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Radiative production of invisible charginos in photon-photon collisions

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Abstract

If in a supersymmetric model, the lightest chargino is nearly degenerate with the lightest neutralino, the former can decay into the latter alongwith a soft pion (or a lepton-neutrino pair). Near degeneracy of the chargino and neutralino masses can cause the other decay products (the pion or the lepton) to be almost invisible. Photon-photon colliders offer a possibility of clean detection of such an event through a hard photon tag.

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1 Introduction

Supersymmetry (SUSY) is widely recognized as a possibility in our quest for physics beyond the Standard Model. The particle spectrum of a particular SUSY model depends on the dynamics of the SUSY breaking. Determination of the SUSY breaking mechanism thus constitutes an integral part of any new accelerator proposal. Different sectors of the supersymmetric model can often be independently investigated in such contexts. An important component of such a model is the chargino-neutralino sector, where the physical states are formed as linear superpositions of the charged (or neutral, as the case may be) gauginos and Higgsinos respectively. A clear and unambiguous observation of this sector is important, as this will not only tell us about the SUSY breaking parameters lending themselves as gaugino masses, but also reveal crucial details of the mixing process operative here, driven by quantities such as the Higgsino mass μ and the ratio of the vacuum expectation values of the two Higgs doublets present in the model.

Direct searches at the Large Electron Positron (LEP) collider have already set lower limits of about 99 GeV [1] on the mass of the lightest chargino. Charginos and neutralinos with higher masses can be explored at Run II of the Tevatron or at the Large Hadron Collider (LHC), where a useful signal arises from the 'hadronically quiet' trilepton final states formed by decays of the $\chi_2^0\chi_1^\pm$ pair [2]. Here $\chi_1^\pm$ is the lightest chargino and χ_2^0, the next-to-lightest neutralino. However, this mode becomes hard to tag to whenever the leptonic decay channels of the $\chi_1^\pm$ get suppressed. This can happen, for example, in theories with anomaly mediated SUSY breaking (AMSB) [3]. Anomaly mediated models attempt to link the dynamics of SUSY breaking to models with extra compactified dimensions. The SUSY breaking sector is confined to a 3-brane separated from the one on which the standard model fields reside. SUSY breaking is conveyed to the observable sector by a super-Weyl anomaly. One remarkable feature of these models is the proportionality of the soft gaugino masses to the corresponding gauge beta-function coefficient (b_i). Since b_2 (corresponding to $SU(2)$) is smaller than b_1 (same for $U(1)_Y$), the $SU(2)$ gaugino mass M_2 is smaller than M_1. As a result, the lightest neutralino (χ_1^0), the LSP, is practically degenerate with $\chi_1^\pm$ and both are wino-like. With the mass separation being a few hundred MeVs at best, the $\chi_1^\pm$ decays dominantly into a χ_1^0 (which is stable and invisible so long as lepton and baryon numbers are conserved) and a soft pion. The lower limit on the mass of such a chargino is around 87 GeV [1]. Thus a $\chi_1^0\chi_1^\pm$ pair, as opposed to $\chi_2^0\chi_1^\pm$ whose production rate is suppressed, essentially escapes undetected. One way out here is to look for macroscopic tracks left by the chargino [4], but the success of this strategy is not guaranteed. Alternative signals for such a scenario have been proposed and studied in detail in the context of a high-energy electron-positron collider, *via* an analysis of the 'single photon plus missing energy' signals arising from $e^+e^- \longrightarrow \chi_1^+\chi_1^-\gamma$ [5–7]. In this paper, we suggest another possibility, namely, the radiative production of chargino pairs in photon-photon collision, triggered by laser back-scattering in a linear e^+e^- or e^-e^- collider.

