Is the LHC Observing the Pseudo-scalar 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 pseudo-scalar 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 $ZZ^{(*)} \rightarrow 4\ell$ and $WW^{(*)} \rightarrow 2\ell 2\nu$ 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 of 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 $h \rightarrow \gamma\gamma$ channel, which would point to a Higgs mass of $m_h \simeq 126$ GeV. For ATLAS [3] the local significance of this excess is of 2.8 $\sigma$, whereas for CMS [4] it is about 3.0 $\sigma$. ATLAS [5] and CMS [6] also observe modest excesses in the $h \rightarrow ZZ^{(*)} \rightarrow 4\ell$ channel with local significance not much above 2.0 $\sigma$, as well as 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 2.7 $\sigma$ excess. Finally, ATLAS [5] has also search for similar signals, which are enhanced by a cut in the transverse momentum of the diphoton system. They found an excess of about 3 $\sigma$. Thus, taken all combined, both experiments appear to coincide in the existence of a Higgs excess somewhere around (124 – 126) GeV, while differing on the invariant mass of the two-significant excesses observed in $ZZ^{(*)} \rightarrow 4\ell$, as well as in $WW^{(*)} \rightarrow \ell^+\nu\ell^-\bar{\nu}$.

In this paper we consider the possibility that the excess in the $\gamma\gamma$ channel is real, but that it is caused by the pseudo-scalar 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] and [8], which find excesses in regions that favor VBF and disfavor gluon fusion. However, the excesses are still small (3 $\sigma$ and 2.7 $\sigma$ for CMS and ATLAS respectively), and the techniques used new. In any case, the pure gluon fusion interpretation is still compatible with both the ATLAS and CMS data at about the 3 $\sigma$ level [12]. So more data is 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 of where un-confined fermions condense to break the electroweak symmetry dynamically [13], due to the existence of an approximate Peccei-Quinn symmetry which keeps the pseudo-scalar state light in comparison t the rest of the scalars.

Whatever the origin of this scalar spectrum, the first signal of it would be the observation of $A \rightarrow \gamma\gamma$ if $m_A < 130$ GeV, just as is the case of the SM Higgs. The production cross section times branching ratio for $gg \rightarrow A \rightarrow \gamma\gamma$ need not be the same as that of the SM Higgs. In fact the ATLAS signal is somewhat larger that 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).

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 [14, 15], depending on the choice of doublet responsible of 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 $\alpha$ between the neutral CP-even scalars. However, the couplings of $A$ to fermions only depend on $\beta$. For instance, the couplings of $A$ to all up-type quarks in all four schemes is cot $\beta$. On the other hand, its couplings to down-type, go like cot $\beta$ in the type I and lepton-specific case, 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 $\sigma \times BR(gg \rightarrow A \rightarrow \gamma \gamma)$ in the inverted THDM. The production cross section for $gg \rightarrow A$ is given by [16]

$$\sigma_A = \frac{9}{4} \left| \cot \beta I_A(\tau_b) + \xi(\beta) I_A(\tau_b) \right|^2 / |I_S(\tau_b)|^2, \quad (1)$$

where $\sigma_{h}^{SM}$ is the SM $gg \rightarrow h$ production cross section, the functions in $[1]$ are given by:

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

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

and

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

expressed in terms of the variable $\tau_f \equiv m_A^2 / m_h^2$. 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 $A \rightarrow \gamma \gamma$ decay width normalized by the SM $h \rightarrow \gamma \gamma$ width, obtaining

$$\frac{\Gamma(A \rightarrow \gamma \gamma)}{\Gamma^{SM}(h \rightarrow \gamma \gamma)} = 4 \left| \frac{N_c q_f^2 \cot \beta I_A(\tau_t) + N_c q_f^2 \xi(\beta) I_A(\tau_b)}{N_c(4/3)(q_U^2 I_S(\tau_t) + q_D^2 I_S(\tau_b))} \right|^2, \quad (5)$$

