Boosted scalar confronting 750 GeV di-photon excess

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ABSTRACT: We consider the di-photon signal arises from two bunches of collimated photon jets emitting from a pair of highly boosted scalars. Following the discussion of detecting the photon jets at the collider, we extend the two-Higgs-doublet model (2HDM) by adding a gauge singlet scalar. To explain the di-photon excess which is recently observed at the first 13 TeV run of the LHC, the mixing between the heavy doublet state and the newly added singlet is crucially needed. After the mixing, one can have a heavy Higgs state $Y_2$ at 750 GeV and a very singlet-like scalar $Y_1$ of sub-GeV, which would be highly boosted through the $Y_2$ decay. Both real singlet and complex singlet extension are studied. It turns out that only the complex model can yield the 1-10 fb cross section in the di-photon final state in accompany with the decay length of the order of 1 m for the $Y_1$. This complex model parametrically predicts the width of 750 GeV resonance $\gtrsim 1$ GeV. In addition, the pseudoscalar component of the singlet in this model is naturally stable and hence could be a dark matter candidate.
1 Introduction

Recently the ATLAS [1] and CMS [2] collaborations at the LHC presented results for the di-photon search at energy $\sqrt{s} = 13$ TeV. Both collaborations have found an excess over the background at $m_{\gamma\gamma} \sim 750$ GeV. The local significances of this signal are $3.9 \sigma$ and $2.6 \sigma$ for the ATLAS and CMS collaborations, respectively. We summarize the LHC data in Table 1 for the excess observed above expected background over an interval centered on $m_{\gamma\gamma} \approx 750$ GeV. Meanwhile, it is important to note that no excess is observed in any other channels, including the $t\bar{t}$, $hh$, $WW$, $ZZ$ and di-jet final states.

Since the announcement of this di-photon excess, a plethora of scenarios have been studied for interpreting this excess through the direct decay of a heavy resonance (to a pair of photon) in the context of both model-independent approaches and concrete models. The simple and natural candidates for the 750 GeV resonance include the heavy CP-even field $H$ and the CP-odd field $A$ of the popular two-Higgs-doublet model (2HDM) or the
minimal supersymmetric SM (MSSM). However, it has been shown in [3] that having the $H$ and/or $A$ at $750$ GeV, the cross section in the di-photon final state via the gluon-fusion production $\sigma(gg \rightarrow H/A)BR(H/A \rightarrow \gamma\gamma)$ approaches at most $10^{-2}$ fb, which is three orders smaller than the level observed at the LHC $13$ TeV run. This results from the fact that the branching ratio of their decay into $\gamma\gamma$ (mediated by the top quark loop and, for the case of the $H$, the $W^\pm$ loop) is maximally of order of $\sim 10^{-5}$ due to the dominance of tree-level decays into $t\bar{t}$ and/or $W^+W^-/ZZ$ (the latter being possibly significant away from the alignment limit for the SM-like $h$ at $m_h \simeq 125$ GeV). The contribution of the charged Higgs to the loop-induced $H\gamma\gamma$ coupling is negligible. Therefore, the minimal 2HDM cannot explain the di-photon excess. To enhance the di-photon signal, one considered to extend the 2HDM by adding extra heavy colored states [4–16] that would give additional loop contributions to the $Xgg$ and/or $X\gamma\gamma$ couplings of the $750$ GeV resonance, which we generically denote by $X$. There also exists many other works related to direct di-photon production [17–145].

Other than the direct decay of a resonance $X(750)$ into di-photon, alternative topologies that could explain the excess have been studied. These include the ideas that i) the di-photon arises from a three-body decay [146–151] in association with an invisible particle or ii) two bunches of collimated photon jets emitting from a pair of highly boosted (pseudo)scalars [152–161]. For the latter case, if there is a (pseudo)scalar $Y$ of sub-GeV scale, its decay into jets will be enormously suppressed and thereby leads to a substantial branching ratio to di-photon. Another element of achieving this scenario is to require a heavy scalar $X$ at $750$ GeV which decays into two sub-GeV scalar particles. Due to the huge mass difference, the light scalar particles will be highly boosted so each photon pair could be identified as photon jet. The CP-odd Higgs $A$ of the 2HDM could have been a good candidate for this light state. However, the presence of a very light $A$ would prevent the non-SM CP-even Higgs $H$ from being heavier than $\simeq 630$ GeV due to the electroweak precision observables (EWPO) once the lightest CP-even Higgs $h$ is identified as the SM one observed at $125$ GeV [162]. \footnote{This conclusion also holds for $m_A$ in the scenario where the heavier CP-Higgs $H$ being $125$ GeV is SM-like.} Therefore, we extend the 2HDM by adding a gauge singlet scalar field. Both real scalar and complex scalar models will be studied in this paper. For our purpose, we allow the mixing between the singlet field $S$ and the heavy Higgs $H$ and identify the heavy mass eigenstate $Y_2$ as the $750$ GeV resonance, while the light mass eigenstate $Y_1$ produces photon jet. Since no additional colored fermion is introduced, the

\begin{table}[h]
\centering
\caption{Summary of the LHC di-photon excess at the invariant mass of $\sim 750$ GeV from the ALTAS [1] and CMS [2] collaborations at energy $\sqrt{s} = 13$ TeV ($\sigma$ is the local statistical significance).}
\begin{tabular}{|l|l|l|}
\hline
 & ATLAS @ $\sqrt{s} = 13$ TeV [1] & CMS @ $\sqrt{s} = 13$ TeV [2] \\
\hline
Excess Events & 14 with 3.9$\sigma$ & 10 with 2.6$\sigma$ \\
$\sigma(pp \rightarrow \gamma\gamma)$ & $(10 \pm 3)$ fb & $(6 \pm 3)$ fb \\
\hline
\end{tabular}
\end{table}
$Y_2$ has to contain sizeable doublet component, otherwise it will be difficult to be produced via gluon fusion. On the contrary, the light state $Y_1$ tends to be singlet-like. We find that it is possible to explain the 750 GeV di-photon excess in the complex singlet model, while the real singlet model is strongly limited by the problem of a too long decay length of the boosted light scalar. Very recently, this idea has been applied to interpret the 750 GeV di-photon excess in the context of NMSSM [157–159].

This paper is organized as follows. In Sec. 2, we introduce properties of photon jet and its implications on di-photon excess. In Sec. 3, we discuss di-photon signal from gluon fusion and its decay. Two attempts to explain this di-photon excess are made in Sec. 4 and in Sec. 5 with real and complex singlet scalar extensions to 2HDM respectively. The collider prospects are discussed in Sec. 6. Finally, Sec. 7 summarizes our main results of this work.

2 Detecting the boosted photon jets in the collider

The photon-jet is a special object that consists of a cluster of (nearly) collinear photons which have a signature similar to that of a single photon. This idea was initially proposed in [163] and has recently been applied to explain this di-photon excess [154–161]. It could be generated from the decay of a highly boosted light particle (sub-GeV). In this section we discuss two technical issues, angular separation and decay length, which are very crucial in detecting a photon-jet in the collider machine.

