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Asymmetric Di-Higgs Signals of the N2HDM-$U(1)$

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The two-Higgs-doublet model with a $U(1)_H$ gauge symmetry (N2HDM-$U(1)$) has several advantages compared to the Z_2 version (N2HDM-Z_2): It is purely based on gauge symmetries, involves only spontaneous symmetry breaking and is more predictive because it contains one less free parameter in its Higgs potential which ensures CP conservation at the same time. However, the phenomenology of its Higgs sector has been barely studied. After pointing out that a second, so far unknown version of the N2HDM-$U(1)$ exists, we examine the phenomenological consequences of the differences in the scalar potentials. In particular, we find that while the N2HDM-Z_2 predicts suppressed branching ratios for decays into different Higgses for small scalar mixing (as suggested by Higgs coupling measurements), both versions of the N2HDM-$U(1)$ allow for sizable rates. This is particularly important in light of the CMS excess in Higgs pair production at around 650 GeV decaying a Standard Model Higgs decaying to photons and a new scalar with a mass of ≈ 90 GeV decaying to bottom quarks (i.e., compatible with the CMS $\gamma\gamma$ excess at 95 GeV), which, as we will show, can be explained within the N2HDM-$U(1)$, predicting an interesting and unavoidable $Z + b\bar{b}$ signature.

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I. INTRODUCTION

The discovery of the Brout-Englert-Higgs boson [1–6] by ATLAS [7] and CMS [8] established, for the first time, the existence of a fundamental scalar particle within the Standard Model (SM). This observation motivates the existence of more scalar particles and in turn the experimental search for them. While the 125 GeV Higgs (h) has approximately SM-like properties [9–14], this does not exclude the existence of additional scalar bosons, as long as their role in the breaking of the SM electroweak (EW) gauge symmetry and the mixing with the SM-like Higgs is sufficiently small.

In this context, strong constraints on new physics (NP) models are provided by the ρ parameter that relates the electroweak gauge couplings to the W and Z masses and is defined to be unity in the SM at tree level. This singles out models with $SU(2)_L$-singlet or $SU(2)_L$-doublet extensions of the SM Higgs sector whose vacuum expectation values (VEVs) conserve the custodial symmetry, such that the additional scalars only give loop-level effects in the ρ parameter.[1]

The most studied extensions of the SM scalar sector are the two-Higgs-doublet models (2HDMs) [16–19]. Here, usually a discrete Z_2 symmetry is imposed to both solve the problem of flavour changing neutral currents [20] (resulting in four different versions with natural flavour conservation [21, 22]) and to provide (accidental) CP conservation in the Higgs potential. However, for phenomenological reasons, i.e., to give VEV-independent masses to the additional scalars, the Z_2 symmetry must be broken. In order to avoid domain wall problems caused by a spontaneous discrete symmetry breaking [23], the Z_2 symmetry is usually softly broken by a dimension-two term [24]. This operator (in case of a non-vanishing and non-aligned λ_5 term) in general gives rise to CP violation within the Higgs potential.

Reference [25] proposed to solve these problems by replacing the discrete Z_2 symmetry with a $U(1)_H$ gauge symmetry, which can mimic the effect of the Z_2 symmetry but forbids the explicit soft-breaking term. However, if the Z' boson originating from the $U(1)_H$ gauge is required to be heavier than the EW scale, one has to supplement the model with an additional scalar charged under $U(1)_H$; minimally a complex scalar ϕ that is a singlet under the SM gauge group. Because its CP-odd component becomes (in the limit without scalar mixing) the longitudinal component of the Z', the scalar potential effectively resembles the one of the Next-to-Minimal

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[1] Also larger $SU(2)_L$ multiplets are possible in case their VEV is very small, or if a global custodial $SU(2)$ symmetry protects the ρ parameter, which however entails larger and more complex Higgs sectors [15].

2HDM (N2HDM) with a real scalar (see, e.g., Refs. [26–44]). In particular, the VEV of ϕ gives rise to the m_{12}^2 term that softly breaks the Z_2 symmetry.

While in the N2HDM with two discrete Z_2 symmetries (N2HDM-Z_2)[2] Higgs-boson-related collider observables have been studied in detail, including loop effects [45, 46], and even automated codes exist [47–50],[3] studies of the N2HDM with a $U(1)_H$ gauge symmetry (N2HDM-$U(1)$) have not focused on the collider phenomenology of the additional scalars but mostly considered dark matter [54–56], muon $g-2$ [57], neutrino masses [58–62], $b \to s\ell^+\ell^-$ anomalies [63, 64], Z' searches [7, 65, 66], and Higgs signal strengths [67]. However, also the (effective) scalar potential of this N2HDM-$U(1)$ is different from the one of the usual N2HDM-Z_2, which, in particular, leads to different Higgs self-interactions and therefore different decay rates for heavy scalars into light ones.

