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Neutrino mass hierarchy and Majorana CP phases within the Higgs triplet model at the LHC

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Abstract: Neutrino masses may be generated by the VEV of an $SU(2)_L$ Higgs triplet. We assume that the doubly charged component of such a triplet has a mass in the range of several 100 GeV, such that it is accessible at LHC. Its decay into like-sign leptons provides a clean experimental signature, which allows for a direct test of the neutrino mass matrix. By exploring the branching ratios of this decay into leptons of various flavours, we show that within this model the type of the neutrino mass spectrum (normal, inverted or quasi-degenerate) might actually be resolved at the LHC. Furthermore, we show that within the Higgs triplet model for neutrino mass the decays of the doubly charged scalar into like-sign lepton pairs at the LHC provide a possibility to determine the Majorana CP phases of the lepton mixing matrix.

Keywords: Neutrino Physics, Beyond Standard Model.

JHEP03(2008)009

Contents

1. Introduction

Recent developments in neutrino physics demand for an extension of the Standard Model in order to give mass to the neutrinos. A popular way to achieve this goal is to introduce right-handed singlet neutrinos. An alternative, equally valid and rather economical possibility is to extend the scalar sector of the Standard Model. In addition to the Higgs doublet, scalar representations consistent with the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ gauge group and the possible fermionic bilinears are [1] a triplet, a singlet with charge +1, or a singlet with charge +2, see refs. [2–4] for corresponding models. In this work we focus on the first mentioned possibility, namely a scalar $\mathrm{SU}(2)_L$ triplet. Such a triplet arises naturally in many extensions of the Standard Model, for example in left-right symmetric models [5], or in Little Higgs theories [6, 7]. When the neutral component of the triplet acquires a vacuum expectation value (VEV), v_T, a Majorana mass term for neutrinos is generated at tree level, proportional to v_T. In order to obtain small neutrino masses this VEV and/or the corresponding Yukawa couplings have to be very small. If a very high energy scale $M \gg v = 246\,\mathrm{GeV}$ is associated to the triplet, one obtains the well-known seesaw (type-II) relation $v_T \sim v^2/M$ as explanation for the smallness of neutrino masses [8–10].

Here we consider a different scenario, assuming that the triplet states have masses not too far from the electroweak scale. The Higgs potential of the Standard Model Higgs doublet ϕ and the triplet Δ contains a term $\mu\phi\Delta\phi$, which breaks lepton number explicitly. Assuming that all other mass parameters in the potential are of the electroweak scale v, the minimisation of the potential leads to the relation for the triplet VEV $v_T \sim \mu$, see e.g. [11]. The hierarchy $\mu \ll v$ may find an explanation for example through extra dimensions [12]. A

Higgs triplet slightly below the TeV scale is the generic situation in Little Higgs theories [6], see ref. [7] for a discussion of neutrino masses in this framework. Other examples for models with TeV scale triplets responsible for neutrino masses can be found, e.g., in refs. [13 – 15]. Our phenomenological analysis does not rely on a specific model realisation, apart from the assumption that neutrino masses arise from a triplet with masses in the TeV range.

The hypothesis of such a Higgs triplet can be tested at collider experiments. In particular, if kinematically accessible, the doubly charged component of the triplet H^{++} will be produced in high energy collisions, and its decay into two equally charged leptons provides a rather spectacular signature, basically free of any Standard Model background. This process has been studied extensively in the literature (see refs. [16 – 24] for an incomplete list), and has been used to look for doubly charged scalars at LEP [25] and Tevatron [26]. These searches resulted in lower bounds for the mass of the order $M_{H^{++}} \gtrsim 130\,\text{GeV}$. Therefore, we will consider in the following masses in the range $130\,\text{GeV} \lesssim M_{H^{++}} \lesssim 1\,\text{TeV}$, above the present bound but still in reach for LHC.

If the Higgs triplet is responsible for the neutrino mass the decay rate for $H^{++} \to \ell_a^+ \ell_b^+$ is proportional to the modulus of the corresponding element of the neutrino mass matrix $|M_{ab}|^2$. This opens a phenomenologically very interesting link between neutrino and collider physics,[1] and by the observation of like-sign lepton events at LHC a direct test of the neutrino mass matrix becomes possible. In this work we assume that a doubly charged Higgs is indeed discovered at LHC, and we use the information from the decays $H^{++} \to \ell_a^+ \ell_b^+$ to learn something about neutrinos, under the hypothesis that the neutrino mass matrix is dominantly generated by the triplet VEV.

Current neutrino data leave some ambiguities for the neutrino mass spectrum. The neutrino mass states can be ordered normally or inverted, and the masses can be hierarchical or quasi-degenerate. We will show that under the above assumptions actually LHC might play a decisive role in distinguishing these possibilities. Furthermore, we show that it might be possible to determine the Majorana phases [9, 33] in the lepton mixing matrix, which in general is a very difficult task. Implications of the different possibilities of the neutrino mass spectrum for the decay of a doubly charged scalar in the Higgs triplet model have been considered previously in ref. [34], see also [22]. Building upon the results obtained there, we perform a full parameter scan including all complex phases, which — as we will see — play a crucial role for the relevant observables.

The outline of the paper is as follows. In section 2 we present the general framework, where in section 2.1 we review how the neutrino mass matrix arises in the Higgs triplet model, and in section 2.2 we discuss the signature of the model at LHC. Section 3 contains the main results of our work. After describing our analysis in section 3.1, we discuss in section 3.2 how the branching ratios of the doubly charged scalar depend on the parameters of the neutrino mass matrix. In section 3.3 we investigate the possibility to determine the type of the neutrino mass spectrum from like-sign lepton events at LHC, whereas in

[1] Such a link exists also in other classes of models, see for example [27 – 32]. However, in most cases the connection between collider signals and the neutrino mass matrix is much less direct as in the Higgs triplet model.

section 3.4 we show that within this framework indeed Majorana phases can be determined. Concluding remarks follow in section 4.

2. Framework

2.1 Neutrino masses from a Higgs triplet

If an $\mathrm{SU}(2)_L$ Higgs triplet with hypercharge $Y = 2$ is present in the theory the following renormalisable term appears in the Yukawa sector of the Lagrangian:

$$\mathcal{L}_\Delta = f_{ab}\, L_a^T C^{-1}\, i\tau_2 \Delta\, L_b + \text{h.c.}\,, \tag{2.1}$$

where the indices $a, b = e, \mu, \tau$ label flavours, L_a are the lepton doublets, C is the charge conjugation matrix, τ_2 is the Pauli matrix, Δ denotes the scalar triplet, and f_{ab} is a symmetric complex Yukawa matrix. Without loss of generality we work in the mass basis of the charged leptons. The components of the triplet are given by:

$$\Delta = \begin{pmatrix} H^+/\sqrt{2} & H^{++} \\ H^0 & -H^+/\sqrt{2} \end{pmatrix}. \tag{2.2}$$

The VEV of the neutral component $\langle H^0 \rangle \equiv v_T/\sqrt{2}$ induces a Majorana mass term for the neutrinos:

$$\frac{1}{2}\nu_{La}^T C^{-1} M_{ab}\, \nu_{Lb} + \text{h.c.} \qquad \text{with} \qquad M_{ab} = \sqrt{2}\, v_T\, f_{ab}\,. \tag{2.3}$$

