Probing Exotic Triple Higgs Couplings for Almost Inert Higgs Bosons at the LHC

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Abstract

In extended Higgs sectors that exhibit alignment without decoupling, the additional scalars are allowed to have large couplings to the Standard Model Higgs. We show that current nonresonant di-Higgs searches can be straightforwardly adapted to look for additional Higgses in these scenarios, where pair production of non-SM Higgses can be enhanced. For concreteness, we study pair production of exotic Higgses in the context of an almost inert two Higgs doublet model, where alignment is explained through an approximate $\mathbb{Z}_2$ symmetry under which the additional scalars are odd. In this context, the smallness of the $\mathbb{Z}_2$ violating parameter suppresses single production of exotic Higgses, but it does not prevent a sizeable trilinear coupling $hHH$ between the SM Higgs ($h$) and the additional states ($H$). We study the process $pp \rightarrow h^* \rightarrow HH$ in the final states of $b\bar{b}b\bar{b}$, $b\bar{b}\gamma\gamma$, and multi-leptons. We find that at the HL-LHC these modes could be sensitive to masses of the additional neutral scalars in the range $130 \text{ GeV} \lesssim m_H \lesssim 290 \text{ GeV}$. 
1 Introduction

The properties of the observed 125 GeV Higgs boson are in close agreement with the Standard Model (SM) predictions. Therefore, any extension of the SM Higgs sector needs to reproduce the SM-like Higgs couplings to the weak gauge bosons and fermions. This can be easily satisfied in a scenario where the observed Higgs boson comes mainly from a SM doublet $H_1$ plus a small deviation from it:

$$H_{SM} = H_1 + \epsilon H_2$$  \hspace{1cm} (1.1)

The second doublet $H_2$ introduces four new degrees of freedom, which can be represented by the following physical states: a charged Higgs $H^\pm$, a neutral scalar $\phi$, and a neutral pseudoscalar $A$. In this scheme, they all have $\epsilon$-suppressed interactions with the SM particles, therefore they are difficult to look for experimentally.

A SM-like Higgs as depicted in Eq (1.1) can be naturally realized in a Two Higgs Doublet Model (2HDM) with an approximate $Z_2$ symmetry, where $H_2$ is odd, and $H_1$ together with the SM fields are even. In the exact $Z_2$-symmetry limit, one possible configuration for the vacuum, that respects the symmetry, is $\langle H_1 \rangle \neq 0$, $\langle H_2 \rangle = 0$. In this case, known as the Inert Doublet Model (IDM), the $Z_2$-odd fields $H^\pm, A, \phi$ are forbidden to decay to final states with only $Z_2$ even fields. \cite{1, 2} In particular, the lightest $Z_2$-odd neutral scalar is completely stable and can provide a dark matter candidate \cite{3–5}.

In IDM, electroweak symmetry breaking (EWSB) directly contributes to the masses of the $Z_2$-odd higgses. In particular, if we take the limit where the mass term of the inert doublet, $m^2 H_2^2$, approaches zero, the masses of the $Z_2$-odd Higgses are solely generated via EWSB, and $m^2_{\lambda, \phi, H^\pm} \sim \lambda v^2$, where $\lambda$ represents the quartic couplings for $Z_2$-conserving terms such as $|H_1|^2 |H_2|^2$. Therefore, large $Z_2$-preserving quartic couplings between $H_1$ and $H_2$ can be responsible for the mass hierarchy between SM and non-SM Higgs bosons. This is an example of non-decoupling electroweak symmetry breaking and motivates the study of triple-higgs couplings $g_{hhH}$, where $h$ is the SM Higgs boson and $H = H^\pm, A, \phi$.

One way to study the coupling $g_{hhH}$ is to look at the decays of the SM-Higgs boson $h$. If the mass of the lightest $Z$-odd Higgs $H$ is less than half the mass of the SM Higgs, $h$ can decay to a pair of on shell $Z$-odd Higgses. Studies on exotic decays of SM Higgs \cite{6–9} largely constrain $g_{hhH}$ in this case. Here we want to focus on the scenario where $H$ is heavy, so that exotic $h$ decays cannot probe $g_{hhH}$ efficiently. Instead, we study $g_{hhH}$ by pair producing $H$ via an off-shell $h$ at the Large Hadron Collider (LHC). Depending on which of the additional scalars ($H^\pm, A, \phi$) is the
lightest $Z$-odd state, different search strategies are required to constrain $g_{hHH}$. If $H^\pm$ is long-lived, searches looking for heavy stable charged particles (HSCP) \cite{10,11} can be sensitive. If the neutral scalar or pseudoscalar is long-lived, searches using missing transverse energy (MET) recoiling against initial state radiation \cite{12–14} can be sensitive. In this work, we relax the exact $Z_2$-symmetry condition, so that small explicit $Z_2$-breaking terms are allowed. With this assumption, the approximately $Z_2$-odd states can start to decay promptly inside the detector, providing a much richer phenomenology at the LHC. \cite{15}

In general, $Z_2$ breaking can be introduced in the Higgs potential via odd powers of $H_2$ or in the Yukawa sector through terms proportional to $H_2$. In Ref. \cite{15}, we introduced $Z_2$ breaking terms in the Higgs potential as well as in the Yukawa sector, so this two sources of $Z_2$-breaking were independent. In this work, we only consider $Z_2$ breaking terms in the scalar potential, for the purpose of staying within the context of the type-I 2HDM. The $Z_2$ breaking terms will introduce a VEV $\langle H_2 \rangle \sim \epsilon v$ and induce $O(\epsilon)$ suppressed couplings for interactions between one $Z_2$-odd field and SM fields.

