect to its alternative. The supersymmetry is a symmetry between fermions and bosons. In the SUSY models, the quadratic divergence can be cancelled out owing to both contributions of boson and fermion loops. This is a key point for the fact that the hierarchy problem can be solved in the SUSY models. However, we should note that fermions and bosons, in other words, the ordinary particles and their SUSY partners (sparticles), must be degenerate in mass in the exact SUSY limit. Obviously there appears to be no evidence in nature for such a situation. Therefore, in order to apply SUSY to particle physics, we must consider models in which SUSY is broken. In this case masses of all sparticles must be less than about 1 TeV in order to solve the hierarchy problem. At first sight, it seems to be unnatural that we need a large number of new particles, sparticles, undiscovered yet. However, such new particles play an essential role in some phenomenologically favourable properties in the model. The most impressive evidence in favour of SUSY may be the unification of gauge couplings in SUSY Grand Unified Theories (GUTs) [8]. The dark matter issue [9] in astrophysics also encourages SUSY proponents. Most of us consider the "SUSY world" as a plausible scenario for future particle physics.

HERA at DESY is the world first electron-proton collider whose $ep$ center of mass energy is 314 GeV produced by 30 GeV electron (positron) and 820 GeV proton. The design luminosity is $1.5 \times 10^{34}$ cm$^{-2}$s$^{-1}$. Polarized electron and positron beams are available. Two experimental groups H1 and ZEUS are now engaged in various experiments. Needless to say, HERA is expected to be the powerful machine to serve us utmost information on the nucleon structure. In addition, HERA has a large potentiality to discover signatures of the new physics beyond the SM. This is because of higher energies available than LEP and cleaner events than hadron colliders such as TEVATRON. Also the unique signature at the $ep$ collider which is not obtained from electron-positron colliders nor hadron colliders reveals us a striking feature of new physics.

The purpose of the present paper is to review the search for sparticles at the $ep$ collider HERA. Our emphasis will be paid upon the scalar top quark (stop), the partner of top quark, characterized by its lighter mass than the top quark and other scalar quarks (squarks) in a model described later [10, 11]. The existence of the light stop could give a clue to the issue on $R_b$ [12]. The expected mass of the stop is clearly within the reach...
of HERA. Particularly, the squarks could be singly produced in $ep$ collisions \cite{13, 14, 15} in the $R$-parity breaking model (RBM) \cite{16} because of the existence of the electron-quark-squark couplings. Clearly HERA is the best machine to search for the stop in such models, because we expect the remarkable peak structure in the Bjorken parameter $x$ distribution. We calculate cross sections together with detectable signals expected to observe at HERA. Although we have studied production processes for selectron, sneutrino or photino at HERA \cite{17} we do not enter their details because of their discovery potential to be not so large.

The organization of the paper is as follows. In Sec. 2 we develop the theoretical framework of Minimal SUSY Standard Model (MSSM) and the origins of the light stop mass is discussed. Decay modes of the stop with mass less than other squarks and gluino are also exploited. Experimental status of SUSY particle search is extensively presented in Sec. 3. Not only status on HERA, but also data on the search from LEP, TEVATRON and TRISTAN are concisely summarized. Various analyses on the sparticle production and their expected signatures from decay processes at HERA are reviewed in Sec. 4. Stop production by boson-gluon fusion and $R$-parity breaking single stop production are presented in detail. The concluding remarks are given in the final section 5.

2 Theoretical framework of MSSM

2.1 Particle content

The Minimal SUSY Standard Model (MSSM) includes the minimal particle content. That is, there should be a new particle (sparticle) for each known particle in the SM and the two Higgs doublets. The additional Higgs doublet must be included in order to produce masses of up- and down-type quarks and to realize the chiral anomaly cancellation \cite{6}. The list of particles in the model is shown in Table I (weak eigenstates).

For quarks and charged leptons there should exist two scalar partners per species. These scalars are called squarks and sleptons, generically sfermions. Before $SU(2) \times U(1)$ breaking, the left- and right-sfermions do not mix since they have different $SU(2) \times U(1)$ quantum numbers. After the breaking, however, they can mix each other. Actually, this mixing effect is substantial only for the third-generation sfermions, especially for the stops (superpartners of the top quark). The left-right mixing of the stops will be discussed in the following subsection.

The new fermions are either the superpartner of spin-1 gauge bosons (gauginos) or that of spin-0 bosons (higgsinos). They mix each other when $SU(2) \times U(1)$ symmetry is broken. The mass eigenstates (inos) are usually some complicated mixture of gauginos and higgsinos. Electrically neutral-inos and charged-inos are called neutralinos $\tilde{Z}_k (k = 1 \sim 4)$ and charginos $\tilde{W}_i (i = 1, 2)$, respectively. Gluinos $\tilde{g}$ are free from the mixing since the color $SU(3)$ is not broken.

We need two Higgs doublets in the MSSM \cite{18}. We have five physical Higgs bosons which remain after $SU(2) \times U(1)$ breaking. They are two CP even ($h^0$, $H^0$) and one CP odd ($A^0$) neutral scalars and remaining one charged scalar ($H^\pm$).
2.2 Basic parameters of MSSM

The standard model has 18 fundamental parameters to be determined by experiments. The MSSM has a somewhat larger number of parameters. They are classified as (i) gauge couplings, (ii) superpotential parameters and (iii) soft-breaking parameters.

2.2.1 gauge couplings

The three gauge coupling parameters corresponding to SU(3), SU(2) and U(1) gauge groups are the same as in the SM. All gauge interactions are governed by these couplings. They determine the fermion - sfermion - gaugino interactions and four-point scalar interactions as well as ordinary fermion - fermion - gauge-boson interactions.

2.2.2 superpotential parameters and $R$-parity

In the MSSM ordinary Yukawa interactions are generalized to the terms in superpotential $W(\hat{\phi})$, where $\hat{\phi}$ denote arbitrary chiral superfields. Renormalizability restricts the functional form of the superpotential to

\[ W(\hat{\phi}) = m_{ij} \hat{\phi}_i \hat{\phi}_j + \lambda_{ijk} \hat{\phi}_i \hat{\phi}_j \hat{\phi}_k \] (1)

The parameters $m_{ij}$ and $\lambda_{ijk}$ are further constrained by the gauge symmetry and some discrete symmetries.

A well-known discrete, multiplicative symmetry is the $R$-parity defined by

\[ R = (-)^{3(B-L)+2S}, \] (2)

where $B$, $L$ and $S$ denote the baryon number, the lepton number and the spin, respectively. This formula implies that all ordinary SM particles have even $R$-parity, whereas the corresponding superpartners have odd $R$-parity. Usually we impose the $B - L$ conservation on the MSSM and then the MSSM possesses the $R$-parity invariance. In this case the superpotential has the form:

\[ W = \mu \tilde{H}_1 \tilde{H}_2 + f_\ell \tilde{H}_1 \tilde{E}^c \tilde{L} + f_d \tilde{H}_1 \tilde{D}^c \tilde{Q} + f_u \tilde{H}_2 \tilde{U}^c \tilde{Q}, \] (3)

where we have used the superfield notation in Table I. Here parameters $f$’s correspond to usual Yukawa couplings and $\mu$ is a new supersymmetric parameter which has no counterpart in the SM.

The $R$-parity conservation in scattering and decay processes has a crucial impact on the SUSY phenomenology at high energy colliders. First, the SUSY particles must be produced in pairs. Second, the lightest SUSY particle (LSP) is absolutely stable. On the other hand, the cosmological constraints tell us that the LSP should be electrically and also color neutral. Consequently, it is weakly interacting in ordinary matter. The heavy unstable SUSY particles must finally decay into the LSP. Therefore the canonical signature for the $R$-parity conserving SUSY models at collider experiments is the large missing (transverse) energy.
It should be emphasized that the $R$-parity conservation is not automatically imposed in the MSSM. In fact the $R$-parity breaking superpotential is allowed by the supersymmetry as well as the gauge symmetry \cite{16};

$$W_R = \lambda_{ijk} \tilde{L}_i \tilde{L}_j \tilde{E}^c_k + \lambda'_{ijk} \tilde{L}_i \tilde{Q}_j \tilde{D}^c_k + \lambda''_{ijk} \tilde{U}_i \tilde{D}_j \tilde{D}^c_k. \quad (4)$$

The first two terms violate $L$ and the last term violates $B$. If we want to explain some unresolved problems such as (i) the cosmic baryon number violation, (ii) the origin of the masses and the magnetic moments of neutrinos and (iii) some interesting rare processes in terms of the $L$ and/or $B$ violation, the $R$-parity breaking terms must be incorporated in the MSSM.

In such a model the signatures for SUSY particles observed at collider experiments should be very different from the canonical one. We will discuss the typical $R$ signals at HERA in Sec.4.

\subsection*{2.2.3 soft-breaking parameters}

SUSY is not an exact symmetry of nature since our world is not manifestly supersymmetric. In the MSSM the SUSY breaking is induced by the soft-SUSY breaking terms, which do not introduce quadratic divergences. Hence the solution of the naturalness problem remains intact. There are four types of soft breaking terms ; (i) gaugino masses $M_i$ ($i = 1 \sim 3$), (ii) masses for the sfermions $\tilde{m}_f$, (iii) trilinear term $A_f$ and (iv) three scalar Higgs mass terms. These three mass parameters can be re-expressed in terms of two Higgs vacuum expectation values, $v_1$ and $v_2$, and one physical Higgs mass. Here $v_1$ and $v_2$ respectively denote the vacuum expectation values of the Higgs field coupled to $d$-type and $u$-type quarks. $v_1^2 + v_2^2 = (246 \text{ GeV})^2$ is determined from the experimentally measured $W$-boson mass, while the ratio

$$\tan \beta = \frac{v_2}{v_1} \quad (5)$$

is a free parameter of the model.

A generally accepted assumption is that all the three gaugino mass parameters $M_i$ are equal at some grand unification scale $M_X$. Then the gaugino mass parameters can be expressed in terms of one of them, for instance, $M_2$. The other two gaugino mass parameters are given by

$$M_1 = \frac{5}{3} \tan^2 \theta_W M_2 \quad (6)$$

$$M_3 = \frac{-\sin^2 \theta_W}{\alpha} M_2, \quad (7)$$

where $\alpha$ and $\alpha_3$ denote the QED and the QCD coupling constants, respectively.

\subsection*{2.3 Scalar top in MSSM}

\subsubsection*{2.3.1 left-right mixing of stops}

In the framework of the MSSM \cite{6}, the stop mass matrix in the ($\tilde{t}_L, \tilde{t}_R$) basis is expressed by

$$\mathcal{M}_t^2 = \begin{pmatrix} m_{\tilde{t}_L}^2 & a_t m_t \\ a_t m_t & m_{\tilde{t}_R}^2 \end{pmatrix}, \quad (8)$$
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 \[19\] :

\begin{align*}
  m_{\tilde{t}_L}^2 &= \tilde{m}_{Q_3}^2 + m_Z^2 \cos 2\beta \left( \frac{1}{2} - \frac{2}{3} \sin^2 \theta_W \right) + m_t^2, \quad (9) \\
  m_{\tilde{t}_R}^2 &= \tilde{m}_{U_3}^2 + \frac{2}{3} m_Z^2 \cos 2\beta \sin^2 \theta_W + m_t^2, \quad (10) \\
  a_t &= A_t + \mu \cot \beta. \quad (11)
\end{align*}

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

\begin{align*}
  \tilde{m}_{Q_3}^2 &= \tilde{m}_{Q_1}^2 - \tilde{I}, \quad (12) \\
  \tilde{m}_{U_3}^2 &= \tilde{m}_{U_1}^2 - 2\tilde{I}, \quad (13)
\end{align*}

where $\tilde{I}$ is a function proportional to the top quark Yukawa coupling $\alpha_t$ and is determined by the renormalization group equations (RGEs) \[20\] in the minimal supergravity GUT (MSGUT). Throughout this paper we adopt the notation of Ref.\[21\].