We assume in the following that this is the sole source for neutrino masses (or at least the dominant contribution). As usual the neutrino mass matrix M_{ab} is diagonalised by:

$$M = U \mathrm{diag}(m_1, m_2, m_3) U^T\,. \tag{2.4}$$

For the PMNS matrix U we adopt the parametrisation

$$\begin{gathered} U = V\, \mathrm{diag}\left(e^{i\frac{\alpha_1}{2}}, e^{i\frac{\alpha_2}{2}}, e^{i\frac{\alpha_3}{2}}\right) \qquad \text{with} \\ V = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -c_{23}s_{12} - s_{13}s_{23}c_{12}e^{i\delta} & c_{23}c_{12} - s_{13}s_{23}s_{12}e^{i\delta} & s_{23}c_{13} \\ s_{23}s_{12} - s_{13}c_{23}c_{12}e^{i\delta} & -s_{23}c_{12} - s_{13}c_{23}s_{12}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \end{gathered} \tag{2.5}$$

where δ is the so-called Dirac CP violating phase which is in principle measurable in neutrino oscillation experiments, and α_i are the Majorana phases [9, 33]. Note that only relative phases $\alpha_{ij} \equiv \alpha_i - \alpha_j$ are physical, and therefore there are only two independent Majorana phases. Neutrino oscillation data determine the so-called solar and atmospheric oscillation parameters [35]:

$$\begin{aligned} \sin^2\theta_{12} &= 0.32 \pm 0.023\,, & \Delta m^2_{21} &= (7.6 \pm 0.20) \times 10^{-5}\,\mathrm{eV}^2\,, \\ \sin^2\theta_{23} &= 0.50 \pm 0.063\,, & |\Delta m^2_{31}| &= (2.4 \pm 0.15) \times 10^{-3}\,\mathrm{eV}^2\,, \end{aligned} \tag{2.6}$$

where we give 1σ errors and $\Delta m^2_{ij} \equiv m_i^2 - m_j^2$. For the mixing angle θ_{13} there is only an upper bound,

$$\sin^2\theta_{13} < 0.05 \quad \text{at } 3\sigma\,, \tag{2.7}$$

whereas nothing is known about the phases δ, α_{ij}. The ordering of the mass states is determined by the sign of Δm^2_{31}: for normal hierarchy (NH) $\Delta m^2_{31} > 0$, whereas for inverted hierarchy (IH) we have $\Delta m^2_{31} < 0$. We denote the lightest neutrino mass by m_0, hence,

$$m_0 = \begin{cases} m_1 & \text{(NH)} \\ m_3 & \text{(IH)} \end{cases} . \tag{2.8}$$

If $m_0 \gtrsim \sqrt{|\Delta m^2_{31}|} \simeq 0.05\,\text{eV}$ the neutrino mass spectrum is quasi-degenerate (QD). The most stringent bound on the absolute scale of the neutrino mass comes from cosmology, which is sensitive to the sum of the three masses. In a recent analysis [36] the upper bound $\sum_i m_i < 0.5\,\text{eV}$ at 95% CL has been obtained, which translates into $m_0 < 0.16\,\text{eV}$. Since this corresponds to the QD regime the bound is the same for NH and IH. Taking into account eq. (2.3), the constraint from cosmology applies directly to the product of triplet VEV and Yukawas:

$$v_T f_{ab} \lesssim 10^{-10}\,\text{GeV}\,. \tag{2.9}$$

2.2 Doubly charged scalars at the LHC

At the LHC the process

$$pp \to H^{++}H^{--} \to \ell^+\ell^+\,\ell^-\ell^- \tag{2.10}$$

provides a very spectacular signature, namely two like-sign lepton pairs with the same invariant mass and no missing transverse momentum, which has essentially no Standard Model background. The pair production of the doubly charged scalar occurs by the Drell-Yan process $q\overline{q} \to \gamma^*, Z^* \to H^{--}H^{++}$, with a sub-dominant contribution also from two-photon fusion $\gamma\gamma \to H^{--}H^{++}$. The cross section is not suppressed by any small quantity (such as the Yukawas or the triplet VEV) and depends only on the mass $M_{H^{++}}$, see e.g. [18, 24]. QCD corrections at next-to-leading order have been calculated [20].[2] The cross section for $H^{--}H^{++}$ pair production at the LHC ranges from 100 fb for a Higgs mass $M_{H^{++}} = 200\,\text{GeV}$ to 0.1 fb for $M_{H^{++}} = 900\,\text{GeV}$ [24]. Hence, if the doubly charged scalar is not too heavy a considerable number of them will be produced at LHC assuming an integrated luminosity of order $100\,\text{fb}^{-1}$.

The rate for the decay $H^{++} \to \ell^+_a \ell^+_b$ is given by

$$\Gamma\left(H^{++} \to \ell^+_a \ell^+_b\right) = \frac{1}{4\pi(1+\delta_{ab})}\,|f_{ab}|^2 M_{H^{++}}\,, \tag{2.11}$$

with $\delta_{ab} = 1\,(0)$ for $a = b\,(a \neq b)$. Hence, the rate is proportional to the corresponding element of the neutrino mass matrix $|M_{ab}|^2$. This observation is the basis of our analysis. Using eqs. (2.3) and (2.11) the branching ratio can be expressed as

$$\text{BR}_{ab} \equiv \text{BR}\left(H^{++} \to \ell^+_a \ell^+_b\right) \equiv \frac{\Gamma(H^{++} \to \ell^+_a \ell^+_b)}{\sum_{cd}\Gamma(H^{++} \to \ell^+_c \ell^+_d)} = \frac{2}{(1+\delta_{ab})}\,\frac{|M_{ab}|^2}{\sum_{cd}|M_{cd}|^2}\,, \tag{2.12}$$

[2] Let us note that — depending on the mass splitting between the double and single charged components of the triplet — also the channel $q'\overline{q} \to H^{\pm\pm}H^{\mp}$ may significantly contribute to the production of doubly charged scalars, see e.g. [19, 22].

JHEP03(2008)009

and from eq. (2.4) and the unitarity of U follows

$$\sum_{cd} |M_{cd}|^2 = \sum_{i=1}^{3} m_i^2 = \begin{cases} 3m_0^2 + \Delta m_{21}^2 + \Delta m_{31}^2 & \text{(NH)} \\ 3m_0^2 + \Delta m_{21}^2 + 2|\Delta m_{31}^2| & \text{(IH)} \end{cases} . \tag{2.13}$$

In addition to the lepton channel the doubly charged Higgs can in principle decay also into the following two-body final states including singly charged Higges and/or the W:

$$H^{++} \to H^+ H^+ , \quad H^{++} \to H^+ W^+ , \quad H^{++} \to W^+ W^+ . \tag{2.14}$$

The first two decay modes depend on the mass splitting within the triplet. We assume in the following that they are kinematically suppressed. The rate for the WW mode is given by

$$\Gamma\left(H^{++} \to W^+ W^+\right) \approx \frac{v_T^2 M_{H^{++}}^3}{2\pi v^4} , \tag{2.15}$$

where $v = 246\,\text{GeV}$ is the VEV of the Standard Model Higgs doublet, and we have used $M_{H^{++}} \gg M_W$, see e.g., ref. [24] for full expressions and a discussion of possibilities to observe this process at LHC. Hence, the branching ratio between $\ell^+\ell^+$ and W^+W^+ decays is controlled by the relative magnitude of the triplet Yukawas f_{ab} and the VEV v_T. The requirement $\Gamma(H^{++} \to W^+W^+) \lesssim \Gamma(H^{++} \to \ell_a^+ \ell_b^+)$, together with the constraint from eq. (2.9) implies:

$$\frac{v_T}{v} \lesssim 10^{-6} \left(\frac{100\,\text{GeV}}{M_{H^{++}}}\right)^{1/2} . \tag{2.16}$$

The triplet VEV contributes to the ρ parameter at tree level as [2] $\rho \approx 1 - 2(v_T/v)^2$. The constraint from electroweak precision data $\rho = 1.0002^{+0.0024}_{-0.0009}$ at 2σ [37] translates into $v_T/v < 0.02$, which is savely satisfied by requiring eq. (2.16).

In this model contributions to lepton flavour violating processes, $g_\mu - 2$, and in principle also to the electron electric dipole moment are expected, see e.g. [34, 38, 39] and references therein. Following refs. [34, 39], the most stringent constraint on the Yukawa couplings f_{ab} comes from $\mu \to eee$, a process which occurs at tree level via eq. (2.1). The branching ratio for this decay is given by [39]:

$$\text{BR}(\mu \to eee) = \frac{1}{4G_F^2} \frac{|f_{\text{ee}}^* f_{e\mu}|^2}{M_{H^{++}}^4} \approx 20 \left(\frac{M_{H^{++}}}{100\,\text{GeV}}\right)^{-4} |f_{\text{ee}}^* f_{e\mu}|^2 . \tag{2.17}$$

Hence, the experimental bound $\text{BR}(\mu \to eee) < 10^{-12}$ [37] constrains the combination $|f_{\text{ee}}^* f_{e\mu}| \lesssim 2 \times 10^{-7} (M_{H^{++}}/100\,\text{GeV})^2$. Assuming that all f_{ab} have roughly the same order of magnitude we obtain an estimate for the interesting range of the Yukawa couplings:

$$4 \times 10^{-7} \left(\frac{M_{H^{++}}}{100\,\text{GeV}}\right)^{1/2} \lesssim f_{ab} \lesssim 5 \times 10^{-4} \left(\frac{M_{H^{++}}}{100\,\text{GeV}}\right) , \tag{2.18}$$

where the lower bound emerges from eq. (2.16) assuming that the bound (2.9) is saturated. We see that several orders of magnitude are available for the Yukawa couplings. For f_{ab} close to the lower bound of eq. (2.18) the decay $H^{++} \to W^+W^+$ will become observable

JHEP03(2008)009

at LHC, whereas close to the upper bound a signal in future searches for lepton flavour violation is expected, where the details depend on the structure of the neutrino mass matrix [34, 39]. The interval for the Yukawas from eq. (2.18) implies a triplet VEV roughly in the keV to MeV range.

