PHYSICAL REVIEW D **91,** 015005 (2015)

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

Gautam Bhattacharyya* and Dipankar Das†

Saha Institute of Nuclear Physics, 1/AF Bidhan Nagar, Kolkata 700064, India
(Received 28 August 2014; published 7 January 2015)

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 analyze 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 noninert 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.

DOI: 10.1103/PhysRevD.91.015005 PACS numbers: 12.60.Fr, 14.80.Fd

I. 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 ρ parameter remains at unity at tree level. Second, it is straightforward to find a combination, namely,

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

[v_i is the vacuum expectation value (vev) of the ith 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 toward 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 [1–6]. We investigate the role of symmetries that 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 *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 the 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 the $h \to \gamma\gamma$ branching ratio can put constraints on the number of additional noninert 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

*gautam.bhattacharyya@saha.ac.in
†d.das@saha.ac.in

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 multidoublet 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

$$\mu_{\gamma\gamma} \equiv \frac{\sigma(pp \to h)}{\sigma^{\rm SM}(pp \to h)} \cdot \frac{{\rm BR}(h \to \gamma\gamma)}{{\rm BR}^{\rm SM}(h \to \gamma\gamma)} = \frac{\Gamma(h \to \gamma\gamma)}{\Gamma^{\rm SM}(h \to \gamma\gamma)}. \tag{2}$$

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

$$g_{hH_i^+H_i^-} = \kappa_i \frac{g m_{i+}^2}{M_W}, \tag{3}$$

where m_{i+} is the mass of the ith charged scalar ($H_i^\pm$). As we will see later, the decoupling or nondecoupling behavior of the ith charged scalar from $\mu_{\gamma\gamma}$ is encoded in κ_i. The expression of the diphoton decay width of the Higgs is given by [8]

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

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

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

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

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

In the limit the charged scalar is very heavy, the quantity $\mathcal{A}_{i+}$ saturates to $1/3$. If κ_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 very close to the SM prediction, then such nondecoupling scenarios will be disfavored. The decoupling would happen only if κ_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.

II. TWO HIGGS-DOUBLET MODELS

We consider a 2HDM with ϕ_1 and ϕ_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-changing neutral current in the fermionic sector. The expression of the scalar potential is displayed below [8]

$$\begin{aligned} V_{\rm 2HDM} = {}& \lambda_1 \left(\phi_1^\dagger \phi_1 - \frac{v_1^2}{2} \right)^2 + \lambda_2 \left(\phi_2^\dagger \phi_2 - \frac{v_2^2}{2} \right)^2 \\ & + \lambda_3 \left(\phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2 - \frac{v_1^2 + v_2^2}{2} \right)^2 \\ & + \lambda_4 ((\phi_1^\dagger \phi_1)(\phi_2^\dagger \phi_2) - (\phi_1^\dagger \phi_2)(\phi_2^\dagger \phi_1)) \\ & + \lambda_5 \left({\rm Re}\phi_1^\dagger \phi_2 - \frac{v_1 v_2}{2} \right)^2 + \lambda_6 ({\rm Im}\phi_1^\dagger \phi_2)^2, \end{aligned} \tag{7}$$

where the λ_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 eight 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 λ_5, can be traded for four physical scalar masses $[m_h, m_H, m_A, m_{1+}(\equiv m_{H^+})]$ and the rotation angle (α) in the neutral CP even sector. The lambdas can be expressed in terms of physical masses as

