Nondecoupling of charged scalars in Higgs decay to two photons and symmetries of the scalar potential

Gautam Bhattacharyya\(^1\), Dipankar Das\(^2\)

Saha Institute of Nuclear Physics, 1/AF Bidhan Nagar, Kolkata 700064, India

Abstract

A large class of two- and three-Higgs-doublet models with discrete symmetries has been employed in the literature to address various aspects of flavor physics. We analyse how the precision measurement of the Higgs to diphoton signal strength would severely constrain these scenarios due to the nondecoupling behavior of the charged scalars, to the extent that the number of additional scalar doublets can be constrained no matter how heavy the nonstandard scalars are. We demonstrate that if the scalar potential is endowed with appropriate global continuous symmetries together with soft breaking parameters, decoupling can be achieved thanks to the unitarity constraints on the mass-square differences of the heavy scalars.

The behavior of the scalar boson observed at the CERN Large Hadron Collider (LHC) is tantalizingly close to that of the Standard Model (SM) Higgs boson. A very timely and relevant question is whether this scalar is the only one of its type as predicted by the SM, or it is the first to have been discovered in a family of more such species arising from an underlying extended scalar sector. A natural extension of the SM is realized by adding more SU(2) scalar doublets, which we consider in this paper. There are two advantages for choosing doublets. First, the \(\rho\)-parameter remains unity at tree level. Second, it is straightforward to find a combination, namely,

\[
h = \frac{1}{v} \sum_{i=1}^{n} v_i h_i, \quad \text{with} \quad v^2 = \sum_{i=1}^{n} v_i^2 = (246 \text{ GeV})^2,
\]

\((v_i\) is the vacuum expectation value (vev) of the \(i\)-th doublet and \(h_i\) is the corresponding real scalar field), which has SM-like couplings with fermions and gauge bosons. We identify this state with the scalar observed at the LHC having a mass \(m_h \approx 125\) GeV. Recovering the SM-like Higgs boson in multi-Higgs models, i.e. to find an alignment limit, is crucial as the LHC data on the Higgs boson branching ratios are showing increasing affinity towards the SM predictions. In this paper we pay specific attention to the \(h \to \gamma\gamma\) process. Though this process is loop driven and has a small branching ratio, it played an important role in the Higgs discovery. Importantly, this branching ratio is expected to be measured in LHC-14 with much greater accuracy. Now, additional SU(2) scalar doublets would bring in additional states, both charged and neutral, in the spectrum. Here our primary concern is how those charged scalars couple to \(h\) and how much they contribute to the \(h \to \gamma\gamma\) rate as virtual states in loops. This leads to the observation that even when the masses of the charged scalars floating in the loop are taken to very large values, they do not ‘necessarily’ decouple from this process. Deciphering the underlying reasons behind this constitutes the motive of this paper. Although this has been noted in the past in the context of two-Higgs-doublet models (2HDM), only some cursory remarks were made on it without exploring its full implications [1–5]. We investigate the rôle of symmetries which are imposed on the scalar potential in figuring out under what conditions the decoupling of heavy charged scalars in the \(h \to \gamma\gamma\) loop takes place. The upshot is that if the potential has an exact \(Z_2\) symmetry, which is the case for a large class of 2HDM scenarios [6], the contribution of the charged scalar does not decouple. If \(Z_2\) is softly broken by a term in the potential then decoupling can be achieved at the expense of fine-tuning of parameters. On the other hand, a global U(1) symmetry followed by its soft breaking can ensure decoupling. For simplicity, we first demonstrate

\[^1\text{gautam.bhattacharyya@saha.ac.in}\]
\[^2\text{d.das@saha.ac.in}\]
this behavior in the context of 2HDM. We then address the same question, for the first time, in the context of three-Higgs-doublet models (3HDM). It is not difficult to foresee what happens if we add more doublets, which leads us to draw an important conclusion: unless decoupling is ensured, e.g. as we did by imposing a global U(1) symmetry in the 2HDM potential, precision measurements of \( h \to \gamma \gamma \) branching ratio can put constraints on the number of additional scalar doublets regardless of how heavy the charged scalars are. We recall that only lower bounds on charged scalar masses have been placed from processes like \( b \to s \gamma \), as the effects decouple when their masses are heavy for all such flavor observables. Thus, precision measurements of \( h \to \gamma \gamma \) would provide complementary information. Incidentally, whatever we comment on \( h \to \gamma \gamma \) applies for \( h \to Z \gamma \) as well at least on a qualitative level.