Suppose a particle $X(750)$, once produced in the hadron collider, instantaneously decays to two highly boosted $Y$'s. For each $Y$, it decays into a pair of photons with a branching ratio close to 100%. The angular separation of two outgoing photons (which forms a photon-jet) will be $\theta \sim \frac{2m_Y}{p_Y} \sim \frac{2m_Y}{375}$, assuming 375 GeV is the expected momentum for $X(750) \rightarrow YY$ process. If this angular separation is smaller than the resolution of the ECAL in the CMS/ATLAS, the photon jet could probably be mistagged as a single photon. Hence, the naive estimation from the parameters of the ECAL segmentation ($\Delta \phi = 0.0174$ in the CMS [164] and $\Delta \phi = 0.025$ at the ATLAS [165, 166]) places a bound on $Y$ mass: $m_Y \lesssim (0.0174 \times 375)/2 = 3.26$ GeV so that two photons will hit the same ECAL segment and eventually be recorded as a single photon. Of course, the actual photon jet conversion could be much more complicated and gives a stronger bound. However, this is not a big issue as $m_Y$ stays well below this bound in the following analysis.

The other concern about this topology in the language of photon-jet is the proper decay length of the light (pseudo)scalar $Y$. This is because a very light but highly boosted particle could have an unexpectedly long decay length. If its decay takes place after passing through the ECAL detector, it cannot be technically detected. Quantitatively, the proper average decay length of such a highly boosted particle is given by

\[ L_{\text{decay}} = c\tau \gamma \approx \frac{375c}{\Gamma_Y m_Y} \tag{2.1} \]

in which $\tau = \Gamma^{-1}$ and $\gamma = \frac{E_Y}{m_Y} \approx \frac{p_Y}{m_Y}$ are employed. $\Gamma_Y$, the total decay width of $Y$, is generally proportional to $m_Y$ for the light $Y$ state. Hence, a lighter $m_Y$ leads to a longer
Proper decay length $\beta_m \Delta$.

**Figure 1**: Assuming that the typical $R_{\text{ECAL}} = 1.5$ m from the beam, the signal loss ratio in percentage versus the decay length ($L_{\text{decay}}$) of light particle $Y$. When the decay length of $Y$ exceeds 2 m, about half of the $Y$ generated would not be detected.

$L_{\text{decay}}$. Once the decay length is comparable with the size of detector radius ($R_{\text{ECAL}}$ in this case), the probability of $Y$ escaping the detection soars. The correlation of the proper decay length versus the percentage of signal loss is illustrated in Fig. 1.

We can see that for a decay length longer than $\sim 1$ m, less than 80% of the $Y_1$ decay events would happen inside (or before reaching) the ECAL and eventually be captured by the detector. In this case, a larger production rate is needed to compensate the significant event loss. However, such a long $L_{\text{decay}}$ will cause another problem. For a considerable portion of $Y$ which pass through the ECAL, they will decay within the HCAL and their decay products will be tagged as displaced jets. We expect that future experiments searching for resonances using the construction of displaced jets and/or photons [167–169] could place an upper bound on $L_{\text{decay}}$.

On the other hand, it was argued in [164] that, if $L_{\text{decay}}$ is close to $R_{\text{ECAL}}$, $Y$ would decay in a position near the ECAL layer, thereby, the two photons produced from a not-so-light $Y$ ($m_Y \gtrsim 3.26$ GeV) can still hit the same ECAL segment simultaneously. Nonetheless, the probability of $Y$ decay takes place drops off exponentially along the distances away from the beam. This infers a larger probability of $Y$ decay occurring close to the beam rather than near the ECAL layer, giving rise to a pair of photons distinguishable when reaching the ECAL layer. If this were true, one would have seen a significant amount of 3-photon or 4-photon signals.

In short, $L_{\text{decay}}(Y) \lesssim 1$ m and $m_Y \lesssim 3$ GeV are viable assumptions in our scenario in order to have sufficient di-photon signal detected in the ECAL.
Figure 2: The correlation of the predicted di-photon signal ratio (in-box numbers) to the size of the branching ratio of $X(750) \to 2\gamma + \,...$. Such decay could be either a direct or cascade decay. The gluon fusion production cross section (indicated in the unit of fb on the top axis) is obtained in the assumption that $X(750)$ resonance behaves like a doublet whose Yukawa couplings given in Type I (left) and Type II (right) models.

3 Gluon fusion production and di-photon decay

3.1 Doublet $X(750)$ and its production at the collider

Let us begin with a numerical estimate to assess the possibility of realizing this scenario in the context of the 2HDM with the inclusion of a singlet scalar. For simplicity, we assume that the $X(750)$ resonance behaves like the heavy CP-even Higgs such as the couplings and the production modes. Thus, the $X(750)$ is dominantly produced via gluon fusion or $b-$quark associated production (we focus on $\tan \beta \lesssim 4$ and thus do not consider this production mode in the present paper). It is shown in [3] that the gluon-fusion production cross section decrease as $\tan \beta$ grows in Type I, while this cross section minimizes at a modest value of $\tan \beta \sim 8$ in Type II. Hence, a maximal value of $\tan \beta$ that could yield the di-photon signal can be estimated by assuming the $X(750)$ decays inclusively into two photons with 100% branching ratio. As shown in Fig. 2, we evaluate the di-photon signal, $\sigma(gg \to X) \text{BR}(X \to \gamma\gamma)$, as a function of $\tan \beta$ at different levels of BR($X \to 2\gamma + \,...$).

As expected, the obtainable di-photon signal diminishes as $\tan \beta$ increases, while maintaining the level of BR($X \to 2\gamma + \,...$). (This is also true for $\tan \beta \lesssim 20$ in Type II.) This is the result of suppressing the $X(750)$ production at high $\tan \beta$ due to the fact the heavy Higgs coupling normalized to the SM value with the top quark, which dominantly mediates the gluon-fusion, is inversely proportional to $\tan \beta$. To offset this drop, the decay branching ratio BR($X \to 2\gamma + \,...$) has to be increased to acquire a larger signal rate. For instance, in order to yield a 10 fb signal, $t_\beta \leq 8(10)$ and $t_\beta \leq 6(7)$ are required in Type I (Type II) in the assumption of BR($X \to 2\gamma + \,...$) at 100% and more practically at 60%, respectively. In particular, such an upper limit on the $\tan \beta$ would be relaxed in Type II for a lower cross section of 6 fb. In both Type I and Type II, a relative large cross session ($\approx 700$ fb in the
of producing the $X(750)$ state is achieved when $t_β \lesssim 1$ as a result of the enhanced coupling to top-quark. Meanwhile, this will maximize the cross section in the $t\bar{t}$ and di-jet final states. Comparing with the current experiment bounds, $\sigma(pp \to X \to gg) \lesssim 10$ pb and $\sigma(pp \to X \to t\bar{t}) \lesssim 700$ fb \[170, 171\], we find that our model stays well below the di-jet bound while the di-top bound could be marginal in the case of $t_β \geq 1$. Thus, we shall limit $t_β \geq 1$ in the following analysis.

