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Production of invisible charginos with a photon tag 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.

Keywords: Supersymmetry Phenomenology, Beyond Standard Model, Supersymmetric Standard Model.

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1. Introduction

Supersymmetry (SUSY) is widely recognized as perhaps the most promising possibility in our quest for physics beyond the Standard Model. Apart from providing a natural solution to the hierarchy problem, it also fits in very nicely with candidate theories operative at very high energies (viz. grand unification and/or superstrings). 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 supersymmetry search itself. Different sectors of a supersymmetric model can often be independently investigated in such contexts. An important component of a generic model is the chargino-neutralino sector, where the physical states are linear superpositions of the charged (or neutral, as the case may be) gauginos and Higgsinos respectively. A clear and unambiguous observation of this sector 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 lighter chargino when it decays into visible charged particles. 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]. In their most common form, 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 the super-Weyl anomaly. In addition, similar possibilities have also been suggested in the four-dimensional context [4].

One remarkable feature of these models is the proportionality of the soft gaugino masses to the corresponding gauge beta-function coefficients (b_i). Since b_2 (corresponding to SU(2)) is smaller than b_1 (corresponding to $\mathrm{U}(1)_Y$), the SU(2) gaugino mass M_2 turns out to be significantly 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] when the sneutrino mass is well above that of the chargino. In such a event, the $\chi_1^+ \chi_1^-$ or $\chi_1^0 \chi_1^\pm$ final states are likely to be as invisible as χ_1^0 pairs. (While the same criticism does not apply to the $\chi_2^0 \chi_1^\pm$ final state at a hadronic collider, the associated production cross-section is generically much smaller.) One way out here is to look for macroscopic tracks left by the chargino [5, 6], 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' signal arising from $e^+e^- \longrightarrow \chi_1^+\chi_1^-\gamma$ [7, 8, 9, 10]. In this paper, we suggest another possibility, namely, chargino pair production in association with a photon in $\gamma\gamma$ collisions 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 [8], 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

$$\sigma_{ijkln} \equiv \sigma(\gamma_i\gamma_j \to \chi_k^+\chi_l^-\gamma_n)\,, \tag{2.1}$$