This aspect is now more relevant in light of the ongoing and intensified LHC searches for new scalar bosons (see, e.g., Ref. [68] for a recent review). While no unequivocal evidence for a new particle has been observed, interesting hints for new scalars with masses around 95 GeV [44, 69–82], 151 GeV [83–87] and 670 GeV [72, 88–90] have been reported.[4] In particular, the CMS excess for a ≈ 650 GeV scalar decaying into a ≈ 90 GeV one (i.e., compatible with the 95 GeV hints mentioned above due to the limited detector resolution for bottom jets) and a SM Higgs [72] needs multiplet new Higgses with a mass hierarchy (which is only possible if they are not within the same $SU(2)_L$ multiple with respecting perturbativity bounds). Furthermore, as it is an asymmetric di-Higgs decay, it requires sizable self-interactions of the three different scalars. While the rates of such asymmetric di-Higgs decays are in general small in the MSSM [95, 96], 2HDMs [97] and also in the N2HDM-Z_2 [98], we will show that for the N2HDM-$U(1)$ they are naturally sizable.

II. THE MODEL

As outlined in the introduction, a Z_2 symmetry is commonly used to construct the four versions of the 2HDMs with natural flavour conservation and, at the same time, constrains the scalar potential. In the N2HDM, even two Z_2 symmetries are usually employed to prevent tree-level flavour changing neutral currents and eliminate most sources of CP violation. We want to use instead a single $U(1)_H$ gauge symmetry under which at least two of the scalar fields are charged.

[2] Under the first Z_2 symmetry only H_2 is odd, while for the second Z_2 only the real scalar is odd.

[3] The program ewN2HDECAY [48] is based on HDECAY [51, 52] and 2HDECY [53].

[4] In addition, anomalies in multi-lepton final states exist, which can also be explained by new scalars [83, 91–94].

We start with the scalar potential for the two $SU(2)_L$ doublets H_1 and H_2 with hypercharge 1/2 (where according to the usual 2HDM conventions H_2 contains most of the SM Higgs). If the $U(1)_H$ charges of H_1 and H_2 are different, operators with an odd number of these fields are forbidden, leading to

$$\begin{aligned}\mathcal{V}_H =& m_{11}^2|H_1|^2 + m_{22}^2|H_2|^2 + \frac{\lambda_1}{2}(H_1^\dagger H_1)^2 + \frac{\lambda_2}{2}(H_2^\dagger H_2)^2 \\ &+ \lambda_3(H_1^\dagger H_1)(H_2^\dagger H_2) + \lambda_4(H_1^\dagger H_2)(H_2^\dagger H_1)\,. \end{aligned} \tag{1}$$

This potential is CP conserving as it does not contain the soft-breaking term $m_{12}^2 H_1^\dagger H_2$ nor the term $\frac{\lambda_5}{2}(H_1^\dagger H_2)^2$ contained in the 2HDM with the (softly-broken) Z_2 symmetry.

Next, we add a complex scalar SM singlet ϕ that is charged under the $U(1)_H$ gauge symmetry. Its self-interactions, as well as the ones with two identical doublets

$$\mathcal{V}_\phi = |\phi|^2\left(m_\phi^2 + \frac{\lambda_\phi}{2}|\phi|^2 + \lambda_{\phi 1}|H_1|^2 + \lambda_{\phi 2}|H_2|^2\right)\,, \tag{2}$$

are allowed independently of the $U(1)_H$ charges. In addition, there are two options for charge assignments under the $U(1)_H$ symmetry:
(a) If $|Q_H(\phi)| = |Q_H(H_1) - Q_H(H_2)|$, one has the term

$$\mathcal{V}_{\phi H}^{\mathrm{a}} = \sqrt{2}\mu H_1^\dagger H_2 \phi + \mathrm{h.c.}\,, \tag{3}$$

or ϕ replaced by $\phi^\dagger$, depending on the sign on the $U(1)_H$ charge.
(b) If $|Q_H(\phi)| = |Q_H(H_1) - Q_H(H_2)|/2$, the term

$$\mathcal{V}_{\phi H}^{\mathrm{b}} = \lambda_{\phi 12}(H_1^\dagger H_2)\phi^2 + \mathrm{h.c.}\,, \tag{4}$$

is gauge invariant. Case (a) was already proposed in Ref. [25], while case (b) is novel, to the best of our knowledge.