The smallness of the $Z_2$-violating couplings will not prevent a sizeable triple Higgs coupling $g_{hHH}$, since the latter is dominated by the $Z_2$-preserving quartics. Consequently, as aforementioned, the production mode of interest is di-Higgs production $pp \to h^* \to HH$, where $H = A, \phi, H^\pm$. Since the coupling between $H$ and SM fermions is suppressed by $\epsilon$, the box diagram contributing to this process, dominated by a top loop, can be safely ignored. This leaves di-Higgs production via off-shell SM Higgs as the only important diagram, unlike for SM-Higgs pair production, where the interference between the box diagram and the triangle diagram is important. Given that there exist CMS and ATLAS searches \cite{16} targeting SM di-Higgs production via final states $bb\bar{b}b$ \cite{17,18}, $bb\gamma\gamma$ \cite{19,21}, $bbWW$ \cite{22,23}, and $bb\tau\tau$ \cite{24,25}, thus constraining the SM triple Higgs coupling $g_{hhh}$, an enhanced $g_{hHH}$ coupling could also be constrained by these searches. Here we study the sensitivity of these analyses to the non-SM triple Higgs couplings $g_{hHH}$.

Another reason to look for di-Higgs production via off-shell SM Higgs is that this production mode can be a complementary probe to electroweak pair production (EWPP) in searches looking for almost inert Higgs bosons. \cite{15} Exotic Higgses produced through EWPP (via $W$ or $Z$) are always distinct pair of physical states: $\phi H^\pm$, $AH^\pm$, $\phi A$ or $H^+H^-$. Therefore, if the neutral scalar $\phi$ is the lightest approximately $Z_2$-odd state, pairs of $\phi$ cannot be produced directly via electroweak processes. In this case, $pp \to h^* \to \phi\phi$ could become the most important production channel for exotic Higgses. Since pairs of carged Higgses $H^+H^-$ can be produced through EWPP, in
this work we focus on pair production of the neutral states ($\phi$ and $A$).

The paper is organized in the following way. In Section 2 we present the model, that corresponds to a Type I 2HDM with small $Z_2$-breaking terms, and discuss the production and decay modes of the approximately inert Higgses. In Section 3 and 4 we find current constraints on $g_{h\phi\phi}$ and $g_{hAA}$ coming from SM di-Higgs searches. Section 5 contains our conclusions and outlook.

2 Almost inert Higgs bosons

In this section we present the model and introduce the relevant parameters for the phenomenology. Consider two Higgs doublets that have the following $Z_2$ symmetry:

$$H_1 \mapsto H_1, \quad H_2 \mapsto -H_2.$$  

(2.1)

The most general $Z_2$-conserving potential is

$$V_0 = m_1^2 H_1^\dagger H_1 + \frac{1}{2} \lambda_1 (H_1^\dagger H_1)^2 + m_2^2 H_2^\dagger H_2 + \frac{1}{2} \lambda_2 (H_2^\dagger H_2)^2 + \lambda_3 (H_1^\dagger H_2)(H_2^\dagger H_1) + \lambda_4 (H_2^\dagger H_1)(H_1^\dagger H_2) + \frac{1}{2} [\lambda_5 (H_1^\dagger H_1)^2 + h.c.].$$  

(2.2)

We assume that $m_1^2 < 0$ and $m_2^2 > 0$, such that in the exact $Z_2$-symmetry limit, only $H_1$ gets a VEV. It is instructive to trade the parameters in the Higgs potential with the phenomenological parameters. Assuming all the parameters are real (so there is no explicit CP violation), we trade the parameters in the potential for the following parameters:

$$v, m_h, m_\phi, m_A, m_{H^\pm}, \lambda_2, m_2^2,$$  

(2.3)

where $v$ and $m_h$ are the VEV and mass of the SM Higgs (which resides in $H_1$), while $m_\phi$, $m_A$, and $m_{H^\pm}$ are the masses of the additional neutral scalar, pseudoscalar and charged scalar, all of which reside in the inert doublet $H_2$. This vacuum configuration does not break the symmetry $H_2 \to -H_2$, so the lightest $Z_2$-odd particles are stable. This is also known as the Inert Doublet Model (IDM).

A large mass splitting between $m_{H^\pm}$ and $m_A$ would result in large violations of the $SU(2)$ custodial symmetry, which is very well constrained by electroweak precision tests. Therefore, in this work we assume $m_A = m_{H^\pm}$.

2.1 Di-Higgs production at LHC

We investigate the non-SM di-Higgs production cross section, i.e. $gg \to h^* \to \phi\phi(AA)$, in the exact $Z_2$-symmetry limit (IDM). First, we extract the triple-Higgs
coupling $g_{hHH}$, where $H = \phi, A, H^+$. As discussed in the Introduction, these couplings arise from $\mathbb{Z}_2$-preserving quartic terms in the Higgs potential:

\begin{align}
    g_{hH^+H^-} &= v\lambda_3, \\
    g_{hAA} &= 2v(\lambda_3 + \lambda_4 - \lambda_5), \\
    g_{h\phi\phi} &= 2v(\lambda_3 + \lambda_4 + \lambda_5),
\end{align}

