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Two Higgs doublets to explain the excesses

$$pp \rightarrow \gamma\gamma(750 \text{ GeV}) \text{ and } h \rightarrow \tau^\pm \mu^\mp$$

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ABSTRACT: The two Higgs doublet model emerges as a minimal scenario in which to address, at the same time, the $\gamma\gamma$ excess at 750 GeV and the lepton flavour violating decay into $\tau^\pm \mu^\mp$ of the 125 GeV Higgs boson. The price to pay is additional matter to enhance the $\gamma\gamma$ rate, and a peculiar pattern for the lepton Yukawa couplings. We add TeV scale vector-like fermions and find parameter space consistent with both excesses, as well as with Higgs and electroweak precision observables.

KEYWORDS: Beyond Standard Model, Higgs Physics

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1 Introduction

The recently presented indications for a diphoton excess at ATLAS and CMS at an invariant mass of 750 GeV [1, 2] have caused much excitement in the high-energy phenomenology community [3–97]. At the same time, some hints of anomalies persist in the LHC run-I data. Notably there is a 2.4σ excess at CMS in the $h \rightarrow \tau^\pm\mu^\mp$ decay of the 125 GeV Standard Model-like Higgs boson h [98], corresponding to a best-fit branching ratio $\text{BR}(h \rightarrow \tau^\pm\mu^\mp) = 0.84^{+0.39}_{-0.37}\%$. This is compatible with the ATLAS analysis which finds $\text{BR}(h \rightarrow \tau^\pm\mu^\mp) = 0.77 \pm 0.62\%$ [99].

One of the simplest renormalisable models allowing for a $h \rightarrow \tau^\pm\mu^\mp$ branching ratio of the order of a percent is the two-Higgs-doublet model (2HDM) [100–103] with lepton flavour violating (LFV) Yukawa couplings. This model has been studied before [104–110], and, with renewed interest, after the CMS excess was announced [111–117]. The aim of the present paper is to study whether a simple 2HDM could explain both the diphoton excess and the LFV Higgs decay. The LHC excess has been studied in the 2HDM in [4–12], and several authors have combined it with other observables, such as the dark matter abundance [64–73], or B -physics anomalies [95–97]. In the 2HDM, candidates for the 750 GeV resonance are the heavier scalar Higgs h_2 and the pseudoscalar A . They can reproduce the observed cross-section times branching ratio into photons if they couple to heavy vector-like charged fermions, as has been discussed by several authors, e.g. [3–6, 13, 14]. The data [1, 2] suggest a broad resonance, which could be due to the exchange of nearly degenerate h_2 and A [15, 16].

We consider a CP-conserving 2HDM of type I in the decoupling limit [121], where the second doublet has a mass ~ 750 GeV. We work in the “Higgs basis”, where $H_1 = [0, (v + h_1)/\sqrt{2}]$ denotes the doublet which gets a vacuum expectation value (vev) $v \simeq 246$ GeV, and which has Standard Model Yukawa couplings. The second doublet $H_2 = [H^+, (h_2 + iA)/\sqrt{2}]$ does not couple to Standard Model fermions, except for a LFV Yukawa to $\tau^\pm\mu^\mp$. The physical Higgs bosons are the CP-even h and H , the pseudoscalar A and the charged

Higgses H^\pm . In the decoupling limit, the light h is almost aligned on the vev, making it the Standard-Model-like Higgs of 125 GeV. In section 2 we show how to enhance the H and A couplings to gluons and photons, by introducing new vector-like charged fermions, while respecting the bounds from electroweak precision tests and h signal strengths. We neglect the charged Higgs H^\pm because it contributes little to $H, A \rightarrow \gamma\gamma$. A small mixing with h_2 allows the LFV decay $h \rightarrow \tau^\pm \mu^\mp$, as discussed in section 3. In section 4, we demonstrate that one can accommodate the 750 GeV excess from the decays of H and A , in agreement with the LFV excess.

2 Two Higgs doublets coupling to extra matter

In this section we neglect the misalignment between the CP-even mass basis, and the ‘‘Higgs’’ basis, and focus on the Higgs couplings to new fermions. That is, we consider the limit where the Standard Model Higgs boson h is identified with h_1 , and the second Higgs doublet H_2 does not couple to the Standard Model, except for its gauge interactions. Therefore, $H = h_2$ and A cannot decay to Standard Model particles at tree-level. We include the misalignment in the following section, in order to obtain $h \rightarrow \tau^\pm \mu^\mp$.

In order for H and/or A to play the role of the 750 GeV resonance, we need to introduce a large effective coupling to $\gamma\gamma$, as well as to gg , in the hypothesis that the resonance is produced via gluon fusion. If the production is dominated by quarks, that have a smaller parton density function, one needs an even larger coupling to $\gamma\gamma$. We will discuss quantitatively these two possibilities in section 4.

