Alignment, reverse alignment, and wrong sign Yukawa couplings in two Higgs doublet models

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(Dated: August 8, 2016)

Abstract

We consider two Higgs doublet models with a softly broken U(1) symmetry, for various limiting values of the scalar mixing angles $\alpha$ and $\beta$. These correspond to the Standard Model Higgs particle being the lighter CP-even scalar (alignment) or the heavier CP-even scalar (reverse alignment), and also the limit in which some of the Yukawa couplings of this particle are of the opposite sign from the vector boson couplings (wrong sign). In these limits we impose a criterion for naturalness by demanding that quadratic divergences cancel at one loop. We plot the allowed masses of the remaining physical scalars based on naturalness, stability, perturbative unitarity and constraints coming from the $\rho$ parameter. We also calculate the $h \to \gamma\gamma$ decay rate in the wrong sign limit.
I. INTRODUCTION

The discovery of a new boson in July 2012 by the ATLAS [1] and CMS Collaborations [2] at the Large Hadron Collider (LHC) is a landmark in the history of Particle Physics. This scalar is most likely the Higgs boson which is the last missing block in the Standard Model (SM). Although it answers most of the questions concerning fundamental particles, the SM has a few shortcomings, thus encouraging a search for theories beyond the Standard Model. Among the inadequacies are the lack of clear answers on the questions of the origins of neutrino mass and dark matter. It also cannot provide the observed matter-antimatter asymmetry of the universe.

One of the simplest ways to go beyond the SM is by extending the scalar sector. This of course affects the $\rho$ parameter, whose deviation from the tree level value of unity is a measure of new physics. The general expression for the tree level $\rho$ parameter for an SU(2)$\times$U(1) gauge theory with $N$ scalar multiplets is [3]

$$\rho \equiv \frac{m_W^2}{\cos^2 \theta_W m_Z^2} = \sum_{i=1}^{N} \left[ T_i (T_i + 1) - \frac{1}{4} Y_i^2 \right] v_i^2,$$

(1)

where $T_i$ and $Y_i$ denote the weak isospin and hypercharge of the $i^{th}$ scalar multiplet respectively, and $v_i$ is the vacuum expectation value (vev) of the neutral component of that multiplet. If the scalar sector contains only SU(2) singlets with $Y = 0$ and doublets with $Y = \pm 1$, then $\rho = 1$ is automatically satisfied without requiring any fine tuning among the vevs. This conforms with the experimental value of $\rho$, which is very close to unity [4]. We therefore confine our discussions to the doublet extensions, specifically the two Higgs-doublet models (2HDMs) [5], which have received a lot of attention mainly because the Type II 2HDM arises as part of minimal supersymmetry.

In this paper we consider the restrictions imposed on the scalar masses by a criterion of naturalness, embodied in the Veltman conditions, in various limits of 2HDMs of all types. The alignment limit and the reverse alignment limit are two scenarios in which the lighter and the heavier CP-even neutral scalar, respectively, correspond to the observed Higgs particle. We also consider the cases where these occur in conjunction with the wrong sign limit, in which the Yukawa coupling of at least one type of fermion is of the opposite sign as the vector coupling. Using the naturalness conditions we analyze the parameter space of masses of scalars in 2HDMs of different types. The parameter space is further restricted
by constraints arising from the $\rho$-parameter, global stability of the scalar potential, and requirement of perturbative unitarity. Section II gives a brief review of 2HDM. Sections III and IV deal with various limits of two Higgs doublet models and their permutations. In section V we calculate the Higgs-diphoton decay width for one of the scenarios and section VI concludes with a discussion of the results.

II. BRIEF REVIEW OF 2HDMS

We will work with the scalar potential [6, 7] considered under the imposition of a U(1) symmetry which forbids flavor-changing neutral currents (FCNCs),

$$V = \lambda_1 \left( |\Phi_1|^2 - \frac{v_1^2}{2} \right)^2 + \lambda_2 \left( |\Phi_2|^2 - \frac{v_2^2}{2} \right)^2 + \lambda_3 \left( |\Phi_1|^2 + |\Phi_2|^2 - \frac{v_1^2 + v_2^2}{2} \right)^2 + \lambda_4 \left( |\Phi_1|^2 |\Phi_2|^2 - |\Phi_1^* \Phi_2|^2 \right) + \lambda_5 \left| \Phi_1^* \Phi_2 - \frac{v_1 v_2}{2} \right|^2,$$

with real $\lambda_i$. This potential is invariant under the symmetry $\Phi_1 \to e^{i\theta} \Phi_1, \Phi_2 \to \Phi_2$, except for a soft breaking term $\lambda_5 v_1 v_2 \Re(\Phi_1^* \Phi_2)$. Additional dimension-4 terms, including one allowed by a softly broken $Z_2$ symmetry [8] are also set to zero by this U(1) symmetry. This is the same U(1) symmetry which prevents FCNC by having left- and right-handed fermions transform differently under it, leading to the four types of 2HDMs.

The scalar doublets are parametrized as

$$\Phi_i = \begin{pmatrix} w_i^+(x) \\ v_i + h_i(x) + i z_i(x) \end{pmatrix} \sqrt{2}, \quad i = 1, 2,$$

where the VEVs $v_i$ may be taken to be real and positive without any loss of generality. Three of these fields get “eaten” by the $W^\pm$ and $Z^0$ gauge bosons; the remaining five are physical scalar fields. There is a pair of charged scalars denoted by $\xi^\pm$, two neutral CP-even scalars $H$ and $h$, and one CP-odd pseudoscalar denoted by $A$. The two CP-even scalars have distinct masses, and $m_h < m_H$. With

$$\tan \beta = \frac{v_2}{v_1},$$
the scalar fields are given by the combinations
\[
\begin{pmatrix}
\omega^+ \\
\xi^+
\end{pmatrix} = \begin{pmatrix}
c_\beta & s_\beta \\
-s_\beta & c_\beta
\end{pmatrix} \begin{pmatrix}
w_1^+ \\
w_2^+
\end{pmatrix},
\]
(5)

\[
\begin{pmatrix}
\zeta \\
A
\end{pmatrix} = \begin{pmatrix}
c_\beta & s_\beta \\
-s_\beta & c_\beta
\end{pmatrix} \begin{pmatrix}
z_1 \\
z_2
\end{pmatrix},
\]
(6)

\[
\begin{pmatrix}
H \\
h
\end{pmatrix} = \begin{pmatrix}
c_\alpha & s_\alpha \\
-s_\alpha & c_\alpha
\end{pmatrix} \begin{pmatrix}
h_1 \\
h_2
\end{pmatrix},
\]
(7)

where \(c_\alpha \equiv \cos \alpha\), etc. We will assume, without loss of generality, that \(0 \leq \beta \leq \frac{\pi}{2}\), and \(-\frac{\pi}{2} \leq \alpha \leq \frac{\pi}{2}\).

