Supersymmetric Signals in Electron-Photon Collisions

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ABSTRACT

Associated selectron-neutralino production in the process $e^-\gamma \rightarrow \tilde{e}^-\tilde{\chi}^0$ provides a striking supersymmetric signal: events with a single high $p_\perp$ electron and otherwise only invisible particles. For $e^-\gamma$ collisions obtained at high energy linear colliders through back-scattering of a laser beam, this reaction is shown to be complementary to selectron pair production in the processes $e^+e^- \rightarrow \tilde{e}^+\tilde{e}^-$ and $e^-e^- \rightarrow \tilde{e}^-\tilde{e}^-$, and to be a probe of heavy selectrons beyond the kinematical limit of pair production. The standard model background from $e^-\gamma \rightarrow e^-Z^0$ and $W^-\nu$ is studied and substantially reduced by rapidity and transverse momentum cuts. The minimum required integrated luminosities for observing this supersymmetric signal are given as functions of several model parameters and collider energies.

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1. Introduction

Supersymmetry is often considered the best motivated and most predictive candidate for new physics at the TeV energy scale. Recently, the attractive features of supersymmetry have again been underlined by the observation that the minimal supersymmetric extension of the standard model is consistent with grand unification and proton stability, provided the symmetry breaking scale is not too high [1]. In contrast, the plain standard model is not. Of course, this only provides a very indirect argument in favour of supersymmetry, but it encourages the search for direct evidence of the existence of supersymmetric partners of the standard model particles at the next generation of high energy colliders.

Strongly interacting supersymmetric particles could be observed at the Large Hadron Collider (LHC) and Superconducting Supercollider (SSC), where the existence of squarks and gluinos can be probed up to masses of the order of 1 TeV. In comparison, the potential in discovering sleptons, charginos and neutralinos is quite limited [2]. Ideal machines to search for electroweak superpartners are $e^+e^-$ colliders. At least in the case of charged particles, masses $\tilde{m} \lesssim \sqrt{s_{ee}}/2$ can be directly produced, $\sqrt{s_{ee}}$ being the $e^+e^-$ centre of mass energy. Thus, in order to reach the TeV mass range in the electroweak sector as well, one needs an $e^+e^-$ linear collider such as the CERN Linear Collider (CLIC) [3].

However, this does not necessarily mean that $e^+e^-$ colliders in the energy range, say, $\sqrt{s_{ee}} = 500 - 1000$ GeV are not competitive or uninteresting from the point of view of supersymmetry searches. The point is that supersymmetric particles which only have electroweak interactions may well be lighter than the strongly interacting squarks and gluinos. In fact, such a pattern is expected in a minimal supergravity model [4] where the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gaugino mass parameters $M_3$, $M_2$ and $M_1$, respectively, are related to a single gaugino mass parameter $m_{1/2}$ by renormalization group equations. Assuming $M_3 = M_2 = M_1 = m_{1/2}$ at the grand unification scale, one finds at low energies $M_3 > $
$M_2 > M_1$. The effective gluino mass is given by $M_3$, while $M_2$ and $M_1$ enter the chargino and neutralino mass matrices. The model also provides renormalization group equations for scalar masses involving essentially a gravitino mass parameter $m_0$ and $M_{1,2,3}$. Unless $m_0 \gg M_{1,2,3}$, sleptons are predicted to be lighter than squarks, $m_{\tilde{l}} < m_{\tilde{q}}$, and partners of right-handed standard model fermions tend to be somewhat lighter than the partners of the corresponding left-handed ones, $m_{\tilde{f}_R} \lesssim m_{\tilde{f}_L}$. As a side remark, the $\tilde{t}$ squark plays a special role and may not fit in this pattern. As a consequence, an experimental bound on the selectron mass puts similar constraints on the basic mass parameters of the model as a considerably higher bound on the squark mass, for example. In this sense, searches for sleptons at $e^+e^-$ colliders in the few hundred GeV range are competitive to squark searches at multi-TeV hadron colliders.

