Scalar sector properties of two-Higgs-doublet models with a global U(1) symmetry

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Abstract

We analyze the scalar sector properties of a general class of two-Higgs-doublet models which has a global U(1) symmetry in the quartic terms. We find constraints on the parameters of the potential from the considerations of unitarity of scattering amplitudes, the global stability of the potential and the $\rho$-parameter. We concentrate on the spectrum of the non-standard scalar masses in the decoupling limit which is preferred by the Higgs data at the LHC. We exhibit charged-Higgs induced contributions to the diphoton decay width of the 125 GeV Higgs boson and its correlation with the corresponding $Z\gamma$ width.

1 Introduction

A new boson with a mass of nearly 125 GeV has been observed by the CMS and ATLAS collaborations at the LHC \cite{1,2}. It is too early to comment on its spin. Preliminary analyses hint towards a $0^+$ state \cite{3}. Different production mechanisms and decay rates of this new scalar are compatible with what we expect for the Standard Model (SM) Higgs boson. Currently it is a subject of intense study \cite{4-15} whether the observed scalar is alone or is a part of a richer scalar sector. The simplest scenario entertaining this possibility is the class of two-Higgs-doublet models (2HDMs) \cite{16}. In these extensions, the value of the electroweak $\rho$-parameter remains unity at tree level as in the SM.

A general problem that one encounters with 2HDMs is the presence of Higgs-mediated flavor changing neutral currents (FCNCs). To ensure natural flavor conservation, one usually relies on a symmetry that completely eliminates tree level Higgs mediated FCNC. One example is the imposition of a discrete $Z_2$ symmetry, assigned in such a way that all fermions with the same electric charge obtain their masses from the vacuum expectation value (VEV) of one particular scalar doublet \cite{17,19}. A larger symmetry of this kind is a global U(1) symmetry, employed in a different context in Peccei-Quinn model. Alternatively, one can admit tree level Higgs mediated FCNCs which are suppressed by small factors involving small entries of the CKM matrix \cite{20,22}. Such a model with a continuous global U(1) symmetry was proposed Branco, Grimus and Lavoura \cite{23}, which and its later generalizations \cite{24,25} fall under the category of models with minimal flavor violation \cite{26}.

We however concentrate only on the scalar sector with a global U(1) symmetry, regardless of the transformation properties of the fermions under this symmetry. Thus our observations extend well beyond the scope of any such individual flavor model. Ferreira and Jones \cite{27} have done a thorough analysis of stability and perturbativity bounds on the parameter space of such models together with the consideration of non-observation of any scalar at LEP. Our motivation in this paper is to explore the scalar potential with a softly broken U(1) symmetry in view of the observation of a 125 GeV Higgs boson ($h$) at the LHC, using constraints from unitarity of scalar scattering cross-sections, stability.
of the potential and electroweak precision tests. These considerations restrict the spectrum of the non-standard scalars. Since the LHC Higgs data seem to be compatible with the SM expectations, we restrict ourselves to the decoupling limit, explained later. In this limit, we study the charged-Higgs induced contribution to $h \rightarrow \gamma\gamma$ decay width and its correlation to $h \rightarrow Z\gamma$.

## 2 Bounds on masses from stability and unitarity

We take the scalar potential as follows [28]:

$$V(\phi_1, \phi_2) = \lambda_1 \left( \phi_1 \phi_1 - \frac{v_1^2}{2} \right)^2 + \lambda_2 \left( \phi_2 \phi_2 - \frac{v_2^2}{2} \right)^2$$

$$+ \lambda_3 \left( \phi_1^\dagger \phi_1 + \phi_2^\dagger \phi_2 - \frac{v_1^2 + v_2^2}{2} \right) + \lambda_4 \left( (\phi_1^\dagger \phi_1) (\phi_2^\dagger \phi_2) - (\phi_1^\dagger \phi_2) (\phi_2^\dagger \phi_1) \right)$$

$$+ \lambda_5 \left( \frac{1}{2} (\phi_1^\dagger \phi_2 + \phi_2^\dagger \phi_1 - v_1 v_2) \right)^2 + \lambda_6 \left( \frac{1}{2v} (\phi_1^\dagger \phi_2 - \phi_2^\dagger \phi_1) \right)^2,$$  