### 3.2 Singlet $Y$: scalar vs. pseudoscalar

Concerning the light state $Y$, there are two important requirements that must be satisfied to realize this scenario. These include the presence of decay into di-photon and the substantially large total width. To accomplish the first goal, unlike the situation in the NMSSM, this state cannot be a pseudoscalar in the present study. As explained in the Introduction, the CP-odd scalar from the doublet field cannot be too light, while the pseudoscalar arising from the newly introduced singlet field does not couple to SM particles in the assumption of CP-conservation and thus has no di-photon decay mode. On the other hand, requiring $Y$ a singlet scalar is not sufficient. Though $Y$ can decay to di-photon via charged Higgs loop, or to 4$f$ via two off-shell $H^*, A^*$ or $H^{\pm*}$, these decay channels are highly suppressed by kinematics. Tree-level decays to light quarks are absent due to its decoupling to the SM particles. The combination of these effects result in a extremely narrow decay width. The proper decay length for a singlet scalar is typically at the order of kilometer. Therefore, $Y$ state must be a scalar and also gains an amount of doublet composition from the mixing. The content to which the mixing is needed will be analyzed in Sec. 4.3. Meanwhile, the singlet pseudoscalar, if present, could serve as a dark matter candidate.

In addition, the mass of $Y$ is very crucial in determining the products of $Y$ decay. When $m_Y$ is $\mathcal{O}(1)$ GeV, its main decay products are jets. In this case, the decay to di-photon is only mediated via loop diagrams and its branching ratio $\lesssim 0.1 - 1\%$ (see an example of $m_Y = 5$ GeV in Fig. 3). In contrast, the notorious jet background is well suppressed in the decay for a sub-GeV particle. In particular, considering a $m_Y$ below the $s\bar{s}$ and $\mu^+\mu^-$ thresholds, it will only decay to $u\bar{u}, d\bar{d}$ and $e^+e^-$ at tree level. Notice that $m_Y < \Lambda_{\text{QCD}}$, so the outgoing quark pairs would not develop into di-jet but hadronize and cascade decay into di-photon. On the other hand, $BR(Y \to e^+e^-)$ is negligible compared to quark pair production due to the small electron Yukawa coupling. Therefore, one can expect the branching ratio of the decay to di-photon for such a light $Y$ is nearly $\sim 100\%$.

Of course, the situation becomes rather subtle when $m_Y$ exceeds the $\mu^+\mu^-$ threshold, $\approx 210$ MeV. Since the muon has a much larger Yukawa coupling than $u$ and $d$ quarks, there will be a considerable portion of $Y$ decaying to muon pairs. This seems to generate a heavy suppression in the branching ratio of the decay to di-photon. However, strong dynamics may have a large correction to the decay width of $Y \to \gamma\gamma$ when this decay is mediated via hadronic states. Above the $\mu^+\mu^-$ threshold, once $m_Y$ reaches the $\pi\pi$ threshold, $\approx 280$ MeV, $Y$ will hadronically decay into either $\pi^+\pi^-$ or $\pi^0\pi^0$ since $\pi$ is the lightest hadron. By virtue of isospin conservation, only about $1/3$ of outgoing pion pairs are $\pi^0\pi^0$ pair and cascade decay to di-photon for each $\pi^0$. Whereas, the remaining $2/3$ will be $\pi^+\pi^-$ and
eventually decay to $\mu^+\mu^- + 4\nu$, which can be detected by the muon chamber, leaving a large missing $p_T$.

To avoid the subtlety, we will consider the case throughout the paper in which $m_Y$ is below the di-muon threshold. In the numerical analysis unless specified, $m_Y = 150$ MeV is chosen and the assumption of $\text{BR}(Y \rightarrow \gamma\gamma) = 100\%$ is globally adopted.

4 Model I: real singlet

We first consider adding to the 2HDM a real scalar gauge singlet $S$. To eliminate the substantial FCNC, we assume a $Z_2$ symmetry under which $\Phi_2$ is odd. For the singlet $S$ we impose a $Z'_2$ symmetry under which $S$ is the only odd field. The gauge invariant and renormalizable Lagrangian for this model is

$$\mathcal{L}_{2HDMS} = \mathcal{L}_{2HDM} + \mathcal{L}_S$$ (4.1)

where the Lagrangian of CP-conserving but soft $Z_2$ breaking 2HDM can be found in [3, 162, 172] in which two Higgs-doublet fields are expanded as

$$\Phi_a = \left(\begin{array}{c} H_a^+ \\ \frac{1}{\sqrt{2}}(v_a + \rho_a + i\eta_a) \end{array}\right) \quad (a = 1, 2), \quad S = v_s + \chi$$ (4.2)

with the ratio of two vevs given by $\tan\beta \equiv v_2/v_1$.

While the $S$-associated part reads

$$\mathcal{L}_S = \frac{1}{2}m_s^2S^2 + \frac{1}{4!}\lambda_sS^4 + \kappa_1\Phi_1\Phi_1S^2 + \kappa_2\Phi_2\Phi_2S^2$$ (4.3)

In contrast to [172], in the present study we allow the $S$ to acquire a vev so that it will mix with the Higgs doublets. Three mass-squared parameters can be replaced by three vevs through the corresponding minimization conditions such that

$$m_{11}^2 = m_{12}^2\frac{v_2}{v_1} - \left(\frac{1}{2}\lambda_1v_1^2 + \frac{1}{2}\lambda_{345}v_2^2 + \kappa_1v_s^2\right)$$ (4.4)

$$m_{22}^2 = m_{12}^2\frac{v_1}{v_2} - \left(\frac{1}{2}\lambda_2v_1^2 + \frac{1}{2}\lambda_{345}v_2^2 + \kappa_2v_s^2\right)$$ (4.5)

$$m_s^2 = -\frac{1}{6}\lambda_sv_s^2 - \kappa_1v_1^2 - \kappa_2v_2^2$$ (4.6)

4.1 Mass eigenstates and spectrum

For the CP-even neutral states ($\rho_1, \rho_2, \chi$) in the $Z_2$ basis, the mass matrix can be written as

$$\mathcal{M}^2 = \begin{pmatrix}
      m_{12}^2\frac{v_2}{v_1} + \lambda_1v_1^2 & -m_{12}^2 + \lambda_{345}v_1v_2 & 2\kappa_1v_1v_s \\
      -m_{12}^2 + \lambda_{345}v_1v_2 & m_{12}^2\frac{v_1}{v_2} + \lambda_2v_2^2 & 2\kappa_2v_2v_s \\
      2\kappa_1v_1v_s & 2\kappa_2v_2v_s & m_s^2
    \end{pmatrix}$$ (4.7)

---

This model has been already constructed in earlier literatures, see an example [172] in which the singlet does not acquire a vev and also is $Z_2$ odd for the sake of having a dark matter candidate.
where \( m_\chi^2 = \frac{1}{2} \lambda_s v_s^2 \) given Eq. (4.6) has been used. In the presence of non-negligible off-diagonal elements \((v_s \neq 0)\) in the above mass matrix, \((\rho_1, \rho_2, \chi)\) are apparently not the mass eigenstates. We first rotate two doublet components (upper \(2 \times 2\) block) into the basis \((\hat{h}, H)\) via an angle \(\alpha\).

\[
\begin{pmatrix}
\hat{h} \\
H
\end{pmatrix} = \begin{pmatrix}
-s_\alpha & c_\alpha \\
c_\alpha & s_\alpha
\end{pmatrix} \begin{pmatrix}
\rho_1 \\
\rho_2
\end{pmatrix}
\]

(4.8)

In fact, they are mass eigenstates in the pure 2HDM but no longer true in the model we consider here due to the doublet-singlet mixture induced by \(v_s \neq 0\). This can be seen explicitly from the full \(3 \times 3\) mass matrix under the unitary rotation.

\[
\hat{M}^2 = \begin{pmatrix}
-s_\alpha & c_\alpha & 0 \\
c_\alpha & s_\alpha & 0 \\
0 & 0 & 1
\end{pmatrix} \begin{pmatrix}
M_1^2 & 0 & \Delta \\
0 & M_2^2 & D \\
\Delta & D & m_\chi^2
\end{pmatrix}
\]

(4.9)

where

\[
D = 2v_s(\kappa_1 c_\alpha c_\beta + \kappa_2 s_\alpha s_\beta)
\]

(4.10)

\[
\Delta = 2v_s(-\kappa_1 s_\alpha c_\beta + \kappa_2 c_\alpha s_\beta)
\]

(4.11)

Clearly, both off-diagonal elements, \(\Delta\) and \(D\), are not vanishing due to the presence of non-zero \(v_s\).