To provide an explicit realization for such effective couplings, we introduce two vector-like fermions, that transform under $SU(3)_c \times SU(2)_w \times U(1)_Y$ as $D \sim (R_c, 2, Q + 1/2)$ and $S \sim (R_c, 1, Q)$, with interactions

$$-\mathcal{L} = M_D \overline{D}_L D_R + M_S \overline{S}_L S_R + \lambda_i^D \overline{D}_L H_i S_R + \lambda_i^S \overline{S}_L H_i^\dagger D_R + \text{h.c.} . \quad (2.1)$$

The state of electric charge $Q + 1$ has mass M_D and no Yukawa couplings. The two states of charge Q couple to the Higgs doublets, and their mass matrix is non-diagonal because of the vev of H_1 . We will denote the mass eigenvalues by $M_1 \leq M_2$. Note that, in order to induce the couplings $H\gamma\gamma$ and $A\gamma\gamma$ (and analogously for gluons), one needs either $\lambda_1^S \lambda_2^D \neq 0$ or $\lambda_1^D \lambda_2^S \neq 0$. This is illustrated diagrammatically in figure 1, and it amounts to generate the effective operator $H_2^\dagger H_1 F_{\mu\nu} F^{\mu\nu}$ via a fermion loop.

The couplings $\lambda_1^{D,S}$ are constrained as they contribute to the h -decays into $\gamma\gamma$ and gg , as well as by the precision electroweak parameters S and T . Indeed, vector-like charged fermions were employed in the past to explain the transient excess in the $h \rightarrow \gamma\gamma$ channel, see e.g. [123, 124]. A detailed analysis of the allowed parameter space is provided in ref. [125]. Here we describe two illustrative cases:

- (1) Degenerate fermion masses, $M_1 = M_2$. This is the case for $M_D = M_S$ and $\lambda_1^S = -\lambda_1^D$. Choosing $M_{1,2} = 1 \text{ TeV}$, $R_c = 3$ and $|Q| \leq 2$, one finds an upper bound $(\lambda_1^S v)/(\sqrt{2}M_1) \lesssim 0.25$. This bound is determined essentially by the T parameter, that is proportional to $N_c \equiv \dim(R_c)$ and independent from Q . When $|Q| > 2$ a

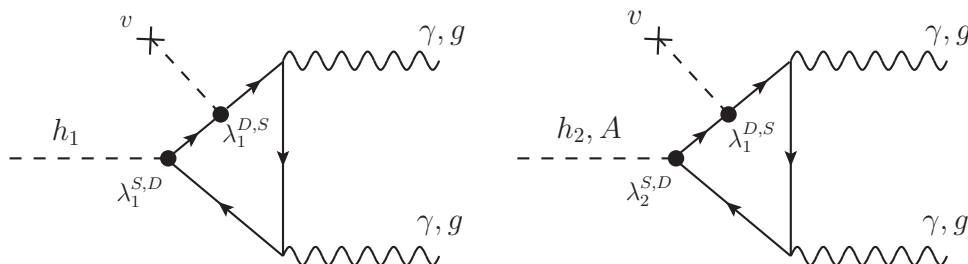


Figure 1. The contributions of two vector-like fermions, a doublet D and a singlet S , to the effective couplings between the Higgs doublets and two photons or gluons. As we work in the basis where the vev resides in H_1 , one finds that the h_1 couplings (left-hand side) are proportional to $\lambda_1^D \lambda_1^S$, while the h_2 and A couplings (right-hand side) are proportional to either $\lambda_1^S \lambda_2^D$ or $\lambda_1^D \lambda_2^S$.

stronger bound comes from the Higgs signal strengths. For $R_c = 8$ and $|Q| \leq 3$, one needs $(\lambda_1^S v)/(\sqrt{2}M_1) \lesssim 0.12$. In this case the bound comes from the hgg coupling.

- (2) One vanishing Yukawa coupling, e.g. $\lambda_1^D = 0$. This pattern strongly suppresses the correction to the couplings $h\gamma\gamma$ and hgg , because, in the limit of heavy fermions, they are proportional to $\lambda_1^D \lambda_1^S$. However, an upper bound on λ_1^S still exists, coming from the T parameter, $(\lambda_1^S v)/(\sqrt{2}M_1) \lesssim 0.35$ (0.25) for $R_c = 3$ (8) and $M_1 = 1$ TeV. Note that T does not depend on the hypercharge, therefore it turns out that one can take it very large, say $Q \sim 10$, without violating the constraints.

Let us now turn to the heavy Higgs doublet H_2 . Its couplings to the fermion mass eigenstates are easily derived [125] in terms of the parameters in the Lagrangian of eq. (2.1). Then, one can compute the decay width into two photons for the scalar H and the pseudoscalar A . The result is particularly compact in the limit $M_H \ll 2M_{1,2}$, since in this case the loop form factor $A_{1/2}[M_H^2/(4M_i^2)]$ is the same for both fermions in very good approximation, $A_{1/2}(0) = 4/3$. Similarly, for A we use the loop form factor $\tilde{A}_{1/2}(0) = 2$. Then, one obtains

$$\frac{\Gamma(H \rightarrow \gamma\gamma)}{M_H} = \frac{\alpha^2}{256\pi^3} \left| \frac{2vM_H}{3M_1M_2} N_c Q^2 (\lambda_2^D \lambda_1^S + \lambda_2^S \lambda_1^D) \right|^2, \tag{2.2}$$

$$\frac{\Gamma(A \rightarrow \gamma\gamma)}{M_A} = \frac{\alpha^2}{256\pi^3} \left| \frac{vM_A}{M_1M_2} N_c Q^2 (\lambda_2^D \lambda_1^S - \lambda_2^S \lambda_1^D) \right|^2, \tag{2.3}$$