Note that we have normalized the prefactors of these potentials in such a way, that once we decompose

$$\phi = (v_S + \hat{S} + i\eta_S)/\sqrt{2}\,, \tag{5}$$

η_S (mostly) becomes the longitudinal mode of the Z' and the terms involving $\hat{S}$ match the N2HDM-Z_2. Here, v_S is the VEV of ϕ and one can choose it to be real and positive without loss of generality. Therefore, disregarding the Z' boson, which could be heavy or weakly coupled, the N2HDM-$U(1)$ resembles the N2HDM-Z_2 with the important differences that the m_{12}^2 and λ_5 terms are only effectively generated by v_S and Z'-exchange, respectively, similar to the μ term in the Next-to-Minimal Supersymmetric Standard Model (NMSSM) [99–104]. Importantly, this leads to the absence of CP violation in the scalar potential (even when the Z' is integrated out), while this, in general, cannot be avoided in both the N2HDM-Z_2 and NMSSM.

We know from Higgs signal strength measurements that the mixing among the SM-like Higgs and the other two CP-even scalars should be rather small. Therefore, we will label the CP-even mass eigenstates, contained in the absence of mixing within H_2, H_1, and ϕ as h, H, and S, respectively.[5] Importantly, the mixing between H, h and S in the N2HDM-$U(1)$ is related to the masses m_H, $m_{H^\pm}$, and m_A (where $H^\pm$ and A denote the charged and CP-odd Higgs, respectively) because they all involve the effective m_{12}^2 term originating from μ or $\lambda_{\phi 12}$. This means that the effective m_{12}^2 term automatically leads to H-S mixing as can be inferred from the CP-even mass matrix (at leading order in $\tan\beta$)

$$M_\rho^2 \approx \begin{pmatrix} -\mu v_S \tan\beta & \mu v_S & \mu v \\ \mu v_S & \lambda_2 v^2 & \lambda_{\phi 2} v v_S \\ \mu v & \lambda_{\phi 2} v v_S & \lambda_\phi v_S^2 \end{pmatrix}, \qquad (6)$$

where $\tan\beta = v_2/v_1$ and $\langle H_i \rangle = v_i/\sqrt{2}$. Note that the effects of $\lambda_1, \lambda_3, \lambda_4, \lambda_{\phi 1}$ on the masses become negligible in the large $\tan\beta$ region. The full expressions for the minimization, the mass matrices, etc., can be found in the appendix.

Concerning the fermion sector, the most natural choice is probably to assume that they are uncharged under $U(1)_H$, or to assign equal charges to left-handed and right-handed fields (such as $B-L$ or $L_\mu - L_\tau$) in order to avoid gauge anomalies. In this setup, the doublet H_2 would be $U(1)_H$ neutral, while H_1 carries some $U(1)_H$ charge Q_H. This then leads to a type-I Yukawa sector, which also has the advantage of being quite unconstrained in the large $\tan\beta$ and small α (small Higgs mixing) limit. However, also the other three types of 2HDMs with natural flavour conservation, as well as the generic type-III model,[6] can be obtained even in an anomaly-free fashion if the fermion sector is extended [25].

III. PHENOMENOLOGY

The N2HDM-$U(1)$ is in general more predictive than the N2HDM-Z_2 as it contains one less free parameter and has no sources of CP violation in the Higgs potential. However, what is the most striking difference regarding LHC observables between the different N2HDMs, even when the Z' predicted by the N2HDM-$U(1)$ is disregarded, as it might be heavy or weakly coupled?

To answer this, let us consider the limit of vanishing mixing among the neutral CP-even scalars, in which h is purely SM-like, H only couples to $W^\pm H^\mp$ and ZA, and S is sterile. Now, μ ($\lambda_{\phi 12}$) in Eq. (3) (Eq. (4)) has to be non-vanishing to give masses to H, A and $H^\pm$ that are above the EW scale, i.e.,

$$m_H^2 \approx m_A^2 \approx m_{H^\pm}^2 \approx -\mu v_S \tan\beta\,. \qquad (7)$$

