

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

Gautam Bhattacharyya<sup>1</sup>, Dipankar Das<sup>2</sup>

*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 in the presence of exact discrete symmetry, the number of additional non-inert 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.

## 1 Introduction

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,

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

( $v_i$  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. This is not in general a mass eigenstate. But when we demand that this is indeed the physical state observed at the LHC with a mass  $m_h \approx 125$  GeV, we are automatically led to the so called *alignment limit*. This limit is motivated by the LHC data on the Higgs boson signal strengths in different channels which are showing increasing affinity towards the SM predictions. In this paper we pay specific attention to the  $h \rightarrow \gamma\gamma$  process. Though this process is loop driven and has a small branching ratio, it played an important rôle 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  $h \rightarrow \gamma\gamma$  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

---

<sup>1</sup>gautam.bhattacharyya@saha.ac.in

<sup>2</sup>d.das@saha.ac.in

implications [1–6]. 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  $h \rightarrow \gamma\gamma$  loop takes place. The upshot is that if the potential has an exact  $Z_2$  symmetry *and* both the scalars receive vevs, which is the case for a large class of 2HDM scenarios [7], 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 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 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  $h \rightarrow \gamma\gamma$  branching ratio can put constraints on the number of additional non-inert 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  $b \rightarrow s\gamma$ , as the effects decouple when their masses are heavy for all such flavor observables. Thus, precision measurements of  $h \rightarrow \gamma\gamma$  would provide complementary information. Incidentally, whatever we comment on  $h \rightarrow \gamma\gamma$  applies for  $h \rightarrow Z\gamma$  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  $h \rightarrow \gamma\gamma$  and  $h \rightarrow Z\gamma$  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

$$\mu_{\gamma\gamma} \equiv \frac{\sigma(pp \rightarrow h)}{\sigma^{\text{SM}}(pp \rightarrow h)} \cdot \frac{\text{BR}(h \rightarrow \gamma\gamma)}{\text{BR}^{\text{SM}}(h \rightarrow \gamma\gamma)} = \frac{\Gamma(h \rightarrow \gamma\gamma)}{\Gamma^{\text{SM}}(h \rightarrow \gamma\gamma)}. \quad (2)$$

For convenience, we parametrize the coupling of  $h$  to the charged scalars as

$$g_{hH_i^+H_i^-} = \kappa_i \frac{gm_{i+}^2}{M_W}, \quad (3)$$

where  $m_{i+}$  is the mass of the  $i$ -th charged scalar ( $H_i^\pm$ ). 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 [8] :

$$\Gamma(h \rightarrow \gamma\gamma) = \frac{\alpha^2 g^2}{2^{10} \pi^3} \frac{m_h^3}{M_W^2} \left| \mathcal{A}_W + \frac{4}{3} \mathcal{A}_t + \sum_i \kappa_i \mathcal{A}_{i+} \right|^2, \quad (4)$$

where, using  $\tau_x \equiv (2m_x/m_h)^2$ , the expressions for  $\mathcal{A}_W$ ,  $\mathcal{A}_t$  and  $\mathcal{A}_{i+}$  are given by

$$\mathcal{A}_W = 2 + 3\tau_W + 3\tau_W(2 - \tau_W)f(\tau_W), \quad \mathcal{A}_t = -2\tau_t[1 + (1 - \tau_t)f(\tau_t)], \quad \mathcal{A}_{i+} = -\tau_{i+}[1 - \tau_{i+}f(\tau_{i+})]. \quad (5)$$

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

$$\mu_{\gamma\gamma} = \frac{\left| \mathcal{A}_W + \frac{4}{3} \mathcal{A}_t + \sum_i \kappa_i \mathcal{A}_{i+} \right|^2}{\left| \mathcal{A}_W + \frac{4}{3} \mathcal{A}_t \right|^2}. \quad (6)$$

In the limit the charged scalar is very heavy, the quantity  $\mathcal{A}_{i+}$  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  $h \rightarrow \gamma\gamma$  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.

