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Is the LHC observing the pseudoscalar state of a two-Higgs-doublet model?

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The ATLAS and CMS collaborations have recently shown data suggesting the presence of a Higgs boson in the vicinity of 125 GeV. We show that a two-Higgs-doublet model spectrum, with the pseudoscalar state being the lightest, could be responsible for the diphoton signal events. In this model, the other scalars are considerably heavier and are not excluded by the current LHC data. If this assumption is correct, future LHC data should show a strengthening of the $\gamma\gamma$ signal, while the signals in the $ZZ^{(*)} \to 4\ell$ and $WW^{(*)} \to 2\ell 2\nu$ channels should diminish and eventually disappear, due to the absence of diboson tree-level couplings of the CP-odd state. The heavier CP-even neutral scalars can now decay into channels involving the CP-odd light scalar which, together with their larger masses, allow them to avoid the existing bounds on Higgs searches. We suggest additional signals to confirm this scenario at the LHC, in the decay channels of the heavier scalars into AA and AZ. Finally, this inverted two-Higgs-doublet spectrum is characteristic in models where fermion condensation leads to electroweak symmetry breaking. We show that in these theories it is possible to obtain the observed diphoton signal at or somewhat above the prediction for the standard model Higgs for the typical values of the parameters predicted.

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INTRODUCTION

Recently, the ATLAS and CMS collaborations have reported important exclusions in the Higgs mass with about 5 fb^{-1} of accumulated luminosity [1,2]. However, both experiments observe excess signal events at low masses. The most significant of these is in the diphoton channel $h \to \gamma\gamma$, which would point to a Higgs mass of $m_h \simeq$ 126 GeV. For ATLAS [3] the local significance of this excess is of 2.8σ, whereas for CMS [4] it is about 3.0σ. ATLAS [5] and CMS [6] also observe modest excesses in the $h \to ZZ^{(*)} \to 4\ell$ channel with local significance not much above 2.0σ, and CMS [7] has also searched for an enhanced $\gamma\gamma jj$ signal coming from vector boson fusion (VBF), which could be associated with fermiophobic Higgs models. With cuts that significantly reduce the gluon fusion contribution, CMS has found a 2.7σ excess. Finally, ATLAS [8] has also searched for similar signals, which are enhanced by a cut in the transverse momentum of the diphoton system. They found an excess of about 3σ. Thus, combined, both experiments appear to coincide in the existence of a diphoton excess somewhere around (124–126) GeV, while differing on the invariant mass of the less significant excesses observed in $ZZ^{(*)} \to 4\ell$, as well as in $WW^{(*)} \to \ell^+\nu\ell^-\bar{\nu}$.

In this paper we consider the possibility that the excess in the $\gamma\gamma$ channel is real, but that it is caused by the pseudoscalar state A in a two-Higgs-doublet model (THDM). We assume that the spectrum of the THDM is inverted with respect to what is usually considered, for instance, in the minimal supersymmetric standard model (MSSM) [9]: A is the lightest state, with $(h, H, H^{\pm})$ much heavier and with small splittings among them [10,11]. In considering this possibility, we must address the VBF-enriched samples in Refs. [7,8], which find excesses in regions that favor VBF and disfavor gluon fusion. However, the excesses are still small (3σ and 2.7σ for CMS and ATLAS, respectively), and the techniques used are new. In any case, the pure gluon fusion interpretation is still compatible with both the ATLAS and CMS data at about the 3σ level [12]. So more data are needed.

The inverted spectrum can naturally appear in a generic THDM, just by choosing the right set of parameters in the potential, even after demanding tree-level unitarity and stability, as well as the correct minimum for electroweak symmetry breaking [10,13]. On the other hand, this spectrum of the THDM is typical in models where unconfined fermions condense to break the electroweak symmetry dynamically [14], due to the existence of an approximate Peccei-Quinn symmetry which keeps the pseudoscalar state light in comparison to the rest of the scalars.