In this paper, we study the prospects for producing and detecting selectrons at linear $e^+e^-$ colliders in the energy range $\sqrt{s_{ee}} \approx 500 - 2000$ GeV. More specifically, we compare single $\tilde{e}^-$ production in $e^-\gamma$ collisions, using Bremsstrahlung photons or back-scattered laser beams, with $\tilde{e}^+\tilde{e}^-$ ($\tilde{e}^-\tilde{e}^-$) pair production in $e^+e^-$ ($e^-e^-$) collisions provided by the same linear collider.

2. The Photon Beam

Bremsstrahlung has a rather soft spectrum, given by the familiar equivalent photon function

$$P(y) = \frac{\alpha}{2\pi} \frac{1 + (1 - y)^2}{y} \ln \frac{s_{ee}}{m_e^2}, \quad (2.1)$$

where $\alpha$ is the fine structure constant and $y$ is the photon energy fraction $E_\gamma/E_e$. Since the energy of the effective $e^-\gamma$ collision is substantially degraded in comparison with the $e^+e^-$ collision energy, this option is not expected to be an efficient way to produce heavy selectrons [5].
A more energetic photon beam can be produced by back-scattering a high intensity laser ray on a high energy electron beam [6]. In principle, the entire electron beam can be converted into photons. These photons are then on-shell and their spectrum is hard. The distribution \( P(y) \) of the energy fraction \( y \) of the electron transferred to a photon, \( y = E_\gamma / E_e \), is given by [6]

\[
P(y) = \frac{1}{N} \left( 1 - y + \frac{1}{1 - y} - \frac{4y}{x(1 - y)} + \frac{4y^2}{x^2(1 - y)^2} \right),
\]

where

\[
0 \leq y \leq \frac{x}{x + 1}
\]

and

\[
x = \frac{4E_eE_{\text{laser}}}{m_e^2}.
\]

The factor \( N \) normalizes \( \int dy \, P(y) \) to 1. The electron and laser beams are taken to be aligned and their respective energies are \( E_e \) and \( E_{\text{laser}} \). In what follows, we assume a 100% conversion efficiency and neglect the angular dispersion of the back-scattered photons. We also choose to take \( x = 2(1 + \sqrt{2}) \approx 4.83 \), which is the threshold for electron pair creation in reactions of the back-scattered and laser photons [6]. For larger values of \( x \) the pair creation process rapidly becomes so important that the conversion efficiency drops considerably. Nevertheless, the effective \( e^-\gamma \) energy is now comparable to the \( e^+e^- \) collision energy and, equally important, the \( e^-\gamma \) luminosity is high, to wit \( \mathcal{L}_{e\gamma} \approx \mathcal{L}_{ee} \approx 10^{33} \text{ cm}^{-2}\text{s}^{-1} \). In the following, results are presented for both types of photon beams.

3. Signal and Backgrounds

We concentrate on a particularly striking signal: a single high \( p_\perp \) electron. This signal is obtained in minimal supersymmetric extensions of the standard model [4] where the
The lightest supersymmetric particle (LSP) is the lightest neutralino \( \tilde{\chi}_1^0 \). A selectron-neutralino pair \( \tilde{e}^- \tilde{\chi}_1^0 \) is produced with the selectron subsequently decaying into an electron-neutralino pair \( e^- \tilde{\chi}_1^0 \). This scenario assumes R-parity to be conserved, so that the LSP is stable and remains unobserved.

One can also consider more complicated scenarios involving the production and cascading decays of neutralinos, charginos, sneutrinos and selectrons. Provided at the end of such a shower we are left with a single electron and only invisible neutrinos and LSPs, such a cascade mechanism yields a similar single electron signal, with a less pronounced \( p_{\perp} \) though.

Here we limit ourselves to the simplest process of this kind:

\[
\begin{align*}
e^- \gamma & \to \tilde{e}^- \tilde{\chi}_1^0 \to e^- \tilde{\chi}_1^0 \tilde{\chi}_1^0.
\end{align*}
\]

Our calculations thus really provide a lower bound on the total cross sections for producing the required signal from supersymmetry at \( e^-\gamma \) colliders. In that sense, our results are conservative. The standard model background to the signal arises from the following two reactions:

\[
\begin{align*}
e^- \gamma & \to e^- Z^0 \to e^- \bar{\nu} \nu, \\
e^- \gamma & \to W^- \nu \to e^- \bar{\nu} \nu.
\end{align*}
\]

In this feasibility study we only consider tree level processes in the narrow width approximation. So, the \( \tilde{e}^- \), \( Z^0 \) and \( W^- \) are produced on shell and are left to decay into respectively \( e^- \tilde{\chi}_1^0 \), \( \bar{\nu}\nu \) and \( \bar{\nu}e^- \) with the corresponding branching ratios.