There are two origins for lightness of the stop compared to the other squarks and sleptons, i) smallness of the diagonal soft masses $m_{\tilde{t}_L}^2$ and $m_{\tilde{t}_R}^2$ and ii) the left-right stop mixing. Both effects are originated from the large Yukawa interaction of the top. The origin i) can easily be seen from Eqs.(9) \sim (13). The diagonal mass parameters $m_{\tilde{t}_L}^2$ and $m_{\tilde{t}_R}^2$ in Eq.(8) have possibly small values owing to the negative large contributions of $\tilde{I}$ proportional to $\alpha_t$ in Eqs.(12) and (13). It should be noted that this contribution is also important in the radiative $\text{SU}(2) \times \text{U}(1)$ breaking in the MSGUT. The Higgs mass squared has an expression similar to Eqs.(12) and (13) ;

\begin{equation}
  \tilde{m}_{H_2}^2 = \tilde{m}_{L_1}^2 - 3\tilde{I},
\end{equation}

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 (8). The mass eigenvalues are obtained by

\begin{equation}
  m_{\tilde{t}_1, \tilde{t}_2}^2 = \frac{1}{2} \left[ \left( m_{\tilde{t}_L}^2 + m_{\tilde{t}_R}^2 \right) \pm \left( \left( m_{\tilde{t}_L}^2 - m_{\tilde{t}_R}^2 \right)^2 + (2a_t m_t)^2 \right)^{1/2} \right].
\end{equation}

and the corresponding mass eigenstates are expressed by

\begin{equation}
  \begin{pmatrix}
    \tilde{t}_1 \\
    \tilde{t}_2
  \end{pmatrix} = \frac{1}{\sqrt{\tilde{t}_L \cos \theta_t - \tilde{t}_R \sin \theta_t}} \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},
\end{equation}

where $\theta_t$ denotes the mixing angle of stops :

\begin{equation}
  \tan \theta_t = \frac{a_t m_t}{m_{\tilde{t}_R}^2 - m_{\tilde{t}_L}^2}.
\end{equation}

We see that if the SUSY mass parameters and the top mass are of the same order of magnitude, small $m_{\tilde{t}_1}$ is possible owing to the cancellation in the expression \[15\], \[10\], \[11\].
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 ($\tilde{t}_1\tilde{t}_1^\dagger Z$) depends sensitively on the mixing angle $\theta_t$. More specifically, it is proportional to

$$C_{\tilde{t}_1} = \frac{2}{3} \sin^2 \theta_W - \frac{1}{2} \cos^2 \theta_t.$$  \hspace{1cm} (18)

Note that for the special value of $\theta_t \sim 0.98$, the $Z$-boson coupling completely vanishes \[22\].

### 2.3.2 R-parity breaking interactions of stop

Interesting properties of the lighter stop $\tilde{t}_1$ described in the last sub-section are not modified even in the $R$-parity breaking models. In this case the stop could have interactions with the ordinary leptons and/or quarks via the $R$-breaking Lagrangian;

$$L_1 = \lambda'_{i3j} \cos \theta_t \tilde{t}_1 \bar{d}_j P_L \ell_i + h.c.,$$  \hspace{1cm} (19)

which is originated from the superpotential \[4\]. In particular the interaction with the electron;

$$L_2 = \lambda'_{i31} \cos \theta_t (\tilde{t}_1 \bar{d} P_L e + \tilde{t}_1^\dagger \bar{e} P_R d)$$  \hspace{1cm} (20)

will be most suitable for the $e\bar{p}$ collider experiments at HERA because the stop will be produced in the $s$-channel in $e$-$q$ sub-processes \[13\]. Note that the stop can not couple to any neutrinos via $R$-breaking interactions. This is a unique property of the stop which could be useful for us to distinguish the stop from some leptoquarks.

### 2.3.3 decay modes of stop

Here 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:

\[
\begin{align*}
\tilde{t}_1 & \rightarrow t \tilde{Z}_1, \\
& \rightarrow b \tilde{W}_i, \\
& \rightarrow b \ell \tilde{\nu}, \\
& \rightarrow b \nu \tilde{\ell}, \\
& \rightarrow bW \tilde{Z}_k, \\
& \rightarrow b f \tilde{f} \tilde{Z}_k, \\
& \rightarrow c \tilde{Z}_1, \\
& \rightarrow e d,
\end{align*}
\]

where $\tilde{Z}_k (k = 1 \sim 4)$, $\tilde{W}_i (i = 1, 2)$, $\tilde{\nu}$ and $\tilde{\ell}$, respectively, denote the neutralino, the chargino, the sneutrino and the charged slepton. (a) $\sim$ (g) are the $R$-parity conserving decay modes, while (h) is realized by the RB couplings \[20\].

If we consider the stop with mass small enough in the $R$ conserving case, the first five decay modes (a) to (e) are kinematically forbidden due to the observed top mass $m_t \simeq 175$ GeV \[4\] as well as the model independent lower mass bounds for sparticles; $m_{\tilde{W}_1} \gtrsim 45$.
GeV, \( m_{\tilde{\ell}} > 45 \) GeV and \( m_{\tilde{\nu}} > 40 \) GeV. So (f) and (g) survive. Hikasa and Kobayashi [11] have shown that the one-loop mode \( \tilde{t}_1 \rightarrow c \tilde{Z}_1 \) (g) dominates over the four-body mode \( \tilde{t}_1 \rightarrow b f f' \tilde{Z}_1 \) (f). 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. On the other hand, if we consider the RB coupling \( \lambda'_{131} > 0.01 \), which roughly corresponds to the coupling strength detectable at HERA, the decay modes (c) to (g) are negligible due to their large power of \( \alpha \) arising from multiparticle final state or one loop contribution. So there the two body modes (a), (b) and (h) survive in this case.

3 Experimental status of SUSY particle search

Most of the theoretical models assume \( R \)-parity conservation, which has important physical implication: (1) SUSY particles must be produced in pairs, (2) heavy SUSY particles decay to lighter SUSY particles, and (3) the LSP is stable. The phenomenology of \( R \)-parity violating interactions differs from that of the MSSM in two main aspects [23]:

1. the LSP is no longer stable since it is not protected by symmetry, that is, it can decay within the detector; and
2. it is possible to have single production of supersymmetric particles, since the final state is no longer restricted to be \( R \)-parity even.

We present the experimental status of SUSY particle search in \( R \)-parity conserving case and \( R \)-parity violating case separately.

3.1 \( R \)-parity conserving case

3.1.1 Sneutrinos and sleptons

Once enough energy is available, the most efficient way to produce superpartners is to produce them in pair in high-energy leptonic or hadronic interactions. Sneutrinos, which are most likely invisible like neutrinos, can be produced in pair in the \( Z^0 \) boson decay at LEP. The partial decay width for sneutrinos is related to that for neutrinos by [24]

\[
\Gamma(Z^0 \rightarrow \tilde{\nu}_\ell \tilde{\nu}_\ell) = \frac{1}{2} \beta^3 \Gamma(Z^0 \rightarrow \nu_\ell \bar{\nu}_\ell), \quad \ell = e, \mu \text{ or } \tau,
\]

where \( \beta = v_{\tilde{\nu}}/c \). The invisible decay width of the \( Z^0 \) measured at LEP is in good agreement with the standard model of \( N_\nu = 3 \) within errors. Therefore any additional contribution to the invisible decay width of the \( Z^0 \) is limited by

\[
\Delta \Gamma_{\text{invis.}}^{\text{Z}} < 8.8 \text{MeV},
\]

at the 95% CL. This leads to the mass limit of \( m_{\tilde{\nu}} > 43 \) GeV if the three sneutrinos are degenerate in mass, while \( m_{\tilde{\nu}} > 40 \) GeV for one generation sneutrino [25].
Events of the charged slepton pair production through the $Z^0$ decay

$$Z^0 \rightarrow \ell^+ \ell^- \rightarrow \ell^+ \ell^- + 2 \text{ unobserved neutralinos}$$

are characterized by non-coplanar lepton pairs with large missing energy -momentum in the detector. The expected main backgrounds are $Z^0 \rightarrow \tau^+ \tau^- (\gamma)$, $Z^0 \rightarrow e^+ e^-\ell^+ \ell^-$, $Z^0 \rightarrow e^+ e^- + \text{hadrons}$. The 95% CL lower mass limit on $\tilde{e}, \tilde{\mu}$ and $\tilde{\tau}$ obtained at LEP is about 45 GeV almost independent of slepton species [25].

### 3.1.2 Gluinos and squarks

Since gluinos and squarks are strongly interacting, they would be the SUSY particles with the largest cross sections at $\bar{p}p$ colliders. The branching ratios of gluinos and squarks decaying into various chargino and neutralino states depend on respective masses and mixing angles. In a scenario of a low mass gluino and massless photino ($\tilde{\gamma}$) as the LSP, the decay modes become very simple. If $m_{\tilde{g}} < m_{\tilde{q}}$, the squark decays dominantly into $\tilde{g}q$ and the main gluino decay is $\tilde{g} \rightarrow q\bar{q}\tilde{\gamma}$. If $m_{\tilde{q}} < m_{\tilde{g}}$, then the gluino decays dominantly into $\tilde{q}q$ and the main squark decay is $\tilde{q} \rightarrow q\tilde{\gamma}$. Since the photinos escape the detectors without any signature, the SUSY events would have two or more jets with a large amount of imbalanced transverse momenta.

The limits on $m_{\tilde{g}}$ and $m_{\tilde{q}}$, presented by the CDF experiment, are based on the comparison of the observed missing transverse energy ($E_T$) distribution with predictions for the standard model background and the QCD background plus SUSY contribution estimated by the ISAJET Monte Carlo samples [24]. Squark and gluino mass limits obtained at the CDF experiment, under the condition of SUSY parameters $\mu = -250$ GeV, $\tan \beta = 2$ and $m_H = 500$ GeV, are shown in Fig.1 together with those of D0, LEP and UA1/UA2 experiments [27]. A search for squarks and gluinos was made by D0 in the three or more jets plus $E_T$ channel. The number of events observed in the data sample from an integrated luminosity of $13.4 \pm 1.6$ pb$^{-1}$ was consistent with background. The limits on $m_{\tilde{g}}$ and $m_{\tilde{q}}$ plane are shown in Fig.1. For heavy squarks, a lower gluino mass limit of 146 GeV was obtained, and for equal squark and gluino masses a mass limit of 205 GeV was obtained at the 95% CL [27].

