HUNTING VIRTUAL LSPs AT LEP 200

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Abstract

Relatively light sneutrinos, which are experimentally allowed, may significantly affect the currently popular search strategies for supersymmetric particles by decaying dominantly into an invisible channel. In certain cases the second lightest neutralino may also decay invisibly leading to two extra carriers of missing energy – in addition to the lightest supersymmetric particle (LSP) \(\tilde{Z}_1\) – the virtual LSPs (VLSPs). It is shown that these VLSPs are allowed in supergravity models with common scalar and gaugino masses at the unification scale for a sizable region of parameter space and are consistent with all constraints derived so far from SUSY searches. The pair production of right handed sleptons, which can very well be the lightest charged SUSY particles in this scenario, at LEP 200 and their decay signatures are discussed. The signal survives kinematical cuts required to remove the standard model background. Charginos are also pair

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produced copiously if kinematically accessible; they also decay dom-
imantly into hadronically quiet di–lepton + $E_T$ modes leading to in-
teresting unlike sign dilepton events which are again easily separable
from the Standard Model backgrounds at LEP 200 energies.
I. Introduction

Supersymmetry (SUSY), the symmetry which relates bosons and fermions, is theoretically an elegant concept [1]. In addition it can solve the notorious fine-tuning problem which haunts the standard model (SM) [1], if the supersymmetric partners of the known particles have masses of the order of one TeV or less. The lower end of the interesting mass spectrum has already been probed at some of the presently operating accelerators like the Fermilab Tevatron or LEP–I at CERN. Planned accelerators like LEP–II or the LHC at CERN can probe even higher mass scales. The search for supersymmetry (SUSY) is therefore a high priority programme.

In most searches for SUSY particles it is assumed that there is a single, stable, weakly interacting particle, the so-called lightest supersymmetric particle (LSP). This particle, if produced, easily escapes detection. It is further assumed that by virtue of a conserved quantum number (R–parity), all other superparticles eventually decay into the LSP. The LSP, therefore, carries missing transverse energy $\not{E}_T$ which is traditionally regarded as the most distinctive signature of SUSY particles.

The minimal supersymmetric extension of the standard model (MSSM) contains four new spin 1/2 neutral particles. They are the super–partners of the photon, the $Z$ boson and the two neutral Higgs bosons. Linear combinations of these four states, the four neutralinos ($\tilde{\chi}_i$, $i=1–4$), are the physical states. In the currently favoured models, the lightest among them ($\tilde{\chi}_1$) is assumed to be the LSP [1].

Recently it has been emphasised [2, 3] that there may exist SUSY particles which, though unstable, decay dominantly into invisible channels. This occurs if the sneutrinos ($\tilde{\nu}$) (the super–partners of the neutrinos), though heavier than the LSP, are lighter than the lighter chargino ($\tilde{W}_1$) or the second lightest neutralino ($\tilde{Z}_2$) and are much lighter than all other SUSY particles. As a consequence, the invisible two–body decay mode $\tilde{\nu} \to \nu \tilde{Z}_1$ opens up and completely dominates over others, being the only kinematically allowed two–body decay of the sneutrinos. The other necessary condition for this scheme to work is that the $\tilde{Z}_1$ has a substantial Zino component. This, however, is almost always the case as long as the gluino (the super–partner of the gluon) has a mass above the lower bound obtained by the SUSY searches at the Tevatron [4]. In such cases the $\tilde{Z}_2$, which also has a substantial Zino component, decays primarily through the process $\tilde{Z}_2 \to \nu \tilde{\nu}$. These particles
decaying primarily into invisible channels, hereafter called virtual LSP’s (VLSP’s), may act as additional sources of $\not{E}_T$ and can significantly affect the strategy for SUSY searches [2].

Another important consequence of relatively light sneutrinos is that $\tilde{W}_1$ decays leptonically with a branching ratio (BR) $\approx 1$. This occurs via the mode $\tilde{W}_1 \rightarrow l\bar{\nu}, l = e, \mu$ or $\tau$, which is the only kinematically allowed two-body decay. This is to be contrasted with the conventional scenario where chargino branching ratios closely follow those of the $W$ bosons, in which case a mixed (leptons + jets + $\not{E}_T$) final state offers the best signal for chargino pair production at $e^+e^-$ colliders [5].

In ref.[2] sparticle masses were treated as free phenomenological parameters, although it was commented briefly that it is not unlikely that the VLSP scenario can be accommodated in the N=1 SUGRA models with common squark and gaugino masses at the GUT scale [6]. In this work we show in detail that this indeed is the case for a reasonably large region of the SUSY parameter space, as described in section II. We also discuss in some detail the impact of this scenario on SUSY searches at LEP 200. In this model there are two viable candidates for the lightest charged SUSY particle: i) the ‘right- handed’ sleptons $\tilde{l}_R$ and ii) the lighter chargino $\tilde{W}_1$. In this scenario both chargino and slepton pair production lead to final states consisting of two oppositely charged leptons and missing $p_T$. In section III the size of the signals for both the above cases is calculated, with detailed discussion of the separation of the signal from the standard model background. Our conclusions are spelt out in section IV. Explicit expressions for the production and decay of chargino pairs, including polarization effects, are listed in the Appendix.

II. The Allowed Region of Parameter Space

We assume a minimal N=1 supergravity model [6] with a common sfermion mass $m_0$ and a common gaugino mass $m_{1/2}$ at the GUT scale $M_X$. We also assume minimal particle content. The neutralino mass matrix in the basis
$(\bar{B}, \bar{W}_3, \tilde{h}_1^0, \tilde{h}_2^0)$ is then given by, in the convention of ref.\cite{1}:

$$
\mathcal{M}^0 = \begin{pmatrix}
M_1 & 0 & -M_Z \sin \theta_W \cos \beta & M_Z \sin \theta_W \sin \beta \\
0 & M_2 & M_Z \cos \theta_W \cos \beta & -M_Z \cos \theta_W \sin \beta \\
-M_Z \sin \theta_W \cos \beta & M_Z \cos \theta_W \cos \beta & 0 & -\mu \\
M_Z \sin \theta_W \sin \beta & -M_Z \cos \theta_W \sin \beta & -\mu & 0
\end{pmatrix}.
$$

(1)

Here, $M_1$ and $M_2$ are SUSY breaking $U(1)$ and $SU(2)$ gaugino masses, $\mu$ is the supersymmetric Higgs(ino) mass and $\tan \beta \equiv \langle H_2^0 \rangle / \langle H_1^0 \rangle$ is the ratio of the vacuum expectation values (vevs). The assumed unification of gaugino masses leads to the following relation at the weak scale:

$$M_1 = \frac{5}{3} \tan^2 \theta_W M_2 \simeq 0.5 M_2,$$

(2)

where

$$M_2(M_Z) = 0.82 m_{1/2}$$

(3)

gives the connection to the GUT scale gaugino mass. The same parameters that enter eq.\cite{1} also determine the chargino masses.