The quartic couplings are related to the physical Higgs masses by [9, 10]:

\[
\lambda_1 = \frac{1}{2v^2 c_\beta^2} \left[ c_\alpha^2 m_H^2 + s_\alpha^2 m_h^2 - \frac{s_\alpha c_\alpha}{\tan \beta} (m_H^2 - m_h^2) \right] - \frac{\lambda_5}{4} (\tan^2 \beta - 1),
\]
(8)

\[
\lambda_2 = \frac{1}{2v^2 s_\beta^2} \left[ s_\alpha^2 m_H^2 + c_\alpha^2 m_h^2 - s_\alpha c_\alpha \tan \beta (m_H^2 - m_h^2) \right] - \frac{\lambda_5}{4} \left( \frac{1}{\tan^2 \beta} - 1 \right),
\]
(9)

\[
\lambda_3 = \frac{1}{2v^2 s_\beta c_\beta} (m_H^2 - m_h^2) - \frac{\lambda_5}{4},
\]
(10)

\[
\lambda_4 = \frac{2}{v^2} m_\xi^2,
\]
(11)

\[
\lambda_5 = \frac{2}{v^2} m_A^2.
\]
(12)

Let us now turn our attention to the fermion couplings. The scalar doublets couple to the fermions in the theory via the Yukawa Lagrangian

\[
L_Y = \sum_{i=1,2} \left[ -\bar{l}_L \Phi_i G_e^i e_R - \bar{Q}_L \Phi_i G_u^i u_R - \bar{Q}_L \Phi_i G_d^i d_R + h.c. \right].
\]
(13)

Here \(l_L, Q_L\) are 3-vectors of isodoublets in the space of generations, \(e_R, u_R, d_R\) are 3-vectors of singlets, \(G_e^i\) etc. are complex 3 \(\times\) 3 matrices in generation space containing the Yukawa coupling constants, and \(\Phi_i = i\tau_2 \Phi_i^*\).

When the fermions are in mass eigenstates, the Yukawa matrices are automatically diagonal if there is only one Higgs doublet as in the Standard Model. But in the presence of a second scalar doublet, the two Yukawa matrices will not be simultaneously diagonalizable in general. Thus the Yukawa couplings will not be flavor diagonal, and neutral Higgs scalars will mediate FCNCs [11–13]. The necessary and sufficient condition for the absence of FCNCs at tree level is that all fermions of a given charge and helicity transform according to
the same irreducible representation of SU(2), corresponding to the same eigenvalue of $T_3$, and that a basis exists in which they receive their contributions in the mass matrix from a single source [14, 15].

For the fermions of the Standard Model, this theorem implies that all right-handed singlets of a given charge must couple to the same Higgs doublet. This can be ensured by using the global U(1) symmetry mentioned earlier, which generalizes a $Z_2$ symmetry more commonly employed for this purpose. The left handed fermion doublets remain unchanged under this symmetry, $Q_L \rightarrow Q_L, l_L \rightarrow l_L$. The transformations of right handed fermion singlets determine the type of 2HDM. There are four such possibilities, which may be identified by the right-handed fields which transform under the U(1): type I (none), type II ($d_R \rightarrow e^{-i\theta}d_R, e_R \rightarrow e^{-i\theta}e_R$), lepton specific ($e_R \rightarrow e^{-i\theta}e_R$), flipped ($d_R \rightarrow e^{-i\theta}d_R$).

The scalar masses get quadratically divergent contributions which require very large fine-tuning of parameters. We will impose a criterion of naturalness on the scalar masses, viz., the cancellation of these quadratic divergences. This gives rise to four mass relations, which we may call the Veltman conditions for the 2HDMs being considered [16],

$$2 \text{Tr} G_1^e G_1^e + 6 \text{Tr} G_1^u G_1^u + 6 \text{Tr} G_1^d G_1^d = \frac{9}{4} g^2 + \frac{3}{4} g'^2 + 6 \lambda_1 + 10 \lambda_3 + \lambda_4 + \lambda_5 ,$$

$$2 \text{Tr} G_2^e G_2^e + 6 \text{Tr} G_2^u G_2^u + 6 \text{Tr} G_2^d G_2^d = \frac{9}{4} g^2 + \frac{3}{4} g'^2 + 6 \lambda_2 + 10 \lambda_3 + \lambda_4 + \lambda_5 ,$$

$$2 \text{Tr} G_1^e G_2^e + 6 \text{Tr} G_1^u G_2^u + 6 \text{Tr} G_1^d G_2^d = 0 ,$$

and another one which is the complex conjugate of the third equation. Here $g, g'$ are the $SU(2)$ and $U(1)_Y$ coupling constants, respectively.

The fermion mass matrix is diagonalized by independent unitary transformations on the left and right-handed fermion fields. In any of the 2HDMs, the U(1) symmetry implies that either $G_{1f}$ or $G_{2f}$ must vanish for each fermion type $f$. For example, in the Type II model $\Phi_1$ couples to down-type quarks and charged leptons, while $\Phi_2$ couples to up-type quarks, so $G_{2e} = G_{2d} = G_{1u} = 0$. Thus Eq. (16) is automatically satisfied in each 2HDM, and the relevant mass relations come from the first two equations above. The non-vanishing Yukawa matrices are related to the fermion masses by

$$\text{Tr} [G_{1f}^\dagger G_{1f}] = \frac{2}{v^2 \cos^2 \beta} \sum m_f^2 ,$$

$$\text{Tr} [G_{2f}^\dagger G_{2f}] = \frac{2}{v^2 \sin^2 \beta} \sum m_f^2 ,$$

5
where $f$ stands for charged leptons, up-type quarks, or down-type quarks, and the sum is taken over generations. These and the scalar mass relations of Eqs. (8) – (12) allow us to write the Veltman conditions in terms of the physical masses of particles.