The basic assumption in our analysis is that a sufficient number of like-sign leptons is observed. If some of the decay modes of eq. (2.14) are present the number of dilepton events will be reduced according to the branching. If enough events from both types of decay (leptonic and non-leptonic) were observed in principle an order of magnitude estimate for the Yukawa couplings f_{ab} and the triplet VEV v_T might be possible [18, 22]. Here we do not consider this case and use only dilepton events, and therefore, we do not obtain any information on the overall scale of the f_{ab} in addition to eq. (2.18).

3. Numerical analysis and results

3.1 Description of the analysis

As mentioned above, we focus in our analysis on the process (2.10), which provides the clean signal of four leptons, where the like-sign lepton pairs have the same invariant mass, namely the mass of the doubly charged Higgs. Given the fact that the branching $H^{++} \to \ell_a^+ \ell_b^+$ is proportional to the neutrino mass matrix, one expects all possible flavour combinations of the four leptons to occur, including lepton flavour violating ones. In ref. [24] simple cuts have been defined for final states consisting of electrons and muons, eliminating essentially any Standard Model background.

In general tau reconstruction is experimentally more difficult because of the missing transverse energy from neutrinos. However, in the case of interest enough kinematic constraints should be available to identify also events involving taus. It turns out that the inclusion of such events significantly increases the sensitivity for neutrino parameters. Therefore, following ref. [23], we assume that events where one of the four leptons is a tau can also be reconstructed.[3] This should be possible efficiently, despite the complications involving the tau reconstruction, since the invariant mass is known from decays without tau, which can be used as kinematic constraint for events of the type $\ell^\pm \ell^\pm \ell^\mp \tau^\mp$ for $\ell = e$ or μ. Furthermore, one can adopt the assumption that the neutrinos carrying away the missing energy are aligned with the tau.

In principle it is difficult to distinguish a primary electron or muon from the ones originating from leptonic tau decays. Since here we are interested in investigating the flavour structure of the decays, leptonically decaying taus might be a "background" for the Higgs decays into electrons and muons, and vice versa. However, due to the energy carried away by the two neutrinos from the leptonic tau decay, a cut on the invariant mass of the like-sign leptons should eliminate such a confusion very efficiently. It is beyond the scope of this work to perform a detailed simulation and event reconstruction study. The above arguments suggest that our assumptions are suitable to estimate the sensitivity of the Higgs decays for neutrino parameters by the procedure outlined in the following.

[3] To be conservative we do not include events with more than one tau, since already the inclusion of events with one tau provides enough information for our purposes.

We define as our five observables the number of like-sign lepton pairs with the flavour combinations

$$x = (ee), (e\mu), (\mu\mu), (e\tau), (\mu\tau)\,. \tag{3.1}$$

Note that these five branchings contain the full information, since $\mathrm{BR}_{\tau\tau}$, which we do not use explicitly, is fixed by $\mathrm{BR}_{\tau\tau} = 1 - \sum_x \mathrm{BR}_x$. Taking into account the number of occurrences of the combinations (3.1) in four leptons where at most one tau is allowed, the number of events in each channel is obtained as:

$$\begin{aligned} N_{ab} &= 2N_{2H}\,\epsilon\,\mathrm{BR}_{ab}\sum_x \mathrm{BR}_x && \text{for} \quad (ab) = (ee), (e\mu), (\mu\mu)\,, \\ N_{ab} &= 2N_{2H}\,\epsilon\,\mathrm{BR}_{ab}(\mathrm{BR}_{\mathrm{ee}} + \mathrm{BR}_{e\mu} + \mathrm{BR}_{\mu\mu}) && \text{for} \quad (ab) = (e\tau), (\mu\tau)\,, \end{aligned} \tag{3.2}$$

where N_{2H} is the total number of doubly charged scalar pairs decaying into four leptons, and ϵ is the detection efficiency for the four lepton events. For simplicity we assume here a flavour independent efficiency. The branching ratios are given in eq. (2.12). To illustrate the sensitivity to neutrino parameters we will use $\epsilon N_{2H} = 10^3$ or $\epsilon N_{2H} = 10^2$ events. For an integrated luminosity of $100\,\mathrm{fb}^{-1}$ at LHC these event numbers will be roughly obtained for $M_{H^{++}} \simeq 350\,\mathrm{GeV}$ and $M_{H^{++}} \simeq 600\,\mathrm{GeV}$, respectively [24].

To carry out the analysis we define a χ^2 function from the observables in eq. (3.2). For given ϵN_{2H} they depend only on neutrino parameters. We consider five continuous parameters: the lightest neutrino mass m_0, s_{13}, the Dirac phase δ, and the two Majorana phases $\alpha_{12} = \alpha_1 - \alpha_2$ and $\alpha_{32} = \alpha_3 - \alpha_2$, plus the discrete parameter h = NH or IH describing the mass ordering. The remaining neutrino parameters, the two mass-squared differences and the mixing angles s_{12} and s_{23}, are fixed to their experimental best fit values given in eq. (2.6). The χ^2 is constructed as:

$$\begin{aligned} \chi^2(m_0, s_{13}, \delta, \alpha_{12}, \alpha_{32}, h) &= \sum_{xy} V_x\, S^{-1}_{xy}\, V_y + \left(\frac{s^2_{13}}{\sigma_{s^2_{13}}}\right)^2 \qquad \text{with} \\ V_x &= N^{\mathrm{pred}}_x(m_0, s_{13}, \delta, \alpha_{12}, \alpha_{32}, h) - N^{\mathrm{exp}}_x \end{aligned} \tag{3.3}$$

where x and y run over the five combinations given in eq. (3.1). For the "data" N^{exp}_x we use the prediction for N_x at some assumed "true values" of the parameters, $(m_0, s_{13}, \delta, \alpha_{12}, \alpha_{32}, h)^{\mathrm{true}}$. Then the statistical analysis tells us the ability to reconstruct these true values from the data. For the covariance matrix S we assume the following form:

$$S_{xy} = N^{\mathrm{exp}}_x \delta_{xy} + \sigma^2_{\mathrm{norm}} N^{\mathrm{pred}}_x N^{\mathrm{pred}}_y + S^{\mathrm{osc}}_{xy}\,. \tag{3.4}$$

It includes statistical errors, a fully correlated normalisation error σ_{norm}, and the uncertainty introduced from the errors on the oscillation parameters S^{osc}. The normalisation error σ_{norm} arises from the uncertainty on the luminosity and the efficiency. Moreover, the possibility that the non-leptonic decays of H^{++} of eq. (2.14) might occur at a sub-leading level and are not observed introduces an uncertainty in the number of leptonic decays. We adopt a value of $\sigma_{\mathrm{norm}} = 20\%$. We have checked that even an analysis with free normal-

JHEP03(2008)009

ization (i.e., $\sigma_{\rm norm} \to \infty$) leads to very similar results. This means that the information is fully captured by the ratios of branchings.[4]

Via the covariance matrix $S^{\rm osc}$ we account for the fact that the parameters Δm^2_{21}, $|\Delta m^2_{31}|$, s_{12} and s_{23} have a finite uncertainty. We include the errors from eq. (2.6) and take into account the correlations which they introduce between the observables N_x. The last term in eq. (3.3) takes into account the constraint on s_{13} from present data according to eq. (2.7). Let us note that within the time scale of a few years the errors on oscillation parameters are likely to decrease. In particular, also the bound on s_{13} will be strengthened or eventually a finite value could be discovered by upcoming reactor or accelerator experiments, see for example ref. [40]. To be conservative we include only present information, although at the time of the analysis better constraints might be available. We have checked that the precise value of s_{13} within the current limits as well as its uncertainty have a very small impact on our results, and a better determination may lead at most to a marginal improvement of the sensitivities.