(1)

where the $\lambda$'s are real because of hermiticity of the Lagrangian. All terms in the potential are invariant under the discrete symmetry:

$$\phi_1 \rightarrow \phi_1, \quad \phi_2 \rightarrow -\phi_2,$$  

(2)

except a term

$$-\frac{1}{2} \lambda_5 v_1 v_2 (\phi_1^\dagger \phi_2 + \phi_2^\dagger \phi_1)$$  

(3)

which breaks the symmetry only softly. We use the following parametrization for the Higgs doublets:

$$\phi_i(x) = \left( \frac{w_i^+(x)}{\sqrt{2}(v_i + h_i(x) + iz_i(x))} \right),$$  

(4)

where the $v_i$ denote the vacuum expectation values (VEVs) of the two doublets, assumed to be real and positive without any loss of generality. It is useful to define the following parameters,

$$\tan \beta = \frac{v_2}{v_1} \quad (0 \leq \beta \leq \frac{\pi}{2}),$$  

(5a)

$$v = \sqrt{v_1^2 + v_2^2} = 246 \text{ GeV}.$$  

(5b)

There will be five physical Higgs bosons in this model. We denote the charged ones by $\xi^\pm$, the CP-odd one by $A$, and use the symbols $H$ and $h$ to denote the heavy and light CP even Higgs bosons respectively. There will also be the combinations $\omega^\pm$, $\zeta$ which are the three would-be Goldstone bosons eaten by the gauge bosons. These combinations are given by

$$\begin{pmatrix} \omega^\pm \\ \xi^\pm \end{pmatrix} = \begin{pmatrix} c_\beta & s_\beta \\ -s_\beta & c_\beta \end{pmatrix} \begin{pmatrix} w_1^\pm \\ w_2^\pm \end{pmatrix},$$  

(6)

$$\begin{pmatrix} \zeta \\ A \end{pmatrix} = \begin{pmatrix} c_\beta & s_\beta \\ -s_\beta & c_\beta \end{pmatrix} \begin{pmatrix} z_1 \\ z_2 \end{pmatrix},$$  

(7)

$$\begin{pmatrix} H \\ h \end{pmatrix} = \begin{pmatrix} c_\alpha & s_\alpha \\ -s_\alpha & c_\alpha \end{pmatrix} \begin{pmatrix} h_1 \\ h_2 \end{pmatrix},$$  

(8)
where \(c_\alpha \equiv \cos \alpha\), \(s_\alpha \equiv \sin \alpha\), and likewise for \(\beta\). The angle of rotation for the matrix in Eq. (5) is given by

\[
\tan 2\alpha = \frac{2(\lambda_3 + \frac{1}{4}\lambda_5)v_1v_2}{\lambda_1 v_1^2 - \lambda_2 v_2^2 + (\lambda_3 - \frac{1}{4}\lambda_5)(v_1^2 - v_2^2)}, \quad (-\frac{\pi}{2} \leq 2\alpha \leq \frac{\pi}{2}).
\] (9)

We have \(v_1\), \(v_2\), and six \(\lambda\)-parameters as independent parameters in the potential of Eq. (11). They are related to the four physical Higgs boson masses by the following relations:

\[
\begin{align*}
\lambda_1 &= -\frac{1}{2v^2c_\beta^2}\left[c_\alpha^2m_H^2 + s_\alpha^2m_h^2 - \frac{s_\alpha c_\alpha}{\tan \beta}m_H^2 - m_h^2\right] - \frac{\lambda_5^5}{4}(\tan^2\beta - 1),
\end{align*}
\] (10a)

\[
\begin{align*}
\lambda_2 &= \frac{1}{2v^2s_\beta^2}\left[s_\alpha^2m_H^2 + c_\alpha^2m_h^2 - s_\alpha c_\alpha \tan \beta(m_H^2 - m_h^2)\right] - \frac{\lambda_5^5}{4}\left(\frac{1}{\tan^2\beta} - 1\right),
\end{align*}
\] (10b)

\[
\begin{align*}
\lambda_3 &= -\frac{1}{2v^2s_\beta^2}c_\alpha(m_H^2 - m_h^2) - \frac{\lambda_5^5}{4},
\end{align*}
\] (10c)

\[
\begin{align*}
\lambda_4 &= \frac{2}{v^2}m_\xi^2,
\end{align*}
\] (10d)

\[
\begin{align*}
\lambda_6 &= \frac{2}{v^2}m_A^2.
\end{align*}
\] (10e)