There are some additional conditions on the parameters which further constrain the scalar masses. One is the perturbativity condition, which puts a constraint on the quartic coupling constants, $\lambda_i \leq 4\pi$ [17]. Another set comes from the condition that the potential is bounded from below. This was examined for more general potentials in 2HDM under U(1) symmetry in [18, 19], and for the potential given in Eq. (2) these conditions become

\begin{align*}
\lambda_1 + \lambda_3 &> 0, \\
\lambda_2 + \lambda_3 &> 0, \\
2\lambda_3 + \lambda_4 + 2\sqrt{(\lambda_1 + \lambda_3)(\lambda_2 + \lambda_3)} &> 0, \\
2\lambda_3 + \lambda_5 + 2\sqrt{(\lambda_1 + \lambda_3)(\lambda_2 + \lambda_3)} &> 0.
\end{align*}

(19) 
(20) 
(21) 
(22)

These conditions put lower bounds on the above combinations of quartic couplings, but there are also upper bounds on these couplings arising from the considerations of perturbative unitarity [20]. These conditions are

\begin{align*}
|2\lambda_3 - \lambda_4 + 2\lambda_5| &\leq 16\pi, \\
|2\lambda_3 + \lambda_4| &\leq 16\pi, \\
|2\lambda_3 + \lambda_5| &\leq 16\pi, \\
|2\lambda_3 + 2\lambda_4 - \lambda_5| &\leq 16\pi, \\
|3(\lambda_1 + \lambda_2 + 2\lambda_3) \pm \sqrt{9(\lambda_1 - \lambda_2)^2 + (4\lambda_3 + \lambda_4 + \lambda_5)^2}| &\leq 16\pi, \\
|(\lambda_1 + \lambda_2 + 2\lambda_3) \pm \sqrt{(\lambda_1 - \lambda_2)^2 + (\lambda_4 - \lambda_5)^2}| &\leq 16\pi, \\
|(\lambda_1 + \lambda_2 + 2\lambda_3) \pm (\lambda_1 - \lambda_2)| &\leq 16\pi.
\end{align*}

(23) 
(24) 
(25) 
(26) 
(27) 
(28) 
(29)

There is another condition that we need to take into account when we calculate bounds on the scalar masses. The oblique electroweak correction $T$, which measures deviations from the standard model due to new physics, is related to the deviation of the $\rho$ parameter from its SM value of unity by

$$
\delta \rho \equiv \rho - 1 = \alpha T,
$$

(30)

where $\alpha = e^2/4\pi$ is the fine structure constant. The effect of the general 2HDM on the $\rho$
parameter is known to be \([22, 23]\)

\[
\delta \rho = \frac{g^2}{64\pi^2 m_Z^2} \left( F(m_{\xi}^2, m_A^2) + \sin^2(\beta - \alpha) F(m_{\xi}^2, m_H^2) + \cos^2(\beta - \alpha) F(m_{\xi}^2, m_h^2) \right.
\]

\[
- \sin^2(\beta - \alpha) F(m_A^2, m_H^2) - \cos^2(\beta - \alpha) F(m_A^2, m_h^2)
\]

\[
+ 3 \cos^2(\beta - \alpha) \left[ F(m_Z^2, m_H^2) - F(m_W^2, m_H^2) \right]
\]

\[
+ 3 \sin^2(\beta - \alpha) \left[ F(m_Z^2, m_h^2) - F(m_W^2, m_h^2) \right]
\]

\[
- 3 \left[ F(m_Z^2, m_{h_{SM}}^2) - F(m_W^2, m_{h_{SM}}^2) \right],
\]

where \(F(x, y)\) is a function of two non-negative arguments \(x\) and \(y\), symmetrical under the exchange of the arguments and vanishes only if \(x = y\). The function has the property that it grows linearly with \(\max(x, y)\), i.e., quadratically with the heaviest scalar mass when that mass becomes very large. The current experimental bound on the total new physics contribution to \(\rho\) is given by \(\delta \rho = -0.00011\) \([4]\).

### III. LIMITS OF 2HDMS

In order to relate a 2HDM to the Higgs sector of the Standard Model, we need to identify some combination of the neutral scalar particles in the theory as the observed Higgs particle. This can be done in several ways, by considering different combinations of the angles \(\alpha\) and \(\beta\). In this section we will consider the different limits for which part of the 2HDM matches the Standard Model, and calculate the allowed range of masses for the additional scalars.

A crucial parameter of the 2HDMS is \(\tan \beta\). Its value is larger than one, based on constraints coming from \(Z \to b\bar{b}\) and \(B_q \bar{B}_q\) mixing \([24]\). A large \(\tan \beta\) is suggested by muon \(g - 2\) in lepton specific 2HDM \([25]\), by using \(b \to s\gamma\) in type I and flipped models \([26]\), which also suppresses the \(t \to bH^+\) branching ratio to a rough agreement with 95\% CL limits from the light charged Higgs searches at the LHC \([27, 28]\). We will assume that \(\tan \beta\) is large, and certainly larger than unity, specific values will be considered for the plots as needed.

#### A. Alignment Limit

If we rotated the neutral \((h_1, h_2)\) doublet by the angle \(\beta\),

\[
\begin{pmatrix}
H^0 \\
R
\end{pmatrix}
= \begin{pmatrix}
c_\beta & s_\beta \\
-s_\beta & c_\beta
\end{pmatrix}
\begin{pmatrix}
h_1 \\
h_2
\end{pmatrix},
\]
we would find that $H^0$ has exactly the Standard Model Higgs couplings with the fermions and gauge bosons [11, 18]. The physical scalar $h$ is related to $H^0$ and $R$ via

$$h = \sin(\beta - \alpha)H^0 + \cos(\beta - \alpha)R.$$  

(33)

Thus in order for $h$ to be the Higgs boson of the Standard Model, we require $\sin(\beta - \alpha) \approx 1$, which has been called the SM-like or alignment limit [21].

There remain three unknown mass parameters, namely $m_H, m_\xi$ and $m_A$, which span the parameter space. By fixing $\tan \beta$ at some specific value, we can use the Veltman conditions to plot the accessible region of the $m_H - m_\xi$ plane corresponding to the allowed range of values for $m_A$. On the other hand, constraints from perturbative unitarity and the oblique correction $T$ also restrict the accessible region on this plane. The intersection of all these regions provides the allowed ranges for $m_H$ and $m_\xi$.