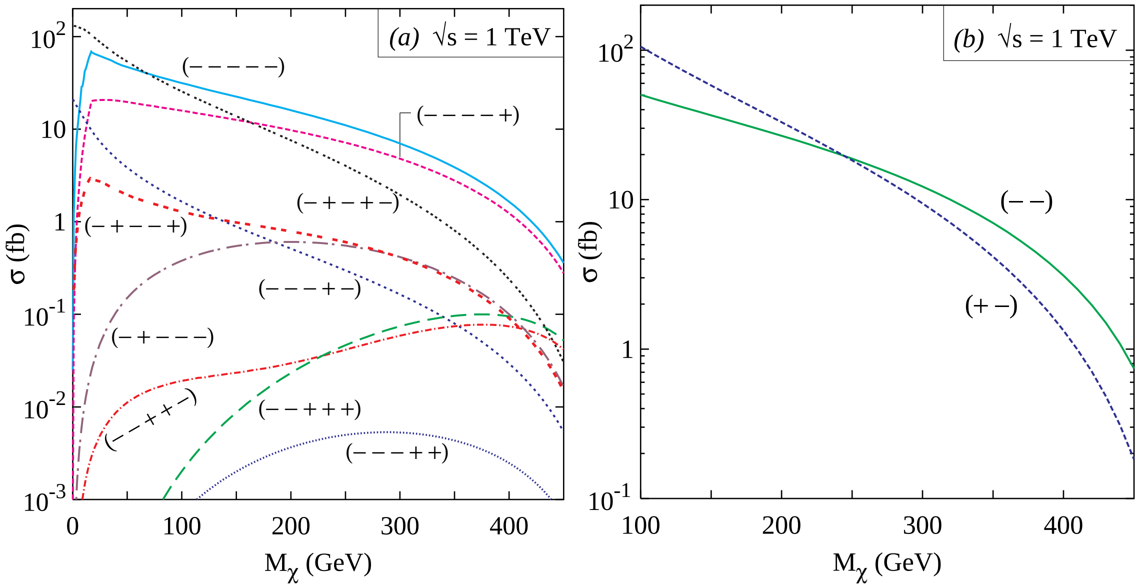

Figure 1: Cross-section for $\gamma\gamma \to \chi^+\chi^-\gamma$ for monochromatic photon beams as a function of the chargino mass. Panel (a) shows the cross-sections for each of the nine independent helicity channels. To avoid kinematical singularities, the following cuts have been imposed: $2^\circ \le \theta_\gamma, \theta_{\chi^\pm} \le 178^\circ$ and $\theta_{\chi^\pm\gamma} > 5^\circ$. Panel ($b$) gives the two independent combinations when the final state polarisations have been summed over. The cuts applied are $E_\gamma > 5\,\mathrm{GeV}$ and $5^\circ \le \theta_\gamma \le 175^\circ$.

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

$$\begin{aligned}
\sigma_{-----} &= \sigma_{+++++} \\
\sigma_{----+} &= \sigma_{++++-} \\
\sigma_{---+-} &= \sigma_{+++-+} = \sigma_{--+--} = \sigma_{++-++} \\
\sigma_{---++} &= \sigma_{+++--} = \sigma_{--+-+} = \sigma_{++-+-} \\
\sigma_{--++-} &= \sigma_{++--+} \\
\sigma_{--+++} &= \sigma_{++---} \\
\sigma_{-+---} &= \sigma_{+-+++} = \sigma_{-++++} = \sigma_{+----} \\
\sigma_{-+--+} &= \sigma_{+-++-} = \sigma_{-+++-} = \sigma_{+---+} \\
\sigma_{-+-+-} &= \sigma_{+-+-+} = \sigma_{-++-+} = \sigma_{+--+-} \\
&= \sigma_{-++--} = \sigma_{+--++} = \sigma_{-+-++} = \sigma_{+-+--}\,.
\end{aligned} \tag{2.2}$$

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. Consequently, we have chosen to show, in figure 1a, 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. (Although the limit itself is an unphysical one, it serves to highlight the structure of the matrix element. And, while, for a massive chargino, the chargino width itself acts as a natural cutoff for the photon energies, this neither invalidates the argument above nor is of sufficient numerical importance to be visible in the figure.)

- 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 1a 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 final state photon with the charginos. This cut is imposed to extrapolate the plots to the limit of vanishing chargino masses which, though unrealistic, are retained to clarify the theoretical issues discussed in the next paragraph.

It is instructive to examine figure 1a 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

[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 (and therefore immaterial) as long as the decay products are invisible.

independent cases to only two:

$$\begin{aligned}\sigma_{--} &\equiv \sum_{k,l,n} \sigma_{--kln} = \sigma_{++} \\ \sigma_{-+} &\equiv \sum_{k,l,n} \sigma_{-+kln} = \sigma_{+-}\end{aligned} \tag{2.3}$$

In figure 1b, 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 $E_\gamma > 5\,\mathrm{GeV}$ and $5^\circ \leq \theta_\gamma \leq 175^\circ$. 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 figure 1a. 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 the electron beam and the laser. The quantity

$$z \equiv \frac{4E_bE_l}{m_e^2}\cos^2\frac{\theta}{2} \tag{3.1}$$

determines the maximal fraction of the electron energy carried by the scattered photon [11]. 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 normally considered to be an optimal choice [11] in this respect, and such a value of z has been adopted in this calculation.[2] Expectedly, the said cross-sections have 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

$$\chi^\pm \to \chi^0 + \pi^\pm \tag{3.2}$$

[2]Note, though, that recent studies [12] suggest that neither the simplistic spectra discussed in ref. [11] nor $z \approx 4.83$ might be feasible options. However, such changes as argued for in these references are expected to lead to only relatively minor corrections in our numerical results. A more detailed investigation would be welcome though.

when the pion is too soft to be detected. Thus, the signal is

$$\gamma\gamma \to \gamma + \text{missing energy-momentum (including heavily ionising tracks)}\,. \tag{3.3}$$