The relevant slepton masses\cite{6} at the weak scale are determined by $M_2$, $m_0$ and $\tan \beta$:

$$m_{\tilde{l}_R}^2 = m_0^2 + 0.223 M_2^2 + \sin^2 \theta_W D_Z,$$

(4a)

$$m_{\tilde{l}_L}^2 = m_0^2 + 0.773 M_2^2 + (0.5 - \sin^2 \theta_W) D_Z,$$

(4b)

$$m_{\tilde{\nu}}^2 = m_0^2 + 0.773 M_2^2 - 0.5 D_Z,$$

(4c)

where

$$D_Z = M_Z^2 \frac{\tan^2 \beta - 1}{\tan^2 \beta + 1} > 0$$

(5)

for $\tan \beta > 1$. Notice that $m_{\tilde{l}_L} > m_{\tilde{l}_R}$ always.

Our free parameters are thus $m_0$, $M_2$ (which we traded for $m_{1/2}$), $\mu$ and $\tan \beta$\cite{3}. There are two sets of constraints on the allowed parameter space:

\footnote{We do not extend the assumption of degenerate scalar masses at the GUT scale into the Higgs sector, since Higgs bosons play no role in our analysis; hence we cannot use the mechanism of radiative symmetry breaking to determine the parameter $\mu$ for given SUSY breaking parameters and top mass. We note, however, that the minimal boundary condition for Higgs masses at the GUT scale, $m_{H_2}^2 = m_0^2 + \mu^2$, does allow for VLSP scenarios.}

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direct experimental constraints (primarily from LEP–I [7]), and requirements
for having VLSP’s. We took the following experimental constraints into
account [7]:

\[ m_{\tilde{l}_R} > 45 \text{ GeV}, \quad m_{\tilde{\nu}} > 40 \text{ GeV}, \quad m_{\tilde{W}_1} > 46 \text{ GeV}, \]
\[ \Gamma(Z \rightarrow \tilde{Z}_1 \tilde{Z}_1) < 12 \text{ MeV}, \quad \sum_{i,j} \Gamma(Z \rightarrow \tilde{Z}_i \tilde{Z}_j) < 0.25 \text{ MeV}, \]

(6)

where the sum does not include \((i, j) = (1, 1)\). The exact bounds on the
invisible and visible \(Z\) decay widths vary a bit with time, but this does not
change the excluded region significantly. It was noted in ref.[2] that in the
presence of VLSPs the current lower bound on the gluino mass \(m_{\tilde{g}}\) is
likely to be relaxed since the lighter chargino arising from gluino produc
tion decays primarily into (soft) leptons rather than jets. We, therefore, do not
take this bound at its face value. Assuming the gluino mass to be unifie
d with the other gaugino masses, a bound \(m_{\tilde{g}} > 120 \text{ GeV}\) would corre
cord to something like \(M_2 > 35\) to 40 GeV, depending on \(\alpha_S\); our explicit examples
discussed in sec.III respect this lower bound.

The defining property of the VLSP scenario is that both the sneutrino
and \(\tilde{Z}_2\) decay invisibly. This implies the following constraints:

\[ m_{\tilde{\nu}} < m_{\tilde{Z}_2} < m_{\tilde{l}_L}, \quad m_{\tilde{l}_R} < m_{\tilde{\nu}} \]
\[ m_{\tilde{l}_R} < m_{\tilde{W}_1} < m_{\tilde{l}_L}. \]

(7)

The second constraint in (7) is included in order to make \(\tilde{W}_1\) decay via the
two body mode discussed above. It does not involve \(m_{\tilde{l}_R}\) since \(\tilde{l}_R\) does not
couple to \(\tilde{W}_1\).

Notice that (7) implies \(m_{\tilde{\nu}} < m_{\tilde{l}_R}\); eqs.(4) then give an upper bound on \(M_2\):

\[ |M_2| \leq 1.15\sqrt{D_Z}. \]

(8)

This in turn requires \(M_2\) to be rather small, very likely yielding a gluino
within the striking range of the Tevatron. For example, with \(\tan\beta = 2\)
the upper bound on the gluino mass is about 245 (285) GeV. For each value of
\(M_2\) there is an upper bound on the lighter chargino mass; the above bounds
imply \(m_{\tilde{W}_1} \leq 95\) (102) GeV for \(\tan\beta = 2\) (10). Then (7) implies that \(m_0\)
cannot be too large, either, hence sleptons will also be light. The lightest
charged sparticle mass is then expected to be around $M_Z$. The masses are, therefore, in the region of interest for LEP–II.

Figures 1a,b,c, show the allowed region in more detail. Here we fixed $\tan\beta$, and plotted the allowed region in the $(\mu, M_2)$ plane for various values of $m_0$. The dotted curve delineates the region excluded by sparticle searches at LEP–I. The allowed region for fixed $m_0$ is not very large, although the fraction of the plane with a VLSP scenario for some value of $m_0$ is substantial, given only that $M_2$ satisfies (8). The results are summarised below.

For $\tan\beta = 2, \mu > 0$, the left boundary of the allowed region comes from the requirement $m_{\tilde{\nu}} < m_{\tilde{W}_1}$; the upper boundary comes from $m_{\tilde{\nu}} > m_{\tilde{Z}_2}$; and the lower boundary comes from $m_{\tilde{\nu}} > 40 \text{ GeV}$. Of course, the LEP searches also take a bite out of the parameter space, as indicated.

For $\tan\beta = 2, \mu < 0$: The almost vertical left boundary for small $m_0$ comes from $m_{\tilde{Z}_2} < m_{\tilde{\nu}}$; the top–right boundaries come from $m_{\tilde{\nu}} < m_{\tilde{Z}_2}$; and the lower boundaries usually come from $m_{\tilde{\nu}} > 40 \text{ GeV}$. The situation for $\tan\beta = 10$ is similar.

Almost any value of $\mu$ can accomodate VLSPs if $M_2$ and $m_0$ are chosen properly. In fig.2 we have plotted the allowed region in the $(m_0, M_2)$ plane after scanning over all $\mu$; in other words, the curves enclose the region where a value of $\mu$ can be found so that a valid VLSP scenario emerges. The upper bound on $M_2$ is basically just given by (8), i.e. this bound can be saturated. However, for a given $m_0$, there’s also a lower bound on $M_2$: If $m_0$ is small one needs a sizable $M_2$ to get the sneutrino to be sufficiently heavy (the negative $D$–term has to be compensated). For larger $m_0$, one needs $M_2$ sufficiently large so that $m_{\tilde{W}_1}, m_{\tilde{Z}_2} > m_{\tilde{\nu}}$; indeed, this gets into conflict with (8) for $m_0 > 80 \text{ GeV}$.

We also searched for the maximal and minimal allowed values, within the VLSP scenario, of certain (differences of) sparticle masses. The lightest charged sparticle turns out to be either a $\tilde{l}_R$ or a $\tilde{W}_1$, if we ignore the possibility of a light stop. The bounds on their masses depend on $\tan\beta$, due to the constraint (8). We find $49 \text{ GeV} \leq m_{\tilde{e}_R} \leq 95 \text{ GeV}$ for $\tan\beta = 2$, and $59 \text{ GeV} \leq m_{\tilde{e}_R} \leq 103 \text{ GeV}$ for $\tan\beta = 10$. We have already given the upper bounds on $m_{\tilde{W}_1}$; we find that, unlike for sleptons, the experimental lower bound of 46 GeV can always be saturated for charginos. We conclude that LEP–II will only be able to probe the entire VLSP region if the centre–of–mass energy is raised substantially above the currently foreseen value of 176 GeV. The bounds on the masses of left–handed sleptons lie about 20 to 30 GeV above
those for $m_{\tilde{e}_R}$.