In the same approximation, the widths into two gluons read

$$\frac{\Gamma(H \rightarrow gg)}{M_H} = \frac{\alpha_s^2}{32\pi^3} \left| \frac{2vM_H}{3M_1M_2} C(R_c) (\lambda_2^D \lambda_1^S + \lambda_2^S \lambda_1^D) \right|^2, \tag{2.4}$$

$$\frac{\Gamma(A \rightarrow gg)}{M_A} = \frac{\alpha_s^2}{32\pi^3} \left| \frac{vM_A}{M_1M_2} C(R_c) (\lambda_2^D \lambda_1^S - \lambda_2^S \lambda_1^D) \right|^2, \tag{2.5}$$

where $C(R_c)$ is the index of the color representation. Note that the ratio of H -rates over A -rates is given by a factor $(2|\lambda_2^D \lambda_1^S + \lambda_2^S \lambda_1^D|)^2 / (3|\lambda_2^D \lambda_1^S - \lambda_2^S \lambda_1^D|)^2$.

For definiteness, consider the case (2) described above, $\lambda_1^D = 0$, and take $M_H \simeq M_A \simeq 750$ GeV. Then, one obtains

$$\frac{\Gamma_{A\gamma\gamma}}{M_A} \simeq \frac{9}{4} \frac{\Gamma_{H\gamma\gamma}}{M_H} \simeq 1.4 \cdot 10^{-6} \left(\frac{1 \text{ TeV}}{M_2} \right)^2 \left(\frac{N_c}{3} \right)^2 \left(\frac{Q}{2} \right)^4 \left(\frac{\lambda_2^D}{3} \right)^2, \quad (2.6)$$

$$\frac{\Gamma_{A_{gg}}}{M_A} \simeq \frac{9}{4} \frac{\Gamma_{H_{gg}}}{M_H} \simeq 4.4 \cdot 10^{-6} \left(\frac{1 \text{ TeV}}{M_2} \right)^2 \left(\frac{C(R_c)}{1/2} \right)^2 \left(\frac{\lambda_2^D}{3} \right)^2, \quad (2.7)$$

where we chose $(\lambda_1^S v)/(\sqrt{2}M_1) \simeq 0.35$, that is the largest value allowed by the T parameter for $R_c = 3$. In the case of a colour octet, $N_c = 8$ and $C(R_c) = 3$, there is a slightly stronger upper bound, $(\lambda_1^S v)/(\sqrt{2}M_1) \simeq 0.25$: therefore, one gains a factor ~ 3 in $\gamma\gamma$ and a factor ~ 20 in gg .

Note that one can reproduce the same rates with smaller Yukawa couplings: taking N pairs of vector-like fermions, all with equal charges and coupling λ_2^D , the rates scale as $(N\lambda_2^D)^2$. From a theoretical point of view, it may be more justified to introduce several vector-like fermions, but with charges related to the Standard Model ones, such as one or more vector-like families, composed of t, b and τ partners. Adding over their contributions one could obtain a qualitative similar effect.

One should also remark that the heavy fermion loops also induce decays of H and A to $Z\gamma$, ZZ and WW , with width of the same order as (or slightly smaller than) for $\gamma\gamma$. However, the upper bounds from the 8 TeV LHC are weaker than the one on $\gamma\gamma$, as discussed e.g. in ref. [55]. Therefore, they are presently unconstraining. At run 2, the better perspective appears to be the observation of the $Z\gamma$ channel.

3 The $\tau^\pm \mu^\mp$ decay of the 125 GeV Higgs boson

Flavour-changing Higgs couplings are generic in the 2HDM, but their effects are not seen in low energy precision experiments searching for lepton or quark flavour change. So a discrete symmetry, which forbids flavour-changing Yukawa couplings, is usually imposed on the 2HDM. To allow for LFV h decays, without generating undesirable flavour-changing processes, we suppose that our 2HDM almost has a discrete symmetry: all the Standard Model fermions have the usual Yukawa couplings to H_1 (“type I” model), and the only two couplings of H_2 to Standard Model fermions are the $\mu\tau$ LFV ones,

$$\mathcal{L} = -\rho_{\tau\mu} \overline{L}_\tau H_2 \mu_R - \rho_{\mu\tau} \overline{L}_\mu H_2 \tau_R + \text{h.c.} \quad (3.1)$$

(see refs. [105, 122] for a more formal analysis). Recall that the diagonalisation of the fermion mass matrices diagonalises the Yukawa couplings of H_1 , which carries the vev. Therefore, the LFV couplings are attributed, by definition, to the doublet H_2 with zero vev. Note that eq. (3.1) amounts to assume that μ and τ numbers are not conserved, while electron number remains a good symmetry at the renormalizable level. Such symmetry has to be slightly broken to allow for viable neutrino masses. In general, this breaking will propagate radiatively to the H_2 Yukawa couplings, however the size of this effect can be sufficiently small, as it strongly depends on the specific neutrino mass model. In section 4

we will also consider a scenario where H_2 is produced from an additional Yukawa coupling to b quarks, that can be added without phenomenological problems.