3.2 Branching ratios

In figure 1 we show the branching ratios for NH and IH as a function of the lightest neutrino mass m_0. For fixed m_0, the interval for the branching emerges due to the dependence on the phases $\alpha_{12}, \alpha_{32}, \delta$, and also the uncertainty on solar and atmospheric oscillation parameters contributes to the interval. In the plots one can identify the regions of hierarchical neutrino masses, $m_0 < 10^{-3}\,$eV, and QD masses, $m_0 > 0.1\,$eV, where NH and IH become indistinguishable. In the limiting cases $m_0 = 0$ and $m_0 \to \infty$ the analytic expressions for the branchings are rather simple. For NH and $m_0 = 0$ one finds to leading order in the small quantities $r \equiv \Delta m^2_{21}/|\Delta m^2_{31}| \approx 0.03$ and $s^2_{13} < 0.05$ (at 3σ):

$$\mathrm{BR}^{\mathrm{NH},m_0=0}_{\rm ee} \approx s^4_{12} r + 2 s^2_{12} s^2_{13} \sqrt{r} \cos(\alpha_{32} - 2\delta)\,, \tag{3.5}$$

$$\mathrm{BR}^{\mathrm{NH},m_0=0}_{e\mu} \approx 2\left[s^2_{12} c^2_{12} c^2_{23} r + s^2_{23} s^2_{13} + 2 s_{12} c_{12} s_{23} c_{23} s_{13} \sqrt{r} \cos(\alpha_{32} - \delta)\right]\,, \tag{3.6}$$

$$\mathrm{BR}^{\mathrm{NH},m_0=0}_{\mu\mu} \approx s^4_{23} + 2 s^2_{23} c^2_{23} c^2_{12} \sqrt{r} \cos\alpha_{32} + c^4_{23} c^4_{12} r - 4 s^3_{23} c_{23} s_{12} c_{12} s_{13} \sqrt{r} \cos(\alpha_{32} - \delta), \tag{3.7}$$

$$\mathrm{BR}^{\mathrm{NH},m_0=0}_{e\tau} \approx 2\left[s^2_{12} c^2_{12} s^2_{23} r + c^2_{23} s^2_{13} - 2 s_{12} c_{12} s_{23} c_{23} s_{13} \sqrt{r} \cos(\alpha_{32} - \delta)\right]\,, \tag{3.8}$$

$$\mathrm{BR}^{\mathrm{NH},m_0=0}_{\mu\tau} \approx 2 s^2_{23} c^2_{23} \left(1 - 2 c^2_{12} \sqrt{r} \cos\alpha_{32} + c^4_{12} r\right)\,. \tag{3.9}$$

For IH and $m_0 = 0, s_{13} = 0$ we have

$$\mathrm{BR}^{\mathrm{IH},m_0=0}_{\rm ee} = \frac{1}{2}\left(1 - \sin^2 2\theta_{12} \sin^2 \frac{\alpha_{12}}{2}\right)\,, \tag{3.10}$$

$$\mathrm{BR}^{\mathrm{IH},m_0=0}_{e\mu} = c^2_{23} \sin^2 2\theta_{12} \sin^2 \frac{\alpha_{12}}{2}\,, \tag{3.11}$$

$$\mathrm{BR}^{\mathrm{IH},m_0=0}_{\mu\mu} = \frac{c^4_{23}}{2}\left(1 - \sin^2 2\theta_{12} \sin^2 \frac{\alpha_{12}}{2}\right)\,, \tag{3.12}$$

$$\mathrm{BR}^{\mathrm{IH},m_0=0}_{e\tau} = s^2_{23} \sin^2 2\theta_{12} \sin^2 \frac{\alpha_{12}}{2}\,, \tag{3.13}$$

$$\mathrm{BR}^{\mathrm{IH},m_0=0}_{\mu\tau} = \frac{1}{4} \sin^2 2\theta_{23} \left(1 - \sin^2 2\theta_{12} \sin^2 \frac{\alpha_{12}}{2}\right)\,, \tag{3.14}$$

[4] This is true as long as all branchings from eq. (3.1) are used; if the events containing taus are omitted our results depend to some degree on the value adopted for $\sigma_{\rm norm}$.

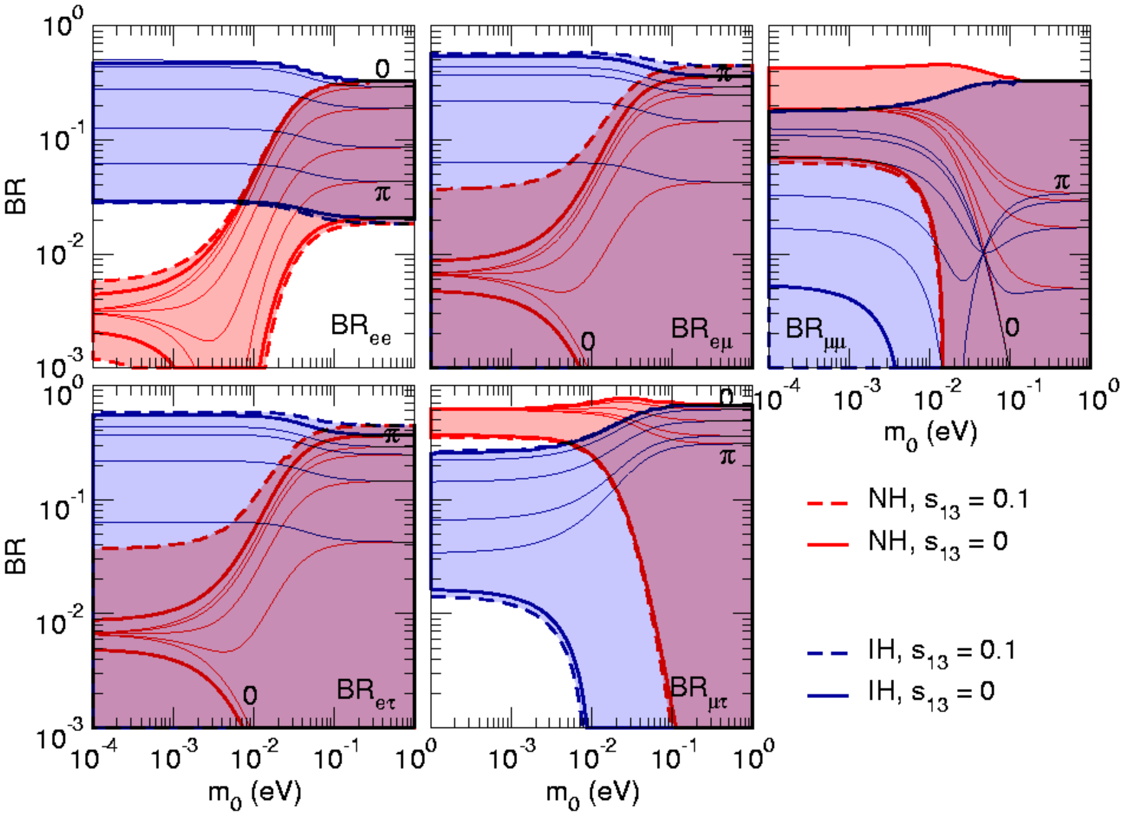

Figure 1: Branching ratios $\mathrm{BR}(H \to \ell_a \ell_b)$ as function of the lightest neutrino mass m_0 for NH (light-red) and IH (dark-blue). The thick solid lines are for $s_{13} = 0$, and the thick dashed lines for $s_{13} = 0.1$, where the dependence on phases as well as the uncertainty of solar and atmospheric oscillation parameters at 2σ are included. The thin solid lines show the branchings for oscillation parameters fixed at the best fit points eq. (2.6), $s_{13} = 0$, $\alpha_{32} = \pi$, and $\alpha_{12} = 0, \pi/4, \pi/2, 3\pi/4, \pi$.

