Extended scalar sectors, effective operators and observed data

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Parametrization of new physics to observe $h \rightarrow \gamma\gamma$ and $h \rightarrow Z\gamma$ decays via new physics effects

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\[
\begin{align*}
(\tilde{g}_{hVV})_{new} &= \kappa_v \times (g_{hVV})_{SM} \\
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  $$(\tilde{g}_{hVV})_{new} = \kappa_v \times (g_{hVV})_{SM} \quad (1)$$
  $$(\tilde{g}_{ht\bar{t}})_{new} = \kappa_t \times (g_{ht\bar{t}})_{SM} \quad (2)$$
  $$(\tilde{g}_{hbb})_{new} = \kappa_b \times (g_{hbb})_{SM} \quad (3)$$
  $$(\tilde{g}_{h\tau\bar{\tau}})_{new} = \kappa_{\tau} \times (g_{h\tau\bar{\tau}})_{SM} \quad (4)$$

- There can be heavy states running in the loop modifying Higgs couplings, can be parametrized to express it in terms of SM gauge-invariant higher dimensional effective operators:
\[ O_{BB} = \frac{f_{BB}}{\Lambda^2} \Phi \dagger \hat{B}_{\mu \nu} \hat{B}^{\mu \nu} \Phi \]

\[ O_{WW} = \frac{f_{WW}}{\Lambda^2} \Phi \dagger \hat{W}_{\mu \nu} \hat{W}^{\mu \nu} \Phi \]

\[ O_{B} = \frac{f_{B}}{\Lambda^2} D_\mu \Phi \dagger \hat{B}^{\mu \nu} D_\nu \Phi \]

\[ O_{W} = \frac{f_{W}}{\Lambda^2} D_\mu \Phi \dagger \hat{W}^{\mu \nu} D_\nu \Phi \]

(a)
These constitute the most general set of dimension-6 effective operators which give rise to the $h\gamma\gamma$ and $hZ\gamma$ vertices and as there are no tree level diagrams exist in this two cases, the effective vertices can make some useful contributions, structured as,

$$-i\mathcal{M}_{\gamma\gamma/Z\gamma}^{(BB/WW/B/W)} = C_{\gamma\gamma/Z\gamma}^{(BB/WW/B/W)} \frac{f_{(BB/WW/B/W)}}{\Lambda^2} \times (k_1.k_2g_{\mu\nu} - k_{1\mu}k_{2\nu})\epsilon^{*\mu}(k_2)\epsilon^{*\nu}(k_1)$$ (5)

Figure 1: Feynman diagrams for $h \rightarrow \gamma\gamma$ and $h \rightarrow Z\gamma$ in the most general situation. Contributions mediated by fields other than those in SM are lumped in the blob.
The global fit

After parametrizing new physics effects, we investigate the region of parameter space favored by the 8 and 13 TeV results at the LHC. Our eight-dimensional parameter space, spanned by the four scale factors $\kappa_V, \kappa_t, \kappa_b, \kappa_\tau$ and $f_{BB}, f_{WW}, f_B, f_W$, the Wilson coefficients in the dimension-6 $hVV$ operators.
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- In order to constrain new physics from more than one independent measurements,

$$L(\mu) = \prod_{i=1}^{n} L_i(\mu) \Rightarrow \chi^2(\mu) = \sum_{i=1}^{n} \chi^2_i(\mu) = \sum_{i=1}^{n} \left( \frac{\mu - \hat{\mu}_i}{\Delta \mu_i} \right)^2$$

For correlated experimental searches, these correlations affect the Log-likelihood function as,

$$-2 \log L(\mu) = \chi^2(\mu) = (\mu - \hat{\mu}_i)^T C_{ij}^{-1} (\mu - \hat{\mu}_j)$$

Where $C^{-1}$ is the inverse of the covariance matrix $C_{ij} = \text{cov}(\hat{\mu}_i, \hat{\mu}_j)$.
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- We consider all the correlations between gluon fusion and vector-boson fusion production for each of the major Higgs decay channels.
This $\chi^2$ is then minimized to get the region allowed by the experimental data at the 1- and 2$\sigma$ levels in each two-parameter subspaces, where all remaining parameters have been marginalized.

**Figure 2:** Allowed regions at 1$\sigma$ (red) and 2$\sigma$ (blue) levels in the some of the parameter space of scale factors and dimension-6 couplings.
Extended Higgs models and dimension-6 operators
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• Various new physics models predict extended electroweak symmetry breaking sectors. It is naturally of interest to link the model-independent analysis presented above to specific theoretical scenarios.

• With this in view, we now translate the results of the global fit to those pertaining to extended Higgs models which give some contribution to new physics (like 2HDM, Higgs triplet models), taking into account the additional constraints that connect model parameters in each case.
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- With this in view, we now translate the results of the global fit to those pertaining to extended Higgs models which give some contribution to new physics (like 2HDM, Higgs triplet models), taking into account the additional constraints that connect model parameters in each case.