Whatever the origin of this scalar spectrum, the first signal of it would be the observation of $A \to \gamma\gamma$ if $m_A <$ 130 GeV, just as in the case of the SM Higgs. The production cross section times branching ratio for $gg \to A \to \gamma\gamma$ need not be the same as that of the SM Higgs. In fact, the ATLAS signal is somewhat larger than the SM prediction for a Higgs of $\simeq$ 125 GeV. On the other hand, since A has no tree-level couplings to ZZ and WW, these channels would not be observed for the lightest mass peak. More importantly, the current excesses in the VBF channels studied by ATLAS and CMS should disappear. If this is the case, other channels will have to be studied to confirm the THDM explanation. In what follows, we specify the parameter space of this scenario as well as its predictions for future LHC data samples. We also point out the strategy for finding the other states that would confirm the THDM hypothesis. Finally, we speculate on the possible dynamical origin of this peculiar THDM spectrum and its relation to theories of electroweak symmetry breaking (EWSB).

II. PREDICTIONS IN THE INVERTED THDM

In order to specify the phenomenology of the THDM we must define the scheme of fermion couplings. There are four possible choices that would avoid flavor changing neutral currents (FCNCs) at tree level: the so-called type I, type II, lepton-specific, and flipped schemes [13,15], depending on the choice of doublet responsible for the masses of right-handed fermions. The couplings of the scalar sector are determined by $\tan\beta \equiv v_2/v_1$, the ratio of vacuum expectation values of the two doublets, and the mixing angle α between the neutral CP-even scalars. However, the couplings of A to fermions only depend on β. For instance, the couplings of A to all up-type quarks in all four schemes are $\cot\beta$. On the other hand, its couplings to down-type, go like $\cot\beta$ in the type I and lepton-specific cases, whereas they go like $\tan\beta$ for the type II and flipped schemes. Finally, the couplings of charged leptons to A in type I and flipped scenarios go like $-\cot\beta$, while the same couplings in type II and lepton-specific schemes go like $\tan\beta$.

We are interested in calculating the $\sigma\times$ BR$(gg \to A \to \gamma\gamma)$ in the inverted THDM. The production cross section for $gg \to A$ is given by [16]

$$\sigma_A = \frac{9}{4}\frac{|\cot\beta I_A(\tau_t) + \xi(\beta) I_A(\tau_b)|^2}{|I_S(\tau_t)|^2}\sigma_h^{\rm SM}, \qquad (1)$$

where $\sigma_h^{\rm SM}$ is the SM $gg \to h$ production cross section, and the functions in (1) are given by

$$I_A(\tau) = \frac{f(\tau)}{\tau}, \qquad (2)$$

$$I_S(\tau) = \frac{3}{2}[\tau + (\tau - 1)f(\tau)]/\tau^2, \qquad (3)$$

and

$$f(\tau) = \begin{cases} \arcsin^2\sqrt{\tau} & \tau \le 1 \\ -\frac{1}{4}\left[\log\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi\right]^2 & \tau > 1, \end{cases} \qquad (4)$$

expressed in terms of the variable $\tau_f \equiv m_{A,h}^2/4m_f^2$. The factor $\xi(\beta)$ in (1) depends on the choice of THDM scheme: for type I and lepton-specific models, $\xi(\beta) = -\cot\beta$, while for type II and flipped cases we have $\xi(\beta) = \tan\beta$. For moderate values of $\tan\beta$ the top quark contribution dominates the ggA vertex. The b quark term is important for large $\tan\beta$ in the type II and flipped schemes.