The advantage of the photon-photon collision mode is that the production is controlled only by electromagnetic interaction. Thus the chargino production rates are independent of the mixing mechanism. Also, in an electron-positron collider the single photon signals are plagued with backgrounds from $e^+e^- \longrightarrow \nu\bar{\nu}\gamma$. A similar background is virtually nonexistent in photon-photon collisions. Since only charged particles are formed *via* $\gamma\gamma$ collisions at the tree level, the single photon visible final states do not get any contribution from χ^0 pairs, unlike in the case of e^+e^- collisions. And lastly, as we shall discuss in further detail in section 4, the suggested signals in an e^+e^- collider can be strongly affected by the sneutrino mass [6], at least in a very significant region of the parameter space. Such model dependence is completely avoided in the $\gamma\gamma$ mode.

In section 2 we point out some features of the process $\gamma\gamma \longrightarrow \chi^+\chi^-\gamma$ in the two-photon centre-of-mass frame. The more realistic case of a laser back-scattering experiment, and possible ways of eliminating backgrounds, are taken up in section 3. We summarise our numerical results and conclude in section 4.

2 $\gamma\gamma \longrightarrow \chi^+\chi^-\gamma$ for monochromatic photon beams

The production, governed solely by quantum electrodynamics, proceeds through six Feynman diagrams. To understand the signal profile, it is useful to consider the individual contributions from each of the various helicity combinations. Clearly, not all of the 32 possible amplitudes are independent, related as they are by discrete symmetries (charge conjugation and parity). With

$$\sigma_{ijkln} \equiv \sigma(\gamma_i\gamma_j \to \chi_k^+\chi_l^-\gamma_n) \tag{1}$$

where the subscripts (taking values $\pm$) refer to the respective particle helicities, we have

$$\begin{array}{lllllllll}
\sigma_{-----} & = & \sigma_{+++++} & & & & & & \\
\sigma_{----+} & = & \sigma_{++++-} & & & & & & \\
\sigma_{---+-} & = & \sigma_{+++-+} & = & \sigma_{--+--} & = & \sigma_{++-++} & & \\
\sigma_{---++} & = & \sigma_{+++--} & = & \sigma_{--+-+} & = & \sigma_{++-+-} & & \\
\sigma_{--++-} & = & \sigma_{++--+} & & & & & & \\
\sigma_{--+++} & = & \sigma_{++---} & & & & & & \\
\sigma_{-+---} & = & \sigma_{+-+++} & = & \sigma_{-++++} & = & \sigma_{+----} & & \\
\sigma_{-+--+} & = & \sigma_{+-++-} & = & \sigma_{-+++-} & = & \sigma_{+---+} & & \\
\sigma_{-+-+-} & = & \sigma_{+-+-+} & = & \sigma_{-++-+} & = & \sigma_{+--+-} & & \\
 & = & \sigma_{-++--} & = & \sigma_{+--++} & = & \sigma_{-+-++} & = & \sigma_{+-+--}.
\end{array} \tag{2}$$

The above relations are true not only for the total cross-sections but also for any partial sum, as long as the two charginos are subjected to identical phase space constraints. Thus we have chosen to show, in Fig. 1(a), the cross-sections corresponding to the nine independent helicity combinations appearing in the first column of each of the above equations.

The following observations are in order here.

- Each of the cross-sections is beset with kinematical singularities. One of them corresponds to the final state photon being a soft one. The other corresponds to mass and collinear singularities in the event of a vanishing chargino mass.

- The electromagnetic vertex is helicity preserving. As all the final states with identical polarisation states for the chargino pair must have encountered a helicity flip, the corresponding amplitudes must be proportional to at least M_χ. Thus, in the limit of vanishing chargino mass, these particular amplitudes should approach zero.

