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Observing Doubly Charged Higgs Bosons

in Photon-Photon Collisions

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Abstract

We discuss the possibility of observing doubly charged Higgs bosons in $\gamma\gamma$ colli-
sions. We find that one can increase the range of observability close to the kinematic
limit by a judicious choice of the polarisations of the initial e^-/e^+ beams as well
as the initial laser beam photons which are made to back-scatter from the former.
We also note that, in a large region of parameter space, the generally used lepton
number violating decay mode is dominated by the decay into a singly charged Higgs
and a real or virtual W-boson, giving rise to a large number of fermions in the fi-
nal state. This can qualitatively alter the strategy for discovering doubly-charged
scalars.

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1 INTRODUCTION

The symmetry breaking sector of the electroweak theory is yet to be confirmed experimentally. The standard one Higgs-doublet scenario is by no means the only, though certainly the simplest, option. As a matter of fact, an extended Higgs sector appears quite naturally in many options beyond the Standard Model(SM), which have been suggested to handle some of the theoretical problems posed by the SM iteself. Hence models with a richer scalar sector are constantly under investigation. A further motivating factor in these studies is the fact that the up-and-coming experiments are capable of testing many of the extensions of the Higgs sector. Of course, an immediate generalisation of the Glashow-Salam-Weinberg scheme is to introduce additional scalar doublet(s), leading to singly- charged physical scalar(s). A further step is to consider the possibility of doubly charged scalar bosons, and the phenomenology entailed by them.

Doubly charged scalars occur in many extensions of the standard model [1]. They usually come as members of triplet Higgs bosons, the most common examples of such models being those with left-right symmetry [2]. All the models containing Higgs triplets have the merit of being able to provide lepton-number violating Yukawa couplings, thereby enabling one to introduce left-handed Majorana masses for neutrinos [2,3].

It is needless to say that the phenomenology of these exotic Higgs bosons [4] is quite interesting. A number of new channels open up when they decay. Also, these scalars may be endowed with lepton number violating couplings, leading to additional physics possibilities. The signals for the singly charged Higgs bosons occuring in these models at the current and future e^+e^- colliders have been studied in some detail [5–7]. Constraints on the lepton flavour violating couplings of the doubly charged Higgs bosons possible from low energy processes as well as the e^+e^- experiments at PEP/PETRA [8,9] as well as the region that can be explored at future linear colliders [10] have also been discussed in the literature. The search possibilities for these at the e^+e^-/e^-e^- colliders [11,12] as well as at the hadronic colliders [13–15] have been investigated to some extent. In this paper we want to point out that photon-photon collisions [16], implemented through laser back-scattering in e^+e^- machines, are particularly suitable for studying doubly charged Higgs bosons. This is because the production mechanism is model-independent and depends entirely upon electromagnetic interaction. The production rate is enhanced due to the 'double' electric charge of the pair-produced scalars. The recent observation of *photon photon* scattering using backscattered lasers [17] has raised the hope that $\gamma\gamma$ colliders [18] can be made to work.

After studying the polarisation dependence of the total production cross-section, we go on to investigate the decays of such scalars in the most common forms of models. Our conclusion is that though like-sign dileptons are frequently referred to as the most striking

signals of a doubly-charged Higgs, other channels are often of greater importance. This is particularly noteworthy in view of the fact that the $\Delta L = 2$ (dilepton) channel is allowed *only in a triplet Higgs scenario.* If for some reason a doubly charged scalar scalar occurs (together with a singly charged companion) in a higher representation of $SU(2)$, then the dileptons will not arise, whereas other decays will still be possible, most important among them being $H^{++} \longrightarrow H^+ W^+$ driven by SU(2) gauge coupling. Our purpose is to draw attention to these other channels in the context of search strategies adopted for the doubly charged Higgs scalars in future experiments.

In section 2 we briefly outline the framework within which we operate in this study, and the constraints on the various parameters in this framework. Section 3 reports on the calculated cross-sections for the pair-production of doubly-charged Higgs bosons is $\gamma - \gamma$ collision. A study of the various possible decay channels is taken up in section 4. Our conclusions are summarised in section 5.