## 2 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,  $\phi_1 \rightarrow \phi_1$  and  $\phi_2 \rightarrow -\phi_2$ , to avoid Higgs mediated flavor-changing neutral current in the fermionic sector. The expression of the scalar potential is displayed below [8] :

$$V_{2\text{HDM}} = \lambda_1 \left( \phi_1^\dagger \phi_1 - \frac{v_1^2}{2} \right)^2 + \lambda_2 \left( \phi_2^\dagger \phi_2 - \frac{v_2^2}{2} \right)^2 + \lambda_3 \left( \phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2 - \frac{v_1^2 + v_2^2}{2} \right)^2 \\ + \lambda_4 \left( (\phi_1^\dagger \phi_1)(\phi_2^\dagger \phi_2) - (\phi_1^\dagger \phi_2)(\phi_2^\dagger \phi_1) \right) + \lambda_5 \left( \text{Re} \phi_1^\dagger \phi_2 - \frac{v_1 v_2}{2} \right)^2 + \lambda_6 \left( \text{Im} \phi_1^\dagger \phi_2 \right)^2, \quad (7)$$

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. It is also implicitly assumed that both the scalar doublets receive vevs.

First, it is important to count the number of free parameters. As we have assumed the parameters to be real, there are only 8 free parameters. Two of them  $v_1$  and  $v_2$  can be traded for  $v$  and  $\tan \beta \equiv v_2/v_1$ . All the remaining parameters, except  $\lambda_5$ , can be traded for four physical scalar masses ( $m_h, m_H, m_A, m_{1+} (\equiv m_{H^+})$ ) and the rotation angle ( $\alpha$ ) in the neutral CP even sector. The lambdas can be expressed in terms of physical masses as:

$$\lambda_1 = \frac{1}{2v^2 \cos^2 \beta} \left[ m_H^2 \cos^2 \alpha + m_h^2 \sin^2 \alpha - \frac{\sin \alpha \cos \alpha}{\tan \beta} (m_H^2 - m_h^2) \right] - \frac{\lambda_5}{4} (\tan^2 \beta - 1), \quad (8a)$$

$$\lambda_2 = \frac{1}{2v^2 \sin^2 \beta} \left[ m_h^2 \cos^2 \alpha + m_H^2 \sin^2 \alpha - \sin \alpha \cos \alpha \tan \beta (m_H^2 - m_h^2) \right] - \frac{\lambda_5}{4} (\cot^2 \beta - 1), \quad (8b)$$

$$\lambda_3 = \frac{1}{2v^2} \frac{\sin \alpha \cos \alpha}{\sin \beta \cos \beta} (m_H^2 - m_h^2) - \frac{\lambda_5}{4}, \quad (8c)$$

$$\lambda_4 = \frac{2}{v^2} m_{1+}^2, \quad (8d)$$

$$\lambda_6 = \frac{2}{v^2} m_A^2. \quad (8e)$$

The quantities that appear on the right-hand-side of Eq. (8) are all *independent* parameters. Since we work in the alignment limit, it follows that  $\alpha (= \beta - \pi/2)$ . This means that we deal with 7 independent parameters, out of which 5 are unknown, namely,  $m_H, m_A, m_{1+}, \tan \beta$  and  $\lambda_5$ , while 2 are known, namely,  $v = 246$  GeV and  $m_h \approx 125$  GeV. As mentioned earlier, the SM-like Higgs boson is recovered in this limit as shown in Eq. (1).

The charged scalar contribution to  $\mu_{\gamma\gamma}$  is controlled by (putting  $i = 1$  in  $\kappa_i$ ) [1, 2, 4, 5, 9]

$$\kappa_1 = -\frac{1}{m_{1+}^2} \left( m_{1+}^2 + \frac{m_h^2}{2} - \frac{\lambda_5 v^2}{2} \right). \quad (9)$$

Clearly,  $\kappa_1$  saturates to  $-1$  as the charged scalar becomes excessively heavy. Decoupling can be achieved by tuning  $m_{1+}^2 \simeq \lambda_5 v^2/2$  [10]. Recalling our counting of independent parameters, any adjustment between the charged scalar mass and  $\lambda_5$  is nothing short of fine-tuning. 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] :

$$\kappa_1 = -\frac{1}{m_{1+}^2} \left( m_{1+}^2 - m_A^2 + \frac{m_h^2}{2} \right). \quad (10)$$