Thus an alternative way of counting the independent parameters is through the four masses, the two angles \(\alpha\) and \(\beta\), the electroweak VEV \(v\), and the parameter \(\lambda_5\), which appear on the right-hand sides of Eq. (10). In this set of eight parameters, \(v\) is known as in Eq. (5), and so is the lightest CP-even Higgs mass as 125 GeV.

The other thing that we will need from the scalar potential is the cubic coupling involving \(h\xi^+\xi^-\), denoted by \(g_{h\xi\xi}\), which will be employed in the \(h \to \gamma \gamma\) and \(h \to Z\gamma\) decay rates of the Higgs boson. This coupling is given by

\[
g_{h\xi\xi} = v\left[(\lambda_1 + \lambda_2)s_\alpha s_\beta s_{2\beta} - \lambda_2 s_{2\beta}c_\beta -\alpha - (2\lambda_3 + \lambda_4)s_{\beta -\alpha} + \frac{1}{2}\lambda_5 s_{2\beta}c_{\alpha +\beta}\right].
\] (11)

In terms of physical masses, this coupling can be written as

\[
g_{h\xi\xi} = -\frac{1}{v}\left[(2m_H^2 - \lambda_5^5v^2)\frac{\cos(\alpha + \beta)}{\sin 2\beta} + (2m_\xi^2 - m_h^2)\sin(\beta - \alpha)\right].
\] (12)

In the present work, we explore the consequences of a global symmetry that is larger than that given in Eq. (2):

\[
\phi_1 \to \phi_1, \quad \phi_2 \to e^{i\theta}\phi_2.
\] (13)

On the quartic terms of the scalar potential, this symmetry is realized by putting

\[
\lambda_5 = \lambda_6,
\] (14)

which means that the potential now reads

\[
V(\phi_1, \phi_2) = \lambda_1 \left(\phi_1^+\phi_1 - \frac{v_1^2}{2}\right)^2 + \lambda_2 \left(\phi_2^+\phi_2 - \frac{v_2^2}{2}\right)^2 + \lambda_3 \left(\phi_1^+\phi_1 + \phi_2^+\phi_2 - \frac{v_1^2 + v_2^2}{2}\right)^2
\]

\[
+ \lambda_4 \left((\phi_1^+\phi_1)(\phi_2^+\phi_2) - (\phi_1^+\phi_2)(\phi_2^+\phi_1)\right) + \lambda_5 \left|\phi_1^+\phi_2 - \frac{v_1v_2}{2}\right|^2.
\] (15)
The term of Eq. (5), which breaks the discrete symmetry softly, is a soft explicit breaking term of the global U(1) symmetry as well. In the limit $\lambda_5 = 0$, the U(1) symmetry is exact in the potential, and the spontaneous breaking of it through $v_2 \neq 0$, since only $\phi_2$ undergoes a nontrivial U(1) phase transformation, gives rise to the Goldstone boson $A$. Since $\lambda_5 \to 0$ corresponds to an enhanced symmetry, it is expected to remain small a la 't Hooft, and consequently, the CP-odd $A$ can remain naturally light. We will later see under what conditions we can entertain a light $A$ boson, and what are its consequences.