The CP-even mass eigenstates h and H are misaligned with respect to the vev by an angle that is commonly parametrized as $\beta - \alpha$:

$$\begin{aligned} h &= \sin(\beta - \alpha)h_1 + \cos(\beta - \alpha)h_2, \\ H &= \cos(\beta - \alpha)h_1 - \sin(\beta - \alpha)h_2. \end{aligned} \quad (3.2)$$

In the decoupling limit [121, 122], $\sin(\beta - \alpha) \simeq 1$ and

$$\cos(\beta - \alpha) = -\frac{\Lambda_6 v^2}{M_H^2}, \quad (3.3)$$

where the Higgs potential contains a term $\Lambda_6 H_2^\dagger H_1 H_1^\dagger H_1 + \text{h.c.}$, in the basis where H_1 has no vev. The coupling of h to $\tau^\pm \mu^\mp$ is therefore proportional to $\cos(\beta - \alpha)\rho$, and one obtains

$$BR(h \rightarrow \tau^\pm \mu^\mp) \simeq \frac{m_h}{16\pi\Gamma_h} \cos^2(\beta - \alpha) \left(|\rho_{\tau\mu}|^2 + |\rho_{\mu\tau}|^2 \right). \quad (3.4)$$

The CMS best-fit is $BR(h \rightarrow \tau^\pm \mu^\mp) = 0.0084$ [98], which gives

$$\cos(\beta - \alpha) (|\rho_{\tau\mu}|^2 + |\rho_{\mu\tau}|^2)^{1/2} \simeq 0.0037, \quad (3.5)$$

where the width was taken at its Standard Model value, $\Gamma_h \simeq 4.1$ MeV.

In the 2HDM, the CMS excess in $h \rightarrow \tau^\pm \mu^\mp$ is consistent with the current upper bound $BR(\tau \rightarrow \mu\gamma) \leq 2.6 \times 10^{-7} BR(\tau \rightarrow \mu\nu\bar{\nu})$ [118, 119]. However, the extra fermions which enhance $H, A \rightarrow \gamma\gamma$ as in eqs. (2.2)–(2.3), will also enhance the rate for $\tau \rightarrow \mu\gamma$ [104]: if a neutral Higgs is exchanged between its $\gamma\gamma$ and $\bar{\tau}\mu$ vertices, and one of the photons connects to the lepton line, a diagram for $\tau \rightarrow \mu\gamma$ is obtained. Such diagrams with a top loop were calculated in the 2HDM in [120]. From their results, the combined contribution of H and A can be estimated, for $M_1 \simeq M_2$ and $\lambda_1^D = 0$, as

$$\frac{m_\tau}{v^2} A_L \simeq \frac{e\alpha}{128\pi^3} \frac{v}{\sqrt{2}M_1^2} N_c Q^2 \lambda_2^D \lambda_1^S \rho_{\tau\mu}^*, \quad (3.6)$$

where the experimental bound is $384\pi^2 (A_L^2 + A_R^2) \leq 2.6 \times 10^{-7}$. With the definition of Yukawa couplings given in eq. (2.1), it turns out that choosing a large λ_2^D (λ_2^S) leads to a destructive (constructive) interference among the H and A amplitudes. This was taken into account in eq. (3.6), where the difference in loop integral functions was chosen $\simeq 1/2$, as given in [120] for $M_1^2/M_H^2 \simeq 2$. A similar estimate can be made for A_R . We neglect the h contribution to $\tau \rightarrow \mu\gamma$, because its coupling to $\gamma\gamma$ is not enhanced, see scenario (2) in section 2. So the Babar-Belle bound on $\tau \rightarrow \mu\gamma$ could be satisfied for

$$\frac{N_c}{3} \left(\frac{Q}{2} \right)^2 \frac{\lambda_2^D}{3} \lambda_1^S \rho_{\tau\mu} \lesssim 0.07, \quad (3.7)$$

which sets a lower bound on $\cos(\beta - \alpha)$ when combined with eq. (3.5):

$$\cos^2(\beta - \alpha) \gtrsim 0.003 \left(\frac{N_c}{3} \right)^2 \left(\frac{Q}{2} \right)^4 \left(\frac{\lambda_2^D}{3} \right)^2 (\lambda_1^S)^2. \quad (3.8)$$

If the masses and couplings were purposefully tuned, it might be possible to suppress the $\tau \rightarrow \mu\gamma$ amplitude even further, so we will consider eq. (3.8) to be a preference but not an exclusion.

4 Reproducing the 750 GeV excess

Let us discuss the decay widths of H and A as a function of the Higgs mixing $\cos(\beta - \alpha)$ and of the LFV couplings $\rho_{\mu\tau, \tau\mu}$.