and in the limit $m_0 \to \infty$ with $s_{13} = 0$ the branchings become

$$\mathrm{BR}^{\mathrm{QD}}_{ee} = \frac{1}{3}\left(1 - \sin^2 2\theta_{12}\,\sin^2\frac{\alpha_{12}}{2}\right) = \frac{2}{3}\,\mathrm{BR}^{\mathrm{IH},m_0=0}_{ee}\,, \tag{3.15}$$

$$\mathrm{BR}^{\mathrm{QD}}_{e\mu} = \frac{2}{3}\,c^2_{23}\,\sin^2 2\theta_{12}\,\sin^2\frac{\alpha_{12}}{2} = \frac{2}{3}\,\mathrm{BR}^{\mathrm{IH},m_0=0}_{e\mu}\,, \tag{3.16}$$

$$\mathrm{BR}^{\mathrm{QD}}_{\mu\mu} = \frac{1}{3}\left[1-\frac{1}{2}\sin^2 2\theta_{23}\left(1-s^2_{12}\cos\alpha_{31}-c^2_{12}\cos\alpha_{32}\right)-c^4_{23}\sin^2 2\theta_{12}\sin^2\frac{\alpha_{12}}{2}\right], \tag{3.17}$$

$$\mathrm{BR}^{\mathrm{QD}}_{e\tau} = \frac{2}{3}\,s^2_{23}\,\sin^2 2\theta_{12}\,\sin^2\frac{\alpha_{12}}{2} = \frac{2}{3}\,\mathrm{BR}^{\mathrm{IH},m_0=0}_{e\tau}\,, \tag{3.18}$$

$$\mathrm{BR}^{\mathrm{QD}}_{\mu\tau} = \frac{1}{3}\sin^2 2\theta_{23}\left(1 - s^2_{12}\cos\alpha_{31} - c^2_{12}\cos\alpha_{32} - \frac{1}{2}\sin^2 2\theta_{12}\sin^2\frac{\alpha_{12}}{2}\right). \tag{3.19}$$

Note that for a vanishing lightest neutrino mass, $m_0 = 0$, there is only one physical Majorana phase, α_{32} for NH, and α_{12} for IH, as clear from eqs. (2.4) and (2.5).

In the following we will explore the parameter dependencies of these branchings to obtain information on the neutrino mass spectrum and on Majorana phases. The rather wide ranges for the branchings in the cases of IH and QD spectrum suggest a strong dependence on the phases, and as we will see in section 3.4 these are the cases where Majorana phases can be measured very efficiently. The determination of the mass spectrum is somewhat more subtle.

A clear signature for the NH with small m_0 is provided by $\mathrm{BR}_{\mathrm{ee}}$.[5] Eq. (3.5) shows that for NH and $m_0 = 0$, $\mathrm{BR}_{\mathrm{ee}}$ is suppressed by r and/or s_{13}^2, and there is the upper bound $\mathrm{BR}_{\mathrm{ee}} < 5.3 \times 10^{-3}$ for the largest value of s_{12}^2 allowed at 2σ and $s_{13}^2 = 0.01$, in agreement with figure 1. In contrast, for IH with $m_0 < 0.01\,\mathrm{eV}$ and for QD spectrum, eqs. (3.10) and (3.15) give the lower bounds $\mathrm{BR}_{\mathrm{ee}} > (1 - \sin^2 2\theta_{12})/2 \approx 0.03$ and $\mathrm{BR}_{\mathrm{ee}} > (1 - \sin^2 2\theta_{12})/3 \approx 0.02$, respectively. Therefore, the characteristic signature of normal hierarchical spectrum is the suppression of Higgs decays into two electrons.

From a first glance at figure 1 one could expect that it might be difficult to distinguish IH and QD spectra, since there is always overlap between the allowed regions in the branchings. Indeed, if only branchings involving electrons and muons ($\mathrm{BR}_{\mathrm{ee}}$, $\mathrm{BR}_{e\mu}$, $\mathrm{BR}_{\mu\mu}$) are considered there is some degeneracy between IH and QD, especially if s_{13} is allowed to be close to the present bound. However, as we will show, due to the complementary dependence on the phases of all the BR_{ab} including also taus, the degeneracy is broken and these two cases can be disentangled. Consider, for example, $\mathrm{BR}_{\mu\mu}$ and $\mathrm{BR}_{\mu\tau}$: in the case of IH with $m_0 = 0$ they behave very similar as a function of α_{12}, see eqs. (3.12) and (3.14), whereas for QD they show opposite dependence, compare eqs. (3.17) and (3.19), and phases which give $\mathrm{BR}_{\mu\mu}^{\mathrm{QD}} = 0$ maximise $\mathrm{BR}_{\mu\tau}^{\mathrm{QD}}$.

Note that for $s_{13} = 0$ and $s_{23}^2 = 0.5$, $\mathrm{BR}_{e\mu}$ and $\mathrm{BR}_{e\tau}$ are identical. Nevertheless there is important complementariness between them. First, the uncertainty on s_{23}^2, see eq. (2.6), affects each of them significantly, and it reduces the final sensitivity if only $\mathrm{BR}_{e\mu}$ is used in the analysis. But since $\mathrm{BR}_{e\mu}$ and $\mathrm{BR}_{e\tau}$ are related by the transformation $s_{23} \to c_{23}$, $c_{23} \to -s_{23}$ this uncertainty is cancelled if both of them are included in the fit. Second, it can be shown that the leading order term in s_{13} is the same for $\mathrm{BR}_{e\mu}$ and $\mathrm{BR}_{e\tau}$, apart from an opposite sign. Therefore, also the impact of s_{13} is strongly reduced if information from both of them is taken into account. One can observe from figure 1 that for small m_0 and NH, $\mathrm{BR}_{e\mu}$ and $\mathrm{BR}_{e\tau}$ show a significant dependence on s_{13}, while in the other cases the dependence is mild. The reason is a leading term linear in $\sqrt{r}s_{13}$ in eqs. (3.6) and (3.8), whereas in all other cases s_{13} appears either in sub-leading terms or at least at second order.

3.3 Determination of the neutrino mass spectrum

Let us now quantify the ability to determine the neutrino mass spectrum by performing a χ^2 analysis as described in section 3.1. In figure 2 we show the χ^2 by assuming that "data" are generated by a hierarchical spectrum with normal ordering (left), a hierarchical spectrum with inverted ordering (middle), or a QD spectrum (right). These data are fitted with both possibilities for the ordering (NH, light-red curves, and IH, dark-blue curves) and a value for m_0 shown on the horizontal axis. We minimise the χ^2 with respect to the other parameters, taking into account the current bound on s_{13}. The results are shown

[5] Note that the behaviour of $\mathrm{BR}_{\mathrm{ee}}$ is the same as the effective neutrino mass probed in neutrino-less double beta-decay, which is also proportional to $|M_{\mathrm{ee}}|$, see for example ref. [41].

JHEP03(2008)009

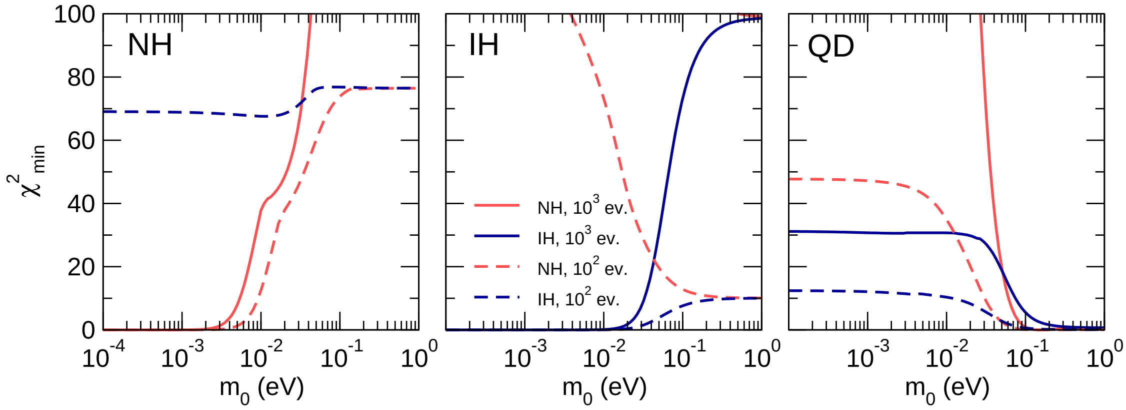

Figure 2: $\chi^2_{\rm min}$ vs m_0 assuming a true hierarchical spectrum with NH (left) and IH (middle), and a true QD spectrum (right). The χ^2 is shown for $\epsilon N_{2H} = 100$ (dashed) and 1000 (solid) events, and $\sigma_{\rm norm} = 20\%$. In the fit we assume either NH (light-red) or IH (dark-blue), and we minimise with respect to s_{13} and the phases. We adopt the following true parameter values. Left: $m_0 = 0$, NH, $\alpha_{32} = \pi$; middle: $m_0 = 0$, IH, $\alpha_{12} = 0$; right: $m_0 = 0.15\,$eV, $\alpha_{12} = 0.1\pi$, $\alpha_{32} = 1.6\pi$; and always $s_{13} = 0$.

for a total number of doubly charged scalars decaying into like-sign leptons of $\epsilon N_{2H} = 10^3$ (solid) and 10^2 (dashed).

