750 GeV resonance at the LHC and perturbative unitarity

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If the diphoton excess at 750 GeV hinted by the 2015 data at the LHC is explained in terms of a scalar resonance participating in the breaking of the electroweak symmetry, this resonance must be accompanied by other scalar states for perturbative unitarity in vector-boson scattering to be preserved. The simplest setup consistent with perturbative unitarity and with the data of the diphoton excess is the Georgi-Machacek model. The custodial singlet of the model is responsible for the diphoton excess; it is mainly produced in the diphoton fusion channel, and its loop-induced coupling to the photon pairs is enhanced by the doubly charged scalar with its large (dimensionful) coupling to the custodial singlet.

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I. MOTIVATIONS

Irrespective of whether it will stay or not, the recent excess in the 2015 LHC data with two photons in the final state at an invariant mass of about 750 GeV [1] reminds us that even after the discovery of the Higgs boson we may still not know all the details of the breaking of the electroweak (EW) symmetry.

Let us interpret the LHC diphoton excess as a new scalar resonance.

The simplest (although perhaps least interesting) possibility is that this resonance takes no part in the breaking of the EW symmetry. In this case, it is possible to reproduce the diphoton excess by coupling the resonance—in a generic fashion—to extra scalar or fermionic degrees of freedom (see, for instance, Refs. [2] and [3]). If this is the case, the rationale of such new physics is bound to remain rather mysterious and we might be justified in thinking that it would be for the best if the diphoton excess were to disappear from the new data in 2016.

On the other hand, if this resonance takes part in the EW symmetry breaking, its existence would tell us something new about such a mechanism, in particular that it is not realized by the vacuum expectation value (VEV) of the Higgs boson alone. Moreover—and more importantly for the present work—the presence of such a state necessarily affects the high-energy behavior of the theory: to the extent that the perturbative unitarity of vector-boson scattering is to be preserved, such a resonance cannot come by itself or with arbitrary couplings [4].

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Let us classify states after symmetry breaking according to their properties under custodial $SU(2)_C$ and take the new resonance to be a singlet. There are two possibilities. This custodial singlet either

(i) comes from one or more doublets [this choice leads to the two-Higgs-doublet model (2HDM) [5] and related constructions] and its coupling to the gauge bosons is fixed by gauge invariance to combine with that of the Higgs boson to cancel the unitarity-violating growth with the center-of-mass (CM) energy; or

(ii) its coupling to the gauge bosons does not combine with that of the Higgs boson to cancel the unitarity violations, and we must also include a quintuplet of custodial $SU(2)_C$—the only scalar with a contribution in the high-energy amplitudes of the opposite sign with respect to that of the Higgs boson and other singlets [6]—in order for unitarity to be preserved.

The inclusion of a custodial singlet resonance arbitrary coupled to the gauge bosons therefore leads naturally to the Georgi-Machacek (GM) model [7]—the simplest model that contains a custodial quintuplet and in which symmetry breaking is achieved by three scalar fields: one doublet (with hypercharge $1/2$) and two triplets (with hypercharges 1 and 0).

If neither of the above is the case, perturbative unitarity cannot be preserved and the singlet resonance must belong to a nonperturbative regime. This would imply the exciting discovery of a new interaction that is strong at the EW scale. A fit of the diphoton excess in terms of a nonperturbative resonance is possible and has been already discussed in the literature (for instance, see Ref. [3]).

In this paper we expand on the reasoning above. We discuss to what extent a singlet resonance can take part in the EW symmetry breaking and still belong to a perturbative regime in which reliable computations can be performed. The GM model seems to emerge as the simplest model satisfying these requirements that also explains the diphoton excess at the LHC for a realistic choice of its parameters.

A. Perturbative unitarity

Perturbative unitarity limits the possible models in which the leading orders of perturbation theory are expected to be a reliable guide to physics [6]. If perturbative unitarity is satisfied, EW interactions are described by a renormalizable gauge theory and the strength of the interactions among the particles remain weak at all energies. If this is not the case, unitarity is recovered by the inclusion of higher-order terms; these, however, cannot be small and a nonperturbative regime is entered.

The requirement of perturbative unitarity is stated in terms of partial-wave amplitudes $a_J(s)$ where the amplitude of vector-boson scattering is

$$a_{VV}(s,t) = 16\pi \sum (2J+1) a_J(s) P_J(\cos\theta), \quad (1)$$

and s and t are the Mandelstam variables. Unitarity requires that

$$|a_0(s)| < 1. \quad (2)$$

In general, the partial-wave amplitude in vector-boson scattering is given by

$$a_J(s) = A\left(\frac{\sqrt{s}}{m_W}\right)^4 + B\left(\frac{\sqrt{s}}{m_W}\right)^2 + C\left(\frac{\sqrt{s}}{m_W}\right)^0, \quad (3)$$

with terms growing as the fourth power and the square of the CM energy, and a constant, respectively. A vanishes by gauge invariance which implies $g_{4V} = g_{3V}^2$. B vanishes in the standard model (SM) because of the Higgs boson h contribution and the relationship

$$m_V^2 g_{4V} - \frac{3}{4} m_V^2 g_{3V}^2 = \frac{1}{4} g_{hVV}^2 \quad (4)$$

among the couplings (with self-explanatory notation). The constant terms in C set a limit on the Higgs boson mass in the SM and on the masses of other states in its extensions.

If there are more singlets—for instance two, H_1 and H_1'—their couplings must satisfy

$$m_V^2 g_{4V} - \frac{3}{4} m_V^2 g_{3V}^2 = \frac{1}{4}(g_{H_1 VV}^2 + g_{H_1' VV}^2) \quad (5)$$

in order for the coefficient B in Eq. (3) to vanish. This is realized in the 2HDM and variations of the same.

