Implications of LHC data on 125GeV Higgs-like boson for the Standard Model and its various extensions

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Abstract: Recent data on 125 GeV Higgs-like boson at the LHC starts to constrain the electroweak symmetry breaking (EWSB) sector of the SM and its various extensions. If one imposes the local gauge symmetry of the Standard Model (SM) \( SU(3)_C \times SU(2)_L \times U(1)_Y \) to the SM and any possible new physics scenarios, the SM Higgs properties will be modified by intrinsically two different ways: by new physics either coupling directly to the SM Higgs boson \( h \), or affecting indirectly the SM Higgs properties through the mixing of \( h \) with a SM singlet scalar \( s \). (Here \( s \) is a singlet under the SM gauge group, but may be charged under a new gauge charge and can have nonrenormalizable couplings to non-SM particles.) The models of two Higgs doublet, extra sequential and mirror fermions belong to the first category, whereas the models with a hidden sector dark matter, extra vector-like fermions and new charged vector bosons, which can enhance the diphoton rate of the SM Higgs-like resonance, belong to the second category. We perform a global fit to data in terms of the effective Lagrangian description of two interaction eigenstates of scalar bosons, a SM Higgs and a singlet scalar, and their mixing. This framework is more suitable to study singlet-extended scenarios discussed above compared to other approaches based on the Lagrangian of mass eigenstates. With fairly model-independent assumptions, the effective Lagrangian contains at most four free parameters still encompassing the majority of models in the literature. Interestingly, the SM gives the best fit if all data from ATLAS and CMS are used, whereas various singlet extensions can fit better to individual ATLAS or CMS data. Without further assumptions, an upper bound on the total width (or, non-standard branching ratio) is generically obtained. Furthermore, global fit based on our parameterization can be used to probe interactions of the singlet scalar if the singlet resides below \( 2m_W \).

Keywords: Higgs boson, singlet scalar, hidden sector, dark matter, invisible Higgs decay, global fit, diphoton
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1 Introduction

After the discovery of a new boson of mass around 125 GeV at the LHC \cite{1, 2}, the most prompt question regarding this new particle is to identify its nature in particle physics context, namely to verify if it is the SM Higgs boson that have been long sought for, or something else. For this purpose, its spin and parity \cite{3–12} and its couplings to the SM particles should be determined as accurately as possible (see Ref. \cite{13} for the review on the SM Higgs boson and references therein).

The recent results from ATLAS and CMS already tell us that the $J^P$ quantum number of 125 GeV boson is consistent with the SM prediction, namely $J^P = 0^+$ \cite{14, 15}. Other assignments of $J^P$ quantum number such as $J^P = 0^-$ or $2^+, 2^-$ yield worse $\chi^2$ fits to the data compared to the $J^P = 0^+$ assignment, and thus disfavored.

Determining the couplings of the 125 GeV resonance to the SM fermions and weak gauge bosons ($W^\pm, Z^0$) as well as gluons and photons has been studied by many groups \cite{16–50} One usually considers a general case where the Higgs properties may be modified by new particles which can be described by dim-5 and dim-6 operators in the effective Lagrangian approaches \cite{51}. In most of these approaches, it is assumed that all the new particles are heavy enough and can be integrated out, resulting in the higher dimensional operators and/or radiative corrections to renormalizable SM interactions. This assumption works for the inert scalar doublet model, the fourth generation sequential or mirror fermions, or new colored and/or charged scalar fields.

However, in the presence of EW-scale singlet scalar boson that can mix with the SM Higgs boson, it may not be a good approximation to integrate out the singlet scalar bosons. For example, in Refs. \cite{52–54}, the authors demonstrated that the effective Lagrangian approaches based on the unbroken subgroup $(SU(3)_C \times U(1)_Y)$ of the SM gauge group produce erroneous results, especially on the direct detection cross section through the Higgs couplings to the dark matter particles. In the renormalizable versions of hidden sector (or Higgs-portal) DM models with the full SM gauge group $(SU(3)_C \times SU(2)_L \times U(1)_Y)$, one always have an extra singlet scalar $s$ which couples to the hidden sector DM particles at renormalizable level. And the mixing of this singlet scalar $s$ with the SM Higgs $h$ thermalize the hidden sector DMs. Due to the mixing between the $h$ and $s$, there are always at least two neutral Higgslike scalar bosons, which are mixtures of the SM Higgs $h$ (remnant of the $SU(2)_L$ doublet Higgs field) and a singlet scalar $s$. The singlet scalar $s$ not only thermalizes the hidden sector DM efficiently, but also improves the stability of the EW vacuum up to Planck scale \cite{53}, unlike the SM case. It is important to realize that one can not integrate out the singlet scalar $s$ simply assuming that $s$ is very heavy. The mixing should be taken into account properly, the discussion of which can be found in Sec. 3.3 of Ref. \cite{54}. Furthermore, the usual effective Lagrangian based on $SU(3)_C \times U(1)_Y$ is not proper approach.
when one attempts to deduce information on new light particles close to the observed 125 GeV resonance and related underlying physics.

In fact, an extra neutral scalar boson does appear in many interesting extensions of the SM, and it generally mixes with the SM Higgs boson unless the mixing is forbidden by some exact symmetry. Introducing an extra singlet scalar boson is not only the simplest extension of the SM in terms of the number of new degrees of freedom, but also has various virtues of leaving the $\delta\rho$ parameter intact at the tree level, and improving the vacuum stability [55] as well as helping achieve the stronger first order electroweak phase transition [56]. Also, SM singlet scalars can be charged under a new gauge symmetry spontaneously breaking the symmetry or playing an role of cold dark matter (CDM). Furthermore, during the last few years or so, there have been a lot of works which considered singlet scalars coupling to extra vectorlike fermions or new charged vector bosons in order to enhance the diphoton rate of 125GeV resonance. Most models involve the extra singlet scalar that mixes with the SM Higgs boson as described in Sec. 3.

Let us list some example BSM’s (with a few references) where at least one neutral scalar boson appears and mixes with the SM Higgs boson $h$:

- Pure singlet extension of the SM [55–59]
- Hidden (dark) sector DM with hidden (dark) gauge symmetry [52–54, 60–67]
- Dilaton or radion [68–80]
- Enhancing diphoton rate from vector-like (VL) fermions or new charged vector bosons [81–89]

Determination of the Higgs couplings in the presence of a (light) singlet scalar boson has not been discussed properly in the recent literature (but, for some earlier attempts, see Refs.[16, 27, 40]). In fact, in most phenomenological analyses of 125 GeV boson, the usual approach actually does not distinguish the 125 GeV boson discovered at the LHC being the pure SM Higgs boson or a mixture with a singlet scalar boson for the reasons described in Sec. 2.

It is the purpose of this paper to introduce a proper methodology for analyzing the LHC Higgs data in the presence of new physics that affects the SM Higgs couplings either directly or indirectly through the mixing of the SM Higgs boson with a singlet scalar that can couple to new physics.

To this end, instead of parameterizing BSM effects on the mass-eigenstate of 125GeV resonance directly, we introduce separate parameterizations for each interaction eigenstates of SM Higgs ($h$) and a singlet scalar boson ($s$), and their mixing through the mixing angle $\alpha$. In other words, we introduce

$$\alpha, \quad b_i, \quad c_i$$

for the mixing angle, corrections to SM Higgs couplings and the singlet scalar couplings to the SM fields $i = WW, ZZ, gg, \gamma\gamma, Z\gamma, f\bar{f}$, respectively (see Sec. 2.1 and 2.2 for detail). With these parameterizations, physical interpretation of data and global fits will be different from
usual approaches where only final modified Higgs couplings can be measured. Usefulness of our approach is clear; one can directly extract information of Lagrangian parameters and thus possibly underlying physics responsible for them. This parameterization also makes it clear that signal strengths can be modified from several independent sources simultaneously so that it is difficult to understand the origin of modification without considering a proper Lagrangian; see Fig. 1 for illustration of this important statement. By the same reason, data may be parameterized by sometimes redundant parameters implying not only that no unique solution of global fit may be found, but also that some parameters maybe hard to be constrained. We discuss how one can handle this difficulty, and how one can still deduce upper bound on non-standard branching ratio, which often causes this difficulty, in a generic way.

This framework is versatile enough to capture various interesting models of singlet scalar bosons ubiquitous in BSMs (as well as the new physics scenarios without a singlet scalar boson). Each case we consider in this paper has corresponding models and motivations. We also discuss how the fit results on these parameters can further be used to constrain new particles that may exist nearby electroweak scale. This will prove the usefulness and the prospects of our approach.

This paper is organized as follows. In Sec. 2, we introduce effective Lagrangians useful to analyze the Higgs properties in the presence of a possible mixing between the SM Higgs boson and a singlet scalar boson nearby, including the cases where the singlet scalar boson couples to new particles such as vectorlike fermions or new charged vector bosons or scalar bosons. In Sec. 3, we list a number of BSM’s which have extra SM singlet scalar fields. Then in Sec. 4, we explain our fit procedure and the LHC data we use. In Sec. 5, we perform global fits to the LHC Higgs data under various assumptions, encompassing all the models described in Sec. 3. In Sec. 6, the fit results are combined with other collider
searches of Higgs-like resonances to constrain another scalar boson that may be present nearby electroweak scale. Finally the results are summarized in Sec. 7.

2 Effective Lagrangian for the SM Higgs and a singlet scalar bosons

In order to extract the Higgs couplings to the SM fermions and gauge bosons in a model independent fashion, it is useful to construct the most general effective Lagrangian describing the interactions between the SM Higgs boson $h$ or a singlet scalar boson $s$ with the SM fields such as $b, t, W^\pm, Z^0, \gamma$ and $g$. In this section, we show the effective Lagrangian for this purpose up to dim-6 operators. We impose the full SM local gauge symmetry $SU(3)_c \times SU(2)_L \times U(1)_Y$ to the effective Lagrangian, not its unbroken subgroup $SU(3)_c \times U(1)_{em}$. By doing so, we can separate two different sources of the modified Higgs properties: one from direct couplings of new particles to the SM Higgs boson ($b_i \neq 1$ in Fig. 1), and the other from the mixing with a singlet scalar boson ($\alpha \neq 0$ in Fig. 1). There could be new particles that have gauge invariant renormalizable couplings to the singlet scalar $s$ ($c_i \neq 0$ in Fig. 1), but not to the SM Higgs boson $h$. Therefore studying the Higgs properties in the 3-dimensional space (ignoring the dimensionality associated with the index $i$) as depicted in Fig. 1 can be justified, and its importance could be appreciated.

2.1 Effective Lagrangian for the SM Higgs boson $h$

Let us assume that the SM Higgs boson couplings are modified due to some new physics effects even without the mixing with a singlet scalar $s(x)$. This could happen if there are additional sequential or mirror fermions (chiral), or extra inert scalar doublet, for example. Integrating out the new heavy particles, one can construct the effective Lagrangian up to dim-5 and dim-6 operators, all of which have been identified by Buchmüller and Wyler sometime ago [90]. We do not reproduce all the operators involving the Higgs fields, but list only some of them just for illustration:

$$H^\dagger H \sigma^\mu \sigma^\nu G^\mu \nu, \quad (H^\dagger D^\mu H)(D^{\mu} H^\dagger), \quad H^\dagger H \overline{Q_{3L}} t \tilde{H} t_R,$$

relegating the complete list to the original paper [90].