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

$$I_W(\tau) = -\frac{2\tau^2 + 3\tau + 3(2\tau - 1)f(\tau)}{\tau^2}. \quad (6)$$

We are now ready to compute the $\sigma \times BR(gg \rightarrow A \rightarrow \gamma \gamma)$ 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 Figure 1 we plot the ratio $R = \sigma \times BR(gg \rightarrow A \rightarrow \gamma \gamma) / \sigma \times BR(gg \rightarrow h \rightarrow \gamma \gamma)$ 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 $\sigma \times BR(gg \rightarrow A \rightarrow \gamma \gamma)$ equal to the prediction for this process mediated by the SM Higgs.

![Figure 1](image-url)
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]. It is also generally compatible with precision electroweak constraints. In Figure 2 we show the contributions to the $S$ and $T$ parameters of 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 $m_{h^\text{eff}} = 117$ GeV. We can see from Figure 2 that there are many solutions within the inverted THDM spectrum that are compatible with electroweak bounds. The points falling out of the allow 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

\[
(h, H) \to AA \quad (h, H) \to AZ
\]

that 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 Figure 3 the branching fractions of the lightest CP-even neutral scalar $h$ for the type II case. The branching ratios are computed for $m_A = 126$ GeV and $m_H = 125$ GeV, which represent a diphoton signal at about the same level of the SM Higgs, as can be seen in Figure 4. 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 the branching fractions for the decays of the heavier CP-even state, $H$, in Figure 4. Thus, we see that the search for the neutral states in the inverted THDM should include, in addition to the $ZZ$ and $WW$ channels, the $(h, H) \to AA \to 4b\ell^+\ell^-$ 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 pseudo-scalar with \( m_A \sim 125 \text{ GeV} \), and heavy masses for the scalars \( (h, H, H^\pm) \) in the \( (500 = 600) \text{ GeV} \) 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 \( 16\pi \), 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.

**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 spectrum can be obtained dynamically. The typical MSSM scalar spectrum requires a light scalar, so the pseudo-scalar is rather heavy, or much lighter as in the NMSSM \[21\]. A techni-pion in techni-color models could be light \[22\], but the rest of the THDM is missing. In Ref. \[13\] 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. \[13\].

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 pseudo-scalar \( A \) 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 \[13\], 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

\[
\frac{\sigma_A}{\sigma_{SM}} |I_S(\tau_I)|^2 = \frac{9}{4} \left| \cot \beta I_A(\tau_t) + \tan \beta I_A(\tau_b) \right|^2,
\]

where the last two terms are the contributions of the \( t' \) and \( b' \) fourth-generation quarks, respectively. As pointed out in Ref. \[13\], 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

\[
N_c q^2 \cot \beta I_A(\tau_{t'}) + N_c q_D^2 \tan \beta I_A(\tau_{b'}).
\]

As a result the \( A \to \gamma \gamma \) branching ratio is also enhanced. All in all, the \( \sigma \times BR(gg \to A \to \gamma \gamma) \) in the fourth generation model is larger than the SM analogous rate by a factor of 10 for \( \tan \beta = 1 \) and \( m_A = 126 \text{ GeV} \), 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 \( \sigma \times BR(gg \to A \to \gamma \gamma) \) to the corresponding SM prediction for the Higgs is plotted as the dashed line in Fig. \[3\]. We see 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 [13]. 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 $\Gamma(A \rightarrow \gamma\gamma)$ by just using (5) supplemented by (9) with $N_f$ substituting $N_c$. As an example, in the solid line of Figure 5 we plot 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 Figure 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.

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 coming increased data samples. In this letter we have studied the possibility that the diphoton signal observed in ATLAS and CMS is due to the pseudo-scalar, CP-odd, state $A$ of a THDM where it is the lightest scalar in the spectrum. If this were the case, no $ZZ$ and $WW$ 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 to 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 is 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 Figure 1. Given these values of $\tan \beta$, we predict the branching fractions for the CP-even neutral gauginos should be accounted for only from the gluon fusion process. 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 is needed.

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