### 3.1.3 Charginos and neutralinos

Charginos are the mass eigenstates corresponding to the linear combination of winos and charged Higgsinos. Charginos have been searched for at LEP via their pair production and decays

$$e^+ e^- \rightarrow Z^0 \rightarrow \tilde{W}_1^+ \tilde{W}_1^-,$$

$$\tilde{W}_1^+ \rightarrow q\bar{q} Z_1, \ell \nu Z_1,$$

$$\tilde{W}_1^+ \rightarrow \ell^+ \bar{\nu}_\ell.$$

The mass limit at L3 comes from the line-shape measurement constraint: $m_{\tilde{W}_1} > 44$ GeV. This result is independent from the chargino decays and the field contents(wino and charged Higgsino mixing) [28]. Searches for the signatures of acoplanar leptons, acoplanar
jets and isolated particles were combined at ALEPH. Irrespective of the field content and of the mass of the LSP, masses for chargeno below 45.2 GeV are excluded at the 95% CL\[29\].

Neutralinos, denoted by $\tilde{Z}_1, \tilde{Z}_2, \tilde{Z}_3, \tilde{Z}_4$ in order of increasing mass, are linear combinations of the photino, the zino, and the Higgsinos. The neutralino pair production at LEP is searched for via the following processes:

$$e^+e^- \rightarrow Z^0 \rightarrow \tilde{Z}_1\tilde{Z}_1, \tilde{Z}_1\tilde{Z}_2, \tilde{Z}_2\tilde{Z}_2.$$  

In most cases, the main $\tilde{Z}_2$ decay mechanism is

$$\tilde{Z}_2 \rightarrow \tilde{Z}_1Z^* \rightarrow \tilde{Z}_1f\bar{f} \text{ or } \tilde{Z}_2 \rightarrow \tilde{Z}_1\gamma.$$  

Since the $\tilde{Z}_1$ escapes undetected, the signature of the neutralino pair production would be excess in $\Gamma_{\text{invis.}}$. The signature of the above $\tilde{Z}_2$ decay processes are missing energy due to the undetected $\tilde{Z}_1$ and one or two photons, two or four acollinear and acoplanar leptons, or one to four hadronic jets from the primary quarks [31].

From the absence of any candidate of $\tilde{Z}_1$ and $\tilde{Z}_2$ at DELPHI\[30\] and L3\[31\], limits on the probability that the $Z^0$ would decay into these channels were set. Combining these results with those from the $Z^0$ widths within the MSSM, one can calculate, for every combination of $M, \mu$ and $\tan \beta$ values, the total contribution of $Z^0$ decays into neutralinos and charginos, and can deduce the excluded regions in the ($M, \mu$) plane. The excluded regions for four values of $\tan \beta$, obtained by the L3 experiment through search for $e^+e^- \rightarrow Z^0 \rightarrow \tilde{Z}_1\tilde{Z}_2$ or $\tilde{Z}_2\tilde{Z}_2$ are shown in Fig.2 [31]. For moderate or high values of $\tan \beta$, a significant part of the accessible parameter space is excluded. All neutralino masses are functions of the parameters $M, \mu$ and $\tan \beta$. Therefore, constraints on the MSSM parameter space translate into limits on these masses which are shown in Fig.3 as a function of $\tan \beta$ [31].

Data taken at the TOPAZ and VENUS with the TRISTAN $e^+e^-$ collider have been analyzed to study single photon events. The TOPAZ , based on data of an integrated luminosity of 164 pb$^{-1}$ at the energy $\sqrt{s} = 59$ GeV, observed 4 single photon candidates remained after event selection, which are consistent with the prediction of the standard model plus background [32].

The VENUS , based on data of 225 pb$^{-1}$ at $\sqrt{s} = 58$ GeV, measured the single photon cross section , which is consistent with expectation. No anomalous signal has been observed leading to a lower limit on the mass of SUSY particles under assumption of radiative pair production of photinos. For massless photinos, the scalar electrons of the mass degenerate case are excluded below 51.9 GeV at the 90% CL [33]. Combining this value with other single photon experiments, the lower limit on the scalar electron was determined to be 72.6 GeV at the 95% CL. Fig.4 shows the excluded region on the ($m_{\tilde{\gamma}}, m_{\tilde{\eta}}$) plane together with the results from other experiments compiled in Ref.[33].

Bounds on the gaugino parameters $\mu, \tan \beta$ and $M_2$ have been calculated from experimental data[34]. Fig.5 shows the region in the ($\mu, M_2$ ) plane for $\tan \beta = 2$ excluded by the experimental data on

A lower bound on the mass of lighter chargino $m_{\tilde{W}_1} > 45$ GeV,
upper bound on the branching ratio of the visible neutralino mode $Br(Z^0 \to \text{vis.}) < 5 \times 10^{-6}$,

upper bound on the invisible width of the $Z^0$, and

accepted gluino mass bound at CDF $m_{\tilde{g}} > 150$ GeV (90\% CL),

where the hatched regions of each contour have been excluded. We see that the stringent bound comes from the constraint on the visible width of the $Z^0$ at LEP (B), or from the gluino search at CDF (D).

3.1.4 Stop

Among squarks, it may be that the stop is very light because of the high mass of the top. For a particular value of the mixing angle between $\tilde{t}^R$ and $\tilde{t}^L$, the lighter mass eigenstate $\tilde{t}_1$ decouples from the $Z_0$. Pair production of light stop has been extensively searched for by OPAL in $e^+e^-$ collisions at LEP\cite{35}. The decay mode $\tilde{t}_1 \to c \tilde{Z}_1$ was searched for in the data of the integrated luminosity of 69.1 pb$^{-1}$, which corresponds to $1.68 \times 10^6$ produced $Z^0 \to \tilde{q}\tilde{q}$ events. With no $\tilde{t}_1$ candidates, the region where more than 3.0 events are expected is excluded at 95\% CL. Fig.6 shows the excluded region in the $(\theta_t, m_{\tilde{t}_1})$ plane for the case of the mass difference $\Delta m$ between $\tilde{t}_1$ and $\tilde{Z}_1$ larger than 5 GeV together with the region excluded by lower energy experiments\cite{35}. The structure in the OPAL limits on $\theta_t$ in Fig.6 is due to decoupling of the $\tilde{t}_1$ from the $Z^0$ for $\theta_t \sim 0.98$. The OPAL excludes the existence of the $\tilde{t}_1$ with a mass below 45.1 GeV at the 95\% CL, where $\theta_t$ of $\tilde{t}_R$ and $\tilde{t}_L$ is smaller than 0.85 or greater than 1.15, and $\Delta m$ is greater than 5 GeV. The exclusion regions in the $(m_{\tilde{t}_1}, m_{\tilde{Z}_1})$ plane are shown in Fig.7 for various $\theta_t$. The DELPHI has excluded significant regions in the $(m_{\tilde{t}_1}, m_{\tilde{Z}_1})$ plane for different mixing angle $\theta_t$ by assuming that $\tilde{t}_1 \to c \tilde{Z}_1$, the most likely decay mode allowed by existing limits\cite{36}. The stop below 46 GeV has been excluded for any value of the mixing angle, except for a small window in $\theta_t \sim 1$ \cite{36}.

A search for the light stop has been carried out by the VENUS at TRISTAN\cite{37}. A data sample of 210 pb$^{-1}$ has been analyzed to find events with large acoplanar particle groups, assuming the two-body decay $\tilde{t}_1 \to c \tilde{Z}_1$. The observed number of events was consistent with that expected from the known processes. They obtained the mass limits in the $(m_{\tilde{t}_1}, m_{\tilde{Z}_1})$ plane at 95\% CL, where $\tilde{t}_1$ was excluded for the mass region from 7.6 GeV to 28.0 GeV for massless LSP. The inclusive $D^{\ast\pm}$ production cross section in two-photon processes was measured with the TOPAZ detector\cite{38}. The differential cross sections $d\sigma(D^{\ast\pm})/dp_T$ obtained were compared with theoretical predictions, such as those involving direct and resolved photon processes. A discussion was made on interpretation of the data of $d\sigma(D^{\ast\pm})/dp_T$ assuming stop pair production decaying into a charm quark and a neutralino\cite{38}.

A constraint from $b \to s\gamma$ process to the MSSM has been derived in the light stop region\cite{39}. It was pointed out that although some region in the parameter space is excluded from this process there remains a large parameter space where the amplitude of the $b \to s\gamma$ is suppressed due to partial cancellation between different diagrams. Stops as light as 20 GeV are still viable from the $b \to s\gamma$ constraint. It is also pointed out by Fukugita et
al. that the light stop with a mass as small as $m_{\tilde{t}} \sim 20$ GeV is still an allowed possibility within the standard, minimal SUSY GUTs with the SUSY breaking induced by minimal supergravity \[41\].

Motivated by the fact that the stop may be considerably lighter than other squarks, H. Baer, J. Sender and X. Tata have reinvestigated its signals at the Fermilab Tevatron with simulations using ISAJET 7.07 under the assumption that it decays via $\tilde{t}_1 \to b\tilde{W}_1$ or $\tilde{t}_1 \to c\tilde{Z}_1$. They have shown that experiments should be able to probe stop mass up to $\sim 100$ GeV with an integrated luminosity of 100 pb$^{-1}$ \[42\].

3.2 $R$-parity violating case

3.2.1 Constraints from low-energy processes and neutrino physics

We are particularly interested in the case in which the term $\lambda'_{ijk} \hat{L}_1 \hat{Q}_j \hat{D}_k$ in Eq. (4) in Sec.2 is dominant, since this leads to resonant squark production through the $s$-channel in the $e-q$ subprocess at HERA \[43\].

The term $\lambda'_{11k} \hat{L}_1 \hat{Q}_1 \hat{D}_k$, $(k = 1, 2, 3)$ contributes to the semileptonic decays of quarks, and the term $\lambda'_{12k} \hat{L}_1 \hat{Q}_2 \hat{D}_k$, $(k = 1, 2, 3)$ contributes to forward-backward asymmetries measured in $e^+e^-$ collisions \[13\]. Existing limits for these coupling constants are given in \[14\] as a function of the mass of squarks. The term $\lambda'_{1j1} \hat{L}_1 \hat{Q}_j \hat{D}_1$, $(j = 1, 2, 3)$ also contributes to atomic parity violation \[16\]. The 1σ bound for the coupling of this term is

$$\lambda'_{1j1} < 0.26(m_{\tilde{q}_L}/100\text{GeV}),$$

where the effect of the radiative corrections is taking into account \[16\]. A general analysis of the constraints derived from neutrino physics on explicit $R$-parity breaking in supersymmetric models is presented \[44\]. The upper bounds for $\lambda_{ijk}$ and $\lambda'_{ijk}$ were obtained through (i) neutrino oscillations, (ii) neutrinoless double beta decay and (iii) neutrino cosmology and astrophysics, which were stringent in particular for lepton number violation involving the third generation \[44\].