2   THE SCENARIO AND THE PARAMETERS

There are many theoretical models which necessitate the existence of doubly charged Higgs bosons. Here we assume only the most common feature of such models, namely, that the doubly-charged scalar occurs in an exotic representation of SU(2) (a triplet in most cases). The usual SU(2) doublet Higgs mixes with this exotic multiplet through the scalar potential, the mixing being parametrized by an angle θ_H. Furthermore, if the exotic multiplet is a triplet, it can have lepton-number violating couplings. We accommodate this feature also in our scenario, and parametrize the strengh of the $\Delta L = 2$ coupling by h_{ll}.

When θ_H is specified, one automatically fixes the ratio of the vacuum expectation values (VEV) of the triplet and the doublet:

$$s_H = \frac{\sqrt{2}w}{\sqrt{v^2 + 2w^2}} \tag{1}$$

where v and w are respectively the VEV's of the doublet and the triplet, and $s_H = Sin\theta_H$.

An immediate consequence of a non-zero VEV of the triplet is a contribution to the ρ-parameter, defined as $\rho = m_W^2/m_Z^2 Cos^2\theta_W$. With $x = w/v$,

$$\rho = \frac{1 + 2x^2}{1 + 4x^2} \tag{2}$$

The experimental constraint on ρ thus severely restricts the value of ω and, in turn, θ_H. It can be seen that the current data lead to the limit $s_H \leq 0.0056$ at 95% confidence level [19].

The couplings h_{ll} ($l = e, \mu\tau$) are also subject to constraints from lepton-number violating processes. These include the contribution of $\Delta L = 2$ vertices to Bhabha scattering, anomalous magnetic moment of the muon, muonium-antimuonium conversion, the decay $K^+ \longrightarrow \pi^- \mu^+ \mu^+$ etc [8–10]. But by far the strongest limits are obtained using the upper bounds on the (Majorana) masses of the three neutrinos implying that $h_{\tau\tau} \leq 1.4 \times 10^{-4}/s_H$. The couplings for the first two generations have to be down by 3 and 8 orders respectively compared to this value. Essentially, this means that the lepton-number violating couplings of the doubly charged Higgs can perhaps be phenomenologically significant for the third generation only. Also, the smaller is s_H, the larger becomes the possibility of having a large $\Delta L = 2$ interaction.

Finally, there are the masses of the doubly-and singly charged scalars (m_{++}, m_+), which we use as free parameters in the present study. The only other parameter which is potentially relevant for the our calculations is the self-coupling involved in the $H^{++}H^-H^-$ interaction. Although this coupling is determined once the masses and the triplet VEV are specified, we choose to emphasize our points by going to a region of the parameter space where the decay $H^{++} \longrightarrow H^+H^+$ is not kinematically allowed. This eliminates the need of using the self-coupling here.

Before we end this section, two comments are in order. First, the introduction of Higgs triplets (or any multiplets higher than doublets) in general implies the existence of a tree-level interaction involving the W, the Z and a charged physical scalar, a feature not present in any scenario containing only doublets. Such an interaction is proportional to the parameter s_H, and is often important in the decays of the singly-charged component H^+ [5, 6]. Side by side, the upper limit on the lepton number violating coupling is proportional to the inverse of s_H. This shows a complementarity between the signals that can be driven by either of these two types of interactions, a feature that should be kept in mind when the signals for such scenarios are being explored.

Secondly, there are some models where the constraint on s_H from ρ is avoided by postulating complex and real triplets together [20,21], and assuming a custodial symmetry protecting the equality of the VEV's. We are not assuming anything of that kind in this paper.

3 PRODUCTION IN PHOTON-PHOTON COLLISIONS

As has been already stated above, $\gamma\gamma$-collsions provide a model-independent, 'democratic' way of producing doubly charged Higgs bosons, with the advantage that the cross-section is going to be 16 times larger than a corresponding case with a singly charged scalar.