The heavily ionising character of the track follows from the fact that the invisible chargino can have a decay length of the order of centimeters [7, 13, 14], and therefore can be exploited in eliminating backgrounds, as we shall clarify later in this section. Also, to ensure that such a photon is visible, we demand that it be 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:

$$\begin{aligned} 10^\circ \le \theta_\gamma \quad & \le 170^\circ \\ \not{p}_T = p_T(\gamma) \; & \ge 10 \text{ GeV} \\ E(\gamma) \; & < 100 \text{ GeV} \\ E(\pi^\pm) \; & < 1 \text{ GeV} \end{aligned} \tag{3.4}$$

where the last condition reflects the invisibility of the pions.[3] To put it differently, we consider any pion with energy less than 1 GeV to be essentially invisible. At first sight, it might seem that this requirement is too restrictive and that we would lose too many events. However, we have checked explicitly that, for an AMSB-type scenario, the pions are always very soft [10]. It makes virtually no difference in the ultimate event rates if we increase the last cut to 5 GeV, where it may be visible but leaves only a soft track. As will be seen below, these criteria not only ensure visibility of the signals 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.

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 eq. (3.3) arise from each of the following processes:

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

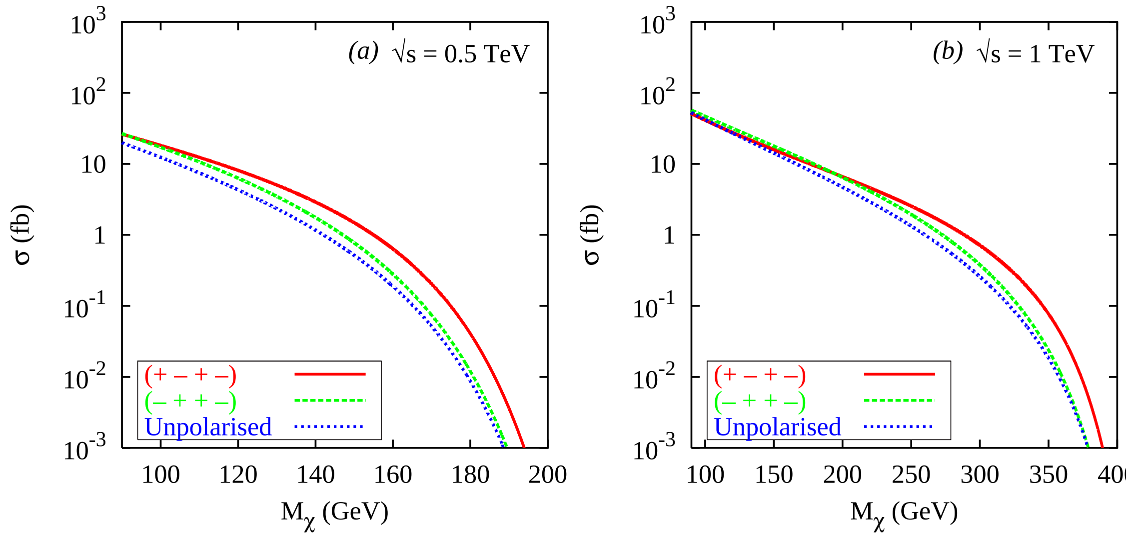

Figure 2: Signal cross-sections as a function of chargino mass for two specific electron-electron centre-of-mass energies. The cuts of eq. (3.4) have been imposed. The parenthetic combinations reflect the initial polarisations: (electron$_1$, laser$_1$, electron$_2$, laser$_2$) with 100% polarisations for the lasers and 90% for the electrons.

- $\gamma\gamma \to \mathcal{B}^+\mathcal{B}^-\gamma$, where $\mathcal{B}$ is a light charged boson (π, K, ρ *etc.*);
- $\gamma\gamma \to f\bar{f}\gamma$ (where f is e, μ or a light quark) with the fermions escaping detection;
- $\gamma\gamma \to f\bar{f}\gamma Z$ (where f is any light fermion escaping detection) and the Z decays invisibly;
- $\gamma\gamma \to W^+W^-\gamma$ with the W decay products escaping detection;
- $\gamma\gamma \to \tau^+\tau^-\gamma$ with both the taus having soft one-prong decays;
- $\gamma\gamma \to \gamma Z$ with the Z decaying into neutrinos.

