Vector leptons in the Higgs triplet model

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We analyze the phenomenological implications of introducing vector-like leptons on the Higgs sector in the Higgs triplet model. We impose only a parity symmetry which disallows mixing between the new states and the ordinary leptons. If the vector leptons are allowed to be relatively light, they enhance or suppress the decay rates of loop-dominated neutral Higgs boson decays $h \to \gamma\gamma$ and $h \to Z\gamma$, and affect their correlation. An important consequence is that, for light vector leptons, the decay patterns of the the doubly charged Higgs boson will be altered, modifying the restrictions on their masses. We study the implications for signals at the LHC, for both $\sqrt{s} = 7$ TeV and for 13 TeV, and show that doubly charged boson decays into same-sign vector leptons could be observed.

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I. INTRODUCTION

With the discovery of the Standard Model (SM) Higgs-like scalar at the LHC [1], the SM particle content seems complete. In particular, the mass and couplings of the neutral Higgs boson seem to disfavor an additional chiral generation of quarks and leptons [2]. However, additional vector-like fermions, in which an SM generation is paired with another one of opposite chirality, and with identical couplings, are less constrained, as there is no quadratic contribution to their masses. These states appear as natural extensions of the SM particle content in theories with warped or universal extra dimensions, as Kaluza-Klein excitations of the bulk field [3], in nonminimal supersymmetric extensions of the SM [4], in composite Higgs models [5], in the little Higgs model [6] and in gauged flavor groups [7]. Vector fermions have identical left- and right-handed couplings and can have masses which are not related to their couplings to the Higgs bosons [8]. Depending on the dominant decay mode, the limits on new vector-like fermions range from ∼100–600 GeV [9], rendering them observable at the LHC.

Vector quarks can modify both the production and decays of the Higgs boson at the LHC, while vector leptons do not carry $SU(3)_c$ charge and can only modify the decay patterns of the Higgs. The study of the latter would be a sensitive probe for new physics. The lepton states contribute to self-energy diagrams for electroweak gauge boson masses and precision observables, and consistent limits on their masses and mixings have been obtained [10,11].

Vector leptons have been studied in the context of the SM [11–13], but less so for models beyond the SM, where they can also significantly alter the phenomenology of the model. In the SM, introducing heavy fermions provides a contribution of the same magnitude and sign as that of the top quark and interferes destructively with the dominant W contribution, reducing the $h \to \gamma\gamma$ rate with respect to its SM value. Recent studies indicate that cancellations between scalar and fermionic contributions allow a wide range of Yukawa and mass mixings among vector states [14]. An investigation of vector leptons in the two-Higgs-doublet model [14] showed that the presence of additional Higgs bosons alleviates electroweak precision constraints. Introducing vector leptons into supersymmetry [15] can improve vacuum stability and enhance the diphoton rate by as much as 50%, while keeping new particle masses above 100 GeV and preserving vacuum stability conditions.

In the present work, we investigate vector leptons in the context of the Higgs triplet model (HTM). We do not deal with the LHC phenomenology (pair production and decay) of the extra leptons, which has been discussed extensively in the literature [16], choosing instead to focus on signature features of this model. We have previously shown that in the Higgs triplet model, an enhancement of the $h \to \gamma\gamma$ rate is possible only for the case where the doublet and triplet neutral Higgs fields mix considerably [17]. We extend our analysis to include additional vector-like leptons in the model and investigate how these affect the Higgs diphoton decay rate, with or without significant mixing in the neutral Higgs sector. Originally, both the CMS and ATLAS experiments at LHC observed an enhancement of the Higgs diphoton rate, while the diboson rates ($h \to WW^*$, ZZ^*) have been roughly consistent with SM expectations. At present CMS observes $\sigma(pp \to h) \times BR(h \to \gamma\gamma) = 0.77 \pm 0.27$ times the SM rate, while ATLAS observes $\sigma(pp \to h) \times BR(h \to \gamma\gamma) = 1.55^{+0.33}_{-0.28}$ times the SM rate [18]. Given these numbers, it is possible that either the SM value will be proven correct, or a modest enhancement will persist. A further test of the SM is the correlation of the decay $h \to Z\gamma$ with one for $h \to \gamma\gamma$. We also include the prediction for the vector-lepton effect on the branching ratio of $h \to Z\gamma$ and comment on its relationship with the diphoton decay.

We have an additional reason to investigate the effects of vector leptons in the Higgs triplet model. The model

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includes doubly charged Higgs bosons, predicted by most models to be light. Being pair produced, the doubly charged Higgs bosons are assumed to decay into a pair of leptons with the same electric charge through Majorana-type interactions [19]. Assuming a branching fraction of 100% decays into leptons, i.e., neglecting possible decays into W-boson pairs, the doubly charged Higgs mass has been constrained to be larger than about 450 GeV or more, depending on the decay channel. However, if the vector leptons are light enough, which they can be, the doubly charged Higgs bosons can decay into them and thus evade the present collider bounds on their masses. We investigate this possibility in the second part of this work.

Our work is organized as follows. We introduce the Higgs triplet model with vector leptons in Sec. III. In Sec. III we analyze the effect of the vector leptons on the decays of the neutral Higgs bosons, and discuss the constraints on the parameter space coming from requiring agreement with present LHC data. We include both loop-dominated decays: $h \to \gamma\gamma$ in Sec. III A and $h \to Z\gamma$ in Sec. III B. In Sec. IV we analyze the effect of the vector leptons on the production and decay mechanisms of the doubly charged Higgs at the LHC. We summarize our findings and conclude in Sec. V.

II. THE MODEL

The Higgs triplet model has been studied previously [20]. Here we concentrate on the effect of extending the model by incorporating additional vector leptons. For the purpose of our analysis, vector quarks either do not exist, or are much heavier and decouple from the spectrum. The model contains a vector-like fourth generation of leptons, namely the $SU(2)_L$ left-handed lepton doublets $L'_L = (\nu'_L, e'_L)$, right-handed charged and neutral lepton singlets, ν'_R and e'_R, and the mirror right-handed lepton doublets, $L''_R = (\nu''_R, e''_R)$ and left-handed charged and neutral lepton singlets ν''_L and e''_L. The vector-like leptons have the following quantum numbers under $SU(3)_C \times SU(2)_L \times U(1)_Y$:

$$L'_L = (\mathbf{1}, \mathbf{2}, -1/2), \qquad L''_R = (\mathbf{1}, \mathbf{2}, -1/2), \qquad e'_R = (\mathbf{1}, \mathbf{1}, -1), \qquad e''_L = (\mathbf{1}, \mathbf{1}, -1), \qquad \nu'_R = (\mathbf{1}, \mathbf{1}, 0), \qquad \nu''_L = (\mathbf{1}, \mathbf{1}, 0), \tag{2.1}$$

with the electric charge given by $Q = T_3 + Y$, where T_3 the weak isospin. The Lagrangian density of this model is

$$\mathcal{L}_{\mathrm{HTM}} = \mathcal{L}_{\mathrm{kin}} + \mathcal{L}_Y + \mathcal{L}_{\mathrm{VL}} - V(\Phi, \Delta), \tag{2.2}$$

where $\mathcal{L}_{\mathrm{kin}}$, $\mathcal{L}_Y$, $\mathcal{L}_{\mathrm{VL}}$ and $V(\Phi, \Delta)$ are the kinetic term, the Yukawa interaction for the ordinary SM fermions, the mass and Yukawa interaction for the vector leptons, and the scalar potential, respectively. The Yukawa interactions for the ordinary SM leptons are [17]

$$\mathcal{L}_Y = -[\bar{L}^i_L h^{ij}_e \Phi e^j_R + \mathrm{H.c.}] - [h_{ij}\overline{L^{ic}_L} i\tau_2 \Delta L^j_L + \mathrm{H.c.}], \tag{2.3}$$

where $\tilde{\Phi} = i\tau_2\Phi^*$, h_e is a 3×3 complex matrix, and h_{ij} is a 3×3 complex symmetric Yukawa matrix. Additionally, with the vector-like family of leptons as defined above, the vector-lepton part of the Lagrangian density is

$$\begin{aligned}\mathcal{L}_{\mathrm{VL}} = -\Big[& M_L \bar{L}'_L L''_R + M_E \bar{e}'_R e''_L + M_\nu \bar{\nu}'_R \nu''_L + \frac{1}{2} M'_\nu \bar{\nu}'^c_R \nu'_R + \frac{1}{2} M''_\nu \bar{\nu}''^c_L \nu''_L + h'_E (\bar{L}'_L \Phi) e'_R + h''_E (\bar{L}''_R \Phi) e''_L + h'_\nu (\bar{L}'_L \tau \Phi^\dagger) \nu'_R \\ & + h''_\nu (\bar{L}''_R \tau \Phi^\dagger) \nu''_L + h'_{ij} \overline{L'^c_L} i\tau_2 \Delta L'_L + h''_{ij} \overline{L''^c_R} i\tau_2 \Delta L''_R + \lambda^i_E (\bar{L}'_L \Phi) e^i_R + \lambda^i_L (\bar{L}^i_L \Phi) e'_R + \lambda'_{ij} \overline{L^{ic}_L} i\tau_2 \Delta L'_L \\ & + \lambda''_{ij} \overline{L^{ic}_R} i\tau_2 \Delta L''_R + \mathrm{H.c.}\Big],\end{aligned} \tag{2.4}$$

where we include explicit mass terms, Yukawa interactions among vector leptons, and Yukawa interactions between vector leptons and ordinary leptons. The scalar potential is

$$\begin{aligned} V(\Phi, \Delta) = {} & m^2 \Phi^\dagger \Phi + M^2 \,\mathrm{Tr}(\Delta^\dagger \Delta) + [\mu \Phi^T i\tau_2 \Delta^\dagger \Phi + \mathrm{H.c.}] + \lambda_1 (\Phi^\dagger \Phi)^2 + \lambda_2 [\mathrm{Tr}(\Delta^\dagger \Delta)]^2 + \lambda_3 \,\mathrm{Tr}[(\Delta^\dagger \Delta)^2] \\ & + \lambda_4 (\Phi^\dagger \Phi) \mathrm{Tr}(\Delta^\dagger \Delta) + \lambda_5 \Phi^\dagger \Delta \Delta^\dagger \Phi, \end{aligned} \tag{2.5}$$

where m and M are the Higgs bare masses, μ is the lepton-number-violating parameter, and λ_1–λ_5 are the Higgs coupling constants. The expressions for the parameters λ_1–λ_5 in terms of Higgs masses are given in Ref. [17].

The electroweak gauge symmetry is broken by the vacuum expectation values (VEVs) of the neutral components of the doublet and triplet Higgs fields,

$$\langle \Phi^0 \rangle = v_\Phi / \sqrt{2}, \qquad \langle \Delta^0 \rangle = v_\Delta / \sqrt{2}, \tag{2.6}$$

where Φ and Δ are the doublet Higgs field and the triplet Higgs field, with $v^2 \equiv v_\Phi^2 + 2v_\Delta^2 = (246\,\text{GeV})^2$. Higgs masses and coupling constants in the presence of nontrivial mixing in the neutral sector have been obtained previously [17].

One can invoke new symmetries to restrict the interaction of the vector leptons with each other or with the ordinary leptons, or disallow the presence of bare mass terms in the Lagrangian. For instance,

(1) If there is an additional U(1) symmetry under which the primed, double primed and the ordinary leptons have different charges, this would forbid the explicit masses M_L, M_E, M_ν and M'_ν from the Lagrangian. Vector leptons would get masses only through couplings to the Higgs doublet fields [13,21].