- To make some bridge between model dependent and model independent approach to study new BSM physics we need some consistency between this two approaches.
Validation to make connection

The various models of our interest take part to \( h \to \gamma\gamma \), \( h \to Z\gamma \) processes via loop induced processes.

The matrix element for the processes \( h \to \gamma\gamma \) coming from extra charged scalar contributions is always of the form

\[
-\frac{\mathcal{M}}{\sqrt{2}} = C_{\text{vertex}} \times \left( p \cdot q_{\gamma} \mu \nu - p \mu q \nu \right) \epsilon^* \mu (q) \epsilon^* \nu (p) \times F(\gamma\gamma) s(m_{H^{\pm}})
\]

Where \( F(\gamma\gamma) \) is the integral factor, which should in general depend on \( m_{H^{\pm}} \), the gauge invariant mass and not coming from EWSB.
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![Diagram showing process](image-url)
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$$-i\mathcal{M} = C_{\text{vertex}} \times (p \cdot q g_{\mu \nu} - p_\mu q_\nu) \epsilon^{* \mu}(q) \epsilon^{* \nu}(p) \times F_s(\gamma \gamma)(m_{H^\pm})$$

Where $F_s(\gamma \gamma)$ is the integral factor which should in general depend on $m_{H^\pm}$, which is gauge invariant mass and not coming from EWSB.
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\]

Where \( F_s(Z\gamma) \) is the integral factor which should in general depend on \( m_Z \) also.
• But in the case of $h \to Z\gamma$ the situation is quite different.

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Where $F_s(Z\gamma)$ is the integral factor which should in general depend on $m_Z$ also.

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• The matrix element for the processes $h \rightarrow Z\gamma$ is always of the form

$$-i\mathcal{M} = C_{\text{vertex}} \times (p.qg_{\mu\nu} - p_\mu q_\nu)\epsilon^*\mu(q)\epsilon^*\nu(p)\times F_s^{(Z\gamma)}(m_{H^\pm}, m_Z)$$

Where $F_s^{(Z\gamma)}$ is the integral factor which should in general depend on $m_Z$ also.
• But decay amplitude coming from effective theory has no dependence of EWSB. As a consequence of this Eqn-(5) is independent of $m_{EWSB} \approx m_Z$.
• A natural way of establishing consistency between the two amplitudes, therefore, is to have no $m_Z$-dependence in the loop amplitude for $h \rightarrow Z\gamma$ in Eqn-(10) as well.
Figure 4: Dependence of the additional scalar loop integral on the $Z$ boson mass in $h \to Z\gamma$. 
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- It is, therefore, legitimate to encapsulate the contributions to these loop amplitudes as Wilson coefficients of dimension-6 gauge-invariant operators, so long as the charged scalars responsible for these amplitude are at least a factor of two heavier than the $Z$. 
Allowed parameter spaces in new physics in different model via $h \rightarrow \gamma\gamma$ and $h \rightarrow Z\gamma$ channel, constrained by model independent approach

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- For different kinds of CP-conserving 2HDMs (Type-I, Type-II, Lepton-Specific, Flipped) scalar potential is similar with different Yukawa couplings for different cases.
Allowed parameter spaces in new physics in different model via $h \to \gamma \gamma$ and $h \to Z \gamma$ channel, constrained by model independent approach

1) 2HDMs

- For different kinds of CP-conserving 2HDMs (Type-I, Type-II, Lepton-Specific, Flipped) scalar potential is similar with different Yukawa couplings for different cases.
- They contribute to the new physics effects to $h \to VV$ process via $H^\pm$ with

$$\Gamma(h \to VV) = \text{Constant} \times \sum_i |N_{ci} R_i h^2 F_i|^2$$

where $i$ denoting by $H^\pm$, top, bottom and $W$ – boson.
We study allowed model independent parameter spaces and reflect them to the allowed model dependent parameter spaces of different 2HDMs as $R_i^h \equiv \tilde{g}_{hH^+H^-}$, $\sin(\beta - \alpha)$ and $f(m_{H^\pm})$ related to $F_i$. 

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Where the left side of yellow line are restricted from experimental bound on $m_{H^\pm}$. 

**Type-I**

**Type-II**

**Lepton-Specific**

**Flipped**
2) Higgs Triplet-models

Similar procedure are maintain here with $m_{H_i}^\pm$ and $m_{H_i}^{\pm\pm}$.

**Single Triplet-model**

\[ g_{hH^+H^-}^{\text{eff}} = \tilde{g}_{hH^+H^-} + 4\tilde{g}_{H^++H^-} \]
two-Triplet-model

\[ g_{\text{eff}}^{hH_1^+H_1^-} = \tilde{g}_1^{hH_1^+H_1^-} + 4\tilde{g}_1^{hH_1^{++}H_1^{--}} \]  
(10)

\[ g_{\text{eff}}^{hH_2^+H_2^-} = \tilde{g}_2^{hH_2^+H_2^-} + 4\tilde{g}_2^{hH_2^{++}H_2^{--}} \]  
(11)

with
With this in view we consider the ratio \( r = \frac{\mu_{\gamma\gamma}}{\mu_{Z\gamma}} = \frac{\Gamma(h \rightarrow \gamma\gamma)}{\Gamma(h \rightarrow Z\gamma)} / \frac{\Gamma(h \rightarrow \gamma\gamma)_{SM}}{\Gamma(h \rightarrow Z\gamma)_{SM}} \) and show color-coded regions in the \( f_{BB}-f_{WW} \) plane.