Expanding the Higgs field in the effective Lagrangian constructed by Buchmüller and Wyler around the EW vacuum with

$$H(x) = \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix},$$
we obtain the following effective operators of interaction eigenstate $h(x)$ field up to dim-6:

$$-\mathcal{L}_{h,\text{int}} = \sum_f m_f \left\{ b_f \frac{h}{v} + \frac{1}{2} b'_f \left( \frac{h}{v} \right)^2 \right\} \bar{f} f$$

$$- \left\{ 2b_W \frac{h}{v} + b'_{\gamma} \left( \frac{h}{v} \right)^2 \right\} m_W^2 W_\mu W_{\mu} - \left\{ b_Z \frac{h}{v} + \frac{1}{2} b'_Z \left( \frac{h}{v} \right)^2 \right\} m_Z^2 Z_\mu Z^\mu$$

$$+ \frac{\alpha}{8\pi} r_{\text{sm}} \left\{ b_\gamma \frac{h}{v} + \frac{1}{2} b'_\gamma \frac{h}{v} \right\} F_{\mu\nu} F^{\mu\nu} + \frac{\alpha_s}{16\pi} r_{\text{sm}} \left\{ b_t \frac{h}{v} + \frac{1}{2} b'_t \frac{h}{v} \right\} G^a_{\mu\nu} G^{a\mu\nu}$$

$$+ \frac{\alpha_2}{\pi} \left\{ 2b_d W h \frac{h}{v} + b_d W \left( \frac{h}{v} \right)^2 \right\} W_{\mu\nu} W^{\mu\nu} + \frac{\alpha_2}{\pi} \left\{ 2b_d Z h \frac{h}{v} + b_d Z \left( \frac{h}{v} \right)^2 \right\} Z_{\mu\nu} Z^{\mu\nu}$$

$$+ \frac{\alpha_2}{\pi} \left\{ 2b_d W \frac{h}{v} + b_d W \left( \frac{h}{v} \right)^2 \right\} W_{\mu\nu} W^{\mu\nu} + \frac{\alpha_2}{\pi} \left\{ 2b_d Z \frac{h}{v} + b_d Z \left( \frac{h}{v} \right)^2 \right\} Z_{\mu\nu} Z^{\mu\nu}$$

$$+ \frac{\alpha}{\pi} \left\{ 2b_d Z h \frac{h}{v} + b_d Z \left( \frac{h}{v} \right)^2 \right\} F_{\mu\nu} Z^{\mu\nu}$$

(2.1)

where $f$ in the first term of the Lagrangian denotes the SM fermions. The Higgs field $h(x)$ is defined after the EWSB: $H(x) = v + h(x)$, and before any possible mixing with a singlet scalar $s$ which will be introduced shortly.

Most of dim-6 operators lead to the definite relation, $b_i = b'_i$, since they involve $H^1 H$ which yields $(v + h)^2$. But this is not the case for $b_f$ and $b'_f$. For example, the following operators ($q_L \equiv (t_L, b_L)$), which are invariant under the full SM gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$,

$$\bar{q}_L D_\mu b R D^\mu H, \quad \bar{q}_L t_R D^\mu H,$$

contribute to the $b_f \sim m_f^2/\Lambda^2$, but not to $b'_f$. Thus the relation $b_f = b'_f$ is no longer true for the Higgs couplings to the SM chiral fermions.

Modification to the SM Higgs Lagrangian is parameterized by multiplicative constants $b_i$ and $b'_i$, and the SM is recovered when all $b_i = b'_i = 1$. We are interested in

$$b_f, b_W, b_Z, b_\gamma, b_g$$

(2.2)

among coefficients $b_i$ because these are most constrained by the current LHC data.

Loop-induced couplings of the SM higgs to photons and gluons involve loop functions $r_{\text{sm}}$ defined in the SM as

$$r_{\text{sm}}^\gamma = A_1(\tau_W) + N_c Q_t^2 A_{1/2}(\tau_t)$$

(2.3)

$$r_{\text{sm}}^g = A_{1/2}(\tau_t)$$

(2.4)

where we follow definitions of $A_1$ and $A_{1/2}$ as in Ref.[13], and $\tau_t = m_t^2/4m_W^2$ and $v = 246$ GeV. Loop effects of new physics is conveniently incorporated as additive shifts $\Delta b_\gamma$ and $\Delta b_g$ defined as

$$b_g = b_0 C_t + \Delta b_g$$

(2.5)
\[ b_\gamma = b_W B_W + b_b B_t + \Delta b_\gamma \]  
(2.6)

Note that \( b_{g,\gamma} \) (and their \( \Delta b \)) are normalized to the corresponding SM couplings. \( b_t \) and \( b_W \) parts describe effects from modification of top and \( W \) boson couplings to Higgs that are involved in loop diagrams. The relative loop-functions of \( W \) boson and top quark for \( m_h = 125 \) GeV are given by

\[ C_t = 1 \]  
(2.7)

\[ B_W = \frac{A_1(\tau_W)}{A_1(\tau_W) + N_c Q_t^2 A_{1/2}(\tau_t)} \simeq 1.283, \]  
(2.8)

\[ B_t = \frac{N_c Q_t^2 A_{1/2}(\tau_t)}{A_1(\tau_W) + N_c Q_t^2 A_{1/2}(\tau_t)} \simeq -0.283. \]  
(2.9)

These parameters \( b_i \)'s may have momentum(mass) dependence. We define these variables at 125GeV relevant to the global fit to 125GeV resonance data.

In most part of this paper, we work on the 125GeV resonance, thus we can conveniently assume that \( b_i \) do not have mass dependence. However, for \( b_g \) and \( b_\gamma \) which are loop-induced couplings, we will discuss mass dependence in Sec. 6 when we study constraints on other particles. We also assume that these parameters are real. This assumption would be good as long as the loop diagram does not develop unitarity phase from the case where the loop particles are on-shell. Considering various constraints on new charged or colored particles, it would be reasonable to assume that there are no new charged or colored particles with mass less than \( m_H/2 \simeq 63 \) GeV.

In the presence of new particles with nonzero EW gauge charges (e.g., another Higgs doublet as in 2 Higgs doublet model, extra sequential fermions or mirror fermions), both tree level processes \( h \to W^+W^- , Z^0Z^0 \) and the loop process \( h \to gg, \gamma\gamma \) can be modified, resulting in \( b_V \neq 1 \) and \( b_\gamma \neq 1 \) and \( b_g \neq 1 \). Except for the 2HDM case, these new physics effects will appear at one loop level, and we would expect that

\[ b_i \sim \text{“1”} + \frac{g^2 m^2}{(4\pi)^2 M^2}, \quad \text{or “1”} + \frac{g^2 m^2}{M^2} \]

where \( m \) is the external SM particle mass, \( M \) is the mass of new particles in the loop, and \( g \) is the couplings between them.

### 2.2 Effective Lagrangian for a singlet scalar boson \( s \)

As in the case of the effective Lagrangian of the SM Higgs field \( H(x) \) up to dim-6, one can construct effective Lagrangian involving a singlet \( S(x) \) and the SM fields up to dim-6, imposing the SM gauge symmetry \( SU(3)_C \times SU(2)_L \times U(1)_Y \). Note that there are only a few operators describing interactions between \( S \) and the SM Higgs boson at renormalizable level:

\[ S H^\dagger H, \quad S^2 H^\dagger H, \]

in addition to the singlet self couplings: \( S^3 \) and \( S^4 \), which lead to the modified self couplings of two Higgs-like scalar bosons \( H_1 \) and \( H_2 \) after the EWSB and the mass mixing between \( h \) and \( s \), as described in Sec. 2.3 below.
Interactions between the singlet scalar $S$ and the SM chiral fermions and the SM gauge bosons occur only at the nonrenormalizable level due to the full SM gauge symmetry\textsuperscript{1}, $SU(3)_C \times SU(2)_L \times U(1)_Y$. As an example, we list a few of them:

\begin{align*}
S \, G_{\mu \nu}^a \, G^{a \mu \nu}, \quad S^2 \, G_{\mu \nu}^a \, G^{a \mu \nu}, \quad S \, D_\mu H^1 D^\mu H, \quad S^2 \, D_\mu H^1 D^\mu H, \\
S \, \overline{Q}_{3L} H_R, \quad S^2 \, \overline{Q}_{3L} H_R,
\end{align*}

etc.. We considered the most general Lagrangian without any symmetry such as $Z_2$ symmetry under $S \to -S$ which is often invoked in the real singlet scalar DM models. It would be a separate question what kind of new underlying physics would generate such dim-5 or dim-6 operators, which we don’t address in this paper.

The singlet scalar field $S(x)$ may develop a nonzero VEV independent of the EWSB:

$$S(x) = v_S + s(x).$$

Expanding around $v_S$, we define the physical singlet scalar $s(x)$ in the interaction basis. Then, the effective Lagrangian for the singlet interaction eigenstate scalar boson $s$ could be written as

\begin{align}
-\mathcal{L}_{s,\text{int}} &= \sum_f c_f \frac{m_f}{v} s \bar{f} f - \left\{ 2c_W \frac{s}{v} + c_W' \left( \frac{s}{v} \right)^2 \right\} m_W^2 W^+_{\mu \nu} W^{-\mu \nu} - \left\{ c_Z \frac{s}{v} + \frac{1}{2} c_Z' \left( \frac{s}{v} \right)^2 \right\} m_Z^2 Z_{\mu \nu} Z^{\mu \nu} \\
&\quad + \frac{\alpha}{8 \pi} r_{\text{SM}}^g \left\{ c_g \frac{s}{v} + \frac{1}{2} c_g' \left( \frac{s}{v} \right)^2 \right\} F_{\mu \nu} F^{\mu \nu} + \frac{\alpha s}{16 \pi} r_{\text{SM}}^g \left\{ c_g \frac{s}{v} + \frac{1}{2} c_g' \left( \frac{s}{v} \right)^2 \right\} G_{\mu \nu}^{a \mu \nu} \\
&\quad + \frac{\alpha_2}{\pi} \left\{ 2 c_{dW} \frac{s}{v} + c_{dW'} \left( \frac{s}{v} \right)^2 \right\} W^+_{\mu \nu} W^{-\mu \nu} + \frac{\alpha_2}{\pi} \left\{ 2 c_{dZ} \frac{s}{v} + c_{dZ'} \left( \frac{s}{v} \right)^2 \right\} Z_{\mu \nu} \tilde{Z}^{\mu \nu} \\
&\quad + \frac{\alpha}{\pi} \left\{ 2 c_{\gamma} \frac{s}{v} + c_{\gamma'} \left( \frac{s}{v} \right)^2 \right\} F_{\mu \nu} \tilde{Z}^{\mu \nu} - \mathcal{L}_{\text{nonSM}} \tag{2.11}
\end{align}

The newly introduced couplings $c_i$’s parameterize the couplings of $s$ to the SM particles in a similar way to the SM Higgs ($h$) couplings to the SM particles. The singlet interaction eigenstate $s(x)$ is defined after the symmetry breaking due to possible nonzero VEV of a singlet scalar field $S(x)$ but before mixing with the SM Higgs field $h$. The last term $\mathcal{L}_{\text{nonSM}}$ represents possible interactions of the singlet scalar $s$ with non-SM particles such as dark matter in some dark matter models such as hidden sector dark matter models. We do not specify this Lagrangian, but we will parameterize this effect by non-standard branching ratio in later sections.

Since all the couplings $c_i$’s are from nonrenormalizable interactions between the singlet scalar $S$ and the SM fields (except for the Higgs fields), one can assume that $c_i$’s are all suppressed by heavy mass scale and/or the loop suppression factors:

$$c_i \sim \frac{g^2 m_s^2}{(4\pi)^2 M^2}, \quad \frac{g^2 m_s^2}{M^2},$$

\textsuperscript{1}As discussed in Sec. 1, the singlet scalar could have renormalizable interactions $S \bar{f} f$ with the SM fermions if we imposed only the unbroken part of the SM gauge symmetry. However this can lead to erroneous results as demonstrated in Refs. [52–54] in the context of Higgs portal DM models.
where \( m^2 \) indicates the mass scales of the external SM particles or the singlet scalar. Therefore the natural scale for the \( c_i \)'s would be parametrically suppressed relative to “1”.

Also, these parameters \( c_i \)'s may have momentum (or mass) dependence. We define these parameters at 125GeV which is relevant to global fit to 125GeV resonance data. In most part of this paper, we work on 125GeV resonance, thus we can conveniently assume that \( c_i \) do not have mass dependence. However, for \( c_g \) and \( c_\gamma \) which are loop-induced couplings, we will discuss mass dependence in Sec. 6 when we study constraints on other particles. We also assume that these parameters are real as before in case of \( b_g \) and \( b_\gamma \).

Note that \( c_f = c_V = 0 \) (with \( V = WW, ZZ \)) at renormalizable level in models with a singlet scalar boson because of the local gauge symmetry of the SM \( SU(3)_c \times SU(2)_L \times U(1)_Y \) forbids renormalizable couplings of the singlet scalar \( s \) to the SM fields except for the \( H^1H \) operator. In case of the hidden sector dark matter models (or Higgs-portal DM models), the singlet \( s \) can couple to the SM singlet particles such as hidden sector dark matters (Sec. 3.2), but not to the SM fermions or weak gauge bosons at renormalizable level. However \( c_f \) and \( c_V \) could be nonzero if \( s \) is a dilaton in spontaneously broken scale symmetric scenario or the radion in the Randall-Sundrum scenario where all the SM fields are on TeV brane (see Sec. 3.3). Unlike the couplings \( c_f \) or \( c_V \), the singlet couplings to the \( \gamma\gamma, gg, Z\gamma \) can be nonzero if there are mixings between the SM fermions and extra singlet fermions (see Sec. 3.4), or if there are an extra charged vector bosons in the theory which get massive through new Higgs fields (Sec. 3.5). In Sec. 3.6, we give an explicit example with (colored and) charged scalar fields, in which contributes \( b_g \) and \( b_\gamma \) can be modified from the SM values 1, and also \( c_g \) and \( c_\gamma \) can be generated.