Conditions for the potential being bounded from below were examined for more general potentials in 2HDM [30][31]. Using the relation of Eq. (14), these conditions read

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

While the conditions in Eq. (16) put lower bounds on certain combinations of the quartic couplings, there exist upper bounds on these couplings arising from the consideration of perturbative unitarity [32]. Scattering amplitudes involving longitudinal gauge bosons and Higgs bosons comprise the elements of an $S$-matrix, having 2-particle states as rows and columns. The eigenvalues of this matrix are restricted by $|a_0| < 1$, where $a_0$ is the $l = 0$ partial wave amplitude. These conditions translate into upper limits on combinations of Higgs quartic couplings, which for multi-Higgs models have been derived by different authors [33][36]. Imposing the condition of Eq. (14), these constraints assume the following form:

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

It is worth noting [23] at this stage that if we rotate the basis $h_1$-$h_2$ by the same angle $\beta$ which appears in Eq. (7), we obtain the states

\begin{equation}
\begin{pmatrix}
H^0 \\
R
\end{pmatrix} = \begin{pmatrix}
c_\beta & s_\beta \\
-s_\beta & c_\beta
\end{pmatrix} \begin{pmatrix} h_1 \\
h_2
\end{pmatrix},
\end{equation}

which has the property that $H^0$ has the exact SM couplings with the fermions and gauge bosons. The state $R$ does not have any cubic gauge coupling at the tree level. It can however have flavor changing Yukawa couplings. The lighter CP even mass eigenstate, $h$, is related to $H^0$ and $R$ via the transformation

\begin{equation}
h = \sin(\beta - \alpha)H^0 + \cos(\beta - \alpha)R.
\end{equation}
If it is eventually settled that the Higgs boson observed at the LHC has SM-like gauge and Yukawa couplings, then we will require

\[ \sin(\beta - \alpha) \approx 1, \tag{20} \]

which has been referred to as the decoupling limit \[31\]. Unless otherwise stated, we make the following assumptions in all our subsequent analysis:

- \( \beta - \alpha = \pi/2; \)
- Eq. (14) holds;
- \( m_h = 125 \text{ GeV}; \)
- \( m_\xi > 100 \text{ GeV}, \) which is a rough lower bound from direct searches \[37\].

The resulting constraints on the scalar masses imposed by the stability and unitarity conditions in Eqs. (16) and (17) have been plotted in Fig. 1 for \( \tan \beta = 1, 5 \) and 10 by performing random scan over all non-standard scalar masses.

We note at this point that the splitting between the heavy scalar masses is also constrained by the oblique electroweak \( T \)-parameter. In the present case, the expression of the \( T \)-parameter in the decoupling limit is given by \[38, 39\]

\[
T = \frac{1}{16\pi\sin^2 \theta_w M_W^2} \left[ F(m_\xi^2, m_H^2) + F(m_\xi^2, m_A^2) - F(m_H^2, m_A^2) \right],
\tag{21}
\]

with

\[
F(x, y) = \frac{x + y}{2} - \frac{xy}{x - y} \ln(x/y).
\tag{22}
\]

Taking the new physics contribution to the \( T \)-parameter as \[40\]

\[
T = 0.05 \pm 0.12,
\tag{23}
\]

we project the 2\( \sigma \) constraints in Fig. 1 for \( \tan \beta = 5 \) and 10. The following salient features emerge from the plots.

1. There is a correlation between \( m_A \) and \( m_H \) which gets stronger for larger values of \( \tan \beta \), to the extent that they become nearly degenerate once \( \tan \beta \) crosses 10. To understand this, we note that Eqs. (16a) and (16b), along with Eq. (17c) for example, imply the inequality

\[
0 \leq \lambda_1 + \lambda_2 + 2\lambda_3 \leq \frac{16\pi}{3}.
\tag{24}
\]

In terms of the scalar masses in Eq. (10), this reads

\[
0 \leq (m_H^2 - m_A^2)(\tan^2 \beta + \cot^2 \beta) + 2m_h^2 \leq \frac{32\pi v^2}{3}.
\tag{25}
\]

Clearly, for \( \tan \beta \) away from unity, \( H \) and \( A \) are almost degenerate.
2. There is a similar correlation between $m_H$ and $m_\xi$, but this time without any dependence on $\tan \beta$. This can again be seen from the inequality

$$\left| 2m_\xi^2 - m_H^2 - m_A^2 + m_h^2 \right| \leq 16\pi v^2 ,$$

which follows from Eq. (17b).