The other possibility is to have a negative contribution: this can only come from a quintuplet (see Refs. [6] and [8]) of custodial $SU(2)_C$. In fact, for interactions

$$\frac{g_{H_1'} v}{2} H_1 \mathrm{Tr} D_\mu \Sigma^\dagger D^\mu \Sigma \quad (6)$$

and

$$-\frac{g_{H_5} v}{2} H_5 \left[D_\mu \Sigma^\dagger D^\mu \Sigma - \frac{\sigma^{aa}}{6} \mathrm{Tr} D_\mu \Sigma^\dagger D^\mu \Sigma \right] \quad (7)$$

between the longitudinal components of the vector-boson fields $\Sigma = \exp[-i/v \sum \sigma^a \pi^a]$ and the singlet in Eq. (6) and quintuplet in Eq. (53), the amplitudes for singlet scalars are always

$$a(s,t)|_{H_1'} = -\frac{g_{H_1}^2}{v^2} \frac{s^2}{s - m_H^2} \quad (8)$$

with the same sign as the Higgs boson, while

$$a(s,t)|_{H_5} = -\frac{g_{H_5}^2}{v^2} \left[\frac{t^2}{t - m_{H_5}^2} + \frac{u^2}{u - m_{H_5}^2} - \frac{2}{3} \frac{s^2}{s - m_{H_5}^2} \right], \quad (9)$$

gives a (repulsive) negative contribution.

Considering the limit $s \gg m_W^2, m_{H_1}, m'_{H_1} m_{H_5}$—and having the Higgs boson contribution already cancel the contribution from the vector bosons to the coefficient B in Eq. (3)—an exact cancellation between Eq. (8) and Eq. (9) requires

$$\frac{5}{6} g_{H_5}^2 = g_{H_1'}^2. \quad (10)$$

As shown below, such a cancellation, and the unitarity of the theory, are automatically implemented in the GM model.

II. THE FIRST POSSIBILITY: THE 2HDM

The first possibility considered in the introduction section is the simplest: perturbative unitarity is maintained by having the scalar resonance coupling at a special value fixed by gauge invariance [see Eq. (5)].

This would be the first choice in trying to incorporate the resonance within a model. Unfortunately, the parameters of the 2HDM model must be pushed to rather unrealistic values in order to accommodate the diphoton data [9]. These values are particularly worrisome in the light of the required size of the Yukawa couplings, the renormalized values of which bring the theory into a nonperturbative regime [10].

We therefore consider the other case discussed in Sec. I.

III. THE GM MODEL

The GM model contains a complex $SU(2)_L$ doublet field ϕ ($Y = 1$), a real triplet field ξ ($Y = 0$), and a complex $SU(2)_L$ triplet field χ ($Y = 2$). The scalar content of the theory can be organized in terms of the $SU(2)_L \otimes SU(2)_R$ symmetry, and we define the following multiplets:

$$\Phi_{(\mathbf{2},\mathbf{2})} \equiv \begin{pmatrix} \phi^{0*} & \phi^+ \\ \phi^- & \phi^0 \end{pmatrix}, \tag{11}$$

$$\Delta_{(\mathbf{3},\mathbf{3})} \equiv \begin{pmatrix} \chi^{0*} & \xi^+ & \chi^{++} \\ \chi^- & \xi^0 & \chi^+ \\ \chi^{--} & \xi^- & \chi^0 \end{pmatrix}, \tag{12}$$

whose VEVs are

$$\langle\Phi\rangle = \frac{v_\phi}{\sqrt{2}}\hat{I}_{2\times 2} \quad \text{and} \quad \langle\Delta\rangle = v_\Delta \hat{I}_{3\times 3}, \tag{13}$$

with $v_\phi^2 + 8v_\Delta^2 = v^2 = 1/\sqrt{2}G_F \simeq (246\,\text{GeV})^2$. The VEVs of the two triplets must be the same in order to preserve custodial $SU(2)_C$.

The doublet and the two triplet states can be written in components:

$$\phi = \begin{pmatrix} \phi^+ \\ (v_\phi + \phi_r^0 + \imath\phi_i^0)/\sqrt{2} \end{pmatrix}, \tag{14}$$

$$\xi = \begin{pmatrix} \xi^+ \\ v_\Delta + \xi^0 \\ \xi^- \end{pmatrix}, \tag{15}$$

$$\chi = \begin{pmatrix} \chi^{++} \\ \chi^+ \\ v_\Delta + (\chi_r^0 + \imath\chi_i^0)/\sqrt{2} \end{pmatrix}, \tag{16}$$

with $\phi^- = -(\phi^+)^*$, $\xi^- = -(\xi^+)^*$, $\chi^- = -(\chi^+)^*$.

The most general potential that conserves $SU(2)_C$ is given by

$$\begin{aligned} V(\Phi,\Delta) = {}& \frac{\mu_2^2}{2}\mathrm{Tr}\Phi^\dagger\Phi + \frac{\mu_3^2}{2}\mathrm{Tr}\Delta^\dagger\Delta \\ &+ \lambda_1[\mathrm{Tr}\Phi^\dagger\Phi]^2 + \lambda_2\mathrm{Tr}\Phi^\dagger\Phi\mathrm{Tr}\Delta^\dagger\Delta \\ &+ \lambda_3\mathrm{Tr}\Delta^\dagger\Delta\Delta^\dagger\Delta + \lambda_4[\mathrm{Tr}\Delta^\dagger\Delta]^2 \\ &- \lambda_5\mathrm{Tr}(\Phi^\dagger\tau^a\Phi\tau^b)\mathrm{Tr}(\Delta^\dagger t^a\Delta t^b) \\ &- M_1\mathrm{Tr}(\Phi^\dagger\tau^a\Phi\tau^b)(U\Delta U^\dagger)_{ab} \\ &- M_2\mathrm{Tr}(\Delta^\dagger t^a\Delta t^b)(U\Delta U^\dagger)_{ab}, \end{aligned} \tag{17}$$

where τ and t are the $SU(2)$ generators in the doublet and triplet representation respectively, and U is a matrix that rotates Δ into the Cartesian basis.