### 2.3 Mixing between \( h \) and \( s \) and physical amplitudes

In general, there would be a mass mixing between \( h(x) \) and \( s(x) \) after \( H(x) \) and \( S(x) \) develop nonzero VEV’s. Relevant nonlinear interactions among them are given by

\[
-L_{\text{bilinear}} = \frac{1}{2}m^2_h h^2 + \frac{1}{2}m^2_s s^2 + m^2_{hs} hs ,
\]

\[
-L_{\text{scalar int}} = \frac{a_{3.0}}{3!} h^3 + \frac{a_{4.0}}{4!} h^4 + \frac{a_{2.1}}{2!} h^2 s + \frac{a_{3.1}}{3!} h^3 s + \frac{a_{1.2}}{2!} h s^2 + \frac{a_{2.2}}{(2!)^2} s^2 h^2 + \frac{a_{0.3}}{3!} s^3 + \frac{a_{1.3}}{3!} s^3 h .
\]  

We included the scalar self couplings for completeness, although they are not relevant to the discussions in this work. Let us parameterize the mixing effects by a mixing angle \( \alpha \) which defines the physical mass eigenstates \( H_1 \) and \( H_2 \) as

\[
H_1 = h \cos \alpha - s \sin \alpha
\]

\[
H_2 = h \sin \alpha + s \cos \alpha
\]
where we conveniently denote 125GeV resonance by $H_1$ although it can be heavier or lighter than $H_2$. Their partial widths to the SM particles $F(\neq H_{i=1,2})^2$ are written as

$$\frac{\Gamma(H_1 \rightarrow F)}{\Gamma(h \rightarrow F)_{SM}^{m_{H_1}}} = (b_F \cos \alpha - c_F \sin \alpha)^2 \quad (2.16)$$

$$\frac{\Gamma(H_2 \rightarrow F)}{\Gamma(h \rightarrow F)_{SM}^{m_{H_2}}} = (c_F \cos \alpha + b_F \sin \alpha)^2 \quad (2.17)$$

Note that we normalize the decay widths of two physical scalar boson with respect to corresponding SM width at the mass of the scalar boson. We treat $b_i$ and $c_i$ are mass-independent if they are generated at tree-level; thus, their values fitted at 125GeV are also applied to other mass region. We discuss how loop-induced couplings which are mass-dependent can be treated in Sec. 6 when we study constraints on other particles.

Another possible effect of mixing is that the heavier eigenstate can decay to the lighter one if it is kinematically allowed. We will parameterize this effect by introducing non-standard branching ratio in Sec. 4.

### 2.4 Comparison with other approaches

Before proceeding further, let us compare our approach with others. Most papers use the effective Lagrangian eq.(2.1) as the starting point, assuming that the $h$ in eq.(2.1) is the SM Higgs boson derived from the SM Higgs double $H$ and imposing the unbroken part of the SM gauge group, namely imposing only $SU(3)_C \times U(1)_{em}$. There is nothing wrong about this, since it would be the most general effective Lagrangian up to dim-6 when we impose local $SU(3)_C \times U(1)_{em}$ symmetry. However one has to be careful since the Higgs field $h$ in eq.(2.1) with local $SU(3)_C \times U(1)_{em}$ symmetry could be different from the genuine SM Higgs field, the remnant of the $SU(2)_L$ doublet scalar fields after EWSB. If we consider local $SU(3)_C \times U(1)_{em}$ symmetry, then the Higgs field in eq.(2.1) could be a mixture of the SM Higgs field and any number of electrically neutral scalar fields, some of them could be EW singlets and others could carry nontrivial EW gauge charges. Therefore there is no way one can tell whether the observed 125GeV boson is the SM Higgs boson or a mixture with a singlet scalar boson within the usual approach.

In contrast, we are proposing to separate $h$ and $s$ in the effective Lagrangian from the beginning by their EW gauge quantum numbers. Therefore one can interpret the global fit results under various assumptions on the underlying new physics models and tell which models are favored and which are not. At the moment, the data currently available is not good enough to constrain or exclude some BSM’s definitely. However in the future when more data is available with better information on the production channels, our approach would be useful for constraining various BSM’s as well as verifying the SM Higgs scenario.

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2Note that in our definition $F$ denotes the SM fields only, so that interaction eigenstate $s$ does not have couplings to $F$ except for the case $F = h$. 

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3 Models with an extra singlet scalar

In this section, we consider a number of models where extra singlet scalar bosons appear in a natural manner. The scalar $S(x)$ is assumed to be a singlet under the SM gauge group, but could be either neutral or charged under a new gauge group. The hidden sector Higgs fields \[54\] or a new singlet scalar charged under $U(1)_{H}$ in the Type-II two Higgs doublet models with gauged Higgs flavor proposed in Ref. \[91\] make good examples.

In many cases, the singlet scalar $S$ develops nonzero VEV $v_{S}$, the physical scalar $s(x)$ is defined as a fluctuation around a nonzero VEV of the original singlet field $S$:

$$S(x) = v_{S} + s(x).$$

In this case, $s$ can not have tree level couplings to the SM fermions or weak gauge bosons: $c_{V} = c_{f} = 0$, but it may have couplings to new charged particles such as vector-like(VL) fermions (either colored fermion $Q$’s or colorless leptons $L$’s) \[86, 88\] or new charged vector bosons such as $W^{\pm}$ \[81, 83\]. In either case, these new particles ($Q, L, W^{\pm}$) can modify the Higgs phenomenology only through the mass mixing between $h$ and $s$, since they couple only to the singlet $s$ but not to the SM Higgs boson $h$. They will induce the nonzero amplitudes for $s \rightarrow \gamma \gamma$ and $s \rightarrow gg$ through loop diagrams, so that $c_{\gamma} \neq 0$ for all of them and $c_{g} \neq 0$ for $Q$’s only. We assume these new charged particles are heavy enough that the 125 GeV resonance discovered at the LHC cannot decay into those particles. Then it would be possible to use the effective Lagrangian for $H(125) \rightarrow \gamma \gamma$ with the couplings $c_{\gamma}$ and $c_{g}$ being real.

3.1 The singlet extension of the SM

The simplest extension of the SM Higgs sector is to add a singlet scalar $S$ to the SM Lagrangian:

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}_{\text{new}},$$

where

$$\mathcal{L}_{\text{new}} = \frac{1}{2} (\partial_{\mu}S \partial^{\mu}S - m_{S}^{2}S^{2}) - \mu_{S}^{2}S^{2} - \frac{\lambda_{S}}{4}S^{4} - \mu_{H}S H^{\dagger}H - \frac{\lambda_{HS}}{2}S^{2}H^{\dagger}H,$$

This model has been studied in various context in \[56, 58, 59\]. After the singlet $S(x)$ develops a nonzero VEV $v_{S}$, a physical singlet scalar $s(x)$ is defined as

$$S(x) = v_{S} + s(x).$$

In this case, $b_{F} = c_{F} = 0$ with $F = g, W, Z, \gamma, b, \tau$, and the SM Higgs signal is diluted by the mixing angle $\alpha$. Only possible modification for the Higgs phenomenology is in the Higgs self coupling, which is beyond the scope of this paper however.

3.2 Hidden sector DM models

In this case, the singlet $s$ comes either from the messengers between the SM and the hidden sectors, or from the remnant of the hidden sector Higgs fields after the spontaneous breaking of hidden sector local gauge symmetry.
3.2.1 Singlet fermion DM

Let us consider a singlet fermion DM model with a singlet Dirac dark matter \( \psi \) and a singlet scalar \( S \) as a messenger between the SM sector and the dark matter sector. Then, the model Lagrangian has 3 pieces, the hidden sector and Higgs portal terms in addition to the SM Lagrangian \([52, 53]\):

\[
L = L_{\text{SM}} + L_{\text{hidden}} + L_{\text{portal}},
\]

(3.3)

where

\[
L_{\text{hidden}} = L_S + L_\psi - \lambda S \overline{\psi} \psi,
L_{\text{portal}} = -\mu H S H^\dagger H - \frac{\lambda H S}{2} S^2 H^\dagger H,
\]

(3.4)

with

\[
L_S = \frac{1}{2} (\partial_\mu S \partial^\mu S - m_S^2 S^2) - \mu_S S - \frac{\mu_3}{3} S^3 - \frac{\lambda_S}{4} S^4,
L_\psi = \overline{\psi} (\partial_\mu \beta - m_\psi) \psi.
\]

(3.5)

It is important to introduce a singlet scalar \( S \) as a messenger in order to have correct thermal relic density of singlet fermion DM and to relax the stringent upper bounds on the Yukawa coupling \( \lambda \) from direct detection experiments \([52, 53]\). The model without the dark sector, namely the SM plus an additional singlet scalar field has been studied in detail in \([56, 59]\). After the singlet \( S(x) \) will develop VEV \( v_S \), a physical singlet scalar \( s(x) \) is defined as

\[
S(x) = v_S + s(x).
\]

The detailed study of this singlet fermion DM with Higgs portal has been presented in Ref. \([52, 53]\).

3.2.2 Higgs portal Abelian vector DM

For the case of hidden sector vector dark matter, one has to introduce a hidden sector Higgs field in order to generate the vector DM mass \([54]\). Let us consider a vector boson dark matter, \( X_\mu \), which is assumed to be a gauge boson associated with Abelian dark gauge symmetry \( U(1)_X \). The simplest model will be without any matter fields charged under \( U(1)_X \) except for a complex scalar, \( \Phi \) (a SM singlet), whose VEV will generate the mass for \( X_\mu \) by the conventional Higgs mechanism:

\[
L_{VDM} = -\frac{1}{4} X_{\mu \nu} X^{\mu \nu} + (D_\mu \Phi)^\dagger (D^\mu \Phi) - \frac{\lambda_\Phi}{4} \left( \Phi^\dagger \Phi - \frac{v_\Phi^2}{2} \right)^2
- \lambda_H \Phi \left( H^\dagger H - \frac{v_H^2}{2} \right) \left( \Phi^\dagger \Phi - \frac{v_\Phi^2}{2} \right),
\]

(3.6)

in addition to the SM lagrangian. The covariant derivative is defined as

\[
D_\mu \Phi = (\partial_\mu + ig_X Q_\Phi X_\mu) \Phi.
\]
where $Q_\Phi \equiv Q_X(\Phi)$ is the $U(1)_X$ charge of $\Phi$.

Assuming that the $U(1)_X$-charged complex scalar $\Phi$ develops a nonzero VEV, $v_\Phi$, and thus breaks $U(1)_X$ spontaneously,

$$\Phi = \frac{1}{\sqrt{2}} (v_\Phi + \varphi(x)).$$

Therefore the Abelian vector boson $X_\mu$ get mass $m_X^2 = g_X^2 Q_\Phi^2 v_\Phi^2$, and the hidden sector Higgs field (or dark Higgs field) $\varphi(x)$ will mix with the SM Higgs field $h(x)$ through Higgs portal of the $\lambda_{H\Phi}$ term. The detailed phenomenological analysis of this model is presented in Ref. [54]. Here, it would suffice to mention that this higgs portal vector DM model is a viable model for CDM and the EW vacuum is stable upto Planck scale without any more new fields. In this paper, we presented this model lagrangian as an example where a singlet scalar appears naturally, by identifying $\Phi(x) \rightarrow S(x)$ and $\varphi(x) \rightarrow s(x)$.