3.2.2 HERA

$R$-parity violating SUSY particles as well as leptoquarks, leptogluons and excited leptons are being searched for at the HERA $ep$ collider by H1. In a data sample of $\approx 0.5$ pb$^{-1}$ no evidence was found for any squark production in the mass between 45 GeV up to 275 GeV \[45, 46\]. In the supersymmetric model, the single production of squark is allowed through $R$-parity violating couplings. If the model is maximally $R$-parity violating or if the photino mass is larger than the squark mass, these states are indistinguishable from leptoquarks. It implies the rejection limits for a search of squarks from $R$-parity violating supersymmetry. The rejection limits on the coupling $\lambda'_{111}$ are shown in Fig.8 as a function of the mass of squark for various fixed photino masses \[16\].

The new particles are being searched for at HERA by ZEUS. No evidence for leptoquark, leptogluon, squark, or excited electron production was found in a data sample of 26 nb$^{-1}$ \[47\]. Limits on the production of squarks were determined for masses above 25 GeV. Using the leptoquark limits, one can derive limits on the $R$-parity violating coupling at the 98% CL as shown in Fig.9. For electroweak coupling $(\sqrt{4\pi/\alpha})$, this yields lower limits on squark masses at the 95% CL: $m_{\tilde{q}} > 168$ GeV and $m_{\tilde{u}} > 92$ GeV \[47\].
4 Production processes at HERA

4.1 Scalar electron production processes

HERA is the world first \( ep \) collider, at which we can observe various inelastic processes, i.e., \( eg, eg, \gamma g \) collisions and so on, in addition to the elastic processes \( ep \rightarrow Xp \). For the purpose of the sparticle search, each collision processes could be useful for us to detect signatures from different kind of sparticles. In this subsection we discuss the scalar electron (selectron) production at HERA.

4.1.1 \( R \)-parity conserving case

The simplest sparticle production process at HERA would be the SUSY neutral current process \[ ep \rightarrow \tilde{e} \tilde{q} X, \tag{21} \]
where \( \tilde{e} \) and \( \tilde{q} \) denote the selectron and squark, respectively. Even in such a simple process, \( \tilde{e} \) and \( \tilde{q} \) could be produced at the same time and we could get information of both masses of \( \tilde{e} \) and \( \tilde{q} \) from analyses of the process. This is a remarkable advantage of the \( ep \) collider HERA. Phenomenological analyses for the process have been given by many authors [18, 19] and we can find that the total cross section of the process will be \( \sigma \gtrsim 0.1 \) pb for \( m_{\tilde{e}} + m_{\tilde{q}} \lesssim 150 \) GeV. Unfortunately, however, the TEVATRON experiment has been excluded a light squark, \( m_{\tilde{q}} \lesssim 150 \) GeV, by the negative search for the squark pair production \( pp \rightarrow \tilde{q} \tilde{q}^* X \) [26, 27], where they assume 5 flavor squarks are degenerated in mass.

Therefore, we should say that the search for \( \tilde{e} \) and \( \tilde{q} \) by (21) will be not so hopeful.

From the same reason the SUSY charged current process [18, 21]
\[ ep \rightarrow \tilde{\nu} \tilde{q} X \tag{22} \]
and the squark pair production via boson-gluon fusion [52]
\[ ep \rightarrow eqq^* X \quad (\tilde{q} \neq \tilde{t}) \tag{23} \]
will have too small cross section for us to extract from some large background processes. Note that the large lower mass bound from the TEVATRON could not be applicable to the stop.

The SUSY bremsstrahlung process
\[ ep \rightarrow \tilde{e} \tilde{Z}_1 q X \tag{24} \]
has also been analyzed by several authors [18, 54, 49] since this process has the dominant contribution from the \( t \)-channel photon exchange diagram which is not affected by the existence of heavy squark. However, the total cross section for this process will be rather small \( \sigma \gtrsim 0.01 \) pb even for the light selectron, \( m_{\tilde{e}} \lesssim 50 \) GeV, which has been almost excluded by the LEP.
Up to now we have discussed the inelastic processes only. On the other hand, the elastic selectron production
\[ ep \rightarrow \bar{e}\bar{Z}_1 p \] (25)
has been known to have rather large and viable cross section [54, 55]. In particular, Lopez et al. [56] have pointed out recently that the Leading Proton Spectrometer at HERA would be efficient detector, which can measure momentum of the leading proton scattered in very forward region.

The slepton - squark production in the NC and CC processes at LEP \( \otimes \) LHC whose total available energies are larger than HERA is also discussed in Refs.[57, 58].

4.1.2 R-parity violating case

If we consider the R-parity breaking moel, the experimental signatures of sparticles will be completely different from those in the R-parity conserving case as has been shown. First, the lightest sparticle will not be stable and will decay into the ordinary particles via the R-parity interaction. Second, some sparticles could be singly produced in eq sub-processes at HERA via the R-parity coupling \( \lambda' \) in Eq.(11).

In the former aspect, Dreiner et al. [59] have shown that the SUSY neutral current process Eq.(21) will have remarkable signature and the mass reach at HERA will become about \( m_{\tilde{e}} + m_{\tilde{q}} \lesssim 200 \text{ GeV} \) even for very small R-breaking couplings \( \lambda' \sim 10^{-6} \).

The extensive study for the latter situation has been given by Butterworth et al. [14] and by ourself [13, 60, 61]. The single production of the SUSY partners of light quarks has been discussed in Refs.[14, 15]. On the other hand, we have published some papers in which considered the single stop production at HERA. Our works will be summarized in next subsection. Butterworth et al. have analyzed extensively the single production of squarks with RB couplings in the first and second generation [14]. For example,
\[ e^- u \rightarrow \bar{d} \] (26)
could be possible and the d-squark will decay via R-conserving couplings
\[ \bar{d} \rightarrow d\bar{Z}_1. \] (27)
Moreover, the lightest neutralino \( \bar{Z}_1 \) will decay into \( u\bar{d}^- \) or \( \bar{u}d^+ \) through the R-breaking interaction proportional to \( \lambda' \). Here the interesting point is that the decay product could contain not only the electron but also the positron. This is because the neutralino is the Majorana fermion. In fact, the positron signature from the electron-proton collision will be distinguishable from almost standard background. In their model of cascade decays of the squarks to the LSP, the squark mass reach was formed to be at HERA
\[ m_{\tilde{q}} \lesssim 270 \text{ GeV} \text{ for } \lambda' \gtrsim 0.08 \]
and the reach in the Yukawa coupling is
\[ \lambda' \sim 5.3 \times 10^{-3} \text{ for } m_{\tilde{q}} \simeq 100 \text{ GeV}. \]
4.2 Stop production processes

4.2.1 Boson-gluon fusion

The purpose of the present subsection is to discuss the production of the light stop at HERA in the framework of the MSSM with the $R$-parity conservation. The possible existence of the light stop $\tilde{t}_1$ has been discussed by some people [10, 11, 22, 34, 41]. The most promising production process in $ep$ collisions will be the boson-gluon fusion (BGF) [63, 52, 62]:

$$e^\pm p \to e^\pm \tilde{t}_1 \tilde{t}_1^* X.$$ \hspace{1cm} (28)

Its Feynman diagrams are depicted in Fig.10. Although it will be shown that the Weizsäcker-Williams approximation (WWA) is very appropriate to our purpose, we first carry out the exact calculations

$$d\sigma = \frac{1}{2\eta s} G(\eta, \tilde{s}) \int d\eta dPS^{(3)} \frac{1}{8} \sum_{\text{spin}} |M_a + M_b + M_c|^2.$$ \hspace{1cm} (29)

Here $\eta$ denotes the momentum fraction of gluon in the proton and $\tilde{s} \equiv (p_f + p_f')^2$. The gluon distribution function $G(\eta, \tilde{s})$ is given by the set 1 of Ref. [64] and the three-body phase space volume is expressed as

$$dPS^{(3)} = \frac{1}{128\pi^3} \Theta(s + 2m_e^2 - 2\sqrt{m_e^2(s + m_e^2) - W_1^2}) dy dQ^2 dz d\Phi.$$ \hspace{1cm} (30)

with

$$W_1^2 \equiv (2m_{\tilde{t}_1} + m_p)^2 - m_p^2,$$
$$Z \equiv p \cdot p_f / p \cdot q,$$
$$\cos \Phi \equiv \frac{(p \times l_e) \cdot (p \times p_f)}{|p \times l_e| \cdot |p \times p_f|}.$$

The invariant amplitudes corresponding to Fig.10 are, respectively, given by

$$M_a = \frac{i e^2 g_s (q - 2p_f') \mu (p - 2p_f)\nu}{Q^2} T_{\mu\nu},$$ \hspace{1cm} (31)
$$M_b = \frac{i e^2 g_s (q - 2p_f) \mu (p - 2p_f')\nu}{Q^2} T_{\mu\nu},$$ \hspace{1cm} (32)
$$M_c = -\frac{2i e^2 g_s}{Q^2} g_{\mu\nu} T_{\mu\nu}$$ \hspace{1cm} (33)

with $\hat{t} = (p_f - p)^2$, $\hat{u} = (p_f' - p)^2$ and

$$T_{\mu\nu} = \left[ \bar{u}(l) \gamma^\mu \left( \frac{2}{3} + C_{\tilde{t}_1} \frac{Q^2}{Q^2 + m_Z^2} (A_L^e P_L + A_R^e P_R) \right) u(l_e) \right] \epsilon_\nu(p),$$
$$C_{\tilde{t}_1} = \frac{\cos^2 \theta_W - \frac{4}{3} \sin^2 \theta_W}{2 \cos \theta_W \sin \theta_W},$$
$$A_L^e = \frac{1}{2} (\cot \theta_W - \tan \theta_W),$$
$$A_R^e = -\tan \theta_W,$$
$$P_L = \frac{1}{2} (1 + \gamma_5).$$
It is worth mentioning that the constant $C_{t_1}$ representing the strength of $\bar{t}_1 t_1 Z$ coupling depends sensitively on the mixing angle $\theta_t$. In particular, $C_{t_1} = 0$ if $\theta_t = \cos^{-1}(\frac{2}{\sqrt{3}} \sin \theta_W) \simeq 0.98$. On the other hand, $\bar{t}_1 t_1 \gamma$ and $\bar{t}_1 t_1 g$ couplings do not depend on $\theta_t$. Therefore, the $Z$ boson contribution to the cross section depends on $\theta_t$, while that of the photon is independent from $\theta_t$.