The mass ranges were studied for the alignment limit in [29], where it was found that if we set $m_h = 125$ GeV, and allowed $m_A$ to run over its entire range of $0 < m_A \lesssim 617$ GeV as determined by the condition of perturbativity, the two unknown masses $m_H$ and $m_\xi$ became restricted to ranges of $550 \text{ GeV} \lesssim m_\xi \lesssim 700 \text{ GeV}$, $450 \text{ GeV} \lesssim m_H \lesssim 620 \text{ GeV}$. The value of $\tan \beta$ used in these calculations was $\tan \beta = 5$, and it was also found that a higher value of $\tan \beta$ pushed the ranges to higher values and also made them narrower. These mass ranges are in agreement with bounds found by analysing experimental data [30].

B. Reverse Alignment Limit

Let us rearrange the equations described in the previous section. Using Eqs. (7) and (32) we obtain $H$ in terms of $H^0$ and $R$,

$$H = H^0 \cos(\beta - \alpha) - R \sin(\beta - \alpha)$$  

(34)

Had $H$ been the SM-like Higgs boson, it would have to resemble the properties of $H^0$, and for that $\beta$ would have to approximately equal $\alpha$ or $\pi + \alpha$. The ultimate results with $\beta \approx \alpha$ and $\beta \approx \pi + \alpha$ are identical, so in what follows we will work with $\beta \approx \alpha$ and call it the Reverse Alignment Limit.
Eqs. (8-12) become, in the reverse alignment limit,
\[
\begin{align*}
\lambda_1 &= \frac{m_h^2}{2v^2}(\tan^2 \beta + 1) - \frac{\lambda_5}{4}(\tan^2 \beta - 1), \\
\lambda_2 &= \frac{m_h^2}{2v^2}(\cot^2 \beta + 1) - \frac{\lambda_5}{4}(\cot^2 \beta - 1), \\
\lambda_3 &= \frac{1}{2v^2}(m_H^2 - m_h^2) - \frac{\lambda_5}{4}, \\
\lambda_4 &= \frac{2}{v^2}m_\xi^2, \\
\lambda_5 &= \frac{2}{v^2}m_A^2.
\end{align*}
\]  
(35)  
(36)  
(37)  
(38)  
(39)

Let us write the Veltman conditions defined in Eqs. (14) and (15) using the above equations.

We will write the equations explicitly for one case, that of the Type II 2HDM, for which the two Veltman conditions read, in the reverse alignment limit,
\[
\begin{align*}
m_h^2 (3 \tan^2 \beta - 2) + 2m_\xi^2 &= 4 \left[ \sum m_{\ell \ell}^2 + 3 \sum m_d^2 \right] \sec^2 \beta - 6M_W^2 - 3M_Z^2 - 5m_H^2 + \lambda_5 \frac{3v^2}{2} \tan \beta, \\
m_h^2 (3 \cot^2 \beta - 2) + 2m_\xi^2 &= 12 \left[ m_u^2 \csc^2 \beta - 6M_W^2 - 3M_Z^2 - 5m_H^2 + \lambda_5 \frac{3v^2}{2} \cot^2 \beta \right].
\end{align*}
\]  
(40)  
(41)  

We have plotted the above equalities on the \( m_h - m_\xi \) plane for several values of \( \lambda_5 \) for a fixed value of \( \tan \beta \) and with \( m_H = 125 \) GeV, with \( m_h \leq m_H \). On the same plane, we have also plotted the region allowed by stability, perturbative unitarity, and constraints from \( \delta \rho \). The conditions of stability and perturbative unitarity, Eq. (19) – Eq. (29), produce the following two inequalities in the reverse alignment limit relevant to this plot:
\[
\begin{align*}
0 &\leq (m_h^2 - m_A^2) (\tan^2 \beta + \cot^2 \beta) + 2m_H^2 \leq \frac{32\pi v^2}{3}, \\
|2m_\xi^2 - m_h^2 - m_A^2 + m_H^2| &\leq 16\pi v^2.
\end{align*}
\]  
(42)  
(43)

These are analogous to similar inequalities found in [29] in the alignment limit.

For \( \tan \beta = 5 \), the plots for all four types of 2HDM are shown in Fig. 1. The gray region covers the points which satisfy the inequalities (42) and (43) in addition to the constraints from \( \delta \rho \), the first Veltman condition provides the curves (ellipses) which cross this region, and the second Veltman condition provides the nearly flat hyperbolas above the gray region.

As we can see from the plots in Fig. 1, there is no region on the \( m_h - m_\xi \) plane where all the constraints are obeyed. In other words, if we insist on naturalness, as embodied by the Veltman conditions, the reverse alignment limit is not a valid limit for any of the 2HDMs, i.e. the observed Higgs particle cannot be the heavier CP-even neutral scalar in any of the 2HDMs.
FIG. 1. Allowed mass range (in GeV) for the charged Higgs and the light CP even Higgs in Reverse alignment limit for (a) type I (b) type II (c) lepton specific and (d) flipped 2HDM for $|\lambda_5| \leq 4\pi$ and $\tan\beta = 5$.

It should be mentioned here that allowed mass ranges of scalars in both the alignment limit and the reverse alignment limit were studied in [31]. However, that paper considered an unbroken $Z_2$ symmetry, not a softly broken symmetry as we have considered. As a result the mass ranges of scalars, as well as the allowed range of $\tan\beta$ found in that paper, are different from the ones we have found.

IV. WRONG SIGN YUKAWA COUPLINGS

The wrong-sign Yukawa coupling regime [21, 32, 33] is defined as the region of 2HDM parameter space in which at least one of the couplings of the SM-like Higgs to up-type and down-type quarks is opposite in sign to the corresponding coupling of SM-like Higgs
to vectors bosons. This is to be contrasted with the Standard Model, where the couplings of \( h_{SM} \) to \( \bar{f}f \) and vector bosons are of the same sign. The *wrong sign limit* needs to be considered in conjunction with either the alignment limit or the reverse alignment limit. We will now calculate the regions of parameter space when each of these two limits are combined with the wrong sign limit.