Clearly, wherever applicable, the two effects mentioned above pull the cross-sections in different directions, one enhancing it, the other suppressing. However, if we impose a minimal energy requirement on the final state particles as also demand that they be (at least slightly) away from the beam pipe and also not collinear with each other, then the kinematic singularities mentioned above are no longer present in the partial cross-section. In such an event, the helicity reversal argument, wherever applicable, would prevail and pull the cross-section down for sufficiently small chargino mass. Figure 1(a) confirms this. The cuts applied here are $2^\circ \leq \theta_\gamma, \theta_{\chi^\pm} \leq 178^\circ$ and $\theta_{\chi^\pm\gamma} > 5^\circ$, where $\theta_{\chi^\pm\gamma}$ is the angular separation of the photon with the charginos.

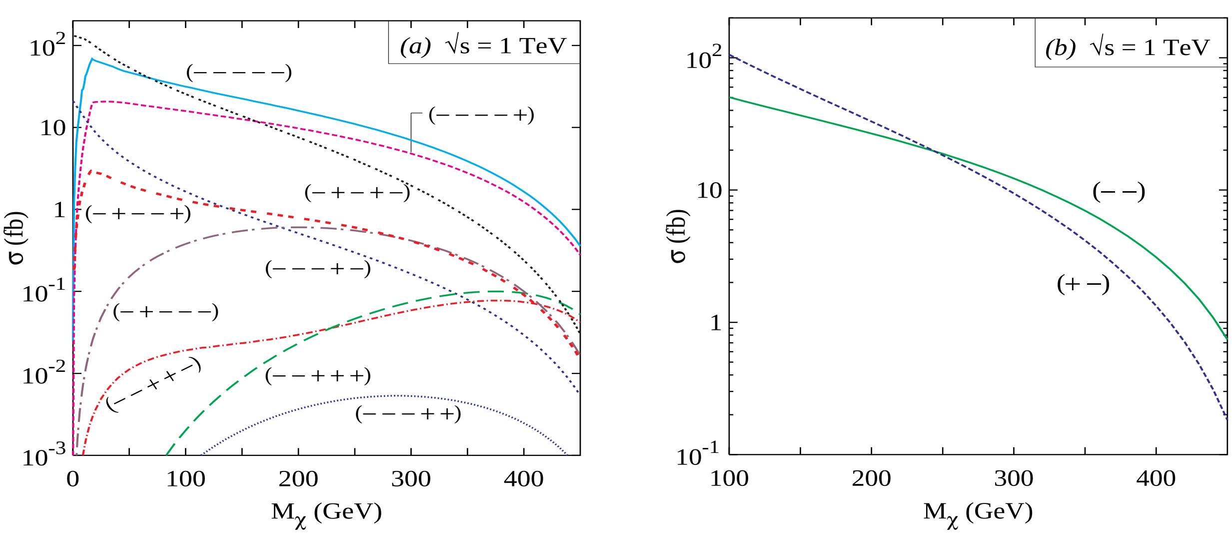

Figure 1: *Cross-section for $\gamma\gamma \to \chi^+\chi^-\gamma$ for monochromatic photon beams as a function of the chargino mass. Panel (a) shows the cross-sections for each of the nine independent helicity channels, while panel (b) gives the two independent combinations when the final state polarisations have been summed over. The cuts applied are as in the text.*

It is instructive to examine Fig. 1(a) in some detail. The fall-off at large chargino masses ($M_\chi \to \sqrt{s_{\gamma\gamma}}/2$) is but a reflection of phase space suppression. In the small M_χ regime, the rapid fall in σ_{ij--n} is a consequence of the aforementioned helicity-flip in the corresponding amplitudes.

Similarly the continued growth at small masses for σ_{-+-+-} and σ_{---+-} is symptomatic of the collinear singularity. The case of σ_{---++} might seem counterintuitive. However, this particular amplitude is essentially governed by a double-helicity flip with the consequent M_χ^4 behaviour of the cross-section for small masses. Some additional insight can be obtained if one considers the partial wave decomposition of the amplitudes into different angular momentum states. For example, the $J = 2$ initial state prefers to go to a final state where the chargino pair can be thought of to be in $J = 1$ state and hence the final state charginos have opposite polarisations (in other words, $\sigma_{-+-+-} = \sigma_{-+-++} \gg \sigma_{-+--\pm}$). Similar arguments can be constructed for the other helicity amplitudes as well.