The mixing does not affect the couplings of the pseudoscalar A , for which the discussion of section 2 applies. On the other hand, the misalignment parametrised in eq. (3.2) implies that the Yukawa couplings to h and H become

$$\lambda_h^{D,S} = \sin(\beta - \alpha)\lambda_1^{D,S} + \cos(\beta - \alpha)\lambda_2^{D,S}, \quad \lambda_H^{D,S} = \cos(\beta - \alpha)\lambda_1^{D,S} - \sin(\beta - \alpha)\lambda_2^{D,S}. \quad (4.1)$$

The H decay widths into photons and gluons are obtained by replacing $\lambda_2^{D,S}$ with $\lambda_H^{D,S}$ in eqs. (2.2) and (2.4). Similarly, for the corrections to $h \rightarrow \gamma\gamma$ and $h \rightarrow gg$ due to the heavy fermions, one has to replace $\lambda_1^{D,S}$ with $\lambda_h^{D,S}$. In addition, all the h couplings to the Standard Model particles n 's are modified, $g_{hn\bar{n}} = \sin(\beta - \alpha)g_{hn\bar{n}}^{SM}$. Since several Higgs signal strengths have been tested at LHC-8 TeV with 10% precision, the Higgs mixing is bounded from above

$$\cos^2(\beta - \alpha) \lesssim 0.1. \quad (4.2)$$

This is consistent with eq. (3.8). As discussed in section 2, the corrections to $h \rightarrow \gamma\gamma$ and $h \rightarrow gg$ may lead to a slightly stronger upper bound on $\cos(\beta - \alpha)$, if the couplings $\lambda_2^{D,S}$ are very large. However, such bound drops for $\lambda_1^S \cdot \lambda_1^D \rightarrow 0$, see case (2) in section 2. Finally, the contributions to S and T from scalar loops are small in the 2HDM close to the decoupling limit [105, 127], as we explicitly checked for our choice of the parameters.

The mixing has an important effect on the total width of H , since the latter can decay to Standard Model particles n 's, with coupling $g_{Hn\bar{n}} = \cos(\beta - \alpha)g_{hn\bar{n}}^{SM}$. The dominant contributions read, at the tree level,

$$\frac{\Gamma(H \rightarrow t\bar{t}, W^+W^-, ZZ)}{M_H} \simeq \frac{\cos^2(\beta - \alpha)}{8\pi v^2} \left[3m_t^2 + \frac{M_H^2}{2} + \frac{M_H^2}{4} \right] \simeq 0.33 \cos^2(\beta - \alpha), \quad (4.3)$$

where, for the latter numerical estimate, we used the accurate values of the widths for $M_H \simeq 750$ GeV, as given in ref. [128]. Here we neglected the channel $H \rightarrow hh$, because the corresponding trilinear scalar coupling may be suppressed, by conveniently choosing the scalar potential parameters. Recall that the cross-section for $pp \rightarrow H \rightarrow \gamma\gamma$ is proportional to $\Gamma(H \rightarrow gg)/\Gamma_H^{\text{tot}}$, where the numerator corresponds to the assumed dominant H production mode, and the denominator is the total width of H . Therefore, the contribution of H to the excess degrades as soon as $\Gamma(H \rightarrow gg)/M_H \lesssim 0.33 \cos^2(\beta - \alpha)$.

The LFV couplings $\rho_{\mu\tau, \tau\mu}$ also open an additional decay channel for both H and A , with a width

$$\frac{\Gamma(H \rightarrow \tau^\pm \mu^\mp)}{M_H} \simeq \frac{\Gamma(A \rightarrow \tau^\pm \mu^\mp)}{M_A} \simeq \frac{1}{16\pi} \left(|\rho_{\tau\mu}|^2 + |\rho_{\mu\tau}|^2 \right) \simeq \frac{3 \cdot 10^{-7}}{\cos^2(\beta - \alpha)}, \quad (4.4)$$

where the last equality comes from eq. (3.5).

One should also mention that the presently preferred width of the excess, $\Gamma \sim 45$ GeV, could be mimicked by two narrow resonances close in mass. Indeed, the mass split between H and A is given, in the decoupling limit, by $M_H^2 - M_A^2 \simeq \Lambda_5 v^2$, where the term

$\frac{1}{2}\Lambda_5(H_1^\dagger H_2)^2 + \text{h.c.}$ appears in the Higgs potential. This is naturally of the correct order of magnitude for $\Lambda_5 \simeq 1$. Note, however, that the H -mediated cross-section tends to be suppressed relatively to the A -mediated one by two factors: the additional Higgs width in eq. (4.3), and the factor $4/9$ from the loop form factors, see eqs. (2.6)–(2.7).