The CP-even neutral scalars couple to the up-type and down-type quarks in the various 2HDMs as shown in Table I with the SM couplings of the quarks to the SM Higgs field normalized to unity.

| 2HDMs     | \( h\bar{U} \) | \( h\bar{D} \) | \( H\bar{U} \) | \( H\bar{D} \) |
|-----------|----------------|----------------|----------------|----------------|
| Type I    | \( \cos \alpha \) \( \sin \beta \) | \( \cos \alpha \) \( \sin \beta \) | \( \sin \alpha \) \( \sin \beta \) | \( \sin \alpha \) \( \sin \beta \) |
| Type II   | \( \cos \alpha \) \( \sin \beta \) | \( - \sin \alpha \) \( \cos \beta \) | \( \sin \alpha \) \( \sin \beta \) | \( \cos \alpha \) \( \sin \beta \) |
| Lepton Specific | \( \cos \alpha \) \( \sin \beta \) | \( \cos \alpha \) \( \sin \beta \) | \( \sin \alpha \) \( \sin \beta \) | \( \sin \alpha \) \( \sin \beta \) |
| Flipped   | \( \cos \alpha \) \( \sin \beta \) | \( - \sin \alpha \) \( \cos \beta \) | \( \sin \alpha \) \( \sin \beta \) | \( \cos \alpha \) \( \sin \beta \) |

TABLE I. Yukawa couplings for the different 2HDMs

### A. Wrong Sign and Reverse alignment limit

Let us first consider the case of wrong sign Yukawa couplings in the reverse alignment limit. The heavier CP-even neutral scalar \( H \) corresponds to the SM Higgs in the reverse alignment limit, with a coupling to vector bosons which is \( \cos(\beta - \alpha) \) times the corresponding SM value. In the convention where \( \cos(\beta - \alpha) \geq 0 \), the \( HVV \) couplings in the 2HDM are always non-negative. To analyze the wrong-sign coupling regime, we write the Yukawa couplings in the type-II and Flipped 2HDMs in the following form:

\[
H\bar{D}D : \quad \frac{\cos \alpha}{\cos \beta} = \cos(\beta + \alpha) + \sin(\beta + \alpha) \tan \beta , \quad (44)
\]

\[
H\bar{U}U : \quad \frac{\sin \alpha}{\sin \beta} = -\cos(\beta + \alpha) + \sin(\beta + \alpha) \cot \beta . \quad (45)
\]

In the case when \( \cos(\beta + \alpha) = -1 \), the \( H\bar{D}D \) coupling normalized to its SM value is equal to \(-1\), whereas the normalized \( H\bar{U}U \) coupling is \(+1\). Thus in this case, when the reverse alignment limit is taken in conjunction with the wrong sign limit, we have \( \alpha \approx \beta \approx \frac{\pi}{2} \). It turns out there is no point on the \( m_h - m_\xi \) plane which satisfies the Veltman conditions as
FIG. 2. Veltman conditions are not satisfied for any \((m_h, m_\xi)\) satisfying unitarity and other bounds, in the reverse alignment limit with wrong sign Yukawa couplings.

well as the bounds coming from unitarity, stability and the \(\rho\)-parameter. In Fig. 2 only the first Veltman condition has been plotted, and it does not cross the grey region corresponding to the bounds. The other Veltman condition does not show up in this picture at all, it is not satisfied for any point in this plot.

On the other hand, in the case when \(\cos(\beta + \alpha) = 1\), the \(H\bar{U}U\) coupling normalized to its SM value is equal to \(-1\), while the normalized \(H\bar{D}D\) coupling is \(+1\). In this limiting case, \(\cos(\beta - \alpha) = \cos 2\beta\), which implies that the wrong-sign \(H\bar{U}U\) couplings can only be achieved for \(\tan \beta < 1\) for the type II and Flipped 2HDMs.

In the type-I and lepton specific 2HDMs, both the \(H\bar{D}D\) and \(H\bar{U}U\) couplings are given by Eq. (45). Thus, for \(\cos(\beta + \alpha) = 1\), both the normalized \(H\bar{D}D\) and \(H\bar{U}U\) couplings are equal to \(-1\), which is only possible if \(\tan \beta < 1\).

Since \(\tan \beta > 1\), we see that the wrong-sign Yukawa coupling is incompatible with the reverse alignment limit in all of the four types of 2HDMs.

**B. Wrong sign in the Alignment limit**

Let us now look at what happens if some Yukawa couplings are of the wrong sign, in the alignment limit. In this case \(h\) is the SM Higgs, and its coupling to the vector bosons is \(\sin(\beta - \alpha)\) times the corresponding SM value. Then in the convention where \(\sin(\beta - \alpha) \geq 0\), the \(hVV\) couplings in the 2HDM are always non-negative. As in the previous case, we write the type-II and Flipped Higgs-fermion Yukawa couplings, normalized with respect to the
Standard Model couplings, in the following form:

\begin{align}
    h\bar{D}D : & \quad \frac{\sin \alpha}{\cos \beta} = -\sin(\beta + \alpha) + \cos(\beta + \alpha) \tan \beta, \\
    h\bar{U}U : & \quad \frac{\cos \alpha}{\sin \beta} = \sin(\beta + \alpha) + \cos(\beta + \alpha) \cot \beta.
\end{align}

In the case when \(\sin(\beta + \alpha) = 1\), the \(h\bar{D}D\) coupling normalized to its SM value is equal to \(-1\), while the normalized \(h\bar{U}U\) coupling is \(+1\). Note that in this limiting case, \(\sin(\beta - \alpha) = -\cos 2\beta\), which implies that the wrong-sign \(h\bar{D}D\) Yukawa coupling can only be achieved for values of \(\tan \beta > 1\).

Likewise, in the case of \(\sin(\beta + \alpha) = -1\), the \(h\bar{U}U\) coupling normalized to its SM value is equal to \(-1\), whereas the normalized \(h\bar{D}D\) coupling is \(+1\). Then \(\sin(\beta - \alpha) = \cos 2\beta\), which implies that the wrong-sign \(h\bar{U}U\) couplings can occur only if \(\tan \beta < 1\). In the type-I and lepton specific 2HDM, both the \(h\bar{D}D\) and \(h\bar{U}U\) couplings are given by Eq. (47). Thus for \(\sin(\beta + \alpha) = -1\), both the normalized \(h\bar{D}D\) and \(h\bar{U}U\) couplings are equal to \(-1\), which is only possible if \(\tan \beta < 1\). Thus realistically only the \(h\bar{D}D\) coupling of the type-II and flipped 2HDM can be of the wrong sign, since \(\tan \beta > 1\).