While the discussion above helps us in understanding the dynamics of the process under consideration, the helicity of the photon in the final state is virtually immeasurable. Furthermore, since we would be primarily interested in the case of the chargino decaying into a neutralino and a pion, the chargino polarisation information is also lost.[1] Thus, we might as well sum over the final state polarisations, thereby reducing the number of independent cases to only two:

$$\begin{aligned} \sigma_{--} &\equiv \sum_{k,l,n} \sigma_{--kln} &= \sigma_{++} \\ \sigma_{-+} &\equiv \sum_{k,l,n} \sigma_{-+kln} &= \sigma_{+-} \end{aligned} \tag{3}$$

In Fig. 1(b), we explicitly demonstrate the M_χ-dependence of these cross-sections, again for the ideal case of monochromatic beams. The cuts applied in drawing the figures are $E_\gamma > 5$ GeV and $5^\circ \le \theta_\gamma \le 175^\circ$. Note that σ_{+-} dominates for small values of M_χ ceding way to σ_{--} at large chargino masses. The turnover point could also be inferred from a study of Fig. 1(a). The exact location approaches $M_\chi \sim \sqrt{s_{\gamma\gamma}}/4$ from below as the energy increases.

3 Signals from backscattered photons and background elimination

While an almost 100% polarised photon beam is possible, obtaining an intense high energy monochromatic beam is a near impossibility. In an actual experiment, one proposes to scatter laser beams off the e^+e^- or e^-e^- pair in the linear collider, and make the scattered photons collide against each other. The cross-section for a subprocess with a given centre-of-mass energy has to be folded with the energy distributions of both the photons, which are essentially Compton spectra. Not only has this spectrum a spread in the photon energy, it is not even in a pure polarisation state. The exact shape of the spectrum and the polarisation density matrix is determined by the polarisation of the initial laser beam and the electron (positron).

The basic parameters in the calculation are the energies of the primary electron (positron) and the laser beam (E_b and E_l respectively), their polarisations, and the angle of incidence (θ) between

[1] In fact, even if we were to consider fermionic decay modes such as $\chi^+ \to \chi^0 \ell^+ \nu$, the effects of chargino polarisation are essentially averaged over as long as the decay products are invisible.

the electron beam and the laser. The quantity

$$z = \frac{4E_b E_l}{m_e^2} \cos^2 \frac{\theta}{2} \tag{4}$$

determines the maximal fraction of the electron energy carried by the scattered photon [8]. An arbitrary increase in this fraction results in pair production through the interaction of the incident and scattered photons. $z = 2(1 + \sqrt{2})$ is considered to be an optimal choice [8] in this respect, and such a value of z has been adopted in this calculation. Expectedly, the said cross-sections are beset with kinematical singularities and hence are well-defined only when regulated (in other words, when phase space constraints are imposed).

Let us start by focusing on the invisible decay

$$\chi^{\pm} \to \chi^0 + \pi^{\pm} \tag{5}$$

when pion is expected to be soft. Thus, the signal is

$$\gamma\gamma \to \gamma + \text{missing energy-momentum}. \tag{6}$$

To ensure that such a photon is visible, we demand that it is emitted sufficiently away from the beam pipe and that it carries sufficiently large transverse momentum. To be specific, the events are subjected to the following cuts:

$$\begin{array}{rcccl} 10^\circ & \leq & \theta_\gamma & \leq & 170^\circ \\ \not{p}_T & = & p_T(\gamma) & \geq & 10 \text{ GeV} \\ & & E(\gamma) & < & 100 \text{ GeV} \\ & & E(\pi^{\pm}) & < & 5 \text{ GeV} \end{array} \tag{7}$$

where the last condition has been put to ensure invisibility of the pions[2]. As will be seen below, these criteria not only ensure visibility of the signals and and ward off singularities, but also serve to eliminate practically all of the SM background. The restriction on the maximal photon energy might seem puzzling at this moment, but the need to impose such a requirement would become clearer as we proceed.