To fit the LHC Higgs data, we expect the SM Higgs \(h\) with \(m_h = 125\) GeV to be nearly pure doublet. This demands the mixing parameter \(\Delta\) very small. For simplicity we will take \(\Delta = 0\) in the following discussion, which then gives us

\[
\kappa_1 = \kappa_2 t_\beta / t_\alpha
\]

(4.12)

Whereas, we allow an arbitrary mixing between \(H\) and \(\chi\) for our purpose. Applying diagonalization between them, we find three resulting mass eigenstates which are formed by:

\[
\begin{pmatrix}
\hat{h} \\
Y_2 \\
Y_1
\end{pmatrix} = \begin{pmatrix}
1 & 0 & 0 \\
0 & c_\theta & s_\theta \\
0 & -s_\theta & c_\theta
\end{pmatrix} \begin{pmatrix}
\hat{h} \\
H \\
\chi
\end{pmatrix}
\]

(4.13)

where the states \((\hat{h}, H)\) defined in Eq. (4.8) are expressed in the \(Z_2\) basis. The mixing angle between \(H\) and \(\chi\) is given by

\[
s_{2\theta} = \frac{2D}{\sqrt{4D^2 + (m_\chi^2 - m_H^2)^2}} \quad , \quad c_{2\theta} = \frac{m_\chi^2 - m_H^2}{\sqrt{4D^2 + (m_\chi^2 - m_H^2)^2}}
\]

(4.14)

Alternatively, one can parameterize the mixing angle Eq. (4.14) as,

\[
c_\theta = \sqrt{\frac{m_Y^2 - m_\chi^2}{m_Y^2 - m_{Y_1}^2}} \quad , \quad s_\theta = \text{Sign}(\kappa_1 v c_\beta c_\alpha + \kappa_2 v s_\beta s_\alpha) \sqrt{\frac{m_Y^2 - m_{Y_1}^2}{m_Y^2 - m_{Y_1}^2}}
\]

(4.15)
Table 2: The couplings of scalars normalized to the SM values except for those to the charged Higgs.

| Scalar states | SM gauge bosons | SM fermions | charged Higgs |
|---------------|-----------------|-------------|---------------|
| $h$           | $s_{\beta-\alpha}$ | $C_f^H$       | $g_{hH^+H^+}$ |
| $Y_1$         | $-c_{\beta-\alpha}s_\theta$ | $-C_f^H s_\theta$ | $-s_\theta g_{H^+H^+} + c_\theta g_{\chi H^+H^+}$ |
| $Y_2$         | $c_{\beta-\alpha}c_\theta$ | $C_f^H c_\theta$ | $c_\theta g_{H^+H^+} + s_\theta g_{\chi H^+H^+}$ |

As discussed in the Introduction, to realize the idea of photon jet we consider the scenario where $Y_1$ is a sub-GeV singlet-like state and $Y_2$ a heavy doublet-like resonance at 750 GeV. To this end, we examine the lower $2 \times 2$ block in the mass matrix, Eq. (4.9). Since both $m_\chi$ and $D$ are proportional to the singlet VEV $v_s$, the mixing $s_\theta$ is small for most of the parameter space if $Y_1$ is required to be very light. (The exception occurs in the region where the potential stability is violated.) This suppression in $s_\theta$ is a fatal weakness for the phenomenology of the di-photon excess in this model as will be shown in Sec. 4.4.

Finally, we present the masses for three scalar mass eigenstates

$$m^2_h = m^2_{h_0}$$

$$m^2_{Y_{1,2}} = \frac{1}{2} \left[ m^2_H + m^2_\chi \pm \sqrt{4D^2 + (m^2_\chi - m^2_H)^2} \right]$$

where the expression of $m^2_h$ and $m^2_H$ for two CP-even Higgs states in the 2HDM can be found in [3]. $m_{Y_2}$ and $m_{Y_1}$ in Eq. (4.17) take + and − signs, respectively.

### 4.2 Higgs couplings and decays

Though the $\chi$ component does not directly couple to any SM fermions, both $Y_1$ and $Y_2$ that arise from the $H - \chi$ mixing couple to SM particles as well as the charged Higgs. Some relevant couplings for the three CP-even mass eigenstates, normalized to SM values, are listed in Table 2, where the couplings of $h, H$ to the charge Higgs are defined in the 2HDM context as

$$g_{hH^+H^+} = -\frac{1}{2v} \left( m^2_h + 2(m^2_H - \bar{m}^2) \right) s_{\beta-\alpha} + 2 \cot 2\beta (m^2_h - \bar{m}^2)c_{\beta-\alpha}$$

$$g_{H^+H^+} = -\frac{1}{2v} \left( m^2_H + 2(m^2_H - \bar{m}^2) \right) c_{\beta-\alpha} - 2 \cot 2\beta (m^2_H - \bar{m}^2)s_{\beta-\alpha}$$

$$g_{\chi H^+H^+} = 2v(s_\beta^2 + \kappa_2 c_\beta^2)$$

with $\bar{m}^2 = \frac{2m^2_1}{\kappa_2}$. It is noticeable that the coupling of the Higgs $h$ to gauge bosons is proportional to $\sin(\beta - \alpha)$ while that for the other two states $Y_1$ and $Y_2$, after mixing, display the $\cos(\beta - \alpha)$ dependence. This implies that most relations in the 2HDM, in particular the alignment limit, still holds in this extended model. To fit the 125 GeV Higgs data, for simplicity we take the alignment limit ($c_{\beta-\alpha} \to 0$), in which the $h$ has SM-like couplings while the $Y_1$...
and $Y_2$ decouple with gauge bosons. This reduces Eq. (4.12) to

$$\kappa_1 \sim -\kappa_2 \tan^2 \beta \quad (4.21)$$

Taken this relation, the coupling $hY_1Y_1$ vanishes at small mixing limit. Hence, the decay of SM Higgs $h$ to new light scalars such as $h \rightarrow Y_1Y_1 \rightarrow 4\gamma$, which is already constrained by the Higgs invisible decay search [173], is switched off automatically.