(2) If one imposes a symmetry under which all the new SU(2) singlet fields are odd, while the new SU(2) doublets are even, this forces all Yukawa couplings involving new leptons to vanish, $h'_E = h''_E = h'_\nu = h''_\nu = h'_{ij} = h''_{ij} = 0$, and the masses arise only from explicit terms in the Lagrangian [11].

(3) Finally one can impose a new parity symmetry which disallows mixing between the ordinary leptons and the new vector-lepton fields, under which all the new fields are odd, while the ordinary leptons are even [22], i.e., such that $\lambda_E^i = \lambda_L^i = \lambda'_{ij} = \lambda'_{ij} = \lambda''_{ij} = 0$; alternatively, one might choose these couplings to be very small.

In this analysis the focus will be on Higgs decays. We investigate the model subjected to symmetry condition 3, as we would like to neglect mixing between the ordinary and the new vector leptons. When allowed, stringent constraints exist on the masses and couplings with ordinary leptons. New vector leptons are ruled out when they mediate flavor-changing neutral current processes, generate SM neutrino masses or contribute to neutrinoless double beta decay. Recent studies of models which allow mixing between the ordinary leptons and the new ones exist [21,22], but restrictions from lepton-flavor-violating decays either force the new leptons to be very heavy, $M_L, M_E \sim 10$–100 TeV, or reduce the branching ratios for $h \to \tau^+\tau^-$, $\mu^+\mu^-$ and $h \to \gamma\gamma$ decays to 30–40% of the SM prediction, neither of which are desirable features for our purpose here. In the Higgs triplet model, distinguishing signals would be provided by lighter vector leptons. Since imposing no mixing between ordinary and new leptons allows new lepton masses to be as light as $\sim$100 GeV—perhaps without a reduction in the Higgs diphoton branching ratio—we investigate this scenario here. In addition, we also investigate the effect of imposing condition 1, that is, requiring the explicit mass terms in the Lagrangian to be 0.

In the charged sector, the 2×2 mass matrix $\mathcal{M}_E$ is defined as [11,13]

$$(E'_L \;\; e''_L)(\mathcal{M}_E)\begin{pmatrix} e'_R \\ E''_R \end{pmatrix}, \quad \text{with } \mathcal{M}_E = \begin{pmatrix} m'_E & M_L \\ M_E & m''_E \end{pmatrix}, \tag{2.7}$$

where $m'_E = h'_E v_\Phi/\sqrt{2}$ and $m''_E = h''_E v_\Phi/\sqrt{2}$, with v_Φ the VEV of the neutral component of the Higgs doublet. The mass matrix can be diagonalized as follows:

$$V_L^\dagger \mathcal{M}_E V_R = \begin{pmatrix} M_{E_1} & 0 \\ 0 & M_{E_2} \end{pmatrix}. \tag{2.8}$$

The mass eigenvalues are

$$M^2_{E_1,E_2} = \frac{1}{2}\Big[(M_L^2 + m_E'^2 + M_E^2 + m_E''^2) \pm \sqrt{X^2 + Y^2}\Big], \tag{2.9}$$

with

$$X = (M_L^2 + m_E'^2 - M_E^2 - m_E''^2), \quad Y = 2(m''_E M_L + m'_E M_E). \tag{2.10}$$

By convention, $M_{E_1} > M_{E_2}$. For simplicity we assume that the lepton Yukawa couplings h'_E and h''_E are real so that the transformations that diagonalize the mass matrix are real orthogonal matrices,

$$V_L = \begin{pmatrix} \cos\theta_L & \sin\theta_L \\ -\sin\theta_L & cos\theta_L \end{pmatrix}, \tag{2.11}$$

$$V_R = \begin{pmatrix} \cos\theta_R & \sin\theta_R \\ -\sin\theta_R & cos\theta_R \end{pmatrix}. \tag{2.12}$$

The angles $\theta_{L,R}$ are given by

$$\tan\theta_L = \frac{m''_E M_L + m'_E M_E}{M^2_{E_2} - M_L^2 - m_E'^2}, \tag{2.13}$$

$$\tan\theta_R = \frac{m'_E M_L + m''_E M_E}{M^2_{E_2} - M_L^2 - m_E''^2}. \tag{2.14}$$

The eigenstates of the vector-lepton mixing matrix enter in the evaluation of $h \to \gamma\gamma$ and $h \to Z\gamma$ in the next section.

III. PRODUCTION AND DECAYS OF THE NEUTRAL HIGGS BOSON

The presence of the vector leptons affects the loop-dominated decays of the neutral Higgs, $h \to \gamma\gamma$ and $h \to Z\gamma$, and the possible relationships between them. In the Higgs triplet model, singly and doubly charged bosons also enter in the loops, creating a different dynamic than in the SM. We analyze these decays in turn, and look for possible correlations between them.

A. $h \to \gamma\gamma$

Recently, the Higgs triplet model has received renewed interest because of attempts to reconcile the excess of events in $h \to \gamma\gamma$ observed at the LHC over those predicted by the SM. Such an enhancement hints at the presence of additional particles—singlets under $SU(3)_c$ but charged under $U(1)_{\text{em}}$—which affect only the loop-dominated decay branching ratio, while leaving the production cross section and tree-level decays largely unchanged. Vector leptons are prime candidates for such particles, so we study their contribution to the Higgs decay branching. The decay width $h \to \gamma\gamma$ is

$$\begin{aligned}
&[\Gamma(h\to\gamma\gamma)]_{\rm HTM} \\
&\quad = \frac{G_F\alpha^2 m_h^3}{128\sqrt{2}\pi^3}\bigg|\sum_f N_c^f Q_f^2 g_{hff} A_{1/2}(\tau_f^h) \\
&\qquad + g_{hWW}A_1(\tau_W^h) + \tilde{g}_{hH^\pm H^\mp}A_0(\tau_{H^\pm}^h) \\
&\qquad + 4\tilde{g}_{hH^{\pm\pm}H^{\mp\mp}}A_0(\tau_{H^{\pm\pm}}^h) + \frac{\mu_{E_1} g_{hff}}{M_{E_1}}A_{1/2}(\tau_{E_1}^h) \\
&\qquad + \frac{\mu_{E_2} g_{hff}}{M_{E_2}}A_{1/2}(\tau_{E_2}^h)\bigg|^2, \qquad (3.1)
\end{aligned}$$

with M_{E_1}, M_{E_2} given in Eq. (2.9), and where the loop functions for spin-0, spin-1/2 and spin-1 are given by

$$A_0(\tau) = -[\tau - f(\tau)]\tau^{-2}, \qquad (3.2)$$

$$A_{1/2}(\tau) = -\tau^{-1}[1 + (1-\tau^{-1})f(\tau^{-1})], \qquad (3.3)$$

$$A_1(\tau) = 1 + \frac{3}{2}\tau^{-1} + 4\tau^{-1}\left(1 - \frac{1}{2}\tau^{-1}\right)f(\tau^{-1}), \qquad (3.4)$$

and the function $f(\tau)$ is given by

$$f(\tau) = \begin{cases} \arcsin^2\sqrt{\tau}, & 0<\tau\le 1, \\ -\frac{1}{4}\left[\log\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi\right]^2, & \tau > 1, \end{cases} \qquad (3.5)$$

with $\tau_i^h = \frac{m_h^2}{4m_i^2}$, and m_i is the mass of the particle running in the loop [11]. In Eq. (3.1) the first contribution is from the top quark, the second is from the W boson, the third is from the singly charged Higgs boson, the is fourth from the doubly charged Higgs boson, and the last two are from the vector leptons. We use the following expressions for the couplings of the Higgs bosons with charged vector leptons:

$$\begin{aligned}
\mu_{E_1} &= -\cos\theta_L\cos\theta_R(m_E'\tan\theta_R + m_E''\tan\theta_L), \\
\mu_{E_2} &= \cos\theta_L\cos\theta_R(m_E'\tan\theta_L + m_E''\tan\theta_R).
\end{aligned} \qquad (3.6)$$

The couplings of h to the vector bosons and fermions are as follows:

$$\begin{aligned}
g_{hff} &= \cos\alpha/\cos\beta_\pm, \\
g_{hWW} &= \cos\alpha + 2\sin\alpha v_\Delta/v_\Phi,
\end{aligned} \qquad (3.7)$$

with $ff = t\bar{t}$, $E_1\bar{E}_1$, $E_2\bar{E}_2$, and the scalar trilinear couplings are parametrized as

$$\begin{aligned}
\tilde{g}_{hH^{++}H^{--}} &= \frac{m_W}{g m_{H^{\pm\pm}}^2}[2\lambda_2 v_\Delta\sin\alpha + \lambda_4 v_\Phi\cos\alpha], \\
\tilde{g}_{hH^+H^-} &= \frac{m_W}{2g m_{H^\pm}^2}\{[4v_\Delta(\lambda_2+\lambda_3)\cos^2\beta_\pm \\
&\quad + 2v_\Delta\lambda_4\sin^2\beta_\pm - \sqrt{2}\lambda_5 v_\Phi\cos\beta_\pm\sin\beta_\pm]\sin\alpha \\
&\quad + [4\lambda_1 v_\Phi\sin\beta_\pm^2 + (2\lambda_4+\lambda_5)v_\Phi\cos^2\beta_\pm \\
&\quad + (4\mu - \sqrt{2}\lambda_5 v_\Delta)\cos\beta_\pm\sin\beta_\pm]\cos\alpha\}.
\end{aligned} \qquad (3.8)$$

Since the new leptons do not affect the Higgs production channels, the effect on the diphoton search channel at the LHC is expressed by the ratio

$$\begin{aligned}
R_{h\to\gamma\gamma} &\equiv \frac{\sigma_{\rm HTM}(gg\to h\to\gamma\gamma)}{\sigma_{\rm SM}(gg\to\Phi\to\gamma\gamma)} \\
&= \frac{[\sigma(gg\to h)\times BR(h\to\gamma\gamma)]_{\rm HTM}}{[\sigma(gg\to\Phi)\times BR(\Phi\to\gamma\gamma)]_{\rm SM}} \\
&= \frac{[\sigma(gg\to h)\times\Gamma(h\to\gamma\gamma)]_{\rm HTM}}{[\sigma(gg\to\Phi)\times\Gamma(\Phi\to\gamma\gamma)]_{\rm SM}}\times\frac{[\Gamma(\Phi)]_{\rm SM}}{[\Gamma(h)]_{\rm HTM}},
\end{aligned} \qquad (3.9)$$

where Φ is the SM neutral Higgs boson and where the ratio of the cross sections by gluon fusion is

$$\frac{\sigma_{\rm HTM}(gg\to h\to\gamma\gamma)}{\sigma_{\rm SM}(gg\to\Phi\to\gamma\gamma)} = \cos^2\alpha. \qquad (3.10)$$

Here α is the mixing angle in the CP-even neutral sector,

$$\begin{pmatrix}\varphi\\ \delta\end{pmatrix} = \begin{pmatrix}\cos\alpha & -\sin\alpha\\ \sin\alpha & \cos\alpha\end{pmatrix}\begin{pmatrix}h\\ H\end{pmatrix}, \qquad (3.11)$$

with

$$\tan 2\alpha = \frac{v_\Delta}{v_\Phi}\frac{2v_\Phi^2(\lambda_4+\lambda_5) - 4v_\Phi^2\mu/\sqrt{2}v_\Delta}{2v_\Phi^2\lambda_1 - v_\Phi^2\mu/\sqrt{2}v_\Delta - 2v_\Delta^2(\lambda_2+\lambda_3)}. \qquad (3.12)$$

In Ref. [17] we investigated the Higgs boson decay branching ratio into $\gamma\gamma$ with respect to the SM—assuming that the lightest Higgs boson is the 2.3σ signal excess observed at the Large Electron Positron Collider (LEP) at 98 GeV, while the heavier Higgs boson is the boson observed at the LHC at 125 GeV—in a Higgs triplet model without vector leptons, and found that this is the only scenario which allows for an enhancement of the $h \to \gamma\gamma$ branching fraction. We thus choose the values 125 GeV and 98 GeV for the h and H masses, respectively, and adjust the parameters λ_1–λ_5 accordingly.