3. As regards the non-standard scalars, the unitarity conditions essentially apply on the difference of their squared masses. Thus, any individual mass can be arbitrarily large without affecting the unitarity conditions. This conclusion crucially depends on the existence of a U(1) symmetry of the potential. When the symmetry of the potential is only a discrete $Z_2$, considerations of unitarity do restrict the individual non-standard masses \cite{34,36}.

4. To provide an intuitive feel on the constraints from the $T$-parameter, we assume $m_H = m_A$, which is anyway dictated by the unitarity constraints for $\tan \beta$ somewhat away from unity. It then follows from Eq. (21) that the splitting between $m_\xi$ and $m_H$ is approximately 50 GeV, for
\( |m_\xi - m_H| \ll m_\xi, m_H \). It turns out from Fig. [1] that the constraints from the \( T \)-parameter are stronger than that from unitarity and stability.

For \( \tan \beta = 1 \), unitarity and stability do not compel \( m_H \) and \( m_A \) to be very close. In this case, the \( T \)-parameter cannot give any definitive constraints in the planes of the heavy scalar masses, unlike the unitarity and stability constraints. For this reason, we have shown only the latter constraints in Fig. [1] for \( \tan \beta = 1 \).

5. Eq. (21) shows that the contribution to the \( T \)-parameter from scalar loops is vanishingly small irrespective of the value of \( m_A \) as long as \( m_H \approx m_\xi \). As Fig. [1] shows, even \( m_A = 0 \) is allowed for \( \tan \beta = 1 \) and \( \tan \beta = 5 \). A light pseudoscalar is experimentally allowed, and various aspects of its phenomenology in the context of 2HDM have been discussed in the literature [41–43]. The value \( \tan \beta = 10 \) is already large enough so that the conclusion of item [1] applies, and \( m_H \) drags \( m_A \) with it.

6. Thus, for moderate or large \( \tan \beta \), the unitarity and stability constraints, together with the constraints coming from the \( T \)-parameter, imply that all three heavy scalar states are nearly degenerate in the decoupling limit.

3 Modifications in Higgs decay width

Since we are working in the decoupling limit, the couplings of \( h \) with the fermions and gauge bosons will be exactly like in the SM. The production cross section of \( h \) will therefore be as expected in the SM. All the tree level decay widths of \( h \) will also have the SM values for the same reason. Loop induced decays like \( h \rightarrow \gamma \gamma \) and \( h \rightarrow Z \gamma \) will however have additional contributions from virtual charged scalars (\( \xi^\pm \)). Since the branching fractions of such decays are tiny, the total decay width is hardly modified.

The contribution of the \( W \)-boson loop and the top loop diagrams to \( h \rightarrow \gamma \gamma \) and \( h \rightarrow Z \gamma \) are same as in the SM. As regards the charged scalar induced loop, we first parametrize the cubic coupling \( g_h \xi \xi \), given in Eq. (11), in the following way:

\[
g_h \xi \xi = \kappa \frac{g m_\xi^2}{M_W}, \tag{27}
\]

where \( \kappa \) is dimensionless. The diphoton decay width is then given by [44, 45]:

\[
\Gamma(h \rightarrow \gamma \gamma) = \frac{\alpha^2 g^2 m_h^2}{2^{10} \pi^3 M_W^2} \left[ F_W + \frac{4}{3} F_t + \kappa F_\xi \right]^2, \tag{28}
\]

where, introducing the notation

\[
\tau_x \equiv (2m_x/m_h)^2, \tag{29}
\]

the values of \( F_W, F_t \) and \( F_\xi \) are given by

\[
F_W = 2 + 3\tau_W + 3\tau_W (2 - \tau_W) f(\tau_W), \tag{30a}
\]

\[
F_t = -2\tau_t [1 + (1 - \tau_t) f(\tau_t)], \tag{30b}
\]

\[
F_\xi = -\tau_\xi [1 - \tau_\xi f(\tau_\xi)]. \tag{30c}
\]
With our assumptions about the masses declared earlier, \( \tau_x > 1 \) for \( x = W, t, \xi \). Then

\[
f(\tau) = \left[ \sin^{-1}\left(\sqrt{1/\tau}\right) \right]^2.
\]  