For small $\tan\beta$, the $\tilde{l}_R - \tilde{Z}_1$ mass difference can be quite small, down to about 10 GeV for $\tan\beta = 2$, but for $\tan\beta = 10$ it amounts to at least 23 GeV; this has immediate bearing on the spectrum of leptons in slepton pair events. In contrast, the bounds on the $\tilde{W}_1 - \tilde{\nu}$ mass difference are almost independent of $\tan\beta$; the lower bound is always zero, indicating that leptons produced in $\tilde{W}_1$ decays can be very soft, while the upper bound is around 30 GeV in the VLSP scenario. Finally, while the upper bounds on $m_{\tilde{e}_R}$ and $m_{\tilde{W}_1}$ are very similar, we find that substantial mass splittings between these states are possible; for $\tan\beta = 2$, right–handed sleptons could lie more than 20 GeV below or more than 25 GeV above the light charginos, while for large $\tan\beta$, right–handed sleptons are almost always heavier, by as much as 37 GeV. In order to cover the entire parameter space one therefore has to search for both $\tilde{l}_R$ and $\tilde{W}_1$ pair production, which we discuss in the next section.

III. The Di–lepton Cross–section

a) Slepton pair production

The differential cross–section $d\sigma(e^+e^- \rightarrow \tilde{l}_R\tilde{l}_R)/dt$ is well–known [8]. Of course, the cross–sections for selectrons and smuons differ, since in the former case one has neutralino $t$–channel exchange diagrams which do not exist for smuons. For this reason the smuon pair cross–section depends only on the smuon mass, while the selectron pair cross–section also depends on neutralino masses and mixings.

This is demonstrated in figs 3, which shows the dependence of the two pair cross–sections on $M_2$ [or equivalently on the slepton mass through eq.(4a)] for $\tan\beta = 2$ and various combinations of $m_0$ and $\mu < 0$. The selectron pair cross–section is even larger for $\mu > 0$, and depends only weakly on $\tan\beta$. Both the selectron and the smuon pair cross–section depend strongly on the mass (there is a $v^3$ factor, where $v$ is the slepton’s cms velocity), but the selectron cross–section is always larger, at least at those rather high energies. Overall, the cross–section away from the threshold is in the pb region.

One characteristic feature of the VLSP scenario is that the charged sleptons can always decay into at least two types of neutralinos ($\tilde{Z}_1$ and $\tilde{Z}_2$), since $m_{\tilde{l}_R} > m_{\tilde{Z}_2}$. Expression for the corresponding partial widths can be
found in ref. [8]. The decay of $\tilde{l}_R$ into a lepton and $\tilde{Z}_2$ can in principle provide a nice test of this scenario. Unfortunately an explicit calculation reveals that the corresponding branching ratio is always $\leq 1\%$, and thus too small to be observable.

Having computed total cross-sections, we have to look at the signal for slepton pair production. Since $BR(\tilde{l}_R \to \tilde{Z}_1 + l) \simeq 1$, the existence of VLSPs has actually little effect here, although it does restrict the parameter space as discussed in the previous section. For the purpose of illustration we have chosen two points in parameter space where $\tilde{l}_R$ is indeed the lightest charged sparticle. The beam energy is chosen such that no other charged sparticles can be produced.

The first choice is $m_0 = 40$ GeV, $M_2 = -\mu = 70$ GeV, $\tan\beta = 2$, which gives $m_{\tilde{l}_R} = 62$ GeV, $m_{\tilde{Z}_1} = 39.5$ GeV and $m_{\tilde{Z}_2} = 58.9$ GeV. In figs. 4 a–c we present some distributions for the final state leptons coming from $\tilde{l}_R$ decay. In these figures, the solid histogram is for selectrons, the dashed one for smuons, and the dotted one for staus, where in the latter case only leptonic $\tau$ decays have been included. We have applied some cuts to get rid of $\gamma\gamma$ and $\tau^+\tau^-$ backgrounds: we require each final state lepton to have at least 2 GeV transverse momentum, and require the missing transverse momentum to exceed 5 GeV; this should reduce $\gamma\gamma \to l^+l^-$ backgrounds to an insignificant level. Finally, we required the opening angle of the two leptons in the transverse plane, $\phi_{l^+l^-}$, to be less than $160^\circ$; this value has been chosen such that $e^+e^- \to \tau^+\tau^-$ backgrounds are removed entirely. About 85% of selectron and smuon pair events pass those cuts, so that the signal is largely unaffected for those cases. However, only 35% of leptonic stau pair events pass the cuts, making it quite doubtful whether stau pairs will be observable at all in this channel.

Fig. 4a shows the energy distribution of the charged leptons in the lab frame. Before cuts this distribution is totally flat for selectrons and smuons, since scalars decay isotropically. The cut on the transverse opening angle tends to remove events where both leptons are emitted in the same direction as the sleptons are going, which gives the maximal energy for the leptons. Hence the distributions slope downwards a little bit. It should still be quite easy to determine the endpoints of the spectrum, however, which allows one to determine both $m_{\tilde{l}_R}$ and $m_{\tilde{Z}_1}$. Notice the little solid and dashed histograms at low energies; they come from $\tilde{l}_R \to \tilde{Z}_2$ decays, the Br for which is about
Unfortunately these events are totally swamped by stau pair events, so even with almost infinite statistics one can probably not see $\tilde{l}_R \to \tilde{Z}_2$ decays in our scenario. The reason is that the lepton spectrum for the stau pair events is quite soft. This is mostly due to the additional neutrinos, of course, which originate from the leptonic decays of the $\tau$ leptons. However, it is also important to include the fact that the $\tau$ leptons are produced with right–handed polarization, which means that the charged lepton in their decay is dominantly emitted in the direction opposite to the $\tau$ momentum, giving a soft spectrum. Given the small size of the leptonic stau pair cross–section (remember that for any one channel, say $e^+e^−$, an additional factor of $1/4$ has to be applied), stau pairs should not obscure the lower edge of the lepton spectrum from selectrons and smuons, at least.

Fig. 4b shows the missing $p_T$ spectrum. Obviously one could make this cut harder, if necessary, without losing much signal (except staus again). Finally, Fig. 4c shows the one distribution where selectron and smuon pair events differ in shape as well as normalization: The distribution in $\cos\theta_l$, which is the angle between the final and initial negative lepton. For selectrons this is peaked at small angles (large $\cos\theta_l$), due to the $t$–channel diagrams which favor the negative selectron to go in the direction of the initial electron. In contrast, smuon pair production has the typical $p$–wave angular distribution for the smuons; this is largely smeared out by smuon decay, however, giving a rather flat distribution. The distribution in the transverse opening angle $\phi_{l^+l^-}$ (not shown) is also quite flat, and has almost the same shape in $\tilde{e}_R$ and $\tilde{\mu}_R$ events.