Armed with the above requirements, we can convolute the cross-section with the photon spectra. As has been mentioned, the latter are determined by the polarisations of the incident laser and the electron-positron pair. Clearly, 16 such distinct combinations are possible. However, with the final state polarisations having been summed over, it can easily be seen that only 6 of these combinations are independent. Rather than consider all six, we shall focus our attention on the two particular combinations (apart from the simplest case, viz., the unpolarised one), that result in the largest cross-sections for relatively heavy charginos.

[2] Were we to consider the decay mode $\chi^+ \to \chi^0 \ell^+ \nu$, a similar restriction on the lepton energy would be imposed instead.

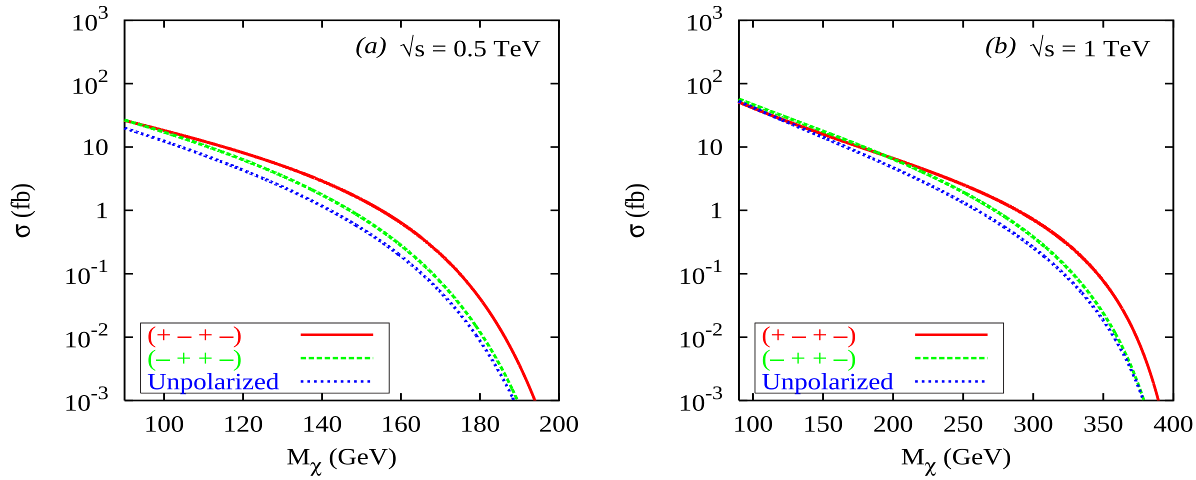

Figure 2: *Signal cross-sections as a function of chargino mass for two specific electron-electron centre-of-mass energies. The cuts of Eqn. (7) have been imposed. The parenthetic combinations reflect the initial polarisations: (electron$_1$, laser$_1$, electron$_2$, laser$_2$) with 100% polarisations for the lasers and 90% for the electrons.*

Figure 2 shows the final cross-sections after applying the above set of cuts, for two values of the electron-positron (or electron-electron) centre-of-mass energy. Polarisation efficiencies of 90% for electrons and 100% for photons have been assumed. As the graphs show, the cross-sections have a significant dependence on the beam polarisation choice for the entire range of chargino masses considered. In particular, for large M_χ, the dependence is quite sizable with the different cross-sections varying by almost an order of magnitude.

While we postpone comments on the numerical results till the next section, it is essential at this stage to enter into a discussion on the sources of backgrounds. The SM backgrounds to the signal of Eqn. (6) arise from each of the following processes:

- $\gamma\gamma \to \mathcal{B}^+\mathcal{B}^-\gamma$, where $\mathcal{B}$ is a light charged boson (π, K, ρ *etc.*);
- $\gamma\gamma \to f\bar{f}\gamma$ (where f is a lepton or a light quark) with the leptons escaping detection;
- $\gamma\gamma \to W^+W^-\gamma$ with the W decay products escaping detection;
- $\gamma\gamma \to \gamma Z$ with the Z decaying into neutrinos.