- **As an example we take Type-I 2HDM**
Summary

• We parametrize new physics beyond SM in terms of higher dimensional effective operators and also tree-level SM-coupling modifiers.
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Summary

- We parametrize new physics beyond SM in terms of higher dimensional effective operators and also tree-level SM-coupling modifiers.
- We impose all the available experimental results to constrain these parameters simultaneously and find the favored region.
- Then we make an observation that even in various specific models, the extra loop integrals can be compared with the tree-level vertex factors arising out of higher dimensional effective operators.
- Hence we tried to make a connection between these two approaches and then establish a connection between allowed parameter regions in the model-independent approach with those of specific new physics models.
Thank You
Covariance matrix

\[
\chi^2 = \sum_{i=1}^{N} \frac{(d_i - t_i - \sum_{j=1}^{K} \alpha_j s_{ji})^2}{\sigma_i^2} + \sum_{j=1}^{K} \alpha_j^2
\]

- \(d_i\) = content of data bin \(i\);
- \(t_i\) = model prediction for bin \(i\);
- \(s_{ij}\) = systematic uncertainty from source \(j\) on the contents of bin \(i\);
- \(\sigma_i\) = statistical uncertainty on the contents of bin \(i\);
- \(\alpha_j\) = fit parameters.

When this \(\chi^2\) is minimised w.r.t \(\alpha_j\)'s then

\[
\chi^2_{min} = \Delta^T C^{-1} \Delta
\]

- \(\Delta\) is a column matrix with elements \((d_i - t_i)\) and \(C\) is covariance metrix \(C_{pl} = \text{cov}(d_p, d_l)\) of the measurements \(d_i\).

\[
C_{ij} = \sigma_i^2 \delta_{ij} + \sum_{m=1}^{K} s_{mi} s_{mj}
\]
Different 2HDMs

- **Type-I 2HDM:**
  Here all fermions are assumed to couple to the same doublet, so that

  \[ \mathcal{L}_{\text{Yukawa}} = y_{ij}^1 \bar{Q}_i L \Phi_2 d_j R + y_{ij}^2 \bar{Q}_i L \bar{\Phi}_2 u_j R + y_{ij}^5 \bar{L}_i L \Phi_2 e_j R \]  

  (12)

  This can be achieved by imposing the discrete symmetry on the \( \mathcal{L}_{\text{Yukawa}} \), \( \Phi_1 \rightarrow -\Phi_1 \).

- **Type-II 2HDM:**
  Here up-type quarks to couple to one doublet, and down-type quarks and leptons to couple to another. Under this assumption

  \[ \mathcal{L}_{\text{Yukawa}} = y_{ij}^1 \bar{Q}_i L \Phi_1 d_j R + y_{ij}^2 \bar{Q}_i L \bar{\Phi}_2 u_j R + y_{ij}^5 \bar{L}_i L \Phi_1 e_j R \]  

  (13)

  This can be enforced by demanding that the \( \mathcal{L}_{\text{Yukawa}} \) remains invariant under \( \Phi_1 \rightarrow -\Phi_1 \) and \( d_R \rightarrow -d_R \) and \( e_R \rightarrow -e_R \).
Different 2HDMs

- **Lepton-Specific 2HDM:**
  Here the up-and down-type quarks to couple to one doublet and leptons to the other doublet.

\[
\mathcal{L}_{\text{Yukawa}} = y_{ij}^1 \bar{Q}_i L \Phi_2 d_R + y_{ij}^2 \bar{Q}_i L \Phi_2 u_R + y_{ij}^5 \bar{L}_i L \Phi_1 e_R
\]  

(14)

This can be achieved by imposing the symmetry of the Lagrangian under the discrete symmetry $\Phi_1 \rightarrow -\Phi_1$ and $e_R \rightarrow -e_R$. $\mathcal{L}_{\text{Yukawa}}$ in this case becomes

- **Flipped 2HDM:**
  Here the up-type quarks and leptons couple to the same doublet while the other doublet couples to down-type quarks. The Yukawa Lagrangian becomes

\[
\mathcal{L}_{\text{Yukawa}} = y_{ij}^1 \bar{Q}_i L \Phi_1 d_R + y_{ij}^2 \bar{Q}_i L \Phi_2 u_R + y_{ij}^5 \bar{L}_i L \Phi_2 e_R
\]  

(15)

which is achieved by imposing $\Phi_1 \rightarrow -\Phi_1$ and $d_R \rightarrow -d_R$. 
Back up slides 2
Figure: 1(red) and 2σ(blue) allowed regions.