It should be noted that in multi-doublet scalar models, the production cross section as well as the tree-level decay widths of the Higgs boson remain unaltered from their respective SM expectations in the alignment limit. Only the loop induced decay modes like \( h \to \gamma \gamma \) and \( h \to Z \gamma \) will pick up additional contributions induced by virtual charged scalars. However, the branching ratios into these channels are too tiny compared to other dominant modes. As a result, the total Higgs decay width will be hardly modified. Considering all these, the expression for the diphoton signal strength is simplified to

\[
\mu_{\gamma \gamma} = \frac{\sigma(pp \to h)}{\sigma^{\text{SM}}(pp \to h)} \cdot \frac{\text{BR}(h \to \gamma \gamma)}{\text{BR}^{\text{SM}}(h \to \gamma \gamma)} = \frac{\Gamma(h \to \gamma \gamma)}{\Gamma^{\text{SM}}(h \to \gamma \gamma)}. \tag{2}
\]

For convenience, we parametrize the coupling of \( h \) to the charged scalars as

\[
g_{H^+_i H^-_i} = \kappa_i \frac{g m_{H^+_i}^2}{M_W}, \tag{3}
\]

where \( m_{H^+_i} \) is the mass of the \( i \)-th charged scalar (\( H^+_i \)). As we will see later, the decoupling or nondecoupling behavior of the \( i \)-th charged scalar from \( \mu_{\gamma \gamma} \) is encoded in \( \kappa_i \). The expression of the diphoton decay width of the Higgs is given by [7]:

\[
\Gamma(h \to \gamma \gamma) = \frac{\alpha^2 g^2}{128 \pi^3} \frac{m_h^3}{M_W^2} \left| A_W + \frac{4}{3} A_t + \sum_i \kappa_i A_{i+} \right|^2, \tag{4}
\]

where, using \( \tau_x \equiv (2 m_x / m_h)^2 \), the expressions for \( A_W, A_t \) and \( A_{i+} \) are given by

\[
A_W = 2 + 3 \tau_W + 3 \tau_W (2 - \tau_W) f(\tau_W), \quad A_t = -2 \tau_t \left[ 1 + (1 - \tau_t) f(\tau_t) \right], \quad A_{i+} = -\tau_i \left[ 1 - \tau_i + f(\tau_i) \right]. \tag{5}
\]

Since we are concerned with heavy charged scalars, we can take \( \tau_x > 1 \) for \( x = (W, t, H^+_i) \), and then \( f(\tau) = \left[ \sin^{-1} \left( \sqrt{1/\tau} \right) \right]^2 \). Now plugging Eq. (4) into Eq. (2), we obtain

\[
\mu_{\gamma \gamma} = \frac{\left| A_W + \frac{4}{3} A_t + \sum_i \kappa_i A_{i+} \right|^2}{\left| A_W + \frac{4}{3} A_t \right|^2}. \tag{6}
\]

In the limit the charged scalar is very heavy, the quantity \( A_{i+} \) saturates to \( 1/3 \). If \( \kappa_i \) also saturates to some finite value in that limit then the charged scalar would not decouple from the \( h \to \gamma \gamma \) loop. Then no matter how heavy the charged scalar is, \( \mu_{\gamma \gamma} \) will differ from its SM value. If the experimental value of \( \mu_{\gamma \gamma} \) eventually settles on very close to the SM prediction then such nondecoupling scenarios will be disfavored. The decoupling would happen only if \( \kappa_i \) falls with increasing charged scalar mass. In what follows, we will illustrate these features by considering some popular doublet extensions of the SM scalar sector.