As for Yukawa couplings, they display an overall dependence on $c_\theta$ and $s_\theta$ for the $Y_2$ and $Y_1$ states, respectively. This means the mixing between them is crucial in determining their decays to the SM fermions. Other than the couplings listed in Table 2, the $Y_2Y_1$ coupling is most relevant to our discussion, which reads in the alignment limit

$$g_{Y_2Y_1} = 3c_\theta s_\theta^2 g_{HHH} + 2\kappa_2 v c_\theta s_\theta^2 s_\beta c_\beta (t_3^2 + 1) - \kappa_2 v s_\beta c_\beta s_\theta^2 (t_3^2 + 1) + \kappa_2 v s_\theta^3 (c_\beta^2 - s_\beta^2 t_3^2) + \frac{1}{2} c_\theta^2 s_\theta v_s (-4\kappa_2 c_\beta^2 + \lambda_s + 4\kappa_2 s_\beta^2 s_3^2) \quad (4.22)$$

In addition, $Y_2 \rightarrow hh$ should have been another important decay mode considering the fact that in the 2HDM context the heavy Higgs $H$ generally has a sizable decay branching ratio into a pair of SM-like Higgs $hh$. However, we examine that this coupling $g_{Y_2hh}$ is vanishing at the exact alignment limit [3]. Therefore, we do not consider this decay mode in the present analysis. It is also necessary to notice that the couplings of the singlet $\chi$ to massless Goldstone bosons are proportional to: $v_s (\kappa_1 c_3^2 + \kappa_2 s_3^2)$. As a result, nearly singlet-like $Y_1$ does not couple to the longitudinal modes of $W^\pm / Z$ in the limit $\Delta \rightarrow 0$.

### 4.3 The decay of $Y_1$

We now turn to study $Y_1$ decay. Recall that $Y_1$ is a mixture of CP-even heavy Higgs $H$ and real singlet $\chi$, so it has Yukawa couplings as a doublet (the fraction of which is described by $s_\theta$) and also couplings with scalars present in the Higgs sector. As a result, it can decay to SM light quarks. The presence of these tree level decays greatly increases the total width and in turns shorten the decay length $L_{\text{decay}}$ within the range of $R_{\text{ECAL}}$ scale. The influence of mixing $s_\theta$ on the decay length of $Y_1$ is illustrated in Fig. 3 for two choices of $m_{Y_1}$. In order to maximally enhance the decay width of $Y_1$, we consider the Yukawa patterns obeying the Type II model in the following numerical analysis.

First, as expected, $\text{BR}(Y_1 \rightarrow \gamma\gamma) \lesssim 0.1\%$ for $m_{Y_1} = 5$ GeV. This is the result of the presence of tree-level decay to light quarks. To better understand the $Y_1$ decay, it is useful to analyze the explicit form of $Y_1$ coupling to fermions, $g_{Y_1ff}$ (c.f. Table 2). This coupling is proportional to $s_\theta$, which depends monotonically on $D = -2v_s \kappa_2 (1 + \tan^2 \beta)$. Thus, as $s_\theta$, or essentially $\kappa_2$, increases, the total decay width of $Y_1$ grows up. This can be easily understood from the fact that $Y_1$ acquires more doublet component in the large mixing. The remaining factors $C_d^H \sim t_\beta$, $C_u^H \sim t_\beta^{-1}$ are specified at the exact alignment limit in the Type II model. In this case $Y_1 \rightarrow d\bar{d}$ decay is enhanced for $t_\beta > 1$. Of course, $Y_1 \rightarrow u\bar{u}$ is simultaneously suppressed but not so efficiently as an offset since $m_d/m_u \simeq 2$. Therefore, it could be expected that the total decay width increases as $t_\beta$ becomes large, which in turn leads to a shorter decay length. All these behaviors are clearly reflected in Fig. 3.
Figure 3: Decay length of $Y_1$ as a function of the doublet fraction in $Y_1$ for $m_{Y_1} = 200$ MeV by choosing various values of $\tan \beta$. Red, orange and green curves correspond to $\tan \beta = 2, 4, 6$, respectively.

4.4 Di-photon cross section

According to the preceding discussion, there are five independent parameters in the numerical analysis, which include $m_H$, $v_s$, $\lambda_s$, $\kappa_2$ and $t_\beta$. However, since $m_{Y_2} = 750$ GeV $\gg m_{Y_1}$, the mixing effect on the mass of the heavy eigenstate would be very small, allowing us to take $m_H = 750$ GeV as a good approximation. While $v_s$ can be determined by using Eq. (4.17) once the value of $m_{Y_1}$ is chosen. In the end, we are left with three free parameters $\lambda_s$, $\kappa_2$ and $t_\beta$.

To compute the di-photon cross section of our interest, we examine the decay modes of $Y_2$. As discussed in Sec. 4.3, we know that a moderate mixing (small $s_\theta$) is required to render the decay length of $Y_1$ sufficiently short (also see Fig. 3). In this case, $|c_\theta| \approx \pm 1$ and the leading contribution to $g_{Y_2 Y_1 Y_1}$ in Eq. (4.22) reads

$$g_{Y_2 Y_1 Y_1} = -\text{sign}(s_\theta)v_{C_\beta} s_\beta \kappa_2 (1 + t_\beta^2) \quad (4.23)$$

It implies that this coupling, or equivalently the decay $Y_2 \to Y_1 Y_1$ tends to be important as either $|\kappa_2|$ or $\tan \beta$ becomes large. However, one cannot achieve both two large because of the constraint from the stability, as you will see shortly. In the alignment limit, $Y_2$ decays mainly into $Y_1 Y_1$ and $t \bar{t}$, the latter one is sensitive to $t_\beta$ as in the 2HDM. In particular, when $t_\beta \lesssim 6$, $Y_2 \to t \bar{t}$ is a predominant decay channel.

The viability of this scenario is illustrated in Fig. 4. There, we vary two free parameters $t_\beta$ and $\kappa_2$, and draw the contours of the production cross section $\sigma(gg \to Y_2)\text{BR}(Y_2 \to Y_1 Y_1)$, assuming the BR($Y_1 \to \gamma \gamma$) = 100% for each $Y_1$ as argued in Sec. 3.2. The decay length of $Y_1$, $L_{\text{decay}}$, and the total width of $Y_2$ are also presented. In this figure, $m_{Y_1} = 150$ MeV is adopted as a typical value and $\lambda_s = 2\pi$ is chosen to make our model compatible with the unitarity conditions \cite{172}. 

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Figure 4: The di-photon cross section $\sigma(gg \rightarrow Y_2)BR(Y_2 \rightarrow Y_1Y_1)$ (in the unit of fb) produced by the 750 GeV resonance $Y_2$ under the assumption that $BR(Y_1 \rightarrow \gamma\gamma) = 100\%$, see the contours with white boxes. Only the region covered by blue scattering points is allowed by the potential stability condition for $\lambda_s = 2\pi$. The dashed green and red lines show the contours of the proper decay length of $Y_1$ (in the unit of meter) and the total decay width of $Y_2$ (in the unit of GeV), respectively.