(31)

The decay width for \( h \to Z\gamma \) can analogously be written as:

\[
\Gamma(h \to Z\gamma) = \frac{\alpha^2 g^2 m_h^3}{2^9 \pi^3 M_W^2} |A_W + A_t + \kappa A_\xi|^2 \left(1 - \frac{M_Z^2}{m_h^2}\right)^3,
\]

(32)

where, introducing

\[
\eta_x = \left(\frac{2m_x}{M_Z}\right)^2,
\]

(33)

the values of \( A_W, A_t \) and \( A_\xi \) are given by \[28\]

\[
A_W = \cot \theta_w \left[ 4(\tan^2 \theta_w - 3)I_2(\tau_W, \eta_W) \right. \\
+ \left. \left\{ \left( \frac{5}{2} + \frac{2}{7W} \right) - \left( 1 + \frac{2}{7W} \right) \tan^2 \theta_w \right\} I_1(\tau_W, \eta_W) \right],
\]

(34a)

\[
A_t = \frac{4\left(\frac{1}{2} - \frac{4}{3} \sin^2 \theta_w\right)}{\sin \theta_w \cos \theta_w} \left[I_2(\tau_t, \eta_t) - I_1(\tau_t, \eta_t)\right],
\]

(34b)

\[
A_\xi = \frac{(2 \sin^2 \theta_w - 1)}{\sin \theta_w \cos \theta_w} I_1(\tau_\xi, \eta_\xi).
\]

(34c)

The functions \( I_1 \) and \( I_2 \) are given by

\[
I_1(\tau, \eta) = \frac{\tau \eta}{2(\tau - \eta)} + \frac{\tau^2 \eta^2}{2(\tau - \eta)^2} \left[ f(\tau) - f(\eta) \right] + \frac{\tau^2 \eta}{(\tau - \eta)^2} \left[ g(\tau) - g(\eta) \right],
\]

(35a)

\[
I_2(\tau, \eta) = -\frac{\tau \eta}{2(\tau - \eta)^2} \left[ f(\tau) - f(\eta) \right],
\]

(35b)

where the function \( f \) has been defined in Eq. (31). Since \( \tau_x, \eta_x > 1 \) for \( x = W, t, \xi \), the function \( g \) assumes the following form:

\[
g(a) = \sqrt{a - 1} \sin^{-1}\left(\sqrt{1/a}\right).
\]

(36)

In the decoupling limit, the parameter \( \kappa \) which appears in Eqs. (27), (28) and (32) is given by

\[
\kappa = \frac{1}{m^2_\xi} (m^2_A - m^2_\xi - \frac{1}{2} m^2_h).
\]

(37)

The appearance of \( m_A \) in Eq. (37) is merely an artefact of the U(1) symmetry in the scalar potential which enforces Eq. (14). In the more general potential of Eq. (1), the expression for \( \kappa \) involves \( \lambda_5 \), which has nothing to do with \( m_A \). The decoupling behaviour of \( \kappa \) for large \( m_\xi \) is not then guaranteed, as also noted in Refs. [29, 46]. However, in the present scenario, unitarity conditions of Eqs. (25) and (26) bound the splitting between heavy scalar masses, ensuring smooth decoupling of \( \kappa \) with increasing \( m_\xi \). As we have noticed, this splitting is also controlled by the \( T \)-parameter.
Figure 2: The left panels show allowed regions in $m_\xi$-$\mu_{\gamma\gamma}$ plane for three values of $\tan\beta$, while the right panels show correlation between $\mu_{\gamma\gamma}$ and $\mu_{Z\gamma}$ for the same choices of $\tan\beta$. The regions shown in red are allowed by unitarity and stability, while the black regions additionally pass the $T$-parameter test. We put $m_\xi > 100$ GeV and $m_H > m_h$. 
In our case, the quantities $\mu_{\gamma\gamma}$ and $\mu_{Z\gamma}$, defined through the equations