We consider each in turn.

$\underline{\gamma\gamma \to \mathcal{B}^+\mathcal{B}^-\gamma}$: Events such $\pi^+\pi^-\gamma$ production could be approximately studied assuming naive scalar electrodynamics. While this is expected to work fairly well at low momentum transfers (where the pion structure is not resolved), one would be justified in questioning the applicability for high energy processes such as the one we are concerned with. Notwithstanding this criticism, this method is expected to yield at least an order of magnitude estimate of the rates. Using this approach, we find that the admittedly large differential cross-section is heavily biased towards very energetic pions. Once the criteria of Eqn. (7) are imposed, this source of backgrounds is essentially eliminated. Analogous statements hold for other possible mesons and hadrons with the comprehensive result that this background is negligible.

$\underline{\gamma\gamma \to f\bar{f}\gamma}$: In order to fake our signal, the pair of charged particles have to travel close enough to the beam axis so as to escape detection. The exact criteria for this to happen are, of course, dependent on the (as yet unknown) detector design. However, it is reasonable to impose the following event detection criteria [5].

- Any charged fermion emitted at an angle greater than 10° with the beam axis is detected as long as its energy is at least 5 GeV.
- Instrumentation in the region between 10° and 1.2° with the beam axis would make detection possible there, provided the particle has a minimum energy of 50 GeV.
- At angles less than 1.2°, no detection would be possible.

Any particle that does not satisfy at least one of the above criteria would then escape undetected and hence contribute to the missing energy-momentum. Explicit calculation shows that no such event satisfies the event selection criteria of Eqn. (7) and hence the corresponding background is eliminated completely.

$\underline{\gamma\gamma \to W^+W^-\gamma}$: Quite analogous to the previous case, this particular process would contribute only if the decay products of the W's escape detection. It is quite obvious that this particular background would be even smaller than the $f\bar{f}\gamma$ one, and again of no concern.

$\underline{\gamma\gamma \to \gamma Z}$: Interestingly, this one-loop process proves to be the largest background. The corresponding cross-section, calculated in [9], is of the order of 30 (50) fb for $\sqrt{s_{e^+e^-}} = 500$ (1000) GeV. Folding in the invisible decay width of the Z, this would imply an "irreducible" background of 6 (10) fb. A naive comparison with Figs. 2 thus suggests a severely limited mass reach of our signal. Fortunately, event kinematics comes to our rescue. Thinking, for the moment, in terms of monochromatic photon beams, the outgoing photon, too, would be monochromatic[3] with an energy of $E_\gamma = (s_{\gamma\gamma} - m_Z^2)/2\sqrt{s_{\gamma\gamma}}$. On the other hand, the photon in the signal process has a wide energy distribution peaking, typically, at smaller E_γ. In the ideal case, then, eliminating photons lying within a relatively narrow energy band would have solved our problem. In reality though, one has to convolute this result with the photon spectrum. Two factors come into play here: (i) the cross-section $\sigma(\gamma\gamma \to \gamma Z)$ falls off very fast below $\sqrt{s_{\gamma\gamma}} \lesssim 170$ GeV; and (ii) with our favoured choice of beam polarisations (namely, opposite polarisations for incoming electron and laser), the photon beam is strongly peaked at large momentum fractions. Together, these imply that, even on

[3] The energy spread due to a possibly off-shell Z is negligible.

folding with the spectrum, the bulk of the contribution arises from very energetic photons leading to typically large values of the outgoing photon energy. The signal profile, on the other hand, is strongly peaked at small outgoing photon energies (see Fig. 3). That the spread is wider for smaller chargino masses is, of course, easily understood. However, since the cross-section for a smaller M_χ is larger, a larger fractional loss in the signal is still not too costly. More importantly, as Fig. 3 shows, a restriction of $E_\gamma < 100$ GeV (see Eqn. 7) leads to a very small loss of signal, while suppressing the γZ background to a small fraction of a femtobarn[4]. To summarise, we have established that the selection criteria of Eqn. (7) serve to make our signal almost background-free.