Similarly, we can compute the $A \to \gamma\gamma$ decay width normalized by the SM $h \to \gamma\gamma$ width, obtaining

$$\frac{\Gamma(A\to\gamma\gamma)}{\Gamma^{\rm SM}(h\to\gamma\gamma)} = 4\frac{|N_c q_U^2 \cot\beta I_A(\tau_t) + N_c q_D^2 \xi(\beta) I_A(\tau_b)|^2}{|N_c(4/3)(q_U^2 I_S(\tau_t) + q_D^2 I_S(\tau_b)) I_W(\tau_W)|^2}, \qquad (5)$$

where $q_{U,D}$ are the charges of up- and down-type quarks and the W contribution to the SM Higgs decay to diphotons is accounted for by the function

$$I_W(\tau) = -[2\tau^2 + 3\tau + 3(2\tau - 1)f(\tau)]/\tau^2. \qquad (6)$$

We are now ready to compute the $\sigma\times$ BR$(gg \to A \to \gamma\gamma)$ in the inverted THDM normalized by the analogous SM Higgs process. We have not included QCD corrections to the quark contributions since they largely cancel in the ratio. In Fig. 1 we plot the ratio $R = \sigma \times {\rm BR}(gg \to A \to \gamma\gamma)/\sigma \times {\rm BR}(gg \to h \to \gamma\gamma)$ vs $\tan\beta$. The dashed line corresponds to the type I and lepton-specific schemes, with the solid curve being for the type II and flipped cases. We see that excesses over the SM prediction for $gg \to h \to \gamma\gamma$ can only be obtained for $\tan\beta < 0.9$ or so. Larger values of $\tan\beta$ would imply $R < 1$, in contradiction with the ATLAS data, as long as we interpret it as purely coming from the signal (i.e. not aided by a significant upward fluctuation). The top Yukawa coupling only becomes nonperturbative for $\tan\beta < 0.3$, so these values are safe [13]. Thus, we see that the inverted THDM can explain the diphoton signal at the LHC for these small values of $\tan\beta$.

This region of parameter space is allowed by all other data. In addition to the absence of FCNC at tree level, the loop contributions to FCNC processes are safely below bounds, given that the charged states of the model, $H^\pm$, are assumed to have masses well in excess of the 95% C.L. bound $m_{H^\pm} > 316$ GeV which is mostly driven by $b \to s\gamma$ [17]. This region is also generally compatible with precision electroweak constraints. In Fig. 2 we show the con-

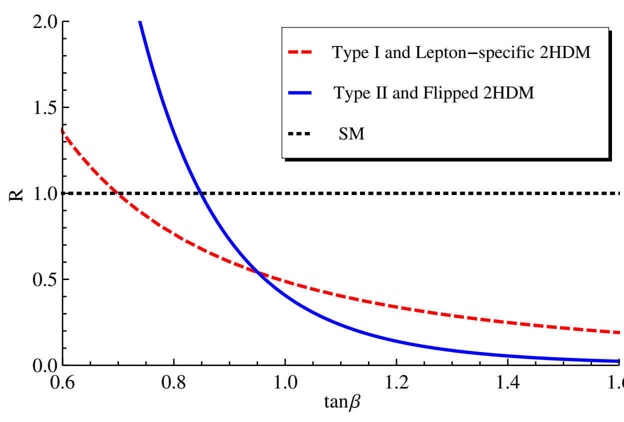

FIG. 1 (color online). The ratio $R = \sigma \times {\rm BR}(gg \to A \to \gamma\gamma)/\sigma \times {\rm BR}(gg \to h \to \gamma\gamma)$ vs $\tan\beta$. The dashed line corresponds to the type I and lepton-specific schemes, with the solid curve being for the type II and flipped cases. The horizontal dotted line corresponds to $\sigma \times {\rm BR}(gg \to A \to \gamma\gamma)$, equal to the prediction for this process mediated by the SM Higgs.