### 3.2.3 Scalar DM with local $Z_2$ symmetry

Let us assume the dark sector has a local $U(1)_X$ gauge which is spontaneously broken into local $Z_2$ symmetry. This can be achieved with two complex scalar fields $\phi_X$ and $X \equiv X_R + i X_I$ in the dark sector with the $U(1)_X$ charges equal to 2 and 1, respectively, in the following lagrangian [92]:

$$\mathcal{L} = \mathcal{L}_{SM} - \frac{1}{4} X_{\mu \nu} X^{\mu \nu} - \frac{1}{2} \epsilon X_{\mu \nu} B^{\mu \nu} + D_\mu \phi_X^* D^\mu \phi_X - \frac{\lambda_X}{4} (\phi_X^* \phi_X - v_\phi^2)^2$$
$$+ D_\mu X^\dagger D^\mu X - m_X^2 X^\dagger X - \frac{\lambda_X}{4} (X^\dagger X)^2 - \left(\mu X^2 + H.c.\right)$$
$$- \frac{\lambda_X H}{4} X^\dagger X H^\dagger H - \frac{\lambda_{\phi X H}}{4} \phi_X^\dagger \phi_X H^\dagger H - \frac{\lambda_{X H}}{4} X^\dagger X \phi_X^\dagger \phi_X (3.7)$$

After the $U(1)_X$ symmetry breaking by nonzero $\langle \phi_X \rangle = v_\phi \neq 0$, the $\mu$–term generates

$$(X^2 + H.c.) = 2(\bar{X}_R^2 - \bar{X}_I^2)$$

which lifts the mass degeneracy between $X_R$ and $X_I$. One could also consider fermion DM with local $Z_2$ in a similar manner. The detailed phenomenology of this model will be presented in Ref. [92].

For the purpose of this paper, it is enough to show that the model above is a well defined scalar DM model with local $Z_2$ symmetry stabilizing the CDM, and improves the usual scalar DM based on $Z_2$ symmetry $S \rightarrow -S$ [93]:

$$\mathcal{L}_{SDM} = \frac{1}{2} \partial_\mu S \partial^\mu S - \frac{1}{2} m_S^2 S^2 - \frac{\lambda_{S H}}{2} S^2 H^\dagger H - \frac{\lambda_S}{4!} S^4 .$$

(3.8)

Note that the scalar CDM model with DM being stabilized by local $Z_2$ is much more complicated than the usual $Z_2$ scalar CDM models. In particular, there appears a singlet scalar $s(x)$ that mixes with the SM Higgs boson after the hidden $U(1)_X$ symmetry breaking. Therefore the Higgs phenomenology will be necessarily modified due to the mixing with the singlet scalar $s(x)$.
3.2.4 Parametrizations of hidden sector DM models

In all the cases described in the previous subsubsections, we end up with a singlet scalar $s$ that mixes with the SM Higgs boson. Therefore $c_F = 0$ for $F$ being the SM fields $F = V, f, g, \gamma$, and only $c_\chi$ and $c_{h^2}$ are nonzero, where $\chi$ denotes CDM field (scalar, fermion or vector CDM in the hidden sector with Higgs portal). In this case, there would be universal signal reductions in all the decay channels of $H_i$, whether the $H_i \rightarrow \chi \chi$ is kinematically open or not. Also note that nonzero $c_\chi$ can lead to invisible decays of $H_i$ after the mixing, if kinematically allowed. These two simple predictions of the hidden sector DM models can be tested at the LHC if more data on the Higgs boson is accumulated in the future.

3.3 Dilaton and Radion

Both dilaton in technicolor models with approximate scale symmetry and the radion in the RS scenario couple to the scale anomaly of the theory, which is nothing but the trace of energy-momentum tensor $T_{\mu}$. Usually it is assumed that

$$T_{\mu} = 2\mu_H H^\dagger H + \sum_f m_f \bar{f} f - 2m_W^2 W^+ W^- - m_Z^2 Z^\mu Z^\mu + \frac{\beta}{g} G_{\mu \nu} G^{\mu \nu}$$  \hspace{1cm} (3.9)

namely the trace of the energy-momentum tensor after EWSB and only the unbroken subgroup of the SM gauge symmetry (namely $SU(3)_C \times U(1)_{em}$) is imposed. This ansatz might be fine if one were interested in the case without fundamental Higgs such as technicolor models or other models with dynamical EWSB.

On the other hand, if the dilaton or the radion appears before the EWSB, it would be more appropriate to impose the full SM gauge symmetry for the dilaton/radion couplings to the SM fields. If we impose the full $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge symmetry, the SM Lagrangian has only one operator that breaks the scale symmetry explicitly, namely the Higgs boson mass term:

$$T_{\mu} = 2\mu_H H^\dagger H + \frac{\beta}{g} G_{\mu \nu} G^{\mu \nu}$$  \hspace{1cm} (3.10)

Therefore, the dilaton will couple to the $H^\dagger H$ operator only at classical level, and to the scale anomaly at quantum level. It turns out that two prescriptions for $T_{\mu}$, eq.(3.9) vs. eq.(3.10), have vastly different phenomenological consequence in the Higgs-dilaton sector [80]. Also note that dilaton couplings have extra pieces which are different from other singlet scalar cases, namely nonzero $b_{dW}$ and $b_{dZ}$. Phenomenological analysis with the SM Higgs mixing with dilation or radion will be presented in detail elsewhere [80], and not covered in this paper.

3.4 Extra vectorlike fermions for enhanced $H \rightarrow \gamma\gamma$

In this case, there could be an extra singlet scalar field in order to have renormalizable interactions between the vector like fermions.
3.4.1 \( SU(2)_L \) singlet vectorlike lepton

If the vectorlike fermions are colorless \( SU(2)_L \) singlets \( S^\pm_L \) and \( S^\pm_R \) with electric charge \( Q_e = -1 = Y \), it can not directly couple to the SM Higgs doublet. Therefore one has to introduce a singlet scalar field \( S \):

\[
L = S^\dagger_L i \, \not{D} S^\dagger_L + S^\dagger_R i \, \not{D} S^\dagger_R - \left\{ S_L^\dagger \left( m_S + \lambda S \right) S_R^\dagger + y_{S_i} \, \overline{t}_L H S_R^\dagger + H.c. \right\} \quad (3.11)
\]

Note that \( S^- \) will mix with the \( e_{Rj} \) after EWSB, and \( S^- \) will decay to \( hl^- \) through \( y_{Si} \) couplings, and its contribution to \( H \rightarrow \gamma \gamma \) is suppressed by the Yukawa coupling \( y_{Si} \) that is presumably small. Also \( S^- \) can not carry color charge, since we cannot introduce \( y_{Si} \) terms which make \( S^- \) decay.

Note that the vectorlike charged scalar \( S^\pm \) will generate \( s \rightarrow \gamma \gamma \) and \( s \rightarrow Z\gamma \) at one loop level. Therefore both \( c_\gamma \) and \( c_{Z\gamma} \) are nonzero, and they would affect \( H \rightarrow \gamma \gamma \) and \( H \rightarrow Z\gamma \) through the mixing between the SM Higgs boson and the singlet scalar \( s \).

3.4.2 \( SU(2)_L \) doublet plus singlet vectorlike leptons

If we consider just one vectorlike lepton \( L \) and \( R \) in the \( SU(2)_L \) doublet representation with \( Y = -1/2 \),

\[
L = (N^0, E^-)^T_L, \quad R = (N^0, E^-)^T_R \quad (3.12)
\]

one has to introduce two singlet fermions \( S_L^- \) and \( S_R^- \) in order to write couplings to SM Higgs field \([86, 88]\):

\[
\mathcal{L}_{\text{mass}} = -m_1 S_L^- S_R^- - m_2 L R - \sqrt{2} y_{12} S_L^\dagger H^\dagger R - y_{21} L H S_R^- + H.c. \quad (3.13)
\]

The model contains one neutral Dirac lepton \( N^0 \) and two Dirac leptons \( S^- \) and \( E^- \) with electric charge \( Q_e = -1 \). The mass matrix of these new leptons is written as

\[
\mathcal{M}(v) = \begin{pmatrix} m_1 & y_{12} v \\ y_{21} v & m_2 \end{pmatrix} \quad (3.14)
\]

in the \((S^-, E^-)\) basis. The new Yukawa couplings \( y_{12} \) and \( y_{21} \) can be complex in general, and there could be CP violation in \( H \rightarrow \gamma \gamma \) as pointed out by M. Voloshin in Ref. [86].

In this model, however, the neutral component \( N^0 \) which would be absolutely stable and could overclose the universe. This trouble can be resolved if we introduce a real singlet scalar \( S \) as in the previous subsection, adding the following Yukawa couplings to the above lagrangian:

\[
\delta \mathcal{L}_{\text{mass}} = -\lambda_S S \, \not{D} R - \lambda_S S L \, \not{D} R + H.c. \quad (3.15)
\]

After \( S(x) \) gets a nonzero VEV, \( s \) will mix with the SM Higgs boson, and \( N^0 \overline{N}^0 \) can pair annihilate into the SM fields through Higgs portal:

\[
N^0 \overline{N}^0 \rightarrow s^* \rightarrow h^* \rightarrow \text{SM particles}
\]

and can be thermalized.
In this model for vectorlike leptons, both $c_γ$ and $c_{Zγ}$ can be nonzero, depending on the SM quantum numbers. Also $c_{f'} \neq 0$, where $f'$ denotes the vector like fermions. The mixing due to the SM charged lepton and extra $S^-$ can induce nonzero $b_γ$ and $b_{Zγ}$, proportional to $y_{Si}$, but we ignore it here for simplicity assuming $y_{Si}$ is small enough.

### 3.5 Extra charged vector bosons

In this subsection, let us describe some models with extra charged vector bosons that can contribute to $H \to γγ$.

In Ref. [81, 83], extra charged vector bosons $W'$ in the $SU(2)_L$ triplet representation were considered as a possible explanation for the enhancement of $H \to γγ$ using the effective operator:

$$O_{W'} = \frac{1}{2} c_{W'} g^2 H^\dagger H W_{\mu}^\dagger W_{\mu}^-$$

with $m_{W'}^2 = m_0^2 + c_{W'} g^2 v^2$. However, this operator is not really renormalizable, since we need to introduce a new Higgs field $H'$ that gives mass to the $W'$. After symmetry breaking from $\langle H' \rangle = v' \neq 0$, there may be a remnant of symmetry breaking in terms of a SM singlet scalar $s$ that couples to $W'_{\mu} W_{\mu}^-$. In principle, these $W'$ can mix with the SM weak gauge bosons $W^\pm$ in general. For example, this is well known in the $SU(2)_R$ extension of the SM. In this paper we will ignore this mixing between the $W'$ and $W^\pm$ for simplicity.

Another simple case would be the $SU(2)_R$ extension of the SM, where there appear 3 more gauge bosons, $W_{R}^\pm$ and $W_{R}^0$, that would mix with the SM weak gauge bosons $W^\pm$ and $Z^0$. In this case, there will be a singlet scalar $s$ that couples to $W_{R}^\pm$ after $SU(2)_R$ symmetry breaking. Also this $s$ will mix with the SM Higgs boson.

We assume that $W'$ is significantly heavier than the SM $W$ boson and the physical Higgs boson of mass 125 GeV. In this case, all the $c_F$'s are zero, except for $c_γ$ due to the $s \to γγ$ through $W'$ loop and $c_{h2}$.

### 3.6 Extra charged scalar bosons

One can consider extra charged scalar bosons. In this case, it is not mandatory for an extra singlet scalar to appear in the models, but we can introduce them if one wishes, with the following operators:

$$H^\dagger H \phi_\alpha^1 \phi_\alpha^\dagger, \quad S \phi_\alpha^1 \phi_\alpha^\dagger, \quad S^2 \phi_\alpha^1 \phi_\alpha^\dagger,$$

where $\phi_\alpha$ are new scalar bosons with nonzero electric charge and/or color charge. Then the SM Higgs properties can modified only by the higher dimensional operators through modified $b_F$ with $F = g$ or $γ$ depending on the SM charges of the scalar bosons, where as $c_γ$ and/or $c_γ$ can be nonzero.

### 3.7 Summary

We assumed that the extra singlet scalar $s$ is an EW singlet, and does not participate in EWSB. The resulting possible new physics contributions are parametrized in terms of $c_γ$ and $c_{Zγ}$.

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3For vectorlike quarks, both $c_γ$ could be also nonzero in addition to $c_γ$ and $c_{Zγ}$.

4Here, charged scalar boson means scalar bosons in nontrivial representation of the SM gauge group.
$c_i$’s, and we tabulate the nonzero $c$’s in Table 1. If the singlet $s$ was coming from another Higgs doublet which break EW symmetry, we should have introduced another parameters (analogous to $\tan \beta$ in 2HDMs) in addition to the mixing angle $\alpha$, and the analysis will be more involved than presented in this paper. There are a number of analysis within 2HDMs in the literature [94–106], and we do not pursue this possibility in this paper.

| Model                        | Nonzero $c_F$’s                                      |
|-------------------------------|-----------------------------------------------------|
| Pure Singlet Extension        | $c_{h^2}$                                           |
| Hidden Sector DM              | $c_x, c_{h^2}$                                      |
| Dilaton                      | $c_y, c_W, c_Z, c_\gamma, c_{h^2}$                  |
| Vectorlike Quarks             | $c_y, c_\gamma, c_Z\gamma, c_{h^2}$                |
| Vectorlike Leptons            | $c_\gamma, c_{Z\gamma}, c_{h^2}$                   |
| New Charged Vector bosons     | $c_\gamma, c_{h^2}$                                 |
| Extra charged scalar bosons   | $c_y, c_\gamma, c_Z\gamma, c_{h^2}$                |

Table 1. Nonvanishing $c_F$’s in various BSM’s with an extra singlet scalar boson.