The second set of parameters we looked at in some detail is $m_0 = 65$ GeV, $M_2 = 80$ GeV, $\mu = -105$ GeV and $\tan\beta=2$, which gives $m_{\tilde{l}_R}=82.5$ GeV, $m_{\tilde{Z}_1}=44.5$ GeV, and $m_{\tilde{Z}_2} = 82.4$ GeV. Here the $B_r(\tilde{l}_R \to \tilde{Z}_2)$ is about $10^{-9}$, so those decays happen basically never. We chose $\sqrt{s} = 180$ GeV to stay below the thresholds for other sparticle pair production; as a result the total cross–sections are quite low here. The primary interest in this study stems from the fact that we are now above the $W^+W^−$ threshold, so that leptonic $W$ decays are a potentially severe background. It however transpires that this background can be handled adequately since in $e^+e^− \to W^+W^− \to l^+l^−\nu\bar{\nu}$ events one can reconstruct the $W$ momenta, using the invariant mass and beam energy constraints. We have kept only those signal events where this reconstruction is not possible; this removes the $W^+W^−$
The cuts on $p_T$ and missing $p_T$ are increased to 3 and 6 GeV, respectively, while the cut on $\phi_{l^+l^-}$ is relaxed to $\phi_{l^+l^-} \leq 165^\circ$, in order to account for the increased beam energy. The overall efficiency for selectron and smuon pairs is still about 43%. The leptons from stau pairs are almost always too soft to possibly come from $W^+W^-$ events, so the efficiency for leptonic stau events is now 41%, almost the same as for the other sleptons.

The resulting lepton energy spectrum and angular distribution are shown in Figs. 5a,b. We see that the cut against $W^+W^-$ backgrounds distorts the lepton spectrum quite a bit; events with harder leptons are more likely to look like $W^+W^-$ events. At least for selectron pairs there should still be sufficiently many events near the upper edge of the spectrum to determine it (and hence the masses) with good precision. The missing $p_T$ spectra (not shown) actually became a bit harder due to the cuts. The angular distribution of the leptons is now quite flat even for selectron events; this is again due to the cut against $W^+W^-$ pairs, which preferrably removes events with very hard leptons. In those events the lepton and the slepton tend to go in the same direction, i.e., the lepton “remembers” the direction of the slepton. For the same reason the $\phi_{l^+l^-}$ distribution (not shown) now peaks at small values of the opening angle. The effects of the cut against the $W^+W^-$ background depend strongly on $\sqrt{s}$, $m_{\tilde{e}_R}$ and $m_{\tilde{Z}_1}$. In particular, combinations of parameters with small $\tilde{l}_R - \tilde{Z}_1$ mass difference, and hence rather soft leptons, now actually have higher efficiencies and less distorted distributions.

b) Chargino pair production

The cross-section for $\tilde{W}_1$ pair production at LEP 200 energies is typically of the order of a few pb$^4$, compared to about 16 pb for $W^+W^-$ production. The branching ratio of chargino decays in the leptonic channel in the MSSM with heavy squarks and sleptons is roughly the same as that of the $W$’s. In this case chargino search in the dilepton channel faces a severe background, and the best signal can be found in the mixed channel where one chargino decays hadronically and the other leptonically$^5$. However, in the VLSP sce-

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$^4$This analysis ignores smearing of the effective beam energy due to initial state radiation. However, not much radiation is expected as long as the beam energy is not much larger than the masses of the $W$, slepton or chargino.

$^5$However, in the VLSP sce-
nario a chargino pair decays into a pair of stable, oppositely charged leptons (e or \(\mu\)) with a large branching ratio (approximately 4/9). The corresponding branching ratio for \(W^+W^-\) pairs is only about 0.05. Thus the prospect of \(\tilde W_1\) search in the dilepton channel is more promising in this scenario. (Of course, now charginos can only produce hadrons from tau lepton decays, making the chargino signal in the mixed mode much more difficult to find.)

Since the dilepton final state in this case arises from the decay of spin 1/2 charginos, the distributions of the final state leptons depend on the polarizations of the charginos. These polarizations can be conveniently taken into account by calculating the di-lepton cross-section and related distributions using the helicity projection technique \([11]\). The calculations are tedious but straightforward and the results are given in the Appendix. The spin-averaged cross-section for chargino pair production is well-known \([9]\). We have checked that the total number of dilepton events resulting from our calculations agree well with that calculated from the cross-section of ref.\([9]\) multiplied by the appropriate branching ratio (4/9). However, polarization affects the distributions of the final state leptons and their response to kinematical cuts significantly.

As discussed in Sec.II, the lighter chargino is one of the two candidates for the lightest charged sparticle within the VLSP scenario. Chargino production can be most conveniently studied by choosing the beam energy below the threshold for the production of other charged sparticles. Di-lepton final states consisting of any combination of e and \(\mu\), after removing the \(W^+W^-\) and other backgrounds by the kinematical cuts discussed in the previous subsection, then provide an unambiguous signal of chargino pair production. For beam energies above the slepton threshold, \(e-\mu\) pairs still provide a clean signature for chargino production. Such final states cannot arise from slepton pairs whose decays conserve flavour. In the following we shall present two illustrative examples of chargino signals in the VLSP model.

As our first case we have chosen \(m_0 = 67.5\) GeV, \(M_2 = 53\) GeV, \(\tan\beta = 2.0\) and \(\mu = -150\) GeV, giving \(m_{\tilde \nu} = 65.4\) GeV, \(m_{\tilde \mu_R} = 79.6\) GeV, \(m_{\tilde \mu_L} = 90.0\) GeV and \(m_{\tilde W_1} = 73.7\) GeV. In this case the dilepton cross-section arising out of direct chargino decays at the c.m. energy \(\sqrt{s} = 160\) GeV is 0.15 pb. The kinematical cuts used are the same as the ones discussed in the last section (the cut against the \(W^+W^-\) background is of course irrelevant in this case). Thus for an integrated luminosity of 500 pb\(^{-1}\) we expect about 75 events.
Since $2m_{\tilde{\nu}} > \sqrt{s}$, any combination of $e$ and $\mu$ contributes to the signal.

Since the charginos can also decay into $\tau$'s, stable leptons from the decays $\tau \rightarrow (e \text{ or } \mu) \nu \bar{\nu}$ can enhance the signal. The cross-section for this process, however, is not significant when both the charginos decay into $\tau$'s due to the small leptonic branching ratio of the $\tau$ as well as the reduced efficiency for these rather soft events. On the other hand, the contribution from events where only one of the produced charginos decays into a $\tau$ while the other directly decays into $e$ or $\mu$ is not negligible. The resulting dilepton cross-section for the choice of parameters and cuts given above is 0.02 pb.