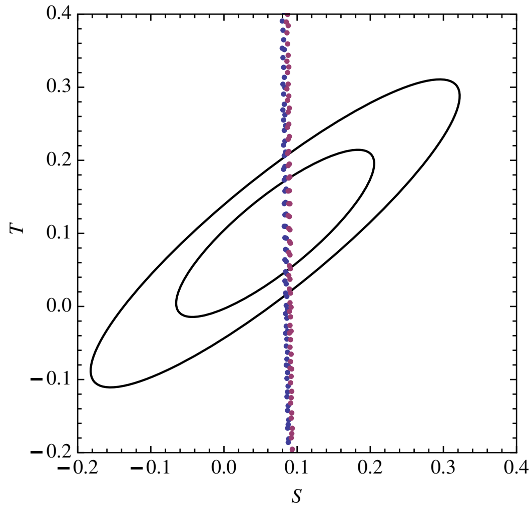

FIG. 2 (color online). Electroweak precision bounds on the inverted THDM. We take $m_A = 125$ GeV and $m_h = (500, 550)$, with m_H and $m_{H^\pm}$ varying in the range (550–650) GeV. A scan over this parameter space results in the points in the figure. The ones outside the 95% C.L. contour come from larger values of $|m_H - m_{H^\pm}|$, which result in large custodial breaking. We take $m_h^{\mathrm{ref.}} = 117$ GeV [19].

tributions to the S and T parameters within the inverted THDM. In particular, we take $m_A = 125$ GeV and a scalar spectrum that is heavy enough not to have been seen at the LHC [18] so far. For the two choices $m_h = 500$ GeV and $m_h = 550$ GeV, we vary m_H and $m_{H^\pm}$ in the range (550–650) GeV. The points represent the results of scanning over this range. We also show the 68% and 98% C.L. intervals obtained from a fit of electroweak data as performed in Ref. [19] using $m_h^{\mathrm{ref.}} = 117$ GeV. We can see from Fig. 2 that there are many solutions within the inverted THDM spectrum that are compatible with electroweak bounds. The points falling out of the allowed region are those with large values of $|m_H - m_{H^\pm}|$, which represent a breaking of custodial symmetry. As long as this mass difference is not large compared to m_W, the inverted THDM spectrum is within the allowed values of the electroweak parameters S and T.

The bounds on Higgs searches from ATLAS and CMS suggest that the rest of the scalar spectrum of the THDM is quite heavy. A detailed study of the exclusion is left for Ref. [18]. But we can see that the neutral states, h and H, are not excluded by the SM Higgs bounds obtained using the standard channels, due to a combination of them being heavy plus the fact that other decay channels are open for their decays. In particular, the decay channels

$$(h, H) \to AA, \qquad (h, H) \to AZ \tag{7}$$

are competitive with $(h, H) \to WW$ and $(h, H) \to ZZ$, the channels that drive the bounds on large SM Higgs masses. In order to illustrate this we show in Fig. 3 the branching fractions of the lightest CP-even neutral scalar h for the

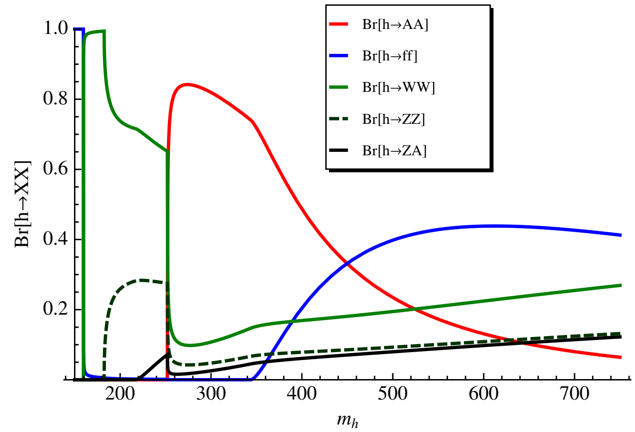

FIG. 3 (color online). Branching ratios for the lightest CP-even h vs m_h, for $\tan\beta = 0.8$, $\alpha = -0.006$, and $m_A = 126$ GeV. Here f refers to all SM fermions, and the couplings to them are type II.

type II case. The branching ratios are computed for $m_A = 126$ GeV and $\tan\beta = 0.8$, which would correspond to a diphoton signal at about the same level of the SM Higgs, as can be seen in Fig. 1. We used a negligibly small value of the mixing angle $\alpha = -0.006$ which comes from the typical parameter space studied here and also is typical in the models presented in the next section. This results in an important coupling of h to the top quark, making $t\bar{t}$ an important decay channel above threshold. More importantly, we see that the ZZ and WW channels used by ATLAS and CMS to put bounds on the mass of the Higgs now have to compete with the AA and AZ channels. In fact, the LHC bounds on Higgs searches do not apply to the THDM h once its mass is above AA threshold or about 250 GeV. The ZZ only becomes more important than the AA channel for rather large masses.