(2.4)

where the factor of 2 in $g_{hAA}$ and $g_{h\phi\phi}$ will compensate the symmetry factor in the relevant terms of the Lagrangian where identical particles appear\footnote{Similarly, the SM triple higgs coupling $g_{hhh}$ will be $3\lambda_1 v$.}. To see what range of values $\lambda_3$, $\lambda_4$, and $\lambda_5$ can take, it is instructive to rewrite them as:

\begin{align}
    \lambda_3 v^2 &= 2(m_H^2 - m_2^2), \\
    \lambda_4 v^2 &= \lambda_5 v^2 + 2(m_A^2 - m_{H^\pm}^2), \\
    \lambda_5 v^2 &= m_\phi^2 - m_A^2.
\end{align}

(2.5)

Therefore, in the IDM we have:

\begin{align}
    g_{hH^+H^-} &= 2(m_H^2 - m_2^2)/v, \\
    g_{hAA} &= 4(m_A^2 - m_2^2)/v, \\
    g_{h\phi\phi} &= 4(m_\phi^2 - m_A^2)/v,
\end{align}

(2.6)

where $m_2^2 > 0$. In the case where $m_2^2 = 0$, the masses of the exotic Higgs are solely generated by the quartic couplings upon EWSB, therefore these couplings must scale with the mass squared of the additional scalars. This will partially compensate for the drop in the gluon parton distribution function in the large mass regime. As an example, Fig. 1 shows the cross section of $gg \rightarrow h^* \rightarrow \phi\phi$ for $m_2^2 = 0$ (red solid) and $m_2^2 = (100 \text{ GeV})^2$ (blue solid). One can see that for different choices of $m_2^2$, the cross section can change for orders of magnitude at a give mass. However, the coupling soon reaches $4\pi$ when $m_\phi \gtrsim 400$ GeV. Therefore, the theory reaches a non-perturbative regime and a leading order estimate of the cross section is no longer reliable. Therefore, we restrict ourselves in considering

\begin{equation}
    125 \text{ GeV} < m_{\phi^0} < 400 \text{ GeV}
\end{equation}

(2.7)

For comparison, we also plot the cross section with the coupling set to be a constant equal to the SM triple higgs coupling (black dashed). Its steep fall is due to the fast drop in the gluon parton distribution function.
2.2 Decay of approximately inert Higgs

In IDM, the lightest $Z_2$ odd neutral state can only appear as missing energy at the LHC. However, if the symmetry is not exact, it can decay within the detector, which is what we investigate in this section.

Starting with the Higgs potential, Eq. (2.2), one can introduce $Z_2$ breaking via including odd powers of $H_1$:

$$\Delta V = \Delta m^2(H_2^\dagger H_1 + h.c.) + \Delta \lambda(H_2^\dagger H_2)(H_1^\dagger H_1 + h.c.) + \Delta \lambda'(H_1^\dagger H_1)(H_2^\dagger H_1 + h.c.),$$

where $Z_2$-breaking terms are all of $O(\epsilon) \ll 1$. Now the ‘inert’ doublet $H_2$ gets a VEV of $O(\epsilon)$:

$$\langle H_2 \rangle \equiv \tilde{v}_2 = -\frac{\Delta m^2/v^2 + \Delta \lambda/2}{m_2^2/v^2 + (\lambda_3 + \lambda_4 + \lambda_5)/2}v + O(\epsilon^2),$$

The ratio of the VEVs of the two doublets defines the conventionally adopted parameter $\tan \beta$ in 2HDM. In the approximate $Z_2$ limit, it is given by $v/\tilde{v}_2$, and tend to be large.\footnote{In the almost inert limit, $H_{1,2} \approx \Phi_{2,1}$, where $\Phi$ represents the usually defined $Z_2$ basis. Consequently, $\tan \beta = v_2/v_1 \approx \tilde{v}_1/\tilde{v}_2$ in this limit. See the appendix of Ref. \cite{15} for more details.}

The fields can be parametrized around their VEVs as

$$H_i = \begin{pmatrix} H_i^+ \\ \frac{1}{\sqrt{2}}(\tilde{v}_1 + h_i^0 + iA_i) \end{pmatrix}.$$
All the physical fields receive an $\mathcal{O}(\epsilon)$ correction with respect to their $Z_2$-symmetric limit. The neutral pseudoscalar and charged components of the doublets mix to generate the physical fields $A$ and $H^\pm$:

$$A = A_2 + \epsilon_A A_1 + \mathcal{O}(\epsilon^2),$$

$$H^+ = H_2^+ + \epsilon_A H_1^+ + \mathcal{O}(\epsilon^2),$$

where

$$\epsilon_A = -\frac{\bar{v}_2}{v} = \mathcal{O}(\epsilon).$$

Note that $\epsilon_A$ is related to the usual $\tan \beta$ parameter through $\epsilon_A \approx -1/\tan \beta$. The physical scalar fields are

$$\left(\begin{array}{l} h \\ \phi \end{array}\right) = \left(\begin{array}{cc} 1 & \epsilon_h \\ -\epsilon_h & 1 \end{array}\right) \left(\begin{array}{c} H_1^0 \\ H_2^0 \end{array}\right) + \mathcal{O}(\epsilon^2),$$

with

$$\epsilon_h = \frac{1}{m_h^2 - m_\phi^2} \left[\frac{\bar{v}_2}{v} \left(m_\phi^2 - 2m_{H^\pm}^2 + \lambda_3 v^2\right) + \Delta \lambda v^2\right] = \mathcal{O}(\epsilon).$$

Note that $\epsilon_h$ is equivalent to $\alpha (\alpha - \pi/2)$ when $m_h < m_\phi$ ($m_h > m_\phi$), where $\alpha$ is another conventionally adopted parameter in 2HDM. From the kinetic terms, we find how $Z_2$-odd fields couple to SM fields:

$$g_{\phi V} \propto \epsilon_V, \quad g_{AZh} \propto \epsilon_V, \quad g_{H^\pm W^\mp h} \propto \epsilon_V,$$

where

$$\epsilon_V \equiv \epsilon_A + \epsilon_h.$$

The limit $\epsilon_V \to 0$ corresponds to the so-called alignment limit in 2HDM.