**Two Higgs-doublet models**: We consider a 2HDM with \( \phi_1 \) and \( \phi_2 \) as the two scalar doublets. Then we impose a \( Z_2 \) symmetry in the potential, namely, \( \phi_1 \to \phi_1 \) and \( \phi_2 \to -\phi_2 \), to avoid Higgs mediated flavor-
In the left panel (a) we display the constraints on $\kappa_1$ in 2HDM coming from the measured values of $\mu_{\gamma\gamma}$ at 95% C.L. by the CMS ($1.14^{+0.28}_{-0.23}$ [9]) and ATLAS ($1.57^{+0.33}_{-0.29}$ [10]) Collaborations. In the right panel (b) we show what would be the 95% C.L. allowed range of $\kappa_1$ if $\mu_{\gamma\gamma}$ is hypothetically measured to be as $1 \pm 0.1(0.05)$ in future colliders. In both panels we have plotted Eq. (8) for two different values of $\lambda_5$. Note that the choice $\lambda_5 = 0$ signifying exact $Z_2$ symmetry is disfavored by the current ATLAS data at 95% C.L.

The expression of the scalar potential is displayed below [7]:

$$V_{2HDM} = \lambda_1 \left( \phi_1^\dagger \phi_1 - \frac{v_1^2}{2} \right)^2 + \lambda_2 \left( \phi_2^\dagger \phi_2 - \frac{v_2^2}{2} \right)^2 + \lambda_3 \left( \phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2 - \frac{v_1^2 + v_2^2}{2} \right)^2 + \lambda_4 \left( (\phi_1^\dagger \phi_2)(\phi_2^\dagger \phi_1) - (\phi_1^\dagger \phi_2)(\phi_2^\dagger \phi_1) \right) + \lambda_5 \left( \text{Re} \phi_1^\dagger \phi_2 - \frac{v_1 v_2}{2} \right)^2 + \lambda_6 \left( \text{Im} \phi_1^\dagger \phi_2 \right)^2,$$

where the $\lambda_5$ term arises due to soft breaking of $Z_2$. We assume all the lambdas to be real, i.e., CP is not broken explicitly. We can recover the SM-like Higgs boson in the alignment limit as shown in Eq. (1). Its coupling to the charged Higgs boson is controlled by (putting $i = 1$ in $\kappa_i$) [1,2,4,5,8]

$$\kappa_1 = -\frac{1}{m_{1+}^2} \left( m_{1+}^2 + m_A^2 - \frac{\lambda_5 v^2}{2} \right).$$

Clearly, $\kappa_1$ saturates to $-1$ as the charged scalar becomes excessively heavy. Decoupling can be achieved at the cost of some unnatural fine tuning between $\lambda_5$ and $m_{1+}$. On the other hand, if the $Z_2$ symmetry in the scalar potential is exact, i.e. $\lambda_5 = 0$, then the charged scalar will never decouple and will cause $\mu_{\gamma\gamma}$ to settle below its SM prediction. In Fig. 1 we have plotted the allowed range of $\kappa_1$ in 2HDM from the present LHC data as well as from an anticipation of future sensitivity.

An interesting possibility arises when we employ a U(1) symmetry, rather than the usual $Z_2$ symmetry, in the potential. The choice $\lambda_5 = \lambda_6$ will ensure U(1) symmetry in the quartic terms. The bilinear term involving $\lambda_5$ still breaks the U(1) symmetry softly. Then the mass of the pseudoscalar gets related to the soft breaking parameter $\lambda_5$ as $m_A^2 = \lambda_5 v^2/2$. In this case, the expression for $\kappa_1$ reads [4]:

$$\kappa_1 = -\frac{1}{m_{1+}^2} \left( m_{1+}^2 - m_A^2 + \frac{m_{1+}^2}{2} \right).$$

What is interesting is that the splitting between the charged scalar and the pseudoscalar mass ($|m_{1+}^2 - m_A^2|$) is restricted from the combined constraints of unitarity, stability and the electroweak T-parameter [4]. Consequently, the numerator in Eq. (9) cannot grow indefinitely with increasing $m_{1+}$. Thus $\kappa_1$ becomes very small.
in that limit and \( \mu_{\gamma\gamma} \) reaches the SM predicted value. The key issue is that the \( Z_2 \) symmetry breaking \( \lambda_5 \) term was not related to the mass of any particle in the spectrum, and hence its adjustment \textit{vis-à-vis} the charged scalar mass was nothing short of fine-tuning. Now, the global \( U(1) \) breaking \( \lambda_5 \) is related to the pseudoscalar mass whose splitting with the charged scalar mass is restricted from above.