The results depend on four supersymmetry parameters [4]:

- the ratio \( \tan \theta_v = v_2/v_1 \) of the Higgs vacuum expectation values;

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the soft supersymmetry breaking mass parameters $M_2$ and $\mu$ associated with the $SU(2)_L$ gauginos and higgsinos, respectively (for the $U(1)_Y$ gaugino mass parameter $M_1$ we assume $M_1 = M_2 5/3 \tan^2 \theta_w$, where $\theta_w$ is the weak mixing angle, in accordance with the renormalization group evolution from a common value $M_1 = M_2$ at the GUT scale);

- the mass of the selectron $m_{\tilde{e}}$ (for simplicity we assume the supersymmetric partners of the left- and right-handed electrons to have equal masses: $m_{\tilde{e}_L} = m_{\tilde{e}_R}$).

For arbitrary values of $M_2$ and $\mu$ the gauginos and higgsinos mix to form neutralino and chargino mass eigenstates [4]. Roughly, for $|\mu| \gtrsim M_2/2$ the higgsino admixture to the lightest neutralino is small [7]. This is an essential condition to obtain measurable cross sections for the process under study since the higgsino-electron Yukawa coupling is suppressed by the mass of the electron. In this region of parameter space the dependence of $\sigma(e^-\gamma \to \tilde{e}^-\tilde{\chi}_1^0)$ on $\theta_v$ remains very small. If, in addition, $M_2 \gtrsim m_Z$, one finds roughly $m_{\tilde{\chi}_1^0} \approx M_2/2$ and $m_{\tilde{\chi}_2^0} \approx m_{\tilde{\chi}_1^+} \approx M_2$ [7]. For consistency with the assumption that $\tilde{\chi}_1^0$ is the LSP we must in this case require $M_2 \lesssim 2m_{\tilde{e}}$. To get a feeling for the boundaries in parameter space which can be probed, the reader may refer to Fig. 6, where the parameter regions yielding favourable cross sections are clearly displayed.

In what follows we have considered four different scenarios, which are summarized in Table 1. For the first scenario we have chosen $\tan \theta_v = 200$, whereas for the three others we have chosen $\tan \theta_v = 2$. The similarity between the output parameters of scenario 0 and scenario 1 illustrates that $\theta_v$ is not an essential parameter (as long as $|\mu| \gtrsim M_2/2$).

The lowest order Feynman diagrams leading to the $\tilde{e}^-\tilde{\chi}_1^0$, $e^-Z^0$ and $\nu W^-$ final states in $e^-\gamma$ collisions are shown in Figs 1. The matrix elements and differential cross sections $d\sigma/dt$ derived from these diagrams are given in Ref. 8. For the total cross sections one finds
\[\sigma (\bar{e}_{L,R}\gamma \rightarrow \bar{e}_{L,R}\chi_0^0) = \frac{\pi \alpha^2 |G_{L,R}|^2}{s^3} \]

\[\left[ (s - 7m_{\chi_1^0}^2 + 7m_e^2)\sqrt{\lambda} - 4(m_e^2 - m_{\chi_1^0}^2)(s + m_e^2 - m_{\chi_1^0}^2) \ln \left( \frac{s - m_{\chi_1^0}^2 + m_e^2 + \sqrt{\lambda}}{s - m_{\chi_1^0}^2 + m_e^2 - \sqrt{\lambda}} \right) \right], \tag{3.3}\]

\[\sigma (e^-\gamma \rightarrow e^-Z^0) = \frac{\pi \alpha^2 (1 - 4 \sin^2 \theta_w)^2 + 1}{s^3 \sin^2 \theta_w} \]