In figure 6a we show the energy distribution of charged leptons from direct chargino decays, for the above choice of parameters (lower solid line). The end points of this spectrum are particularly interesting, since they determine both $m_{\tilde{\nu}}$ and $m_{\tilde{W}_1}$. Notice that the spectrum rises with increasing $E_l$; this is due to the polarization of the produced charginos. As discussed above, stable leptons arising from $\tau$ decays can in principle obscure the lower edge of the spectrum. The distortion of the energy distribution in the presence of a single tau decay is also plotted in the same figure (upper solid line). It is clear that the number of the latter events is not large enough to obscure completely the characteristics of the energy spectrum of the leptons arising from direct chargino decays. Therefore, one should be able to determine $m_{\tilde{W}_1}$ and $m_{\tilde{\nu}}$ from the energy distribution with reasonable accuracy.

We note in passing that in the above example the mass difference between the chargino and the sneutrino is rather small. The clean environment provided by an $e^+e^-$ collider nevertheless allows us to derive a viable signal. This is to be contrasted with chargino searches in this channel at the Tevatron, where the signal can be completely washed out if the mass difference is small [11]. This is due to the fact that much stronger $p_T$ cuts on the final state leptons are required in this case to eliminate backgrounds.

The second set of parameters we have considered in some detail is $m_0 = 40$ GeV, $M_2 = 70$ GeV, $\tan\beta = 2.0$ and $\mu = -75$ GeV. This yields the following sparticle spectrum: $m_{\tilde{\nu}} = 53.8$ GeV, $m_{\tilde{l}_R} = 61.9$ GeV, $m_{\tilde{l}_L} = 82.0$ GeV and $m_{\tilde{W}_3} = 79.9$ GeV. This case is interesting for two reasons. Firstly, the beam energy has to be above the $W^+W^-$ threshold and the cuts against the $W^+W^-$ background are necessary. Secondly, since $m_{\tilde{l}_R} < m_{\tilde{W}_1}$, only $e - \mu$ final states should be considered for the chargino signal. A calculation using the standard cuts yields a dilepton cross-section of 0.13 pb from direct decays.
at c.m. energy $\sqrt{s} = 180$ GeV. Thus for an integrated luminosity of 500 pb$^{-1}$ we expect about 65 events. The cross-section for the same final state arising due to intermediate $\tau$ decays is 0.04 pb. However, we again find that such decays cannot obscure the characteristics of the $e - \mu$ energy spectrum (see the dashed histograms in Fig. 6a).

The angular distribution of the negative lepton relative to the electron beam for the two cases is shown in figure 6b. Case 2 has higher beam energy and a lighter sneutrino; this enhances the contribution of the sneutrino $t-$channel exchange diagram, leading to a peak at $\cos \theta = 1$. This peak would be even more pronounced without our cut against the $W^+W^-$ background, which again mostly removes events with hard leptons, which dominantly go in the direction of the parent chargino. However, by comparing Fig. 6a with Fig. 5a we see that here the leptons are significantly softer than in the slepton production example we discussed above, so that the distortion due to the cut against the $W^+W^-$ background is less severe. In contrast, in case 1 the angular distribution is peaked in the backward direction. This is due to destructive interference between $s-$ and $t-$channel diagrams in the forward direction. Notice that the light chargino has a significantly larger higgsino component in case 2 than in case 1; this changes the relative importance of the various contributions to the cross-section. Finally, the effect of events where one chargino decays into a tau lepton which in turn decays into an $e$ or $\mu$ is now mostly an increase of the total rate, with only minor effects on the shape of the distribution.

We have also checked the cross-section after cuts for several representative points of the parameter space for $\tan \beta = 10$. Results at $\sqrt{s} = 150$ and 180 GeV are shown in table 1. Here $N_A$ and $N_B$ refer to the number of direct and $\tau$ mediated di-lepton events, respectively, for an integrated luminosity of 500 pb$^{-1}$. If $m_{\tilde{W}_1} > m_{\tilde{\tau}_R}$, only $e - \mu$ events are considered. We see that a viable signal can be expected for this value of $\tan \beta$ as well. We have also checked that, inspite of the $\tau$ mediated events, the end points of the lepton energy spectrum in each case are sufficiently well defined to allow a determination of $m_{\tilde{W}_1}$ and $m_{\tilde{\nu}}$. 

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IV. Conclusions

In conclusion, we reiterate that the VLSP scenario can drastically alter search strategies for SUSY particles at future colliders, due to the presence of additional carriers of $\not{E}_T$ and the enhanced leptonic branching ratio of the lighter chargino. One purpose of this paper is to show that this scenario is consistent even with the $N = 1$ SUGRA models with highly constrained mass spectra for the SUSY particles. In this scenario either the right–handed slepton or the lighter chargino turns out to be lightest charged SUSY particle.

It has been known for some time that both can be pair produced with sizable cross–sections at LEP-II, if they are not too heavy. In the VLSP scenario, both slepton and chargino production are signalled by a pair of charged leptons and missing $p_T$, which can be easily disentangled from the standard model background by the kinematical cuts discussed in the text. In particular, we suggested to remove $W^+W^−$ backgrounds by only accepting events which cannot be interpreted as coming from $W$ pairs. In principle this removes this background completely, leaving at least 40% of the signal behind. We found that the presence of a relatively light second neutralino has little effect on the production or decay of right–handed sleptons. Light sneutrinos play no role at all in the slepton signal, but they do have dramatic effects on the signal for the pair production of light charginos, which now nearly always decay into a sneutrino and a charged lepton. The cleanest signal again comes from di–lepton final states containing only $e$ and $\mu$, where the contribution from $\tilde{W}_1 \rightarrow \tau \rightarrow e, \mu$ decays is also significant. The most straightforward way to distinguish between sleptons and charginos in this scenario is that the former always give lepton pairs of the same flavour (up to a very small contribution from $\tilde{\tau}$ pair production), while chargino pairs are equally likely to result in unlike–flavoured lepton pairs. Finally, while LEP–II will almost certainly probe a significant region of parameter space where VLSPs exist, a full exploration is only possible if the centre–of–mass energy can be raised to 210 GeV or so.

Unfortunately the presence of a light $\tilde{Z}_2$ has no direct impact on the chargino signal. Within the MSSM one might be able to infer its presence from slepton production by trying to fit the parameters $M_2$, $\mu$ and $\tan\beta$ describing the neutralino sector using the measured values of the lightest neutralino mass, the total $\tilde{e}_R$ production cross–section, and the angular distribution of the electrons. A careful analysis of signal and backgrounds,
including detector simulation, is necessary to decide to what extent this is feasible.

In this paper we focussed on di-lepton final states. In the VLSP scenario we also expect enhanced signals for events with single photons and missing energy [12]. Not only \( \tilde{Z}_1 \) pairs, but also \( \tilde{Z}_1 \tilde{Z}_2 \), \( \tilde{Z}_2 \) pairs, and sneutrino pairs contribute to this signal. Unfortunately the contribution from neutralino production turns out to be quite small compared to the SM and sneutrino contributions; the single photon signal is therefore also not very sensitive to the presence of a light \( \tilde{Z}_2 \).

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Appendix

In the following we summarise the derivation of the matrix element squared for the process \( e^+ e^- \rightarrow \tilde{W}_1^+ \tilde{W}_1^- \), with subsequent decay \( \tilde{W}_1^+ (\tilde{W}_1^-) \rightarrow l^+ \bar{\nu} (l^- \nu^*) \), where \( l = e \) or \( \mu \).