Let us put all the constraints together to identify the possible windows of parameters that allow to reproduce the 750 GeV excess in agreement with the preferred $h \rightarrow \tau^\pm \mu^\mp$ rate. The resonant LHC total cross-section, in the crude zero-width approximation, reads

$$\sigma(pp \rightarrow H(A) \rightarrow \gamma\gamma) = \sum_i P_i \frac{\Gamma(H(A) \rightarrow i)\Gamma(H(A) \rightarrow \gamma\gamma)}{s \Gamma_{\text{tot}} M_{H(A)}}, \quad (4.5)$$

where $s = (13 \text{ TeV})^2$, $M_{H(A)} \simeq 750 \text{ GeV}$, and the P_i coefficients are the integrals for convoluting over parton densities, that define the parton luminosities for each species i :

$$\begin{aligned} P_{gg} &\equiv \frac{\pi^2}{8} \int_{\frac{M^2}{s}}^1 \frac{dx}{x} g(x)g\left(\frac{M^2}{xs}\right), \\ P_{\bar{q}q} &\equiv \frac{4\pi^2}{9} \int_{\frac{M^2}{s}}^1 \frac{dx}{x} \left[q(x)\bar{q}\left(\frac{M^2}{xs}\right) + \bar{q}(x)q\left(\frac{M^2}{xs}\right) \right]. \end{aligned} \quad (4.6)$$

Consistency with the absence of resonances at 8 TeV favours i to be either gluons or bs , for which the luminosity is $P_{bb} \simeq 14$ and $P_{gg} \simeq 2000$ (we used for eq. (4.6) the latest pdfs from ref. [126]).

We focus first on gluon-gluon fusion as the dominant production mechanism. This channel enjoys the largest parton density functions, so it is sufficient to have $\Gamma(H, A \rightarrow \gamma\gamma)/M_{H,A} \simeq 10^{-6}$ [3], as long as $\Gamma_{\text{tot}}^{H,A} \simeq \Gamma(H, A \rightarrow gg)$. However, the latter is loop-suppressed as shown in eq. (2.7). The total cross-section for some choices of the parameters is shown in figure 2 as a function of $\cos(\beta - \alpha)$. Note that for completeness and cross-check, we have also compared with the more elaborated invariant mass distribution $d\sigma/dM(gg \rightarrow H(A) \rightarrow \gamma\gamma)$ where $M \equiv \sqrt{\hat{s}}$ is the $\gamma\gamma$ invariant mass, that we have calculated taking into account the exact width dependence, and integrating this expression over an appropriate large range for M around the resonance. The numerical differences with the narrow-width approximation expression in eq. (4.5) is at most 2-3 % for all the relevant parameter choices discussed below, as could be intuitively expected since the total width of either A or H remains in all cases sufficiently moderate with respect to the resonance mass, such that the narrow width approximation is justified a posteriori. We can envisage two scenarios:

- (A) For both H and A to contribute significantly to the excess, both the tree-level widths in eq. (4.3) and eq. (4.4) should be small compared with $\Gamma(H, A \rightarrow gg)$. So the optimal value for the Higgs mixing is $\cos^2(\beta - \alpha) \simeq 10^{-3}$, that minimizes the sum of the tree-level widths. Then, to reach a cross-section of a few fb one needs $\Gamma(H, A \rightarrow gg)\Gamma(H, A \rightarrow \gamma\gamma)/M_{H,A}^2 \gtrsim 5 \cdot 10^{-10}$. To reach this value for both H and A requires some stretch in the parameters, e.g. in eqs. (2.2)–(2.5) one should take $R_c = 8$, $Q = 3$, $\lambda_2^S = -1$, $\lambda_2^D = 5$, $M_1 = M_2$ and $\lambda_1^S = -\lambda_1^D$ with the corresponding constraint $(\lambda_1^S v)/(\sqrt{2}M_1) \lesssim 0.12$. In addition, the amplitude for $\tau \rightarrow \mu\gamma$ in this scenario exceeds the indicative bound of eq. (3.7) by about an order of magnitude.