Let us therefore consider a type II model with a wrong sign \(h\bar{D}D\) coupling. The wrong sign limit approaches the alignment limit for \(\tan \beta \approx 17\) as was displayed in [32, 33] for the allowed parameter space of the type II CP-conserving 2HDM, based on the 8 TeV run of the LHC. For this model, we will plot the values of the pair \((m_H, m_\xi)\) allowed by the naturalness conditions as well as the constraints imposed by perturbativity, stability, tree-level unitarity, and the \(\rho\) parameter. We will do this for four different values of \(\tan \beta\) around the ‘critical’ value of 17. By choosing a small enough \(\alpha\) we can ensure that for all these choices, both \(\sin(\beta - \alpha) \approx 1\) and \(\sin(\beta + \alpha) \approx 1\), as needed for the alignment limit and the wrong sign coupling.

In Fig. 3 we have plotted the Veltman conditions on the \(m_H - m_\xi\) plane for Type II 2HDM for the four choices of \(\tan \beta\), for different values of \(m_A\) constrained by \(|\lambda_5| \leq 4\pi\). This plots are further constrained by conditions coming from stability of the potential, perturbative unitarity, and experimental bounds on \(\delta \rho\). We have also taken \(m_h = 125\) GeV. One can estimate from the plots that for \(\tan \beta = 17\) that the range of \(m_H\) is approximately \((250, 330)\) GeV, and that of \(m_\xi\) is approximately \((260, 310)\) GeV. At higher values of \(\tan \beta\), both ranges become narrower and move down on the mass scale.
FIG. 3. Allowed mass range in GeV for the charged Higgs and the heavy CP even Higgs when approaching wrong sign and alignment limits simultaneously for (a) $\tan \beta = 10$ (b) $\tan \beta = 17$ (c) $\tan \beta = 20$ and (d) $\tan \beta = 30$ for $|\lambda_5| \leq 4\pi$ and Type II 2HDM.

V. MODIFICATION OF HIGGS-DIPHOTON DECAY WIDTH

The $h \rightarrow \gamma \gamma$ decay channel is perhaps the most popular channel for Higgs and related searches. The decay width can be enhanced or reduced in the 2HDMs due to loop effects. In the alignment limit, the couplings of the lighter CP even neutral scalar $h$ to gauge bosons are identical to that for the SM Higgs. Then the tree level decay widths of $h$ will be the same as for the SM Higgs. For loop induced decays, such as $h \rightarrow \gamma \gamma$ and $h \rightarrow Z\gamma$, the contribution of the $W$ boson loop and the top loop diagrams are the same as in the SM. But there will have some additional contributions due to the virtual charged scalars $\xi^\pm$ in the loop. Thus the decay widths will be different from the SM in general. Contributions from the fermion loops are the same in this case as for the SM.
On the other hand, suppose \( h \) has wrong sign Yukawa couplings to the down-type quarks. Then the bottom quarks will contribute with a relative negative sign in the loops, and the \( h \to \gamma\gamma \) decay width will be different from the SM, as well as from 2HDMs in the usual alignment limit.

The Higgs-diphoton decay width is calculated using the formula [34]

\[
\Gamma(h \to \gamma\gamma) = \frac{G_\mu \alpha^2 m_h^3}{128 \sqrt{2} \pi^3} \left| \sum_f N_c Q_f^2 g_{hff} A_{1/2}^h(\tau_f) + g_{hVV} A_1^h(\tau_W) + \frac{m_W^2 \lambda_{h\xi^+\xi^-}}{2 c_W^2 M_{\xi^\pm}^2} A_0^h(\tau_{\xi^\pm}) \right|^2.
\]  (48)

In this equation, \( N_c \) is the number color multiplicity, \( Q_f \) is the charge of the fermion \( f \), \( G_\mu \) is the Fermi constant, and the reduced couplings \( g_{hff} \) and \( g_{hVV} \) of the Higgs boson to fermions and \( W \) bosons are \( g_{htt} = \frac{\cos \alpha}{\sin \beta} \), \( g_{hbb} = -\frac{\sin \alpha}{\cos \beta} \) and \( g_{hWW} = \sin(\beta - \alpha) \), while the trilinear \( \lambda_{h\xi^+\xi^-} \) couplings to charged Higgs bosons is given by

\[
\lambda_{h\xi^+\xi^-} = \cos 2\beta \sin(\beta + \alpha) + 2 c_W^2 \sin(\beta - \alpha)
\]  (49)
\[
= \lambda_{hAA} + 2 c_W^2 g_{hVV},
\]  (50)

where \( c_W = \cos \theta_W \), with \( \theta_W \) being the Weinberg angle. The decay rate does not depend on the type of the 2HDM.

The amplitudes \( A_i \) at lowest order for the spin-1, spin-\( \frac{1}{2} \) and spin-0 particle contributions are given by [7]

\[
A_{1/2}^h = -2\tau [1 + (1 - \tau) f(\tau)]
\]  (51)
\[
A_1^h = 2 + 3\tau + 3\tau(2 - \tau) f(\tau)
\]  (52)
\[
A_0^h = \tau [1 - \tau f(\tau)]
\]  (53)

in the case of the CP even Higgs boson \( h \).

Here

\[
\tau_x = 4 m_x^2 / m_h^2
\]  (54)

and

\[
f(\tau) = \begin{cases} 
\arcsin^2 \sqrt{1/\tau}, & \tau \geq 1 \\
-\frac{1}{4} \left[ \log \frac{1 + \sqrt{1 - \tau}}{1 - \sqrt{1 - \tau}} - i\pi \right]^2, & \tau < 1 
\end{cases}
\]  (55)

Using the above definitions in the decay width formula given in Eq. [48], we arrive at a much simplified expression for the decay width,

\[
\Gamma(h \to \gamma\gamma) = \frac{G_\mu \alpha^2 m_h^3}{128 \sqrt{2} \pi^3} \left| g_{hVV} A_W^h + \frac{4}{3} g_{htt} A_t^h \pm \frac{1}{3} g_{hbb} A_b^h + \kappa A_\xi^h \right|^2,
\]  (56)
where the $'+'$ sign before $A_h^h$ is for when the $h\bar{b}b$ Yukawa coupling has the same sign as the $hVV$ coupling and the $'-'$ sign is for the wrong sign of the Yukawa coupling, and $\kappa$ is defined as

$$\kappa = \frac{1}{m_\xi^2} \left( m_\xi^2 + \frac{1}{2} m_h^2 - m_A^2 \right).$$

(57)

The appearance of $m_A$ in Eq. (57) is merely an artefact of U(1) symmetry of the scalar potential. For a more general potential the expression for $\kappa$ involves $\lambda_5$ [35]. In Fig. 4 we have plotted the $h\to\gamma\gamma$ decay width in 2HDMs in the alignment limit, normalized with respect to the SM value, against the mass of the charged Higgs particle, and for different values of the mass of the CP-odd scalar. Fig. 4(a) shows the decay width for the case where the $h\bar{q}q$ Yukawa coupling has the same sign as the $hVV$ coupling, whereas Fig. 4(b) is for the decay width corresponding to the case where the Yukawa coupling of $h$ to the down-type quarks is of the opposite sign to the $hVV$ coupling. We note that the first case has been plotted, albeit for smaller values of $\tan\beta$ and without the use of the Veltman conditions (thus for a much larger range of $m_\xi$), in [36].