Similarly, we plot the branching fractions for the decays of the heavier CP-even state, H, in Fig. 4. Thus, we see that

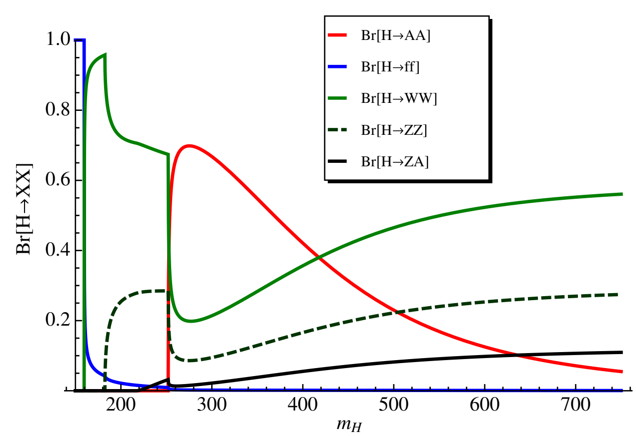

FIG. 4 (color online). Branching ratios for the heaviest CP-even, H, vs m_H, for $\tan\beta = 0.8$, $\alpha = -0.006$, and $m_A = 126$ GeV. Here f refers to all SM fermions, and the couplings to them are type II.

the search for the neutral states in the inverted THDM should include, in addition to the ZZ and WW channels, the $(h, H) \to AA \to 4b's$ and $(h, H) \to AZ \to b\bar{b}\ell^+\ell^-$. In both cases the A should be reconstructed to have the same mass as the one in the $\gamma\gamma$ channel, which would provide a nice confirmation of the model [18].

The spectrum of the inverted THDM presented here, with a light pseudoscalar with $m_A \sim 125$ GeV, and heavy masses for the scalars $(h, H, H^\pm)$ in the (500–600) GeV range, is compatible with tree-level unitarity. To check this we test the parameters of the model by calculating the scattering matrix of scalar interactions. We require that the eigenvalues of this matrix be smaller than 16π, corresponding to saturation. All the points in the parameter space considered here satisfy this constraint [20]. On the other hand, the large masses for the scalars point to the existence of a strongly coupled sector which should appear not too far above the scalar masses. Thus, this scenario for the scalar spectrum of the THDM should be accompanied by new states not too far above the scalar masses, such as new gauge bosons and/or new fermions. We will present such an example in the next section.

III. A MODEL OF THE INVERTED THDM

Although, in principle, the inverted THDM spectrum can always be considered as a possibility, it is interesting to ask whether such a spectrum can be obtained dynamically. The typical MSSM scalar spectrum requires a light scalar, so the pseudoscalar is rather heavy, or much lighter as in the next-to-minimal supersymmetric standard model [21]. A technipion in technicolor models could be light [22], but the rest of the THDM is missing. In Ref. [14] it was shown that this inverted THDM is precisely obtained in theories where the condensation of a fermion sector leads to EWSB. In these models the new chiral fermions feel a strong interaction at the few TeV scale, condensing and breaking the electroweak symmetry. Since both the up-type and down-type right-handed fermions condense with the left-handed doublet, the resulting scalar spectrum at low energies is that of the THDM. Furthermore, the presence of a Peccei-Quinn symmetry in the fermion theory, only broken by the new interaction's instanton effects, guarantees that the CP-odd state A is the lightest state of the scalar spectrum. The rest of the spectrum, given by $(h, H, H^\pm)$, is much heavier, as is expected from the condensation of fermions with a cutoff scale in the multi-TeV region. The scalars are also somewhat degenerate, giving this inverted THDM a very interesting phenomenology at the LHC. The bounds from electroweak precision measurements vary somewhat in the presence of the new fermions, but it is still possible to have agreement with them, as shown in Ref. [14].