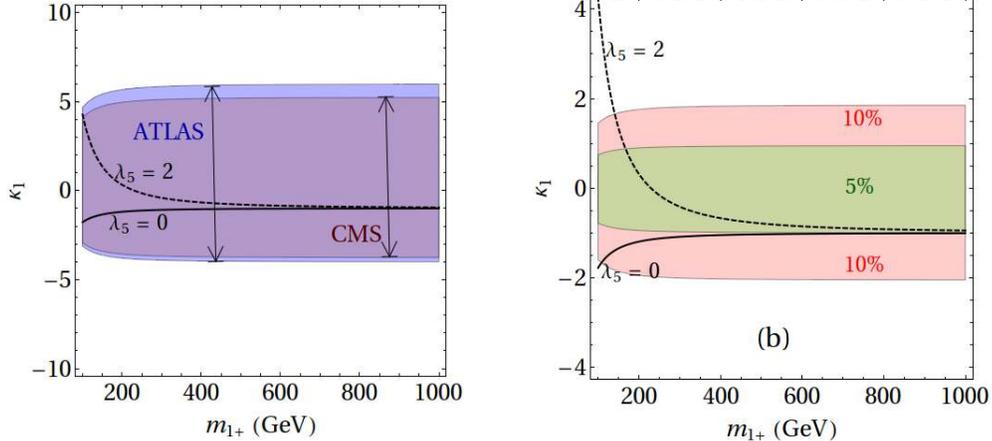


Figure 1: In the left panel (a) we display the constraints on  $\kappa_1$  in 2HDM coming from the measured values of  $\mu_{\gamma\gamma}$  at 95% C.L. by the CMS ( $1.14^{+0.26}_{-0.23}$  [11]) and ATLAS ( $1.17 \pm 0.27$  [12]) Collaborations. In the right panel (b) we show what would be the 95% C.L. 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. (9) for two different values of  $\lambda_5$ .

In a previous paper [4], we provided a detailed analysis on the unitarity and stability constraints on various combination of  $\lambda_i$  couplings when the 2HDM scalar potential has a softly broken U(1) symmetry. We cite some of them here to demonstrate ‘decoupling’ for large individual quartic couplings, as what is constrained from unitarity is only their differences in certain combinations. For example,  $(2\lambda_3 + \lambda_4) \leq 16\pi$  implies  $(2m_{1+}^2 - m_H^2 - m_A^2 + m_h^2) \leq 16\pi v^2$ . Also,  $(\lambda_1 + \lambda_2 + 2\lambda_3) \leq 16\pi/3$  implies  $(m_H^2 - m_A^2)(\tan^2 \beta + \cot^2 \beta) + 2m_h^2 \leq 32\pi v^2/3$ . These relations, together with  $|m_{1+} - m_H| \ll (m_{1+}, m_H)$  arising from the electroweak  $T$  parameter, restrict the splitting between the charged scalar and the pseudoscalar mass ( $|m_{1+}^2 - m_A^2|$ ). As displayed through more such relations among quartic couplings and the associated plots in the plane of non-SM scalar masses in [4], the individual scalar masses can become very large without violating unitarity as long as their mass-square differences are within certain limits. Consequently, the numerator in Eq. (10) 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 *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 unitarity.

## 2.1 Underlying dynamics behind decoupling

We now discuss the underlying reason behind decoupling or nondecoupling of nonstandard scalars from physical processes in the 2HDM context. The conclusion is equally applicable for  $n$ HDM where  $n > 2$ . First, we write down the 2HDM potential using a different notation:

$$\begin{aligned}
V_{2\text{HDM}} = & m_{11}^2 \phi_1^\dagger \phi_1 + m_{22}^2 \phi_2^\dagger \phi_2 - \left( m_{12}^2 \phi_1^\dagger \phi_2 + \text{h.c.} \right) + \frac{\beta_1}{2} \left( \phi_1^\dagger \phi_1 \right)^2 + \frac{\beta_2}{2} \left( \phi_2^\dagger \phi_2 \right)^2 \\
& + \beta_3 \left( \phi_1^\dagger \phi_1 \right) \left( \phi_2^\dagger \phi_2 \right) + \beta_4 \left( \phi_1^\dagger \phi_2 \right) \left( \phi_2^\dagger \phi_1 \right) + \left\{ \frac{\beta_5}{2} \left( \phi_1^\dagger \phi_2 \right)^2 + \text{h.c.} \right\}. \quad (11)
\end{aligned}$$