4 Fitting procedure

4.1 General parameterization and data

Following the LHC Higgs working group [107], we find it useful to introduce model-independent notations to describe signal strength data:

$$\kappa_i^2 = \frac{\Gamma(h \to i)}{\Gamma(h \to i)_{SM}}, \quad \kappa_H^2 = \frac{\Gamma_{tot}}{\Gamma_{tot,SM}}. \quad (4.1)$$

What is actually measured at LHC then can be written as

$$R_j^i = R \left( \sigma(i \to h) \frac{\Gamma(h \to j)}{\Gamma_{tot}} \right) = \frac{\kappa_i^2 \kappa_j^2}{\kappa_H^2} = \tilde{\kappa}_i^2 \tilde{\kappa}_j^2, \quad \mu_j = \sum_i R_j^i \quad (4.2)$$

where we define $\tilde{\kappa}_i^2 = \kappa_i^2 / \kappa_H$. Signal strength in the final state $j$ for Higgs production modes combined is denoted by $\mu_j$. We discuss how our theory parameterization and the $\kappa_i$ notation are related in the later section.

As introduced in previous section, we parameterize effective couplings of Higgs-like boson by constants $b_i$, $c_i$ and the mixing angle $\alpha$ (and possible non-standard decay modes introduced below). We emphasize once again that these parameters do not multiplicatively parameterize signal strengths. For final states to which the Higgs boson couples at tree-level (such as $WW$, $ZZ$ and fermion pairs),

$$\kappa_i^2 = \frac{\Gamma(h \to i)}{\Gamma(h \to i)_{SM}} = (b_i c_\alpha - c_i s_\alpha)^2 \quad \text{for } i = WW, ZZ, f \bar{f} \quad (4.3)$$
where we denote $\cos \alpha = c_\alpha$ and similarly $\sin \alpha = s_\alpha$. For $gg$, $\gamma\gamma$ and $\gamma Z$ final states whose couplings to the Higgs boson are loop-induced,

$$\kappa_g^2 = \frac{\Gamma(h \rightarrow gg)}{\Gamma(h \rightarrow gg)_{SM}} = (b_g c_\alpha - c_g s_\alpha)^2 = (c_\alpha (b_t C_t + \Delta b_g) - c_g s_\alpha)^2$$  \hspace{1cm} (4.4)

$$\kappa_\gamma^2 = \frac{\Gamma(h \rightarrow \gamma\gamma)}{\Gamma(h \rightarrow \gamma\gamma)_{SM}} = (b_\gamma c_\alpha - c_\gamma s_\alpha)^2 = (c_\alpha (b_W B_W + b_t B_t + \Delta b_\gamma) - c_\gamma s_\alpha)^2$$  \hspace{1cm} (4.5)

where the relative loop-functions of $W$ boson and top quark for $m_h = 125$ GeV are given in eq.(2.9). Modifications to loop-induced decay eq.(4.5) have several contributions; (i) from scalar mixing denoted by $\alpha$, (ii) one inherit from singlet scalar couplings denoted by $c_g, c_\gamma$, (iii) from modification of top and/or $W$ boson couplings in the loop denoted by $b_t$ and $b_W$ weighted by loop factors introduced above, and (iv) last from some new physics effects directly modifying Higgs interaction eigenstate coupling denoted by $\Delta b_g, \Delta b_\gamma$.

We also allow Higgs may have non-standard decay modes (such as invisible decay or flavor violating decay mode) which we parameterize by branching ratio into these modes, $BR_{nonSM}$. Total width is then written as

$$\kappa_H^2 = \frac{\Gamma_{tot}}{\Gamma_{tot}_{SM}} = \sum_{i \in SM} \kappa_i^2 BR(h \rightarrow i)_{SM} \left(1 - BR_{nonSM}\right).$$  \hspace{1cm} (4.6)

We use the most up-to-date CMS and ATLAS data for $h \rightarrow ZZ, WW, \gamma\gamma, \tau\tau, b\bar{b}$ which is tabulated in Table 2. We naively use the best-fit signal strengths for each channel obtained with each best-fit $m_h$; refer to official notes of ATLAS [118] and CMS [119] for more dedicated study of signal strengths obtained with the same $m_h$. For the $\gamma\gamma$ and $ZZ$ modes, both ATLAS and CMS have analyzed the contributions from different production modes, but we use the combined results of multiple production modes since the uncertainty on the contribution of different production modes is rather large at the moment and the full covariance matrices on the errors are not available. We assume that all data are from $gg$-fusion and vector-boson fusion (VBF) in proportion to their production rates (except for $bb$ mode which is assumed to be purely from $W/Z$ associate production)

$$R(\sigma(pp \rightarrow h)) = \kappa_g^2 A_g + \kappa_W^2 A_W + \kappa_Z^2 A_Z$$  \hspace{1cm} (4.7)

where weighting factors for $m_h = 125$ GeV are

$$A_g = \frac{\sigma(ggF)}{\sigma(ggF) + \sigma(VBF)} \simeq 0.925, \quad A_W + A_Z = \frac{\sigma(VBF)}{\sigma(ggF) + \sigma(VBF)} \simeq 0.075.$$  \hspace{1cm} (4.8)

7 TeV and 8 TeV production rates are weighted-summed in proportion to luminosities accumulated in data we use. Likewise, $Vh$ associate production proceeds via either $W$ or $Z$ boson,

$$R(\sigma(pp \rightarrow Vh)) = \kappa_W^2 A_W' + \kappa_Z^2 A_Z'$$  \hspace{1cm} (4.9)

where $A_W' + A_Z' = 1$. We numerically checked that small mixed-in of VBF denoted by $A_W + A_Z$ does not significantly affect fit results. In this work, we simply ignore $A_W + A_Z$.
| channel | luminosity ($fb^{-1}$) | $\mu$ | ref. |
|---------|----------------|------|-----|
| CMS     |                  |      |     |
| $\gamma\gamma$ | 24.7 | $0.78^{+0.28}_{-0.26}$ | [108] |
| $ZZ$    | 24.7            | $0.91^{+0.30}_{-0.24}$ | [109] |
| $WW$    | 24.7            | $0.76^{+0.21}_{-0.21}$ | [110] |
| $\tau\tau$ | 24.3 | $1.1^{+0.4}_{-0.4}$ | [111] |
| $\bar{b}b$ | 17             | $1.3^{+0.7}_{-0.6}$ | [112] |
| ATLAS   |                  |      |     |
| $\gamma\gamma$ | 25 | $1.65^{+0.35}_{-0.30}$ | [113] |
| $ZZ$    | 25              | $1.7^{+0.50}_{-0.40}$ | [114] |
| $WW$    | 25              | $1.01^{+0.31}_{-0.31}$ | [115] |
| $\tau\tau$ | 18            | $0.7^{+0.7}_{-0.7}$ | [116] |
| $\bar{b}b$ | 18        | $-0.4^{+1.06}_{-1.06}$ | [117] |

Table 2. Signal strength data we use. 7TeV and 8TeV results are combined and all subcategories of each final states are combined. We use best-fit signal strengths for each channel obtained with each best-fit $m_h$.

| signal strength $\mu_i$ | $pp \rightarrow h \rightarrow \gamma\gamma, WW, ZZ, \tau\tau$ | $pp \rightarrow Vh \rightarrow V\bar{b}b$ |
|--------------------------|---------------------------------------------------|-------------------------------------|
| $(\hat{\kappa}_g^2 A_g + \hat{\kappa}_W^2 A_W + \hat{\kappa}_Z^2 A_Z) \hat{\kappa}_i^2$ | $(\hat{\kappa}_W^2 A_W' + \hat{\kappa}_Z^2 A_Z') \hat{\kappa}_i^2$ |

Table 3. General parameterization of signal strength data that we use in terms of model independent $\hat{\kappa}$ parameters introduced in eq.(4.1) and eq.(4.2). $A_i$ and $A_i'$ are defined in eq.(4.7) and eq.(4.9).

We use $m_h = 125\text{GeV}$ when necessary.

In all, we tabulate general parameterization of Higgs signal strengths we use in terms of $\kappa$ parameters in Table 3.

4.2 SM fit

First of all, pure SM fit is performed. We use MINUIT package [120] to carry out $\chi^2$-fit:

$$\chi^2/\nu = 12.01/10 = 1.20 \quad \text{(both)} \quad (4.10)$$

$$\chi^2/\nu = 2.33/5 = 0.466 \quad \text{(CMS)} \quad (4.11)$$

$$\chi^2/\nu = 9.69/5 = 1.94 \quad \text{(ATLAS).} \quad (4.12)$$

4.3 Preliminary: fits without extra singlet scalar

In this subsection, we use only $b_i$ to carry out best-fit. When new particles directly couple to Higgs boson in a gauge invariant way under the SM gauge group and modify Higgs
couplings, the model falls into this category. This study also allows us to compare our fit results with other results available in literature, and moreover to discuss similarities and differences between individual ATLAS and CMS data.

We simply assume custodial symmetry for the SM Higgs couplings to the weak gauge bosons:

$$b_W = b_Z \equiv b_V.$$  \hfill (4.13)

We also simply treat all fermions couplings universally

$$b_t = b_b = b_\tau \equiv b_f$$  \hfill (4.14)

although we checked that allowing $b_t$ to float independently would not qualitatively modify our statements.

We consider following four cases: (results are also summarized in Table 5)

- $\{\Delta b_\gamma\}$: This may represent some models of extra leptons [87] or $W'$ or extra charged scalar. In this case,

$$\kappa_\gamma^2 = b_\gamma^2 = (1 + \Delta b_\gamma)^2, \quad \kappa_V^2 = \kappa_f^2 = 1, \quad \kappa_H^2 \simeq 1$$  \hfill (4.15)

giving best-fit

$$\Delta b_\gamma = 0.090^{+0.0889}_{-0.0999}, \quad \chi^2/\nu = 11.19/9 = 1.24 \quad \text{(both)}$$  \hfill (4.16)

$$\Delta b_\gamma = -0.117^{+0.147}_{-0.162}, \quad \chi^2/\nu = 1.71/4 = 0.428 \quad \text{(CMS)}$$  \hfill (4.17)

$$\Delta b_\gamma = 0.28^{+0.134}_{-0.118}, \quad \chi^2/\nu = 4.99/4 = 1.25 \quad \text{(ATLAS)}$$  \hfill (4.18)

We show best-fit using both ATLAS and CMS data, as well as best-fit obtained from individual data. $\mu_\gamma = 1.19$ (both).