This then at the same time induces $\tan\beta$ suppressed H-h and H-S mixing while from Eq. (6) we see that H-h mixing can be avoided, at leading order in $\tan\beta$, for $\lambda_{\phi 2} = 0$. This means that for $m_H \gg v$ the only unsuppressed decay of H, in the absence of Yukawa coupling of H_1, is $H \to Sh$ for case (a), and in addition $H \to SS$ for case of (b), if $m_H \gg m_S$ and $m_{H^\pm} \approx m_H$. Therefore, in the large $\tan\beta$ limit, the N2HDM-$U(1)$ predicts sizable branching ratios for $H \to Sh$ (and also $H \to SS$ in case (b)). As this decay in the N2HDM-Z_2 is suppressed by small mixing angles, the discovery of an asymmetric di-Higgs signal would be a smoking gun for the N2HDM-$U(1)$.[7]

Let us now illustrate this observation more quantitatively in the context of the hint for the ≈ 650 GeV boson decaying into a ≈ 90 GeV scalar and the SM Higgs with a global (local) significance of $2.8\,\sigma$ ($3.8\,\sigma$) [72]. Here, the ≈ 90 GeV resonance decays into $b\bar{b}$ and the SM Higgs into $\gamma\gamma$. Because the detector resolution for bottom jets is not so good, this ≈ 90 GeV excess could be compatible with the $\gamma\gamma$ [70], $\tau\bar{\tau}$ [71] and the LEP ZH measurement [69] around ≈ 95 GeV as well as with the $\gamma\gamma$ and ZZ excesses around ≈ 670 GeV. This makes this asymmetric di-Higgs signal particularly interesting and effectively eliminates the look-elsewhere effect of the CMS di-Higgs analysis.

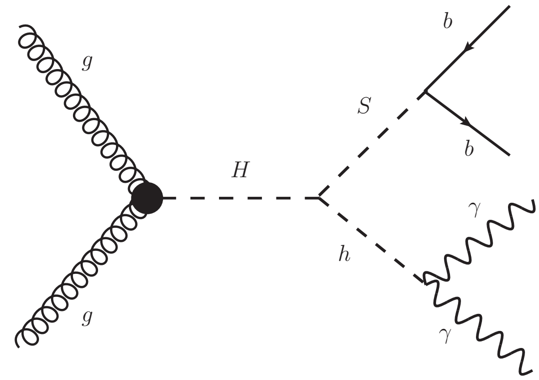

FIG. 1. Feynman diagram showing resonant asymmetric Higgs pair production. The discovery of this process, for which the CMS measurement constitutes a first hint, would be a smoking gun for the N2HMD-$U(1)$. Here, the black circle denotes the loop-induced effective coupling to gluons. However, note that the heavy top limit cannot be used because $m_t \ll m_H$ and we use the expression for a dynamical top quark in our numerical analysis.

[5] As we only consider the case of small mixing, we will label in the main text both the CP-even components of the doublets and the singlet, as well as the mass eigenstates by h, H, and S and use them interchangeably. In the appendix, the full mass matrices in the weak eigenbasis are given.

[6] The type-III model has been comprehensively studied in Ref. [105] as it can (partially) explain the tensions in $R(D^{(*)})$ [106, 107].

[7] This resembles the situation in the NMSSSM where also sizable asymmetric Higgs decays are possible [108, 109].

The target cross section for $pp \to b\bar{b}\gamma\gamma$ is $\approx 0.35^{+0.17}_{-0.13}$ fb. Taking into account that $\mathrm{BR}(h \to \gamma\gamma) \approx 0.23\%$, we need $\sigma(pp \to (650) \to (95)\, h) \times \mathrm{BR}((95) \to b\bar{b}) \approx 150$ fb to explain the CMS excess in $b\bar{b}\gamma\gamma$. However, the CMS analysis of $pp \to b\bar{b}\tau\bar{\tau}$ [72] finds an upper limit on the corresponding cross section of ≈ 4 fb. With $\mathrm{BR}(h \to \tau\bar{\tau})/\mathrm{BR}(h \to \gamma\gamma) \approx 20$, this translates into the limit $\sigma(pp \to (650) \to (95)\, h) \times \mathrm{BR}((95) \to b\bar{b}) \lessapprox 90$ fb. Therefore, the $b\bar{b}\gamma\gamma$ excess cannot be fully explained, but it is still possible to account for it within $2\,\sigma$ and we will aim at

$$\sigma(pp \to (650) \to (95)\, h) \times \mathrm{BR}((95) \to b\bar{b}) \approx 70\,\mathrm{fb}\,. \quad (8)$$

There are two options within the N2HDM-$U(1)$; one can identify the ≈ 95 GeV state with H and the ≈ 650 GeV one with S or vice versa. However, in the case of $pp \to S \to Hh$, the μ term is naturally small because H is light, such that also the branching ratio is suppressed, unless one chooses very small mixing angles among the CP-even scalars. Let us, therefore, consider $pp \to H \to Sh$ in the following, again in the limit of small mixing, i.e., neglecting $\tan\beta$ suppressed terms. To obtain a sufficient production cross section of H we will consider the case of a non-minimal flavour structure and assume that H has an (effective) Yukawa coupling to top quarks originating from the Lagrangian term $-\tilde{Y}^t \bar{Q}_L \tilde{H}_1 t_R$.[8] This coupling then also leads to unsuppressed decays of $H \to t\bar{t}$ (and $A \to t\bar{t}$).