As we have seen in the previous section, simultaneously choosing the alignment limit and the wrong sign limit also sets $\tan\beta$ at a high value. The critical value $\tan\beta = 17$, and a small but non-zero value of $\alpha$, namely $\alpha \simeq 0.035$, was chosen for both the plots. The plots are not noticeably different for other high values of $\tan\beta$ or other similar values of $\alpha$. The decay width does not depend on the type of 2HDM once the masses of the charged Higgs particle and the CP-odd Higgs particle are fixed. However, the range of allowed
masses depends on the type of 2HDM being considered. We have chosen the ranges $225 \text{ GeV} \leq m_\xi \leq 290 \text{ GeV}$ and $200 \text{ GeV} \leq m_A \leq 300 \text{ GeV}$ which cover the allowed ranges for all four types for $\tan \beta = 17$. Although a picture is worth a thousand words, it is perhaps worth pointing out that when $m_A$ is small, for example $m_A \simeq 200 \text{ GeV}$, the diphoton decay width deviates from the SM value by 5-7% for all values of $m_\xi$. The deviation is noticeable for many other values of $m_A$ also, as can be easily seen from the plots. On the other hand, for specific choices of $(m_A, m_\xi)$ the $h \to \gamma\gamma$ decay width is the same as for the SM, so the non-observation of a deviation does not rule out 2HDMs.

The two plots are similar, but not identical. The decay width when the $h\bar{D}D$ Yukawa coupling is of the ‘wrong sign’ is smaller than the decay width for the case when it is of the same sign (as $hVV$ couplings) by about 1.5%, as can be seen from the ratio of the decay widths, displayed in Fig. 5.

![Relative Decay Width](image)

**FIG. 5.** $h\gamma\gamma$ decay width for ‘wrong sign’ $h\bar{D}D$ coupling relative to the case with ‘same sign’ Yukawa couplings

**VI. RESULTS AND CONCLUSION**

In this paper we have looked at how a certain criterion of naturalness, namely the cancellation of quadratic divergences, affect the allowed ranges of masses of the additional scalars in 2HDMs in the alignment or SM-like limit with ‘wrong sign’ Yukawa couplings, and also in the reverse alignment limit. A similar calculation was done in [29] for the alignment limit without the ‘wrong sign’ assumption.

We found that reverse alignment, *i.e.* the scenario in which the heavier CP-even neutral scalar is the Standard Model Higgs particle, is clearly not a viable scenario for 2HDMs. Con-
straints arising from naturalness, stability, perturbative unitarity and experimental bounds on the ρ-parameter completely rule out this scenario. The naturalness criterion is crucial for this conclusion – reverse alignment is an allowed scenario if quadratic divergences are taken care of by some mechanism of fine tuning, for example.

We have also considered a limit where the lighter CP-even neutral scalar corresponds to the SM-like Higgs but where the Yukawa couplings of this particle to D-type quarks are of the wrong sign relative to their gauge couplings. In this scenario we obtain mass ranges for the rest of the physical Higgs bosons for various benchmark values of tan β. In this paper we have shown only the plot for Type II 2HDM, but the results are similar for the other 2HDMs with a small variation of a few GeV.

The Higgs-diphoton decay width in a 2HDM receives additional contributions from loops containing the charged scalar ξ±, so the decay width in a 2HDM is different from the SM value. Further, in the wrong sign limit, loops containing down type quarks contribute with a different sign. We have plotted the h → 2γ decay width against the mass of the charged Higgs, and also for different values of the mass of the CP-odd neutral scalar, and found that the decay width can differ from its SM value by up tp 6% for some values of the parameters.

While this paper was being completed, another paper which investigates what we call the reverse alignment limit appeared as an e-print [37]. However, that paper uses fewer constraints, so limits on the masses of ξ± are less restrictive.

More recently, the ATLAS and CMS collaborations at the LHC have reported an excess corresponding to a diphoton resonance at 750 GeV [38]. We note that according to the naturalness criterion we have used in this paper, this excess cannot be one of the scalar particles in any of the four types of 2HDMs, in agreement with the negative result found in [39] using several other lines of argument.

ACKNOWLEDGEMENT

AB thanks Dipankar Das for useful discussions. The authors thank the anonymous referee for raising questions about the plot of the diphoton decay width in the first version, which
helped us find and correct a mistake in the original plot.

[1] G. Aad et al. [ATLAS Collaboration], “Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC,” Phys. Lett. B 716, 1 (2012).

[2] S. Chatrchyan et al. [CMS Collaboration], “Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC,” Phys. Lett. B 716, 30 (2012).

[3] P. Langacker, “Grand Unified Theories and Proton Decay,” Phys. Rept. 72, 185 (1981).

[4] K. A. Olive et al. [Particle Data Group Collaboration], “Review of Particle Physics,” Chin. Phys. C 38, 090001 (2014). doi:10.1088/1674-1137/38/9/090001

[5] G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, M. Sher and J. P. Silva, “Theory and phenomenology of two-Higgs-doublet models,” Phys. Rept. 516, 1 (2012)

[6] T. D. Lee, “A Theory of Spontaneous T Violation,” Phys. Rev. D 8, 1226 (1973).

[7] J. F. Gunion, H. E. Haber, G. L. Kane and S. Dawson, “The Higgs Hunter’s Guide,” Front. Phys. 80, 1 (2000).

[8] J. F. Gunion, H. E. Haber, G. L. Kane and S. Dawson, “Errata for the Higgs hunter’s guide,” arXiv:hep-ph/9302272.

[9] S. Kanemura, T. Kubota and E. Takasugi, “Lee-Quigg-Thacker bounds for Higgs boson masses in a two doublet model,” Phys. Lett. B 313, 155 (1993).

