Di-Photon excess in the 2HDM: hastening towards the instability and the non-perturbative regime

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Abstract

We challenge the interpretation of the di-photon excess recently observed by both ATLAS and CMS in a two Higgs doublet framework. Due to the large enhancement necessary to obtain the observed di-photon signal, a large number of colored and charged vector-like fermions are called for. We find that even before the hypercharge gauge coupling becomes non perturbative, the one loop effects of these fermions abruptly drive the scalar potential to instability.
1 Introduction

Recently, ATLAS [1] and CMS [2] have reported excesses in the di-photon channel using Run 2 data at $\sqrt{s} = 13$ TeV center of mass energy. The ATLAS collaboration, with an integrated luminosity of $3.2 \text{ fb}^{-1}$, has found an excess at a di-photon invariant mass of $\approx 750$ GeV, with a local significance of $3.9\sigma$ ($2.3\sigma$ after the look-elsewhere effect). Using 2.6 fb$^{-1}$ of data, CMS on its turn has observed an excess peaking at an invariant mass of 760 GeV, with a local significance of $2.6\sigma$ ($1.2\sigma$ global). Although there is a very mild preference for a resonance width of about 45 GeV, the data is yet too insipid to support any claim in this direction.

These intriguing results have been intensively investigated by the community [3]. A simple possibility is that the signal comes from the decay of a spin 0 or 2 (by virtue of the Landau-Yang theorem) resonance decaying into a two photons final state. If the resonance participate in the breaking of electroweak (EW) symmetry, a two Higgs doublet model (2HDM) is one of the simplest scenario to focus on [4]. In this scenario, the resonance can be identified with the heaviest of the CP even scalars (H), with the CP-odd state (A), or even with a superposition of both, in case their masses are degenerate within $O(10 \text{ GeV})$.

Nevertheless, the fact that the decay of scalars to di-photons is loop induced poses many difficulties in building up a compelling model to explain the data. As nothing beyond ordinary was observed consistent with a 750 GeV resonance decaying to other channels, like diboson or $t\bar{t}$, the scalars couplings to standard model (SM) fermions and massive gauge boson need to be quite suppressed, which in turn also suppress the production via gluon fusion and the decay to photons. Therefore, to explain the ATLAS and CMS excesses additional field content is called for. The effective couplings of the resonance to photons and gluons can be enhanced by introducing new vector-like quarks or leptons. Still, to obtain enough enhancement it is necessary to have a large number of VL fermions, electrically charged and possibly colored, with sizeable couplings to the heavy scalars. In this work we investigate the consequences of such profusion of particles for the stability of the EW vacuum, and the evolution of the gauge couplings at high scales.

The paper is organized as follow: in Section 2 we introduce the model and we show the region of the parameter space where the excess can be accommodated. In Section 3 we perform an analysis of the Renormalization Group Equation (RGE) evolution of the relevant couplings. Finally, we conclude in Sec. 4.

2 2HDM and vector-like fermions

The most general 2HDM potential compatible with the gauge symmetries of the Standard Model (SM) is the following (see e.g. [5]):

\[
V(\Phi_1, \Phi_2) = m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - m_{12}^2 (\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1) + \frac{\lambda_1}{2} (\Phi_1^\dagger \Phi_1)^2 + \frac{\lambda_2}{2} (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 \Phi_1^\dagger \Phi_1 \Phi_2^\dagger \Phi_2 + \lambda_4 \Phi_1^\dagger \Phi_2 \Phi_2^\dagger \Phi_1 + \frac{\lambda_5}{2} [(\Phi_1^\dagger \Phi_2)^2 + (\Phi_2^\dagger \Phi_1)^2],
\]

with $\Phi_1$ and $\Phi_2$ two complex SU(2) doublets with hypercharge 1/2. Here, we have assumed a $Z_2$ symmetry ($\Phi_1 \rightarrow -\Phi_1$), except for the soft breaking term proportional to $m_{12}^2$. For each doublet we can define:
In the following we assume that the CP symmetry is respected by the scalar potential, so the vacuum-expectation values $v_1$ and $v_2$ are real numbers. Moreover $\sqrt{v_1^2 + v_2^2} = v_{SM} = 246.2$ GeV. This scalar theory contains five massive states: a charged scalar $H^\pm$, a CP odd state $A$ and two CP even particles $h, H$, the lightest of those ($h$) is identified with the 125 GeV Higgs boson. The mass eigenstates are obtained after performing the following rotations:

$$H = \rho_1 \sin \alpha + \rho_2 \cos \alpha,$$
$$h = \rho_1 \cos \alpha - \rho_2 \sin \alpha,$$
$$A = -\eta_1 \cos \beta + \eta_2 \sin \beta,$$
$$H^+ = -\phi_1^+ \cos \beta + \phi^+_2 \sin \beta,$$

where the angles $\beta$ and $\alpha$ are:

$$\tan \beta = \frac{v_1}{v_2}, \quad \tan 2\alpha = \frac{2(-m_{12}^2 + \lambda_{345} v_1 v_2)}{m_{12}^2 (v_1/v_2 - v_2/v_1) + \lambda_2 v_2^2 - \lambda_1 v_1^2},$$

and $\lambda_{345} = \lambda_3 + \lambda_4 + \lambda_5$.

The quartic couplings of the scalar potential can be expressed in terms of the masses of the scalars, $\tan \beta$ and $M^2 \equiv m_{12}^2 / \sin 2\beta$:

$$\lambda_1 = \frac{1}{v^2} \left( -\tan^2 \beta M^2 + \frac{\sin^2 \alpha}{\cos^2 \beta} m_H^2 + \frac{\cos^2 \alpha}{\cos^2 \beta} m_H^2 \right),$$
$$\lambda_2 = \frac{1}{v^2} \left( -\cot^2 \beta M^2 + \frac{\cos^2 \alpha}{\sin^2 \beta} m_h^2 + \frac{\sin^2 \alpha}{\sin^2 \beta} m_h^2 \right),$$
$$\lambda_3 = \frac{1}{v^2} \left( -M^2 + 2m_{H^\pm}^2 + \frac{\sin 2\alpha}{\sin 2\beta} (m_{H^\pm}^2 - m_h^2) \right),$$
$$\lambda_4 = \frac{1}{v^2} \left( M^2 + m_A^2 - 2m_{H^\pm}^2 \right),$$
$$\lambda_5 = \frac{1}{v^2} \left( M^2 - m_A^2 \right).$$

The potential is unbounded from below if the following conditions are fulfilled:

$$\lambda_{1,2} > 0, \quad \lambda_3 > -(\lambda_1 \lambda_2)^{1/2} \quad \lambda_3 + \lambda_4 - |\lambda_5| > -(\lambda_1 \lambda_2)^{1/2}. \quad (11)$$

The unitarity constraints can be found in ref. [5]. As we will see later, the impact of the VL fermions will be so large that the precise expressions for the unitarity bound will not matter.

Concerning the couplings of the SM fermions to the physical scalars, different configurations are possible. Here we consider the type I 2HDM, a scenario where all the SM fermions couple only to one doublet ($\Phi_1$ by convention). Notice that the $Z_2$ symmetry of the potential suits this case very well, as it can be the reason of why SM fermions do not couple to $\Phi_2$. In this scenario, the couplings between fermions, gauge bosons and the
Scalars, after rotating to the physical basis, are given by

\[ y_f^h = y_{fSM} \left( \sin(\beta - \alpha) + \frac{\cos(\beta - \alpha)}{\tan \beta} \right), \]

(12)

\[ y_f^H = y_{fSM} \left( \cos(\beta - \alpha) - \frac{\sin(\beta - \alpha)}{\tan \beta} \right), \]

(13)

\[ y_f^A = \pm y_{fSM} \frac{1}{\tan \beta}, \]

(14)

\[ g_{HV} = 2 \sin(\beta - \alpha) \frac{m_v^2}{v}, \]

(15)

\[ g_{HV} = 2 \cos(\beta - \alpha) \frac{m_v^2}{v}, \]

(16)

where \( y_{f,A,H,SM} \) are the Yukawa coupling of the fermion \( f \) to the scalars \( h, A, H \), and the standard model Yukawa, respectively. The \( \pm \) sign in \( y_A \) applies to up and down quarks, respectively. In order to have SM-like couplings for the lightest scalar, we must invoke the so-called “alignment limit”, \( \beta - \alpha \sim \pi/2 \). As can be seen, in this limit the couplings of all the heavy scalars are suppressed.