The amplitude for the production of a chargino pair \( \tilde{W}_1^+ (p_3, \lambda_3) \) and \( \tilde{W}_1^- (p_4, \lambda_4) \) with momenta and helicities as indicated are given by (the subscripts in the following expressions refer to the exchanged particles in the s—or t—channel):

\[
A_{\tilde{f}}(\lambda_3, \lambda_4) \equiv A_1(\lambda_3, \lambda_4) = C_{\tilde{f}} \bar{u}_{\tilde{W}_1^-}(p_4, \lambda_4) P_L u_e(p_2, s_2) \bar{e}(p_1, s_1) P_R \nu_{\tilde{W}_1^+}(p_3, \lambda_3) \quad (A.1)
\]

\[
A_{\gamma,Z}(\lambda_3, \lambda_4) \equiv A_{2,3}(\lambda_3, \lambda_4)
\]
In the above the chargino spinors (with subscript $\tilde{W}_1$) are helicity spinors. Spin averaging will be done in the matrix element squared for the initial $e^+e^-$ pair (denoted by the subscript $e$). We have used the following abbreviations:

\begin{align}
C_{\bar{\nu}} &= -ig^2|V_{11}|^2 \quad (A.3) \\
C_Z &= \frac{ig^2}{c_w(s - m^2_Z)} \quad (A.4) \\
C_\gamma &= \frac{ie^2}{s} \quad (A.5) \\
\Gamma^\mu_{\gamma} &= \Gamma^\mu_{\gamma}' = \gamma^\mu \quad (A.6) \\
\Gamma^\mu_{Z} &= \gamma^\mu(O^L_{11}P_L + O^R_{11}P_R) \quad (A.7) \\
\Gamma^\mu_{Z}' &= \gamma^\mu(c_\nu + c_\alpha\gamma_5) \quad (A.8)
\end{align}

where $g$ is the coupling constant of $SU(2)_L$, $s_w(c_w)$ is the sine (cosine) of the weak mixing angle, $s = (p_1 + p_2)^2$, $t = (p_1 - p_3)^2$, $c_\nu = -0.25 + s_w^2$, and $c_\alpha = 0.25$. The factors $|V_{11}|$ and $O^L_{11}R^L$ enter in the $\tilde{W}_1 - l - \bar{\nu}$ and $\tilde{W}_1 - \tilde{W}_1 - Z$ couplings, respectively \cite{11, 9}. They are obtained from the diagonalisation of the 2 x 2 chargino mass matrix. We have followed the notation and prescription of \cite{11}.

In this formalism two–body decay amplitudes of positive and negative charginos are given by

\begin{align}
D_+(\lambda_3) &= -igV^*_{11}\bar{u}_l(p_5, s_5)P_Lu_{\tilde{W}_1}(p_3, \lambda_3) \quad (A.9) \\
D_-(\lambda_4) &= -igV_{11}\bar{v}_{\tilde{W}_1}(p_4, \lambda_4)P_Rv_l(p_7, s_7) \quad (A.10)
\end{align}

In the matrix element squared the final state lepton (denoted by the subscript $l$) spins will be summed.

The full matrix element squared for a particular helicity configuration of the charginos is of the form [after averaging (summing) over the initial (final) lepton spins]:

\begin{align}
M_{ij}(\lambda_3, \lambda'_3; \lambda_4, \lambda'_4) &= \frac{1}{4} \sum_{s_1,s_2} \sum_{s_5,s_7} A_i(\lambda_3, \lambda_4) D_+(\lambda_3)D_-(\lambda_4) \\
&\times A_j^*(\lambda'_3, \lambda'_4) D_+^*(\lambda'_3)D_-^*(\lambda'_4) \quad \text{(A.11)}
\end{align}
We have not included the Breit–Wigner propagators of the two charginos. It is understood that after using the narrow width approximation on them the total phase space for the $2 \to 4$ process can be factorised into a product of three phase spaces: a two-body phase space for the production of the charginos and two two-body phase spaces for their decays.

Using standard manipulations we obtain the following matrix element squared for the production of chargino pairs:

$$\frac{1}{4} \sum_{{s_1, s_2}} |A_1|^2(\lambda_3, \lambda'_3; \lambda_4, \lambda'_4) = \frac{1}{4} |C_{\nu}|^2 \text{Tr} \left[ v_{\tilde{W}_1}^\nu (p_3, \lambda_3) \bar{v}_{\tilde{W}_1} (p_3, \lambda'_3) P_L \not{p}_1 \right]$$

$$\cdot \text{Tr} \left[ u_{\tilde{W}_1} (p_4, \lambda'_4) \bar{u}_{\tilde{W}_1} (p_4, \lambda_4) P_L \not{p}_2 \right]$$

$$\frac{1}{4} \sum_{{s_1, s_2}} |A_{2,3}|^2(\lambda_3, \lambda'_3; \lambda_4, \lambda'_4) = \frac{1}{4} |C_{\nu}|^2 \text{Tr} \left[ \not{p}_1 \Gamma_{\gamma, \gamma, z} \not{p}_2 \Gamma_{\gamma, \gamma, z} \right]$$

$$\cdot \text{Tr} \left[ u_{\tilde{W}_1} (p_3, \lambda'_3) \bar{u}_{\tilde{W}_1} (p_3, \lambda_3) \Gamma_{\mu, \gamma, z} \not{v}_{\tilde{W}_1} (p_4, \lambda') \right]$$

$$\cdot \text{Tr} \left[ \not{v}_{\tilde{W}_1} (p_4, \lambda'_4) \bar{u}_{\tilde{W}_1} (p_4, \lambda_4) \Gamma_{\nu, \gamma, z} \not{p}_{L} \right]$$

(A.12)

$$\frac{1}{4} \sum_{{s_1, s_2}} A_1 A_{2,3}^* (\lambda_3, \lambda'_3; \lambda_4, \lambda'_4) + CT = \frac{1}{4} C_{\gamma} C_{\gamma}^* \text{Tr} \left[ \not{v}_{\tilde{W}_1} (p_4, \lambda') \bar{v}_{\tilde{W}_1} (p_4, \lambda'_4) \Gamma_{\mu, \gamma, z} \not{u}_{\tilde{W}_1} (p_3, \lambda_3) \bar{u}_{\tilde{W}_1} (p_3, \lambda'_3) \not{p}_{R} \right]$$

$$\cdot \not{p}_1 \Gamma_{\mu, \gamma, z} \not{p}_2 \Gamma_{\gamma, \gamma, z}$$

(A.13)

where $\Gamma_{\mu, \gamma, z} = \Gamma_{\mu, \gamma} (c_a \to -c_a)$, $\Gamma_{\mu, \gamma} = \Gamma_{\mu, \gamma}$, and CT stands for the conjugate term.

$$\frac{1}{4} \sum_{{s_1, s_2}} A_3 A_{2}^* (\lambda_3, \lambda'_3; \lambda_4, \lambda'_4) + CT = \frac{1}{4} C_{\gamma} C_{\gamma}^* \text{Tr} \left[ \not{p}_1 \Gamma_{\mu, \gamma, z} \not{p}_2 \gamma_{\nu} \right]$$

$$\cdot \text{Tr} \left[ u_{\tilde{W}_1} (p_3, \lambda'_3) \bar{u}_{\tilde{W}_1} (p_3, \lambda_3) \Gamma_{\mu, \gamma, z} \not{v}_{\tilde{W}_1} (p_4, \lambda_4) \bar{v}_{\tilde{W}_1} (p_4, \lambda'_4) \gamma_{\nu} \right]$$