Vector lepton masses and mixing parameters depend on M_L and M_E, the explicit mass parameters in the Lagrangian, and h_E', h_E'', the vector-lepton Yukawa parameters. In the limit of vanishing Dirac mass terms

M_L and M_E, the prefactors $\frac{\mu_{E_i}}{M_{E_i}}$ in Eq. (11) go to 1. It then follows that there is destructive interference between the dominant W-boson contribution and the charged leptons loops [11]. In this limit, despite possible enhancement from singly and doubly charged Higgs bosons in the loop, we find a large suppression of the diphoton rate. We present the plots for the relative signal strength of $R_{h\to\gamma\gamma}$, defined in Eq. (3.9) as a function of $m'_E = m''_E$ (or equivalently $h'_E = h''_E$), for various values of the doubly charged Higgs boson mass in the left panel of Fig. 1, for $\sin\alpha = 0$. Clearly, for this case (no mixing) the decay of the h is suppressed significantly with respect to the value in the SM over the whole region of the parameter space in m'_E.

Allowing mixing in the neutral Higgs sector changes the relative contributions of the charged Higgs to the diphoton decay. We show decay rates for $h \to \gamma\gamma$ as a function of $\sin\alpha$ for different values of the doubly charged Higgs boson mass, assuming $m'_E = m''_E = 100$ GeV (and 200 GeV), in the middle and right panels of Fig. 1, respectively. Considerations for relative branching ratios are affected by the fact that the total width of the Higgs boson in the HTM for $\sin\alpha \neq 0$ is not the same as in the SM. The relative widths factor is

$$\frac{[\Gamma(h)]_{\rm HTM}}{[\Gamma(\Phi)]_{\rm SM}} = \frac{[\Gamma(h\to\sum_f f\bar{f}) + \Gamma(h\to WW^*) + \Gamma(h\to ZZ^*) + \Gamma(h\to\nu\nu)]_{\rm HTM}}{[\Gamma(\Phi\to\sum_f f\bar{f}) + \Gamma(\Phi\to WW^*) + \Gamma(\Phi\to ZZ^*)]_{\rm SM}}. \tag{3.13}$$

The plots in Fig. 1 correspond to symmetry condition 1 in Sec. II, that is, $M_L = M_E = 0$.

However, if mixing with SM leptons is forbidden, but the vector leptons are still allowed to mix with each other, the prefactors μ_{E_i}/M_{E_i} in Eq. (3.1) are not 1, and can modify the Higgs diphoton decay. In the next plots we investigate the effect of nonzero mass parameters M_L and M_E, for fixed values of the Yukawa couplings. In Fig. 2 we present the contour plots of constant $R_{h\to\gamma\gamma}$ for $h'_E = h''_E = 0.8$ in the plane of the explicit mass terms M_L and M_E, for various values of the doubly charged Higgs boson mass and $\sin\alpha$. The contours are labeled by the value of $R_{h\to\gamma\gamma}$. The vector-lepton masses are restricted to values for which [13] $M_{E_2} \geq 62.5$ GeV, where $M_{E_{1,2}}$ are given in Eq. (2.9). The plots indicate that it is difficult to obtain any significant enhancement of the ratio $R_{h\to\gamma\gamma}$ for $\sin\alpha = 0$, and this does not depend on the chosen values for the doubly charged Higgs mass; however, for $\sin\alpha \neq 0$, enhancements are possible for various values of $m_{H^{\pm\pm}}$. In Fig. 3 we investigate the dependence of $R_{h\to\gamma\gamma}$ on the Yukawa couplings and vector-lepton masses. We show contour plots for fixed $R_{h\to\gamma\gamma}$ in a $h'_E - M_L$ plane, with $h'_E = h''_E$ and $M_L = M_E$, for various values of $\sin\alpha$ and the doubly charged Higgs boson mass. Enhancements are possible here for all values of $\sin\alpha$, but while for $\sin\alpha = 0$ the decay $h \to \gamma\gamma$ is enhanced for large vector-lepton masses and Yukawa couplings, for $\sin\alpha \neq 0$ we observe enhancements for light vector-lepton masses and small Yukawa couplings.

If we wish to study the light vector-lepton parameter space where $h \to \gamma\gamma$ is enhanced, $\sin\alpha \neq 0$ is preferred. The enhancement is affected by mixing in the vector-lepton sector, the various values for the doubly charged Higgs boson mass and the values of $\sin\alpha$.

As the plots cover only a limited range of the parameter space, in the tables below we give the ranges for the values of the ratio $R_{h\to\gamma\gamma}$ for the various scenarios. In Table I, we fix the value of the Yukawa coupling to $h'_E = 0.8$, allow the vector-lepton masses to vary in the (100–500) GeV range, and show the values for $R_{h\to\gamma\gamma}$ for different values of $\sin\alpha$ and the doubly charged Higgs mass. We note that the relative branching ratios are very sensitive to the values of both the doubly charged Higgs mass and $\sin\alpha$. Enhancements in the branching ratio of $h \to \gamma\gamma$ are possible for light values of $m_{H^{\pm\pm}} \lesssim 300$ GeV, and are much more pronounced at large $\sin\alpha$. Note that for $\sin\alpha = 0$, the result is independent of $m_{H^{\pm\pm}}$, in agreement with the results obtained in Ref. [17]. The reason is the following.

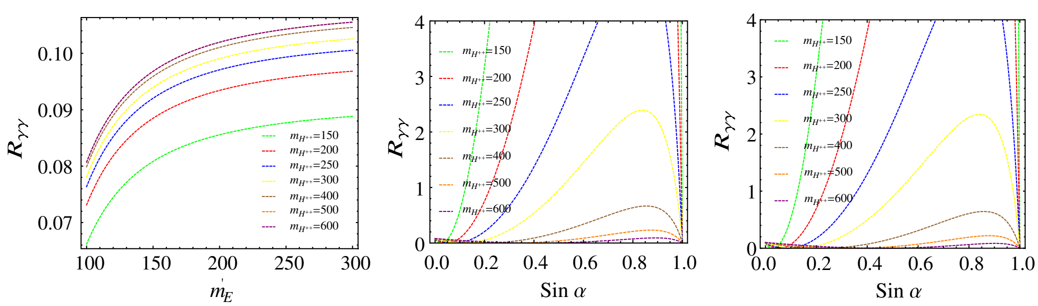

FIG. 1 (color online). Relative decay rate $R_{h\to\gamma\gamma}$ in the limit $M_L = M_E = 0$ (left panel) as a function of $m'_E = m''_E$, for $\sin\alpha = 0$, as a function of $\sin\alpha$ for $m'_E = m''_E = 100$ GeV (middle panel), and as a function of $\sin\alpha$ for $m'_E = m''_E = 200$ GeV (right panel). The colored-coded curves correspond to different values of the doubly charged Higgs mass, given in the attached panels in GeV.

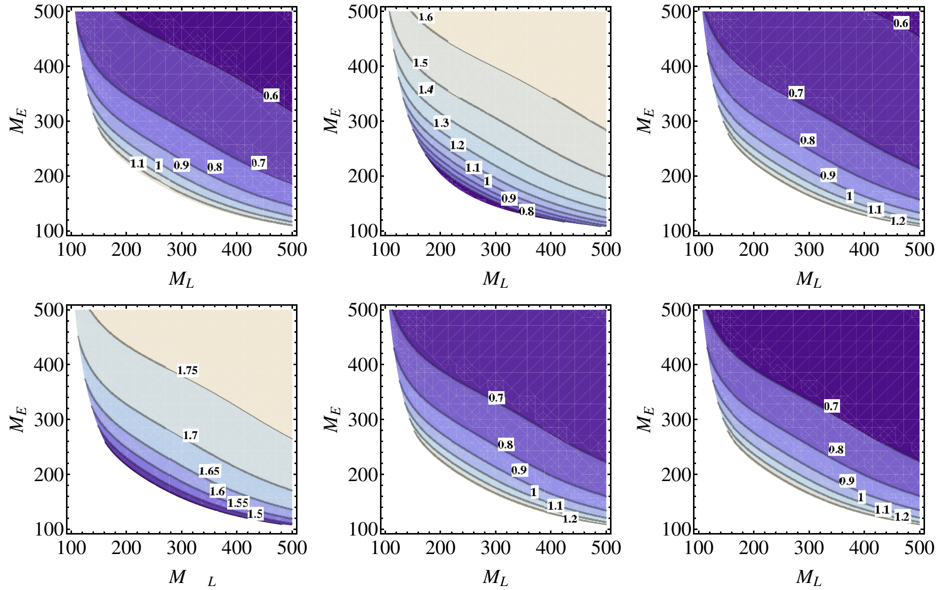

FIG. 2 (color online). Contour plots of constant $R_{h\to\gamma\gamma}$ for mass terms M_E and M_L, for $h'_E = h''_E = 0.8$ and combinations of doubly charged Higgs boson masses and $\sin\alpha$: (upper left panel) $m_{H^{\pm\pm}} = 150$ GeV, $\sin\alpha = 0$; (upper middle panel) $m_{H^{\pm\pm}} = 150$ GeV, $\sin\alpha = 0.2$; (upper right panel) $m_{H^{\pm\pm}} = 300$ GeV, $\sin\alpha = 0$; and (lower left panel) $m_{H^{\pm\pm}} = 300$ GeV, $\sin\alpha = 0.9$; (lower middle panel) $m_{H^{\pm\pm}} = 500$ GeV, $\sin\alpha = 0$; (lower right panel) $m_{H^{\pm\pm}} = 600$ GeV, $\sin\alpha = 0$.

In Eq. (3.8), for $\sin\alpha = 0$, the coupling between the neutral and doubly charged Higgses is

$$\tilde{g}_{hH^{++}H^{--}} = \frac{m_W}{g m^2_{H^{\pm\pm}}}[\lambda_4 v_\Phi] = \frac{m_W}{g m^2_{H^{\pm\pm}}}\left[2\frac{m^2_{H^{\pm\pm}}}{v^2_\Phi} v_\Phi\right] = 2\frac{m_W}{g v_\Phi}, \tag{3.14}$$

where we used the expression for λ_4 from Ref. [17], which is independent of $m_{H^{\pm\pm}}$. In Table II we allow—in addition to mass variations,—variations in the Yukawa coupling $h'_E \in (0\text{–}3)$. This means allowing both explicit (Dirac) masses and additional contributions by electroweak symmetry breaking, m'_E, m''_E. The dependence on the Yukawa coupling h'_E is much weaker than on $\sin\alpha$ or on $m_{H^{\pm\pm}}$.