This parametrization does not *a priori* assume, unlike the one used in this paper given in Eq. (7), that  $\phi_1$  or  $\phi_2$  necessarily acquires any vev. In this parametrization, in the limit when the dimensionless couplings  $\beta_2 = \beta_3 = \beta_4 = \beta_5 = 0$ , the mass mixing parameter  $m_{12}^2 = 0$ , and  $m_{22}^2 > 0$ , the second Higgs doublet  $\phi_2$  does

not acquire any vev, and the SM scalar potential is recovered with the relation  $v^2 = v_1^2 = -m_{11}^2/\beta_1$ . This is the *inert doublet* scenario with a perfectly  $Z_2$  symmetric potential, in which all the nonstandard scalars decouple from physical processes when the parameter  $m_{22}^2$  controlling their masses is taken to infinitely large value. Note that  $m_{22}^2$ , in this case, does not have its origin in spontaneous symmetry breaking (SSB), and this is why its large value could ensure decoupling. On the contrary, if we try to establish the equivalence between the two parametrizations, given in Eq. (11) and Eq. (7), we have to go to a situation when both the scalars receive vevs, and only then we obtain the following relations:

$$\begin{aligned} m_{11}^2 &= -(\lambda_1 v_1^2 + \lambda_3 v^2) ; m_{22}^2 = -(\lambda_2 v_2^2 + \lambda_3 v^2) ; m_{12}^2 = \frac{\lambda_5}{2} v_1 v_2 ; \beta_1 = 2(\lambda_1 + \lambda_3) ; \\ \beta_2 &= 2(\lambda_2 + \lambda_3) ; \beta_3 = 2\lambda_3 + \lambda_4 ; \beta_4 = \frac{\lambda_5 + \lambda_6}{2} - \lambda_4 ; \beta_5 = \frac{\lambda_5 - \lambda_6}{2} . \end{aligned} \quad (12)$$

Note that when both the doublets receive vevs, one can trade the two parameters  $m_{11}^2$  and  $m_{22}^2$  in favor of  $v_1$  and  $v_2$ . Then the magnitude of the third parameter  $m_{12}^2$ , or equivalently  $\lambda_5$ , has nothing to do with SSB, and this parameter provides the regulator whose large value ensures decoupling of all nonstandard scalars from physical processes. However, while employing  $m_{12}^2$  (or equivalently  $\lambda_5$  in our parametrization) for decoupling, one cannot escape from some tuning of parameters for softly broken  $Z_2$  as explained around Eq. (9), but no such tuning is required for softly broken U(1) (discussed before). Nondecoupling would result when the symmetry of the potential is exact ( $m_{12}^2 = 0$ ), and at the same time, both the scalars receive vevs (which implies  $\lambda_5 = 0$ ). In this case all the non-SM physical scalar masses would be proportional to the electroweak vev, and there is no independent mass-dimensional parameter which has non-SSB origin. As illustrated in the inert doublet case, even with exactly symmetric potential, decoupling is achieved in 2HDM.

Admittedly, the parametrization of Eq. (7) is less general than that of Eq. (11). Any connection between the two sets of parameters can be established only when both the scalars receive vevs. The inert doublet scenario can be very easily realized in the parametrization of Eq. (11), while just setting  $v_2 = 0$  in the parametrization of Eq. (7) does not lead us to the same limit. To appreciate this salient aspect, we consider a simpler scenario when we have only one Higgs doublet. Then the potential can be written in two equivalent ways:  $V \sim \mu^2 |\phi|^2 + \lambda |\phi|^4$ , and  $V' \sim \lambda (|\phi|^2 - v^2/2)^2$ . They become truly equivalent when  $\mu^2 < 0$ , and consequently, the scalar receives a vev. But when  $\mu^2 > 0$ , the scalar remains inert. In that case, putting  $v = 0$  in  $V'$  does not take us to the physical situation given by  $V$ , as the latter still contains, in addition to  $\lambda$ , an independent dimensionful parameter  $\mu^2$ . Our Eqs. (7) and (11) are 2HDM generalizations of  $V'$  and  $V$ , respectively.