It is interesting to compare our exact tree level calculation with the Weizsäcker-Williams approximation (WWA);

$$\frac{d\sigma}{dz} = \int dy P(y) \int d\eta G(\eta, \hat{s}') \frac{d\hat{\sigma}}{dz}, \quad (34)$$

where $\hat{s}' \equiv y \eta s$. The WWA factorizes the cross section of the process, shown in Fig. 10, into the probability for emitting photon from lepton;

$$P(y) = \frac{\alpha}{2\pi} \frac{1 + (1-y)^2}{y} \log \frac{Q_{\text{max}}^2}{Q_{\text{min}}^2} \quad (35)$$

and the BGF cross section involving real photon;

$$\frac{d\hat{\sigma}}{dz} = \frac{4}{9} \frac{\pi \alpha s}{\hat{s}'^3 z^2 (1-z)^2} \left[ 2m_{t_1}^4 - 2m_{t_1}^2 s' z (1-z) + s'^2 z^2 (1-z)^2 \right]. \quad (36)$$

The numerical calculations have been performed by using the program packages BASES. The standard model parameters are taken as $\alpha = 1/137$, $\sin^2 \theta_W = 0.23$ and $m_Z = 91.1$ GeV. In Fig. 11 we show the stop mass dependence of the total cross sections $\sigma$ in the exact tree level calculation and in the WWA. It is found that the $\gamma$ contribution is much larger than the $Z$ boson contribution, therefore the WWA is a good approximation. As mentioned already the photon contribution does not depend on the stop mixing angle $\theta_t$. The total cross section, consequently, is insensitive to $\theta_t$. Throughout the calculations in Fig. 11, we have adopted the lower cut for $Q^2$ as $Q^2 > 5$ GeV$^2$ in order to make detectable the scattered electron. Because of the photon dominance, the total cross section significantly increases when smaller $Q^2$ cuts are adopted. If $Q^2$ is cut at the kinematical limit; $Q^2 > m_{t_1}^2 \frac{s'}{1-y}$, the total cross section is about four times as large as that with cut $Q^2 > 5$ GeV$^2$. If we adopt $Q^2$ cut $Q^2 > m_{t_1}^2 \frac{s'}{1-y}$, the detectable cross section turns out to be $\sigma \sim 0.1$ pb in the case of $m_{t_1} \sim 50$ GeV.

### 4.2.2 Resonance production

In the present subsection, we shall concentrate our discussions to $R$-parity breaking (RB) process. We start from a coupling of the stop $\bar{t}_1$

$$\mathcal{L}_{\text{int}} = \lambda'_{131} \cos \theta_t (\bar{t}_1 \bar{d} P_L e + \bar{t}_1 \bar{e} P_R d) \quad (37)$$

originated from the RB superpotential Eq. (4). Here $P_{L,R}$ denote left and right handed chiral projection operators. The coupling Eq. (20) is the most suitable for the $ep$ collider experiments at HERA, because the stop will be produced through the $s$-channel in the $e$-d subprocess as shown in Fig. 12 for example,

$$e^+ p \rightarrow (\bar{t}_1 X) \rightarrow e^+ q X. \quad (38)$$
For simplicity, we will assume $\lambda'_{131}$ to be only non-zero coupling parameter in the following. The upper bound on the strength of coupling has been investigated through the low-energy experiments [10] and the neutrino physics [44]. The most stringent bound $\lambda'_{131} \sim 0.3$ comes from the atomic parity violation experiment [10].

As mentioned in 2.3.3, only two-body decay modes $(a), (b)$ and $(h)$ are targets of our study. The formulae of the decay width for each mode are respectively given by

$$
\Gamma(t_1 \to Z_i) = \frac{\alpha}{2m_{t_1}} \lambda^2 (m_{t_1}^2, m_{t_1}^2, m_{Z_i}^2)
$$

$$
\times \left[ (|F_L|^2 + |F_R|^2)(m_{t_1}^2 - m_{t_1}^2 - m_{Z_i}^2) - 4m_t m_{Z_i} Re(F_R F_L^*) \right],
$$

$$(39)
$$

$$
F_L \equiv \frac{m_t N_{14}^* \cos \theta_t}{2m_W \sin \theta_W \sin \beta} + e_u (N_{11}^* - \tan \theta_W N_{12}^*) \sin \theta_t,
$$

$$(40)
$$

$$
F_R \equiv \left( e_u N_{11}^* + \frac{1/2 - e_u \sin^2 \theta_W}{\cos \theta_W \sin \theta_W} N_{12}^* \right) \cos \theta_t - \frac{m_t N_{14}^* \sin \theta_t}{2m_W \sin \theta_W \sin \beta},
$$

$$(41)
$$

and

$$
\Gamma(t_1 \to W_k) = \frac{\alpha}{4 \sin^2 \theta_W m_{t_1}^2} \lambda^2 (m_{t_1}^2, m_b^2, m_{W_k}^2)
$$

$$
\times \left[ (|G_L|^2 + |G_R|^2)(m_{t_1}^2 - m_b^2 - m_{W_k}^2) - 4m_b m_{W_k} Re(G_R G_L^*) \right],
$$

$$(42)
$$

$$
G_L \equiv -\frac{m_b U_{k2}^* \cos \theta_t}{\sqrt{2} m_W \cos \beta},
$$

$$(43)
$$

$$
G_R \equiv V_{k1} \cos \theta_t + \frac{m_t V_{k2} \sin \theta_t}{\sqrt{2} m_W \sin \beta},
$$

$$(44)
$$

and

$$
\Gamma(t_1 \to ed) = \frac{\lambda_{131}^2}{16\pi} \cos^2 \theta_t m_{t_1}^2,
$$

$$(45)
$$

where $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2xy - 2yz - 2zx$. $N_{ij}, V_{kl}$ and $U_{kl}$ respectively stand for the neutralino and chargino mixing angles [3]. The mixing angles as well as masses of the neutralinos $m_{Z_i}$ and the charginos $m_{W_k}$ are determined from the basic parameters in the MSSM($\mu, \tan \beta, M_2$). As seen from Fig. 13 the branching ratio of the stop depends on the stop mass. If the stop is heavy enough, i.e., $m_{t_1} > m_t + m_{Z_i}$ or $m_{t_1} > m_b + m_{W_k}$ and the RB coupling is comparable with the gauge or Yukawa coupling $\lambda'_{131}/4\pi \sim \alpha, \alpha_t$ there is a parameter region where $BR(t_1 \to Z_i)$ or $BR(t_1 \to W_k)$ competes with $BR(t_1 \to ed)$. For light $t_1$, however, the decay $t_1 \to ed$ predominates over other decay channels and we have $BR(t_1 \to ed) \simeq 100\%$.

(A) The case of light stop : $m_{t_1} < m_t + m_{Z_i}$ or $m_{t_1} < m_b + m_{W_k}$

Assuming the stop has $BR(t_1 \to ed) \simeq 100\%$, we can calculate the inclusive differential cross section for $e^\pm p \to e^\pm qX$ with polarized $e^\pm$ beams (see Fig. 12) as follows:
where the coefficients $T_i(e_{L,R}^\pm q)$ are represented in the Appendix. For $e_L \bar{d} \to \bar{t}_1$ followed by the decay $t_1 \to e_L^\pm \bar{d}$ we have

$$
\frac{d\sigma}{dx dq^2}[e_{L,R}^\pm] = \frac{2\pi\alpha^2}{x^2 s^2} \sum_q [g(x, Q^2)] \sum_{i=1}^4 T_i(e_{L,R}^\pm q) + \bar{q}(x, Q^2) \sum_{i=1}^4 T_i(e_{L,R}^\pm q),
$$

(46)

where the decay width $\Gamma(\bar{t}_1 \to ed)$ is given by (45). From Eq.(47) we can expect a clean signal of the stop as a sharp resonance peak in the distribution of the Bjorken parameter $x$. The position peak corresponds to $x = m_{\bar{t}_1}^2 / s$ for fixed $Q^2$. Figure 14 clearly shows the resonance behavior of the production process. The background curve represents the prediction of SM. It is obviously seen from Fig. 14 that the lower $Q^2$ cuts would be very efficient to suppress the background expected from SM, since the $s$-channel resonance contribution is independent of $Q^2$. It would be worthy of remarking that the similar peak could be expected in the leptoquark production at HERA [56].

We should point out that the stop with the RB couplings will be discriminated from most of the leptoquarks by its distinctive properties: (1) the $x$ peak originated from the stop would exist only in the NC (not exist in the CC) process due to no RB stop couplings to neutrinos, (2) $e^+$ beams are more favorable than $e^-$ beams as will be mentioned later. One of the leptoquarks $\tilde{S}_{1/2}$ with the charge $Q = -\frac{2}{3}$ will give the same signature as the RB stop, if the stop decays into the electron and $d$-quark with $BR(\bar{t}_1 \to ed) \simeq 100\%$. The event rate depends on not only the RB coupling strength $\lambda'_{131}$ but also the kind of beams. Figure 15 shows the $y$ distribution at fixed $x$ for $e^−$ and $e^+$ beams. We can see that the $e^+$ beam is more efficient than the $e^−$ one to distinguish the stop signal from the SM background. This can be understood from the fact that the $e^+$ collides with valence $d$-quark in the proton, while the $e^−$ does only with sea $d$-quarks. The difference of their distribution functions is naturally reflected in the cross sections in Fig. 15. It is hopefully expected that the longitudinally polarized $e^+$ and $e^−$ beams will soon be available at HERA. They could be advantageous to suppress the SM background. In Fig. 16 we show the $y(= Q^2/s)$ dependence of the asymmetries defined by

$$
C_R = \frac{d\sigma(e_R^+)/dx dy - d\sigma(e_R^-)/dx dy}{d\sigma(e_R^+)/dx dy + d\sigma(e_R^-)/dx dy},
$$

(48)

and

$$
A_{e^-} = \frac{d\sigma(e_L^-)/dx dy - d\sigma(e_R^-)/dx dy}{d\sigma(e_L^-)/dx dy + d\sigma(e_R^-)/dx dy}.
$$

(49)

It will be seen that the longitudinally polarized $e^\pm$ beams will be useful to identify the stop signal for a reasonable range of $y$ in $C_R$. Finally we show the searchable parameter region

$^1$H1 collaboration at HERA has given the lower mass bound $m_{\bar{t}_1} \gtrsim 98$ GeV on the RB stop from the negative result for the leptoquark $\tilde{S}_{1/2}$ search at 95% CL or $\lambda'_{131} = 0.3$ [34].
in \((\lambda_{131}, m_{\tilde{t}_1})\) plane at HERA. In Fig. 17 the shaded region is experimentally excluded from the atomic parity violation experiment [10]. The area inside the solid contour is accessible at HERA whose production rate is more than ten signal events above the SM background with \(Q^2 > 10^3\ \text{GeV}^2\) in 100 pb\(^{-1}\) running. From Fig. 17 it is seen that the stop mass reach is

\[ m_{\tilde{t}_1} \lesssim 200(270)\text{GeV} \]

for \(e^- (e^+)\) beams at the RB coupling

\[ \lambda_{131} \simeq 0.1 \]

(B) The case of heavy stop: \(m_{\tilde{t}_1} > m_t + m_{Z_i}\) or \(m_{\tilde{t}_1} > m_b + m_{\tilde{W}_k}\)