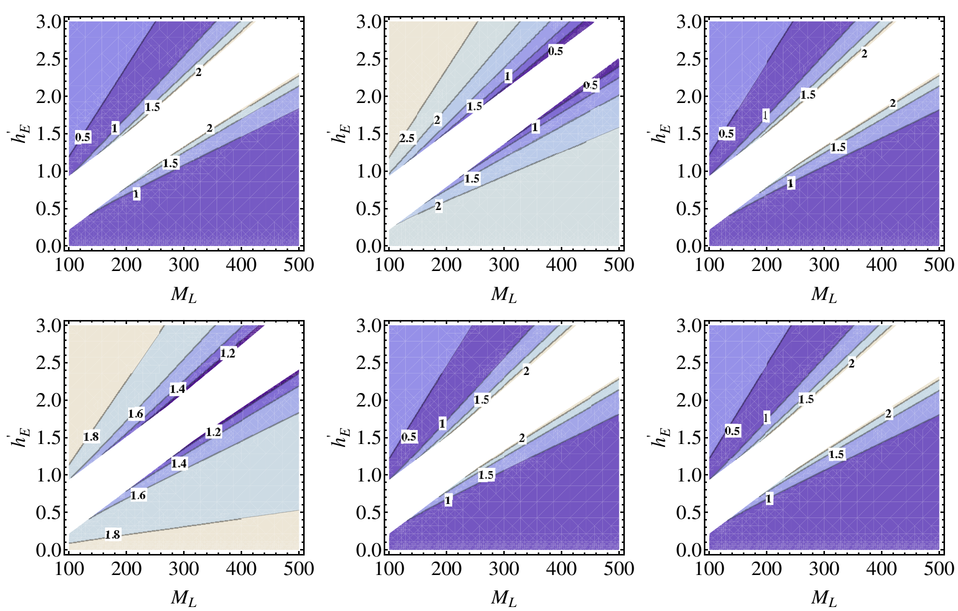

FIG. 3 (color online). Contour plots of constant $R_{h\to\gamma\gamma}$ for mass terms M_L and $h'_E = h''_E$, for various doubly charged Higgs boson masses and $\sin\alpha$: (upper left panel) $m_{H^{\pm\pm}} = 200$ GeV, $\sin\alpha = 0$; (upper middle panel) $m_{H^{\pm\pm}} = 200$ GeV, $\sin\alpha = 0.4$; (upper right panel) $m_{H^{\pm\pm}} = 300$ GeV, $\sin\alpha = 0$; (lower left panel) $m_{H^{\pm\pm}} = 300$ GeV, $\sin\alpha = 0.9$; (lower middle panel) $m_{H^{\pm\pm}} = 500$ GeV, $\sin\alpha = 0$; (lower right panel) $m_{H^{\pm\pm}} = 600$ GeV, $\sin\alpha = 0$.

TABLE I. Range of the ratio $R_{h\to\gamma\gamma}$, as defined in the text, for the doubly charged Higgs mass (in columns) and the neutral Higgs mixing angle $\sin\alpha$ (in rows), for Dirac vector-lepton masses in the range M_E, $M_L \in$ (100–500) GeV, with $h'_E = 0.8$.

$R_{\gamma\gamma}$	$m_{H^{\pm\pm}} = 150$ GeV	200 GeV	250 GeV	300 GeV	500 GeV	600 GeV
$\sin\alpha = 0$	0.6–1.2	0.6–1.2	0.6–1.2	0.6–1.2	0.6–1.2	0.6–1.2
$\sin\alpha = 0.1$	0.02–0.08	0.05–0.23	0.2–0.5	0.3–0.7	0.5–1	0.6–1
$\sin\alpha = 0.2$	0.8–1.6	0.02–0.14	0.01–0.08	0.1–0.3	0.4–0.8	0.5–0.9
$\sin\alpha = 0.3$	4–5.2	0.2–0.9	0.02–0.12	0.01–0.04	0.25–0.55	0.3–0.7
$\sin\alpha = 0.4$	9–10.75	1.4–2.2	0.1–0.5	0.02–0.08	0.15–0.35	0.25–0.55
$\sin\alpha = 0.5$	16–18	3.2–4.2	0.6–1.2	0.05–0.3	0.06–0.22	0.15–0.35
$\sin\alpha = 0.6$	24–26.5	5.6–6.8	1.4–2.1	0.25–0.7	0.01–0.08	0.06–0.2
$\sin\alpha = 0.7$	32.5–34.5	8.2–9.4	2.4–3.1	0.7–1.1	0.002–0.008	0.02–0.08
$\sin\alpha = 0.8$	38.5–40.75	10.4–11.4	3.4–4.1	1.2–1.65	0.005–0.045	0.001–0.006
$\sin\alpha = 0.9$	36.2–37.4	10.2–10.9	3.7–4.1	1.5–1.75	0.04–0.1	0.005–0.02

TABLE II. Same as in Table I, but also allowing $h'_E \in$ (0–3).

$R_{\gamma\gamma}$	$m_{H^{\pm\pm}} = 150$ GeV	200 GeV	250 GeV	300 GeV	500 GeV	600 GeV
$\sin\alpha = 0$	0.5–2	0.5–2	0.5–2	0.5–2	0.5–2	0.5–2
$\sin\alpha = 0.1$	0.05–0.3	0.1–0.6	0.2–1	0.2–1.4	0.5–1.75	0.5–2
$\sin\alpha = 0.2$	0.5–2	0.1–0.4	0.05–0.3	0.1–0.7	0.25–1.5	0.25–1.75
$\sin\alpha = 0.3$	2–6	0.5–1	0.1–0.4	0.05–0.2	0.2–1.2	0.2–1.4
$\sin\alpha = 0.4$	6–11	1–2.5	0.2–0.6	0.05–0.3	0.2–0.8	0.2–1
$\sin\alpha = 0.5$	14–18	2–4	0.5–1.5	0.2–0.4	0.1–0.5	0.2–0.8
$\sin\alpha = 0.6$	20–26	4–7	0.5–2.5	0.25–0.75	0.05–0.25	0.1–0.5
$\sin\alpha = 0.7$	30–36	7–10	1.5–3.5	0.25–1.25	0.01–0.07	0.05–0.2
$\sin\alpha = 0.8$	36–40	9–12	2.5–4	0.5–1.75	0.02–0.08	0.01–0.04
$\sin\alpha = 0.9$	35–38	9.5–11	3.2–4.2	1.2–1.8	0.05–0.1	0.01–0.04

However, one can see from the tables that modest enhancements of the ratio $R_{h\to\gamma\gamma}$ are possible for $\sin\alpha = 0$ for large vector-lepton Yukawa couplings, unlike in the case of the triplet model without vector leptons. This would then be a clear distinguishing feature, namely enhancements of the decay $h \to \gamma\gamma$ in the absence of mixing in the neutral sector. The absence of mixing would manifest itself in observing tree-level decays ($h \to b\bar{b}$, $\tau^+\tau^-$, ZZ^* and WW^*) identical to those in the SM. There seems to be a minimum value of $R_{h\to\gamma\gamma}$ for $\sin\alpha = 0.1$, where the contribution from the doubly charged Higgs bosons is important for small doubly charged masses and counters the contribution from the vector leptons. This is a suppression of the branching ratio for $h \to \gamma\gamma$, which is due to the fact that the vector-lepton contribution interacts destructively with the dominant $W^\pm$ contribution. As a general feature, $R_{h\to\gamma\gamma}$ increases when we lower the doubly charged Higgs mass and increase $\sin\alpha$. This rules out part of the parameter space. For instance, for $m_{H^{\pm\pm}} = 150$ GeV, the mixing cannot be larger than $\sin\alpha = 0.2$, and for $m_{H^{\pm\pm}} = 200$ GeV, mixings larger than $\sin\alpha \geq 0.5$ are ruled out. If the value of $m_{H^{\pm\pm}}$ is increased to 500–600 GeV, only modest enhancements are possible, and only for $\sin\alpha = 0$, for vector-lepton explicit masses in the 100–500 GeV range and $h'_E = 0.8$. Increasing the vector-lepton Yukawa coupling increases the overall ratio $R_{h\to\gamma\gamma}$.

B. $h \to Z\gamma$

In most models, the $h \to \gamma\gamma$ and $h \to Z\gamma$ partial decay widths are correlated or anticorrelated, though usually the enhancement/suppression in the $Z\gamma$ channel is much smaller compared to that in the $\gamma\gamma$ channel. However, as in models with new loop contributions to $h \to \gamma\gamma$, $Z\gamma$, a sensitivity to both is expected; we study the correlation between the two here, in the presence of vector leptons. An investigation of the branching ratio of $h \to Z\gamma$ is also further justified by the recent results from CMS and ATLAS [23], which indicate branching fractions consistent with the SM expectation at 1σ in the Higgs boson h mass region at 95% C.L. The decay width for $h \to Z\gamma$ is given by [24]

$$[\Gamma(h\to Z\gamma)]_{\rm HTM} = \frac{\alpha G_F^2 m_W^2 m_h^3}{64\pi^4}\left(1-\frac{m_Z^2}{m_h^2}\right)^3 \Bigg| \frac{1}{c_W}\sum_f 2N_c^f Q_f (I_3^f - 2Q_f s_W^2) g_{hff} A_{1/2}^h(\tau_f^h, \tau_f^Z)$$
$$+\frac{(I_3^E - 2Q_E s_W^2)(2Q_E)}{c_W}\left[\frac{\mu_{E_1} g_{hff}}{M_{E_1}} A_{1/2}(\tau_{E_1}^h, \tau_{E_1}^Z) + \frac{\mu_{E_2} g_{hff}}{M_{E_2}} A_{1/2}^h(\tau_{E_2}^h, \tau_{E_2}^Z)\right]$$
$$+ c_W g_{hWW} A_1^h(\tau_W^h, \tau_W^Z) - 2 s_W \tilde{g}_{hH^\pm H^\mp} g_{ZH^\pm H^\mp} A_0^h(\tau_{H^\pm}^h, \tau_{H^\pm}^Z) - 4 s_W \tilde{g}_{hH^{\pm\pm}H^{\mp\mp}} g_{ZH^{\pm\pm}H^{\mp\mp}} A_0^h(\tau_{H^{\pm\pm}}^h, \tau_{H^{\pm\pm}}^Z)\Bigg|^2, \tag{3.15}$$

where $\tau_i^h = 4m_i^2/m_h^2$, $\tau_i^Z = 4m_i^2/m_Z^2$ [with $i = f(\equiv t)$, E_1, E_2, W, $H^\pm$, $H^{\pm\pm}$], and the loop factors are given by

$$A_0^h(\tau_h, \tau_Z) = I_1(\tau^h, \tau^Z), \qquad A_{1/2}^h(\tau^h, \tau^Z) = I_1(\tau^h, \tau^Z) - I_2(\tau^h, \tau^Z),$$
$$A_1^h(\tau^h, \tau^Z) = 4(3 - \tan^2\theta_W) I_2(\tau^h, \tau^Z) + [(1 + 2\tau^{h-1})\tan^2\theta_W - (5 + 2\tau^{h-1})] I_1(\tau^h, \tau^Z). \tag{3.16}$$