(A.14)

Similarly the M.E. squared for the decays are

$$\sum_{s_5} D_+ (\lambda_3) D_-^* (\lambda'_3) = g^2 |V_{11}|^2 \text{Tr} \left[ u_{\tilde{W}_1} (p_3, \lambda_3) \bar{v}_{\tilde{W}_1} (p_3, \lambda_3) P_L \not{p}_5 \right]$$

$$= g^2 |V_{11}|^2 \text{Tr} \left[ u_{\tilde{W}_1} (p_3, \lambda_3) \bar{v}_{\tilde{W}_1} (p_3, \lambda'_3) P_R \not{p}_5 \right]$$

(A.16)

$$\sum_{s_7} D_- (\lambda_4) D_-^* (\lambda'_4) = g^2 |V_{11}|^2 \text{Tr} \left[ u_{\tilde{W}_1} (p_4, \lambda_4) \bar{v}_{\tilde{W}_1} (p_4, \lambda_4) P_R \not{p}_7 \right]$$

$$= g^2 |V_{11}|^2 \text{Tr} \left[ u_{\tilde{W}_1} (p_4, \lambda_4) \bar{v}_{\tilde{W}_1} (p_4, \lambda'_4) P_L \not{p}_7 \right]$$

(A.17)
The sum over different helicity configurations can be carried out by the following strategy. It is sufficient to compute each of the above traces for the configuration $\lambda_3 = \lambda'_3 = \lambda_4 = \lambda'_4 = +$. The remaining traces can then be obtained by inspection. Using the outer products of helicity spinors given in ref. $[11]$, Appendix C, p 573, a typical trace in eqs (A12) – (A15) reduces to the form

$$T = A + S_3 \cdot X_3 + S_4 \cdot X_4 + S_3 \cdot Y(S_4)$$  \hspace{1cm} (A.18)

where $S_3, S_4$ are the covariant spin vectors of the charginos. In eq.(A.18), $A$ is a scalar, the four vectors $X_3$ and $X_4$ are independent of the covariant spin vectors and all terms in the four vector $Y$ are linear in $S_4$. Multiplying the traces by the decay terms for this helicity configuration we obtain the following M.E. squared:

$$|M|^2 = T \left[ p_5 \cdot (p_3 - mS_3) \right] \left[ p_7 \cdot (p_4 + mS_4) \right],$$  \hspace{1cm} (A.19)

where $m$ is the chargino mass.

The summation over the helicities can now be carried out using the following observation. For $\lambda_3 = +, \lambda'_3 = -$ ($\lambda_3 = -, \lambda'_3 = +$) and fixed $\lambda_4, \lambda'_4$, the result is obtained from the $S_3$ dependent terms of eqs.(A18) and (A19) by the substitution $S_3 \rightarrow C_3(C_3^*)$ where the four vector $C$ is defined in ref.$[11]$. For $\lambda_3 = \lambda'_3 = -$, the corresponding result can be obtained from eqs.(A18) and (A19) by the substitution $S_3 \rightarrow -S_3$. The summation over $\lambda_3, \lambda'_3$ can now be performed by using the identity

$$2(X \cdot S_i)(Y \cdot S_i) + (X \cdot C_i)(Y \cdot C_i^*) + (X \cdot C_i^*)(Y \cdot C_i) = (2/m^2)(X \cdot p_i)(Y \cdot p_i) - 2(X \cdot Y)$$  \hspace{1cm} (A.20)

where $X$ and $Y$ are two arbitrary four vectors and $p_i$ refers to the momentum of the chargino. After this helicity summation we obtain from eq.(A19):

$$\sum_{\lambda_3, \lambda'_3} |M|^2 = \left\{ 2(A + S_4 \cdot X_4)(p_3 \cdot p_5) - 2m \left[ f(X_3) + f(Y(S_4)) \right] \right\} \left[ p_7 \cdot (p_4 + mS_4) \right],$$  \hspace{1cm} (A.21)

where

$$f(X) = \left[ (X \cdot p_3)(p_3 \cdot p_5)/m^2 \right] - X \cdot p_5.$$  \hspace{1cm} (A.22)

The summation over $\lambda_4, \lambda'_4$ can now be carried out by making substitutions
similar to the ones stated above and using eq. (A20). The final result is:

$$\sum_{\lambda_3, \lambda_3', \lambda_4, \lambda_4'} |M|^2 = 4 \left[ A \ p_{35} \ p_{47} - m f(X_3) \ p_{47} - p_{35} \ g(D_3) + m^2 g(D_5) + p_{35} m \ g(X_4) \right]$$

where

$$g(X) = \frac{(X \cdot p_4) p_{47}}{m^2} - X \cdot p_7;$$

$$p_{ij} = p_i \cdot p_j,$$

and the four vector $D_i$ is defined by

$$p_i \cdot Y(S_4) = S_4 \cdot D_i$$

It is therefore sufficient to calculate $A$, $X_3$, $X_4$ and $Y(S_4)$ for each of the terms in eqs. (A12 - 15). The computation of the trace and the following simplifications can be easily done by using MATHEMATICA.

For the $|A_1|^2$ term (A12) the calculation is simple and one obtains directly from eqs. (A12), (A16) and (A17) the following matrix element squared:

$$|M_1|^2 = \frac{1}{4} \left| C_{\nu} \right|^2 (4p_{24} \ p_{47} - 2m^2 \ p_{27}) (4p_{13} \ p_{35} - 2m^2 \ p_{15})$$

The relevant terms for the $s$–channel $Z$ exchange diagram are:

$$A = 4O_L^2 \left[ (c_a + c_v)^2 p_{14} p_{23} + (c_a - c_v)^2 p_{13} p_{24} \right] + 4O_R^2 \left[ (c_a + c_v)^2 p_{13} p_{24} + (c_a - c_v)^2 p_{14} p_{23} \right] + 8m^2 O_L O_R (c_a^2 + c_v^2) p_{12}$$

$$X_3 = 4m \left\{ -O_L^2 \left[ (c_a + c_v)^2 p_{14} p_2 + (c_a - c_v)^2 p_{13} p_{24} \right] + O_R^2 \left[ (c_a + c_v)^2 p_{14} p_{24} + (c_a - c_v)^2 p_{13} p_2 \right] \right\}$$

$$+ 16m O_L O_R c_a c_v (p_1 p_{23} - p_2 p_{13})$$

$$X_4 = 4m \left\{ O_L^2 \left[ (c_a + c_v)^2 p_{14} p_{13} + (c_a - c_v)^2 p_{12} p_{24} \right] - O_R^2 \left[ (c_a + c_v)^2 p_{14} p_{24} + (c_a - c_v)^2 p_{12} p_{13} \right] \right\}$$

$$+ 16m O_L O_R c_a c_v (p_1 p_{24} - p_2 p_{14})$$

$$Y(S_4) = -4O_L^2 m^2 \left[ (c_a + c_v)^2 (p_1 \cdot S_4) p_2 + (c_a - c_v)^2 (p_2 \cdot S_4) p_1 \right]$$
\[-4O_R^2m^2 \left[(c_a + c_v)^2(p_2 \cdot S_4)p_1 + (c_a - c_v)^2(p_1 \cdot S_4)p_2\right] \]
\[+ 8O_LO_R(c_a^2 + c_v^2)\left[(p_1 \cdot S_4)(p_4p_{23} - p_2p_{34}) + (p_2 \cdot S_4)(p_4p_{13} - p_1p_{34}) \right.
\[\left.+ (p_3 \cdot S_4)(p_1p_{24} + p_2p_{14} - p_4P_{12}) + S_4(p_{12}p_{34} - p_{13}p_{24} - p_{14}p_{23})\right]\]