[10] A. G. Akeroyd, A. Arhrib and E. M. Naimi, “Note on tree level unitarity in the general two Higgs doublet model,” Phys. Lett. B 490, 119 (2000)

[11] G. C. Branco, W. Grimus and L. Lavoura, “Relating the scalar flavor changing neutral couplings to the CKM matrix,” Phys. Lett. B 380, 119 (1996) [hep-ph/9601383].

[12] F. J. Botella, G. C. Branco, A. Carmona, M. Nebot, L. Pedro and M. N. Rebelo, “Physical Constraints on a Class of Two-Higgs Doublet Models with FCNC at tree level,” JHEP 1407, 078 (2014) [arXiv:1401.6147 [hep-ph]].

[13] G. Bhattacharyya, D. Das and A. Kundu, “Feasibility of light scalars in a class of two-Higgs-doublet models and their decay signatures,” Phys. Rev. D 89, 095029 (2014) [arXiv:1402.0364 [hep-ph]].
[14] S. L. Glashow and S. Weinberg, “Natural Conservation Laws for Neutral Currents,” Phys. Rev. D 15, 1958 (1977).

[15] E. A. Paschos, “Diagonal Neutral Currents,” Phys. Rev. D 15, 1966 (1977).

[16] C. Newton and T. T. Wu, “Mass relations in the two Higgs doublet model from the absence of quadratic divergences,” Z. Phys. C 62, 253 (1994).

[17] S. Kanemura, T. Kasai and Y. Okada, “Mass bounds of the lightest CP even Higgs boson in the two Higgs doublet model,” Phys. Lett. B 471, 182 (1999) doi:10.1016/S0370-2693(99)01351-9 [hep-ph/9903289].

[18] J. F. Gunion and H. E. Haber, “The CP conserving two Higgs doublet model: The Approach to the decoupling limit,” Phys. Rev. D 67, 075019 (2003)

[19] M. Sher, “ Electroweak Higgs Potentials and Vacuum Stability,” Phys. Rept. 179, 273 (1989).

[20] B. W. Lee, C. Quigg and H. B. Thacker, “Weak Interactions at Very High-Energies: The Role of the Higgs Boson Mass,” Phys. Rev. D 16, 1519 (1977).

[21] P. M. Ferreira, J. F. Gunion, H. E. Haber and R. Santos, “Probing wrong-sign Yukawa couplings at the LHC and a future linear collider,” Phys. Rev. D 89, no. 11, 115003 (2014) [arXiv:1403.4736 [hep-ph]]

[22] W. Grimus, L. Lavoura, O. M. Ogreid and P. Osland, “A Precision constraint on multi-Higgs-doublet models,” J. Phys. G 35, 075001 (2008) [arXiv:0711.4022 [hep-ph]]

[23] S. Kanemura, Y. Okada, H. Taniguchi and K. Tsumura, “Indirect bounds on heavy scalar masses of the two-Higgs-doublet model in light of recent Higgs boson searches,” Phys. Lett. B 704, 303 (2011).

[24] A. Arhrib, R. Benbrik, C. H. Chen, R. Guedes and R. Santos, “Double Neutral Higgs production in the Two-Higgs doublet model at the LHC,” JHEP 0908, 035 (2009).

[25] J. Cao, P. Wan, L. Wu and J. M. Yang, “Lepton-Specific Two-Higgs Doublet Model: Experimental Constraints and Implication on Higgs Phenomenology,” Phys. Rev. D 80, 071701 (2009).

[26] J. h. Park, “Lepton non-universality at LEP and charged Higgs,” JHEP 0610, 077 (2006).

[27] G. Aad et al. [ATLAS Collaboration], “Search for charged Higgs bosons decaying via $H^+ \rightarrow \tau \nu$ in top quark pair events using pp collision data at $\sqrt{s} = 7$ TeV with the ATLAS detector,” JHEP 1206, 039 (2012).
[28] S. Chatrchyan et al. [CMS Collaboration], “Search for a light charged Higgs boson in top quark decays in pp collisions at $\sqrt{s} = 7$ TeV,” JHEP 1207, 143 (2012).

[29] A. Biswas and A. Lahiri, “Masses of physical scalars in two Higgs doublet models,” Phys. Rev. D 91, no. 11, 115012 (2015) [arXiv:1412.6187 [hep-ph]].

[30] S. Kanemura, H. Yokoya and Y. J. Zheng, “Complementarity in direct searches for additional Higgs bosons at the LHC and the International Linear Collider,” Nucl. Phys. B 886, 524 (2014).

[31] B. Coleppa, F. Kling and S. Su, “Constraining Type II 2HDM in Light of LHC Higgs Searches,” JHEP 1401, 161 (2014).

[32] P. M. Ferreira, R. Guedes, M. O. P. Sampaio and R. Santos, “Wrong sign and symmetric limits and non-decoupling in 2HDMs,” JHEP 1412, 067 (2014).

[33] P. M. Ferreira, R. Guedes, J. F. Gunion, H. E. Haber, M. O. P. Sampaio and R. Santos, “The Wrong Sign limit in the 2HDM,” arXiv:1410.1926 [hep-ph].

[34] A. Djouadi, “The Anatomy of electro-weak symmetry breaking. II. The Higgs bosons in the minimal supersymmetric model,” Phys. Rept. 459, 1 (2008) [hep-ph/0503173].

[35] A. Arhrib, M. Capdequi Peyranere, W. Hollik and S. Penaranda, Phys. Lett. B 579, 361 (2004) doi:10.1016/j.physletb.2003.10.006 [hep-ph/0307391].

[36] G. Bhattacharyya, D. Das, P. B. Pal and M. N. Rebelo, JHEP 1310, 081 (2013) doi:10.1007/JHEP10(2013)081 [arXiv:1308.4297 [hep-ph]].

[37] J. Bernon, J. F. Gunion, H. E. Haber, Y. Jiang and S. Kraml, “Scrutinizing the Alignment Limit in Two-Higgs-Doublet Models. Part 2: $m_H = 125$ GeV,” arXiv:1511.03682 [hep-ph].

[38] ATLAS 13 TeV Results - December 2015. Talk by Marumi Kado at CERN, and ATLAS note: ATLAS-CONF-2015-081; CMS 13 TeV Results - December 2015. Talk by Jim Olsen at CERN, and CMS note: CMS-PAS-EXO-15-004 .

[39] A. Angelescu, A. Djouadi and G. Moreau, “Scenarii for interpretations of the LHC diphoton excess: two Higgs doublets and vector-like quarks and leptons,” arXiv:1512.04921 [hep-ph].