\[\left[ (s - m_Z^2)(s + 3m_Z^2) + 2(s^2 - 2sm_Z^2 + 2m_Z^4) \ln \left( \frac{(s - m_Z^2)^2}{m_Z^2s} \right) \right], \tag{3.4}\]

\[\sigma (e^-\gamma \rightarrow \nu W^-) = \frac{\pi \alpha^2}{s^3} \frac{1}{4 \sin^2 \theta_w} \]

\[\left[ \frac{(s - m_W^2)(4s^2 + 5sm_W^2 + 7m_W^4)}{m_W^2} - 4(2s^2 + sm_W^2 + m_W^4) \ln \left( \frac{s}{m_W^2} \right) \right], \tag{3.5}\]

where

\[\lambda = \lambda(s, m_e^2, m_{\chi_1^0}^2) = s^2 + m_e^4 + m_{\chi_1^0}^4 - 2sm_e^2 - 2m_e^2m_{\chi_1^0}^2 - 2m_{\chi_1^0}^2s. \tag{3.6}\]

The \(e_L\bar{e}_L\chi_1^0\) and \(e_R\bar{e}_R\chi_1^0\) couplings \(G_L\) and \(G_R\) depend on the photino (\(\tilde{\gamma}\)) and zino (\(\tilde{Z}\)) content of the lightest neutralino \(\tilde{\chi}_1^0\) (we can safely ignore its higgsino (\(\tilde{H}\)) admixture, since the \(ee\tilde{H}\) coupling is proportional to the mass of the electron):

\[G_L = U_{11}Q + U_{21} \frac{T_3 - Q \sin^2 \theta_w}{\sin \theta_w \cos \theta_w} \]

\[G_R = U_{11}^*Q + U_{21}^* \frac{-Q \sin \theta_w}{\cos \theta_w}, \tag{3.7}\]

where \(Q = -1\) and \(T_3 = -1/2\) are the electrons charge and third component of the weak isospin and \(\theta_w\) is the weak mixing angle. \(U_{11}\) and \(U_{21}\) are elements of the unitary matrix which diagonalizes the neutralino mass matrix [4,7]. They, as well as the masses of the neutralinos, depend in a non-trivial manner on \(\theta_w, \mu\) and \(M_2\).
In the cross section formulas Eqs (3.3)-(3.5) we have neglected the mass of the electron $m_e$ everywhere, except in Eq. (3.4) where $m_e$ has been kept in the electron propagator of the second diagram of Fig. 1b. Indeed, the $u$-channel pole is only regularized by the finite mass of the electron. As a result, the electron distribution of the $e^{-}\gamma \rightarrow e^{-}Z^0$ channel is very strongly peaked in the photon direction.

Using Eqs (2.1) and (2.2) we fold all cross sections with the energy distribution $P(y)$ of the photon beam. The laboratory frame is thus not the centre of mass frame and the centre of mass energy $\sqrt{s_{e\gamma}}$ is given by

$$s_{e\gamma} = y s_{ee},$$

where $\sqrt{s_{ee}} = 2E_e$ is the collider energy. The convoluted cross sections are obtained from

$$\sigma(s_{ee}) = \int_{y_{min}}^{y_{max}} dy \ P(y) \sigma(s_{e\gamma}),$$

where the upper integration limit $y_{max}$ is given by 1 in the case of Bremsstrahlung and by Eq. (2.3) in the case of a back-scattered laser beam. The lower limit $y_{min}$ is set by the kinematical threshold, $y_{min} = (m_{e^-} + m_{\tilde{\chi}_1^0})^2/s_{ee}$ for the supersymmetric process and $y_{min} = m_{Z,W}^2/s_{ee}$ for the background processes.