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

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

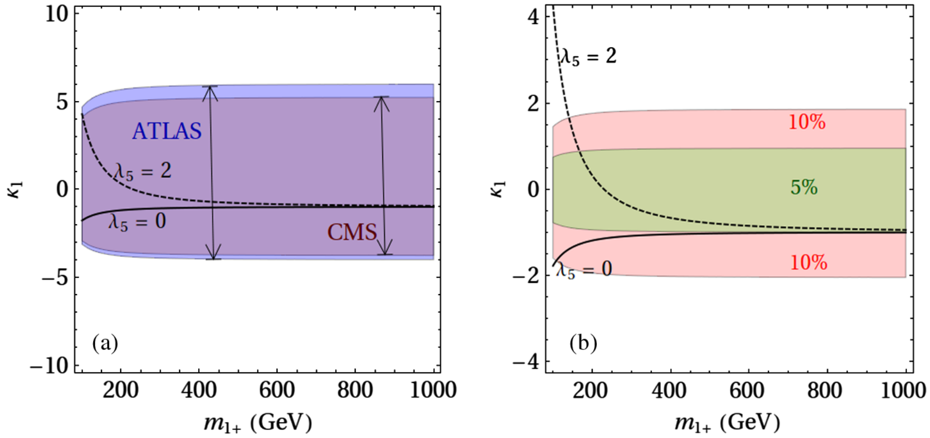

FIG. 1 (color online). (a) We display the constraints on κ_1 in 2HDM coming from the measured values of $\mu_{\gamma\gamma}$ at 95% C.L. by the CMS ($1.14^{+0.26}_{-0.23}$ [11]) and ATLAS (1.17 ± 0.27 [12]) Collaborations. (b) We show what would be the 95% C.L. allowed range of κ_1 if $\mu_{\gamma\gamma}$ is hypothetically measured to be $1 \pm 0.1(0.05)$ in future colliders. In both panels we have plotted Eq. (9) for two different values of λ_5.

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

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

$$\lambda_6 = \frac{2}{v^2}m_A^2. \tag{8e}$$

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 seven independent parameters, out of which five are unknown, namely, $m_H, m_A, m_{1+}, \tan\beta$, and λ_5, while two 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 κ_i) [1,2,4,5,9]

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

Clearly, κ_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 λ_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 κ_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 λ_5 still breaks the U(1) symmetry softly. Then the mass of the pseudoscalar gets related to the soft breaking parameter λ_5 as $m_A^2 = \lambda_5 v^2/2$. In this case, the expression for κ_1 reads [4]

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

In a previous paper [4], we provided a detailed analysis on the unitarity and stability constraints on various combinations of λ_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, $(2\lambda_3 + \lambda_4) \le 16\pi$ implies $(2m_{1+}^2 - m_H^2 - m_A^2 + m_h^2) \le 16\pi v^2$. Also, $(\lambda_1 + \lambda_2 + 2\lambda_3) \le 16\pi/3$ implies $(m_H^2 - m_A^2)(\tan^2\beta + \cot^2\beta) + 2m_h^2 \le 32\pi v^2/3$. 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 κ_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 λ_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 λ_5 is related to the pseudoscalar mass whose splitting with the charged scalar mass is restricted from unitarity.

A. 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 nHDM where $n > 2$. First, we write down the 2HDM potential using a different notation,

$$V_{2\mathrm{HDM}} = m_{11}^2\phi_1^\dagger\phi_1 + m_{22}^2\phi_2^\dagger\phi_2 - (m_{12}^2\phi_1^\dagger\phi_2 + \mathrm{H.c.}) + \frac{\beta_1}{2}(\phi_1^\dagger\phi_1)^2 + \frac{\beta_2}{2}(\phi_2^\dagger\phi_2)^2 + \beta_3(\phi_1^\dagger\phi_1)(\phi_2^\dagger\phi_2) + \beta_4(\phi_1^\dagger\phi_2)(\phi_2^\dagger\phi_1) + \left\{\frac{\beta_5}{2}(\phi_1^\dagger\phi_2)^2 + \mathrm{H.c.}\right\}. \quad (11)$$

This parametrization does not *a priori* assume, unlike the one used in this paper given in Eq. (7), that ϕ_1 or ϕ_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 ϕ_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 an 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 do we obtain the following relations:

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

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 λ_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 λ_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 that has non-SSB origin. As illustrated in the inert doublet case, even with an 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 very easily be 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 λ, an independent dimensionful parameter μ^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 role in addressing phenomenological issues associated with processes like $b \to s\gamma$, $B_s \to \ell^+\ell^-$, and $B \to D^{(*)}\ell\nu$ [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 $h \to \gamma\gamma$. This is because the top quark receives all its mass from SSB, and increasing 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 $M \to \infty$, the top quark contribution can be made to decouple from the diphoton decay width of the Higgs boson.