From the (canonically normalized) kinetic terms

$$\mathcal{L}_{\text{kin}} = |D_\mu^{(\phi)}\phi|^2 + \frac{1}{2}|D_\mu^{(\xi)}\xi|^2 + |D_\mu^{(\chi)}\chi|^2, \tag{18}$$

we can read the interactions with the EW gauge bosons. Considering the neutral components of the scalar fields in Eq. (11), a direct computation gives

$$\begin{aligned} \mathcal{L}_{\text{kin}} \supset {}& (v_\phi + \phi_r^0)^2\left(\frac{g^2}{4}W_\mu^+W^{-,\mu} + \frac{g^2+g'^2}{8}Z_\mu Z^\mu\right) \\ &+ (v_\Delta + \xi^0)^2(g^2 W_\mu^+ W^{-,\mu}) \\ &+ (\sqrt{2}v_\Delta + \chi_r^0)^2\left(\frac{g^2}{2}W_\mu^+W^{-,\mu} + \frac{g^2+g'^2}{2}Z_\mu Z^\mu\right). \end{aligned} \tag{19}$$

The imaginary part of ϕ and χ does not interact with the EW gauge bosons as a consequence of CP invariance. The gauge boson masses are given by

$$m_W^2 \equiv \frac{g^2}{4}(v_\phi^2 + 8v_\Delta^2), \qquad m_Z^2 \equiv \frac{g^2+g'^2}{4}(v_\phi^2 + 8v_\Delta^2). \tag{20}$$

Under $SU(2)_C$ we have the group representations $(\mathbf{2},\mathbf{2}) \sim \mathbf{1}\oplus\mathbf{3}$, and $(\mathbf{3},\mathbf{3}) \sim \mathbf{1}\oplus\mathbf{3}\oplus\mathbf{5}$. One of the two triplets is unphysical, since it represents the Goldstone bosons eaten by the EW gauge bosons. Accordingly, the GM model has ten physical degrees of freedom: two $SU(2)_C$ singlets H_1^0, $H_1^{0\prime}$ (the Higgs and the additional scalar resonance), one $SU(2)_C$ triplet (H_3^+, H_3^0, H_3^-) and one $SU(2)_C$ quintuplet $(H_5^{++}, H_5^+, H_5^0, H_5^-, H_5^{--})$.

If compared with the setup envisaged in Sec. I, the spectrum of the GM model has one additional scalar triplet. However, the triplet H_3 does not interact with the EW gauge bosons.

The mass eigenstates in terms of gauge eigenstates are

$$\begin{aligned} H_5^{++} &= \chi^{++}, \\ H_5^+ &= (\chi^+ - \xi^+)/\sqrt{2}, \\ H_5^0 &= (2\xi^0 - \sqrt{2}\chi_r^0)/\sqrt{6}, \\ H_3^+ &= \cos\theta_H(\chi^+ + \xi^+)/\sqrt{2} - \sin\theta_H\phi^+, \\ H_3^0 &= \imath(-\cos\theta_H\chi_i^0 + \sin\theta_H\phi_i^0), \\ H_1^0 &= \phi_r^0, \\ H_1^{0\prime} &= (\sqrt{2}\chi_r^0 + \xi^0)/\sqrt{3}. \end{aligned} \tag{21}$$

From the Lagrangian in Eq. (19) we find the physical couplings

$$\begin{aligned} \mathcal{L}_{\text{kin}} \supset {}& \cos\theta_H\frac{H_1^0}{v}(2m_W^2W_\mu^+W^{-,\mu} + m_Z^2 Z_\mu Z^\mu) \\ &+ \frac{2\sqrt{2}}{\sqrt{3}}\sin\theta_H\frac{H_1^{0\prime}}{v}(2m_W^2W_\mu^+W^{-,\mu} + m_Z^2 Z_\mu Z^\mu) \\ &+ \frac{2}{\sqrt{3}}\sin\theta_H\frac{H_5^0}{v}(m_W^2W_\mu^+W^{-,\mu} - m_Z^2 Z_\mu Z^\mu), \end{aligned} \tag{22}$$

where the doublet-triplet mixing angle is given by

$$\tan\theta_H \equiv 2\sqrt{2}\frac{v_\Delta}{v_\phi}. \tag{23}$$

As far as the charged interactions are concerned, we find, in the $g' \to 0$ limit,

$$\mathcal{L}_{\rm kin} \supset -2\sin\theta_H \frac{m_W m_Z}{v} H_5^+ W_\mu^- Z^\mu + \text{H.c.} \tag{24}$$

From the interactions in Eqs. (22)–(24) we have

$$g^2_{H_1^0 VV} \equiv \cos^2\theta_H \quad \text{and} \quad g^2_{H_1^{0\prime}} \equiv \frac{8}{3}\sin^2\theta_H \tag{25}$$

for the singlets, and

$$g^2_{H_5} \equiv 2\sin^2\theta_H, \tag{26}$$

for the quintuplet. The cancellation of the coefficient B in the vector-boson scattering amplitude follows from

$$1 - g^2_{H_1^0 VV} - g^2_{H_1^{0\prime}} + \frac{5}{6} g^2_{H_5} = 0. \tag{27}$$

A. Mass spectra and couplings

After EW symmetry breaking, a mixing between the neutral singlet scalar states H_1^0 and $H_1^{0\prime}$ is generated. The corresponding mass matrix is

$$\mathcal{M}^2 = \begin{pmatrix} \mathcal{M}^2_{11} & \mathcal{M}^2_{12} \\ \mathcal{M}^2_{12} & \mathcal{M}^2_{22} \end{pmatrix}, \tag{28}$$

with

$$\begin{aligned} \mathcal{M}^2_{11} &= 8\lambda_1 v_\phi^2, \\ \mathcal{M}^2_{12} &= \frac{\sqrt{3}}{2} v_\phi[-M_1 + 4(2\lambda_2 - \lambda_5)v_\Delta], \\ \mathcal{M}^2_{22} &= \frac{M_1 v_\phi^2}{4v_\Delta} - 6M_2 v_\Delta + 8(\lambda_3 + 3\lambda_4)v_\Delta^2. \end{aligned} \tag{29}$$