In the narrow width approximation, the cross sections for producing a single electron signal are obtained by multiplying the folded cross sections Eq. (3.9) by the appropriate branching ratios. We have taken the $Z^0 \rightarrow \tilde{\nu}\nu$ and $W^- \rightarrow e^-\tilde{\nu}$ branching ratios to be 20% and 10% respectively. For the $\tilde{e}^- \rightarrow e^-\tilde{\chi}_1^0$ decay [11], the branching ratio is 100% if the selectron is lighter than the second lightest chargino or neutralino. If there are charginos or other neutralinos which are lighter than the selectron, the $\tilde{e}_L^- \rightarrow e_L^-\tilde{\chi}_1^0$ branching ratio can be considerably less than 100%. As long as $|\mu|, M_2 \gtrsim m_Z$, however, the $\tilde{e}_R^- \rightarrow e_R^-\tilde{\chi}_1^0$ branching ratio remains close to 100%.
4. Results

In Figs 2 total cross sections are plotted as functions of the selectron mass. We compare the $\tilde{e}^+\tilde{e}^-$ and $\tilde{e}^-\tilde{e}^-$ production rates at $e^+e^-$ [12] and $e^-e^-$ [13] colliders\(^1\) with the $\tilde{e}^-\tilde{\chi}_1^0$ production at the same colliders operating in the $e^-\gamma$ mode. This is done for scenarios 1-3 summarized in Table 1 and correspondingly three different collider energies: $\sqrt{s_{ee}} = 500$, 1000, 2000 GeV. Although the $e^-\tilde{\chi}_1^0\tilde{\chi}_1^0$ channel cross sections are lower for small selectron masses, only this channel remains when $m_{\tilde{e}} > \sqrt{s_{ee}}/2$. Note that the scenarios considered here yield relatively heavy neutralinos. With lighter neutralinos the cross sections extend even further beyond the kinematical limit of pair production. We also plotted the cross sections obtained with Bremsstrahlung photons off the electron beam in the Weizsäcker-Williams approximation [5]. As expected, they are negligible in comparison with the cross sections obtained with back-scattered laser photon beams.

Fig. 3 displays the behaviour of the total $e^-\gamma \rightarrow e^-\tilde{\chi}_1^0\tilde{\chi}_1^0$ cross section as a function of the collider energy. The four scenarios summarized in Table 1 are considered here, and the selectron mass is set equal to $M_2$. For scenario 0 the $\tilde{e}_L^- \rightarrow e_L^-\tilde{\chi}_1^0$ branching ratio is only 75.8\% because the selectron can also decay into a chargino and neutrino. For scenarios 1-3, however, the selectron can only decay into the LSP and electron. One observes that scenarios 0 and 1, which differ only by the value of $\theta_v$, yield very similar results. The standard model $e^-\gamma \rightarrow e^-\bar{\nu}\nu$ backgrounds are also shown. It appears that the supersymmetric signal is completely swamped. However, since the largest portion of the standard model cross sections is due to the u-channel exchange in $e^-\gamma \rightarrow e^-Z^0$ and the t-channel exchange in $e^-\gamma \rightarrow W^-\nu$ (since $m_e, m_W \ll \sqrt{s_{ee}}$), these cross sections are drastically reduced by an angular or rapidity cut.

The final electron transverse momentum and rapidity distributions are shown in Figs

\(^1\) We have computed these cross sections in the approximation where $m_{\tilde{\chi}_i^0} \ll m_{\tilde{\chi}_i^0}$ ($i=2-4$).
4, for the $Z^0$ and $W^-$ channels and the $\tilde{e}^-$ channel in scenario 1. The selectron mass and the electron beam energy have been chosen to be equal: $m_{\tilde{e}} = \sqrt{s_{ee}}/2 = 250$ GeV. One sees that in the $Z^0$ and (to a lesser extent) $W^-$ channels the final electron is preferentially emitted in the direction of the incoming photon. In contrast, the $\tilde{e}^-$ channel displays very little preference, the decay electron being produced centrally.

In order to enhance the supersymmetric signal relative to the standard model background, we have chosen to impose the following cuts:

$$p_\perp > \sqrt{s_{ee}}/10$$

$$0 < \eta < 2.$$  \hfill (4.1)

The cut cross sections are displayed in Figs 5 as functions of the selectron mass. This is done for scenarios 1 and 2 of Table 1 and respectively two different collider energies: $\sqrt{s_{ee}} = 500, 1000$ GeV. With these two cuts the backgrounds drop by more than one order of magnitude, whereas for selectron masses higher than the beam energy the supersymmetric signal is reduced by about a factor 2. For selectron masses much lower than the beam energy, these cuts also dwarf the supersymmetric signal, but then anyway selectron pair production in $e^+e^-$ and $e^-e^-$ collisions is another promising mechanism as is shown in Figs 2.