$$
\mu_{\gamma\gamma} = \frac{\sigma(pp \rightarrow h)}{\sigma^{SM}(pp \rightarrow h)} \frac{\text{BR}(h \rightarrow \gamma\gamma)}{\text{BR}^{SM}(h \rightarrow \gamma\gamma)},
$$

$$
\mu_{Z\gamma} = \frac{\sigma(pp \rightarrow h)}{\sigma^{SM}(pp \rightarrow h)} \frac{\text{BR}(h \rightarrow Z\gamma)}{\text{BR}^{SM}(h \rightarrow Z\gamma)},
$$

assume the following forms:

$$
\mu_{\gamma\gamma} = \frac{\Gamma(h \rightarrow \gamma\gamma)}{\Gamma^{SM}(h \rightarrow \gamma\gamma)} = \left| F_W + \frac{4}{3} F_t + \kappa F_\xi \right|^2 / \left| F_W + \frac{4}{3} F_t \right|^2,
$$

$$
\mu_{Z\gamma} = \frac{\Gamma(h \rightarrow Z\gamma)}{\Gamma^{SM}(h \rightarrow Z\gamma)} = \left| A_W + A_t + \kappa A_\xi \right|^2 / \left| A_W + A_t \right|^2.
$$

In Fig. 2, we show the variation of $\mu_{\gamma\gamma}$ against $m_\xi$ and the correlation between $\mu_{\gamma\gamma}$ and $\mu_{Z\gamma}$ for $\tan \beta = 1, 5, 10$. When we show the variation of $\mu_{\gamma\gamma}$ with $m_\xi$, we take into consideration all values of $m_H$ and $m_A$ which are allowed in Fig. 1. As in Fig. 1, the red points are those which are allowed by unitarity and stability constraints, while the superimposed black points are allowed by the $T$-parameter. The points allowed by unitarity prefer suppression in $\mu_{\gamma\gamma}$ compared to the SM expectation. Note that large suppressions appear near the lower end of $m_\xi$ values in Fig. 2 for $\tan \beta = 1$ and 5, but not for $\tan \beta = 10$. This is because small values of $m_A$, including $m_A = 0$, are allowed for $\tan \beta = 1$ and $\tan \beta = 5$, which give sizable negative values of $\kappa$ through Eq. (37) even for small $m_\xi$. The correlation between $\mu_{\gamma\gamma}$ and $\mu_{Z\gamma}$ can in principle be used for discriminating new physics models with increased sensitivity in the future course of the LHC run.

## 4 Conclusions

In this paper, we have studied quantitative correlation among the non-standard scalar masses in a class of 2HDM with a global U(1) symmetry in the potential. We outline below the salient features of our analysis. We derive our constraints from the consideration of unitarity of scattering amplitudes and the global stability of the potential. We additionally superimpose the constraints from the oblique electroweak $T$-parameter on these plots. We have restricted our analysis to the decoupling limit which entails a relation between the parameters $\beta$ and $\alpha$, where the 125 GeV Higgs boson has SM-like couplings. A crucial observation is that when $\tan \beta$ stays close to unity, the CP-odd $A$ can be light, and it is in this limit that we obtain the maximum deviation (in fact, a suppression) in the diphoton decay width. We also observe that for values of $\tan \beta \sim 5$ or larger, all the three non-standard scalar masses are roughly degenerate. More specifically, in this limit unitarity dictates $m_H$ and $m_A$ to be almost equal and $|m_\xi^2 - m_H^2|$ to be small, while the $T$-parameter restricts $|m_\xi - m_H|$ to be very small. Another interesting observation is that the charged-Higgs induced contribution to the diphoton decay amplitude depends on $(m_\xi^2 - m_A^2)/m_\xi^2$, for which the numerator is of order $v^2$, and therefore the contribution decouples with increasing $m_\xi$. In the absence of a global U(1) symmetry in the potential, the parameter $\lambda_5$ cannot be related to $m_A$, and as a result, the above decoupling behaviour is not apparent. It is also important to note that, thanks to the global U(1) symmetry, unitarity restricts mass-squared differences and not the individual masses of the non-standard scalars.
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