The functions I_1 and I_2 are given by

$$I_1(\tau^h, \tau^Z) = \frac{\tau^h \tau^Z}{2(\tau^h - \tau^Z)} + \frac{\tau^{h2}\tau^{Z2}}{2(\tau^h - \tau^Z)^2}[f(\tau^{h-1}) - f(\tau^{Z-1})] + \frac{\tau^{h2}\tau^Z}{(\tau^h - \tau^Z)^2}[g(\tau^{h-1}) - g(\tau^{Z-1})],$$
$$I_2(\tau^h, \tau^Z) = -\frac{\tau^h\tau^Z}{2(\tau^h - \tau^Z)}[f(\tau^{h-1}) - f(\tau^{Z-1})], \tag{3.17}$$

where the function $f(\tau)$ is defined in Eq. (3.5), and the function $g(\tau)$ is defined as

$$g(\tau) = \begin{cases} \sqrt{\tau^{-1} - 1}\sin^{-1}(\sqrt{\tau}), & (\tau < 1), \\ \frac{1}{2}\sqrt{1 - \tau^{-1}}\left[\log\left(\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}}\right) - i\pi\right], & (\tau \geq 1). \end{cases} \tag{3.18}$$

In Eq. (3.15) we list, in order, the ordinary-lepton, vector-lepton, W-boson, singly charged Higgs, and doubly charged Higgs contributions. The scalar couplings $g_{hf\bar{f}}$ and $g_{hW^+W^-}$ are given in Eq. (3.7), and the scalar trilinear couplings $\tilde{g}_{hH^\pm H^\mp}$ and $\tilde{g}_{hH^{\pm\pm}H^{\mp\mp}}$ are given in Eq. (3.8). The remaining couplings in Eq. (3.15) are given by

$$g_{ZH^+H^-} = -\tan\theta_W, \qquad g_{ZH^{++}H^{--}} = 2\cot 2\theta_W. \tag{3.19}$$

We proceed to perform a similar analysis as in Sec. III A. We show first the variation of the branching ratio $h\to Z\gamma$ with the mass $m_E' = m_E''$, for various values of the doubly charged Higgs mass, for the case of no mixing in the neutral sector, for both $\sin\alpha = 0$ (shown in the left panel of Fig. 4) and as a function of the mixing angle $\sin\alpha$ for $m_E' = 100$ GeV and $m_E' = 200$ GeV in the middle and right panels of Fig. 4, respectively. We have chosen the same parameter values as in Fig. 1 for comparison. It is clear that the branching ratio into $Z\gamma$ is fairly independent of both m_E' and $m_{H^{\pm\pm}}$, and is always just below the SM expectations. Note that the severe suppression seen in $h\to\gamma\gamma$ for $\sin\alpha = 0$ (Fig. 1, left panel) does not occur here, and the results of the left panel of Fig. 4 are consistent with the data at the LHC.

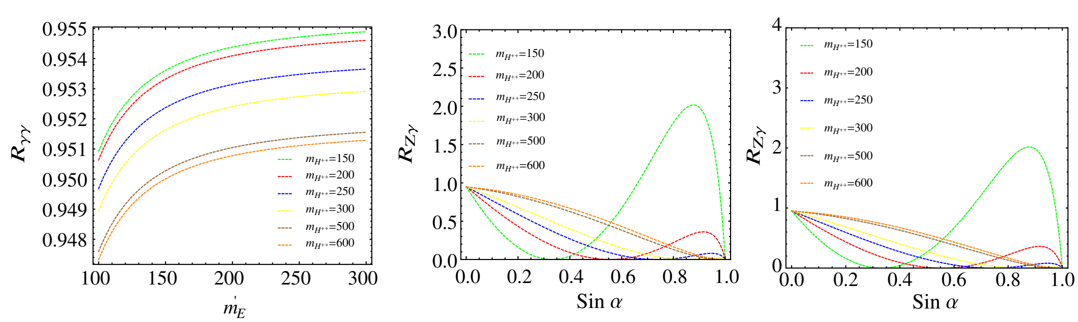

FIG. 4 (color online). Relative decay rate for $R_{h\to\gamma Z}$ as a function of $m_E' = m_E''$ for different values of doubly charged Higgs masses, in the case of no mixing, i.e., $\sin\alpha = 0$ (left panel); and as a function of $\sin\alpha$ for $m_E' = m_E'' = 100$ GeV (middle panel), and $m_E' = m_E'' = 200$ GeV (right panel). The colored-coded curves correspond to different values of doubly charged Higgs masses, given in the attached panels in GeV.

TABLE III. Range of the ratio $R_{h\to Z\gamma}$, as defined in the text, for the doubly charged Higgs mass (in columns) and the neutral Higgs mixing angle $\sin\alpha$ (in rows), for Dirac vector-lepton masses in the range M_E, $M_L \in$ (100–500) GeV, with $h'_E = 0.8$.

$R_{Z\gamma}$	$m_{H^{\pm\pm}} = 150$ GeV	200 GeV	250 GeV	300 GeV	500 GeV	600 GeV
$\sin\alpha = 0$	0.96–1	0.96–1	0.96–1	0.96–1	0.96–1	0.96–1
$\sin\alpha = 0.1$	0.48–0.51	0.68–0.72	0.78–0.82	0.83–0.87	0.91–0.94	0.92–0.96
$\sin\alpha = 0.2$	0.16–0.18	0.44–0.47	0.6–0.63	0.69–0.73	0.83–0.87	0.85–0.89
$\sin\alpha = 0.3$	0.01–0.015	0.24–0.26	0.43–0.45	0.55–0.58	0.74–0.78	0.78–0.81
$\sin\alpha = 0.4$	0.03–0.04	0.09–0.10	0.27–0.3	0.41–0.43	0.63–0.66	0.68–0.71
$\sin\alpha = 0.5$	0.25–0.26	0.01–0.02	0.14–0.16	0.27–0.29	0.52–0.54	0.56–0.59
$\sin\alpha = 0.6$	0.64–0.66	0.005–0.006	0.05–0.06	0.15–0.17	0.39–0.41	0.43–0.46
$\sin\alpha = 0.7$	1.18–1.21	0.07–0.08	0.005–0.007	0.06–0.07	0.25–0.27	0.3–0.32
$\sin\alpha = 0.8$	1.76–1.78	0.21–0.22	0.008–0.01	0.009–0.011	0.13–0.14	0.17–0.18
$\sin\alpha = 0.9$	1.98–1.99	0.35–0.36	0.06	0.004–0.005	0.03–0.04	0.05–0.06

TABLE IV. Same as in Table III, but also allowing $h'_E \in$ (0–3).

$R_{Z\gamma}$	$m_{H^{++}} = 150$ GeV	200 GeV	250 GeV	300 GeV	500 GeV	600 GeV
$\sin\alpha = 0$	0.94–1.04	0.94–1.04	0.94–1.04	0.94–1.04	0.94–1.04	0.94–1.04
$\sin\alpha = 0.1$	0.47–0.54	0.66–0.74	0.76–0.84	0.8–0.9	0.92–0.98	0.9–1
$\sin\alpha = 0.2$	0.16–0.2	0.43–0.49	0.6–0.7	0.68–0.74	0.82–0.9	0.8–0.9
$\sin\alpha = 0.3$	0.01–0.018	0.23–0.27	0.42–0.47	0.54–0.6	0.74–0.8	0.76–0.84
$\sin\alpha = 0.4$	0.03–0.05	0.09–0.12	0.27–0.31	0.4–0.5	0.63–0.68	0.67–0.73
$\sin\alpha = 0.5$	0.23–0.27	0.01–0.02	0.14–0.17	0.27–0.3	0.51–0.56	0.56–0.61
$\sin\alpha = 0.6$	0.6–0.7	0.001–0.009	0.05–0.07	0.15–0.18	0.38–0.42	0.43–0.47
$\sin\alpha = 0.7$	1.16–1.22	0.07–0.08	0.004–0.008	0.06–0.08	0.26–0.28	0.3–0.33
$\sin\alpha = 0.8$	1.73–1.79	0.2–0.23	0.006–0.01	0.008–0.01	0.13–0.15	0.17–0.18
$\sin\alpha = 0.9$	1.95–2	0.34–0.36	0.06–0.07	0.004–0.005	0.03–0.04	0.05–0.06

But the variation with the mixing angle α is pronounced, and the branching ratio can reach almost twice its SM value for $\sin\alpha \sim 0.8$. However—correlated with our predictions from Sec. III A and LHC measurements for $R_{h\to\gamma\gamma}$—the parameter space corresponding to an enhanced $h \to Z\gamma$, for both $m'_E = 100$ GeV and 200 GeV, for a doubly charged Higgs mass $m_{H^{\pm\pm}} = 150$ GeV is ruled out. For all other values considered, the value for $R_{h\to Z\gamma}$ is close to or below the SM expectations. This is a general prediction of the model.

For a large range of parameter space, the decay $h \to Z\gamma$ can be suppressed significantly with respect to the SM. We plot decay rates for $h \to Z\gamma$ as a function of $\sin\alpha$ for different values of the doubly charged Higgs boson mass, assuming $m'_E = m''_E = 100$ GeV and 200 GeV, in the middle and right panels of Fig. 4, respectively. Again, considerations for relative branching ratios are affected by the fact that the total width of the Higgs boson in the HTM is not the same as in the SM. The widths are the same as those in the SM for h for $\sin\alpha = 0$, while for $\sin\alpha \neq 0$ we take into account the relative width factors [Eq. (3.13)].

In Tables III and IV we present the explicit ranges of $R_{h\to Z\gamma}$ for varying values of M_E, M_L and for a range of h'_E parameters. We choose a fixed value for $h'_E = 0.8$ in Table III, as this is the preferred choice from other analyses [11,13] and to facilitate a comparison with Table I. A comparison of Tables I and III shows that the decay $h \to Z\gamma$ is far more stable against variations in masses and values for $\sin\alpha$ than $h \to \gamma\gamma$, making it a less sensitive indicator for the presence of vector-lepton states.

In Table IV we also allow variations in the Yukawa coupling $h'_E \in$ (0–3). As before this amounts to allowing both explicit and contributions from electroweak symmetry breaking, m'_E, m''_E, to vector-lepton masses. A comparison of Tables III and IV indicates that the results are not very sensitive to variations in the Yukawa coupling h'_E or the vector-lepton mass parameters M_E, M_L. However, the relative branching ratios are very sensitive to values of $\sin\alpha$. While the branching ratio into $Z\gamma$ is almost always suppressed with respect to its SM value, there is a small region of the parameter space—with a light $H^{\pm\pm}$ and $\sin\alpha \simeq 0.7$–0.9—where enhancement is possible; however, as discussed before, this region is ruled out by constraints from $h \to \gamma\gamma$ measurements (Table II). Note that for $\sin\alpha = 0$ the branching ratio is (as before) independent of the mass of $H^{\pm\pm}$ and is about the same as in the SM.

IV. PRODUCTION AND DECAYS OF THE DOUBLY CHARGED HIGGS BOSONS

The discovery of the doubly charged Higgs bosons would be one of the most striking signals of physics beyond the SM, and a clear signature for the Higgs triplet model. From theoretical expectations, the decay modes of $H^{\pm\pm}$ depend on the value of the VEV of the neutral triplet Higgs component, v_Δ. When $v_\Delta \leq 0.1$ MeV, the dominant decay mode of $H^{\pm\pm}$ is into lepton pairs. If $v_\Delta \gg 0.1$ MeV, the main decay modes of $H^{\pm\pm}$ are into $W^{\pm(\star)}W^{\pm(\star)}$,[1] and into $H^{\pm}W^{\pm(\star)}$, if kinematically allowed.