- $\{\Delta b_g, \Delta b_\gamma\}$: This may represent some models of extra quarks [87]. In this case,

$$\kappa_g^2 = b_g^2 = (1 + \Delta b_g)^2, \quad \kappa_\gamma^2 = b_\gamma^2 = (1 + \Delta b_\gamma)^2, \quad \kappa_{mix}^2 = 1,$$

$$\kappa_H^2 \simeq \kappa_g^2 \text{Br}(h \to gg) + \text{Br}(h \to \text{others})$$  \hfill (4.19)

giving best-fit ($\mu_\gamma = 1.19$ (both))

$$\Delta b_g = -0.0180^{+0.0559}_{-0.0577}, \quad \Delta b_\gamma = 0.107^{+0.0916}_{-0.1001}, \quad \chi^2/\nu = 11.13/8 = 1.39 \quad \text{(both)}$$  \hfill (4.21)

$$\Delta b_g = -0.078^{+0.0760}_{-0.0784}, \quad \Delta b_\gamma = -0.048^{+0.157}_{-0.175}, \quad \chi^2/\nu = 0.859/3 = 0.286 \quad \text{(CMS)}$$  \hfill (4.22)

$$\Delta b_g = 0.11^{+0.0867}_{-0.0830}, \quad \Delta b_\gamma = 0.17^{+0.117}_{-0.113}, \quad \chi^2/\nu = 4.14/3 = 1.38 \quad \text{(ATLAS)}$$  \hfill (4.23)
• \{ b_V, b_f \}: This may represent some simple composite Higgs models, see e.g. Ref.[25, 121]. In this case,
\[ \kappa_g^2 = b_g^2, \quad \kappa_\gamma^2 = (b_V B_W + b_f B_t)^2, \quad \kappa_V^2 = b_V^2, \quad \kappa_f^2 = b_f^2 \]  
(4.24)

\[ \kappa_H^2 = \sum_i \kappa_i^2 \text{BR}(i)_{SM}. \]  
(4.25)
giving best-fit
\[ b_V = 1.031^{+0.0682}_{-0.0688}, \quad b_f = 0.962^{+0.124}_{-0.124}, \quad \chi^2/\nu = 11.74/8 = 1.47 \quad \text{(both)} \]  
(4.26)
\[ b_V = 0.898^{+0.081}_{-0.081}, \quad b_f = 1.021^{+0.137}_{-0.154}, \quad \chi^2/\nu = 0.808/3 = 0.27 \quad \text{(CMS)} \]  
(4.27)
\[ b_V = 1.345^{+0.162}_{-0.144}, \quad b_f = 0.808^{+0.144}_{-0.117}, \quad \chi^2/\nu = 4.52/3 = 1.51 \quad \text{(ATLAS)} \]  
(4.28)

• \{ b_V, b_u, b_d, (c_\alpha) \}: We relax assumptions of universal fermion couplings, and allow all up-type quark couplings are modified by \( b_u \), and all up-type quarks and charged leptons are modified by \( b_d \). We further assume that
\[ b_V \leq 1. \]  
(4.29)

This ansatz can represent various two-Higgs doublet model when both of two Higgs doublets develop nonzero VEV’s and contribute to the EWSB. When \( c_\alpha \) is allowed, this ansatz may further consider as a doublet+singlet scenario such as the scalar sector of next-to-MSSM (NMSSM). Generally,
\[ \kappa_g^2 = b_g^2 c_\alpha^2, \quad \kappa_\gamma^2 = (b_V B_W + b_u B_t)^2 c_\alpha^2, \quad \kappa_V^2 = b_V^2 c_\alpha^2, \quad \kappa_u^2 = b_u^2 c_\alpha^2, \quad \kappa_d^2 = b_d^2 c_\alpha^2 \]  
(4.30)

\[ \kappa_H^2 = \sum_i \kappa_i^2 \text{BR}(i)_{SM}. \]  
(4.31)

With \( c_\alpha = 1 \) fixed, best-fit is
\[ b_V = 1.0_{-0.0601}, \quad b_u = 0.969^{+0.0632}_{-0.0647}, \quad b_d = 0.938^{+0.0899}_{-0.0788}, \]  
(4.32)
\[ \chi^2/\nu = 11.86/7 = 1.69 \quad \text{(both)} \]  
(4.33)

Interestingly, \( b_V = 1 \) gives the best-fit. Another interesting result is that allowing \( c_\alpha \) doesn’t change best-fit results; \( c_\alpha = 1 \) gives (the same) best-fit. For reference, if we allowed \( b_V \) to vary above 1, we would have obtained \( b_V = 1.05 \) as a best-fit.

• \{ \Delta b_g, \Delta b_\gamma, b_V, b_f \}: In this case,
\[ \kappa_g^2 = b_g^2 = (b_f + \Delta b_g)^2, \quad \kappa_\gamma^2 = b_\gamma^2 = (b_V B_W + b_f B_t + \Delta b_\gamma)^2, \quad \kappa_V^2 = b_V^2, \quad \kappa_f^2 = b_f^2 \]  
(4.34)

\[ \kappa_H^2 = \sum_i \kappa_i^2 \text{BR}(i)_{SM}. \]  
(4.35)
giving best-fit
\[ \Delta b_g = 0.041^{+0.0596}_{-0.0621}, \quad \Delta b_\gamma = 0.117^{+0.0927}_{-0.101}, \quad b_V = 0.941^{+0.0569}_{-0.0581}, \quad b_f = 0.961^{+0.116}_{-0.127} \]  
(4.36)
Table 4. Comparison of our fit results with results available in other literature. Only results based on up-to-date data after Moriond 2013 are compared. We sometimes re-interpret other’s results in accordance with our notation. If only best-fit figure is available, we cite relevant figure and reference. Cases that are not shown here do not have equivalent results in literature.

| | our fits | fits in other refs. |
|---|---|---|
| $(\Delta b_g, \Delta b_\gamma)$ | $(-0.0180^{+0.0559}_{-0.0577}, 0.107^{+0.0916}_{-0.100})$ | $(0.12^{+0.11}_{-0.08}, 0.18^{+0.12}_{-0.07})$ [48] |
| (ATLAS-only) | $(0.11^{+0.0867}_{-0.0830}, 0.17^{+0.117}_{-0.113})$ | $(0.08^{+0.14}_{-0.06}, 0.23^{+0.16}_{-0.13})$ [118] |
| $(b_V, b_f)$ | $(1.031^{+0.0682}_{-0.0688}, 0.962^{+0.124}_{-0.124})$ | $(1.03^{+0.06}_{-0.07}, 0.84^{+0.15}_{-0.12})$ [48] |
| (ATLAS-only) | $(1.345^{+0.162}_{-0.144}, 0.808^{+0.144}_{-0.117})$ | $(1.13^{+0.08}_{-0.06}, 0.96^{+0.17}_{-0.16})$ [118] |

Table 5. Best-fit results using $b_i$ only from both CMS and ATLAS data as well as individual. Errors are shown in text.

| | both | CMS | ATLAS |
|---|---|---|---|
| SM | $\chi^2/\nu = 12.01/10 = 1.20$ | $2.33/5 = 0.466$ | $9.69/5 = 1.94$ |
| $(\Delta b_\gamma)$ | $(-0.090)$ | $(-0.117)$ | $(0.28)$ |
| & $11.19/9 = 1.24$ | $1.71/4 = 0.428$ | $4.99/4 = 1.25$ |
| $(\Delta b_g, \Delta b_\gamma)$ | $(-0.018, 0.107)$ | $(-0.078, -0.048)$ | $(0.11, 0.17)$ |
| & $11.13/8 = 1.39$ | $0.859/3 = 0.286$ | $4.14/3 = 1.38$ |
| $(b_V, b_f)$ | $(1.031, 0.962)$ | $(0.898, 1.021)$ | $(1.345, 0.808)$ |
| & $11.74/8 = 1.47$ | $0.808/3 = 0.27$ | $4.52/3 = 1.51$ |
| $(b_V \leq 1, b_u, b_d)$ | $(1.0, 0.969, 0.938)$ | $11.86/7 = 1.69$ |
| $(\Delta b_g, \Delta b_\gamma, b_V, b_f)$ | $(0.041, 0.117, 0.941, 0.961)$ | $11.07/6 = 1.85$ |

We do not consider fitting to individual ATLAS and CMS data here because there are too small number of degrees of freedom ($\nu = 1$) which may not allow meaningful statistical interpretation of fit results.

We compare our fit results with other results available in literature. For proper comparison, we use other results based on up-to-date data after Moriond 2013. As tabulated in Table 4, we obtain fairly good agreement on central values and sizes of uncertainties. Some differ-
ence may originate from different datasets used; we are using only inclusively combined data while most other literatures uses more individual production modes. Another important difference is that we naively use best-fit signal strengths each fitted with different best-fit $m_h$ values while ATLAS official document [118] uses signal strengths obtained all with the same $m_h = 125.5$ GeV. For example, our $\mu_{ZZ} = 1.7$ is obtained with best-fit $m_h = 124.3$ GeV while ATLAS uses $\mu_{ZZ} = 1.5$ [118] obtained with assumed $m_h = 125.5$ GeV. This causes some difference in best-fit $b_V$ values. We note that official study of coupling fit from CMS [119] is not updated with up-to-date data yet.

From cases (1), (2) and (3) that are also tabulated in Table 5, one can also learn about general trends of present ATLAS and CMS data. Generally, CMS data has better fit to $\{b_i\}$ parameterization than ATLAS data does. CMS data generally prefers the suppression of signal strengths while ATLAS data prefers the enhancement. However, each best-fit parameters are marginally consistent with each other (between CMS-only and ATLAS-only fit results) within combined $2\sigma$ uncertainties although $b_V$ best-fit results are a bit more discrepant. Consequently, best-fits in terms of $\{b_i\}$ to both ATLAS and CMS data simultaneously is not improved significantly from pure SM fit. This may or may not mean that $\{b_i\}$ parameterization is disfavored partly because fits to individual data are indeed improved.

5 Theories of an extra singlet scalar mixed in

We now consider an extra singlet scalar mixing with SM Higgs. In the remainder of the section, we make following simplifying but general enough assumptions:

- Singlet has only loop-induced couplings to SM particles; only $c_g, c_\gamma$ are non-zero among all $c_i$.
- However, singlet may have non-standard decay modes: $BR_{\text{nonSM}} \geq 0$. This parameter also considers a possibility that $m_{H_2} \leq m_{H_1}/2 = 62.5$ GeV.
- Higgs couplings are not directly modified: all $b_i = 1$, $\Delta b_i = 0$. But we discuss briefly the case with non-zero $\Delta b_i$ in Sec.5.5.

With these assumptions, we have following three model-independent parameters that can parameterize signal strength data

$$\{ \hat{\kappa}_g, \hat{\kappa}_\gamma, \hat{\kappa}_{\text{mix}} \} \tag{5.1}$$

where $\kappa_{\text{mix}}$ globally parameterizes all SM decay widths other than $gg$ and $\gamma\gamma$ modes defined as

$$\kappa_{\text{mix}}^2 = \frac{\Gamma(H_1 \to WW)}{\Gamma(h \to WW)_{SM}}. \tag{5.2}$$

Again, $\hat{\kappa}_{\gamma}^2 = \kappa_{\gamma}^2 / \kappa_H$. See Table 6 for how $\kappa$s contribute to signal strengths, where different production modes are weighted-summed with production ratio as discussed previously. This parameterization has not been considered yet by the LHC Higgs working group [107], and this type of model has only been briefly studied in Refs. [16, 26].
In this section, we consider several special cases of this general parameterization. All best-fit results are summarized in Table 7.

5.1 How likely are signals universally modified? : \{ \alpha, BR_{\text{nonSM}} \}

We begin by evaluating how likely that all Higgs signal strengths are universally modified. Universal modification can arise through singlet portal mixing and/or non-SM decay width.

In this case, general parameterization introduced above are simplified, and one common parameter

$$\hat{\kappa}_{\text{univ}}^2 \equiv \frac{\kappa_{\text{univ}}^2}{\kappa_H}$$

universally parameterizes all Higgs signal strengths as

$$\mu_i = \kappa_{\text{univ}}^2 \frac{\kappa_{\text{univ}}^2}{\kappa_H} = \hat{\kappa}_{\text{univ}}^2 \hat{\kappa}_{\text{univ}}^2.$$  (5.4)

This parameterization shows that it is difficult to know whether the universal modification of signal strengths is originated from \(\kappa_{\text{univ}}^2\)(singlet portal mixing) or \(\kappa_H^2\)(non-SM decay mode). Best-fit is

$$\hat{\kappa}_{\text{univ}}^2 = 1.012^{+0.0517}_{-0.0549}, \quad \chi^2/\nu = \frac{11.96}{9} = 1.33 \quad \text{(both)}$$  (5.5)

$$\hat{\kappa}_{\text{univ}}^2 = 0.930^{+0.0675}_{-0.0710}, \quad \chi^2/\nu = \frac{1.25}{4} = 0.31 \quad \text{(CMS)}$$  (5.6)

$$\hat{\kappa}_{\text{univ}}^2 = 1.149^{+0.0762}_{-0.0812}, \quad \chi^2/\nu = \frac{6.53}{4} = 1.63 \quad \text{(ATLAS)}$$  (5.7)

1, 2\(\sigma\) ranges of best-fit parameters are defined as ranges for \(\Delta \chi^2 = 1, 4\) from a single variable cumulative distribution function of \(\chi^2\) assuming no correlations, where \(\chi^2 = \chi_{\text{min}}^2 + \Delta \chi^2\). Again, fit is not improved from pure SM fit if both data are used while fits to individual data are improved. Almost no universal modification is preferred. This is partially resulted from the fact that CMS prefers suppression while ATLAS prefers enhancement.