We note that Eq. (7), where it is implicitly assumed that both scalars receive vevs, covers a large class of 2HDM models (Types I-IV), where a nonvanishing (and often large)  $\tan\beta$  has played an important rôle in addressing phenomenological issues associated with processes like  $b \rightarrow s\gamma$ ,  $B_s \rightarrow \ell^+\ell^-$ ,  $B \rightarrow D^{(*)}\ell\nu$ , etc [7].

To provide further intuition into the argument of decoupling and its close connection to the existence of some non-SSB origin parameter, we draw the following analogy. It is well known that the top quark in the SM does not decouple from  $h \rightarrow \gamma\gamma$ . This is because the top quark receives all its mass from SSB and increasing the its mass will invariably imply enhancing the Yukawa coupling ( $h_t$ ). Now, suppose that the top quark receives part of its mass ( $M$ ) from some non-SSB origin, i.e.  $m_t = h_t v + M$ . Then the top-loop contribution will yield a prefactor  $h_t v / (h_t v + M)$ . In this case, by taking  $M \rightarrow \infty$ , the top quark contribution can be made to decouple from the diphoton decay width of the Higgs boson.

### 3 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. [13, 14], and also references therein for flavor physics discussions both when the  $S_3$  symmetry is exact as

well as when it is softly broken),

$$\begin{aligned}
V_{3\text{HDM}}^{S_3} = & -\mu_1^2(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2) - \mu_3^2\phi_3^\dagger\phi_3 \\
& + \lambda_1(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2)^2 + \lambda_2(\phi_1^\dagger\phi_2 - \phi_2^\dagger\phi_1)^2 + \lambda_3\left\{(\phi_1^\dagger\phi_2 + \phi_2^\dagger\phi_1)^2 + (\phi_1^\dagger\phi_1 - \phi_2^\dagger\phi_2)^2\right\} \\
& + \lambda_4\left\{(\phi_3^\dagger\phi_1)(\phi_1^\dagger\phi_2 + \phi_2^\dagger\phi_1) + (\phi_3^\dagger\phi_2)(\phi_1^\dagger\phi_1 - \phi_2^\dagger\phi_2) + \text{h.c.}\right\} \\
& + \lambda_5(\phi_3^\dagger\phi_3)(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2) + \lambda_6\left\{(\phi_3^\dagger\phi_1)(\phi_1^\dagger\phi_3) + (\phi_3^\dagger\phi_2)(\phi_2^\dagger\phi_3)\right\} \\
& + \lambda_7\left\{(\phi_3^\dagger\phi_1)(\phi_3^\dagger\phi_1) + (\phi_3^\dagger\phi_2)(\phi_3^\dagger\phi_2) + \text{h.c.}\right\} + \lambda_8(\phi_3^\dagger\phi_3)^2.
\end{aligned} \tag{13}$$

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 [14].

Now we write the potential satisfying  $A_4$  symmetry (see e.g. [15]),

$$\begin{aligned}
V_{3\text{HDM}}^{A_4} = & -\mu^2\left(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2 + \phi_3^\dagger\phi_3\right) + \lambda_1\left(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2 + \phi_3^\dagger\phi_3\right)^2 \\
& + \lambda_2\left(\phi_1^\dagger\phi_1\phi_2^\dagger\phi_2 + \phi_2^\dagger\phi_2\phi_3^\dagger\phi_3 + \phi_3^\dagger\phi_3\phi_1^\dagger\phi_1\right) + \lambda_3\left(\phi_1^\dagger\phi_2\phi_2^\dagger\phi_1 + \phi_2^\dagger\phi_3\phi_3^\dagger\phi_2 + \phi_3^\dagger\phi_1\phi_1^\dagger\phi_3\right) \\
& + \lambda_4\left[e^{i\epsilon}\left\{\left(\phi_1^\dagger\phi_2\right)^2 + \left(\phi_2^\dagger\phi_3\right)^2 + \left(\phi_3^\dagger\phi_1\right)^2\right\} + \text{h.c.}\right].
\end{aligned} \tag{14}$$

In one plausible scenario, the minimization conditions require that all the three vevs are equal [16]. 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. (14), 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  $A_4$ :

$$\kappa_i = -\left(1 + \frac{m_h^2}{2m_{i+}^2}\right) \quad \text{for } i = 1, 2. \tag{15}$$