III. THREE-HIGGS-DOUBLET MODELS

S_3 or A_4 symmetric flavor models are typical examples that employ three Higgs doublets. With ϕ_1, ϕ_2, and ϕ_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),

$$V_{3\mathrm{HDM}}^{S_3} = -\mu_1^2(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2) - \mu_3^2\phi_3^\dagger\phi_3 + \lambda_1(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2)^2 + \lambda_2(\phi_1^\dagger\phi_2 - \phi_2^\dagger\phi_1)^2$$
$$+ \lambda_3\{(\phi_1^\dagger\phi_2 + \phi_2^\dagger\phi_1)^2 + (\phi_1^\dagger\phi_1 - \phi_2^\dagger\phi_2)^2\} + \lambda_4\{(\phi_3^\dagger\phi_1)(\phi_1^\dagger\phi_2 + \phi_2^\dagger\phi_1) + (\phi_3^\dagger\phi_2)(\phi_1^\dagger\phi_1 - \phi_2^\dagger\phi_2) + \mathrm{H.c.}\}$$
$$+ \lambda_5(\phi_3^\dagger\phi_3)(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2) + \lambda_6\{(\phi_3^\dagger\phi_1)(\phi_1^\dagger\phi_3) + (\phi_3^\dagger\phi_2)(\phi_2^\dagger\phi_3)\}$$
$$+ \lambda_7\{(\phi_3^\dagger\phi_1)(\phi_3^\dagger\phi_1) + (\phi_3^\dagger\phi_2)(\phi_3^\dagger\phi_2) + \mathrm{H.c.}\} + \lambda_8(\phi_3^\dagger\phi_3)^2. \quad (13)$$

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]),

$$V_{3\mathrm{HDM}}^{A_4} = -\mu^2(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2 + \phi_3^\dagger\phi_3) + \lambda_1(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2 + \phi_3^\dagger\phi_3)^2 + \lambda_2(\phi_1^\dagger\phi_1\phi_2^\dagger\phi_2 + \phi_2^\dagger\phi_2\phi_3^\dagger\phi_3 + \phi_3^\dagger\phi_3\phi_1^\dagger\phi_1)$$
$$+ \lambda_3(\phi_1^\dagger\phi_2\phi_2^\dagger\phi_1 + \phi_2^\dagger\phi_3\phi_3^\dagger\phi_2 + \phi_3^\dagger\phi_1\phi_1^\dagger\phi_3) + \lambda_4[e^{i\epsilon}\{(\phi_1^\dagger\phi_2)^2 + (\phi_2^\dagger\phi_3)^2 + (\phi_3^\dagger\phi_1)^2\} + \mathrm{H.c.}]. \quad (14)$$

In one plausible scenario, the minimization conditions require that all 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 ϵ.

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:

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

Clearly, the charged scalars do not decouple from the diphoton decay width, since κ_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 ϕ_1 and ϕ_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 $(-\mu_{12}^2\phi_1^\dagger\phi_2 + \mathrm{H.c.})$. The real part of μ_{12}^2 softly breaks the SO(2) symmetry and prevents the occurrence of any massless scalar in the theory. In any case, we assume μ_{12}^2 to be real just like any other parameters in the potential. The relevant minimization conditions are given by

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

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

Note that nonzero μ_{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 $h'VV$ $(V = W, Z)$. The 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 (μ_{12}^2) gets related to the mass of h' as

$$m_{h'}^2 = 2\mu_{12}^2. \quad (17)$$

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 Eqs. (18a) and (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, κ_1 becomes vanishingly small. However, this decoupling does not occur in κ_2, which contains the second charged Higgs mass m_{2+}. It is not difficult to intuitively argue that with an extended global symmetry $\mathrm{SO}(2) \times \mathrm{U}(1)$, together with an extra soft breaking parameter which is related to m_{A2}, decoupling in κ_2 can be ensured. Starting from the softly broken SO(2) symmetric potential, this additional U(1) extension $(\phi_3 \to e^{i\alpha}\phi_3)$ and its soft breaking can be realized by putting $\lambda_7 = 0$ in Eq. (13) and introducing a term that softly breaks this 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

TABLE I. 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 an 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.