In the end, there are two theory parameters \{ \alpha, BR_{\text{nonSM}} \} for the universal modification of Higgs signal strengths, with the following relations:

$$\kappa_{\text{univ}}^2 = c_\alpha^2, \quad \kappa_H^2 = \frac{c_\alpha^2}{1 - BR_{\text{nonSM}}}.$$  (5.8)

The 2D contours of \(\Delta \chi^2\) in the theory parameters plane is shown in Fig. 2. If either \(c_\alpha = 1\) or \(BR_{\text{nonSM}} = 0\) is fixed,

$$BR_{\text{nonSM}} = -0.0241^{+0.108}_{-0.107} \quad \text{for } c_\alpha = 1$$  (5.9)

$$c_\alpha = 1.012^{+0.0517}_{-0.0549} \quad \text{for } BR_{\text{nonSM}} = 0.$$  (5.10)
Thus, \( BR_{nonSM} \leq 18.8\% \) at 95\%C.L. if \( c_\alpha = 1 \) fixed as also similarly obtained by Ref.[43, 46, 48] using up-to-date data. Likewise, \( c_\alpha \geq 0.904 \) at 95\%C.L. if \( BR_{nonSM} = 0 \) fixed.

Suppose future data is \( R^i_j = 1.0 \pm 0.2 \) (0.1) for all 10 channels we are considering. This would be a perfect case for the SM Higgs boson. The best-fit then would yield

\[
\hat{\kappa}^2_{univ} = 1.0^{+0.0311(0.0157)}_{-0.0321(0.0159)}.
\]

(5.11)

Corresponding 2D \( \Delta \chi^2 \) contours are also shown in dashed curves in Fig. 2 to illustrate how much we are improved, and how much we are still not able to probe. The best fit would imply \( BR_{nonSM} \leq 12.4\, (6.2)\% \) at 95\%C.L. if \( c_\alpha = 1 \) fixed, or \( c_\alpha \geq 0.94\, (0.97) \) if \( BR_{nonSM} = 0 \) fixed. This discussion may help us grasp how well one can constrain universal-suppression models with future data.

### 5.2 Models with extra leptons or \( W' \) : \{\( \alpha, c_\gamma \)\}

Extra leptons or \( W' \) induce singlet couplings to photons at one-loop. The free parameters to fit are \{\( \alpha, c_\gamma \)\}. In this case,

\[
\kappa^2_\gamma = (c_\alpha - c_\gamma s_\alpha)^2, \quad \kappa^2_g = \kappa^2_{mix} = c_\alpha^2, \quad \kappa^2_H \simeq c_\alpha^2
\]

(5.12)

where we ignored small diphoton decay modes in total width. Therefore, two parameters \( \hat{\kappa}_\gamma \) and \( \hat{\kappa}_{mix} = \hat{\kappa}_g \) can parameterize all signal strengths. Best-fit is

\[
c_\alpha = 0.98_{-0.056}^{+0.06}, \quad c_\gamma = -0.55_{-0.45}^{+0.50}, \quad \chi^2/\nu = 11.1/8 = 1.39
\]

(5.13)

corresponding to \( \mu_\gamma = 1.19 \). Preferred parameter space is shown in Fig. 3 left panel. Although total \( \chi^2 \) of the fit is improved from SM fit, \( \chi^2/\nu \) is not.
Figure 3. $1, 2, 3 \sigma$ ranges of best-fit for $\{\alpha, c_\gamma\}$.

We attempt to interpret the best-fit results in terms of specific underlying models. The example model \cite{88} contains a singlet scalar which couples to two sets of vector-like fermions carrying electric charge 1. The $c_\gamma$ can be written as a function of the lightest fermion mass, $m_L$, and its yukawa coupling to singlet, $x$, as

$$c_\gamma = \frac{v}{\sqrt{2} A_{SM}} \frac{x A_{1/2}(m_{H_1}^2/4m_L^2)}{m_L}. \quad (5.14)$$

Since $c_\gamma$ is fitted from 125GeV data, $m_{H_1} = 125$GeV. The favored parameter space of $m_L$ and $x$ is plotted in Fig. 4 left panel. As SM is consistent with fit and errors are somewhat large, we do not obtain strong constraints on this model. However, one can already observe that the majority of parameter space favored within $\pm 1\sigma$ range involve yukawa couplings stronger than top yukawa and/or fermions lighter than top quarks.

5.3 Models with extra quarks : $\{\alpha, c_\gamma, c_g\}$

Extra quarks can induce singlet couplings to photons as well as gluons at one-loop. The free parameters to fit are $\{\alpha, c_\gamma, c_g\}$, and we have

$$\kappa_{\gamma}^2 = (c_\alpha - c_\gamma s_\alpha)^2, \quad \kappa_g^2 = (c_\alpha - c_g s_\alpha)^2, \quad \kappa_{mix}^2 = c_\alpha^2, \quad \kappa_{H}^2 \simeq 0.0857 \kappa_g^2 + 0.9143 c_\alpha^2 \quad (5.15)$$

Therefore, three parameters $\hat{\kappa}_g$, $\hat{\kappa}_\gamma$ and $\hat{\kappa}_{mix}$ can parameterize all signal strengths. Best-fit is

$$c_g = -0.128^{+0.185}_{-0.174}, \quad c_\gamma = -0.313^{+0.296}_{-0.269}, \quad c_\alpha = 0.947_{-0.0873}, \quad \chi^2/\nu = 11.1/7 = 1.58 \quad (5.16)$$

and various 2D contours of $\Delta \chi^2$ are shown in Fig. 5. Compared to extra lepton models, total $\chi^2$ of the best fit is not improved, whereas somewhat larger mixing is preferred.
Figure 4. ±1σ favored region from best-fits of $c_\gamma$ (black solid) and/or $c_g$ (blue dashed) in models of (Left): vector-like lepton considered in case(2), and (Right): top-like quark considered in case(3).

We take a simple toy model involving a singlet scalar coupling to a single extra quark whose quantum numbers are identical to top quarks. The quarks generate both $c_\gamma$ and $c_g$. With the lightest extra quark mass $m_Q$ and its coupling $x$ to the singlet scalar, loop-induced couplings are

$$c_\gamma = \frac{v}{A_\gamma^{SM}} \frac{x N_e Q_2^2 A_{1/2}(m_{H_1}^2/4m_Q^2)}{m_Q} \quad (5.17)$$

$$c_g = \frac{xv A_{1/2}(m_{H_1}^2/4m_Q^2)}{A_{1/2}(m_{H_1}^2/4m_t^2)m_Q} \quad (5.18)$$

where $m_{H_1} = 125\text{GeV}$. In this simple model, the parameter space favored by $c_\gamma$ and $c_g$ are separately shown in Fig. 4 right panel; two regions barely overlap. The best-fit solution of $c_g$ prefers to negative sign of yukawa coupling $x$ while $c_\gamma$ prefer to positive sign at 1σ level.

5.4 Non-standard decay modes and an upper bound : \{ $\alpha$, $c_\gamma$, $c_g$ $BR_{nonSM}$ \}

As briefly discussed in universal modification case, if both $\alpha$ and $BR_{nonSM}$ are involved, no unique solution of best fit is found. Let us elaborate this difficulty, how to handle this difficulty, and how we generically obtain upper bound on $BR_{nonSM}$ in the singlet extended models. We first consider following most general case based on our assumptions: \{ $\alpha$, $c_\gamma$, $c_g$ $BR_{nonSM}$ \}, with

$$\kappa_g^2 = (c_\alpha - c gs_\alpha)^2,$$  \( \kappa_\gamma^2 = (c_\alpha - c_\gamma s_\alpha)^2, \)

$$\kappa_{mix}^2 = \kappa_\alpha^2, \quad \kappa_H^2 = \frac{0.0857 \kappa_g^2 + 0.9143 \kappa_{mix}^2}{1 - BR_{nonSM}}. \quad (5.19)$$

Therefore, three parameters $\kappa_g$, $\kappa_\gamma$ and $\kappa_{mix}$ can parameterize all signal strengths. We do not obtain unique solution of best-fit because theory has four free parameters while data is
parameterized by three parameters. Alternatively, one can note that \( \hat{\kappa} \)‘s are invariant under the following transformation of theory parameters \( c \rightarrow c' \)

\[
\frac{c'^2}{c_H'} = \frac{c_\alpha^2}{c_H}, \quad c'_g = c_g s_\alpha c'_\alpha s'_\alpha, \quad c'_\gamma = c_\gamma s_\alpha c'_\alpha s'_\alpha \quad \Rightarrow \quad \hat{\kappa}' = \hat{\kappa} \quad (5.20)
\]

which explains redundancy in this four-parameter description of data.

We thus rather fit three hatted theory parameters

\[
\{ \hat{c}_g, \quad \hat{c}_\gamma, \quad \hat{c}_\alpha \} \quad (5.21)
\]

defined as

\[
\hat{\kappa}^2_{\text{mix}} = \frac{c_\alpha^2}{\kappa_H} \equiv \hat{c}_\alpha^2, \quad \hat{\kappa}_g^2 = \frac{(c_\alpha - c_g s_\alpha)^2}{\kappa_H} \equiv (\hat{c}_\alpha - \hat{c}_g \hat{s}_\alpha)^2, \quad \hat{\kappa}_\gamma^2 = \frac{(c_\alpha - c_\gamma s_\alpha)^2}{\kappa_H} \equiv (\hat{c}_\alpha - \hat{c}_\gamma \hat{s}_\alpha)^2.
\]

\[
(5.22)
\]
Hatted cosine and sine are defined to obey $\hat{s}_\alpha^2 + \hat{c}_\alpha^2 = 1$ with a reasonable assumption $\hat{c}_\alpha^2 \leq 1$. This assumption however does not affect the result of best-fit. Solving these in terms of un-hatted parameters,

$$c_\alpha^2 = \kappa_H \hat{c}_\alpha^2, \quad c_g = \hat{c}_g \frac{\hat{s}_\alpha c_\alpha}{\hat{c}_\alpha s_\alpha} = \hat{c}_g \frac{\hat{s}_\alpha}{\hat{c}_\alpha} \frac{\sqrt{\kappa_H \hat{c}_\alpha}}{\sqrt{1 - \kappa_H \hat{c}_\alpha}}, \quad c_\gamma = \hat{c}_\gamma \frac{\hat{s}_\alpha c_\alpha}{\hat{c}_\alpha s_\alpha}$$ (5.23)

which is nothing but the transformation eq.(5.20) between $c_i$ and $\hat{c}_i$ assuming that $\hat{c}_i$ are theory parameters $c_i$ with $\kappa_H = 1$.

Best-fit in terms of hatted parameters is

$$\hat{c}_g = -0.176^{+0.231}_{-0.219}, \quad \hat{c}_\gamma = -0.432^{+0.406}_{-0.374}, \quad \hat{c}_\alpha = 0.971^{0.0451}, \quad \chi^2/\nu = 11.1/7 = 1.58.$$ (5.24)

The previous case (3) with $\{\alpha, c_\gamma, c_g\}$ is one special point of four-parameter fit here. This can be seen from the trajectories of theory parameters in Fig. 6 giving equivalently good fits to data; $BR_{\text{nonSM}}$ vanishes at previous best-fit $c_\alpha = 0.947$ and other parameters there correspond to best-fit values of case (3).

If mixing is too large ($c_\alpha \lesssim 0.947$), the central value of best-fit non-SM width becomes negative. If one requires the positivity of non-SM width, this large mixing region would have poorer fit. On the other hand, if mixing is too small, then required $c_g, \gamma$ become very large. This might indicate either charged particles lighter than $W$ boson or coupling much stronger than electroweak coupling, thus may not be favored or calculable.

Now we see how Higgs total decay width or non-standard branching ratio can be upper bounded from the fact that $c_\alpha^2 \leq 1$. If $\kappa_H$ is too large, even maximally possible $c_\alpha^2 = 1$ cannot make $\hat{c}_\alpha^2$ large enough to satisfy the lower-bound of best-fit result. Using the relation $\kappa_H = c_\alpha^2/\hat{c}_\alpha^2$, this argument leads to the following upper bound on $\kappa_H$:

$$\kappa_H \leq \kappa_H^{\text{max}} \equiv \frac{c_{\alpha}^{\text{max}}}{c_{\alpha}^{2, \text{min}}} = \frac{1}{c_{\alpha}^{2, \text{min}}},$$ (5.25)
If we use $1\sigma$ or 95\%CL lower ranges of $\hat{c}_\alpha$ fit result, we obtain
\begin{equation}
\kappa_H^2 = \frac{\Gamma_{tot}}{\Gamma_{SM}} \leq \begin{cases} 
1.36 \text{ (1}\sigma) \\
1.65 \text{ (95\%CL).}
\end{cases}
\end{equation}
Likewise, we can solve the relation $\kappa_H = \frac{c^2_\alpha}{\hat{c}^2_\alpha}$ for $BR_{nonSM}$ using eq.(5.19), and obtain
\begin{equation}
BR_{nonSM} = 1 - \frac{c^4_\alpha}{\hat{c}^4_\alpha} \left(0.0857\kappa^2_\gamma + 0.9143\kappa^2_\alpha\right) \leq 1 - \hat{c}^4_{\alpha,\min} = \begin{cases} 
0.27 \text{ (1}\sigma) \\
0.39 \text{ (95\%CL).}
\end{cases}
\end{equation}
In obtaining these upper bounds, we have used built-in condition, $c_\alpha \leq 1$. If certain well-motivated theoretical assumptions on vector-boson couplings are made, one may obtain stronger bounds on non-standard decay width, e.g., see Refs.[32, 34].