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  $\text{SO}(2)$  under which  $\phi_1$  and  $\phi_2$  form a doublet. The expression for the scalar potential is similar to Eq. (13), 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  $\text{SO}(2)$  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

$$v_1\mu_1^2 + v_2\mu_{12}^2 = v_1(v_1^2 + v_2^2)(\lambda_1 + \lambda_3) + \frac{1}{2}v_1v_3^2(\lambda_5 + \lambda_6 + 2\lambda_7), \tag{16a}$$

$$v_2\mu_1^2 + v_1\mu_{12}^2 = v_2(v_1^2 + v_2^2)(\lambda_1 + \lambda_3) + \frac{1}{2}v_2v_3^2(\lambda_5 + \lambda_6 + 2\lambda_7). \tag{16b}$$

Note that nonzero  $\mu_{12}^2$  requires  $v_1 = v_2$ . An interchange symmetry ( $1 \leftrightarrow 2$ ) 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 [14, 17, 18]. The soft breaking parameter ( $\mu_{12}^2$ ) gets related to the mass of  $h'$  as

$$m_{h'}^2 = 2\mu_{12}^2. \tag{17}$$

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

$$\kappa_1 = -\frac{1}{m_{1+}^2} \left( m_{1+}^2 - m_{h'}^2 + \frac{m_h^2}{2} \right), \quad (18a)$$

$$\kappa_2 = -\left( 1 + \frac{m_h^2}{2m_{2+}^2} \right). \quad (18b)$$

The similarity between Eq. (18a) and Eq. (10) 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  $m_{2+}$ . It is not difficult to intuitively argue that with an extended global symmetry  $\text{SO}(2) \times \text{U}(1)$ , 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  $\text{SO}(2)$  symmetric potential, this additional  $\text{U}(1)$  extension ( $\phi_3 \rightarrow e^{i\alpha} \phi_3$ ) and its soft breaking can be realized by putting  $\lambda_7 = 0$  in Eq. (13) and introducing a term that softly breaks this  $\text{U}(1)$ . 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. (10) and (18a) – only when they are related to *soft* global symmetry breaking parameters.

Model		Expression for $\kappa_i$	prediction $\mu_{\gamma\gamma}$	prediction $\mu_{Z\gamma}$	Decoupling?
2HDM	Softly broken $Z_2$	$-\left(1 + \frac{m_h^2}{2m_{1+}^2} - \frac{\lambda_5 v^2}{2m_{1+}^2}\right)$	Depends on $\lambda_5$	Depends on $\lambda_5$	Possible
	Exact $Z_2$	$-\left(1 + \frac{m_h^2}{2m_{1+}^2}\right)$	$\leq 0.9$	$\leq 0.96$	No
	Softly broken $\text{U}(1)$	$-\left(1 + \frac{m_h^2}{2m_{1+}^2} - \frac{m_A^2}{m_{1+}^2}\right)$	Depends on $m_A$	Depends on $m_A$	Yes
3HDM	Exact $S_3$	$-\left(1 + \frac{m_h^2}{2m_{i+}^2}\right)$ for $i = 1, 2$	$\leq 0.8$	$\leq 0.93$	No
	Exact $A_4$	$-\left(1 + \frac{m_h^2}{2m_{i+}^2}\right)$ for $i = 1, 2$	$\leq 0.8$	$\leq 0.93$	No
	Softly broken $\text{SO}(2)$	$\kappa_1 = -\left(1 + \frac{m_h^2}{2m_{1+}^2} - \frac{m_{h'}^2}{m_{1+}^2}\right)$ $\kappa_2 = -\left(1 + \frac{m_h^2}{2m_{2+}^2}\right)$	Depends on $m_{h'}$	Depends on $m_{h'}$	Partial

Table 1: Behavior of 2HDM and 3HDM scenarios in the alignment limit strictly when all the doublets receive vevs. 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 say ‘Possible’, we mean that decoupling can be achieved with some tuning, while in the last row ‘Partial’ implies that only the first charged scalar decouples.