As explained in the Introduction, additional matter is necessary to reproduce the di-photon excess. Following the recent literature \[4\], we introduce new vector-like (VL) quarks. The minimal scenario includes one VL SU(2) doublet (Q) and one SU(2) singlet (D) with an appropriate hypercharge assignment. We take the VL in the fundamental representation of SU(3). The Lagrangian includes:

\[ \mathcal{L} \supset y_i^Q \bar{Q}_R \Phi_i D_L + y_i^D \bar{Q}_L \Phi_i D_R. \]

(17)

For simplicity we take \( y_i^Q = y_i^D \equiv y_i \) and a common VL mass \( M_{VLQ} \) for these extra fermions.

### 2.1 Signal and constraints

To account for the ATLAS and CMS excess, the total cross-section in di-photons at 750 GeV should be \[1,2\]

\[ \sigma(pp \to \gamma \gamma)_{ATLAS} = (10 \pm 3) \text{ fb} \quad \sigma(pp \to \gamma \gamma)_{CMS} = (6 \pm 3) \text{ fb}. \]

(18)

We have computed the production cross-section of \( H \) and \( A \) in gluon fusion processes and we have included decays of those resonances into SM fermions and gauge bosons. We get:

\[ \sigma(H/A)_{\gamma \gamma} = \frac{C_{gg} \Gamma_{gg}}{M_S} \text{BR}(H/A \to \gamma \gamma) \]

(19)

The mass of the resonance \( M \) is fixed at 750 GeV and \( \Gamma_{gg} \) is the width in gluons. The partonic integral factor \( C_{gg} \) is obtained employing the set of pdf MSTW2008NLO \[6\] at a scale \( \mu = M \), and it reads \( C_{gg} = 2137 \) at \( \sqrt{s} = 13 \text{ TeV} \). In the following we assume that \( H \) and \( A \) are degenerate in mass, in such a way that they both contribute to the di-photon signal.

A viable scenario to explain the excess can be obtained with the following configuration:
the 2HDM should be close to the alignment limit, $|\cos(\beta - \alpha)| \lesssim 0.3$. This is necessary to reproduce the observed couplings of the mostly SM Higgs $h$ to gauge bosons, Eq. (12). The decays of $H$ to gauge bosons are suppressed, since they are proportional to $\sin(\beta - \alpha)$;

- In the 2HDM type I, the couplings of $H$ and $A$ to the SM fermions are universally suppressed by $1/\tan \beta$, compared to the SM couplings to the Higgs boson. Therefore, focusing on large $\tan \beta$, one can suppress the decays of $H$ and $A$ into SM fermions. Moreover, this condition is necessary to minimize the departure of the $h$ couplings to the SM values;

- In the case of large $\tan \beta$, the contribution to $h \rightarrow \gamma \gamma$ due to the VL quarks of Eq. (17) are proportional to $y_1^2$, while the contributions to $(H,A) \rightarrow \gamma \gamma$ are proportional to the product $y_1 y_2$ [4]. As such, in order to keep under control the deviations of the $h$ couplings to gluons and photons induced by the VL quarks, a hierarchy $y_1 < y_2$ should be imposed.

In Fig. 1 we show some benchmark cases. In the green area, $\sigma(H/A)_{\gamma \gamma} = (3 \div 10)$ fb and the di-photon excess can be reproduced. The regions on the left of the blue and black lines are excluded since the deviations of $h$ couplings to $\gamma \gamma$ and $gg$ are larger than 20% [7,8]. As can be noticed from the left panel of Fig. 1, the measurements of the $h$ couplings strongly constrain this model, and a consistent explanation of the di-photon excess requires a quite extreme configuration. Moreover, the direct searches for VLQ through the decay processes $T^{+5/3} \rightarrow W^+ t$, $T^{+2/3} \rightarrow b W^+$, $B^{-1/3} \rightarrow t W^-$ and $B^{-1/3} \rightarrow b h$ constrain the mass of these fermions to be above 800 GeV [9], 705 GeV [10], 800 GeV [11], and 846 GeV [12], respectively. [7]

For instance, we have explicitly checked that, in order to explain Eq. (18) without incurring into troubles with the direct bounds, for VL quarks with charge $Q = 5/3$ more than $N_f \gtrsim 5$ families are required. Such large multiplicity suggests that the new states could dramatically modify the evolutions of the couplings of the theory, through RGE effects. This is the focus of the next section.