Clearly, there exists parameter space which could yield the 750 GeV di-photon signal, comparable to those observed at the LHC. The yellow shaded strip indicates the $\sigma(gg \rightarrow Y_2)BR(Y_2 \rightarrow Y_1Y_1)$ within 4-10 fb. However, it is not necessarily acceptable because the allowed value of $|\kappa_2|$ for a fixed $\tan\beta$ is constrained by the vacuum stability, which has been shown to play the most important role in eliminating the parameter space [172]. In essence, the upper bound on $|\kappa_2|$ is determined by $t_\beta, \lambda_s$ as well as quartic couplings $\lambda_{1-5}$ in the Higgs sector. A simple derivation to obtain this bound can be found in the Appendix. The allowed region is sketched by blue scattering points in the figure and this band displays a tendency of compression as $\tan\beta$ increases. Though it is still possible to find a value of $\kappa_2$ for $\tan\beta \lesssim 4$ producing the desired signal rate which can fit the data, the decay length of $Y_1$ for such value of $\kappa_2$ within the stability bound is incredibly long due to the insufficient mixing. Therefore, we conclude that this model containing a real singlet is difficult to simultaneously yield the di-photon signal comparable to the observed data and also achieve a reasonable proper decay length.

5 Model II: complex singlet embedding a pseudoscalar dark matter

To remedy the problem of too long decay length present in the real singlet model, we are now introducing a complex scalar gauge singlet field $S$ to the 2HDM in this section.
The $S$-associated part is given by
\[
\mathcal{L}_S = \frac{1}{2}m_0^2S^2 + \frac{1}{4!}\lambda_S S^4 + \kappa_1 \Phi_1^\dagger \Phi_1 S^4 S + \kappa_2 \Phi_2^\dagger \Phi_2 S^4 S + \omega_1 \Phi_1^\dagger (S^\dagger + S) + \omega_2 \Phi_2^\dagger (S^\dagger + S) \tag{5.1}
\]
Here $\omega_1$ and $\omega_2$ are dimensionful parameters. Unlike the real singlet model discussed in Sec. 4, we do not impose the $Z'_{2}$ on the complex field $S$ so that the linear terms of $S$ are present in the above Lagrangian. The singlet scalar can be expanded as
\[
S = \chi_S + i\chi_A \tag{5.2}
\]
We stress that the complex singlet field $S$ cannot acquire a VEV: $\langle S \rangle = 0$, otherwise the CP-odd mode $\chi_A$ would be a massless Goldstone boson. Consequently, the minimization conditions in the scalar potential takes the same form as Eqs. (4.4) and (4.5) with the elimination of the $v_s$ term, while the one with respect to $v_s$ (i.e. Eq. (4.6)) is absent.

After EWSB, this complex singlet model has an additional CP-odd state $\chi_A$, that could be a candidate for dark matter. However, among the three CP-even states there are many similarities between Model I and Model II considered in this paper even if the $S$ in Model II does not acquire a vev. It is useful to comment that the single $\chi_S$ interactions with two doublets appearing in Eq. (5.1) are also present in Model I when the real field $S$ gets vev. This implies that the coupling $g_{\chi_S H^+ H^+}$ can be obtained from Eq. (4.20) with a simple substitution of $\kappa_a v_s$ by $\omega_a$. This replacement is also true in the $(\rho_1, \rho_2, \chi_S)$ mixing. To avoid redundancy, we shall transmit the results which have been derived in the real singlet model to the complex singlet model. A particularly important exception is the $Y_2 Y_1 Y_1$ coupling. In the present model, this coupling depends on both $\omega_a$ and $\kappa_a$ ($a = 1, 2$). As a result, such an interplay provides us the possibility of simultaneously ensuring the potential stability which is crucially determined by $\kappa_a$ and achieving a desired doublet-singlet compound for the light scalar through an essential increment on the $\omega_2$. In the end, the difficulty of too-long decay length can be overcome.

5.1 Spectrum and couplings

Without the $Z'_{2}$ protection, $S$ can singly couple to the doublets. This results in the doublet-singlet mixing which occurs only when the real singlet gets vev as discussed in the previous section. As already argued, the mass matrix in the $(\rho_1, \rho_2, \chi_S)$ basis can be obtained from Eq. (4.7) by replacing $k_a v_s$ with $\omega_a$ and $m_\chi^2 = m_0^2 + \kappa_1 v_1^2 + \kappa_2 v_2^2$ in this model. By virtue of this similarity, one can simply follow the procedures described in Sec. 4 to derive the mixing parameters $\Delta$, $D$, the composition of the resulting mass eigenstates $h, Y_1, Y_2$ and their mass spectrum. The corresponding results are analogous to Eqs. (4.10), (4.11), Eq. (4.13) and Eqs. (4.16), (4.17) with the only substitution of $k_a v_s \rightarrow \omega_a$. This implies that in this scenario the mixing between CP-even fields as described by $s_\theta$ depends on $\omega_2$ as opposed to $\kappa_2$.

In accordance with the LHC data for 125 GeV Higgs, we also employ the alignment limit and require the zero singlet fraction for the $h$ as discussed in Sec. 4. In this context,
$D = -2\nu \omega_2 t_\beta$ and $\Delta = 0$ because $\omega_1 = -\omega_2 t_\beta^2$. In contrast to the CP-even fields, the mixing in the CP-odd sector is absent in this scenario and hence $\chi_A$ itself is a mass eigenstate with mass $m_{\chi_A} = m_\chi$. Once it is light enough, both the SM Higgs $h$ and the 750 GeV resonance state $Y_2$ in the model could decay into a pair of them. The $h$ to $\chi_A \chi_A$ decay, if kinematically allowed, can be switched off by the vanishing relevant coupling $g_{h\chi_A \chi_A}$. This actually gives rise to Eq. (4.21), $\kappa_1 = -\kappa_2 t_\beta^2$, under which the couplings given in Table 2 are maintained except the substitution of $g_{\chi^1 H^+ H^-}$ by

$$g_{\chi^1 S H^+ H^-} = 2\omega_2 s_\beta^2 + 2\omega_2 c_\beta^2 = 2\omega_2 (1 - t_\beta^2)$$

(5.3)

In addition, similar to the real singlet case discussed in Sec. 4, the decay channel $h \to Y_1 Y_1$ is also closed by Eq. (4.21) in this scenario.

### 5.2 $Y_2$ decay

In addition to the main decay channels (to $Y_1 Y_1$ and $tt$), $Y_2$ in this model can invisibly decays into $\chi_A \chi_A$. Of the three relevant couplings, $g_{Y_2 hh}$ is given in Table 2, while the other two can be easily obtained at the exact alignment limit $c_{\beta - \alpha} = 0$.

$$g_{Y_2 Y_1 Y_1} = 3c_\beta s_\beta^2 g_{HHH} - \kappa_2 t_\beta (c_\theta^2 - 2s_\theta^2)c_\beta v + s_\theta (s_\theta^2 - 2c_\theta^2)(1 - t_\beta^2)\omega_2$$

(5.4)

$$g_{Y_2 \chi_A \chi_A} = (\kappa_1 - \kappa_2)c_\beta s_\beta c_{\alpha} v = -\kappa_2 t_\beta c_{\alpha} v$$

(5.5)