The mass matrix can be easily diagonalized by introducing the physical states

$$h = c_\alpha H_1^0 - s_\alpha H_1^{0\prime}, \qquad H = s_\alpha H_1^0 + c_\alpha H_1^{0\prime}, \tag{30}$$

where α is a mixing angle and we used the shorthand notation $c_\alpha \equiv \cos\alpha$, $s_\alpha \equiv \sin\alpha$ from which $\alpha = \pm\sin^{-1}[(1 - c_{2\alpha})/2]$. The mass eigenvalues are

$$2m^2_{h,H} = \mathcal{M}^2_{11} + \mathcal{M}^2_{22} \mp \sqrt{\Delta^2}, \tag{31}$$

with $\Delta^2 \equiv (\mathcal{M}^2_{11} - \mathcal{M}^2_{22})^2 + 4(\mathcal{M}^2_{12})^2$. The mixing angle is defined by

$$s_{2\alpha} = \frac{2\mathcal{M}^2_{12}}{(m_H^2 - m_h^2)}. \tag{32}$$

The masses of the custodial triplet and quintuplet are given by

$$m^2_{H_3} = \left(\frac{M_1}{4v_\Delta} + \frac{\lambda_5}{2}\right)v^2, \tag{33}$$

$$m^2_{H_5} = \frac{M_1}{4v_\Delta}v_\phi^2 + 12M_2 v_\Delta + \frac{3}{2}\lambda_5 v_\phi^2 + 8\lambda_3 v_\Delta^2. \tag{34}$$

Neglecting loop-induced mass splitting, the mass is degenerate within the same custodial multiplet.

As a consequence of the rotation in Eq. (30) and the ratio of VEVs in Eq. (23) the Higgs couplings with gauge bosons and fermions are modified with respect to the corresponding SM values. One finds

$$g_{hW^+W^-} = -\frac{g^2}{6}(8\sqrt{3}s_\alpha v_\Delta - 3c_\alpha v_\phi), \tag{35}$$

$$g_{hf\bar{f}} = -\frac{\imath m_f}{v}\frac{c_\alpha}{\cos\theta_H}, \tag{36}$$

with $g_{hW^+W^-} = c_W^2 g_{hZZ}$.

IV. FITTING THE 750 GEV DIPHOTON EXCESS

There exists a number of constraints that the parameters of the GM model must satisfy in order to reproduce the observed diphoton excess while, at the same time, not be in violation of other known observables.

First of all, for the model to be consistent, its parameters must have the following characteristics.

(i) *Satisfy perturbative unitarity.* Perturbative unitarity on the $2 \to 2$ scalar field scattering amplitudes provides a set of stringent constraints on the parameters of the scalar potential [11,12]:

$$\sqrt{P_\lambda^2 + 36\lambda_2^2} + |6\lambda_1 + 7\lambda_3 + 11\lambda_4| < 4\pi, \tag{37}$$

$$\sqrt{Q_\lambda^2 + \lambda_5^2} + |2\lambda_1 - \lambda_3 + 2\lambda_4| < 4\pi, \tag{38}$$

$$|2\lambda_2 + \lambda_4| < \pi, \tag{39}$$

$$|\lambda_2 - \lambda_5| < 2\pi, \tag{40}$$

with $P_\lambda \equiv 6\lambda_1 - 7\lambda_3 - 11\lambda_4$, $Q_\lambda \equiv 2\lambda_1 + \lambda_3 - 2\lambda_4$. As a consequence, we have

$$\lambda_2 \in \left(-\frac{2}{3}\pi, \frac{2}{3}\pi\right), \qquad \lambda_5 \in \left(-\frac{8}{3}\pi, \frac{8}{3}\pi\right). \tag{41}$$

(ii) *Have a potential bounded from below*. This requirement restricts $\lambda_{1,2,4}$ in the following intervals [12]:

$$\lambda_1 > 0, \tag{42}$$

$$\lambda_4 > \begin{cases} -\frac{\lambda_3}{3} & \text{if } \lambda_3 \geq 0, \\ -\lambda_3 & \text{if } \lambda_3 < 0, \end{cases} \tag{43}$$

$$\lambda_2 > \begin{cases} \frac{\lambda_5}{2} - 2\sqrt{\lambda_1(\frac{\lambda_3}{3} + \lambda_4)} & \text{if } \lambda_5 \geq 0, \lambda_3 \geq 0, \\ \omega_+(\zeta)\lambda_5 - 2\sqrt{\lambda_1(\zeta\lambda_3 + \lambda_4)} & \text{if } \lambda_5 \geq 0, \lambda_3 < 0, \\ \omega_-(\zeta)\lambda_5 - 2\sqrt{\lambda_1(\zeta\lambda_3 + \lambda_4)} & \text{if } \lambda_5 < 0, \end{cases} \tag{44}$$

where we refer to Ref. [12] for the exact definition of ζ and $\omega_\pm(\zeta)$.

In addition, we must verify that, for each choice of parameters, known experimental constraints are satisfied. These are as follows.

(i) *Modification of the SM Higgs couplings*. Higgs coupling measurements [13] strongly constrained the allowed values of v_Δ and α.

(ii) *Electroweak precision tests*. The presence of additional scalar states, charged under the EW symmetry, generates a nonzero contribution to the S parameter [14].

In order to explore the model, we perform a parameter scan by proceeding as follows:

(1) The lightest state h is the physical Higgs boson, with $m_h = 125.09$ GeV, while we identify the second mass eigenstate H with the new resonance at $m_H = 750$ GeV. Equation (31) can be inverted, and one can fix two parameters of the scalar potential. We solve Eq. (31) for λ_1 and M_1.

(2) The parameters $\lambda_{2,3,4,5}$ are randomly generated within the intervals in Eqs. (41)–(44); for each quadruplet, we check that the unitarity constraints are satisfied.

(3) The remaining parameters v_Δ and M_2 are randomly generated within the intervals $v_\Delta \in (0, 50)$ GeV, $|M_2| \in (1, 10^4)$ GeV. The VEV v_ϕ is given by $v_\phi = \sqrt{v^2 - 8v_\Delta^2}$.

(4) For each sample of values the mass matrix in Eq. (28)—and hence the mixing angle α—and the mass eigenstates in Eq. (33) can be computed.

(5) As a final step in our Monte Carlo generation, we check that the values of v_Δ and α are consistent with the Higgs coupling measurements at the 2-σ level. Following Ref. [15], we perform a two-parameter χ^2 fit of the most recent ATLAS and CMS measurements [13]. We show in Fig. 1 the corresponding 1- and 2-σ confidence level contours in the plane (α, v_Δ).