To understand the features of this production mechanism, let us first consider a hypothetical $\gamma\gamma$ machine operating at a *fixed center-of-mass energy* $\sqrt{s}$. The expression for the differential cross-section is given by

$$
\frac{d\sigma}{d\Omega} = \frac{16\alpha^2}{s}\,\beta\left[\frac{1+\xi\cdot\chi}{2} - \frac{\beta^2 s_\theta^2}{1-\beta^2 c_\theta^2}\left\{1+\xi\cdot\chi+(\xi_3+\chi_3)c_{2\phi}+(\xi_1+\chi_1)s_{2\phi}\right\}\right.
$$
$$
\left.+ \left(\frac{\beta^2 s_\theta^2}{1-\beta^2 c_\theta^2}\right)^2 (1+\xi_3 c_{2\phi}+\xi_1 s_{2\phi})\,(1+\chi_3 c_{2\phi}+\chi_1 s_{2\phi})\right] \tag{3}
$$

where $s_\psi \equiv \sin\psi$, $c_\psi \equiv \cos\psi$ and $\beta \equiv (1-4m^2/s)^{1/2}$ is the velocity of the H^{++} in the center-of-mass frame. In eq.(3) $\xi \equiv (\xi_1,\xi_2,\xi_3)$ are the Stokes parameters for one of the photons with ξ_2 denoting the circular polarization (helicity). Similarly, χ_i are the Stokes parameters for the other photon. The cross-section for unpolarized scattering are obtained from eq.(3) by substituting $\chi_i = \xi_i = 0$.

For the sake of simplicity, let us consider only circularly polarised photon beams. Obviously, the differential cross section should now possess an azimuthal symmetry and this is reflected by the expression above. In fact, all the polarisation dependence now appears through the first two terms of eq.(3) and is proportional to the product $P_{\gamma\gamma} \equiv \chi_2\xi_2$. In Fig.1(a), we show the mass-dependence of this cross-section for three different values of this product. The intersection point of the curves is given by the transcendental equation $(1+\beta)/(1-\beta) = exp(2\beta)$. Clearly for a small higgs mass, the $P_{\gamma\gamma} = -1$ mode is preferable, while close to the kinematical limit, the $P_{\gamma\gamma} = 1$ mode wins easily. This feature needs to be remembered in designing the experiment.
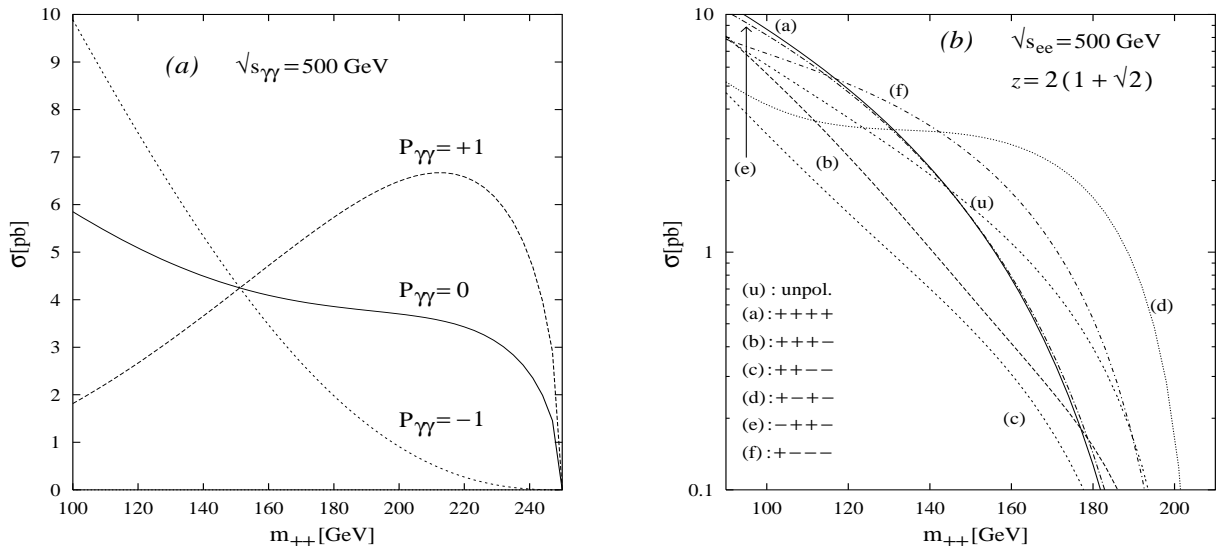

Figure 1: *Production cross-sections as a function of m_{++}.* (a) *For a monochromatic $\gamma\gamma$ machine and three choices of the product of photon helicities.* (b) *For a realistic (see eq.(4)) $\gamma\gamma$ collider operating at a fixed s_{ee}. The successive polarisation signs in the labels correspond to (e_1, laser$_1$, e_2, laser$_2$) in that order.*