**Three-Higgs-doublet models:** \( S_3 \) or \( A_4 \) symmetric flavor models are typical examples which employ three Higgs doublets. With \( \phi_1, \phi_2 \) and \( \phi_3 \) as the three scalar SU(2) doublets, the scalar potential for the \( S_3 \) symmetric case can be written as (see e.g. [11, 12]).

\[
V_{V3HDM}^{S_3} = -\mu_1^2(\phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2 + \phi_3^\dagger \phi_3) + \lambda_1(\phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2 - \phi_3^\dagger \phi_3)^2 + \lambda_2(\phi_1^\dagger \phi_2 - \phi_2^\dagger \phi_1)^2 + \lambda_3(\phi_1^\dagger \phi_3 - \phi_3^\dagger \phi_1)^2 + \lambda_4(\phi_1^\dagger \phi_2)(\phi_2^\dagger \phi_1) + (\phi_1^\dagger \phi_3)(\phi_3^\dagger \phi_1) + \lambda_5(\phi_1^\dagger \phi_3)(\phi_3^\dagger \phi_1) + \lambda_6(\phi_1^\dagger \phi_2)(\phi_2^\dagger \phi_3) + \lambda_7(\phi_1^\dagger \phi_3)(\phi_3^\dagger \phi_2) + (\phi_1^\dagger \phi_2)(\phi_2^\dagger \phi_3) + \lambda_8(\phi_1^\dagger \phi_3)^2.
\]

Assuming the lambdas to be real, potential minimization conditions attribute a relation between two of the three vevs \((v_1 = \sqrt{3}v_2)\). Using this relation, an alignment limit can be obtained for this model also [12].

Now we write the potential satisfying \( A_4 \) symmetry (see e.g. [13]).

\[
V_{V3HDM}^{A_4} = -\mu^2(\phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2 + \phi_3^\dagger \phi_3) + \lambda_1(\phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2 + \phi_3^\dagger \phi_3)^2 + \lambda_2(\phi_1^\dagger \phi_1 \phi_2^\dagger \phi_2 + \phi_2^\dagger \phi_2 \phi_3^\dagger \phi_3 + \phi_3^\dagger \phi_3 \phi_1^\dagger \phi_1) + \lambda_3(\phi_1^\dagger \phi_2 \phi_3^\dagger \phi_1 + \phi_2^\dagger \phi_3 \phi_1^\dagger \phi_2 + \phi_3^\dagger \phi_1 \phi_1^\dagger \phi_3)
\]

\[
+ \lambda_4 \left[ e^{i\epsilon} \left\{ (\phi_1^\dagger \phi_2)^2 + (\phi_2^\dagger \phi_3)^2 + (\phi_3^\dagger \phi_1)^2 \right\} + \text{h.c.} \right].
\]

In one plausible scenario, the minimization conditions require that all the three vevs are equal [14]. This particular choice automatically yields a SM-like Higgs as well as two pairs of complex neutral states with mixed CP properties. Note that for \( \epsilon = 0 \) in Eq. (11), the symmetry of the potential is enhanced to \( S_4 \). However, our conclusions do not depend on the value of \( \epsilon \).