We consider each in turn.

$\gamma\gamma \to \mathcal{B}^+\mathcal{B}^-\gamma$: 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 eq. (3.4) 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.

$\gamma\gamma \to f\bar{f}\gamma$: 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 [7].

- 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 eq. (3.4) and hence the corresponding background is eliminated completely.

$\gamma\gamma \to f\bar{f}\gamma Z$: Although this process occurs at a higher order in perturbation theory as compared to the preceding one, it, nonetheless, has to be considered explicitly as the photon may now recoil off the (invisibly decaying) Z thereby acquiring sufficient transverse momentum. With the visibility (or otherwise) criteria for the fermion pair being the same as for the $f\bar{f}\gamma$ background discussed above, it is clear that the case $f = e$ would constitute the bulk of this background (on account of the small electron mass). Summing over all possible fermion species, we find, on explicit matrix element calculation, this background to be $\lesssim \mathcal{O}(0.1\,\mathrm{fb})$.

$\gamma\gamma \to \tau^+\tau^-\gamma$: This becomes a source of backgrounds when the taus decay into soft one-prong channels [15]. Using the cuts specified here, the total background cross-section emanating from such τ-decays is approximately 17 fb (4 fb) for the e^+e^- centre-of-mass energy of 500 GeV (1 TeV). One way to eliminate such backgrounds is to utilize the large mass of the chargino and demand at least one heavily ionizing track (of length of about 2 cm) for the signals. As will be seen from our numerical results in the next section, the search limits we obtain are still valid, at 2σ or above, even if the above prescription fails, provided that one has an integrated luminosity of at least two hundred fb^{-1}.

$\gamma\gamma \to W^+W^-\gamma$: Quite analogous to the previous case, this particular process would contribute only if the decay products of the W's escape detection. Apart from the one-step decays into lepton-neutrino pairs, one has also to consider the added complications offered by cascade decays involving (an) intermediate tau(s). Explicit computation shows that this particular background is even smaller than the $f\bar{f}\gamma$ one, and again of no concern.

$\gamma\gamma \to \gamma Z$: Interestingly, this one-loop process proves to be the largest background so long as one is having everything but the photon invisible. The corresponding cross-section, calculated in [16], 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 figures 2 thus suggests a severely limited mass

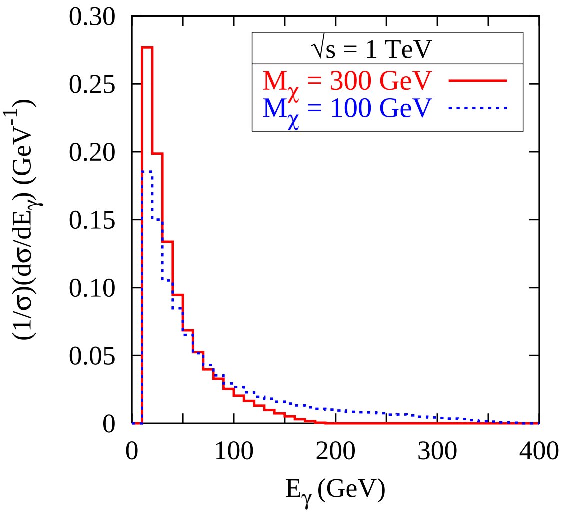

Figure 3: The energy distribution for the outgoing photon for the signal process. Two different chargino masses are considered. The cuts of eq. (3.4) have been imposed.

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[4] 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 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 figure 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 figure 3 shows, a restriction of $E_\gamma < 100$ GeV (see eq. 3.4) leads to a very small loss of signal, while suppressing the γZ background to a small fraction of a femtobarn.[5] To summarise, we have established that the selection criteria of eq. (3.4) serve to make our signal almost background-free. Of course, as has been shown by Gunion and Mrenna (ref. [8]), the accompanying soft tracks could also be used as efficient tags of the events under consideration, in which case one does not have to worry too much about the $Z\gamma$ backgrounds.