First we discuss the sensitivity to hierarchical spectra with a very small lightest neutrino mass m_0. The left panel of figure 2 shows that a NH with small m_0 can be identified with very high significance. An inverted hierarchical spectrum as well as a QD spectrum have $\Delta\chi^2 \gtrsim 60$ already for 100 events. An upper bound on the lightest neutrino mass of $m_0 \lesssim 0.01\,$eV at 3σ can be established by LHC data. As discussed in the previous section this information comes mainly from the suppression of the decay into two electrons, which occurs only for normal hierarchical spectrum. An inverted hierarchical spectrum (middle panel) can be distinguished from a QD one at around 3σ with 100 events, where the χ^2 increases roughly linearly with the number of events. The ability to exclude a QD spectrum in case of a true IH depends on the true value of the Majorana phase α_{12}. The example chosen in figure 2, $\alpha_{12}^{\rm true} = 0$, corresponds to the worst case; for all other values of α_{12} the χ^2 for QD is bigger.

Figure 3 shows the ability to identify a hierarchical spectrum as a function of the true value for the Majorana phase, where for $m_0 = 0$ there is only one physical phase. The shaded regions show that for 1000 events the true spectrum can be identified at 5σ significance, and an upper bound on the lightest neutrino mass $m_0 < 8 \times 10^{-3}\,$eV for NH and $m_0 < 4 \times 10^{-2}\,$eV for IH is obtained, independent of the true phase. For the black contours in figure 3 we do not use the information from decays into taus, i.e., we use only the lepton pairs $(ee), (e\mu), (\mu\mu)$. This analysis illustrates the importance of the tau events. For example, if tau events are not used an IH with $m_0 = 0$ cannot be distinguished from a QD spectrum for $\alpha_{12}^{\rm true} \sim \pi$. Also the sensitivity to a NH is significantly reduced, which becomes even more severe if less events were available.

Now we move to the discussion of a true QD spectrum. As shown in the right panel of

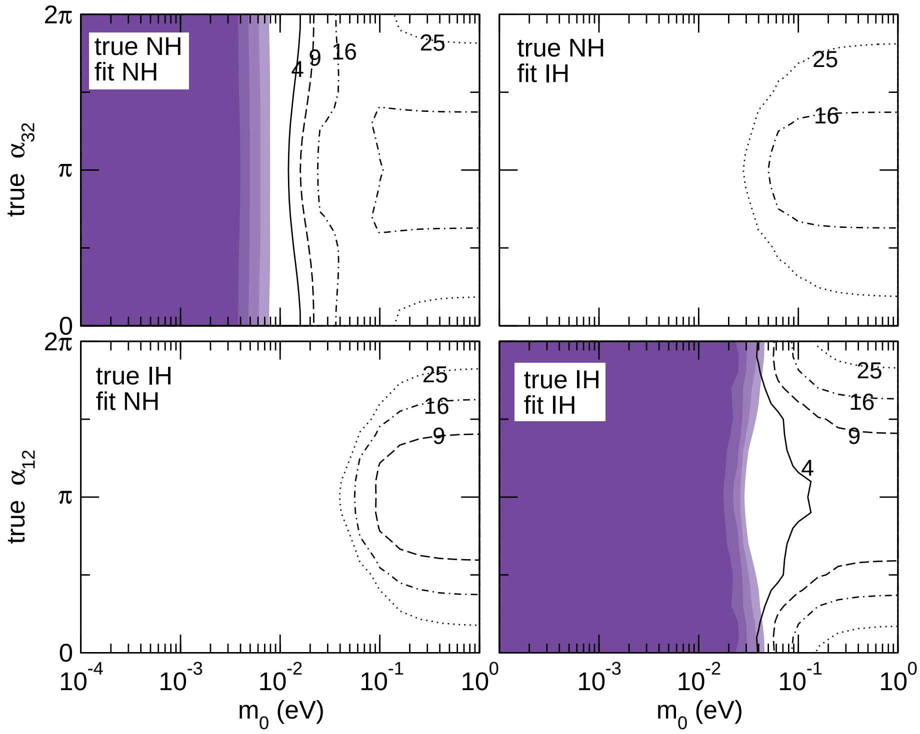

Figure 3: Determination of hierarchical neutrino mass spectra, $m_0^{\rm true} = 0$, assuming 1000 Higgs pair decays. The upper (lower) panels are for a true NH (IH), and for the left (right) panels the fit is performed assuming a NH (IH). As a function of the true value of the Majorana phases we show contours $\chi^2 = 4, 9, 16, 25$ (from dark to light), minimising with respect to all parameters except from m_0. Coloured regions correspond to our standard analysis, whereas for the black contours we do not use decays into tau leptons.

figure 2 also a QD spectrum can be identified quite well, and a lower bound on the lightest neutrino mass of $m_0 > 2\,(6) \times 10^{-2}\,$eV at 3σ can be obtained for 100 (1000) events. Note that for the example shown in figure 2, 100 events give a $\Delta\chi^2 \approx 12.4$ for the IH with $m_0 = 0$, which corresponds roughly to an exclusion at 3.5σ. The potential to exclude a hierarchical inverted spectrum depends on the true values of the Majorana phases, and the true values of α_{12} and α_{32} adopted in figure 2 correspond to the worst sensitivity. In figure 4 we show contours of $\Delta\chi^2$ for IH with $m_0 = 0$ assuming a true QD spectrum, in the plane of the true Majorana phases. For 1000 events we find some islands in the plane of α_{12} and α_{32} where the χ^2 reaches values as low as 30 (compare figure 2), however in most parts of the parameter space the exclusion is at more than 7σ. For 100 events typically a significance better than 4σ is reached, but there are some notable regions ($-\pi/2 \lesssim \alpha_{12} \lesssim \pi/2$ and $\alpha_{32} \sim \pi/2, 3\pi/2$) with χ^2 values between 16 and 9.

Let us add that for the exclusion of an inverted hierarchical spectrum in the case of a true QD spectrum the branchings into tau leptons are crucial. If only electron and muon events are used in most regions of the parameter space an IH with $m_0 = 0$ can fit data from a QD spectrum. For $(ee), (e\mu), (\mu\mu)$ branchings a degeneracy between IH and QD

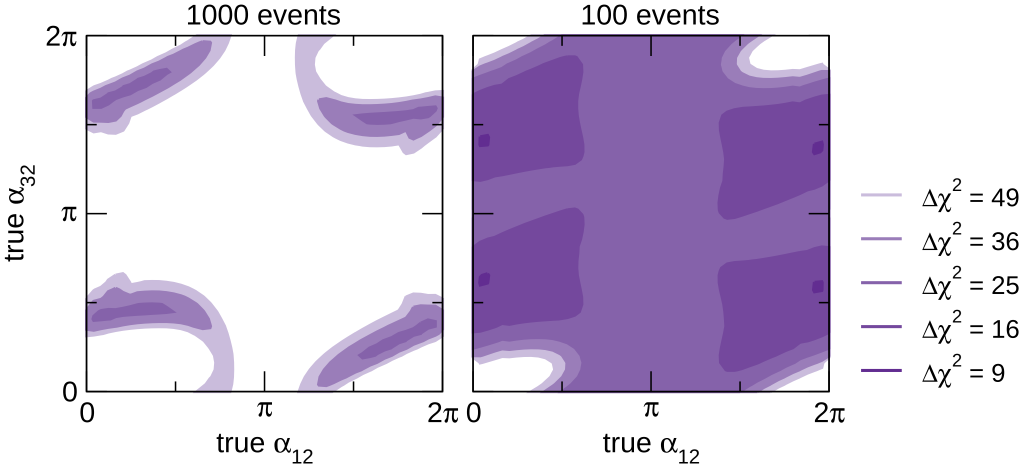

Figure 4: Exclusion of an IH with $m_0 = 0$ in the case of a true QD spectrum. We show χ^2 contours for 1000 events (left) and 100 events (right) in the plane of the true Majorana phases assuming a true QD spectrum ($m_0^{\rm true} = 0.15\,\text{eV}$, $s_{13}^{\rm true} = 0$) fitted with IH and $m_0 = 0$, minimising with respect to all other parameters.