If the mass of the stop is heavy enough, various decay channels as (a) to (g) compete with \((h)\) and a sharp peak at \(x = m_{\tilde{t}_1}^2 / 2\) can no longer be expected. For the case of \(BR(\tilde{t}_1 \to ed) \ll 100\%\), we should take into account of the processes

\[ ep \to t\tilde{Z}_i X \]

and

\[ ep \to b\tilde{W}_k X \]

as mentioned before. Their Feynman diagrams are shown in Fig. 18. Here we should consider the virtual contributions of the selectron, sneutrino and d-squark with the same RB couplings constants \(\lambda_{131}\). The differential cross sections are given by

\[
\frac{d\sigma}{dx dQ^2}(ep \to t\tilde{Z}_i X) = \frac{\alpha \lambda_{131}^2}{8s^2} \left[ |F_e|^2 \frac{(\hat{u} - m_{\tilde{t}_1}^2)(\hat{u} - m_{Z_i}^2)}{(\hat{u} - m_{e_L}^2)^2} + |F_d|^2 \frac{(\hat{t} - m_{\tilde{t}_1}^2)(\hat{t} - m_{Z_i}^2)}{(\hat{t} - m_{d_R}^2)^2} \right. \\
\left. + \frac{\cos^2 \theta_e \hat{s}}{(\hat{s} - m_{\tilde{t}_1}^2)^2 - m_{\tilde{t}_1}^2 m_{Z_i}^2} \left( |F_L|^2 + |F_R|^2 \right) \hat{s} - m_{\tilde{t}_1}^2 - m_{Z_i}^2 - 4m_t m_{\tilde{Z}_i} Re(F_R F_{*L}) \right] \\
+ 2 Re(F_{e_{*d}}(\hat{u} - m_{e_L}^2)(\hat{t} - m_{d_R}^2)) \\
+ \frac{2 \cos^2 \theta_t \hat{s}}{(\hat{s} - m_{\tilde{t}_1}^2)^2 + m_{\tilde{t}_1}^2 m_{Z_i}^2} Re\left( F_e^* (F_R \hat{u} + F_L m_{\tilde{t}_1} m_{Z_i}) \right) \\
+ \frac{2 \cos^2 \theta_t \hat{s}}{(\hat{s} - m_{\tilde{t}_1}^2)^2 + m_{\tilde{t}_1}^2 m_{Z_i}^2} Re\left( F_d^* (F_R \hat{t} + F_L m_{\tilde{t}_1} m_{Z_i}) \right),
\]

where \(\hat{s} = xs, \hat{t} = -Q^2\) and

\[ F_{e_{*d}} = e_e N_{11} - \frac{1/2 + e_e \sin^2 \theta_W}{\cos \theta_W \sin \theta_W} N_{12}, \]

\[ F_{d_{*d}} = e_e N_{11}' - e_d \tan \theta_W N_{12}', \]
would be expected for us a favourable evidence for our arguments on the stop. As seen from Fig. 19 a few events for the transverse momentum distribution of scattered muon from the process (55) under would emerge as final products to be detectable. Figure 21 shows the Monte Carlo events as a characteristic signature of the heavy stop with mass $\tilde{m}_t \sim 150$ GeV for $e^+$ beams. As far as $e^-$ beams are concerned only $e^-p \rightarrow b\tilde{W}_kX$ would be detectable for $m_{\tilde{t}_1} \sim 170$ GeV. In our model the LSP, the lightest neutralino $\tilde{\chi}_1^0$ will decay into $R$-even particles via only non-zero RB coupling $\lambda_{131}^\prime$. A typical decay chain will be (see Fig. 20)

$$ep \rightarrow t\tilde{Z}_1X \rightarrow (bW)(bd\nu)X \rightarrow (b(\ell\nu))(bd\nu)X.$$  

(54)

The chargino $\tilde{W}_1$ will also decay into $R$-even particles through $\tilde{Z}_1$ like

$$ep \rightarrow b\tilde{W}_1X \rightarrow (b\ell\nu\tilde{Z}_1)X \rightarrow b(\ell(\nu bd\nu))X.$$  

(55)

In both processes (54) and (55) 2b-jets+jet+lepton+missing transverse momentum $P_T$ would emerge as final products to be detectable. Figure 21 shows the Monte Carlo events for the transverse momentum distribution of scattered muon from the process (55) under the condition of the integrated luminosity $L = 300$ pb$^{-1}$ for $e^-p$ collisions at HERA. The branching ratio $BR(\tilde{W}_1 \rightarrow \nu\mu\tilde{Z}_1)$ is assumed to be $1/2$ [77]. Shown in Fig. 21 are distributions for $m_{\tilde{t}_1} = 150$ GeV and $m_{\tilde{t}_1} = 100$ GeV together with background(short-dashed line) of muon events coming from charged current process(CC) $e^-p \rightarrow \nu qX$ and $W$-gluon fusion process (WGF) $e^-p \rightarrow \nu s\bar{c}X, \nu bX$. The generators LEPTO [68] and AROMA [59] with JETSET [70] have, respectively, been used for CC and WGF. From Fig. 19 we can see that $ep \rightarrow b\tilde{W}_1X$ will have a reasonably feasible cross section to which the stop contributes from the s-channel for $m_{\tilde{t}_1} \sim 100$ GeV. This is also the case for $ep \rightarrow t\tilde{Z}_1X$ since the same final states are realized. Thus, in both processes, $ep \rightarrow t\tilde{Z}_1X$ and $ep \rightarrow b\tilde{W}_1X$ a possible typical signature of the stop production would be 2b-jets+jet+lepton+missing transverse momentum $P_T$ owing to the LSP decay via RB coupling. One of the signals to be detected at HERA is characterized by the high $P_T$ spectrum of muons where the lower $P_T$ cut makes the event distinctive from its background. Although it is difficult to discriminate the stop from one of the leptoquark $\tilde{S}_{1/2}$ for light enough stop $m_{\tilde{t}_1} < m_t + m_{\tilde{Z}_1}, m_{\tilde{t}_1} < m_b + m_{\tilde{W}_k}$, a characteristic signature of the heavy stop with mass $m_{\tilde{t}_1} > m_t + m_{\tilde{Z}_1}, m_{\tilde{t}_1} > m_b + m_{\tilde{W}_k}$ could clearly be distinguished from the leptoquark $\tilde{S}_{1/2}$.

Recent observation of a single muon of high $P_T$ event in $e^+p \rightarrow \mu^+X$ by H1 [71] gives us a favourable evidence for our arguments on the stop. As seen from Fig. 19 a few events would be expected for $m_{\tilde{t}_1} = 150$ GeV and present luminosity $L \sim 3$ pb$^{-1}$.
Concluding remarks

We have discussed SUSY particle searches at HERA by paying our particular attention upon the single production of stop in the framework of the MSSM with $R$-parity breaking interactions. Observation of the moderately light stop predicted by the model would be a nice target of HERA experiments so far the RB coupling constant $\lambda'_{131} < \sim 0.3$.

In the case of $\text{BR}(\tilde{t}_1 \to ed) \simeq 100\%$ corresponding to $m_{\tilde{t}_1} < m_t + m_{\tilde{Z}_1}$ or $m_{\tilde{t}_1} < m_b + m_{\tilde{W}_1}$, the stop produced via neutral currents can clearly be seen as a resonance peak at the Bjorken parameter $x$ distribution. Heavier stop, $m_{\tilde{t}_1} > m_t + m_{\tilde{Z}_1}$ or $m_{\tilde{t}_1} > m_b + m_{\tilde{W}_1}$, are produced through charged currents emerge from the following chain of decay processes:

$$ ep \rightarrow t\tilde{Z}_1X \rightarrow (bW)(bd\nu)X \rightarrow (b(\ell\nu))(bd\nu)X. $$

The chargino $\tilde{W}_1$ will also decay into $R$-even particles through $\tilde{Z}_1$ like

$$ ep \rightarrow b\tilde{W}_1X \rightarrow (b\ell\nu\tilde{Z}_1)X \rightarrow b(\ell\nu(bd\nu))X. $$

The evidence for the existence of the stop would emerge from the careful identification of $2b$-jets + jet + $\ell + P_T$. Possible backgrounds come from $e^-p \rightarrow \nu qX$ and $e^-p \rightarrow \nu s\bar{c}X$, $\nu b\bar{c}X$. The transverse momentum distribution of scattering muon will distinguish the stop events more clearly from backgrounds. Extensive Monte Carlo events with a variety of SUSY parameter sets would highly be desirable. Since the production cross section of the stop is unfortunately rather low the high luminosity is very much favourable.

Until LEPII or Next Liner Collider will be built we are sure that HERA would play a unique role to open the "SUSY world" through our stop with mass of 100 to 300 GeV. TEVATRON groups are also enthusiastic about sparticle search at high energies. We are very much expecting HERA and TEVATRON complementarily step forward to enter into the novel region of particle physics by making full use of their specialities of detectors.
Appendix

The analytic expression for the cross section Eq. (58) is given as follows.

\[
\frac{d\sigma}{dx dQ^2}(e^{-L,R}) = \frac{2\pi\alpha^2}{x^2 s^2} \left[ \sum_{i=1}^{2} T_i^u(e_{L,R})(x, Q^2)u(x, Q^2) + \sum_{i=1}^{2} T_i^u(e_{L,R})(x, Q^2)\bar{u}(x, Q^2) \right] + \sum_{i=1}^{4} T_i^d(e_{L,R})(x, Q^2)d(x, Q^2) + \sum_{i=1}^{4} T_i^d(e_{L,R})(x, Q^2)d(x, Q^2),
\]

where

\[
T_i^u(e_{L}) = \frac{(Q^2 - sx)^2}{Q^4} (e^e e^u + \frac{Q^2}{Q^2 + m_Z^2} A_L^e A_R^u),
\]

\[
T_i^d(e_{L}) = \frac{(Q^2 - sx)^2}{Q^4} (e^e e^d + \frac{Q^2}{Q^2 + m_Z^2} A_R^e A_L^d),
\]

\[
T_i^u(e_{R}) = \frac{(Q^2 - sx)^2}{Q^4} (e^e e^u + \frac{Q^2}{Q^2 + m_Z^2} A_L^e A_R^u),
\]

\[
T_i^d(e_{R}) = \frac{(Q^2 - sx)^2}{Q^4} (e^e e^d + \frac{Q^2}{Q^2 + m_Z^2} A_R^e A_L^d),
\]

\[
T_1^d(e_{L,R}) = T_1^u(e_{L,R})\big|_{u\to d}, \quad T_2^d(e_{L,R}) = T_2^u(e_{L,R})\big|_{u\to d},
\]

\[
T_3^d(e_{L}) = -F_{RB}(t_1) \frac{(Q^2 - sx)^2}{Q^2(Q^2 - sx - m_{t_1}^2)^2} (e^e e^d + \frac{Q^2}{Q^2 + m_Z^2} A_L^e A_R^d),
\]

\[
T_3^d(e_{R}) = 0,
\]

\[
T_4^d(e_{L}) = \frac{1}{4} F_{RB}(t_1) \frac{(Q^2 - sx)^2}{Q^2 - sx - m_Z^2 t_1^2},
\]

\[
T_4^d(e_{R}) = 0,
\]

\[
T_1^d(e_{L,R}) = T_1^u(e_{L,R})\big|_{u\to d}, \quad T_2^d(e_{L,R}) = T_2^u(e_{L,R})\big|_{u\to d},
\]

\[
T_3^d(e_{L}) = -F_{RB}(t_1) \frac{s^2 x^2 (sx - m_{t_1}^2) (s^2 x^2 (sx - m_{t_1}^2)^2 + m_{t_1}^2 \Gamma_{t_1}^2)}{(Q^2((sx - m_{t_1}^2)^2 + m_{t_1}^2 \Gamma_{t_1}^2))},
\]

\[
T_3^d(e_{R}) = 0,
\]

\[
T_4^d(e_{L}) = \frac{1}{4} F_{RB}(t_1) \frac{s^2 x^2}{(sx - m_{t_1}^2)^2 + m_{t_1}^2 \Gamma_{t_1}^2},
\]

\[
T_4^d(e_{R}) = 0.
\]
for $e^-$ beams. The formula for $e^+$ beams can be obtained by the following replacement in the above formula for the $e^-$ beams:

\[
\begin{align*}
E_{L,R} & \rightarrow E_{R,L} & & (78) \\
q & \rightarrow \bar{q} & & (79) \\
\bar{q} & \rightarrow q & & (80)
\end{align*}
\]