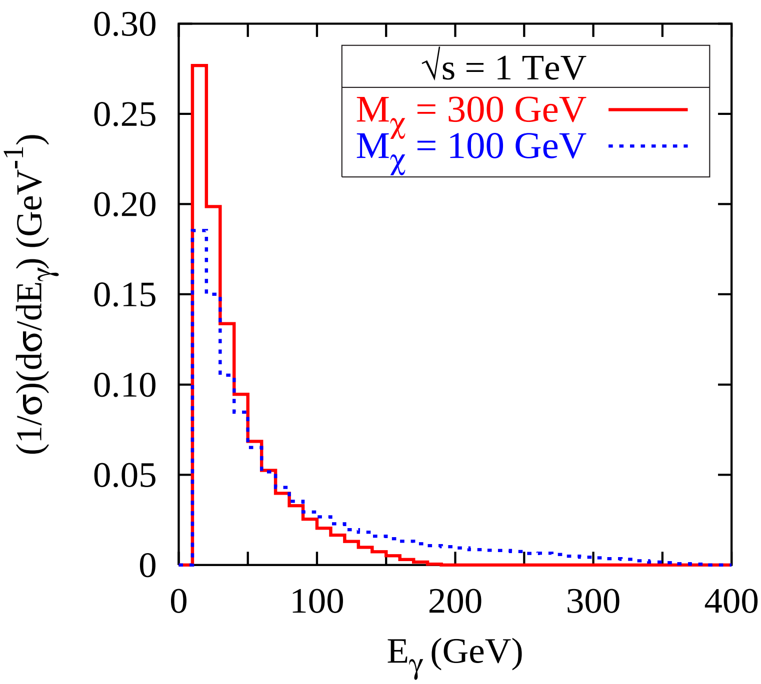

Figure 3: *The energy distribution for the outgoing photon for the signal process. Two different chargino masses are considered. The cuts of Eqn. (7) have been imposed.*

4 Numerical results and conclusions

In the absence of any background, 5 events would constitute a discovery. A perusal of Figures 2 (along with reasonable estimates for detector efficiencies) thus gives us the discovery potential for a particular choice of beam polarisation. Assuming an integrated luminosity of 100 fb^{-1} , an invisible chargino of mass upto about 165–175 GeV can thus be detected at a linear collider operating at

[4]Note, though, that our argument does not hold for γZ production starting from resolved photons. However, as Ref. [9] points out, the corresponding cross-sections by themselves are more than a magnitude smaller.

$\sqrt{s}$ = 500 GeV. For a 1 TeV collider, the limit can easily go up to 370 GeV or so[5]. This should be contrasted with the search limits that can be reached through radiative chargino production in electron-positron collision. Although it has been claimed that charginos can be probed upto the kinematic limit in the latter case, such a claim is tenable only when the sneutrino-induced diagrams do not contribute appreciably to the production rate. As has been shown in Ref. [6], the signal rate suffers a reduction by nearly an order of magnitude due to destructive interference when sneutrino masses are close to the lower limit obtained from LEP. Thus, an unambiguous exploration of the chargino-neutralino sector is not possible from e^+e^- collision data. A further consideration is that of the $\nu\bar{\nu}\gamma$ background. While the part that emanates from $Z\gamma$ background can be largely eliminated by registering events within a certain window in the photon energy, the non-resonant part needs a much more careful reconsideration in terms of photon energy and momentum [6]. The errors involved in such reconstruction, and the efficiency factor coming therein, may considerably reduce the efficacy of invisible chargino searches in this channel. Signals suggested in $\gamma\gamma$ collision can be advantageous in this respect, thanks to both the total absence of the $\nu\bar{\nu}\gamma$ backgrounds as well as the lack of interference, in the signal cross-section, from the sneutrino-induced channels.