The simplest model for the condensing fermions is to assume that they are quarks belonging to a fourth generation [23]. However, this assumption, together with the hypothesis that the diphoton signal comes from the pseudoscalar A, is in great tension with the recent LHC data. In order to see this, we first consider the contributions of the fourth-generation quarks to the ggA vertex. The dynamics of the condensation naturally selects the type II scheme for the fourth generation [14], so it is natural to adopt it for all four generations. With the addition of the fourth-generation quarks, now Eq. (1) for the $gg \to A$ production cross section reads

$$\frac{\sigma_A}{\sigma_h^{\rm SM}}|I_S(\tau_t)|^2 = \frac{9}{4}|\cot\beta I_A(\tau_t) + \tan\beta I_A(\tau_b) + \cot\beta I_A(\tau_{t'}) + \tan\beta I_A(\tau_{b'})|^2, \quad (8)$$

where the last two terms are the contributions of the t' and b' fourth-generation quarks, respectively. As pointed out in Ref. [14], the condensation models typically select $\tan\beta \simeq 1$. Thus, we see that the production cross section is greatly enhanced, by a factor of about 9 at leading order. On the other hand, and unlike the case for the Higgs in the presence of a fourth generation [24], the $A \to \gamma\gamma$ is not suppressed in the presence of the fourth generation since there is no W contribution against which to cancel. To the numerator in the expression (5), we must now add

$$N_c q_U^2 \cot\beta I_A(\tau_{t'}) + N_c q_D^2 \tan\beta I_A(\tau_{b'}). \quad (9)$$

As a result, the $A \to \gamma\gamma$ branching ratio is also enhanced. All in all, the $\sigma \times {\rm BR}(gg \to A \to \gamma\gamma)$ in the fourth-generation model is larger than the SM analogous rate by a factor of 10 for $\tan\beta = 1$ and $m_A = 126$ GeV, which is excluded by the LHC data. Although not motivated by the condensation models, we could consider type I THDM with a fourth generation. The resulting ratio of $\sigma \times {\rm BR}(gg \to A \to \gamma\gamma)$ to the corresponding SM prediction for the Higgs is plotted as the dashed line in Fig. 5. We see

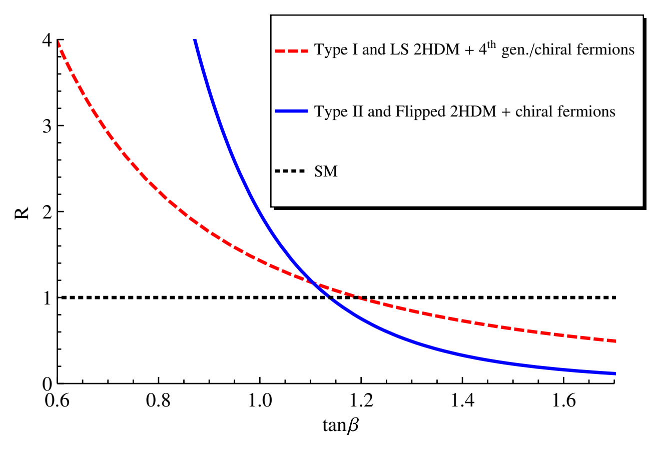

FIG. 5 (color online). The ratio $R = \sigma \times {\rm BR}(gg \to A \to \gamma\gamma)/\sigma \times {\rm BR}(gg \to h \to \gamma\gamma)$ vs $\tan\beta$, for $m_A = 126$ GeV, in a model with colorless chiral fermions. The line is for the type II scheme. The horizontal dashed line corresponds to $\sigma \times {\rm BR}(gg \to A \to \gamma\gamma)$, equal to the prediction for this process mediated by the SM Higgs.

that for this case, it is possible to generate rates consistent with or slightly above the SM rates for a Higgs of the same mass, as long as $\tan\beta \lesssim 1.1$.