Here, $e^f$ denote the electromagnetic charge of matter fermion $f$, and

\[
\begin{align*}
A_f^L & \equiv -\frac{T_3^f - e^f \sin^2 \theta_W}{\cos \theta_W \sin \theta_W}, & & (81) \\
A_f^R & \equiv e^f \tan \theta_W, & & (82)
\end{align*}
\]

where $T_3^f$ are the third component of isospin and $\theta_W$ is the Weinberg angle.
References

[1] S. Dimopoulos, *Proc. of the XXVIIth Int. Conf. on High Energy Physics, Glasgow, 1994*, ed. P. J. Bussey and I. G. Knowles, (Institute of Physics Publishing, Bristol and Philadelphia, 1995), p.93

[2] F. Abe et al. (CDF Collab.), *Phys. Rev. Lett.* 74, 2626 (1995); S. Abachi et al. (D0 Collab.), *Phys. Rev. Lett.* 74, 2632 (1995)

[3] K. Hagiwara, talk at *Int. Symposium on Lepton and Photon Interactions at High Energies*, Beijing, 1995

[4] H. Georgi and S. Glashow, *Phys. Rev. Lett.* 32, 438 (1974); J. C. Pati and A. Salam, *Phys. Rev.* D8, 1240 (1973)

[5] S. Weinberg, *Phys. Rev.* D13, 974 (1979); D19, 1277 (1979); L. Susskind, *Phys. Rev.* D20, 2619 (1979); E. Fahri and L. Susskind, *Phys. Rep.* 74, 277 (1981)

[6] For reviews, see, H. Nilles, *Phys. Rep.* 110, 1 (1984); H. Haber and G. Kane, *Phys. Rep.* 117, 75 (1985)

[7] M. B. Green, J. H. Schwartz and E. Witten, "Superstring Theory I, II", (Cambridge University Press, 1987)

[8] U. Amaldi et al., *Phys. Rev.* D36, 1385 (1987); J. Ellis S. Kelley and D.V. Nanopoulos, *Phys. Lett.* B249, 441 (1990); P. Langacker and M. Luo, *Phys. Rev.* D44, 817 (1991); F. Anselmo et al., *Nuovo Cimento* 104A, (1991) 1817; 105A 1357 (1992); A. Zichichi, Talk at 10 Yeras of SUSY Confronting Experiment, Geneva, Switzerland, 7-9 Sep., 1992

[9] For example, see, L. Roszkowski, "Supersymmetric Dark Matter - A Review", Proceedings of 23rd Workshop of the INFN Eloisatron Project, Erice, 1992, ed. L. Cifarelli and V.A. Khoze, (World Scientific, Singapore, 1993), p. 429

[10] J. Ellis and S. Rudaz, *Phys. Lett.* 128B, 248 (1983); G. Altarelli and R. Rückl, *Phys. Lett.* 144B, 126 (1984); I. Bigi and S. Rudaz, *Phys. Lett.* 153B, 335 (1985)

[11] K. Hikasa and M. Kobayashi, *Phys. Rev.* D36, 724 (1987)

[12] J. D. Wells, C. Kolda and G. L. Kane, *Phys. Lett.* B338, 219 (1994); D. Garcia, R. Jimenez and J. Sola, *Phys. Lett.* B347, 321 (1995); X. Wang, J. L. Lopez and D. V. Nanopoulos, CTP-TAMU-25/95 [hep-ph/9506217]

[13] T. Kon and T. Kobayashi, *Phys. Lett.* B270, 81 (1991)

[14] J. Butterworth and H. Dreiner, *Proc. of the HERA Workshop : "Physics at HERA"* 1991, eds. W. Buchmüller and G. Ingelman, Vol.2, p.1079; *Nucl. Phys.* B397, 3 (1993)

[15] J. L. Hewett, "Research Directions for the Decade", *Proc. of 1990 Summer Study on High Energy Physics, Snowmass, 1990*, ed. E. L. Berger, (World Scientific, Singapore, 1992), p.566
[16] V. Barger, G. F. Giudice and T. Han, *Phys. Rev.* **D40**, 2987 (1989)

[17] A. Bartl et al., *Proceedings of the HERA Workshop: Physics at HERA*, W. Buchmüller and G. Ingelman (eds), Vol2, p.1118 (1991)

[18] J. Gunion and H. Haber, *Nucl. Phys.* **B272**, 1 (1986)

[19] For example, H. E. Haber, *Note on Supersymmetry* in Review of Particle properties, Part II, *Phys. Rev.* **D45**, IX.5 (1992)

[20] I. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, *Prog. Theor. Phys.* **67** (1982) 1889 ; **68**, 927 (1982); J. Ellis, J. Hagelin, D. Nanopoulos and K. Tamvakis, *Phys. Lett.* **125B**, 275 (1983); L. Alvarez-Gaumé, J. Polchinski and M. Wise, *Nucl. Phys. B221*, 495 (1983) ; L. Ibáñez and C. López, *Phys. Lett.* **126B**, 54 (1983)

[21] K. Hikasa, "*JLC Supersymmetry Manual*", unpublished

[22] M. Drees and K. Hikasa, *Phys. Lett.* **B252**, 127 (1990)

[23] H. Dreiner, *Proceedings of 23rd Workshop of the INFN Eloisatron Project*, ed. L.Cifarelli and V.A.Khoze, (World Scientific, Singapore,1993), p.195

[24] P. Fayet, *Proceedings of 23rd Workshop of the INFN Eloisatron Project*, ed. L.Cifarelli and V.A.Khoze, (World Scientific, Singapore,1993), p.1

[25] S. Komamiya. Invited talk at the Recontres du Vietnam, Hanoi Vietnam, 1993 (Preprint UT-ICEPP 94-04, 1994)

[26] Ping Hu, *Proceedings SUSY93 Workshop*, ed. P. Nath, (World Scientific, Singapore,1994), p.229

[27] S. Hagopian, *Proc. of the XXVIIth Int. Conf. on High Energy Physics, Glasgow, 1994*, ed. P. J. Bussey and I. G. Knowles, (Institute of Physics Publishing, Bristol and Philadelphia, 1995), p.809

[28] B. Zhou , *Proceedings SUSY93 Workshop*, ed. P. Nath, (World Scientific, Singapore,1994), p.54

[29] T. Medcalf , *Proceedings SUSY93 Workshop*, ed. P. Nath, (World Scientific, Singapore,1994), p.33

[30] P. Lutz, *Proceedings SUSY93 Workshop*, ed. P. Nath, (World Scientific, Singapore,1994), p.45

[31] L3 Collab. M. Acciarri et al., CERN-PPE/95-14 (1995)

[32] TOPAZ Collab., T. Abe, *Proceedings of the 2nd Workshop on TRISTAN Physics at High Luminosity*, ed. H.Sagawa, T. Tsukamoto, I. Watanabe and Y. Yamada, (KEK, Tsukuba, 1994), p.418

[33] VENUS Collab., N. Hosoda et al., *Phys. Lett.* **B311**, 211 (1994)

[34] T. Kon and T. Nonaka, Seikei University preprint, ITP-SU-94/02 (1994)

[35] OPAL Collab., R.Akers et al., *Phys. Lett.* **B337**, 207 (1994)

25
[36] K. Hultqvist, *Proceedings of the International Europhysics Conference on High Energy Physics, Marseille, 1993*, ed. J. Carr and M. Perrottet, (Editions Frontieres, Gif-sur-Yvette Cedex-France, 1994), p.276

[37] VENUS Collab., J. Shirai, *Proceedings of the 2nd Workshop on TRISTAN Physics at High Luminosity*, ed. H. Sagawa, T. Tsukamoto, I. Watanabe and Y. Yamada, (KEK, Tsukuba, 1994), p.405

[38] TOPAZ Collab., R. Enomoto et al., *Phys. Rev.* **D50**, 1879 (1994)

[39] TOPAZ Collab., R. Enomoto et al., *Proceedings of the 2nd Workshop on TRISTAN Physics at High Luminosity*, ed. H. Sagawa, T. Tsukamoto, I. Watanabe and Y. Yamada, (KEK, Tsukuba, 1994), p.369

[40] Y. Okada, *Phys. Lett.* **B315**, 119 (1993)

[41] M. Fukugita, H. Murayama, M. Yamaguchi and T. Yanagida, *Phys. Rev. Lett.* **72**, 3009 (1994)

[42] H. Baer, J. Sender and X. Tata, *Phys. Rev.* **D50**, 4517 (1994)

[43] J. Butterworth and H. Dreiner, *Nucl. Phys.* **B397**, 3 (1993)

[44] K. Enqvist, A. Masiero and A. Riotto, *Nucl. Phys.* **B373**, 95(1992)

[45] H1 Collab., I. Abt et al., *Nucl. Phys.* **B396**, 3 (1993)

[46] T. Köhler, *Proc. of the XXVIIth Int. Conf. on High Energy Physics*, ed. P. J. Bussey and I. G. Knowles, (Institute of Physics Publishing, Bristol and Philadelphia, 1995), p.801

[47] ZEUS Collab., K. W. McLean, *Proceedings of the International Europhysics Conference on High Energy Physics, Marseille, 1993*, ed. J. Carr and M. Perrottet, (Editions Frontieres, Gif-sur-Yvette Cedex-France, 1994), p.268

[48] S. K. Jones and C. H. Llewellyn Smith, *Nucl. Phys.* **B217**, 145 (1983); P. R. Harrison, *Nucl. Phys.* **B249**, 704 (1985); J. A. Bagger and M. E. Peskin, *Phys. Rev.* **D31**, 2211 (1985); J. Bartels and W. Hollik, *Z. Phys.* **C39**, 433 (1988); H. Komatsu and R. Rückl, *Nucl. Phys.* **B299**, 401 (1988); A. Bartl, H. Fraas and W. Majerotto, *Nucl. Phys.* **B297**, 479 (1988); *Z. Phys.* **C41**, 475 (1988)

[49] T. Kon, K. Nakamura and T. Kobayashi, *Z. Phys.* **C45**, 567 (1990)

[50] R. J. Cashmore et al., *Phys. Rep.* **122**, 275 (1985); Y. Eisenberg et al., *Proceedings of the HERA Workshop : Physics at HERA*, W. Buchmüller and G. Ingelman (eds), Vol 2, p.1124 (1991)

[51] T. Kon, K. Nakamura and T. Kobayashi, *Phys. Lett.* **233**, 461 (1989)

[52] M. Drees and K. Grassie, *Z. Phys.* **C28** (1985) 451; K. Gaemers and M. Janssen, *Z. Phys.* **C48** (1990) 491