For the numerical analysis we use that a SM Higgs with a mass of $\approx$650 GeV would have a gluon fusion production cross section of ≈ 1.35 pb at NNLO [112–117]. This means that a coupling to top quarks is needed, that is around one quarter of the one of the SM Higgs, i.e., $\tilde{Y}_t \approx Y_t/4/(\sqrt{\mathrm{BR}(H \to Sh)}\sqrt{\mathrm{BR}(S \to b\bar{b})})$. Therefore, assuming that S decays SM-like ($\mathrm{BR}(S \to b\bar{b}) \approx 0.8$) results with Eq. (8) in $\sigma(pp \to H) \approx 84\,\mathrm{fb}/\mathrm{BR}(H \to Sh)$. Based on the Goldstone boson equivalence theorem [118, 119], we also expect $\mathrm{BR}(A \to SZ) \approx \mathrm{BR}(H \to Sh)$ leading to $pp \to A \to ZS \to Zb\bar{b}$ (and also $pp \to A \to Zh \to Zb\bar{b}$) with cross sections around 1.5×70 fb,[9] searched for by ATLAS [121–123] and CMS [124, 125]. Note that in fact, Ref. [125] finds a mild excess within the relevant region.

Furthermore, we can predict the cross section of $H \to t\bar{t}$ and $A \to t\bar{t}$, as well as $pp \to t\bar{t}H \to t\bar{t}t\bar{t}$ and $pp \to t\bar{t}A \to t\bar{t}t\bar{t}$ as a function of $\tan\beta$ and v_S (assuming $\lambda_{\phi 2} = 0$ as well as $m_H \approx m_A$) and compare this to the limits on the resonant $t\bar{t}$ production of CMS [127] and ATLAS [128] as well as to $t\bar{t}t\bar{t}$ production measured by CMS [129] and ATLAS [126]. This is illustrated in Fig. 2 where we show the predicted cross section for $pp \to A \to t\bar{t}$ in units of pb as a function of $\tan\beta$ and v_S. The red region is excluded by the $pp \to t\bar{t}t\bar{t}$ search of ATLAS and the yellow region by the requirement of positive eigenvalues of the mass matrix as well as perturbative couplings. Since $\sigma(pp \to H \to Sh) \approx 84$ fb is required, when $\mathrm{BR}(S \to \gamma\gamma) \approx 0.15\%$ (again assuming that S has SM-like branching ratios) we obtain for the inclusive cross section $\sigma(pp \to S + X \to \gamma\gamma + X) \approx 0.1$ fb which is compatible with the current limits but insufficient to explain the $\gamma\gamma$ excess at 95 GeV of ≈ 50 fb [70]. Therefore, direct production of S would be required in addition, to explain the $\gamma\gamma$ excess.

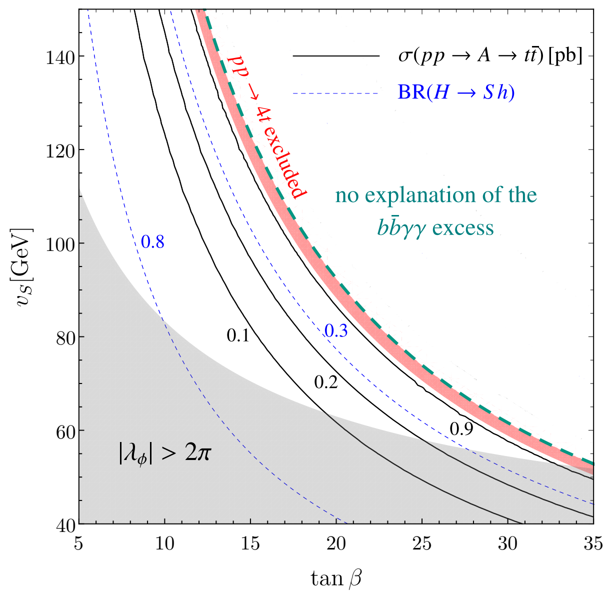

FIG. 2. Predictions for $\sigma(pp \to A \to t\bar{t})$[pb] as a function of $\tan\beta$ and v_S in the N2HMD-$U(1)$ for case (a), assuming that the CMS excess $b\bar{b}\gamma\gamma$ in Eq. (8) is explained. The grey region is excluded by the requirement of perturbative couplings, while the red region is excluded by the $t\bar{t}t\bar{t}$ search [126], assuming $m_A \approx m_H$. Note that the $b\bar{b}\gamma\gamma$ excess cannot be explained in the top-right region of the green dashed line.