We also check that the correction to the S parameter is within 3-σ of the LEP-I and LEP-II fit of the EW precision observables. In Fig. 3 we show the constraint from the EW parameter S on the scan of the parameters v_Δ and α of the GM model.

Having set the scope and range of the parameter scan, we are now in the position to discuss the fit of the diphoton excess.

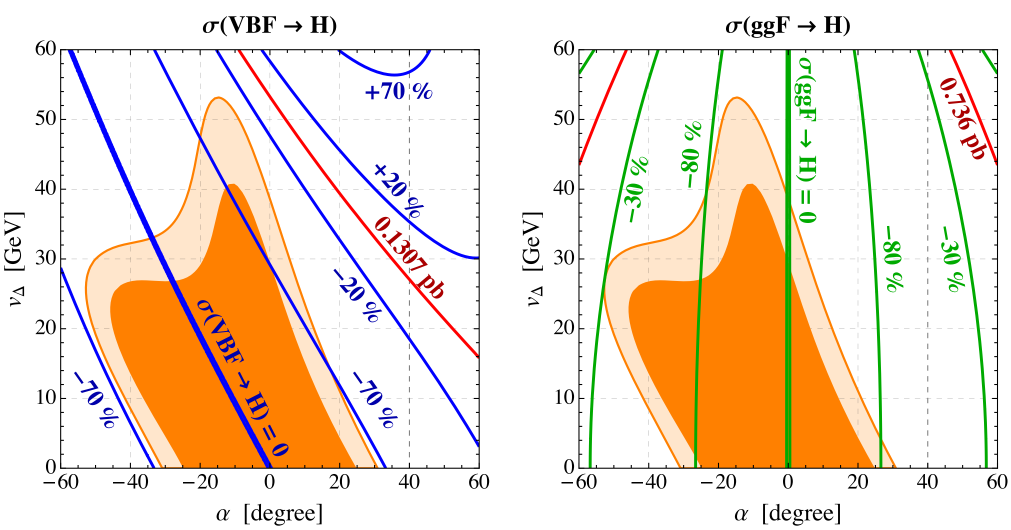

FIG. 1. Contours of the production cross section for the scalar resonance H via VBF (left panel) and ggF (right panel) at $\sqrt{s} = 13$ TeV in the two-dimensional plane (α, v_Δ). In both cases the red line marks the production cross sections for a SM Higgs boson with $m_h = 750$ GeV, which are $\sigma(\text{VBF} \to h)_{m_h = 750\text{ GeV}} \simeq 0.1307$ pb and $\sigma(\text{ggF} \to h)_{m_h = 750\text{ GeV}} \simeq 0.736$ pb. The orange regions represent the 1- and 2-σ confidence levels (darker and lighter orange, respectively) allowed by the Higgs coupling measurements at the LHC.

A. Production cross section

The mixing with the Higgs boson in Eq. (30) and the presence of a nonzero VEV v_Δ automatically allows for H production via both vector-boson fusion (VBF) and gluon fusion (ggF). The former is triggered by tree-level H couplings with the EW gauge bosons, the latter at one loop by H coupling to SM fermions, with the top quark providing the most sizable contribution.

The relevant couplings are

$$g_{HW^+W^-} = \frac{g^2}{6}(8\sqrt{3}c_\alpha v_\Delta + 3s_\alpha v_\phi), \tag{45}$$

$$g_{Ht\bar{t}} = -\frac{\imath m_t}{v}\frac{s_\alpha}{\cos\theta_H}. \tag{46}$$

The H production cross section can be straightforwardly obtained by rescaling the production cross section of a SM Higgs with $m_h = 750$ GeV. At $\sqrt{s} = 13$ TeV we have $\sigma(\mathrm{VBF} \to h)_{m_h=750\ \mathrm{GeV}} \simeq 0.1307$ pb and $\sigma(\mathrm{ggF} \to h)_{m_h=750\ \mathrm{GeV}} \simeq 0.736$ pb [16], and the rescaling is simply given by

$$\sigma(\mathrm{VBF} \to H) = (c_V^H)^2 \times \sigma(\mathrm{VBF} \to h)_{m_h=750\ \mathrm{GeV}},$$
$$\sigma(\mathrm{ggF} \to H) = (c_F^H)^2 \times \sigma(\mathrm{ggF} \to h)_{m_h=750\ \mathrm{GeV}}, \tag{47}$$

where

$$c_V^H = \frac{1}{3}\left[\frac{8\sqrt{3}c_\alpha v_\Delta + 3s_\alpha v_\phi}{v}\right], \tag{48}$$

$$c_F^H = \frac{v s_\alpha}{\sqrt{v^2 - 8v_\Delta^2}}. \tag{49}$$

The rescaled cross sections crucially depend on the values of v_Δ and α. In Fig. 1 we show contours of constant VBF (left panel, blue lines) and ggF (right panel, green lines) H production compared with the reference values of the SM Higgs with $m_h = 750$ GeV (red lines). As is clear from the plot, in the allowed region of the (α, v_Δ) plane we always observe a reduction if compared with the SM case.

In addition to VBF and ggF, we also include—following Ref. [17]—production via photon fusion ($\gamma\gamma$F) for inelastic, partially elastic and elastic collisions.

B. Total decay width and diphoton decay

The diphoton signal strength at $\sqrt{s} = 13$ TeV is given by

$$\mu_H = [\sigma(\mathrm{ggF} \to H) + \sigma(\mathrm{VBF} \to H)] \times \mathcal{BR}(H \to \gamma\gamma) + 10.8\ \mathrm{pb}\left(\frac{\Gamma_H}{45\ \mathrm{GeV}}\right) \times [\mathcal{BR}(H \to \gamma\gamma)]^2, \tag{50}$$

where the last line accounts for production via $\gamma\gamma$F [17].

Given the preliminary status of the experimental analysis, we do not perform any complicated fit. On the contrary, the purpose of this section is to check whether the GM model can account for a diphoton signal strength of the order of a few fb, which is the order of magnitude suggested by present data. As discussed in Sec. II, a positive answer is anything but trivial in weakly coupled theories (in particular without invoking the presence of extra vector-like fermions with either large multiplicities, electric charge or Yukawa couplings) and would be a remarkable result if achieved in the GM model.