The standard model background can be computed with great precision. The accuracy of these calculations can even be checked in the case of the $W^-$ channel, by comparing with the $e^-\gamma \rightarrow \mu^-\bar{\nu}\nu$ signal. In principle, thus, any deviation from these predicted results can be a signal for supersymmetry. Finite statistics, however, preclude too small a supersymmetry signal to emerge from the background statistical fluctuations. Assuming Poisson statistics, to obtain a $R\sigma$ confidence level, the number of signal events $n_{\text{SUSY}}$ needs to be larger than $R$ times the square root of the number of background events $n_{\text{SM}}$: $n_{\text{SUSY}} \geq R\sqrt{n_{\text{SM}}}$. 

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The luminosity required to separate the supersymmetric signal from its standard model background is thus

$$\mathcal{L}_{\text{req}} \geq \frac{R^2 \sigma_{\text{SM}}}{\sigma_{\text{SUSY}}}.$$  \hspace{1cm} (4.2)

Setting $R = 3$ (3$\sigma$ confidence level) we have plotted on the right vertical axis of Figs 5 the minimum required integrated luminosities for having a discovery potential. Also, specializing on a 250 GeV/beam machine we have plotted in Fig. 6, in the $(\mu, M_2)$-plane, the contours of minimum required luminosity for extracting the supersymmetric signal from the background. This plot has been produced with $\tan \theta_v = 2$ and for a selectron with a mass equal to the electron beam energy, so that the selectron cannot be pair produced.

For values of $M_2$ larger than about 380 GeV (for large absolute values of $\mu$) the studied reaction is kinematically forbidden since $m_{\tilde{e}} + m_{\tilde{\chi}_0^1} > \sqrt{s_{ee}}$. The region where $|\mu| \lesssim M_2/2$ is inaccessible because of the large Higgsino content of the LSP. The parameter region to be explored by the Large Electron Positron Collider (LEP 200) is also shown [2,14].

5. Conclusions

It appears possible to convert a high energy electron beam into an almost as energetic and very intense photon beam by back-scattering a suitable laser beam. Linear colliders will thus eventually provide high energy $e\gamma$ and $\gamma\gamma$ collisions, in addition to $ee$ collisions. We have studied to what extent the $e^-\gamma$ option can be exploited for supersymmetry searches. For this purpose, we have concentrated on the most interesting channel, $e^-\gamma \to \tilde{e}^-\tilde{\chi}_1^0 \to \tilde{e}^+\tilde{\chi}_1^0\tilde{\chi}_1^0$, and compared its prospects to the expectations for $\tilde{e}^+\tilde{e}^-$ and $\tilde{e}^-\tilde{e}^-$ pair production in $e^+e^-$ and $e^-e^-$ collisions. For the present study, we assumed the lightest neutralino $\tilde{\chi}_1^0$ to be the lightest supersymmetric particle and, therefore, stable and undetectable.

In summary, our results indicate that the $e^-\gamma$ mode is indeed a valuable option. Not only does the $e^-\chi_1^0\chi_1^0$ channel exhibit an extremely simple and striking signature, but also
the most dangerous background from the standard model processes \( e^-\gamma \rightarrow e^-Z^0 \rightarrow e^-\bar{\nu}\nu \) and \( e^-\gamma \rightarrow W^-\nu \rightarrow e^-\bar{\nu}\nu \) is manageable. Firstly, this background can be calculated with great precision because it only involves leptons and weak gauge bosons. Secondly, it can be substantially reduced experimentally by simple cuts on the rapidity and transverse momentum of the final state electron.

In comparison to selectron searches in \( e^+e^- \) and \( e^-e^- \) collisions at the same linear collider, the situation is as follows. For \( m_{\tilde{e}} < \sqrt{s_{ee}}/2 \), \( \tilde{e}^+\tilde{e}^- \) and \( \tilde{e}^-\tilde{e}^- \) pair production is more favourable than single \( \tilde{e}^- \) production in \( e^-\gamma \) collisions because of a larger signal cross section, while the background problems are similar. Nevertheless, also in this lower mass range, the \( e^-\gamma \) mode is of considerable interest as a cross-check of a signal which might be observed in the \( e^+e^- \) and \( e^-e^- \) collisions. Moreover, \( e^-\gamma \) collisions provide a more direct probe of neutralinos and their properties than \( e^+e^- \) collisions, where one has to consider higher order processes such as \( e^+e^- \rightarrow \tilde{\chi}_1^0\tilde{\chi}_1^0\gamma \).