It should be noted that the heavy Higgs $H$ in the 2HDM context generally has a sizable decay branching ratio into a pair of SM-like Higgs $h$. However, we examine that this coupling $g_{Y_2 hh}$ is vanishing at the exact alignment limit [3]. Therefore, we do not consider this decay mode in the present analysis. For illustration, we display the branching ratios of $Y_2$ decay in Fig. 5 by taking $\tan \beta = 2, 3, 4$ (from left to right) in the range that could yield an observed di-photon cross session as we will see. In each graph, branching ratio curves are drawn in different colors corresponding to $\kappa_2 = 0$ and the maximal value of $|\kappa_2|$ such that the stability condition is obeyed for each $\tan \beta$. As seen from Fig. 5, in the small mixing case when $\omega_2$ is small, $tt$ channel dominates the decay of $Y_2$ for small $\tan \beta$ as long as $\kappa_2$ stays within the stability bound, whereas $BR(Y_2 \to Y_1 Y_1)$ and $BR(Y_2 \to \chi_A \chi_A)$ are not substantial. Reversely, increasing $\omega_2$ will invoke a larger mixing, which leads to the reduction of the doublet fraction in the $Y_2$, while $Y_1$ gains more doublet component. This results in a quick grow on the coupling $g_{Y_2 Y_1 Y_1}$, but little change on the coupling $g_{Y_2 hh}$. Consequently, $BR(Y_2 \to tt)$ drops drastically while $BR(Y_2 \to Y_1 Y_1)$ becomes important as the mixing increases. Thus, large $\omega_2$ is favored in order to accomplish a sufficiently large $BR(Y_2 \to Y_1 Y_1)$ for our purpose. In the plots we also observe that $BR(Y_2 \to Y_1 Y_1)$ vanishes at a certain value of $\omega_2$ when $\kappa_2 < 0$. This is the consequence of the interplay between $\kappa_2$ and $\omega_2$ terms in the coupling $g_{Y_2 Y_1 Y_1}$ given in Eq. (5.4). Besides, the particularly presented decay $Y_2$ into $\chi_A \chi_A$ is always sub-dominant. This can be attributed to the smallness of $\kappa_2$ as demanded by the stability.
Figure 5: The \( \omega_2 \) dependence of \( \text{BR}(Y_2 \rightarrow Y_1 Y_1) \) (solid), \( \text{BR}(Y_2 \rightarrow t\bar{t}) \) (dashed) and \( \text{BR}(Y_2 \rightarrow \chi A \chi A) \) (dotted) are indicated. In each graph, the branching ratio curves are drawn corresponding to \( \kappa_2 = 0 \) (purple) and the minimal/maximal (blue/red) value of \( \kappa_2 \) (which are given on top of each graph) such that the stability condition is obeyed for each \( \tan \beta \).

5.3 Phenomenology of 750 GeV state

We begin our discussion for 750 GeV state by presenting in Fig. 6 the contours of cross section \( \sigma(gg \rightarrow Y_2 \rightarrow Y_1 Y_1) \). For illustration, we take \( m_{Y_1} = 150 \text{ MeV} \) and \( \lambda_s = 2\pi \) \(^4\). The four graphs are produced by choosing \( \omega_2 = 50, 100, 200, 400 \text{ GeV} \) in sequence. The yellow shaded band corresponds to \( \sigma(gg \rightarrow Y_2 \rightarrow Y_1 Y_1) \) within 4-10 fb that could fit the ALTAS+CMS data. In the figure, the decay length of \( Y_1 \) and the total decay width of \( Y_2 \) are also shown in green and red dashed lines, respectively. The contour numbers are uniformly indicated in the corresponding colored boxes. In addition, we examine the stability condition of this model and mark the allowed region which is covered by blue scattering points. The trapezoid shape indicates that a smaller value of \( \tan \beta \) is able to accommodate a larger \( |\kappa_2| \). This can be understood from the fact that \( \kappa_1 = -\kappa_2 \tan \beta \) is employed and the stability condition essentially places upper bounds on \( |\kappa_1| \) and \( |\kappa_2| \) as well as their ratio \([172]\).

The presence of blue scattering points in the yellow shaded band tells us that this complex singlet model could easily yield the observed cross section \( \sigma(gg \rightarrow Y_2 \rightarrow Y_1 Y_1) \), while obeying the stability condition. Recall that \( \text{BR}(Y_1 \rightarrow \gamma\gamma) = 100\% \) is assumed in the estimate. Next, we examine the decay length of \( Y_1 \). This figure shows the decay length of \( Y_1 \) (green dashed line) has no dependence on \( \kappa_2 \) but is very sensitive to \( \tan \beta \) and \( \omega_2 \). Particularly, the decay length of \( Y_1 \) increases as \( \tan \beta \) goes down.

\(^4\)Here we adopt the somewhat more conservative value of \( \lambda_s = 2\pi \) so that our model would remain valid to at least a moderately higher scale before additional new physics would need to be included to obtain a theory valid at all energy scales. The impact of \( \lambda_s \) value on the stability was discussed in Ref. [172] in detail.
Figure 6: The di-photon cross section $\sigma(gg \rightarrow Y_2) BR(Y_2 \rightarrow Y_1 Y_1)$ produced by the 750 GeV resonance $Y_2$ under the assumption that $BR(Y_1 \rightarrow \gamma \gamma) = 100\%$, see the contours with white box. $\omega_2$ is chosen different values specified on the top of each graph. Only the region which is covered by blue (cyan) scattering points for $\lambda_s = 2 \pi (0.2)$ is allowed by the potential stability condition. The dashed lines with green and red box label the contours of the proper decay length of the light singlet $Y_1$ and the total decay width of $Y_2$, respectively.

analogous to the real singlet model presented in Sec. 4.3. Apparently, the requirement that the decay length $L_{\text{decay}} \lesssim 1$ m has significant impact on eliminating the small $\tan \beta$ region. Another important measurement that characterizes the potential 750 GeV di-photon resonance is its total width (red dashed lines), which varies from a few to tens of GeV seen from Fig. 6. Thus, this could be used as a critical signature in examining this scenario or determining the model parameter if confirmed. Two interesting observations regarding the cross section are placed in order. First, the cross section contours display an asymmetry with respect to $\kappa_2 = 0$. This is actually a result of Eq. (5.4). Second, the
Figure 7: The di-photon cross section $\sigma(gg \to Y_2)BR(Y_2 \to Y_1Y_1)$ in the $\omega_2 - \tan \beta$ plane with $\kappa_2 = 0.1, -0.1, 0.01$, under the assumption that $BR(Y_1 \to \gamma\gamma) = 100\%$. The legend is the same as in Fig. 6 except the blue/cyan horizontal dashed line which gives the approximate upper bound on $\tan \beta$ coming from the stability.

magnitude of cross section becomes less sensitive to $\kappa_2$ as $\omega_2$ increases, as seen from the fact that the contours keep tilting to the left.

Alternatively, the result can be projected onto $\kappa_2 - \omega_2$ plane. We show two examples of $\kappa_2 = \pm 0.1$ in Fig. 7. One can easily gain the additional information from this figure regarding the $\omega_2$ dependence. As $|\omega_2|$ goes large, the decay width of $Y_1$ increases because $Y_1$ is composed of more doublet fraction. This eventually results in a shorter decay length. On the other hand, for $\tan \beta \leq 2$, the total width of $Y_2$ is marginally sensitive to $\omega_2$. This is because $t\bar{t}$ channel dominates the decay of $Y_2$ in the small $\tan \beta$ region as seen in Fig. 5. Finally, it is important to mention that the minimal width of $Y_2$ is $\gtrsim 1$ GeV in this scenario. This is in contrast to many models where the width of 750 GeV state is of sub-GeV scale.
Figure 8: The allowed range on $m_A$ and $m_{H^\pm}$. The gray points are compatible with both theoretical constraints and EWPO within the ±3σ range. The green dashed line here indicates the mass relation $m_A = 750 - m_Z$. Thus, $Y_2$ cannot decay to $AZ$ on-shell on the right hand side of this line. Similarly, the orange and blue dashed line represent $m_A = 750 + m_Z$ and $m_A = m_{H^\pm} + m_W$ (The result is generically the same as in both Type I and Type II models.)