The $Z_2$-breaking terms in the Higgs potential also induces $\mathcal{O}(\epsilon)$ $Z_2$-violating terms in the Yukawa sector. In the exact $Z_2$-symmetry limit, only $H_1$ can have Yukawa couplings. Therefore, this is a Type I 2HDM. After $Z_2$ breaking, the approximate $Z_2$-odd scalars couple to SM fermions via $\mathcal{O}(\epsilon)$ VEV and/or mass mixing:

$$g_{\phi f \bar{f}} \propto \epsilon_V - \epsilon_A, \quad g_{A f \bar{f}} \propto \epsilon_A, \quad g_{H^\pm f \bar{f}} \propto \epsilon_A.$$

In summary, the following set of phenomenological parameters can be used to discuss the approximate $Z_2$ Higgses:

$$v, m_h, m_\phi, m_A, m_{H^\pm}, m_2^2, \epsilon_V, \epsilon_A,$$

where we have left out $\lambda_2$ and $\Delta \lambda'$. Both terms have more than two powers of $H_2$. Since $\langle H_2 \rangle = \bar{v}_2 = \mathcal{O}(\epsilon)$, these terms are not important at the tree level. Fig. 2 shows examples of $\phi$'s branching ratios. It can be seen that adjusting the relative size of $\epsilon_V$ to $\epsilon_A$ radically changes $\phi$'s dominant decays.
Fig. 2: Examples of dominant decay modes of the neutral approximately $Z_2$-odd Higgses as a function of their mass. The small $Z_2$ breaking parameters $\epsilon_V, \epsilon_A$ roughly control the size of their decays to SM weak gauge bosons and fermions, respectively. The subplot in 2d shows that $\phi^0$ can still undergo prompt decay even when $\epsilon_V \sim \epsilon_A \to 0$ via loop induced processes, which are controlled by $\Delta \lambda'$. 
A special case: $\phi \to \gamma\gamma$

If $\phi$ is the lightest approximate $\mathbb{Z}_2$-odd Higgs, in the limit that $\epsilon_V \to 0, \epsilon_A \to 0$, we expect that $\phi$ will never decay inside the detector. However, in this scenario, where the tree-level decays of $\phi$ are suppressed, one can no longer ignore loop induced processes. In particular, the quartic interaction $\Delta \lambda' |H_2|^2 (H_1 H_2^\dagger + h.c.)$ contributes a $\mathbb{Z}_2$-violating coupling $g_{\phi H_H^-} \sim \Delta \lambda' v$. In the limit that $\Delta \lambda' \gg \epsilon_V \sim \epsilon_A \to 0$, $\phi$’s dominant decay can become $\gamma\gamma/\gamma Z$, induced by an $H^+$ loop, with one caveat. The inclusion of a term like $\Delta \lambda' |H_2|^2 (H_1 H_2^\dagger + h.c.)$ necessarily introduces the $\mathbb{Z}_2$ breaking term $\Delta m^2 H_1^\dagger H_2$ back at one loop. If $\Lambda$ is the UV cutoff of our model, $\Delta m^2$ can be estimated to be $\sim \frac{\Delta \lambda'}{16\pi^2} \Lambda^2$. From Eq. (2.9) (2.12), the limit $\Delta \lambda' \gg \epsilon_V \sim \epsilon_A \to 0$ is only achievable if $\Lambda \lesssim 4\pi v \sim 3$ TeV.

Fig. 2d shows an example of $\phi$’s branching ratios in this region of the parameter space. In the subplot, we also show the total width of $\phi$. Therefore, even in the limit that $\epsilon_V \sim \epsilon_A \to 0$, $\phi$ can still decay within the detector. More importantly, it decays with the spectacular signal involving multiple photons, which have much smaller SM backgrounds.

3 Constraining $h\phi\phi$ coupling

In this section, we begin detailed benchmark studies under the assumption of different schemes of $\mathbb{Z}_2$ breaking parameters, as studied in Section 2.2. Even though all the $\mathbb{Z}_2$-breaking parameters are assumed to be small, $\phi$’s main decay modes can change drastically depending on the relative size between them (cf. Fig. 2). Since there are no a priori assumptions on the origin of $\mathbb{Z}_2$ breaking, we have to study all possible qualitatively different decays of $\phi$ at the LHC. Fig. 2 shows the main decay modes for different regions of the parameter space. Each subsection focuses on a scenario where $\phi$ dominantly decays to only one or two types of SM particles, based on which, limits of $g_{h\phi\phi}$ as a function of mass are drawn. If the LHC is insensitive to such ideal scenarios, we may conclude that it is not going to be sensitive to $g_{h\phi\phi}$ in more general scenarios.