Another possibility for generating the inverted THDM spectrum is to just consider that the condensing fermions, although chiral, do not carry color [14]. In this case the strong interaction responsible for fermion condensation is not related to $SU(3)_c$. With this assumption, the new heavy fermions do not contribute to the ggA vertex, leaving this to be just as in Eq. (1) for the three-generation case. On the other hand, the fact that the new fermion condensates break the electroweak symmetry implies that they must have couplings to the photon, and therefore they will contribute to the $A\gamma\gamma$ vertex. If for concreteness we assume that the exotic fermions have the same charges as the up and down quarks, and that the new strong interaction requires that there be N_f copies of them [e.g. the new interaction is $SU(N_f)$], we can make a definite prediction for the new fermion contributions to the $\Gamma(A \to \gamma\gamma)$ by just using (5) supplemented by (9), substituting N_c by N_f. As an example, the solid line of Fig. 5 plots the ratio R for a model with colorless chiral fermions, including a doublet and two singlets with the same hypercharges as the fourth-generation quarks, and for $N_f = 3$. Although the new fermions are typically heavy ($m_f \simeq 600$ GeV), the results depend little on the value of their mass as long as they are significantly heavier than the top quark. We see that, although the curve is somewhat shifted upwards with respect to the three-generation THDM case of Fig. 1, it is still possible to obtain cross sections for the $\gamma\gamma$ channel that are of order of the SM Higgs ones, or even somewhat larger as long as $\tan\beta \simeq O(1)$, which is the value selected by the dynamics of the fermion condensation models.

IV. CONCLUSIONS

The ATLAS and CMS data are beginning to probe the scalar sector of the SM, the least known sector of the theory. It is important that we consider alternatives to the one doublet case, and that we study the resulting phenomenology in light of the upcoming increased data samples. In this paper we have studied the possibility that the diphoton signal observed in ATLAS and CMS is due to the pseudoscalar, CP-odd state A of a THDM, where it is the lightest scalar in the spectrum. If this were the case, no ZZ or WW signal should be confirmed at the diphoton invariant mass, $m_A \simeq 125$ GeV. Furthermore, the VBF-enriched channels of diphoton production, studied by both collaborations with different methods, should be accounted for only from the gluon fusion process. These channels are new, and the excesses found are still small. But if they are confirmed with more data, they would put severe constraints on their interpretation as coming purely from gluon fusion. Depending on where the results of ATLAS and CMS settle in these channels, it may or may not be possible to accommodate the THDM interpretation. For this, more data are needed.

We computed the signal cross section into diphotons and showed that it could explain the observations in the diphoton channel at the LHC for small values of $\tan\beta$, as it can be seen in Fig. 1. Given these values of $\tan\beta$, we predict the branching fractions for the CP-even neutral states, shown in Figs. 3 and 4, and show that the new decay modes $(h, H) \to AA$ and $(h, H) \to AZ$ are competitive enough so that these states, if heavy enough, are not affected by the LHC bounds on the SM Higgs. In general, it is enough for CP-even masses to be above the AA threshold, or $\simeq 250$ GeV, for them not to be bound by the existing LHC data up to rather very large values. The confirmation of this scenario requires the search for these decay modes in $AA \to 4b's$ and $AZ \to \bar{b}b\ell^+\ell^-$ [18].

The inverted THDM spectrum can be dynamically generated by models where chiral fermions interact strongly with a new interaction at the TeV scale, leading to their condensation and EWSB [14]. We showed that the LHC diphoton signal can only be accommodated in these models if the new fermions are either colorless, so as not to significantly affect the production vertex, or in type I scenarios which are not motivated in fourth-generation condensation models [14].

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