Model		Expression for κ_i	Prediction $\mu_{\gamma\gamma}$	Prediction $\mu_{Z\gamma}$	Decoupling?
2HDM	Softly broken Z_2	$-(1+\frac{m_h^2}{2m_{1+}^2}-\frac{\lambda_5 v^2}{2m_{1+}^2})$	Depends on λ_5	Depends on λ_5	Possible
	Exact Z_2	$-(1+\frac{m_h^2}{2m_{1+}^2})$	≤ 0.9	≤ 0.96	No
	Softly broken U(1)	$-(1+\frac{m_h^2}{2m_{1+}^2}-\frac{m_A^2}{m_{1+}^2})$	Depends on m_A	Depends on m_A	Yes
3HDM	Exact S_3	$-(1+\frac{m_h^2}{2m_{i+}^2})$ for $i=1, 2$	≤ 0.8	≤ 0.93	No
	Exact A_4	$-(1+\frac{m_h^2}{2m_{i+}^2})$ for $i=1, 2$	≤ 0.8	≤ 0.93	No
	Softly broken SO(2)	$\kappa_1=-(1+\frac{m_h^2}{2m_{1+}^2}-\frac{m_{h'}^2}{m_{1+}^2})$ $\kappa_2=-(1+\frac{m_h^2}{2m_{2+}^2})$	Depends on $m_{h'}$	Depends on $m_{h'}$	Partial

expressions of κ_i—e.g., see Eqs. (10) and (18a)—only when they are related to *soft* global symmetry breaking parameters.

IV. CONCLUSIONS AND OUTLOOK

To our knowledge, this is the first attempt toward 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 multidoublet 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 the 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 superheavy. In fact, $\mu_{\gamma\gamma}$ can constrain the *number* of such doublets. Table I 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 that 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 does not decouple. Employing the soft breaking terms of an extended global continuous symmetry, namely, $\mathrm{SO}(2)\times\mathrm{U}(1)$, 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 *hhh* 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 the multidoublet 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%—noninert 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 role 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 the future, it would be interesting to explore the consequences of global symmetries in the potential for nondoublet scalar extensions in the present context.

ACKNOWLEDGMENTS

D. D. thanks Department of Atomic Energy, India, for financial support.

[1] A. Djouadi, V. Driesen, W. Hollik, and A. Kraft, Eur. Phys. J. C **1**, 163 (1998).
[2] A. Arhrib, M. C. Peyranere, W. Hollik, and S. Penaranda, Phys. Lett. B **579**, 361 (2004).
[3] W.-F. Chang, J. N. Ng, and J. M. Wu, Phys. Rev. D **86**, 033003 (2012).
[4] G. Bhattacharyya, D. Das, P. B. Pal, and M. Rebelo, J. High Energy Phys. 10 (2013) 081.
[5] P. Ferreira, J. F. Gunion, H. E. Haber, and R. Santos, Phys. Rev. D **89**, 115003 (2014).
[6] D. Fontes, J. Romao, and J. P. Silva, Phys. Rev. D **90**, 015021 (2014).
[7] G. Branco, P. Ferreira, L. Lavoura, M. Rebelo, M. Sher, and J. P. Silva, Phys. Rep. **516**, 1 (2012).
[8] J. F. Gunion, H. E. Haber, G. L. Kane, and S. Dawson, Front. Phys. **80**, 1 (2000).
[9] B. Swiezewska and M. Krawczyk, Phys. Rev. D **88**, 035019 (2013).
[10] J. F. Gunion and H. E. Haber, Phys. Rev. D **67**, 075019 (2003).
[11] V. Khachatryan *et al.* (CMS Collaboration), Eur. Phys. J. C **74**, 3076 (2014).
[12] G. Aad *et al.* (ATLAS Collaboration), arXiv:1408.7084.
[13] E. Barradas-Guevara, O. Flix-Beltrn, and E. R. Juregui, Phys. Rev. D **90**, 095001 (2014).
[14] D. Das and U. K. Dey, Phys. Rev. D **89**, 095025 (2014).
[15] E. Ma and G. Rajasekaran, Phys. Rev. D **64**, 113012 (2001).
[16] R. de Adelhart Toorop, F. Bazzocchi, L. Merlo, and A. Paris, J. High Energy Phys. 03 (2011) 035.
[17] G. Bhattacharyya, P. Leser, and H. Pas, Phys. Rev. D **83**, 011701 (2011).
[18] G. Bhattacharyya, P. Leser, and H. Pas, Phys. Rev. D **86**, 036009 (2012).
[19] S. Gabriel and S. Nandi, Phys. Lett. B **655**, 141 (2007).
[20] S.-L. Chen, M. Frigerio, and E. Ma, Phys. Rev. D **70**, 073008 (2004).