For all the following studies, we generate the parton level events for signal with MadGraph5 [27], followed by hadronization and shower by Pythia8 [28] and detector simulation by Delphes3 [29]. Each subsection assumes a distinct relative size between $\epsilon_V$ and $\epsilon_A$ and discusses possible constraints on $g_{h\phi\phi}$ arising from the dominant decays.
3.1 $\epsilon_A \gg \epsilon_V : \phi \phi \rightarrow (b\bar{b})(b\bar{b})$

As shown in Fig (2b), just like the SM Higgs, $\phi$ can dominantly decay to a pair of $b$ quarks. The signature of interest here is four $b$ jets. Since both ATLAS and CMS have performed SM di-Higgs searches in this channel \cite{30,31}, our goal is to see if these standard searches can be sensitive to the almost inert heavy Higgses. We follow the selection and cuts from ATLAS’s di-Higgs search at $36^{-1}$ fb \cite{30}, excepting one cut involving the b-tagging score. The interested reader can find the details of these cuts in Appendix A.1. The right panel of Fig. 3 shows the efficiencies of both signal reconstruction and final selection, where reconstruction simply requires 4 $b$-tagged jets with $p_T > 40$ GeV. We see from the figure that the overall efficiency is maintained at $\sim 5\%$ in spite of different masses.

The SM backgrounds for this channel are multijets, hadronic $t\bar{t}$ and semi-leptonic $t\bar{t}$. Both CMS and ATLAS use mainly data driven methods to model them. Therefore, we simply extrapolate the backgrounds at $36$ fb$^{-1}$ from \cite{30} to draw limits on coupling vs mass for higher luminosities. This is a conservative estimate because a real analysis is expected to be more efficient reducing backgrounds for higher masses. It is clear from Fig. 3 that the 4$b$ channel of SM di-Higgs searches can be important in constraining $g_{h\phi\phi}/v$, even though they are not optimised to target the exotic Higgs.
Upper Bounds on $g_{h\phi\phi}/v$

![Graph showing upper bounds on $g_{h\phi\phi}/v$ vs $m_{\phi}$ with lines for multi-lepton, 300/fb (95% CL) and multi-lepton, 3000/fb (95% CL) at $m_2^2=0$.]

Multi-lepton Signal Efficiency

![Graph showing multi-lepton signal efficiency vs $m_{\phi}$ with efficiency in percentage.]

3.2 $\epsilon_A \lesssim \epsilon_V : \phi \to (WW)(WW) \to (\ell\nu\ell\nu)(\ell\nu jj)$

As shown in Section 2.2, $\phi$'s coupling to SM fermions, $g_{\phi f\bar{f}}$, is proportional to the difference between $\epsilon_V$ and $\epsilon_A$. Since these two parameters are independent, there can be a scenario (Fig. 2a) where they cancel each out identically, leaving $\phi$ mainly decaying to $W^+W^-$. Out of all the possibilities of the $4W$ decay, we focus on the final state with 3 leptons and 2 jets. This is a relatively clean signature, where the additional jets allow triggering. However, one cannot easily reconstruct the decay chain, therefore this search reduces to basic lepton counting, which can be potentially covered by several multi-lepton searches at ATLAS and CMS, for example [32, 33].

Our goal here is to investigate whether these standard multi-lepton searches can be sensitive to the almost inert Higgs. We follow the 36 fb$^{-1}$ three lepton search adopted in [34], of which details of the cuts are spelled out in Appendix A.3.

For this channel, the SM backgrounds are mainly non-prompt leptons, diboson and $t\bar{t}V$. Again, we simply extrapolate the backgrounds at 36 fb$^{-1}$ from [34] to draw 95% CL contours on coupling vs mass for higher luminosities. As can be seen in Fig. 4, this search only starts being effective in probing $g_{h\phi\phi}$ at the high luminosity LHC.
3.3 $\Delta \lambda' \gg \epsilon_V \sim \epsilon_A$ : $\phi \phi \rightarrow (b\bar{b})(\gamma\gamma)$

In the limit where $\Delta \lambda' \gg \epsilon_A \sim \epsilon_V$, tree level decays of $\phi$ are suppressed, and loop-mediated decays become important. The interaction $\Delta \lambda'|H_2|^2 (H_1^\dagger H_2 + h.c.)$ contributes to the $\mathbb{Z}_2$-violating vertex $g_{\phi H^+ H^-} \sim v \Delta \lambda'$, allowing the decays $\phi \rightarrow \gamma\gamma$ and $\phi \rightarrow Z\gamma$. Therefore, the final states $(\gamma\gamma)(\gamma\gamma)$ and $(\gamma\gamma)(Z\gamma)$ could be sensitive in this scenario, but the analysis would require an estimation of the background due to fake photons from multijets, which is beyond the scope of this work. Additionally, if $\epsilon_A$ is small, but not zero, we could also have a sizeable decay to $b\bar{b}$, as shown in Fig. 2d. In this scenario, di-Higgs searches in the $b\bar{b}\gamma\gamma$ final state could be relevant. On the other hand, if $\epsilon_A \sim \epsilon_V \rightarrow 0$, depending on the value of $\Delta \lambda'$, the exotic Higgs could be long lived. In the latter case, delayed photon searches could be sensitive, but these analyses cannot be straightforwardly recast and a dedicated analysis, beyond the scope of this work, would be required. Therefore, for this scenario we focus on the prompt $b\bar{b}\gamma\gamma$ final state.

In particular, we study the potential sensitivity to $g_{h\phi\phi}$ of current ATLAS and CMS di-Higgs searches, in the $b\bar{b}\gamma\gamma$ final state. For concreteness, we consider the scenario depicted in Fig. 2d, where $\Delta \lambda' = 0.1$, $\epsilon_V = 0$ and $\epsilon_A = 10^{-4}$.