In practice, however, the only known way to achive high energy $\gamma\gamma$ collisions is to affect laser back-scattering in a linear e^+e^- or e^-e^- collider [18]. The spectrum of the Compton-scattered photons is determined by the polarization of the electron and the laser and one combination of their energies and the incident angle, namely,

$$z \equiv \frac{4E_bE_\ell}{m_e^2} \cos^2 \frac{\theta_{b\ell}}{2} \ .$$

The quantity z also determines the maximum value of the fraction y of the electron energy E_e carried by the photon ($y = \omega/E_b$)

$$y_{\max} = \frac{z}{1+z}$$

It may seem that z and hence ω_{max} can be enhanced by increasing the energy of the incident laser beam. However, an arbitrary increase in that energy makes the process inefficient because of electron-positron pair production through interaction of the incident and scattered photons. An optimal choice in this respect is $z = 2(1 + \sqrt{2})$, which we have adopted in our calculation.

If the initial laser photons are circularly polarised (P_ℓ) and the electron polarisation is P_e, then the photon number density and average helicity are given by [18]

$$
\begin{aligned}
\frac{\mathrm{d}n}{\mathrm{d}y} &= \frac{2\pi\alpha^2}{m_e^2 z\sigma_C} C(y) \\
\xi_2(y) &= \frac{1}{C(y)} \left[P_e \left\{ \frac{y}{1-y} + y(2r-1)^2 \right\} - P_\ell(2r-1)\left(1 - y + \frac{1}{1-y}\right) \right] \quad (4) \\
C(y) &= \frac{1}{1-y} + (1-y) - 4r(1-r) - 2P_eP_\ell rz(2r-1)(2-y)
\end{aligned}
$$

In eq.(4), σ_C is the total Compton cross-section and $r \equiv y/z/(1-y)$.

The actual cross-sections can then be calculated by folding the cross-section of eq.(3) with the above spectrum for various polarization choices. While it is relatively straight-forward to have a fully polarized laser, it might be difficult to achieve more than 95% polarization for an electron. In the rest of the study, whenever we mention polarized beams, we shall understand it to mean

$$|P_\ell| = 1, \qquad |P_e| = 0.9 \ .$$

In Fig.1(b), we plot the production rates for a collider operating at $\sqrt{s_{ee}} = 500$ GeV. We show results both for unpolarised beams and for configurations with different combinations of electron and laser polarisation. Several comments about the cross sections are in order:

- The functional behaviour is determined by two sources: the dynamical dependence on the product $\xi_2\chi_2$ and the shape of the spectrum.

- The curves remain invariant under simultaneous reversal of all polarizations.

- The unpolarised case is the incoherent average of the other six.

- dn/dy has a high-energy bump for cases (d) and (e). For the others, it is monotonically decreasing.

- For cases (a) and (c), dn/dy are the same, but the product $\xi_2(y)\chi_2(y)$ has opposite values. The same comments hold for the pairs (b, f) and (d, e).

- At small y, $\xi_2(y)\chi_2(y)$ is positive and large for all of (a, b, d). At large y, the product is small. While it is monotonically decreasing for (a), it has one and two nodes for (b) and (d) respectively.

- $m_{++} < 120\,\text{GeV}$: $\qquad$ $(+ + ++)$ gives largest σ;
  $120\,\text{GeV} < m_{++} < 140\,\text{GeV}$: $\qquad$ $(+ - --)$ gives largest σ;
  $m_{++} > 140\,\text{GeV}$: $\qquad$ $(+ - +-)$ gives largest σ.

Thus we see that by a proper choice of polarisation combinations we can actually increase the range in m_H for a given beam energy. Since this is a reflection of the polarisation dependence of production rate for a pair of scalars, one can say by using the equivalence theorem that even the W^+W^- production cross-section should show some interesting interplay between the polarisations of the initial state e^-/e^+ and laser photons and the final state Ws.