IV. CONCLUSIONS

While Higgs physics in the N2HDM with two discrete Z_2 symmetries (N2HDM-Z_2) has been studied in detail in the literature, this phenomenological aspect of the N2HDM with a $U(1)_H$ gauge symmetry has received little to no attention so far. While both versions have desirable features such as natural flavour conservation, there are even several advantages of the N2HDM-$U(1)_H$ over the N2HDM-Z_2:

- Only one $U(1)_H$ gauge symmetry is needed instead

[8] This coupling can be induced at tree-level by vector-like quarks mixing with SM ones via a coupling to S. Alternatively, an effective coupling to gluons could be loop-induced by colored new heavy fermions or scalars. In fact, CMS observed an excess in di-di-jet searches [110] that point towards new colored particles at the TeV scale [111].

[9] Note that at this energy, the gluon fusion production cross section for a pseudo-scalar via top-quark loops is ≈ 1.5 times the one of a CP-even scalar with the same mass and coupling (see, e.g., Refs. [96, 120]).

of two Z_2 symmetries.

- Like the SM, the N2HDM-$U(1)_H$ is built on local gauge symmetries and spontaneous symmetry breaking (i.e., unlike the N2HDM-Z_2 no soft-breaking is needed).

- The N2HDM-$U(1)_H$ symmetry is more predictive than the N2HDM-Z_2 because it contains one less free parameter.

If the Z' boson is decoupled from phenomenology, either because it is heavy or weakly interacting, the scalar sector of the N2HDM-$U(1)_H$ is close to the one of the N2HDM-Z_2, however, there are important differences:

- In the N2HDM-$U(1)$ no λ_5 term is allowed, leading to CP conservation. This feature is even conserved when the Z' is integrated out because of an automatic phase alignment.

- The m_{12}^2 term is absent before spontaneous symmetry breaking and induced by the VEV of ϕ, either from the term $\mu H_1^\dagger H_2 \phi$ or $\lambda_{\phi 12}(H_1^\dagger H_2)\phi^2$, depending on the charge assignment. Please note that the latter option was, to the best of our knowledge, not proposed before in the literature.

- While in N2HDM-Z_2, if H is heavy, only symmetric decays into Higgs pairs, i.e., $H \to SS$ and $H \to hh$ are possible in the limit of zero mixing, in the N2HDM-$U(1)$ one expects naturally large branching ratios for $H \to Sh$. Note that while in case (a), only asymmetric decays are unsuppressed, in case (b) also decays to identical scalars (e.g., $H \to SS$) can be sizable.

The latter has important implications for the asymmetric $\approx 650\,\text{GeV}$ excess in $b\bar{b}\gamma\gamma$. While even if H is equipped with a sufficiently high production cross section (e.g., from direct top-quark Yukawa couplings of H_1), the N2HDM-Z_2 could not account for the preferred central value of the measurement as $\text{BR}(H \to Sh)$ could not be sizable enough, taking into account the limits on the scalar mixing from Higgs coupling strength measurements at the LHC. However, the N2HDM-$U(1)$ can account for this measurement, predicting signatures in $pp \to H(A) \to t\bar{t}$, $pp \to t\bar{t}H(A) \to 4t$ and $pp \to A \to SZ$, not far away from the current experimental limits.

Finally, let us point out that Z–Z' mixing, in general present in this model, can naturally account for the higher than expected value of the W mass [130], as suggested by the measurement of the CDF-II collaboration [131]. Together with the previous arguments this strongly motivates detailed studies of the N2HDM-$U(1)$ (including its Higgs sector) which, in our opinion, should be considered to be (at least) at the same level of interest as the standard N2HDM-Z_2 and therefore be examined with the same scrutiny in the future.

Note Added: During completion of this work, CMS presented updated results for low mass searches for new scalars decaying into $\gamma\gamma$ [132], confirming the previous excess.