We perform a Monte Carlo simulation of the signal, for different masses of the neutral scalar $\phi$, and apply the cuts from Ref. [35], excepting for a modification of the invariant mass cuts. Contrary to the cuts for the $b\bar{b}b\bar{b}$ final state in the previous section, these cuts are optimized for the mass of the SM Higgs, so the efficiency drops rapidly as we depart from the SM Higgs mass. Therefore, we use the following invariant mass cuts:

$$|m_{\gamma\gamma} - m_{\phi}| < 3 \text{ GeV}, \quad |m_{b\bar{b}} - m_{\phi}| < 25 \text{ GeV},$$

which reduce to the cuts from Ref. [35] for $m_{\phi} = 125$ GeV. The right panel in Fig. 5 shows the total signal efficiency as a function of $m_{\phi}$, for $m_{\phi} = 125$ GeV (SM di-Higgs cuts) and for variable $m_{\phi}$ (modified cuts). The efficiencies obtained from our Monte Carlo simulation have been scaled to match the efficiency obtained in Ref. [35] for the point $m_{\phi} = 125$ GeV ($\approx 2.89\%$).

Using the signal efficiency with modified cuts, and the branching fractions shown in Fig. 2d, we find 95% CL limits on $g_{h\phi\phi}$, assuming the same backgrounds as for the SM di-Higgs search from Ref. [35]. This is a conservative estimate because the continuum diphoton background is expected to be much smaller for higher invariant mass regions. We find that current di-Higgs searches in the $b\bar{b}\gamma\gamma$ final state can be straightforwardly adapted to be sensitive to the exotic triple Higgs coupling $g_{h\phi\phi}$ for $m_{\phi} \lesssim 250$ GeV (300 fb$^{-1}$) and $m_{\phi} \lesssim 300$ GeV (3000 fb$^{-1}$).
Fig. 5: Left panel shows the upper limits on $g_{h\phi\phi}/v$ from $pp \rightarrow h \rightarrow \phi\phi \rightarrow (b\bar{b})(\gamma\gamma)$ at the LHC with 300 fb$^{-1}$ and HL-LHC (3000 fb$^{-1}$). We also include the value of $g_{h\phi\phi}/v$ in aIDM for $m_2^2 = 0$, which is given by $4m_\phi^2/v^2$. Right panel shows the signal efficiency with the cuts from Ref. [35] and using $m_\phi$-dependent invariant mass cuts.

4 Constraining $hAA$ coupling

The lightest almost inert neutral state can also be the pseudoscalar $A$. How $A$ decays largely depends on the relative size between the $Z_2$-breaking parameters $\epsilon_A$ and $\epsilon_V$. For instance, as shown in Fig. 2c for $\epsilon_V \approx \epsilon_A$, the decay of $A$ is dominated by $b\bar{b}$ for $m_A < 125$ GeV and $Zh$ for $m_A > 125$ GeV. Therefore, the 4$\bar{b}$ channel (Sec. 3.1) studied previously also constrains $g_{hAA}$. The new ingredient here is that for high masses of $A$, the pair produced $A$s can undergo cascade decays to two $Z$ bosons and two SM $h$s. If all of these particles are produced on shell, it is possible to reconstruct the entire decay chain.

Each subsection below assumes a distinct relative size between $\epsilon_V$ and $\epsilon_A$ and discuss possible constraints on $g_{hAA}$ arising from the dominant decays.

4.1 $\epsilon_A \gtrsim \epsilon_V : AA \rightarrow (b\bar{b})(b\bar{b})$

As mentioned before, for a wide region of the parameter space, $A \rightarrow b\bar{b}$ is one of the main decay modes of $A$. In general, this is the dominant decay mode (almost 100%) for $\epsilon_A \gtrsim \epsilon_V$. Also, for $\epsilon_A \approx \epsilon_V$, this mode dominates for $m_A \lesssim 215$ GeV, as shown in Fig. 2c. Therefore, the limits obtained in Fig. 3 for $g_{h\phi\phi}$, in the $(b\bar{b})(b\bar{b})$ final state, are also valid for $g_{hAA}$. 

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4.2 \( \varepsilon_A \lesssim \varepsilon_V : AA \rightarrow (Zh)(Zh) \)

The final states coming from a pair of cascade decays can be rich and complex. In order to reconstruct the decay chain, we require \( h \rightarrow b\bar{b} \) for both \( hs \) and at least one \( Z \) undergoing leptonic decay. Depending on how the other \( Z \) decays, there arise multiple decay channels that can be potentially interesting.

The cleanest channel is the one with both \( Z \) decaying leptonically, but it also suffers from a tiny signal rate. The second cleanest channel is when \( ZZ \rightarrow (\ell\ell)(\nu\bar{\nu}) \), which has a same final state as the dileptonic decay of \( t\bar{t}h \), where \( h \rightarrow b\bar{b} \). There exist dedicated search on \( t\bar{t}h \) from CMS \(^{38}\), where a lot of the search strategies can be migrated here. However, a crucial difference between \( t\bar{t}h \) and our model is that the former requires a \( Z \) veto which is undesirable in the latter case. In order to constrain the coupling from this channel, we estimate the SM backgrounds based on our own simulation. We also considered \( ZZ \rightarrow (\ell\ell)(jj) \), but the SM background such as \( Z + \text{jets} \) are difficult to reduce.