(A.31)

For the $s-$channel $\gamma$ exchange diagram the corresponding terms are:

\[A = 8(m^2p_{12} + p_{14}p_{24} + p_{13}p_{24})\]  

(A.32)
\[Y(S_4) = 8 \left\{((p_3 \cdot S_4)p_{24} - (m^2 + p_{34})(p_2 \cdot S_4))p_1 + (p_1 \leftrightarrow p_2)\right\} \]
\[+ 8 \left[(p_1 \cdot S_4)p_{23} + p_2 \cdot S_4p_{13} - p_3 \cdot S_4p_{12}\right]p_4 + (p_{12}p_{34} - p_{14}p_{23} - p_{13}p_{24})S_4\]  

(A.33)

For the $\bar{\nu} - Z$ interference the corresponding terms are:

\[A = (c_v - c_a)(2O_Lp_{13}p_{24} + O_Rp_{12}m^2)\]  

(A.34)
\[X_3 = -(c_v - c_a)m[2O_Lp_{24}p_1 + O_R(p_{12}p_{23} - p_{2}p_{13})]\]  

(A.35)
\[X_4 = (c_v - c_a)m[2O_Lp_{13}p_2 - O_R(p_{1}p_{24} - p_{2}p_{14})]\]  

(A.36)
\[Y(S_4) = (c_v - c_a)\left\{-2O_Lm^2(p_2 \cdot S_4)p_1 \right. \]
\[\left. + O_R[-p_{2}p_{34}(p_1 \cdot S_4) - p_{1}p_{34}(p_2 \cdot S_4) + (p_3 \cdot S_4)(p_{1}p_{24} + p_{2}p_{14}) \right. \]
\[\left. + ((p_1 \cdot S_4)p_{23} + (p_2 \cdot S_4)p_{13} - (p_3 \cdot S_4)p_{12})p_4 \right. \]
\[\left. + (p_{12}p_{34} - p_{13}p_{24} - p_{14}p_{23})S_4\right\}\]  

(A.37)

For the $\bar{\nu} - \gamma$ interference the corresponding terms are

\[A = m^2p_{12} + 2p_{13}p_{24}\]  

(A.38)
\[X_3 = -mp_1(p_{23} + 2p_{24}) + mp_2p_{13}\]  

(A.39)
\[X_4 = -mp_1p_{24} + mp_2(2p_{13} + p_{14})\]  

(A.40)
\[Y(S_4) = -2m^2p_1(p_2 \cdot S_4) - p_{2}p_{34}(p_1 \cdot S_4) - p_{1}p_{34}(p_2 \cdot S_4) + p_{1}p_{24}(p_3 \cdot S_4) \]
\[p_{2}p_{14}(p_3 \cdot S_4) + p_{4}p_{23}(p_1 \cdot S_4) \]
\[+ p_{4}p_{13}(p_2 \cdot S_4) - p_{4}p_{12}(p_3 \cdot S_4) - S_4p_{14}p_{23} - S_4p_{13}p_{24} + S_4p_{12}\]  

(A.41)

For the $\gamma - Z$ interference the corresponding terms are

\[A = 4c_v(O_L + O_R)(m^2p_{12} + p_{14}p_{23} + p_{13}p_{24})\]
\[ X_3 = -4c_a m (O_L - O_R)(p_{14} p_{23} - p_{13} p_{24}) \]  
\[ X_4 = 4c_a m (O_L - O_R)(p_{14} p_{23} + p_{13} p_{24}) \]

\[ + 4c_a m (O_L + O_R) [p_1 (p_{24} + p_{23}) - p_2 (p_{14} + p_{13})] \]  
\[ Y(S_4) = 4c_v (O_L + O_R) \left\{ -m^2 [p_2 (p_1 \cdot S_4) + p_1 (p_2 \cdot S_4)] - p_{34} [p_2 (p_1 \cdot S_4) + p_1 (p_2 \cdot S_4)] \right. \]
\[ + (p_3 \cdot S_4)(p_{14} p_{23} + p_{13} p_{24}) + S_4 (p_{12} p_{34} - p_{13} p_{24} - p_{14} p_{23}) \]
\[ + p_4 [p_{23} (p_1 \cdot S_4) + p_{13} (p_2 \cdot S_4) - p_{12} (p_3 \cdot S_4)] \right\} \]
\[ - 4c_a m^2 (O_L - O_R) [p_2 (p_1 \cdot S_4) - p_1 (p_2 \cdot S_4)]. \]
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Table 1: Number of di-lepton events from chargino pair production for \( \tan\beta = 10 \), after the cuts discussed in Sec.III have been imposed, assuming an integrated luminosity of 500 pb\(^{-1}\). \( N_A \) includes only direct \( \tilde{W}_1 \rightarrow e, \mu \) decays, while \( N_B \) is the number of additional events due to \( \tilde{W}_1 \rightarrow \tau \rightarrow e, \mu \) decays. All masses are in GeV.

| \( m_0 \) | \( M_2 \) | \( \mu \) | \( \sqrt{s} \) | \( m_{\tilde{W}_1} \) | \( m_{\tilde{\nu}} \) | \( m_{\tilde{\ell}} \) | \( N_A \) | \( N_B \) |
|---|---|---|---|---|---|---|---|---|
| 20 | 90 | -120 | 150 | 68.85 | 50.93 | 63.85 | 41 | 10 |
| 20 | 90 | -120 | 180 | 68.85 | 50.93 | 63.85 | 100 | 25 |
| 40 | 80 | -120 | 150 | 62.97 | 49.8 | 69.98 | 90 | 22 |
| 60 | 80 | -200 | 150 | 74.45 | 66.93 | 83.05 | 19 | 3 |
| 60 | 80 | -200 | 180 | 74.45 | 66.93 | 83.05 | 55 | 10 |
| 20 | 88 | 140 | 150 | 72.59 | 48.15 | 63.23 | 131 | 27 |
| 40 | 80 | 200 | 150 | 62.32 | 49.8 | 69.98 | 133 | 32 |
| 60 | 80 | 400 | 150 | 73.57 | 66.93 | 83.05 | 50 | 8 |
| 60 | 80 | 400 | 180 | 73.57 | 66.93 | 83.05 | 125 | 20 |
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