4   DECAYS OF THE H^{++}

The signals of the doubly-charged Higgs will depend crucially on which of its decay modes are favoured in the corresponding region of the parameter space. In most studies so far, it has been assumed that the l^+l^+-channel reigns supreme and that the main signals are the very conspicuous like-sign dilepton pairs, with the invariant masses of both the pairs peaking at the same region. However, we would like to point out here that the dileptonic decay channel is by no means always the most favoured one, especially when there exists a singly charged scalar which is lighter than the doubly charged one.

Assuming that the doubly-charged scalar, H^{++}, can decay into the singly-charged one, H^+, the following decay channels are available to H^{++}:

1. $H^{++} \longrightarrow l^+l^+$

2. $H^{++} \longrightarrow W(W^*)W(W^*)$

3. $H^{++} \longrightarrow H^+ W(W^*)$

4. $H^{++} \longrightarrow H^+ H^+$

We have already mentioned that the last mode will be neglected in the study here. Of the remaining three, the level of domination of (1) clearly depends on the lepton-number violating coupling h_{ll} which in turn has a maximum possible value depending on s_H. Therefore, the lower the value of s_H is, the more insignificant we expect process (2) to be, with a corresponding increase in the share of (1).

The above observations, however, have to be balanced against the fact that process (3) is in fact driven by the SU(2) gauge coupling, and is *not* suppressed by any small model parameter. Consequently, this mode and the resulting final states are definitely expected to dominate the signals of a doubly charged Higgs over a substantial region of the parameter space. This dominaton is of course more for cases where $|m_{++} - m_+|$ is large enough for the two-body decay to be allowed. However, even when it is not so, the decay $H^{++} \longrightarrow H^+ W^*$, with the virtual W decaying fermionically, is a strong competitor against the like-sign dilepton channel, and can even dominate over it in certain regons. This fact has been demonstrated in Figs 2(a–d), where we plot the branching fractions of an H^{++} into the various modes. The mass of the singly charged Higgs has been fixed at 100 GeV in all the figures.

In all of these figures, the mode driven by gauge coupling is unmistakably dominant once the two-body decay involving a real W is allowed. Furthermore, especially when h_{ll} is on the smaller side, the three-body decay rate also exceeds that for the ll mode over a substantial region of the parameter space. The various combinations of parameters adequately demonstate that for a phenomenological study of the signals of H^{++}, it is extremely important to give due weightage to channel (3).

The final signals should depend in such a case on the various channels of H^+-decay, which can be two-fermion modes (driven by h_{ll} or s_H and suppressed by m_f/m_W), or four-fermion final states which are controlled by m_H but are not suppressed by any light fermion mass. The relative importance of these different modes have already been discussed in the literature [5–7].