Therefore we focus on a final state \((\ell\ell b\bar{b})(\nu\nu b\bar{b})\), coming from the cascade decays, \( A \rightarrow Zh \), of pair produced \( A \) at the 13 TeV LHC. The dominating backgrounds include \( t\bar{t} \), single top and \( Z + \text{jets} \), as well as other SM rare processes (cf. Table. 1). We first perform a detailed benchmark study, for which we generate the number of background events that is equivalent to an integrated luminosity of 300 fb\(^{-1}\).

| \( AA \rightarrow (Zh)(Zh) \rightarrow (\ell\ell b\bar{b})(\nu\nu b\bar{b}) \) | \( \sigma(\text{fb}) \) | Initial @ 300 fb\(^{-1}\) | 2\( \ell\)4\( b \) on-Z | \( E_T > 120 \) | MT2\( _W \) > 100 | MT2\( _t \) > 200 |
|---|---|---|---|---|---|---|
| BM1: \( m_A = 220 \) | 0.2 | 60 | 2.5 | 1.9 | 1.4 |
| BM2: \( m_A = 315 \) | 0.08 | 24 | 1.1 | 0.8 | 0.6 |
| \( t\bar{t} + jj \rightarrow (\nu\ell b)(\nu\ell b) + jj \) | 19680 | \( 5.9 \times 10^6 \) | 8330 | 182 | 0.9 |
| \( Z + bbjj, Z \rightarrow \ell\ell \) | 19790 | \( 5.9 \times 10^6 \) | 1275 | 86 | 3.6 |
| \( tW + b\bar{b} \rightarrow (\nu\ell b)(\nu\ell) + b\bar{b} \) | 145 | 43500 | 443 | 14 | 0 |
| \( t\bar{t}V \rightarrow (\nu\ell b)(\nu\ell b) + V \) | 43 | 12900 | 42 | 2 | 0 |
| \( tth \rightarrow (\nu\ell b)(\nu\ell b)(bb) \) | 15 | 4500 | 122 | 3 | 0 |
| SM Backgrounds Total | - | - | - | - | 4.5 |

Table 1: Signal cross section is given by \( \sigma(pp \rightarrow AA \rightarrow (Zh)(Zh)) \times 2\text{Br}(Z \rightarrow \ell\ell)\text{Br}(Z \rightarrow \nu\bar{\nu}) \times \text{Br}(h \rightarrow b\bar{b})^2 \), where \( g_{hAA} \) is taken to be \( 4m_A^2/v \) in the signal benchmarks. \( j = u, c, d, s, g, b \).

Table. 1 shows the cutflow for two signal benchmarks and the main SM back-
grounds. For the pre-selection, we require 2 leptons and 4 $b$-tagged jets with standard $p_T$ and $\eta$ cuts. To suppress $Z$+jets, we require $E_T > 120$ GeV. To suppress leptonically-decayed top backgrounds, we further require that the lepton pair has an invariant mass within the 15 GeV window around the $Z$ boson mass (on-Z). Furthermore, we investigate MT2 type of kinematic variables $^{[39,40]}$. In particular, MT2$_W$, formed by two leptons plus $E_T$, and MT2$_t$, formed by two leptons and two leading $b$-tagged jets plus $E_T$, can be used to discriminate the di-leptonic $t\bar{t}$. Even though the efficiency for pre-selection is $\lesssim 5\%$ for signals, the cut efficiency $\gtrsim 50\%$ in spite of different mass values.

**Fig. 6:** Left panel shows the upper limits on $g_{hAA}/v$ from $pp \rightarrow h \rightarrow AA \rightarrow (Zh)(Zh) \rightarrow (\ell\ell bb)(\nu\nu bb)$ at the LHC with 300 fb$^{-1}$ and HL-LHC (3000 fb$^{-1}$). We also include the value of $g_{hAA}/v$ in aIDM for $m_2^2 = 0$, which is given by $4m_A^2/v^2$. Right panel shows the cut efficiency for the signals.

From Fig. 6 we see that the HL-LHC will be sensitive in the region around 215 – 230 GeV. As well as for the other modes, this reach could be improved with a more sophisticated analysis, therefore this an interesting final state to look at.

## 5 Results and Conclusions

In this work, we assume a simple extended Higgs sector (2HDM) that exhibits alignment without decoupling, allowing the additional scalars to have large couplings to the Standard Model Higgs. We investigate the potential of the LHC for constraining the exotic triple Higgs couplings $hHH$, where $h$ represents the SM Higgs and $H$ the lightest exotic neutral Higgs. In particular, the exotic Higgs sector is assumed to be odd under an approximate $Z_2$ symmetry, such that single production of the additional
states is highly suppressed, leaving pair production as the main production mechanism. Motivated by the non-decoupling feature exhibited by this scenario, where the triple Higgs coupling $g_{hHH}$ scales roughly as the mass squared of $H$, we look at the pair production of $H$ via an off-shell SM Higgs: $pp \to h^* \to HH$. This cross section does not drop as fast as the square of the gluon parton distribution function, because of the enhancement of the coupling in the high mass regime.

We identify the main decay modes of pair produced $H$s by varying the size of the small $Z_2$-breaking parameters $\epsilon_V$ and $\epsilon_A$, that control the tree-level couplings of $H$ to SM vector bosons and fermions, respectively. Furthermore, when both $\epsilon_V$ and $\epsilon_A$ become vanishingly small, loop-induced di-photon decay becomes dominant for the lightest exotic neutral Higgs.