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

[5]Note, though, that our argument does not hold for γZ production starting from resolved photons. However, as ref. [16] points out, the corresponding total cross-sections by themselves are more than a magnitude smaller.

4. Numerical results and conclusions

In the absence of any background, even a single event would normally constitute a discovery. In the absence of any systematic uncertainties, this is the criterion that would be applicable to us (especially if the signal were to be defined by the presence of heavily ionizing tracks). We, however, adopt a more cautious standpoint and consider 5 events to 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 $100\,\mathrm{fb}^{-1}$, an invisible chargino of mass up to about 165–170 GeV can thus be detected at a linear collider operating at $\sqrt{s} = 500\,\mathrm{GeV}$. For a 1 TeV collider, the limit can easily go up to 370 GeV or so.[6] 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 up to 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. [8], 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, although signals for invisible charginos can always be obtained at an e^+e^- collider, the possibility as well as the extent of this interference (and, hence, the precise event rates) will depend on the mass of the sneutrino, the mixing in the chargino sector (in other words, the relative sizes of the gaugino mass parameter M_2 and the Higgsino mixing parameter μ) and, to a small extent, the mixing in the sneutrino sector. On the other hand, the $\gamma\gamma$ production channel depends only the chargino mass and, thus, is entirely independent of the model parameters. To this extent, the production mechanism being discussed here may be deemed complementary to that relevant to an e^+e^- collision.

A further consideration, for the e^+e^- mode, 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 consideration in terms of photon energy and momentum. The errors involved in such reconstruction, and the efficiency factor thereof, 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 figures 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, while this is true

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

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. Here we must admit that, for charginos decaying to hard leptons, production and decays of these particles at an e^+e^- collider have been proven to be very useful [17] for extracting the SUSY parameters beyond the chargino mass. But, for an invisibly decaying chargino photon-photon option seems to be better. And while the $\gamma\gamma$ option is not of much use in measuring the remaining parameters in the chargino–neutralino sector, this "disadvantage" needs to be seen in the light of the cleanliness in M_χ measurement and the various cross checks permitted by the use of beam polarisation. This, clearly, would serve to enhance the resolving power of the e^+e^- mode.

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 [5, 6]. Apart from AMSB, almost invisible charginos can occur for $\mu \gg M_{1,2}$ (with M_2 substantially smaller than M_1) or $\mu \ll M_{1,2}$ [7]. 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 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 up to about 170 (340) GeV with e^+e^- centre-of-mass energy of 0.5 (1) TeV and integrated luminosity of 100 fb^{-1}. But then the charginos could be looked for in e^+e^- collisions instead, albeit with somewhat reduced rates. 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.

It is perhaps necessary to remember problems that might affect our prediction under some less ideal circumstances. For example, the 'spent' electron-positron beams with which laser back-scattering takes places can masquerade as soft tracks. The design to keep these spent beams out of the interaction region can impose further constraints on the detection efficiency of the final state under consideration here. Similar complications may accrue from resolved photons as well. While a recent study [18] in the context of Higgs detection suggests that this background could be controlled, some uncertainties persist. These are open issues as of now; to keep such constraints to a minimum is one of the many challenges of detector design for a photon-photon collider.

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 [19], or the analysis of *forward jets + missing* *energy* signals [20] 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 up to 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 (and free from dependencies on parameters other than mass) with the suggested event selection criteria, can thus be complementary to the searches carried out at a hadron collider. To summarise, we have floated the basic idea of probing 'invisible' charginos (and, possibly, other invisible charged particles as well) through a single photon tag in $\gamma\gamma$ collisions. The ultimate efficacy of this probe depends on the details of beam dynamics, the design of the $\gamma\gamma$ interaction region, an accurate prediction of the back-scattered photon spectrum, and lastly, a full detector simulation of all possible backgrounds. Since the results derived here are certainly encouraging, we hope that such an effort, especially on the part of experimentalists, would be forthcoming.

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