For \( m_{\tilde{e}} > \sqrt{s_{ee}}/2 \), the role of \( e^-\gamma \) and \( e^+e^- \) or \( e^-e^- \) collisions is reversed. In the \( e^-\gamma \) mode the selectron can still be singly produced and studied directly. In contrast, in the \( e^+e^- \) and \( e^-e^- \) modes it can only be probed indirectly via virtual effects in the \( e^+e^- \rightarrow \tilde{\chi}_1^0\tilde{\chi}_1^0\gamma \) channel. From Figs 2 and 5 it appears, however, that for a fixed integrated luminosity, say \( \int L \, dt = 20 \, \text{fb}^{-1} \), this advantage can only be exploited experimentally at collider energies \( \sqrt{s_{ee}} < 1 \, \text{TeV} \).

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References

1. U. Amaldi, W. de Boer and H. Fürstenau, *Phys. Lett.* **B 260**, 447 (1991).

2. F. Pauss, in *Large Hadron Collider Workshop*, eds G. Jarlskog and D. Rein (CERN 90-10, ECFA 90-133) Vol. I, p. 118.

3. $e^+e^-$

4. H.E. Haber and G.L. Kane, *Phys. Rep.* **117**, 75 (1985).

   H.P. Nilles, *Phys. Rep.* **110**, 1 (1984).

5. M.K. Gaillard, L. Hall and I. Hinchliffe, *Phys. Lett.* **B 116**, 279 (1982).

6. I.F. Ginzburg, G.L. Kotkin, V.G. Serbo and V.I. Telnov, *Nucl. Instr. Meth.* **205**, 47 (1983).

7. H. Komatsu and R. Rückl, *Nucl. Phys.* **B 299**, 407 (1988).

8. G. Altarelli, G. Martinelli, B. Mele and R. Rückl, *Nucl. Phys.* **B 262**, 204 (1985).

9. I.F. Ginzburg, G.L. Kotkin, S.L. Panfil and V.G. Serbo, *Nucl. Phys.* **B 228**, 285 (1983).

10. J.A. Grifols and R. Pascual, *Phys. Lett.* **B 135**, 319 (1984).

11. A. Bartl, H. Fraas and W. Majerotto, *Nucl. Phys.* **B 297**, 479 (1988).

12. A. Bartl, H. Fraas and W. Majerotto, *Z. Phys.* **C 34**, 411 (1987).

13. W.-Y. Keung and L. Littenberg, *Phys. Rev.* **D 28**, 1067 (1983).

14. A. Bartl et al., in *Large Hadron Collider Workshop*, eds G. Jarlskog and D. Rein (CERN 90-10, ECFA 90-133) Vol. II, p. 1033.
Table 1: Four typical scenarios parametrized in terms of $\tan \theta_v, \mu$ and $M_2$, and the resulting lightest neutralino and chargino masses. The unitary matrix elements $U_{11}$ ($U_{21}$) give the photino (zino) content of the lightest neutralino. The $\tilde{e}_L$ branching ratio is given for $m_{\tilde{e}} = M_2$, while $BR(\tilde{e}_R \to e_R \tilde{\chi}_1^0) \approx 1$ for all cases.

| scenario | $\tan \theta_v$ | $\mu$ [GeV] | $M_2$ [GeV] | $m_{\tilde{\chi}^0_1}$ [GeV] | $m_{\tilde{\chi}^-_1}$ [GeV] | $U_{11}$ | $U_{21}$ | BR $\tilde{e}_L \to e_L \tilde{\chi}_1^0$ |
|----------|----------------|-------------|-------------|----------------|----------------|--------|--------|----------------|
| 0        | 200            | -375        | 250         | 123            | 233            | .855   | -.498  | .76            |
| 1        | 2              | -375        | 250         | 127            | 255            | .891   | -.445  | 1              |
| 2        | 2              | -750        | 500         | 250            | 502            | .881   | -.471  | 1              |
Figure Captions

Figs 1: Lowest order Feynman diagrams of the processes

(a) $e^-\gamma \rightarrow \tilde{e}^-\tilde{\chi}^0_1$,
(b) $e^-\gamma \rightarrow e^-Z^0$,
(c) $e^-\gamma \rightarrow W^-\nu$.