![Fig. 7: Combined 95% CL upper bounds on $g_{h\phi\phi}$ at 13 TeV LHC. The solid (dashed) curve represents the exclusion limit at 300 fb$^{-1}$ (3000 fb$^{-1}$). Examples of $g_{h\phi\phi}$ expected from almost inert double model are plotted on the right panel.](image)

We find that the HL-LHC can be efficient in constraining $g_{hHH}$ in a non-decoupling extended Higgs sector. In particular, the most sensitive channel is 4 $b$-jets when $\epsilon_A > \epsilon_V$, and 3$\ell$ when $\epsilon_A < \epsilon_V$, both of which are covered by existing ATLAS and CMS searches. Fig. 7 shows a summary of the regions, in the $g_{h\phi\phi}$ vs $m_\phi$ plane, excluded by the different final states we analyzed. For $b\bar{b}b\bar{b}$ and $b\bar{b}\gamma\gamma$, we extrapolated the backgrounds from the SM di-Higgs searches that are not optimized for the heavy Higgses. Consequently, the limit on the coupling obtained in these analyses is rather conservative, and has the potential to be considerably improved by a full optimized study.

In this work we have analyzed the scenario where the exotic Higgses decay promptly
inside the detector. However, this is no longer valid if all the $\mathbb{Z}_2$-breaking parameters become $\lesssim \mathcal{O}(10^{-4})$, where searches based on displaced jets, leptons or delayed photons become important. We leave this non-prompt scenario for future work, since it features a very distinctive phenomenology and requires different search strategies.
Appendix A: Details of Benchmark Studies

A.1 4b channel

We require four $b$-tagged jets with $p_T > 40$ GeV. The pairings of $b$-jets with Higgs boson candidates are accepted only if they satisfy

\[
\begin{align*}
\frac{360 \text{ GeV}}{m_{4j}} - 0.5 &< \Delta R_{jj,\text{lead}} < \frac{653 \text{ GeV}}{m_{4j}} + 0.475, \\
\frac{235 \text{ GeV}}{m_{4j}} &< \Delta R_{jj,\text{subl}} < \frac{875 \text{ GeV}}{m_{4j}} + 0.35 \quad \text{if} \quad m_{4j} < 1250 \text{ GeV}, \\
0 &< \Delta R_{jj,\text{lead}} < 1, \\
0 &< \Delta R_{jj,\text{subl}} < 1 \quad \text{if} \quad m_{4j} > 1250 \text{ GeV}. 
\end{align*}
\]

If more than one $b$-jet pairing satisfy these conditions, the one with the smallest value $D_{HH}$ value is chosen, where

\[
D_{HH} = \frac{|m_{2j}^{\text{lead}} - 120 m_{2j}^{\text{subl}}|}{\sqrt{1 + \left(\frac{120}{110}\right)^2}}. \quad (A.3)
\]

Additionally, the $p_T$ of the leading and subleading Higgs candidate is required to have

\[
\begin{align*}
p_T^{\text{lead}} &> 0.5 m_{4j} - 103 \text{ GeV}, \\
p_T^{\text{subl}} &> 0.33 m_{4j} - 73 \text{ GeV}. 
\end{align*}
\]

The pseudorapidity difference between the Higgs boson candidates is required to satisfy $|\Delta \eta_{HH}| < 1.5$. Finally, the Higgs boson candidates are required to satisfy

\[
X_{HH} = \sqrt{\left(\frac{m_{2j}^{\text{lead}} - 120 \text{ GeV}}{0.1 m_{2j}^{\text{lead}}}\right)^2 + \left(\frac{m_{2j}^{\text{subl}} - 110 \text{ GeV}}{0.1 m_{2j}^{\text{subl}}}\right)^2} < 1.6. \quad (A.6)
\]

A.2 $bb\gamma\gamma$ channel

Following the analysis in Ref. [35], we require 2 photons and between 2 and 5 jets with $|\eta| < 2.5$ and $p_T > 30$ GeV. The leading $b$-jet satisfies $p_T > 40$ GeV, while for photons $|\eta| < 2.37$ and $p_T > 30$ GeV. Events containing isolated leptons with $p_T > 25$ GeV are vetoed. The angular separation $\Delta R = \sqrt{\Delta \eta + \Delta \phi}$ satisfies

\[
\begin{align*}
0.4 &< \Delta R_{bb} < 2.0, \\
0.4 &< \Delta R_{\gamma\gamma} < 2.0.
\end{align*}
\]

(A.7)
A.3 multi-ℓ channel

For triggering, we require at least 2 jets with transverse momentum ($p_T$) larger than 30 GeV falling within $|\eta| < 2.4$ and missing transverse energy (MET) greater than 50 GeV. We veto events with $b$-tagged jets.

At least 3 isolated electrons or muons are required to have

$$\begin{aligned}
\text{leading } e (\mu) : p_T &> 25 (20) \text{ GeV} \\
\text{subleading } e (\mu) : p_T &> 15 (10) \text{ GeV}
\end{aligned} \quad \text{if HT < 400 GeV,} \quad (A.8)$$

$$\begin{aligned}
\text{leading } e (\mu) : p_T &> 15 (10) \text{ GeV} \\
\text{subleading } e (\mu) : p_T &> 15 (10) \text{ GeV}
\end{aligned} \quad \text{if HT > 400 GeV.} \quad (A.9)$$

The lowest-$p_T$ lepton must have $p_T > 10$ GeV in all cases. We further veto events containing an opposite-charge, same-flavor lepton pair with an invariant mass within the 15 GeV window around the Z boson mass.

The search is then categorized based on different $E_T$, HT and $M_T$ bins following [34], where $M_T$ is defined to be the minimum of

$$M_T = \sqrt{2p_T^{\ell} p_T^{\text{miss}}[1 - \cos(\phi_\ell - \phi_{\text{miss}})]} \quad (A.10)$$

out of all combinations. The bin that is most sensitive to our signal is that $300 \text{ GeV} > E_T > 150 \text{ GeV}$, HT $> 400 \text{ GeV}$, $M_T < 120 \text{ GeV}$.

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