Thus, a 3HDM can provide an SM-like Higgs along with two pairs of charged scalars, as exemplified with \( S_3 \) and \( A_4 \) scenarios. After expressing the lambdas in terms of the physical masses, we obtain the following expressions for \( \kappa_i \) \((i = 1, 2)\) in the alignment limit, which are the same for both \( S_3 \) and \( A_4 \):

\[
\kappa_i = -\left(1 + \frac{m_{h_i}^2}{2m_{i^+}^2}\right) \quad \text{for } i = 1, 2.
\]

Clearly, the charged scalars do not decouple from the diphoton decay width, since \( \kappa_i \) settles to \(-1\) when \( m_{i^+} \) is very large compared to \( m_h \). Note, both the charged scalars contribute in the same direction to reduce \( \mu_{\gamma\gamma} \).

Now we turn our attention to the case of a global continuous symmetry in 3HDM potential. For illustration, we consider that the symmetry is SO(2) under which \( \phi_1 \) and \( \phi_2 \) form a doublet. The expression for the scalar potential is similar to Eq. (10), only that now \( \lambda_4 = 0 \) and the potential contains an additional bilinear term \((-\mu_{12}^2 \phi_1^\dagger \phi_2 + \text{h.c.})\). The real part of \( \mu_{12}^2 \) softly breaks the SO(2) symmetry and prevents the occurrence of any massless scalar in the theory. In any case, we assume \( \mu_{12}^2 \) to be real just like any other parameters in the potential. The relevant minimization conditions are given by

\[
\begin{align*}
    v_1\mu_1^2 + v_2\mu_1^4 &= v_1(v_1^2 + v_2^2)(\lambda_1 + \lambda_3) + \frac{1}{2} v_1 v_3^2(\lambda_5 + \lambda_6 + 2\lambda_7), \\
    v_2\mu_1^2 + v_1\mu_1^4 &= v_2(v_1^2 + v_2^2)(\lambda_1 + \lambda_3) + \frac{1}{2} v_2 v_3^2(\lambda_5 + \lambda_6 + 2\lambda_7).
\end{align*}
\]
Note that nonzero $\mu_{12}^2$ requires $v_1 = v_2$. An interchange symmetry (1 ↔ 2) is accidentally preserved even after spontaneous symmetry breaking. We will have three CP even scalars ($h', H, h$), two pseudoscalars ($A_1, A_2$) and two pairs of charged scalars ($H_1^\pm, H_2^\pm$). Among these, $h', A_1$ and $H_1^\pm$ are odd under the interchange symmetry and the rest are even under it. Being odd under this interchange symmetry, $h'$ does not couple to gauge bosons as $h'VV$ ($V = W, Z$). Appearance of such an exotic scalar was noted earlier in the context of an $S_3$ symmetric 3HDM [12,15,16]. The soft breaking parameter ($\mu_{12}^2$) gets related to the mass of $h'$ as

$$m_{h'}^2 = 2\mu_{12}^2.$$  

It is straightforward to express the lambdas in terms of the physical masses. We then obtain

$$\kappa_1 = -\frac{1}{m_{1+}^2} \left( m_{1+}^2 - m_{h'}^2 + \frac{m_{h'}^2}{2} \right),$$  

$$\kappa_2 = - \left( 1 + \frac{m_{h'}^2}{2m_{2+}^2} \right).$$

The similarity between Eq. (15a) and Eq. (9) is striking. Note that $|m_{1+}^2 - m_{h'}^2|$ is constrained from unitarity. Therefore, when the first charged Higgs mass $m_{1+}$ is very large, $\kappa_1$ becomes vanishingly small. However, this decoupling does not occur in $\kappa_2$ which contains the second charged Higgs mass $m_{q+}$. It is not difficult to intuitively argue that with an extended global symmetry SO(2)$\times$U(1), together with an extra soft breaking parameter which is related to $m_{A2}$, decoupling in $\kappa_2$ can be ensured. Starting from the softly broken SO(2) symmetric potential, this additional U(1) extension ($\phi_4 \rightarrow e^{i\alpha} \phi_3$) and its soft breaking can be realized by putting $\lambda_7 = 0$ in Eq. (10) and introducing a term that softly breaks this U(1). A crucial observation we make in this paper is that the masses $m_A$ in the 2HDM context and $m_{h'}$ in the 3HDM context enter into the expressions of $\kappa_i$ - e.g. see Eqs. (9) and (15a) - only when they are related to soft global symmetry breaking parameters.