appears due to the freedom in adjusting s_{13}, δ, θ_{23} and the Majorana phases. This effect is also apparent from the black contour lines in figure 3 (lower-right panel). The significance of this degeneracy depends on details such as the errors imposed on s_{13}^2 and s_{23}^2, as well as on the systematical error $\sigma_{\rm norm}$. As discussed in section 3.2, taking into account also decays into $e\tau$ and $\mu\tau$ is crucial to break this degeneracy, and in the full analysis used to calculate figures 2 and 4 the dependence on subtleties such as s_{13} and $\sigma_{\rm norm}$ is small.

3.4 Determination of Majorana phases

Let us now investigate the tantalising possibility to determine the Majorana phases $\alpha_{ij} \equiv \alpha_i - \alpha_j$ from the doubly charged Higgs decays. Since the decay is governed by a single diagram without any interference term the decays are CP conserving, and therefore no explicit CP violating effects can be observed. Nevertheless, the branchings depend (in a CP conserving way) on the phases, which eventually may allow to establish CP violating values for them. In general the measurement of Majorana phases is a very difficult task. Probably the only hope to access these phases will be neutrino-less double beta-decay in combination with an independent neutrino mass determination, where under very favourable circumstances [41] the phase α_{12} might be measurable.

We start by discussing some general properties of the branchings related to the Majorana phases. Using eqs. (2.4) and (2.5) one can write:

$$\mathrm{BR}_{ab} \propto |M_{ab}|^2 = \left| \sum_{i=1}^{3} V_{ai} V_{bi}\, e^{i\alpha_i}\, m_i \right|^2 , \tag{3.20}$$

From this expression it is evident that for a vanishing lightest neutrino mass, $m_0 = 0$, there is only one physical Majorana phase, α_{32} for NH and α_{12} for IH. Next we note that since

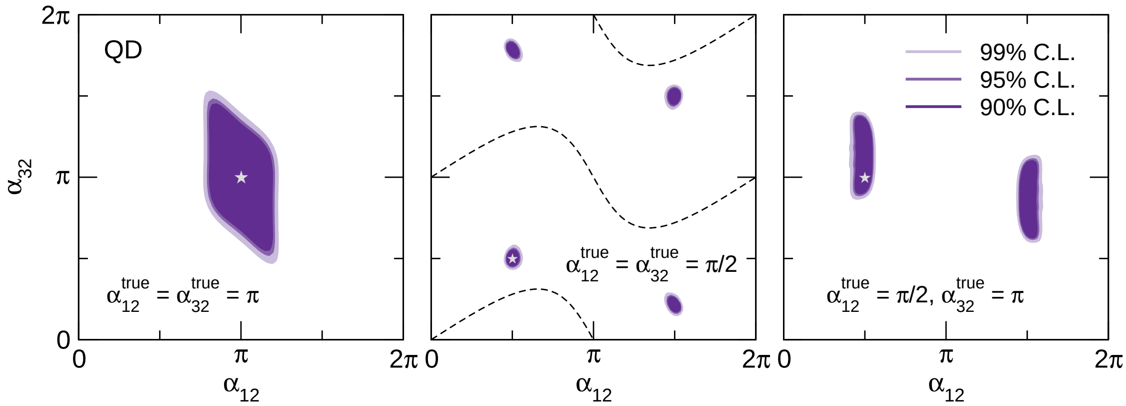

Figure 5: Determination of the Majorana phases for QD spectrum ($m_0 = 0.15\,\text{eV}$) from 1000 doubly-charged Higgs pair events. We assume $s_{13}^{\text{true}} = 0$ and three example points for the true values of the Majorana phases given in each panel. The dashed lines in the middle panel correspond to the true values of the phases for which the degenerate solution according to eq. (3.23) appears at a CP conserving value of α_{32}.

$V_{e3} \propto s_{13}$, it is clear that for $s_{13} = 0$ all branchings involving electrons can only depend on α_{12}.[6] Since the small effects of s_{13} cannot be explored efficiently, the determination of both phases simultaneously necessarily involves $\text{BR}_{\mu\mu}$ and/or $\text{BR}_{\mu\tau}$, see also eqs. (3.5) to (3.19). Furthermore, from eq. (3.20) it can be seen that the branchings are invariant under

$$\alpha_{ij} \to 2\pi - \alpha_{ij}\,, \qquad \delta \to 2\pi - \delta\,. \tag{3.21}$$

This symmetry is a consequence of the fact that there is no CP violation in the decays, and therefore the branchings have to be invariant under changing the signs of all phases simultaneously.

In figure 5 we show that for a QD spectrum the observation of the decay of 1000 doubly-charged Higgs pairs allows to determine both Majorana phases. We assume some true values for the two phases and then perform a fit leaving all parameters free, where for s_{13} we impose the constraint from present data. The actual accuracy to determine the phases depends on their true values, where we show three different examples in the three panels. For $\alpha_{12} = \alpha_{32} = \pi$ (left panel) the allowed region is the largest, however the phases can be constrained to a unique region. In the other two cases the accuracy is better, but some ambiguities are left. The symmetry from eq. (3.21) is apparent in all panels, whereas in the case $\alpha_{12} = \alpha_{32} = \pi$ it does not introduce an ambiguity.

The features of figure 5 can be understood from eqs. (3.15) to (3.19). In addition to the symmetry eq. (3.21) one finds that in the limit $s_{13} = 0$ the phases α_{31} and α_{32} appear only in the particular combination

$$\left(s_{12}^2 \cos\alpha_{31} + c_{12}^2 \cos\alpha_{32}\right) \propto \cos\left(\alpha_{32} - \varphi\right) \qquad \text{with} \qquad \tan\varphi = \frac{s_{12}^2 \sin\alpha_{12}}{c_{12}^2 + s_{12}^2 \cos\alpha_{12}}\,, \tag{3.22}$$

[6]For the same reason only α_{12} can be tested in neutrino-less double beta-decay, where $|M_{\text{ee}}|$ is probed.

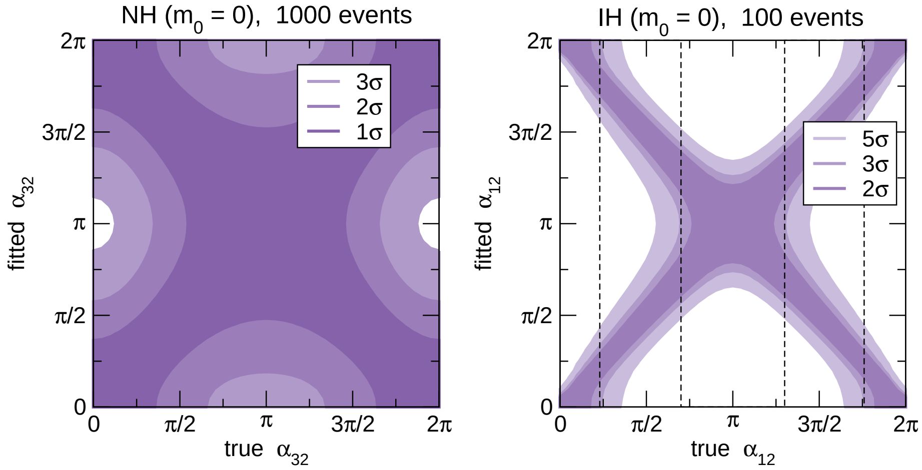

Figure 6: Determination of the Majorana phase for vanishing lightest neutrino mass. We assume $s_{13}^{\rm true} = 0$. Left: 1, 2, 3σ ranges for α_{32} as a function of its true value for NH assuming 1000 doubly-charged Higgs pair events. Right: 2, 3, 5σ ranges for α_{12} as a function of its true value for IH assuming 100 doubly-charged Higgs pair events. The dashed vertical lines indicate the region where CP violating values of α_{12} can be established at 3σ.