Figs 2: Total cross sections as functions of the selectron mass for $e^-\gamma \rightarrow \tilde{e}^-\tilde{\chi}^0_1 \rightarrow e^-\tilde{\chi}^0_1\tilde{\chi}^0_1$ convoluted with the energy spectrum of back-scattered laser photons (full curves), $e^-\gamma \rightarrow \tilde{e}^-\chi_1 \rightarrow e^-\chi_1\chi_1$ convoluted with a Bremsstrahlung spectrum (dotted curves), $e^-e^- \rightarrow \tilde{e}^-\tilde{e}^-$ (dot-dashed curves) and $e^+e^- \rightarrow \tilde{e}^+\tilde{e}^-$ (dashed curves). These cross sections are shown for three different collider energies and supersymmetric scenarios:

(a) $\sqrt{s_{ee}} = 500$ GeV and scenario 1;
(b) $\sqrt{s_{ee}} = 1000$ GeV and scenario 2;
(c) $\sqrt{s_{ee}} = 2000$ GeV and scenario 3.

Fig. 3: Signal and background total cross sections as functions of the collider energy $\sqrt{s_{ee}}$ for $e^-\gamma \rightarrow \tilde{e}^-\tilde{\chi}^0_1 \rightarrow e^-\chi_1\chi_1$ and scenarios 1-3 (full curves) and 0 (dotted curve) with $m_{\tilde{e}} = M_2$, $e^-\gamma \rightarrow e^-Z^0 \rightarrow e^-\bar{\nu}\nu$ (dot-dashed curve) and $e^-\gamma \rightarrow W^-\nu \rightarrow e^-\bar{\nu}\nu$ (dashed curve).
Figs 4: Transverse momentum (a) and rapidity (b) distributions of the final electron in the processes: \( e^- \gamma \rightarrow \tilde{e}^- \chi^0_1 \rightarrow e^- \chi^0_1 \chi_1 \) for scenario 1 and \( m_{\tilde{e}} = 250 \) GeV (diamonds), \( e^- \gamma \rightarrow e^- Z^0 \rightarrow e^- \bar{\nu} \nu \) (crosses) and \( e^- \gamma \rightarrow W^- \nu \rightarrow e^- \bar{\nu} \nu \) (squares), at \( \sqrt{s_{ee}} = 500 \) GeV.

Figs 5: Signal and background cross sections after kinematical cuts as functions of the selectron mass for \( e^- \gamma \rightarrow \tilde{e}^- \chi^0_1 \rightarrow e^- \chi^0_1 \chi_1 \) (full curves), \( e^- \gamma \rightarrow e^- Z^0 \rightarrow e^- \bar{\nu} \nu \) (dotted lines) and \( e^- \gamma \rightarrow W^- \nu \rightarrow e^- \bar{\nu} \nu \) (dot-dashed lines). These cross sections are shown for two different collider energies and supersymmetric scenarios:

(a) \( \sqrt{s_{ee}} = 500 \) GeV and scenario 1;
(b) \( \sqrt{s_{ee}} = 1000 \) GeV and scenario 2;

The luminosities required for extracting the supersymmetric signal from the background at a 3\( \sigma \) confidence level are shown on the right vertical axis.

Fig. 6: Contours in the \((\mu, M_2)\)-plane showing the minimum required integrated luminosities for obtaining a 3\( \sigma \) effect for \( \sqrt{s_{ee}} = 500 \) GeV, \( m_{\tilde{e}} = 250 \) GeV and \( \tan \theta_v = 2 \). The dark grey areas are excluded since there the selectron would be lighter than the lightest neutralino. The light grey region will be explored by LEP 200 [2,14].