6 Discovery prospects at the collider

6.1 Searching for $Y_2(750)$ in other channels

Aside from into the $Y_1Y_1$, the $Y_2$ state of 750 GeV largely decays into $f\bar{f}$ final state. This has been shown in Sec. 5.2. The cross section of producing a $Y_2$ decaying to the $t\bar{t}$ or $b\bar{b}$ final states could be as large as hundreds of fb. Hence, they could be possible channels to search for $Y_2$ at the future run of the LHC. Compared to the tree-level $f\bar{f}$ decays, the branching ratio of the loop-induced decay $Y_2 \rightarrow gg$ is negligible. This channel also suffers from a large QCD background at the LHC. Therefore, it is less likely to be a promising discovery channel.

6.2 Pseudo-scalar Higgs and charged Higgs

The range of the rest of two scalars, the pseudo-scalar Higgs mass $m_A$ and charged Higgs mass $m_{H^\pm}$, is also interesting. In principle, $m_A$ can be above or below either of two CP-even Higgs bosons, and even $m_A < m_h/2$ is possible and consistent with the data [162]. However, once the heavy Higgs mass ($m_{Y_2}$ in our model) is fixed, the allowed range of $m_A$ is limited. The interrelation between $m_A$ and $m_{H^\pm}$ for the case of $m_{Y_2} = 750$ GeV is illustrated in Fig. 8. There, we observe that $m_{H^\pm}, m_A$ are bounded in the 400 – 950 GeV. As expected, this result is identical to what is displayed in the 2HDM [3]. This is not surprising because the introduced complex singlet does not generate the mixing in the CP-odd states.
Due to kinematical suppression, $Y_2 \to AA/H^±H^±$ is always forbidden. The absence of these two decay modes is crucially important in the success of explaining the di-photon excess. Otherwise, they will eat a large amount of branching ratio of the $Y_2$ decay, so that the $Y_2 \to Y_1Y_1$ decay will be heavily suppressed. While $Y_2 \to ZA$ and/or $Y_2 \to W^±H^±$ decay is kinematically possible for $m_{H^±}, m_A \lesssim 670$ GeV, their contributions to the $Y_2$ decay are so small that could be neglected.

As regards for the $A$, $AZh$ coupling is vanishing in the exact alignment limit and thus the exotic decay $A \to Zh$ is not present. However, $A \to Y_2Z$ (right to the orange line) and $A \to W^±H^±$ (above the blue line) could be important in addition to the fermionic decay being the potential discovery channels for the $A$ including the $t\bar{t}$, $b\bar{b}$ and $\tau^−\tau^+$ final states. In general, the cross section of producing the $A$ via gluon-fusion is proportional to $\tan \beta$ and varies from few pb and 10 fb at the 13 TeV, depending on the exact mass. Since the decay into $t\bar{t}$ dominates for moderate $\tan \beta$ we consider here, the cross section $gg \to A \to b\bar{b}$ can at most reach $\sim 10$ fb, an order which hardly enables them to compete with the large QCD background.

7 Conclusions

First, we find that a pure doublet state at 750 GeV generally has a very limited branching ratio for the loop-induced di-photon decay when the tree-level decay to $t\bar{t}$ is present. As a result, the process of $gg \to H/A \to \gamma \gamma$ with $CP$-even or $CP$-odd Higgs $H/A$ being identified as the 750 GeV resonance in the minimal version of 2HDM cannot reproduce the di-photon signal that is comparable to the observed level.

Other than through direct decay, we alternatively consider the di-photon signal that arises from two bunches of collimated photon jets emitting from a pair of highly boosted particles in this work. In particular, we studied the impact of $m_{Y_1}$ and its proper decay length on collider phenomenology. It turns out that $m_{Y_1}$ should be $\lesssim 210$ MeV to efficiently produce a photon jet as narrow as it must be. However, such a light boosted state cannot be the $A$ in the 2HDM, once the $H$ is identified as the 750 GeV resonance. Therefore, we extend the 2HDM by adding a gauge singlet scalar field. Though the singlet scalar field generically mixes with two doublet fields, it is possible to accomplish a SM-like Higgs $h$ with another two mass eigenstates. They include one heavy doublet-like $Y_2$ and one light singlet-like $Y_1$ with mass of sub-GeV. The $Y_2(750)$ can decay to a pair of $Y_1$, each of which can further decay into two photon jets. On the other hand, the presence of this mixing, although rather small, is the key of controlling the proper decay length of $Y_1$.

Two specific models containing an extra real or complex singlet scalar are studied. In both models, the SM-like $h$ decay to $Y_1Y_1$ or invisible final states are switched off. For the real singlet model, the mixing between heavy Higgs and the singlet state is strongly constrained so that it is difficult to simultaneously yield the di-photon signal comparable to the observed data and also achieve a reasonable proper decay length. In contrast, in the absence of the $Z_2$ symmetry the complex model has linear singlet terms, which are irrelevant to the Higgs invisible decay but play an essential role in generating the mixing in the scalar sector. As a result, this model is easy to yield the 1-10 fb cross section in the
di-photon final state with a decay length of $\mathcal{O}(1)$ m for the $Y_1$, meanwhile parametrically predicts the width of 750 GeV resonance $\gtrsim 1$ GeV. In addition, the pseudoscalar component of the singlet is naturally stable and hence could be a dark matter candidate.

Finally, we have discussed the discovery prospects of other scalar states such as the $CP$–odd $A$ and the charged Higgs $H^\pm$ at the future LHC run.

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Appendix: Stability constraint

To derive the stability bound we use the inequalities of $\kappa_{1,2}$ in [172]. Given in Eq. (4.21), $\kappa_1$ and $\kappa_2$ are of the opposite signs: $\kappa_1 = -t_3^2 \kappa_2$. Employing this relation, the upper limit of $|\kappa_2|$ reads

$$|\kappa_2| \leq \text{Min}(A_1, A_2)$$

where $A_1 = \sqrt{\lambda_1 \lambda_s / 12 / t_3^2}$, $A_2 = \sqrt{\lambda_2 \lambda_s / 12}$. When $\lambda_3 < 0$, one should ensure

$$|\kappa_2| \leq \sqrt{\frac{A_1^2 A_2^2 - A_3^2}{A_1^2 + A_2^2 - 2 A_3}} \quad \text{if} \quad A_3 < A_1 A_2$$

$$|\kappa_2| = 0 \quad \text{if} \quad A_3 \geq A_1 A_2$$

where $A_3 = -\lambda_3 \lambda_s / (12 t_3^2)$. Additionally, if $\lambda_3 + \lambda_4 - |\lambda_5| < 0$ is satisfied, one more constraint is required which can be obtained by replacing $A_3$ by $A_4$ with

$$A_4 = -(\lambda_3 + \lambda_4 - |\lambda_5|) \lambda_s / (12 t_3^2).$$

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