We briefly summarize the results of the experimental constraints on doubly charged bosons. Searches for $H^{\pm\pm}$ were performed at LEP [25], the Hadron Electron Ring Accelerator [26] and the Tevatron [27]. The most up-to-date bounds have been recently derived by the ATLAS and CMS collaborations at the LHC. Assuming a Drell-Yan-like pair production, these collaborations have looked for long-lived doubly charged states, and after analyzing 5 fb^{-1} of LHC collisions at a center-of-mass energy $\sqrt{s} = 7$ TeV, and 18.8 fb^{-1} of collisions at $\sqrt{s} = 8$ TeV, they constrained the masses to lie above 685 GeV [28]. The assumption is that the doubly charged Higgs bosons decay 100% into a pair of leptons with the same electric charge through Majorana-type interactions [19], thus neglecting possible decays into W-boson pairs, which is shown to alter the pattern of $H^{\pm\pm}$ branching fractions [29]. In this work, we allow decays into $W^\pm W^\pm$ bosons, and also include the decays into vector leptons, which, if light enough, would modify the decays of the doubly charged Higgs bosons further. We take $v_\Delta = 1$ GeV throughout our considerations.[2]

The main production mode for doubly charged bosons $H^{\pm\pm}$ is the pair production $pp \to \gamma^*, Z^* \to H^{\pm\pm}H^{\mp\mp}$ and the associated production $pp \to W^{\pm *} \to H^{\pm\pm}H^{\mp}$. The production cross sections for both the vector-boson fusion $qQ \to q'Q'H^{\pm\pm}$ and the associated weak-boson production $qQ \to W^{\pm *} \to H^{\pm\pm}W^{\mp}$ are proportional to v_Δ^2, and are much less significant for $v_\Delta \ll v_\Phi$.

At hadron colliders the partonic cross section for the leading order (LO) production cross section for a doubly charged Higgs boson pair is

$$\hat{\sigma}_{\rm LO}(q\bar{q} \to H^{\pm\pm}H^{\mp\mp}) = \frac{\pi\alpha^2}{9Q^2}\beta^3\left[4e_q^2 + \frac{2e_q v_q v_{H^{\pm\pm}}(1 - M_Z^2/Q^2) + (v_q^2 + a_q^2)v_{H^{\pm\pm}}^2}{(1 - M_Z^2/Q^2)^2 + M_Z^2\Gamma_Z^2/Q^4}\right], \tag{4.1}$$

where we have defined

$$v_q = \frac{2I_{3q} - 4e_q \sin^2\theta_W}{\sin 2\theta_W}, \qquad a_q = \frac{2I_{3q}}{\sin 2\theta_W},$$
$$v_{H^{\pm\pm}} = \frac{2I_{3H^{\pm\pm}} - 4\sin^2\theta_W}{\sin 2\theta_W},$$

with I_{3i} being the third component of the isospin for particle i, $Q^2 = \hat{s}$ the square of the partonic center-of-mass energy, $\beta = \sqrt{1 - 4m_{H^{\pm\pm}}^2/Q^2}$, and α the QED coupling constant evaluated at the scale Q. The hadronic cross section is obtained by convolution with the partonic density functions of the proton,

$$\sigma_{\rm LO}(pp \to H^{\pm\pm}H^{\mp\mp}) = \int_{\tau_0}^{1} d\tau \sum_q \frac{d\mathcal{L}^{q\bar{q}}}{d\tau}\hat{\sigma}_{\rm LO}(Q^2 = \tau s), \tag{4.2}$$

where $\mathcal{L}^{q\bar{q}}$ is the parton luminosity and $\tau_0 = 4m_{H^{\pm\pm}}^2/s$ (s is the total energy squared at the LHC). The cross section for pair production, including next-to-leading-order (NLO) corrections, has been evaluated in Ref. [30].

Depending on the mass parameters in the model, the doubly charged Higgs boson can decay into lepton pairs, including vector leptons, $W^\pm$ pairs, or $H^\pm W^\pm$ states. In the Higgs triplet model, the decay rate for $H^{\pm\pm}$ into leptons is

$$\Gamma(H^{\pm\pm} \to l_i^\pm l_j^\pm) = S_{ij}|h_{ij}|^2\frac{m_{H^{\pm\pm}}}{4\pi}\left(1 - \frac{m_i^2}{m_{H^{\pm\pm}}^2} - \frac{m_j^2}{m_{H^{\pm\pm}}^2}\right) \times \left[\lambda\left(\frac{m_i^2}{m_{H^{\pm\pm}}^2}, \frac{m_j^2}{m_{H^{\pm\pm}}^2}\right)\right]^2, \tag{4.3}$$

where m_i is the mass of the ith lepton ($i = e$, μ or τ) and $S_{ij} = 1$ (1/2) for $i \neq j$ ($i = j$). Similarly the decay rate of $H^{\pm\pm}$ into fourth generation vector leptons is, if kinematically allowed,

$$\Gamma(H^{\pm\pm} \to E_i^\pm E_j^\pm) = S_{ij}[|h'_{E_iE_j}|^2 + |h''_{E_iE_j}|^2]\frac{m_{H^{\pm\pm}}}{4\pi}\left(1 - \frac{m_{E_i}^2}{m_{H^{\pm\pm}}^2} - \frac{m_{E_j}^2}{m_{H^{\pm\pm}}^2}\right) \times \left[\lambda\left(\frac{m_{E_i}^2}{m_{H^{\pm\pm}}^2}, \frac{m_{E_j}^2}{m_{H^{\pm\pm}}^2}\right)\right]^2, \tag{4.4}$$

where M_{E_i} is the mass eigenvalue from Eq. (2.9). In addition, we include the decay rates of $H^{\pm\pm}$ into $W^\pm W^\pm$ and $W^\pm H^\pm$,

[1]For the present analysis, the mass of the doubly charged Higgs boson will be such that decays into on-shell $W^\pm$ pairs are kinematically allowed.

[2]This value of v_Δ is small enough to satisfy electroweak precision conditions, but large enough to allow decays into gauge and charged Higgs bosons [17,20].

$$\Gamma(H^{\pm\pm} \to W^{\pm}W^{\pm}) = \frac{g^4 v_\Delta^2 m_{H^{\pm\pm}}^3}{64\pi m_W^4}\left(1 - \frac{4m_W^2}{m_{H^{\pm\pm}}^2} + \frac{12m_W^4}{m_{H^{\pm\pm}}^4}\right)\beta\left(\frac{m_W^2}{m_{H^{\pm\pm}}^2}\right), \tag{4.5}$$

$$\Gamma(H^{\pm\pm} \to W^{\pm}H^{\pm}) = \frac{g^2 m_{H^{\pm\pm}}^3}{16\pi m_W^2}\cos^2\beta_{\pm}\left[\lambda\left(\frac{m_W^2}{m_{H^{\pm\pm}}^2}, \frac{m_{H^{\pm}}^2}{m_{H^{\pm\pm}}^2}\right)\right]^{3/2}, \tag{4.6}$$

where $\cos\beta_{\pm} \simeq 1$ is the mixing angle in the singly charged Higgs sector and

$$\beta(x) = \sqrt{1-4x}, \qquad \lambda(x,y) = 1 + x^2 + y^2 - 2xy - 2x - 2y. \tag{4.7}$$

We investigate the branching ratios of $H^{\pm\pm}$ in two distinct parameter regions:

(i) Condition 1: When $H^{\pm\pm} \to W^{\pm}H^{\pm}$ is not kinematically allowed, $H^{\pm\pm}$ decays into leptons and $W^{\pm}$ pairs only,

$$BR(X_i^{\pm}X_j^{\pm}) = \frac{\Gamma(H^{\pm\pm} \to X_i^{\pm}X_j^{\pm})}{\Gamma_1(H^{\pm\pm})}, \quad \text{where} \quad X_i = l_i^{\pm}, E_i^{\pm}, W^{\pm}, \tag{4.8}$$

with the total width for Condition 1 being

$$\Gamma_1(H^{\pm\pm}) = \Gamma(H^{\pm\pm} \to l_i^{\pm}l_j^{\pm}) + \Gamma(H^{\pm\pm} \to E_i^{\pm}E_j^{\pm}) + \Gamma(H^{\pm\pm} \to W^{\pm}W^{\pm}).$$

(ii) Condition 2: When $H^{\pm\pm} \to W^{\pm}H^{\pm}$ is kinematically allowed,[3] $H^{\pm\pm}$ is able to decay into charged Higgs and gauge bosons as well,

$$BR(X_i^{\pm}X_j^{\pm}) = \frac{\Gamma(H^{\pm\pm} \to X_i^{\pm}X_j^{\pm})}{\Gamma_2(H^{\pm\pm})}, \quad \text{where } X_i = l_i^{\pm}, E_i^{\pm}, W^{\pm}, \quad \text{and} \quad BR(W^{\pm}H^{\pm}) = \frac{\Gamma(H^{\pm\pm} \to H^{\pm}W^{\pm})}{\Gamma_2(H^{\pm\pm})}, \tag{4.9}$$

with the total decay width for Condition 2 being

$$\Gamma_2(H^{\pm\pm}) = \Gamma(H^{\pm\pm} \to l_i^{\pm}l_j^{\pm}) + \Gamma(H^{\pm\pm} \to E_i^{\pm}E_j^{\pm}) + \Gamma(H^{\pm\pm} \to W^{\pm}W^{\pm}) + \Gamma(H^{\pm\pm} \to H^{\pm}W^{\pm}).$$

In what follows, we wish to analyze the decay patterns of $H^{\pm\pm}$ and present plots of the production cross section times the branching fractions for large regions of the allowed parameter space, for the LHC operating at both

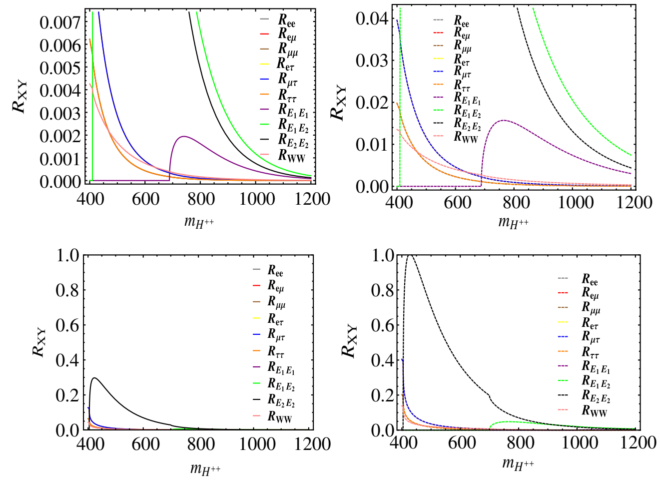

FIG. 5 (color online). $R_{XY} = \sigma(pp \to H^{\pm\pm}H^{\mp\mp})BR(H^{\pm\pm} \to XY)$, in fb, as a function of the doubly charged Higgs boson mass at $\sqrt{s} = 7$ TeV (left panels) and $\sqrt{s} = 13$ TeV (right panels), for Condition 1 (when $H^{\pm\pm} \to W^{\pm}H^{\pm}$ is not kinematically allowed). The upper and lower panels depict the values for Case A and Case B, respectively. Cross sections include the QCD NLO correction factor $K \simeq 1.25$. Throughout we take $h'_{EE} = h''_{EE} = 0.1$ and $h_{ij} = 0.01$.