3 RGE running and the fate of the EW vacuum

Let us now study the RGE evolution for the SM gauge couplings ($g_3, g', g$), the top Yukawa coupling ($y_t$) and the quartic couplings $\lambda_i$ of the 2HDM potential, Eq. (1). The corresponding 1-loop $\beta$ functions are reported in Appendix A. We have approximated the Coleman-Weinberg effective potential with the RGE-improved tree level potential, substituting the bare couplings with the corresponding RGE running quantities. In this way, the stability of the EW vacuum is still given by Eqs. (11).

Our results are shown in Fig. 2. To fix the boundary conditions for the quartic couplings we must choose a possible scalar spectrum. For simplicity, and in order to minimize the contribution to precision observables [5], we take all the heavy scalars to be degenerate at $M_H \simeq M_A \simeq M_{H^\pm} = 750$ GeV, and we fix $m_h = 125$ GeV. We expect our results to be valid in general, and not only for this choice of boundary conditions. Having fixed the boundary

\[\text{These bounds were derived assuming only one VL quark at a time. A large multiplicity } N_f \text{ of VL quarks with similar masses would naively raise the cross sections under consideration by } N_f, \text{ and therefore the bounds would typically be at the TeV scale.} \]
conditions, we can now evolve the couplings to high energy, fixing the parameters in the VL quark sector in such a way to explain the di-photon excess, Fig. 1. As an example, we fix a common VL quark bare mass at $M_{VLQ} = 1050$ GeV, with charge $Q = 5/3$ and couplings $y_1 = 0.25$ and $y_2 = 1.5$. We also fix the number of VL families to $N_f = 10$. As can be seen from Eq. (20), the hypercharge gauge coupling receives a large positive contribution which rapidly drives it to non perturbative values slightly above the VL threshold, as confirmed in Fig. 2. However, another important effect can be inferred from the RGE’s in Eq. (20). Indeed, $\lambda_2$ receives a large negative contribution proportional to $N_f |y_2|^4$, which drives it to negative values slightly above the VL quark threshold. Comparing with the bounds in Eq. (11), we see that this makes the potential unbounded from below already at the TeV scale, before the theory reaches a non perturbative regime.

These findings imply that, in this scenario, new physics should occur at the scale of the VL quarks to stabilize the scalar potential. Moreover the model becomes strongly coupled around these energies, at which new degrees of freedom should emerge and a new description of this theory is necessary. Although, we have not performed a complete scan of the parameter space of the model, we expect our conclusions to be generic for simple 2HDM interpretations of the diphoton excess. Indeed, i) generically large multiplicity and couplings of the VL quarks are needed to explain the excess, ii) this implies large corrections of the RGE evolutions of the gauge couplings, independently on the choice of $\lambda_i$, and iii) starting from a weakly-coupled theory some of the $\lambda_i$ unavoidably runs toward negative values, before being attracted to large positive values by the contribution of gauge couplings.

\footnote{Similar results have been obtained in the case of a scalar singlet \cite{13}.}
Figure 2. Left panel: RGE evolution of $g_3$, $g'$, $g$ and $y_t$ respectively in green, red, black and blue lines. Right panel: RGE evolutions of $\lambda_i$. For these plots, we have chosen $M_{VLQ} = 1050$ GeV, $y_1 = 0.25$, $y_2 = 1.5$, $Q = 5/3$ and $N_f = 10$. The masses of the scalar sector are $M_H = 750$ GeV, $M_A = 751$ GeV and $M_{H^\pm} = 750$ GeV.