It has been assumed in the entire discussion above that the H^{++} is heavier than the H^+. However, in a phenomenological analysis, one should also take into account when it is the other way around. In such a case, the hitherto dominant decay becomes kinematically disallowed, and the channel $H^{++} \longrightarrow l^+ l^+$ competes with the WW mode. In regions where the latter are kinematically suppressed, the former is more or less uniformly dominant, and the like-sign dilepton pairs constitute the most important signal. However, there is a rather interesting possibility where the dileptonic mode is the only viable one, but $h_{\tau\tau}$ is so small that the decay length is very large. (For example, if the tau neutrino
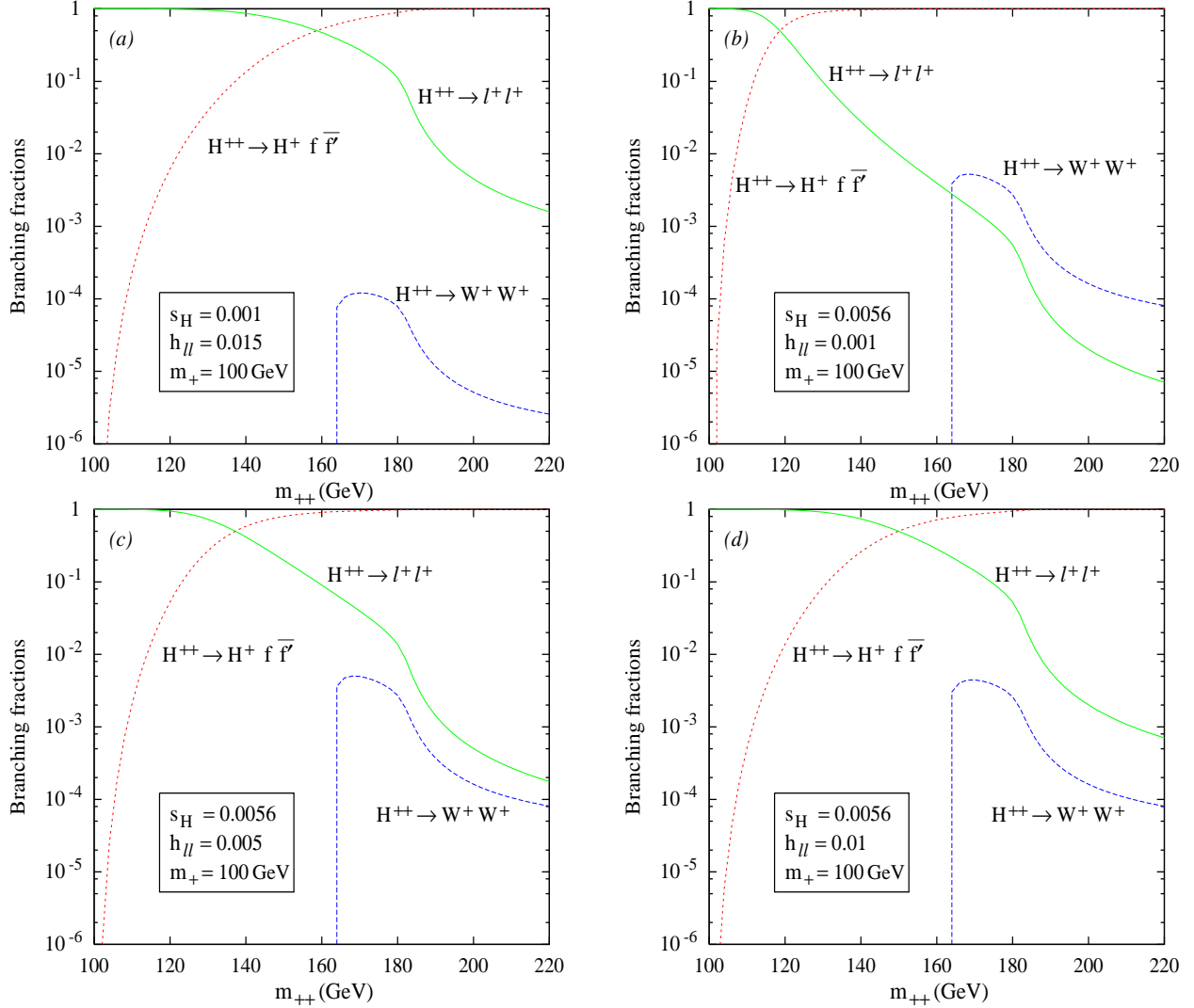
Figure 2: *The branching ratios for the different decay modes of a doubly charged Higgs, for various combinations of parameters*

has to be stable, then the bound on $h_{\tau\tau}$ given in section 2 gets tightened by some six orders of magnitude for a given s_H.) With, say, $h_{\tau\tau} = O(10^{-9})$, and $m_{++} = 150$ GeV, the decay length can be as large as 20 meters. This kind of a situation can give rise to signals in the form of thickly ionised tracks (corresponding to heavy long-lived doubly charged scalars) followed by 'delayed' events, where the ditaus may be registered as far out as the muon detectors.

5 CONCLUSIONS

We have pointed out that the $\gamma\gamma$ option in a high energy electron-positron collider can be a copious source of doubly charged scalars. Our analysis also shows that the production rates are experimentally significant for a rather large range of $m_{H^{++}}$ in a 500 GeV linear e^+e^- collider, with some interesting dependence on the polarisation choice, which can be utilised in increasing the signal strength and hence the range of m_{++} that can be probed.

As for the decay of the doubly charged scalars, we have demonstrated that the like-sign dilepton mode is not always the most dominant one, especially when a singly singly charged scalar is there for the former to decay into. In such cases, the dominant decay channel is mostly the decay into the singly charged scalar together with a real or a virtual W. This should lead to signals with a large multiplicity of fermions in the final state. It is, therefore, extremely important to consider such signals and find ways to eliminate the corresponding backgrounds, if one is interested in observing doubly charged Higgs bosons.

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