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Appendix A: Minimization and Mass Matrices

In this Appendix, we give the minimization conditions and mass matrices of both types, (a) and (b), of the N2HDM-$U(1)$. We define H_1, H_2 and S as,

$$H_1 = \begin{pmatrix} w_1^+ \\ \frac{v_1 + \hat{H} + i\eta_1}{\sqrt{2}} \end{pmatrix}, \qquad H_2 = \begin{pmatrix} w_2^+ \\ \frac{v_2 + \hat{h} + i\eta_2}{\sqrt{2}} \end{pmatrix}, \qquad \phi = \frac{v_S + \hat{S} + i\eta_S}{\sqrt{2}}. \tag{A1}$$

1. Case (a): $|Q_H(\phi)| = |Q_H(H_1) - Q_H(H_2)|$

The minimization conditions are

$$m_{11}^2 + \frac{1}{2}\lambda_1 v_1^2 + \frac{1}{2}\lambda_{345} v_2^2 + \frac{1}{2}\lambda_{\phi 1} v_S^2 + \mu v_S \frac{v_2}{v_1} = 0\,,$$

$$m_{22}^2+\frac{1}{2}\lambda_2 v_2^2+\frac{1}{2}\lambda_{345}v_1^2+\frac{1}{2}\lambda_{\phi 2}v_S^2+\mu v_S\frac{v_1}{v_2}=0\,,$$
$$m_S^2+\frac{1}{2}\lambda_{\phi 1}v_1^2+\frac{1}{2}\lambda_{\phi 2}v_2^2+\frac{1}{2}\lambda_\phi v_S^2+\mu\frac{v_1 v_2}{v_S}=0\,, \tag{A2}$$

where $\lambda_{345}=\lambda_3+\lambda_4+\lambda_5^{\rm eff}$, and $\lambda_5^{\rm eff}\neq 0$ is only generated if the Z' is integrated out. The scalar squared-mass matrices are

$$M_\rho^2=\begin{pmatrix}\lambda_1 v_1^2-\mu v_S\frac{v_2}{v_1} & \lambda_{345}v_1v_2+\mu v_S & \lambda_{\phi 1}v_1v_S+\mu v_2\\ \lambda_{345}v_1v_2+\mu v_S & \lambda_2 v_2^2-\mu v_S\frac{v_1}{v_2} & \lambda_{\phi 2}v_2v_S+\mu v_1\\ \lambda_{\phi 1}v_1v_S+\mu v_2 & \lambda_{\phi 2}v_2v_S+\mu v_1 & \lambda_\phi v_S^2-\mu\frac{v_1v_2}{v_S}\end{pmatrix}, \tag{A3}$$

$$M_\eta^2=\begin{pmatrix}-\mu v_S\frac{v_2}{v_1}-\lambda_5^{\rm eff}v_2^2 & \mu v_S+\lambda_5^{\rm eff}v_1v_2 & \mu v_2\\ \mu v_S+\lambda_5^{\rm eff}v_1v_2 & -\mu v_S\frac{v_1}{v_2}-\lambda_5^{\rm eff}v_1^2 & -\mu v_1\\ \mu v_2 & -\mu v_1 & -\mu\frac{v_1v_2}{v_S}\end{pmatrix}, \tag{A4}$$

$$M_w^2=\begin{pmatrix}-\mu v_S\frac{v_2}{v_1}-(\lambda_4+\lambda_5^{\rm eff})\frac{v_2^2}{2} & \mu v_S+(\lambda_4+\lambda_5^{\rm eff})\frac{v_1v_2}{2}\\ \mu v_S+(\lambda_4+\lambda_5^{\rm eff})\frac{v_1v_2}{2} & -\mu v_S\frac{v_1}{v_2}-(\lambda_4+\lambda_5^{\rm eff})\frac{v_1^2}{2}\end{pmatrix}, \tag{A5}$$

which are defined via the bilinear potential terms

$$V_{m^2}=\frac{1}{2}\begin{pmatrix}\hat H & \hat h & \hat S\end{pmatrix}M_\rho^2\begin{pmatrix}\hat H\\ \hat h\\ \hat S\end{pmatrix}+\frac{1}{2}\begin{pmatrix}\eta_1 & \eta_2 & \eta_S\end{pmatrix}M_\eta^2\begin{pmatrix}\eta_1\\ \eta_2\\ \eta_S\end{pmatrix}+\begin{pmatrix}w_1^- & w_2^-\end{pmatrix}M_w^2\begin{pmatrix}w_1^+\\ w_2^+\end{pmatrix}. \tag{A6}$$