Another matter of particular interest is the determination of the chargino mass. Unlike in the case of an e^+e^- collider, the cross-section here is a function of only M_χ and the centre-of-mass energy. As Figs. 2 show, the dependence is quite a marked one, especially for the $(- + +-)$ polarisation combination. Thus, event counting itself would allow a fairly good determination of M_χ, provided sufficient luminosity is available. Furthermore, the difference in the functional dependence of $\sigma(M_\chi)$ for different polarisation choices could be exploited for consistency checks. In addition to this, one would expect the edge of the photon energy spectrum to give an independent measurement. However, as this is true in principle, in practice this turns out to be of little use. Even for an integrated luminosity of 100 fb^{-1}, the event rates fall to such low values near the edge that the corresponding errors would be too large for this method to be useful.

So far, we have assumed that $\chi_1^\pm$ decay to $\chi_1^0\pi^\pm$ has 100% branching ratio. This, however, is a model dependent feature. The width in this channel primarily depends on the mass difference (Δm) between χ_1^0 and $\chi_1^\pm$. In anomaly mediated models, Δm cannot exceed a few hundreds of MeV. In that case, $\chi_1^0\pi^\pm$ is the overwhelmingly dominant channel in which the chargino can decay. Apart from AMSB, almost invisible charginos can occur for $\mu \gg M_{1,2}$ (with M_2 substantially smaller than M_1) or $\mu \ll M_{1,2}$. The first choice is very similar to the AMSB scenario, where the LSP (as well as the lighter chargino) is basically a wino. In the second case (Higgsino LSP), Δm can be relatively larger so that the three-body decay of a $\chi_1^\pm$ leading to $l^\pm\,\nu\,\chi_1^0$ can take place. The opening up of such a mode reduces the branching ratio of the two-body channel down to $5 - 15\%$. Invisibility of the lighter chargino in such cases is dictated by the softness of the lepton coming from the

[5]The kinematic limit for the two modes are 207 GeV and 414 GeV respectively.

[6]Note, however that right polarising the electron beam could suppress background while eliminating the aforementioned model-dependence of the signal.

three-body channel as well as the pion from the two-body channel. The invisible branching fraction of $\chi_1^\pm$ for such invisible decays can then vary from $33 - 40\%$ depending on the parameters. With such branching ratios, the signals discussed here enable us to probe a lighter chargino mass upto about 175 (340) GeV with e^+e^- centre-of-mass energy of 0.5 (1) TeV and integrated luminosity of 100 fb^{-1}. If $\Delta m < m_\pi$, then $l^\pm \nu \chi_1^0$ is the only possible decay mode of the $\chi_1^\pm$. Here the leptons will be so soft that we can expect a signal of the type (single photon + missing energy) for this case as well.

Before we conclude, it may be relevant to note that ways of probing an invisibly decaying chargino have also been suggested in the context of the LHC. The proposed methods consist in either the detailed study of superparticle cascades and the various $jets + leptons + missing\ energy$ final states [10], or the analysis of $forward\ jets\ +\ missing\ energy$ signals [11] induced by gauge boson fusion into chargino pairs. The search limit for invisible charginos in the first case, however, is again crucially dependent on the slepton (sneutrino) mass parameters. In the second method, the search limits are independent of this parameter, and can extend upto chargino masses of 300–500 GeV at various confidence levels. However, backgrounds cannot be completely removed there, and the success of the analysis depends on the precision of the background estimates. The $\gamma\gamma$ channel of invisible chargino production, being largely background-free with the suggested event selection criteria, can thus be complementary to the searches carried out at a hadron collider.

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