In order to evaluate Eq. (50) we need to compute the total decay width of the singlet, Γ_H, and the diphoton decay width.

At the tree level, H predominantly decays—as far as the SM final states are concerned—into W^+W^-, ZZ, $t\bar{t}$ and hh. The corresponding decay widths can be computed by rescaling those of the SM Higgs boson. We find

$$\Gamma_{VV}^{(H)} = \frac{G_\mu m_H^3 (c_V^H)^2 \delta_V}{16\sqrt{2}\pi}\sqrt{1-4x_V}(1 - 4x_V + 12x_V^2),$$
$$\Gamma_{f\bar{f}}^{(H)} = \frac{G_\mu N_C m_H m_f^2 (c_F^H)^2}{4\sqrt{2}\pi}\left(1 - \frac{4m_f^2}{m_H^2}\right)^{3/2},$$
$$\Gamma_{hh}^{(H)} = \frac{g_{hhH}^2}{32\pi m_H}\sqrt{1 - \frac{4m_h^2}{m_H^2}}, \tag{51}$$

where $\delta_{V=W,Z} = 2(1)$, $x_V = m_V^2/m_H^2$, $G_\mu = 1/(\sqrt{2}v)^{1/2}$. The trilinear scalar coupling is [11]

$$g_{hhH} = 24\lambda_1 c_\alpha^2 s_\alpha v_\phi + 8\sqrt{3}c_\alpha s_\alpha^2 v_\Delta(\lambda_3 + 3\lambda_4) + 2[\sqrt{3}c_\alpha v_\Delta(3c_\alpha^2 - 2) + s_\alpha v_\phi(1 - 3c_\alpha^2)](2\lambda_2 - \lambda_5) - \frac{\sqrt{3}}{2}M_1 c_\alpha(3c_\alpha^2 - 2) - 4\sqrt{3}M_2 c_\alpha s_\alpha^2. \tag{52}$$

The singlet H can also decay into the custodial triplet and quintuplet if the corresponding channels are kinematically allowed. If $m_H > m_{H_5}/2$ ($m_H > m_{H_3}/2$), the new decay channels are $\Gamma^{(H)}_{H_5^+H_5^-}$, $\Gamma^{(H)}_{H_5^{++}H_5^{--}}$, $\Gamma^{(H)}_{H_5^0H_5^0}$ ($\Gamma^{(H)}_{H_3^+H_3^-}$, $\Gamma^{(H)}_{H_3^0H_3^0}$). The decay widths can be computed as in Eq. (51), and the relevant couplings are [11]

$$g_{HH_5^0H_5^0} = 8\sqrt{3}(\lambda_3 + \lambda_4)c_\alpha v_\Delta + (4\lambda_2 + \lambda_5)s_\alpha v_\phi + 2\sqrt{3}M_2 c_\alpha, \tag{53}$$

with $g_{HH_5^0H_5^0} = g_{HH_5^+H_5^-} = g_{HH_5^{++}H_5^{--}}$, and

$$g_{HH_3^0H_3^0} = 64\lambda_1 s_\alpha \frac{v_\Delta^2 v_\phi}{v^2} + \frac{8v_\phi^2 v_\Delta}{\sqrt{3}v^2} c_\alpha(\lambda_3 + 3\lambda_4) - \frac{2\sqrt{3}M_2 v_\phi^2}{v^2} c_\alpha + \frac{16 v_\Delta^3 c_\alpha}{\sqrt{3}v^2}(6\lambda_2 + \lambda_5) + \frac{4v_\Delta M_1}{\sqrt{3}v^2}(c_\alpha v_\Delta + \sqrt{3}s_\alpha v_\phi) + \frac{s_\alpha v_\phi^3}{v^2}(4\lambda_2 - \lambda_5) + \frac{8\lambda_5 v_\Delta v_\phi}{\sqrt{3}v^2}(c_\alpha v_\phi + \sqrt{3}s_\alpha v_\Delta), \quad (54)$$

with $g_{HH_3^0H_3^0} = g_{HH_3^+H_3^-}$.

Finally, H can decay into a vector boson plus a custodial triplet scalar. If $m_H > m_W + m_{H_3}$ and $m_H > m_Z + m_{H_3}$ the corresponding decay channels are $\Gamma^{(H)}_{W^\pm H_3^\mp}$ and $\Gamma^{(H)}_{ZH_3^0}$. We find

$$\Gamma^{(H)}_{VH_3} = \frac{|g_{HVH_3}|^2 m_V^2}{16\pi m_H}\lambda\left(\frac{m_H^2}{m_V^2}, \frac{m_{H_3}^2}{m_V^2}\right)\lambda^{1/2}\left(\frac{m_V^2}{m_H^2}, \frac{m_{H_3}^2}{m_H^2}\right), \quad (55)$$

where the kinematic function λ is $\lambda(x,y) = (1-x-y)^2 - 4xy$. The relevant couplings are

$$g_{HZH_3^0} = \frac{i\sqrt{2}g}{\sqrt{3}c_W}\left(\frac{c_\alpha v_\phi}{v} - \frac{\sqrt{3}v_\Delta s_\alpha}{v}\right), \quad (56)$$

$$g_{HW^\pm H_3^\mp} = -\frac{\sqrt{2}g}{\sqrt{3}}\left(\frac{\sqrt{3}s_\alpha v_\Delta}{v} - \frac{c_{\rm alpha} v_\phi}{v}\right). \quad (57)$$

The sum of the tree-level decay widths reconstructs the total width Γ_H.

The loop-induced diphoton decay width for the scalar singlet $\mathcal{H} = h$, H is therefore

$$\Gamma^{(\mathcal{H})}_{\gamma\gamma} = \frac{G_\mu \alpha^2 m_{\mathcal{H}}^3}{128\sqrt{2}\pi^3}\left|\sum_f N_C Q_f^2 g_{\mathcal{H}f\bar{f}} A^{\mathcal{H}}_{1/2}(\tau_f) + g_{\mathcal{H}W^+W^-}A_1^{\mathcal{H}}(\tau_W) + \sum\nolimits_s \beta_s Q_s^2 A_0^{\mathcal{H}}(\tau_s)\right|^2, \quad (58)$$

where the loop functions are known and can be found, for instance, in Ref. [18]. The last term in Eq. (58) represents the contribution of the electrically charged scalar states, and we have $\beta_s \equiv g_{\mathcal{H}H_sH_s^*} v/2m_s^2$.