| Model                     | Expression for $\kappa_i$                        | prediction $\mu_{2\gamma}$ | prediction $\mu_{Z\gamma}$ | Decoupling?
|----------------------------|-----------------------------------------------|-----------------------------|----------------------------|---------
| 2HDM                      | Softly broken $Z_2$                          | $- \left( 1 + \frac{m_{i+}^2}{2m_{1+}^2} - \frac{\lambda_{A2}^2}{2m_{1+}^2} \right)$ | Depends on $\lambda_5$     | Depends on $\lambda_5$ | No      
|                           | Exact $Z_2$                                 | $- \left( 1 + \frac{m_{1+}^2}{2m_{i+}^2} \right)$ | $\leq 0.9$                 | $\leq 0.96$ | No      
|                           | Softly broken U(1)                          | $- \left( 1 + \frac{m_{1+}^2}{2m_{i+}^2} - \frac{m_{A}^2}{m_{1+}^2} \right)$ | Depends on $m_A$           | Depends on $m_A$ | Yes     
| 3HDM                      | Exact $S_3$                                 | $- \left( 1 + \frac{m_{i+}^2}{2m_{1+}^2} \right)$ for $i = 1, 2$ | $\leq 0.8$                 | $\leq 0.93$ | No      
|                           | Exact $A_4$                                 | $- \left( 1 + \frac{m_{i+}^2}{2m_{1+}^2} \right)$ for $i = 1, 2$ | $\leq 0.8$                 | $\leq 0.93$ | No      
|                           | Softly broken SO(2)                         | $\kappa_1 = - \left( 1 + \frac{m_{i+}^2}{2m_{1+}^2} - \frac{m_{h'}^2}{m_{1+}} \right)$ | Depends on $m_{h'}$        | Depends on $m_{h'}$ | Partial |

Table 1: Behavior of 2HDM and 3HDM scenarios in the alignment limit. In the case of exact discrete symmetries, every charged scalar pair reduces $\mu_{\gamma\gamma}$ approximately by 0.1. Although explicit expression for $\mu_{Z\gamma}$ is not shown in text, its predictions in different scenarios are displayed. In the last column where we refer to ‘nondecoupling’ by ‘No’ and ‘decoupling’ by ‘Yes’, we implicitly assume that there is no fine-tuning, while in the last row ‘Partial’ implies that only the first charged scalar decouples.

Conclusions: To our knowledge, this is the first attempt towards establishing a connection between decoupling or nondecoupling of charged scalars from the diphoton decay of the Higgs with the symmetries of the scalar potential. The main conclusion of our analysis is that charged scalars in multi-doublet scalar extensions of the SM do not necessarily decouple from $\mu_{\gamma\gamma}$. Therefore, a precisely measured $\mu_{\gamma\gamma}$ can smell the presence of such nonstandard scalars even if they are super-heavy. In fact, $\mu_{\gamma\gamma}$ can constrain the number of such doublets. Table 1 shows that each additional pair of charged scalars ($H_i^\pm$) reduces $\mu_{\gamma\gamma}$ approximately by 0.1 when the potential has an exact discrete symmetry. Our illustrations are based on two- and three-Higgs-doublet models.
which are motivated by flavor symmetries. We have explicitly demonstrated how soft breaking of a global U(1) symmetry can ensure decoupling in 2HDM in the alignment limit. In the case of 3HDM, with a softly broken global SO(2) symmetry in the potential, decoupling can be ensured for one pair of charged scalars ($H^+_1$), while the second pair ($H^+_2$) still do not decouple. Employing the soft breaking terms of an extended global continuous symmetry, namely, SO(2) × U(1), the nondecoupling effects of $H^+_2$ can be tamed. If we have more pairs of charged scalars in the theory stemming from additional scalar doublets, even more enhanced or extended global continuous symmetries — only softly broken — would be required to ensure decoupling of all charged scalars from $\mu_{\gamma\gamma}$.