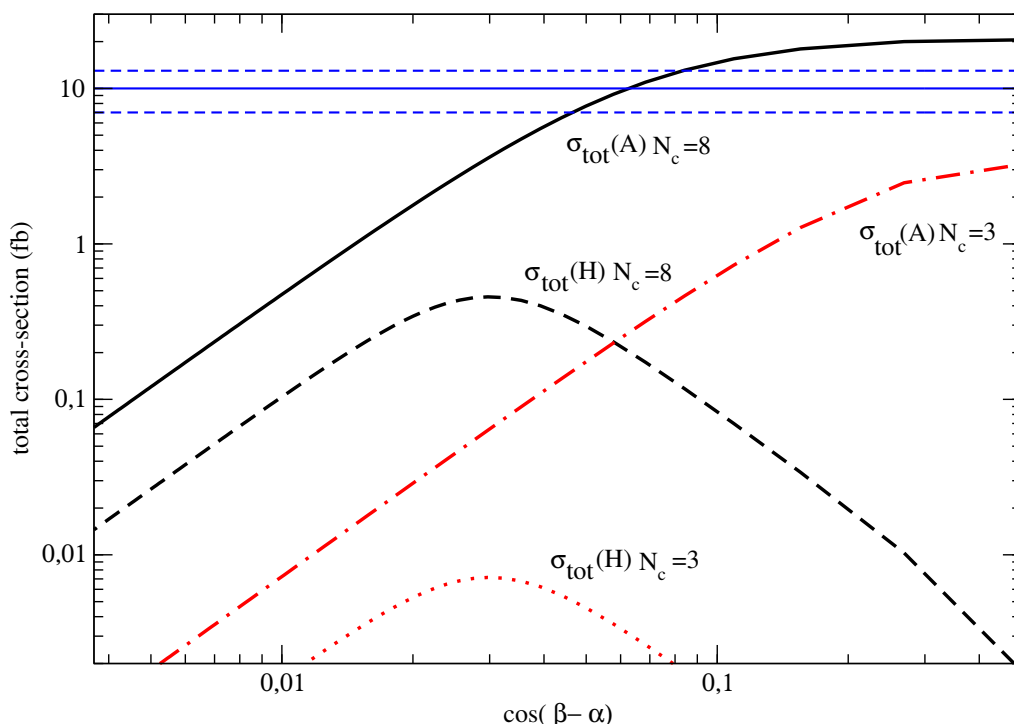


Figure 2. The total cross-section $\sigma(pp \rightarrow H(A) \rightarrow \gamma\gamma)$ in fb, assuming the gluon fusion production channel, as a function of $\cos(\beta - \alpha)$, for a pair of vector-like fermions in the color representation $R_c = 3$ or $R_c = 8$, as indicated. We fixed their charge, $Q = 2$, and their Yukawa couplings to H and A , $\lambda_2^D = 3$ and $\lambda_2^S = 0$. The horizontal band is the preferred cross-section at 1σ for the ATLAS excess [1, 2].

- (B) If one neglects the putative large width of the excess, the H contribution to the signal is no longer necessary. In the case of the pseudoscalar A , only the tree-level width in eq. (4.4) competes with gluon fusion, therefore the signal can be maximized by taking the Higgs mixing as large as allowed by Standard Model constraints, $\cos^2(\beta - \alpha) \simeq 0.1$ (see eq. (4.2)). Then, one can reach a cross-section of a few fb's as long as $\Gamma(A \rightarrow gg)\Gamma(A \rightarrow \gamma\gamma)/M_A^2 \gtrsim 3 \cdot 10^{-12}$, as realized with the reference values in eqs. (2.6)–(2.7). The bound (3.8) from $\tau \rightarrow \mu\gamma$ is satisfied for these parameters.

Let us compare with the alternative possibility that the production of H and A is not dominated by gluon fusion, rather by $b\bar{b} \rightarrow H, A$. The parton density functions give a suppression of order 100 with respect to gluons, so that the excess requires $\Gamma(H, A \rightarrow \gamma\gamma)/M_{H,A} \gtrsim 2 \cdot 10^{-4}$ [3]. The advantage is that a Yukawa coupling $(\rho_b/\sqrt{2})\bar{b}(h_2 + i\gamma_5 A)b$ can easily overcome the other tree-level widths in eq. (4.3) and eq. (4.4),

$$\frac{\Gamma(H \rightarrow b\bar{b})}{\sin^2(\beta - \alpha)M_H} \simeq \frac{\Gamma(A \rightarrow b\bar{b})}{M_A} \simeq \frac{3\rho_b^2}{16\pi} \simeq 0.06\rho_b^2. \quad (4.7)$$

Indeed, one can reproduce the preferred value $\Gamma \simeq 45 \text{ GeV}$ for $\rho_b \simeq 1$. Moreover, there is no constraint from dijet searches at 8 TeV, as the b -quark parton density function is very small. Therefore, one identifies the following scenario:

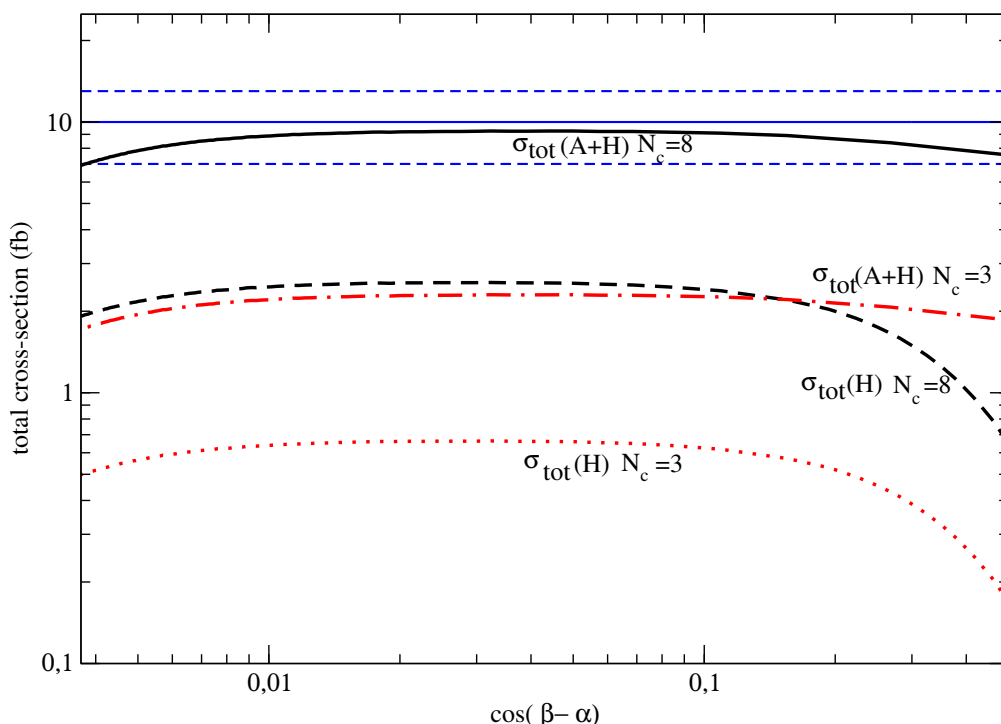


Figure 3. The same as in figure 2, but adding a $b\bar{b}$ production channel with $\rho_b = 1$, see eq. (4.7), and increasing the vector-like fermion charge, $Q = 5$. Here we displayed the cross-section for H only, as well as the cross-section for H plus A .

- (C) When $\Gamma_{\text{tot}}^{H,A} \simeq \Gamma(H, A \rightarrow b\bar{b})$, both H and A contribute to the excess, as long as $\Gamma(H, A \rightarrow \gamma\gamma) \simeq 2 \cdot 10^{-4}$. Confronting with eq. (2.6), one needs a pair of vector-like fermions with $R_c = 3$ and $Q = 7$, or $R_c = 8$ and $Q = 5$. Note that is difficult to avoid such large exotic charges by augmenting the number of multiplets in the loop, as the signal scales with Q^4 . As discussed in section 2, such large Q can be compatible with Higgs decays and the S and T parameters, however the bound of eq. (3.8) from $\tau \rightarrow \mu\gamma$ is exceeded by a factor of few.

The total cross-sections, combining both the gluon fusion and $b\bar{b}$ production channels, are shown in figure 3 as a function of $\cos(\beta - \alpha)$, for $Q = 5$ and other parameters as in figure 2. Here the cross-sections are calculated with the exact width dependence and integrating $d\sigma/dM(gg, b\bar{b} \rightarrow H(A) \rightarrow \gamma\gamma)$. In fact due to the dominant contribution of the $b\bar{b}$ decay to the total width Γ_{tot} in this case, the $b\bar{b}$ production channel largely dominates (for instance the gluon fusion process contributes to the total cross-section by about $\sim 10\%$ only for $R_c = 8$, and much less for $R_c = 3$). Note that in this case the discrepancy with the cross-sections in the narrow width approximation of eq. (4.5) amounts to 7-8 %, for the parameter choices discussed above, that is roughly of order $\Gamma_{\text{tot}}^{H,A}/M_{H,A}$.

5 Final comments

We entertained the possibilities that both the $\gamma\gamma$ excess at 750 GeV and the $h \rightarrow \tau^\pm \mu^\mp$ excess are due to new physics. A minimal way to introduce (renormalisable) flavour viola-

tion and extra bosons to the Standard Model is to add a second Higgs doublet. Its $\tau \leftrightarrow \mu$ coupling may be connected to large 2 – 3 mixing in the neutrino sector, in scenarios where the Yukawa couplings of charged leptons and neutrinos are related.

The neutral scalars H and A can play the role of the 750 GeV resonance, even though the strength of the excess in the early 13 TeV data is significantly larger than the one expected in the 2HDM alone. We take this as a hint that additional states close to the TeV are present in the underlying theory, with large Yukawa couplings to the second Higgs doublet. We have shown that a pair of vector-like fermions is sufficient to reproduce the right cross-section, and respect all other constraints. However, such fermions must have gauge charges larger than the Standard Model fermions: indicatively, for a Yukawa $\simeq 3$ and $R_c \leq 8$, one needs $|Q| \geq 2$ in scenarios (A) and (B), and $|Q| \geq 5$ in scenario (C), see section 4. Alternatively, several pairs of fermions have to be introduced. These are important indications to constrain those well-motivated extensions of the Standard Model that predict vector-like fermions, such as top partners.

Were the heavy Higgses to have no couplings to Standard Model fermions, then $gg \rightarrow H, A \rightarrow \gamma\gamma$ is a natural discovery channel. However, to explain $h \rightarrow \tau^\pm \mu^\mp$, the heavy Higgses must interact with $\tau^\pm \mu^\mp$, and mixing is required between h and H . Both requirements gives Standard Model decay channels to H and A , which reduces $BR(H, A \rightarrow gg, \gamma\gamma)$; nonetheless we find three scenarios that fit both excesses. In addition, the mixing must respect both a lower bound to reproduce the LFV excess, and an upper bound to protect the 125 GeV Higgs couplings: $10^{-3} \lesssim \cos(\beta - \alpha) \lesssim 0.3$.

The decay $\tau \rightarrow \mu\gamma$ is a particular challenge for this model, because the heavy Higgses couple to $\tau^\pm \mu^\mp$ and have an enhanced coupling to $\gamma\gamma$. In combination, these interactions give a “Barr-Zee” contribution to $\tau \rightarrow \mu\gamma$ which is dangerously large. By choosing the Yukawas to obtain destructive interference between A and H , we find that at least two of the scenarios are compatible with the current experimental limit on $\tau \rightarrow \mu\gamma$.

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