$\sqrt{s} = 7$ TeV (where analyses of the existing data still continue) and at $\sqrt{s} = 13$ TeV (the next energy frontier).[4] To cover a wide range of parameter space, we distinguish two cases for each condition, depending on the vector-lepton masses. We set the Yukawa coupling of the vector leptons with the doublet Higgs bosons to be $h'_E = h''_E = 0.8$ for both cases, and impose symmetry condition 3, that is, we only disallow mixing of vector and ordinary leptons.

(i) Case A corresponds to very light vector leptons: $M_E = M_L = 205$ GeV. For this case we obtain for the mass eigenvalues $M_{E_1} = 344.2$ GeV, $M_{E_2} = 65.8$ GeV, the latter of which is close to the allowed minimum.

(ii) Case B corresponds to intermediate-mass vector leptons: $M_E = 400$ GeV, $M_L = 300$ GeV. For this case we obtain for the mass eigenvalues $M_{E_1} = 498$ GeV, $M_{E_2} = 202$ GeV.

In Fig. 5 we plot the graphs corresponding to Condition 1 (when the decay $H^{\pm\pm} \to W^{\pm}H^{\pm}$ is not kinematically allowed), with Case A in the top row and Case B in the bottom row. We plot $R_{XY} = \sigma(pp \to H^{\pm\pm}H^{\mp\mp}) \times BR(H^{\pm\pm} \to XY)$ with X, Y as specified in the attached panels, as functions of the doubly charged masses. In evaluating the cross sections we have included the NLO correction factor $K \simeq 1.25$, as calculated in Ref. [30]. On the left and right sides of the figure we show the results for $\sqrt{s} = 7$ TeV and $\sqrt{s} = 13$ TeV, respectively. We set

[3]In this version of the HTM, the doubly charged Higgs boson is always heavier than the singly charged one [17].

[4]Cross sections and branching ratios at $\sqrt{s} = 8$ TeV and 14 TeV are practically indistinguishable from those at 7 TeV and 13 TeV, respectively.

$h'_{EE} = h''_{EE} = 0.1$ and $h_{ij} = 0.01$ throughout. Similar graphs would be obtained with smaller values of h'_{EE} and h''_{EE}, but the values of R_{XY} would be correspondingly reduced. We have chosen to investigate the cross section times branching ratios for intermediate to high values of the doubly charged Higgs mass (400–1200 GeV), as in this region the decay into vector leptons becomes relevant. If the masses of the vector leptons are low, the doubly charged Higgs boson decays significantly into them; in particular, the decay into the two lightest vector leptons E_2 becomes dominant, and can reach 80–90% when kinematically allowed (in the $m_{H^{\pm\pm}} > 200$ GeV region for $M_{E_2} = 65.8$ GeV, Case A) and overwhelms the other decay modes, which are now below 5%. The branching ratio into $W^\pm W^\mp$ is important for doubly charged masses below the threshold for pair production of vector leptons, $m_{H^{\pm\pm}} <$ 400 GeV, and becomes negligible for $m_{H^{\pm\pm}} > 700$ GeV.

The difference between the figures at $\sqrt{s} = 7$ TeV and $\sqrt{s} = 13$ TeV is in the total cross section for pair production of doubly charged bosons, which is expected to be $\simeq$1.5 fb at $\sqrt{s} = 7$ TeV and $\simeq$ 8 fb at $\sqrt{s} = 13$ TeV (for $m_{H^{\pm\pm}} = 400$ GeV). The integrated luminosity was taken to be $\mathcal{L} = 10\ \mathrm{fb}^{-1}$ [31].

The differences between Case A and Case B in this figure are threshold effects. For Case B, $m_{H^{\pm\pm}} >$ 400 GeV for decay into pairs of E_2 states, as $M_{E_2} =$ 202 GeV; otherwise, the branching ratios are the same. Had we chosen smaller doubly charged Higgs couplings with vector leptons, $h'_{EE} \simeq 0.01$, the branching ratios into ordinary leptons, vector leptons, and $W^\pm$ pairs would be comparable for $m_{H^{\pm\pm}} \gtrsim 600$ GeV.

In Fig. 6 we plot the same quantities for Condition 2 (when $H^{\pm\pm} \to W^\pm H^\pm$ is kinematically allowed), with Case A in the top row and Case B in the bottom row, for $\sqrt{s} = 7$ TeV on the left-hand side and $\sqrt{s} = 13$ TeV on the right-hand side. The decay pattern is very different here, and it is dominated by the decay $H^{\pm\pm} \to W^\pm H^\pm$. For a vector lepton coupling to the doubly charged Higgs set to $h'_{EE} = h''_{EE} = 0.1$, the decay into $H^\pm W^\pm$ dominates throughout the parameter space where it is kinematically allowed and can reach a branching fraction of over 90%, while the decay into vector leptons can have branching ratios of up to 25%. Again, the decay rates into $W^\pm$-boson pairs and ordinary leptons are below 1%, and the only differences between Case A and Case B are, as in Fig. 5, threshold effects. The dominance of the decay mode $H^{\pm\pm} \to W^\pm H^\pm$ persists, and is even more evident for smaller couplings with vector leptons, $h'_{EE} \simeq 0.01$.

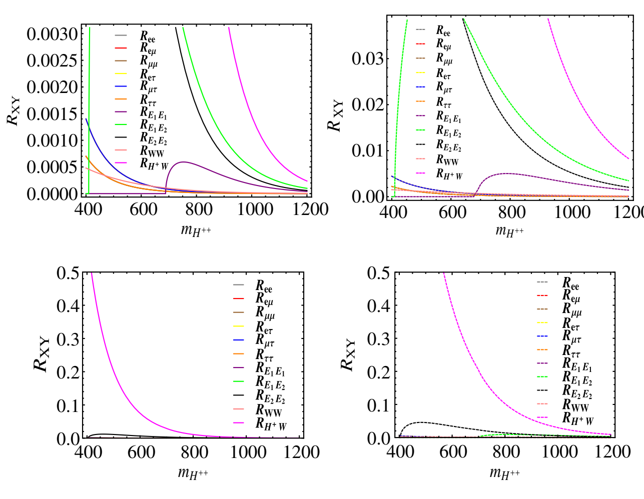

FIG. 6 (color online). Same as Fig. 5, but for Condition 2 (when $H^{\pm\pm} \to W^\pm H^\pm$ is kinematically allowed).

V. CONCLUSION

Despite no new signals of physics beyond the SM at the LHC, the SM cannot be the complete theory of particle interactions. An extension of the SM via additional vector-like leptons is not ruled out experimentally, and has been shown to provide a dark matter candidate. In models beyond the SM, the vector leptons can alter not only the phenomenology of the Higgs, but also that of other additional particles predicted by the models. We provide an example within the Higgs triplet model, where previously we showed that, in the absence of triplet-doublet Higgs mixing in the neutral sector ($\sin\alpha = 0$), there is no enhancement of the rate of decay of $h \to \gamma\gamma$ in this model with respect to the SM expectation. The HTM with vector leptons resolves two outstanding problems in the SM: the existence of neutrino masses and of a dark matter candidate.

Introducing vector leptons does not affect any of the tree-level decays or the production decay of the neutral Higgs boson observed at the LHC. However, loop decays into electroweak particles, such as $h \to Z\gamma$ and $h \to \gamma\gamma$, would be affected. We show that for the no-mixing scenario ($\sin\alpha = 0$) the decays rates into $\gamma\gamma$ and $Z\gamma$ do not depend on the doubly charged Higgs mass, and thus without the additional vector leptons these decays would be unchanged from the SM expectations. With vector leptons modest enhancements or suppressions are possible, most notably for $h \to \gamma\gamma$, where for large Yukawa couplings the rate of decays could even double. Under the same circumstances, the decay width for $h \to Z\gamma$ remains practically unchanged from its SM value. The model thus presents a mechanism for enhancing one loop decay and not the other, which seems to be consistent with the LHC data (so far).

If $\sin\alpha \neq 0$, the effect of the doubly charged Higgs bosons is felt for both $h \to \gamma\gamma$ and $h \to Z\gamma$, most spectacularly so for a very light $m_{H^{\pm\pm}} = 150$ GeV, which is ruled out for $\sin\alpha \neq 0$. Parameter-space regions where light doubly charged Higgs masses (200–250 GeV) *and* significant mixing in the neutral sector coexist are disfavored. In general, there are many parameter combinations for which the decay $h \to \gamma\gamma$ is enhanced, but there are few for an enhanced $h \to Z\gamma$, and these regions are ruled out by the branching ratio for $h \to \gamma\gamma$. However, if the decay $h \to \gamma\gamma$ is (modestly) enhanced, while $h \to Z\gamma$ is the same as the SM to 1σ, small mixing angles and light doubly charged Higgs bosons $m_{H^{\pm\pm}} \lesssim 300$ GeV are preferred. The fact that there are no regions of the parameter space consistent with present measurements of the branching ratio for $h \to \gamma\gamma$ and an enhanced rate for $h \to Z\gamma$ is a feature of this model, valid over the whole explored range of the parameter space. Other than this, there are no definite correlations or anticorrelations between these two loop-dominated decays.

The intermediate-mass doubly charged Higgs boson can decay into light vector leptons, which would alter its decay profile significantly. We explored this possibility and found that, if the singly charged Higgs mass is such that the decay $H^{\pm\pm} \to W^\pm H^\pm$ is not kinematically accessible, dominant branching ratios into vector leptons, if kinematically accessible, are expected for triplet Yukawa couplings $h'_{EE} = 0.1$. If and where the decay $H^{\pm\pm} \to W^\pm H^\pm$ is kinematically accessible, its corresponding branching ratio is the largest, while the branching fraction into vector leptons could reach 20–25% for $h'_{EE} = 0.1$. Under both of these circumstances, the decay patterns of the doubly charged Higgs bosons are changed, raising the hope that they can be found at masses around 200–600 GeV. The analyses presented here show that the cross section times branching ratios into vector leptons is significant enough to considerably alter the decay patterns of the doubly charged bosons, and at $\sqrt{s} = 13$ TeV these decay modes would be observable at the LHC—with cross sections times branching ratios of the order of several fb—and may be a promising way to discover vector leptons.

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[1] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012); S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).

[2] T. Aaltonen *et al.* (CDF Collaboration), Phys. Rev. Lett. **104**, 091801 (2010); A. Lister (CDF Collaboration), arXiv:0810.3349.

[3] K. Agashe, G. Perez, and A. Soni, Phys. Rev. D **71**, 016002 (2005); K. Agashe, G. Perez, and A. Soni, Phys. Rev. D **75**, 015002 (2007); A. Carmona and J. Santiago, J. High Energy Phys. 01 (2012) 100; G.-Y. Huang, K. Kong, and S. C. Park, J. High Energy Phys. 06 (2012) 099; C. Biggio, F. Feruglio, I. Masina, and M. Perez-Victoria, Nucl. Phys. **B677**, 451 (2004); D. E. Kaplan and T. M. P. Tait, J. High Energy Phys. 06 (2000) 020; H.-C. Cheng, Phys. Rev. D **60**, 075015 (1999).