If future measurement of $\mu_{\gamma\gamma}$ is found to be consistent with the SM prediction to a high degree of precision — say better than 10% — we can start disfavoring multi-Higgs-doublet models with exact discrete symmetries, a large class of which is motivated from flavor physics. It is realized for the first time in this paper that the soft breaking terms in the potential, which are often used in the literature, can play an important role in ensuring decoupling, albeit with fine-tuning. To avoid fine-tuning, one must start with a global continuous symmetry in the potential followed by its soft breaking. We believe that the observations made in this paper are going to influence model building in flavor physics. In future, it would interesting to explore the consequences of global symmetries in the potential for nondoublet scalar extensions in the present context.

References

[1] A. Djouadi, V. Driesen, W. Hollik, and A. Kraft, *The Higgs photon - Z boson coupling revisited*, Eur.Phys.J. C1 (1998) 163–175, [hep-ph/9701342].

[2] A. Arhrib, M. Capdequi Peyranere, W. Hollik, and S. Penaranda, *Higgs decays in the two Higgs doublet model: Large quantum effects in the decoupling regime*, Phys.Lett. B579 (2004) 361–370, [hep-ph/0307391].

[3] W.-F. Chang, J. N. Ng, and J. M. Wu, *Constraints on New Scalars from the LHC 125 GeV Higgs Signal*, Phys.Rev. D86 (2012) 033003, [arXiv:1206.5047].

[4] G. Bhattacharyya, D. Das, P. B. Pal, and M. Rebelo, *Scalar sector properties of two-Higgs-doublet models with a global U(1) symmetry*, JHEP 1310 (2013) 081, [arXiv:1308.4297].

[5] P. 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. D89 (2014) 115003, [arXiv:1403.4736].

[6] G. Branco, P. Ferreira, L. Lavoura, M. Rebelo, M. Sher, et al., *Theory and phenomenology of two-Higgs-doublet models*, Phys.Rept. 516 (2012) 1–102, [arXiv:1106.0034].

[7] J. F. Gunion, H. E. Haber, G. L. Kane, and S. Dawson, *The Higgs Hunter’s Guide*, Front.Phys. 80 (2000) 1–448.

[8] B. Swiezewska and M. Krawczyk, *Diphoton rate in the inert doublet model with a 125 GeV Higgs boson*, Phys.Rev. D88 (2013) 035019, [arXiv:1212.4100].

[9] CMS Collaboration, V. Khachatryan et al., *Observation of the diphoton decay of the Higgs boson and measurement of its properties*, arXiv:1407.0558.

[10] ATLAS Collaboration, *ATLAS-CONF-2014-009*, https://twiki.cern.ch/twiki/bin/view/AtlasPublic/CONFnotes.

[11] E. Barradas-Guevara, O. Flix-Beltrn, and E. R. Juregui, *S(3) flavoured Higgs model trilinear self-couplings*, arXiv:1402.2244.

[12] D. Das and U. K. Dey, *Analysis of an extended scalar sector with S3 symmetry*, Phys.Rev. D89 (2014) 095025, [arXiv:1404.2491].
[13] E. Ma and G. Rajasekaran, *Softly broken $A(4)$ symmetry for nearly degenerate neutrino masses*, Phys.Rev. D64 (2001) 113012, [hep-ph/0106291].

[14] R. de Adelhart Toorop, F. Bazzocchi, L. Merlo, and A. Paris, *Constraining Flavour Symmetries At The EW Scale I: The $A_4$ Higgs Potential*, JHEP 1103 (2011) 035, [arXiv:1012.1791].

[15] G. Bhattacharyya, P. Leser, and H. Pas, *Exotic Higgs boson decay modes as a harbinger of $S_3$ flavor symmetry*, Phys.Rev. D83 (2011) 011701, [arXiv:1006.5597].

[16] G. Bhattacharyya, P. Leser, and H. Pas, *Novel signatures of the Higgs sector from $S_3$ flavor symmetry*, Phys.Rev. D86 (2012) 036009, [arXiv:1206.4202].