[53] G. Altarelli et al., *Nucl. Phys.* **B262**, 204 (1985)

[54] H. Tsutsui, K. Nishikawa and S. Yamada, *Phys. Lett.* **245**, 663 (1990)
[55] M. Drees and D. Zeppenfeld, *Phys. Rev.* **D39**, 2536 (1989)

[56] J. L. Lopez et al., *Phys. Rev.* **D48**, 4029 (1993)

[57] A. Bartl et al., *Z. Phys.* **C52**, 677 (1991)

[58] T. Wöhrmann and H. Fraas, *Phys. Lett.* **B336**, 107 (1994) ; *Phys. Rev.* **D52**, 78 (1995)

[59] H. Dreiner and P. Morawitz, *Nucl. Phys.* **B428**, 31 (1994)

[60] T. Kon, T. Kobayashi, S. Kitamura, K. Nakamura and S. Adachi, *Z. Phys.* **C61**, 239 (1994)

[61] T. Kon, T. Kobayashi and S. Kitamura, *Phys. Lett.* **B333**, 263 (1994)

[62] T. Kobayashi, T. Kon, K. Nakamura and T. Suzuki, *Mod. Phys. Lett.* **A7**, 1209 (1992)

[63] G. A. Schuler. *Nucl. Phys.* **B299**, 21 (1988)

[64] E. Eichten et al., *Rev. Mod. Phys.* **56**, 579 (1984) ; **58**, 1065 (1986) (Errata)

[65] S. Kawabata, *Comput. Phys. Commun.* **41**, 127 (1986)

[66] B. Schrempp, *Proceedings of the HERA Workshop : Physics at HERA*, W. Buchmüller and G. Ingelman (eds), Vol 2, p.1034 (1991)

[67] T. Schimert, C. Burgess and X. Tata, *Phys. Rev.* **D32**, 707 (1985)

[68] G. Ingelman, *Proceedings of the HERA Workshop : Physics at HERA*, W. Buchmüller and G. Ingelman (eds), Vol 3, p.1366 (1991)

[69] G. Ingelman and G. A. Schuler, *Proceedings of the HERA Workshop : Physics at HERA*, W. Buchmüller and G. Ingelman (eds), Vol 3, p.1346 (1991)

[70] T. Sjöstrand, *Comput. Phys. Commun.* **39**, 347 (1986)

[71] H1 Collab., DESY preprint, DESY 94-248
Figure Captions

**Figure 1:** D0, LEP and UA1/UA2 squark and gluino mass limits as a function of squark and gluino mass\(^2[27]\).

**Figure 2:** The excluded regions of the MSSM parameter space at 95% CL as a function of the parameters \(M\) and \(\mu\). The kinematical limit corresponds to the sum of the lightest and next to lightest neutralino masses being equal to the center-of-mass energy. The exclusion coming from the lineshape measurement goes beyond the kinematic limit of the direct search\(^3[31]\).

**Figure 3:** Neutralino lower mass limits(95% CL). The lines correspond to the lower mass limit for the different neutralinos, where \(\chi, \chi', \chi''\) and \(\chi'''\) in this figure represent \(\tilde{\chi}^0_1, \tilde{\chi}^0_2, \tilde{\chi}^0_3\) and \(\tilde{\chi}^0_4\) in the text, respectively. The regions below the lines are excluded\(^3[31]\).

**Figure 4:** Excluded region on the \((m_L, m_\gamma)\) plane. VENUS (90% CL, solid line), ASP (90% CL, dashed line), CELLO (90% CL, dotted line), ALEPH (95% CL, dot-dashed line) and combined single photon result (90% CL, solid line)\(^3[33]\).

**Figure 5:** The region in the \((\mu, M_1^2)\) plane for \(\tan \beta = 2\) excluded by the experimental data on

A lower bound on the mass of lighter chargino \(m_{\tilde{W}^+_1} > 45\) GeV,

B upper bound on the branching ratio of the visible neutralino mode \(Br(Z^0 \rightarrow \text{visi.}) < 5 \times 10^{-6}\),

C upper bound on the invisible width of the \(Z^0\), and

D accepted gluino mass bound at CDF \(m_{\tilde{G}} > 150\) GeV (90% CL),

where the hatched regions of each contour have been excluded\(^3[34]\).

**Figure 6:** The excluded region in the \((\theta_L(\varphi_{mix}), m_{\tilde{L}_1})\) plane at 95% CL, for the case of the mass difference \(\Delta m\) between \(\tilde{L}_1\) and \(\tilde{Z}_1\) larger than 5 GeV. The region excluded from the limit on the \(Z^0\) decay width \((\Delta \Gamma \leq 26\) GeV at 95% CL) and limits from previous publications are also shown\(^3[35]\).

**Figure 7:** The excluded regions in the \((m_{\tilde{t}_1}, m_{\tilde{Z}_1})\) plane at 95% CL, where the mixing angle is assumed to be \(\theta_L(\varphi_{mix}) \leq 0.85\) or \(\geq 1.15\) rad (shaded area), and \(\theta_L(\varphi_{mix}) \leq 0.97\) or \(\geq 0.99\) rad (hatched area). The dashed curve shows the contour of the limit from the previous publication\(^3[35]\).

**Figure 8:** Rejection limits at the 95% CL for the coupling \(\lambda_{111}^\prime\) as a function of the squark mass for various fixed photino masses. Regions above the curves are excluded. The limits combine all charged and neutral decays of the \(\tilde{t}\) and \(\tilde{e}\)\(^4[40]\).

**Figure 9:** The 95% CL upper limits on the couplings of squarks versus mass for (a) the \(\tilde{d}\) (in order of decreasing coupling limits based on the NC, CC and combined samples) and (b) the \(\tilde{u}\) and \(\tilde{t}\) using the NC samples only. The sensitivity of the result to the photino mass is also shown\(^4[47]\).
Figure 10: Feynman diagrams for $e^- g \to e^- \tilde{t}_1 \tilde{t}_1^*$ [62].

Figure 11: Total cross-sections for $e^- p \to e^- \tilde{t}_1 \tilde{t}_1^* X$ with cut $Q^2 > 5\text{ GeV}^2$. Solid line, dashed line and dotted line correspond to exact calculation including both $\gamma$ and $Z$, including only $\gamma$ and WWA, respectively [62].

Figure 12: Feynman diagrams for sub-processes $e^\pm q \to e^\pm q$ [63].

Figure 13: $m_{\tilde{t}_1}$ dependence of branching ratio of stop. We take $m_{\tilde{t}_1}=135\text{ GeV}$, $\tan\beta=2$, $\theta_t=1.0$, $\lambda'_{131}=0.1$ and $(M_2 (\text{GeV}), \mu (\text{GeV})) = (50, -100)$ for (a) and $(100, -50)$ for (b) [61].

Figure 14: $x$ distribution at fixed $Q^2$. Adopted parameters are $m_{\tilde{t}_1}=200\text{ GeV}$, $\lambda'_{131}=0.25$ and $\theta_t=0.0$ [60].

Figure 15: $y$ distributions at fixed $x$ using the electron and the positron. $x$ is fixed at 0.2. Adopted parameters are $m_{\tilde{t}_1}=140\text{ GeV}$ and $\theta_t=\pi/4$ [60].

Figure 16: $y$ distribution at fixed $x$ of differential asymmetries $C_R$ (a) and $A_e$ (b). $x$ is fixed at 0.2. Adopted parameters are $m_{\tilde{t}_1}=140\text{ GeV}$ and $\theta_t=0$ [60].

Figure 17: Searchable parameter region at HERA in $(\lambda'_{131}, m_{\tilde{t}_1})$. Kinematical cut is $Q^2 > 10^3\text{ GeV}^2$ [60].

Figure 18: Feynman diagrams for sub-processes $ep \to t\tilde{Z}_i$ and $ep \to b\tilde{W}_k$ [61].

Figure 19: Stop mass dependence of total cross section. We take $m_{\tilde{t}_1}=135\text{ GeV}$, $\theta_t=1.0$, $\tan\beta=2$, $\lambda'_{131}=0.1$, $M_2 = 100\text{ GeV}$, $m_{\tilde{t}} = 200\text{ GeV}$, $m_{\tilde{q}} = 300\text{ GeV}$ and $\mu = -50\text{ GeV}$. Solid, short-dashed, dotted and dashed lines correspond to $e^- p \to b\tilde{W}_1^- X$, $e^+ p \to b\tilde{W}_1^+ X$, $e^- p \to t\tilde{Z}_1 X$ and $e^+ p \to t\tilde{Z}_1 X$, respectively [61].

Figure 20: Feynman diagrams for LSP decay [61].

Figure 21: Monte Carlo events for transverse momentum distribution of scattered muon from $e^- p \to b\tilde{W}_1 X$ (solid lines) together with backgrounds CC and WGF processes (short-dashed line). We take $m_{\tilde{t}}=200\text{ GeV}$, $m_{\tilde{q}}=300\text{ GeV}$, $m_{\tilde{t}}=135\text{ GeV}$, $\theta_t=1.0$, $\tan\beta=2$, $\lambda'_{131}=0.1$, $M_2 = 100\text{ GeV}$, $\mu = -50\text{ GeV}$ and integrated luminosity $L=300\text{ pb}^{-1}$ [61].
| Superfields | Spin-0 | Spin-1/2 | Spin-1 | Color | $T$ | $Y$ | $B$ | $L$ |
|-------------|--------|----------|--------|-------|-----|-----|-----|-----|
| $Q$         | $\tilde{q}_L$ | $q_L$   | $3$    | $\frac{1}{2}$ | $+\frac{1}{2}$ | $+\frac{1}{3}$ | $0$ |
| $\bar{U}^c$ | $\bar{u}_R$   | $(u^c)_L$ | $3^*$  | $0$   | $-\frac{2}{3}$ | $-\frac{3}{3}$ | $0$ |
| $\bar{D}^c$ | $\bar{d}_R$   | $(d^c)_L$ | $3^*$  | $0$   | $+\frac{1}{3}$ | $-\frac{3}{3}$ | $0$ |
| $\tilde{L}$ | $\tilde{\ell}_L$ | $\ell_L$ | $1$    | $\frac{1}{2}$ | $-\frac{1}{2}$ | $0$ | $+1$ |
| $\bar{E}^c$ | $\bar{e}_R$   | $(e^c)_L$ | $1$    | $0$   | $+1$ | $0$ | $-1$ |
| $G$         | $\bar{g}$     | $g$     | $8$    | $0$   | $0$ | $0$ | $0$ |
| $\hat{W}$   | $\hat{W}$     | $W$     | $1$    | $1$   | $0$ | $0$ | $0$ |
| $\hat{B}$   | $\hat{B}$     | $B$     | $1$    | $0$   | $0$ | $0$ | $0$ |
| $\hat{H}_1$ | $H_1$         | $H_{1L}$ | $1$    | $\frac{1}{2}$ | $-\frac{1}{2}$ | $0$ | $0$ |
| $\hat{H}_2$ | $H_2$         | $H_{2L}$ | $1$    | $\frac{1}{2}$ | $+\frac{1}{2}$ | $0$ | $0$ |