The electrically charged scalars affect the diphoton decay of both the new scalar resonance H and the Higgs h [the scalar couplings in Eqs. (53) and (54) for the Higgs boson can be found in Ref. [11]]. The challenge is to explain the diphoton signal strength observed by ATLAS and CMS without introducing a big deviation in the diphoton Higgs decay.

C. Results: μ_H and Γ_H

In Fig. 2 we show the parameter scan in the plane (μ_H, Γ_H). There exists a particular region of the scan where the model reproduces a signal strength with size $\mu_H \sim O(1)$ fb. The red points in Fig. 2, where μ_H is larger, correspond to the right-hand side of the allowed interval in $-M_2 \in (1, 10^4)$ GeV. In this range of values the scalar couplings in Eqs. (53)–(54) are large, thus dominating the loop in $\Gamma^H_{\gamma\gamma}$.

In Fig. 3 we recast the parameter scan in the plane (α, v_Δ). The yellow contours agree with Ref. [19]. We see that points where $\mu_H \sim O(1)$ fb correspond to a small and negative mixing angle, $\alpha \sim -3°$ and triplet VEV $v_\Delta \lesssim 20$ GeV. In this region the dominant contribution to the production cross section is given by $\gamma\gamma$F. Production by means of ggF and VBF contributes up to 20%. As a consequence, the tension between the diphoton excess observed at $\sqrt{s} = 13$ TeV and the absence of such a signal in the data set at $\sqrt{s} = 8$ TeV is alleviated. The production cross section via ggF—going from $\sqrt{s} = 13$ TeV to $\sqrt{s} = 8$ TeV—is reduced by the factor $\sigma(\mathrm{ggF} \to H)_{13\ \mathrm{TeV}}/\sigma(\mathrm{ggF} \to H)_{8\ \mathrm{TeV}} = 4.693$ while the production cross section via $\gamma\gamma$F is reduced by a factor of 2. These scaling factors make the diphoton excess at $\sqrt{s} = 13$ TeV consistent with the bound extracted from the $\sqrt{s} = 8$ TeV data set.

In Fig. 4 we recast the parameter scan in the plane (m_{H_5}, m_{H_3}). Points where $\mu_H \sim O(1)$ fb correspond to $m_{H_5} \sim 400$–600 GeV, $m_{H_3} \sim 650$–700 GeV. This feature is expected because for these values the corresponding loop in the diphoton decay amplitude of H is maximized.

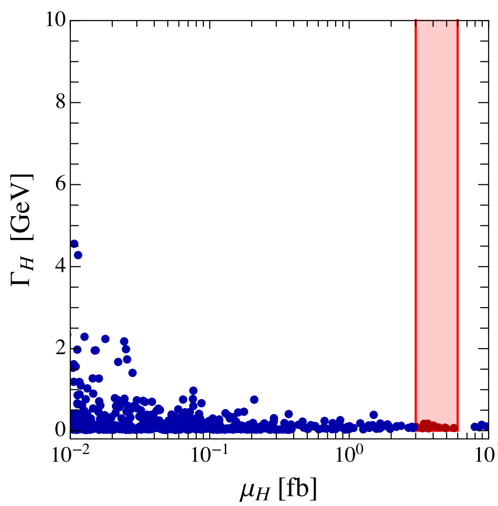

FIG. 2. Result of the parameter scan in terms of the total decay width Γ_H versus the diphoton signal strength μ_H for the new scalar resonance at $m_H = 750$ GeV. We mark in red the points where $\mu_H = [3–6]$ fb, as suggested by experimental data on the diphoton excess.

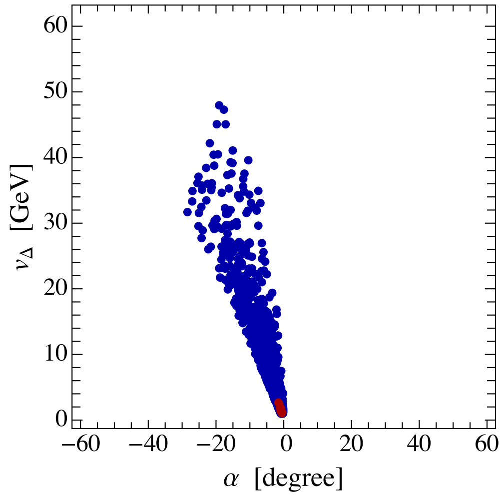

FIG. 3. Result of the parameter scan in terms of the mixing angle α versus the triplet VEV v_Δ. We superimpose the analyzed points on the region allowed by Higgs coupling measurements. The constraint from the EW parameter S on the scan is shown in green (1-, 2- and 3-σ confidence level regions correspond to lighter shades).

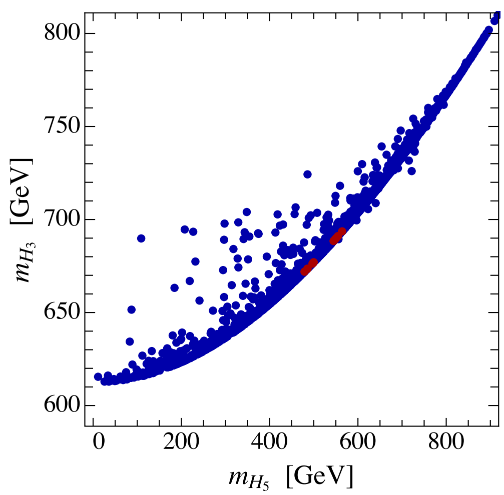

FIG. 4. Result of the parameter scan in terms of the custodial triplet versus quintuplet masses, $m_{H_{5,3}}$.