The eigenvalues of the CP-odd and charged-Higgs masses are then given by

$$M_A^2=-\mu\left(\frac{v_Sv^2}{v_1v_2}+\frac{v_1v_2}{v_S}\right)-\lambda_5^{\rm eff}v^2\,, \tag{A7}$$

$$M_{H^\pm}^2=-\frac{\mu v_Sv^2}{v_1v_2}-\left(\lambda_4+\lambda_5^{\rm eff}\right)\frac{v^2}{2}\,. \tag{A8}$$

2. Case (b): $|Q_H(\phi)|=|Q_H(H_1)-Q_H(H_2)|/2$

The minimization conditions in this case are

$$m_{11}^2+\frac{1}{2}\lambda_1v_1^2+\frac{1}{2}\lambda_{345}v_2^2+\frac{1}{2}\lambda_{\phi 1}v_S^2+\frac{1}{2}\frac{\lambda_{\phi 12}v_2v_S^2}{v_1}=0\,,$$
$$m_{22}^2+\frac{1}{2}\lambda_2v_2^2+\frac{1}{2}\lambda_{345}v_1^2+\frac{1}{2}\lambda_{\phi 2}v_S^2+\frac{1}{2}\frac{\lambda_{\phi 12}v_1v_S^2}{v_2}=0\,,$$
$$m_S^2+\frac{1}{2}\lambda_{\phi 1}v_1^2+\frac{1}{2}\lambda_{\phi 2}v_2^2+\frac{1}{2}\lambda_\phi v_S^2+\lambda_{\phi 12}v_1v_2=0\,, \tag{A9}$$

and the squared-mass matrices are given by

$$M_\rho^2=\begin{pmatrix}\lambda_1v_1^2-\frac{\lambda_{\phi 12}v_2v_S^2}{2v_1} & \lambda_{345}v_1v_2+\frac{1}{2}\lambda_{\phi 12}v_S^2 & \lambda_{\phi 1}v_1v_S+\lambda_{\phi 12}v_2v_S\\ \lambda_{345}v_1v_2+\frac{1}{2}\lambda_{\phi 12}v_S^2 & \lambda_2v_2^2-\frac{\lambda_{\phi 12}v_1v_S^2}{2v_2} & \lambda_{\phi 2}v_2v_S+\lambda_{\phi 12}v_1v_S\\ \lambda_{\phi 1}v_1v_S+\lambda_{\phi 12}v_2v_S & \lambda_{\phi 2}v_2v_S+\lambda_{\phi 12}v_1v_S & \lambda_\phi v_S^2\end{pmatrix}, \tag{A10}$$

$$M_\eta^2=\begin{pmatrix}-\frac{\lambda_{\phi 12}v_2v_S^2}{2v_1}-\lambda_5^{\rm eff}v_2^2 & \frac{1}{2}\lambda_{\phi 12}v_S^2+\lambda_5^{\rm eff}v_1v_2 & \lambda_{\phi 12}v_2v_S\\ \frac{1}{2}\lambda_{\phi 12}v_S^2+\lambda_5^{\rm eff}v_1v_2 & -\frac{\lambda_{\phi 12}v_1v_S^2}{2v_2}-\lambda_5^{\rm eff}v_1^2 & -\lambda_{\phi 12}v_1v_S\\ \lambda_{\phi 12}v_2v_S & -\lambda_{\phi 12}v_1v_S & -2\lambda_{\phi 12}v_1v_2\end{pmatrix}, \tag{A11}$$

$$M_w^2=\frac{1}{2}\begin{pmatrix}-\frac{\lambda_{\phi 12}v_2v_s^2}{v_1}-(\lambda_4+\lambda_5^{\rm eff})v_2^2 & \lambda_{\phi 12}v_s^2+(\lambda_4+\lambda_5^{\rm eff})v_1v_2\\ \lambda_{\phi 12}v_s^2+(\lambda_4+\lambda_5^{\rm eff})v_1v_2 & -\frac{\lambda_{\phi 12}v_1v_s^2}{v_2}-(\lambda_4+\lambda_5^{\rm eff})v_1^2\end{pmatrix}. \tag{A12}$$

The eigenvalues of the CP-odd and charged-Higgs masses are then given by

$$M_A^2 = -\lambda_{\phi 12}\left(\frac{v_S^2 v^2}{2v_1v_2} + 2v_1v_2\right) - \lambda_5^{\rm eff} v^2\,, \tag{A13}$$

$$M_{H^\pm}^2 = -\frac{\lambda_{\phi 12} v_S^2 v^2}{2v_1v_2} - \left(\lambda_4 + \lambda_5^{\rm eff}\right)\frac{v^2}{2}\,. \tag{A14}$$

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