where we have used α_{12} and α_{32} as independent parameters. For constant α_{12} there are two values of α_{32} which leave this combination invariant: for each α_{32} we expect a degenerate solution at

$$\alpha'_{32} = 2\varphi - \alpha_{32}\,. \tag{3.23}$$

For $\alpha_{12} = \pi/2$ one finds $2\varphi \approx 0.28\pi$. In the case of $\alpha_{32} = \pi$ shown in the right panel of figure 5 this degenerate solution appears at $\alpha'_{32} \simeq 1.28\pi$, which cannot be resolved from the original one, and we are left with a two-fold ambiguity, due to eq. (3.21). In the middle panel, for $\alpha_{12} = \alpha_{32} = \pi/2$, the ambiguity (3.23) leads to a separated solution around $\alpha'_{32} \simeq 1.78\pi$ and, together with the symmetry from eq. (3.21) we end up with four degenerate solutions. However, in this case the individual regions are rather small, and the CP violating values of both phases can be established despite the presence of the four-fold ambiguity.

Note that the symmetry (3.21) does not mix CP conserving and violating values of the phases, whereas this can happen for the degeneracy eq. (3.23). The dashed curves in the middle panel of figure 5 correspond to the true values of the phases, for which $\alpha'_{32} = 0$ or π. Hence, along these curves CP violating values for α_{32} cannot be established since the degeneracy is located at a CP conserving value.

Let us now discuss the potential to determine Majorana phases in case of hierarchical spectra. As mentioned above, in this case there is only one physical phase, α_{32} for NH and α_{12} for IH. In figure 6 we show the allowed interval for this phase which is obtained from the data as a function of its true value. In the fit the χ^2 is minimised with respect to all other parameters. The left panel shows that for NH even with 1000 events at most a 2σ indication can be obtained, on whether α_{32} is closer to zero or π. This can be understood from eqs. (3.5) to (3.9), which show that α_{32} appears at least suppressed by $\sqrt{r}$. In contrast, as visible in the right panel, for IH a rather precise determination of α_{12} is possible already for 100 events, apart from the ambiguity $\alpha_{12} \to 2\pi - \alpha_{12}$. For α_{12} around $\pi/2$ or $3\pi/2$ its CP violating value can be established, as marked by the vertical lines in figure 6. The good sensitivity is obvious from eqs. (3.10) to (3.14), which show a strong dependence of the leading terms in the branchings on α_{12}.

4. Summary and concluding remarks

In this work we have adopted the assumptions that (*i*) neutrino masses are generated by the VEV of a Higgs triplet, (*ii*) the doubly charged component of the triplet is light enough to be discovered at LHC, i.e., lighter than about 1 TeV, and (*iii*) it decays with a significant fraction into like-sign lepton pairs. We have shown that under these assumptions LHC will provide very interesting information for neutrino physics. The reason is that the branching ratio of the doubly charged Higgs into like-sign leptons of flavour a and b, $\mathrm{BR}(H^{++} \to \ell_a^+ \ell_b^+)$, is proportional to the modulus of the corresponding element of the neutrino mass matrix, $|M_{ab}|^2$. Hence the flavour composition of like-sign lepton events at LHC provides a direct test of the neutrino mass matrix.

We have shown that the type of the neutrino mass spectrum (normal hierarchical, inverted hierarchical, or quasi-degenerate) can be identified at the 3σ level already with 100 doubly charged Higgs pairs $H^{--}H^{++}$ decaying into four leptons. Typically such a number of events will be achieved for doubly charged scalar masses below 600 GeV and 100 fb^{-1} integrated luminosity, whereas for masses of 350 GeV of order 1000 events will be obtained. We have found that it is possible to decide whether the lightest neutrino mass is smaller or larger than roughly 0.01 eV, which marks the transition between hierarchical and quasi-degenerate spectra. If it is smaller the mass ordering (normal vs inverted) can be identified. A hierarchical spectrum with normal ordering has a distinct signature, namely a very small branching of the doubly charged Higgs decays into two electrons. Therefore, this mass pattern can easily be confirmed or ruled out at very high significance level. The other two possibilities for the neutrino mass spectrum, inverted hierarchical or quasi-degenerate, are somewhat more difficult to distinguish, but also in this case very good sensitivity is obtained, depending on the observed number of events.

In this respect the inclusion of final states involving tau leptons is important, since if only electrons and muons are considered a degeneracy between IH and QD spectra appears. In our analysis we have conservatively assumed that events where one of the four charged leptons is a tau can be reconstructed efficiently, thanks to the kinematic constraints and the information on the invariant mass of the event available from events without a tau. Certainly a more realistic study including detailed simulations and event reconstruction should confirm the assumptions which we have adopted here.

The decay of the doubly charged Higgs in this framework does not show explicit CP violation, since the decay is dominated by a tree-level diagram without any interference term which could induce CP violation. Nevertheless, the CP conserving branching ratios strongly depend on the Majorana CP phases of the lepton mixing matrix. Therefore, the framework considered here opens the fascinating possibility to measure the Majorana phases in the neutrino mass matrix via CP even observables. Our results show that for an inverted hier-

JHEP03(2008)009

archical spectrum as well as for quasi-degenerate neutrinos this is indeed possible. In the first case, there is only one physical phase, α_{12}, which can be determined up to an ambiguity $\alpha_{12} \leftrightarrow 2\pi - \alpha_{12}$ already with 100 events. In the case of a quasi-degenerate spectrum both Majorana phases can be measured, where, depending on the actual values some ambiguities might occur. In many cases CP violating values of the phases can be established.

Certainly the observation of a doubly charged scalar at LHC would be a great discovery of physics beyond the Standard Model. Of course this alone does by no means confirm the Higgs triplet mechanism for neutrino masses, since doubly charged particles decaying into leptons are predicted in many models. Therefore, in case such a particle is indeed found at LHC various consistency checks will have to be performed. It might turn out that the relation $\mathrm{BR}(H^{++} \to \ell_a^+ \ell_b^+) \propto |M_{ab}|^2$ cannot be fulfilled for any neutrino mass matrix consistent with oscillation data. This would signal that a Higgs triplet cannot be the only source for neutrino masses. In this respect the information from decays into leptons of all flavours (including taus) will be important. For example, also in the Zee-Babu model [4] for neutrino masses doubly charged scalars might be found at LHC. However, in this case branchings into tau leptons are suppressed by powers of $(m_\mu/m_\tau)^2$ with respect to muons [27], whereas in the Higgs triplet model they are of similar size because of close to maximal θ_{23} mixing.

If LHC data on $\mathrm{BR}(H^{++} \to \ell_a^+ \ell_b^+)$ will be consistent with a neutrino mass matrix from oscillation data, an analysis as pointed out in this work can be performed. Also in this case it will be of crucial importance to cross check the results with independent measurements, for example the determination of the neutrino mass ordering by oscillation experiments, or the measurement of the absolute neutrino mass in tritium beta-decay, neutrino-less double beta-decay or through cosmological observations. In particular, neutrino-less double beta-decay will provide a crucial test, since it gives an independent determination of the $|M_{\mathrm{ee}}|$ element of the neutrino mass matrix, which — combined with information from oscillation experiments — will further constrain the allowed flavour structure of the di-lepton events at LHC. The next generation of neutrino-less double beta-decay experiments is expected to probe the regime of the QD neutrino spectrum within a timescale comparable to the LHC measurement. Information from searches for lepton flavour violating processes may be used as additional important consistency checks for the model.

In conclusion, a TeV scale Higgs triplet offers an appealing mechanism to provide mass to neutrinos, which can be directly tested at the LHC. Such a scenario opens the possibility to measure the Majorana phases of the lepton mixing matrix, which in general is a very difficult — if not a hopeless task.

Acknowledgments

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JHEP03(2008)009

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JHEP03(2008)009

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JHEP03(2008)009