[4] T. Moroi and Y. Okada, Mod. Phys. Lett. A **07**, 187 (1992); T. Moroi and Y. Okada, Phys. Lett. B **295**, 73 (1992); K. S. Babu, I. Gogoladze, M. U. Rehman, and Q. Shafi, Phys. Rev. D **78**, 055017 (2008); S. P. Martin, Phys. Rev. D **81**, 035004 (2010); P. W. Graham, A. Ismail, S. Rajendran, and P. Saraswat, Phys. Rev. D **81**, 055016 (2010); S. P. Martin, Phys. Rev. D **82**, 055019 (2010); J. Kang, P. Langacker, and B. D. Nelson, Phys. Rev. D **77**, 035003 (2008).

[5] B. A. Dobrescu and C. T. Hill, Phys. Rev. Lett. **81**, 2634 (1998); R. S. Chivukula, B. A. Dobrescu, H. Georgi, and C. T. Hill, Phys. Rev. D **59**, 075003 (1999); H. Collins, A. K. Grant, and H. Georgi, Phys. Rev. D **61**, 055002 (2000); H.-J. He, C. T. Hill, and T. M. P. Tait, Phys. Rev. D **65**, 055006 (2002); C. T. Hill and E. H. Simmons, Phys. Rep. **381**, 235 (2003); **390**, 553(E) (2004); R. Contino, L. Da Rold, and A. Pomarol, Phys. Rev. D **75**, 055014 (2007); **79**, 075003 (2009); K. Kong, M. McCaskey, and G. W. Wilson, J. High Energy Phys. 04 (2012) 079; A. Carmona, M. Chala, and J. Santiago, J. High Energy Phys. 07 (2012) 049; M. Gillioz, R. Grober, C. Grojean, M. Muhlleitner, and E. Salvioni, J. High Energy Phys. 10 (2012) 004.

[6] N. Arkani-Hamed, A. G. Cohen, E. Katz, and A. E. Nelson, J. High Energy Phys. 07 (2002) 034; T. Han, H. E. Logan, B. McElrath, and L.-T. Wang, Phys. Rev. D **67**, 095004 (2003); M. Perelstein, M. E. Peskin, and A. Pierce, Phys. Rev. D **69**, 075002 (2004); M. Schmaltz and D. Tucker-Smith, Annu. Rev. Nucl. Part. Sci. **55**, 229 (2005); M. S. Carena, J. Hubisz, M. Perelstein, and P. Verdier, Phys. Rev. D **75**, 091701 (2007); S. Matsumoto, T. Moroi, and K. Tobe, Phys. Rev. D **78**, 055018 (2008).

[7] A. Davidson and K. C. Wali, Phys. Rev. Lett. **60**, 1813 (1988); K. S. Babu and R. N. Mohapatra, Phys. Rev. D **41**, 1286 (1990); B. Grinstein, M. Redi, and G. Villadoro, J. High Energy Phys. 11 (2010) 067; D. Guadagnoli, R. N. Mohapatra, and I. Sung, J. High Energy Phys. 04 (2011) 093.

[8] A. Azatov, M. Toharia, and L. Zhu, Phys. Rev. D **80**, 035016 (2009).

[9] ATLAS Collaboration, Report No. ATLAS-CONF-2013-019; S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **718**, 348 (2012).

[10] F. del Aguila, J. de Blas, and M. Perez-Victoria, Phys. Rev. D **78**, 013010 (2008); S. Dawson and E. Furlan, Phys. Rev. D **86**, 015021 (2012); S. Dawson, E. Furlan, and I. Lewis, Phys. Rev. D **87**, 014007 (2013).

[11] A. Joglekar, P. Schwaller, and C. E. M. Wagner, J. High Energy Phys. 12 (2012) 064.

[12] J. Kearney, A. Pierce, and N. Weiner, Phys. Rev. D **86**, 113005 (2012); H. Davoudiasl, H.-S. Lee, and W. J. Marciano, Phys. Rev. D **86**, 095009 (2012); K. J. Bae, T. H. Jung, and H. D. Kim, Phys. Rev. D **87**, 015014

(2013); M. B. Voloshin, Phys. Rev. D **86**, 093016 (2012); D. McKeen, M. Pospelov, and A. Ritz, Phys. Rev. D **86**, 113004 (2012); H. M. Lee, M. Park, and W.-I. Park, J. High Energy Phys. 12 (2012) 037; C. Arina, R. N. Mohapatra, and N. Sahu, Phys. Lett. B **720**, 130 (2013); B. Batell, S. Jung, and H. M. Lee, J. High Energy Phys. 01 (2013) 135; H. Davoudiasl, I. Lewis, and E. Ponton, Phys. Rev. D **87**, 093001 (2013); J. Fan and M. Reece, J. High Energy Phys. 06 (2013) 004; A. Carmona and F. Goertz, J. High Energy Phys. 04 (2013) 163; C. Cheung, S. D. McDermott and K. M. Zurek, J. High Energy Phys. 04 (2013) 074; W.-Z. Feng and P. Nath, Phys. Rev. D **87**, 075018 (2013); C. Englert and M. McCullough, J. High Energy Phys. 07 (2013) 168.

[13] K. Ishiwata and M. B. Wise, Phys. Rev. D **84**, 055025 (2011).

[14] S. K. Garg and C. S. Kim, arXiv:1305.4712; L. Wang and X.-F. Han, Phys. Rev. D **86**, 095007 (2012).

[15] A. Joglekar, P. Schwaller, and C. E. M. Wagner, J. High Energy Phys. 07 (2013) 046.

[16] G. F. Giudice and O. Lebedev, Phys. Lett. B **665**, 79 (2008); E. Del Nobile, R. Franceschini, D. Pappadopulo, and A. Strumia, Nucl. Phys. **B826**, 217 (2010); S. P. Martin, Phys. Rev. D **81**, 035004 (2010).

[17] F. Arbabifar, S. Bahrami, and M. Frank, Phys. Rev. D **87**, 015020 (2013).

[18] CMS Collaboration, Report No. CMS-PAS-HIG-13-001; CMS Collaboration Report No. CMS-PAS-HIG-13-005; G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **726**, 88 (2013); G. Aad *et al.* (ATLAS Collaboration), Report No. ATLAS-CONF 2013-029.

[19] S. Chatrchyan *et al.* (CMS Collaboration), Eur. Phys. J. C **72**, 2189 (2012); G. Aad *et al.* (ATLAS Collaboration), Eur. Phys. J. C **72**, 2244 (2012).

[20] A. G. Akeroyd, M. Aoki, and H. Sugiyama, Phys. Rev. D **77**, 075010 (2008); A. G. Akeroyd, M. Aoki, and H. Sugiyama, Phys. Rev. D **79**, 113010 (2009); T. Fukuyama, H. Sugiyama, and K. Tsumura, J. High Energy Phys. 03 (2010) 044; S. T. Petcov, H. Sugiyama, and Y. Takanishi, Phys. Rev. D **80**, 015005 (2009); T. Fukuyama, H. Sugiyama, and K. Tsumura, Phys. Rev. D **82**, 036004 (2010); A. G. Akeroyd, C.-W. Chiang, and N. Gaur, J. High Energy Phys. 11 (2010) 005; A. Arhrib, R. Benbrik, M. Chabab, G. Moultaka, and L. Rahili, J. High Energy Phys. 04 (2012) 136; A. Arhrib, R. Benbrik, M. Chabab, G. Moultaka, and L. Rahili, in *Proceedings of The 2011 International Workshop on Future Linear Colliders (LCWS11), Granada, Spain, 2011* (to be published), arXiv:1202.6621; A. Arhrib, R. Benbrik, M. Chabab, G. Moultaka, M. C. Peyranere, L. Rahili, and J. Ramadan, Phys. Rev. D 84, 095005 (2011); E. J. Chun, H. M. Lee, and P. Sharma, J. High Energy Phys. 11 (2012) 106; A. Melfo, M. Nemevšek, F. Nesti, G. Senjanović, and Y. Zhang, Phys. Rev. D **85**, 055018 (2012); M. Aoki, S. Kanemura, and K. Yagyu, Phys. Rev. D **85**, 055007 (2012); S. Kanemura and K. Yagyu, Phys. Rev. D **85**, 115009 (2012); M. Aoki, S. Kanemura, M. Kikuchi, and K. Yagyu, Phys. Rev. D **87**, 015012 (2013); E. J. Chun and P. Sharma, Phys. Lett. B **722**, 86 (2013); L. Wang and X.-F. Han, Phys. Rev. D **87**, 015015 (2013).

[21] R. Dermisek and A. Raval, Phys. Rev. D **88**, 013017 (2013).

[22] K. Ishiwata and M. B. Wise, Phys. Rev. D **88**, 055009 (2013).

[23] S. Chatrchyan *et al.* (CMS Collaboration), arXiv:1307.5515; ATLAS Collaboration, Report No. ATLAS-CONF-2013-009.

[24] C.-S. Chen, C.-Q. Geng, D. Huang, and L.-H. Tsai, Phys. Rev. D **87**, 075019 (2013); C.-S. Chen, C.-Q. Geng, D. Huang, and L.-H. Tsai, Phys. Lett. B **723**, 156 (2013); P. S. Bhupal Dev, D. K. Ghosh, N. Okada, and I. Saha, J. High Energy Phys. 03 (2013) 150; J. High Energy Phys. 05 (2013) 049(E).

[25] G. Abbiendi *et al.* (OPAL Collaboration), Phys. Lett. B **526**, 221 (2002); G. Abbiendi *et al.* (OPAL Collaboration), Phys. Lett. B **577**, 18 (2003); P. Achard *et al.* (L3 Collaboration), Phys. Lett. B **576**, 18 (2003); J. Abdallah *et al.* (DELPHI Collaboration), Phys. Lett. B **552**, 127 (2003).

[26] J. Sztuk-Dambietz (H1 and ZEUS Collaboration), J. Phys. Conf. Ser. **110**, 072042 (2008); A. Aktas *et al.* (H1 Collaboration), Phys. Lett. B **638**, 432 (2006).

[27] V. M. Abazov *et al.* (D0 Collaboration), Phys. Rev. Lett. **101**, 071803 (2008); V. M. Abazov *et al.* (D0 Collaboration), Phys. Rev. Lett. **108**, 021801 (2012); D. E. Acosta *et al.* (CDF Collaboration), Phys. Rev. Lett. **93**, 221802 (2004); T. Aaltonen *et al.* (The CDF Collaboration), Phys. Rev. Lett. **101**, 121801 (2008).

[28] S. Chatrchyan *et al.* (CMS Collaboration), J. High Energy Phys. 07 (2013) 122; G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **722**, 305 (2013); S. Chatrchyan *et al.* (CMS Collaboration), Eur. Phys. J. C **72**, 2189 (2012); G. Aad *et al.* (ATLAS Collaboration), Eur. Phys. J. C **72**, 2244 (2012).

[29] C.-W. Chiang, T. Nomura, and K. Tsumura, Phys. Rev. D **85**, 095023 (2012); S. Kanemura, K. Yagyu, and H. Yokoya, Phys. Lett. B **726**, 316 (2013).

[30] M. Muhlleitner and M. Spira, Phys. Rev. D **68**, 117701 (2003).

[31] A. G. Akeroyd, C.-W. Chiang, and N. Gaur, J. High Energy Phys. 11 (2010) 005.