## 4 Conclusions and outlook

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. We show that charged scalars in multi-doublet scalar extensions of the SM do not necessarily decouple from physical processes, e.g.  $\mu_{\gamma\gamma}$  in the context of this paper, specifically when the potential has an exact symmetry and all the scalars receive vevs.

Here we give a few examples where phenomenology of such scenarios has been studied. A spontaneously broken  $Z_2$  symmetric potential in 2HDM context, with a tiny but nonvanishing  $v_2$ , has been advocated to account for the smallness of the neutrino mass and the stability of a scalar dark matter on a cosmological scale [19]. A few 3HDM examples are also in order. Novel scalar sector phenomenology with exotic scalar decay properties has been studied with exact  $S_3$  symmetric potential [17, 18]. General flavor physics studies were carried out in  $S_3$  [20] as well as in  $A_4$  [16] symmetric scenarios.

In such scenarios, a precisely measured  $\mu_{\gamma\gamma}$  can smell the presence of 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 ( $H_1^\pm$ ), while the second pair ( $H_2^\pm$ ) still do not decouple. Employing the soft breaking terms of an extended global continuous symmetry, namely, SO(2)  $\times$  U(1), the nondecoupling effects of  $H_2^\pm$  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  $\mu_{\gamma\gamma}$ . Keeping in mind the expected accuracy in the measurement of the  $h h h$  vertex in the high luminosity option of LHC or in the future linear collider, whose tree level expression in the alignment limit remains the same as in SM even for multi-doublet Higgs structure,  $\mu_{\gamma\gamma}$  may offer a better bet for diagnosing the underlying layers of the Higgs dynamics.

*To sum up*, 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% – non-inert type multi-Higgs-doublet models with exact discrete symmetries will be constrained. We have demonstrated in this paper that the soft breaking terms in the potential, which are often used in the literature, can play an important rôle in ensuring decoupling, *albeit* with some tuning. To avoid this, one must start with a global continuous symmetry in the potential followed by its soft breaking. In future, it would be 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] D. Fontes, J. Romao, and J. P. Silva, *A reappraisal of the wrong-sign  $h b \bar{b}$  coupling and the study of  $h \rightarrow Z \gamma$* , *Phys.Rev.* **D90** (2014) 015021, [[arXiv:1406.6080](#)].
- [7] 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](#)].
- [8] J. F. Gunion, H. E. Haber, G. L. Kane, and S. Dawson, *The Higgs Hunter’s Guide*, *Front.Phys.* **80** (2000) 1–448.
- [9] 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](#)].
- [10] J. F. Gunion and H. E. Haber, *The CP conserving two Higgs doublet model: The Approach to the decoupling limit*, *Phys.Rev.* **D67** (2003) 075019, [[hep-ph/0207010](#)].

- [11] **CMS** Collaboration, V. Khachatryan et al., *Observation of the diphoton decay of the Higgs boson and measurement of its properties*, [arXiv:1407.0558](#).
- [12] **ATLAS** Collaboration, G. Aad et al., *Measurement of Higgs boson production in the diphoton decay channel in pp collisions at center-of-mass energies of 7 and 8 TeV with the ATLAS detector*, [arXiv:1408.7084](#).
- [13] E. Barradas-Guevara, O. Flix-Beltrn, and E. R. Juregui, *S(3) flavoured Higgs model trilinear self-couplings*, [arXiv:1402.2244](#).
- [14] D. Das and U. K. Dey, *Analysis of an extended scalar sector with  $S_3$  symmetry*, *Phys.Rev.* **D89** (2014) 095025, [[arXiv:1404.2491](#)].
- [15] E. Ma and G. Rajasekaran, *Softly broken  $A(4)$  symmetry for nearly degenerate neutrino masses*, *Phys.Rev.* **D64** (2001) 113012, [[hep-ph/0106291](#)].
- [16] 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](#)].
- [17] 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](#)].
- [18] 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](#)].
- [19] S. Gabriel and S. Nandi, *A New two Higgs doublet model*, *Phys.Lett.* **B655** (2007) 141–147, [[hep-ph/0610253](#)].
- [20] S.-L. Chen, M. Frigerio, and E. Ma, *Large neutrino mixing and normal mass hierarchy: A Discrete understanding*, *Phys.Rev.* **D70** (2004) 073008, [[hep-ph/0404084](#)].