4 Conclusions

Although still not statistically relevant, it is tantalizing to interpret the recently observed di-photon excess at 750 GeV in terms of extensions of the Standard Model. In this paper we have focused on the case of the 2HDM. As is well known [4], a 2HDM alone cannot reproduce the observed signal, calling thus for a more elaborate SM extension. The simplest possibility, already considered in the literature, is the one of a 2HDM augmented with vector-like quarks. In principle, such a framework could explain the excess through the decays of the heavy $H$ and/or $A$ scalars. Here, we have shown that a consistent realization of this scenario present severe difficulties. Our results can be summarized as follows: a large number of VL quarks is required in order to get the correct cross-section into two photons. The presence of such extra matter strongly affects the RG evolution of the hypercharge gauge coupling, which rapidly reaches non perturbative values a few hundreds of GeV above the VL quarks threshold. However, even before this happens, some of the Higgs quartic couplings become negative, destabilizing the scalar potential, as it becomes unbounded from below. We thus reach the following broad conclusions: new physics is required around the VL quarks threshold in order to stabilize the vacuum, and a strongly coupled description of the theory should emerge at slightly larger scales. Let us now comment on what can be inferred about possible UV completions of this scenario. A natural possibility would be to consider a composite 2HDM scenario [14–17] in which vector like quarks are naturally expected around the TeV scale. However, in order to stabilize the vacuum, additional bosonic degrees of freedom are needed around the same scale. Whether such additional particles should be part of an extended coset, or can be additional vectorial resonances, is still an open question that we may explore in a forthcoming work.

3In principle, a composite 2HDM could also emerge from a fermionic UV completion of the Minimal Composite Higgs Model [18].
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A 1-loop $\beta$ functions

We now list the relevant 1-loop beta functions of the model:

\[
\begin{align*}
\beta_g' &= \beta_{2}^{2\text{HDM}} + \frac{4N_f}{3}\frac{g^3}{16\pi^2}(9Q^2 + 6Q + 3/2), \\
\beta_g &= \beta_{2}^{2\text{HDM}} + 2N_f\frac{g^2}{16\pi^2}, \\
\beta_{g_3} &= \beta_{2}^{2\text{HDM}} + 2N_f\frac{g_3^2}{16\pi^2}, \\
\beta_y_t &= \beta_{2}^{2\text{HDM}} y_t, \\
\beta_{\lambda_1} &= \beta_{2}^{2\text{HDM}} - 3N_f\frac{|y_Q^1|^4 + |y_D^1|^4}{32\pi^2}, \\
\beta_{\lambda_2} &= \beta_{2}^{2\text{HDM}} - 3N_f\frac{|y_Q^1|^4 + |y_D^1|^4 |y_Q^2|^4}{32\pi^2}, \\
\beta_{\lambda_3} &= \beta_{2}^{2\text{HDM}} - 3N_f\frac{|y_Q^2|^4 |y_D^2|^4 + |y_D^2|^4 |y_Q^2|^4}{32\pi^2}, \\
\beta_{\lambda_4} &= \beta_{2}^{2\text{HDM}} - 3N_f\frac{|y_Q^2|^4 |y_D^1|^4 + |y_D^1|^4 |y_Q^2|^4}{32\pi^2}, \\
\beta_{\lambda_5} &= \beta_{2}^{2\text{HDM}} - 3N_f\frac{|y_Q^1|^4 |y_D^2|^4 + |y_D^2|^4 |y_Q^1|^4}{32\pi^2}, \\
\beta_{\lambda_6} &= \beta_{2}^{2\text{HDM}} - 3N_f\frac{|y_Q^2|^4 |y_D^1|^4 + |y_D^1|^4 |y_Q^2|^4}{32\pi^2}, \\
\beta_{\lambda_7} &= \beta_{2}^{2\text{HDM}} - 3N_f\frac{|y_Q^1|^4 |y_D^2|^4 + |y_D^2|^4 |y_Q^1|^4}{32\pi^2}.
\end{align*}
\]

The $\beta$ functions of the 2HDM, $\beta_{2}^{2\text{HDM}}$, can be found in [5]. We have included in the $\beta$ functions the dominant contribution of the VL quarks from the renormalization of the 4-point scalar vertexes, while we have disregarded the sub-leading contribution from the wave-function renormalization. Also note that $Z_2$ breaking terms proportional to $\lambda_6$ and $\lambda_7$ (see [5] for their definition) are radiatively generated.
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