The explanation of the diphoton excess in the context of the GM model predicts the presence of additional light scalar degrees of freedom, including the doubly charged state H_5^{++}. Notice that tree-level decays of H into triplet or quintuplet scalar states are not kinematically allowed at the red points of the scan. The characteristic phenomenology [19,20] of these scalar states represents a signature of the model.

We checked that the model, for the chosen choice of parameter values, is consistent with other searches for resonant production of a pair of SM particles which constrain the tree-level decay modes of H [2].

As it can be seen in Fig. 3, there is a moderate tension with the EW parameter S for which the fit of the diphoton excess (the red dots) only agrees at the 3-σ level. This is to be expected given the presence of the additional charged new states.

The Higgs scaling factor κ_γ is defined as the ratio between the loop-induced $h \to \gamma\gamma$ coupling in the GM model with respect to that of the SM. At the red points in Fig. 2, we find $0.8 \lesssim \kappa_\gamma \lesssim 1.2$. The presence of such a deviation is consistent with the present experimental bound [13].

We find that the other two neutral scalars, H_3^0 and H_5^0 give a negligible contribution to the diphoton cross section.

Concerning the total decay width Γ_H, points where $\mu_H \sim O(1)$ fb correspond to $\Gamma_H \sim 1$ GeV. The value of the total decay width suggested by data represents at the moment the most controversial aspect of the diphoton excess. Since the typical diphoton invariant mass resolution at 750 GeV is estimated to be around 10 GeV, it is natural to expect a large total decay width, $\Gamma_H \lesssim 40$ GeV. At this stage of the experimental analysis no conclusive statements can be made, and the value $\Gamma_H \sim 1$ GeV is perfectly consistent with the data. However, if large values of Γ_H are confirmed by future analysis, an explanation of the diphoton excess in terms of weakly coupled theories will be disfavored.

D. Perturbative reliability

The result above is qualitatively different with respect to both the case in which the resonance is not taking part in the EW symmetry breaking (and one is forced to introduce additional electrically charged vector-like fermions to boost both the production cross section and diphoton decay) and the 2HDM [in which the condition $\mu_H \sim O(1)$ fb requires unrealistically large Yukawa couplings]. In our scan, all the dimensionless couplings of the GM model are kept within the perturbative regime.

This point is better understood in terms of the overall size of the diphoton decay induced by the loop of scalar particles. In full generality, we can consider the effective Lagrangian

$$\mathcal{L}_{\rm eff} = \frac{e^2}{4v} c_{\gamma\gamma} H A_{\mu\nu} A^{\mu\nu}, \qquad (59)$$

where $A_{\mu\nu}$ is the usual photon field strength. The effective operator in Eq. (59) induces the diphoton decay

$$\Gamma_{\gamma\gamma}^{(H)} = \frac{c_{\gamma\gamma}^2 e^4 m_H^3}{64\pi v^2}. \qquad (60)$$

We can recast, for illustrative purposes, the scalar loop contribution in Eq. (58) in terms of the Wilson coefficient $c_{\gamma\gamma}$. Approximating for simplicity the scalar loop function as $A_0(\tau) \sim -1/3$, we find

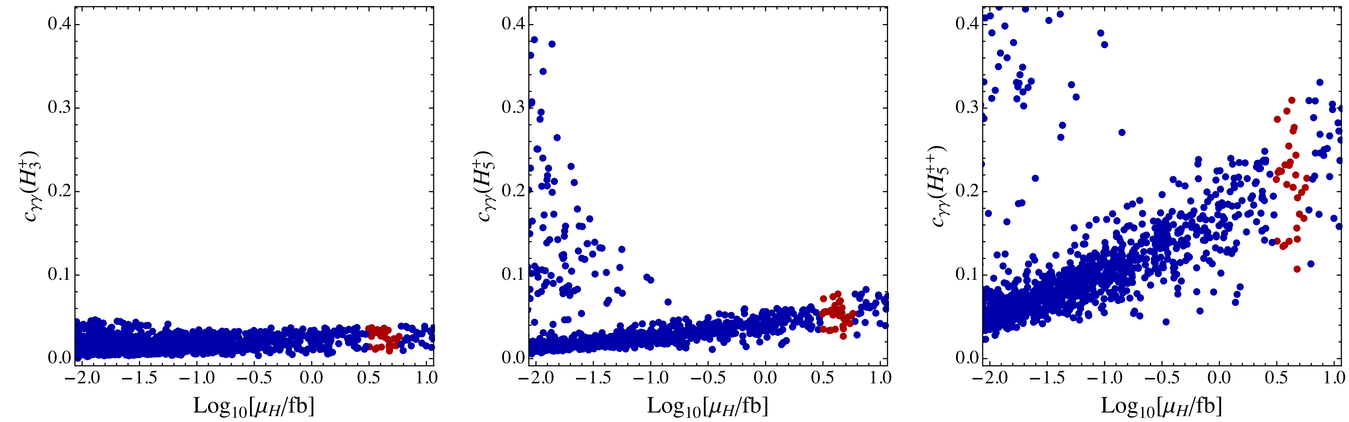

FIG. 5. Result of the parameter scan in terms of the Wilson coefficient in Eq. (61).

$$c_{\gamma\gamma}(s) = \left[\frac{\beta_s^2 Q_s^4}{36(4\pi)^2\pi^2}\right]^{1/2}, \qquad s = H_5^+, H_5^{++}, H_3^+. \tag{61}$$

In Fig. 5 we show the typical size of these coefficients in our parameter scan. The typical size is $c_{\gamma\gamma}(s) \sim 0.05$. The only exception is $c_{\gamma\gamma}(H_5^{++})$, which can reach values $c_{\gamma\gamma}(H_5^{++}) \lesssim 0.4$ (due to the large electric charge, $Q^4_{H_5^{++}} = 16$).

E. Stability of the vacuum

The GM potential (17) as a function of v_Δ has two minima whose depth depends on the choice of the other parameters. The values necessary to explain the diphoton excess give rise to a second minimum which is at a value much larger than the one we used and which is actually deeper. A complete study of the stability of the vacuum would be necessary to verify the metastability of the first (smaller) vacuum against the decay to the deeper one. We postpone such an analysis to further work after the existence of the resonance is better established.

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