Note that the singlet-like eigenstate can still have some non-standard decay modes if those modes are not be kinematically allowed for SM-like Higgs. These decays are not constrained from global fit of 125GeV data.

Let us consider a special case where $c_\gamma = 0$, namely that the singlet scalar has interactions with new charged particles that are colorless : $\{\alpha, c_\gamma, BR_{nonSM}\}$. Then one has
\begin{equation}
\kappa^2_\gamma = (c_\alpha - c_\gamma s_\alpha)^2, \quad \kappa^2_\alpha = \kappa^2_{mix} = c^2_\alpha, \quad \kappa^2_H = \frac{c^2_\alpha}{1 - BR_{nonSM}} \end{equation}
giving the best-fit
\begin{equation}
\hat{c}_\alpha = 0.990_{-0.0292}, \quad \hat{c}_\gamma = -0.786^{+0.714}_{-0.637}, \quad \chi^2/\nu = 11.1/8 = 1.39.
\end{equation}
Upper bound on non-standard branching ratio is also obtained in the same way as above
\begin{equation}
BR_{nonSM} \leq 1 - \hat{c}^4_{\alpha,\min} = \begin{cases} 
0.15 \text{ (1}\sigma) \\
0.24 \text{ (95\%CL).}
\end{cases}
\end{equation}

## 5.5 More dedicated models with multiple leptons or quarks
Models of multiple vector-like leptons can couple both to Higgs and singlet scalar directly. This type of model has been studied to simultaneously enhance diphoton rate and tame resulting vacuum instability via singlet scalar threshold effects [88]. $b_\gamma$ is then modified in addition to $c_\gamma$. Thus one can consider following case : $\{\alpha, c_\gamma, b_\gamma\}$. However, it is clear that no unique solution of $b_\gamma, c_\gamma$ will be obtained because those two free parameters only enter into a single $\kappa_\gamma$. In terms of underlying model parameters, $b_\gamma$ and $c_\gamma$ are induced by different independent couplings. Thus, underlying model is also not well constrained from global fit. We do not study this modes furthermore in this paper.

## 6 Implications on the other scalar boson nearby 125 GeV Higgs boson
Interesting application of our global fit in terms of the coupling constants of interaction eigenstates is to derive further constraints on the extra scalar boson mixing with 125GeV resonance. We use the null results of SM Higgs searches at LEP and the $\gamma\gamma$, $ZZ$ resonance searches in the other mass range at LHC to illustrate the application. The evolution of loop-induced couplings will also be discussed. Recently reported diphoton bump at 136.5GeV observed by CMS will be briefly discussed at the end.
Table 7. Summary of best-fit results with scalar mixing. If $BR_{nonSM}$ is included in fit, no unique solution is found, and its upper bound at 95%CL is presented. Only central values of best-fit are shown, and errors can be found in text.

| Models                        | Best-fit results | $\chi^2/\nu$ |
|-------------------------------|-----------------|--------------|
| SM                            |                 | 12.01/10 = 1.20 |
| universal modification        | $(\tilde{\kappa}_2^{univ})$ | 11.96/9 = 1.33 |
|                              | $(BR_{nonSM})$  | $\leq 18.8\%$ at 95\%CL |
|                              | $(\cos\alpha)$ | $\geq 0.904$ at 95\%CL |
| VL lepton, $W^\prime, S^\prime$ | $(c_{\alpha}, c_{\gamma})$ | 11.1/8 = 1.39 |
| VL quark                      | $(c_{\alpha}, c_{\rho}, c_{\gamma})$ | 11.1/7 = 1.58 |
|                              | $(c_{\alpha}, c_{\gamma}, Br_{nonSM})$ | $BR_{nonSM} \leq 24\%$ at 95\%CL |
|                              | $(c_{\alpha}, c_{\rho}, c_{\gamma}, Br_{nonSM})$ | $BR_{nonSM} \leq 39\%$ at 95\%CL |
| singlet mixed-in $\tilde{\kappa}$ | $(\tilde{\kappa}_g, \tilde{\kappa}_1, \tilde{\kappa}_2^{mix})$ | 11.1/7 = 1.58 |
| singlet mixed-in theory       | $(\tilde{c}_g, \tilde{c}_{\gamma}, \tilde{c}_{\alpha})$ | 11.1/7 = 1.58 |

6.1 Universal modification and LEP bounds

Universal modification scenario in terms of two parameters

$$\{ \alpha, BR_{nonSM} \}$$

is the simplest scenario that can be constrained from LEP searches of light Higgs boson. $BR_{nonSM}$ can be relevant if $m_s \leq m_h/2$. Our discussion in this section, however, does not depend on whether $m_s \leq m_h/2$ or not because only $\tilde{\kappa}_{univ}$ combination (not individual $\alpha$ or $BR_{nonSM}$) is constrained; see Sec.5.1 for discussion. LHC bounds will be discussed in the next subsection using another scenario although similar bounds can be derived in this case. As both production and decay relevant at LEP are proceeded by tree-level induced coupling. Thus, we do not need to discuss the mass dependences of couplings here – we refer to next subsection for this discussion. With these, the signal strengths of $H_2$ is universally modified as that of $H_1$.

LEP1 looked for Higgs via Bjorken process $Z \to Z^\ast h$ followed by $Z \to \ell^+\ell^-, \nu\bar{\nu}, b\bar{b}, \tau\tau$ and $h \to b\bar{b}, \tau\tau$. LEP2 looked for Higgs via $e^+e^- \to Zh$ [122]. Signal strengths in all cases...
are globally parameterized by
\[ \frac{\sigma}{\sigma_{SM}} = \hat{\kappa}_{\text{univ}}^4 = c_\alpha^2 (1 - BR_{\text{nonSM}}). \] (6.2)

The LEP bounds on the signal strength can be interpreted as bounds on the parameter \( \hat{\kappa}_{\text{univ}} \). The 1,2\( \sigma \) favored region of global fit is overlayed with LEP bounds in Fig. 7 left panel. Since SM is consistent with best-fit results, no strong bound can be derived yet. However, best-fit results can already constrain heavy mass regime better than LEP.

### 6.2 Scalars enhancing the diphoton rate with \( c_\gamma \)

This subsection discusses a useful application of our global fits and Lagranigan parameterization introduced in this paper. We will evolve the coupling constants \( b_i \) and \( c_i \) fitted at 125GeV to the mass of the singlet \( m_{H_2} \neq 125 \text{GeV} \) in order to extract physical couplings of \( H_2 \) bosons, and discuss if these couplings can be further constrained from other \( \gamma\gamma \) [123], \( ZZ \) [114] resonance searches at collider. To this end, we will especially have to discuss how mass dependences of loop-induced couplings can be evolved properly. All other couplings are assumed to be tree-level generated and to be constant over the mass range. We explicitly write the mass dependences of loop-induced couplings: \( c_{g,\gamma}(m) \) and \( b_{g,\gamma}(m) \). Recall that these parameters are defined and determined at \( m = m_{H_1} = 125 \text{GeV} \) by global fits.

Focusing on \( m_{H_2} \lesssim 160 \text{GeV} \) allows us to meaningfully extract loop-induced couplings of \( H_2 \). This is technically useful because the mass dependences of \( c_{g,\gamma} \) and \( b_g \) coming from the loops of heavy particles – top quarks and extra new vector-like leptons or quarks – are small. See Fig. 8 that top quark’s loop function does not change much in this mass range; thus, so do heavier particles’ loop functions. This is very useful simplification as \( c_g \) and \( b_g \) enter via \( gg \) fusion production of \( H_2 \), and thus \( gg \) fusion is independent on mass scale – it is simply rescaled by mixing angle regardless of \( m_{H_2} \). New particles, of course, could be lighter than \( H_2 \) or \( 2m_W \), but we do not consider such a case given that collider bounds
Figure 8. (Left): Loop functions of $W$ boson, $A_1(W)$, and top quark, $A_{1/2}(t)$, are evaluated at mass scale $m_{H_2}$. SM Higgs coupling to photons corresponds to black line while coupling to gluon is proportional to $A_{1/2}(t)$. (Right): $b_\gamma(m_{H_2})$ relevant to the coupling of $H_2$ boson to photons as defined in eq.(6.7).

Generally predict heavier charged (and/or colored) particles. On the other hand, $b_\gamma(m_{H_2})$ scales mildly with $m_{H_2}$ in this mass range due to the contribution from lighter $W$ boson as shown in Fig. 8. The $W$ loop function could have behaved more rapidly and obtained imaginary part if $H_2$ were heavier than about $2 m_W$. Also, $H_2$ could have dominantly decayed to $WW^*$ making diphoton modes irrelevant. In all, we assume that $c_{g,\gamma}, b_g$ are mass-independent, whereas $b_\gamma$ is properly evolved from 125GeV to $m_{H_2}$ as will be explained in eq.(6.7). For simplicity, we assume that $H_2$ does not have extra non-standard decay modes.

We consider a two-parameter fit

$$\{ \alpha, c_\gamma \}. \quad (6.3)$$

Decay widths of 125GeV Higgs boson is denoted by $\kappa_\gamma^2$ as usual

$$\kappa_\gamma^2 = \frac{\Gamma(H_1 \rightarrow \gamma\gamma)}{\Gamma(h \rightarrow \gamma\gamma)_{SM}|_{125\text{GeV}}} = (b_\gamma(125)c_\alpha - c_\gamma(125)s_\alpha)^2 = (c_\alpha - c_\gamma s_\alpha)^2. \quad (6.4)$$

where $b_\gamma(125) = 1$ and $c_\gamma(m) = c_\gamma$ constant as discussed. The $H_2$ decay width with respect to SM width at $m_{H_2}$ is written with some care by

$$\zeta_\gamma^2 = \frac{\Gamma(H_2 \rightarrow \gamma\gamma)}{\Gamma(h \rightarrow \gamma\gamma)_{SM}|_{m_{H_2}}} = \frac{\Gamma(H_2 \rightarrow \gamma\gamma)|_{m_{H_2}}}{\Gamma(h \rightarrow \gamma\gamma)_{SM}|_{125}} \cdot \frac{\Gamma(h \rightarrow \gamma\gamma)_{SM}|_{125}}{\Gamma(h \rightarrow \gamma\gamma)_{SM}|_{m_{H_2}}} = (c_\gamma(m_{H_2})c_\alpha + b_\gamma(m_{H_2})s_\alpha)^2 \cdot \left( \frac{b_\gamma(125)}{b_\gamma(m_{H_2})} \right)^2 \quad (6.5)$$

$$= \left( \frac{c_\gamma(m_{H_2})}{b_\gamma(m_{H_2})} c_\alpha + s_\alpha \right)^2 \quad (6.6)$$

where in the second line, we used the fact that $c_i$ and $b_i$ are defined with respect to corresponding 125GeV SM Higgs couplings. $c_\gamma(m_{H_2}) = c_\gamma(m_{H_1}) = c_\gamma$ constant. $b_\gamma(m_{H_2})$ is
given by

\[ b_\gamma(m_{H_2})^2 = \left| \frac{1}{A_{SM}} (A_1(m_{H_2}/4m_W^2) + N_c Q_i^2 A_{1/2}(m_{H_2}/4m_t^2)) \right|^2 \]

\[ = \frac{A_1(m_{H_2}/4m_W^2) + N_c Q_i^2 A_{1/2}(m_{H_2}/4m_t^2)}{A_1(m_{H_1}/4m_W^2) + N_c Q_i^2 A_{1/2}(m_{H_1}/4m_t^2)}^2 \]  

(6.7)

and is plotted in the right panel of Fig. 8. Decay widths of other decay modes are simply rescaled by mixing angle

\[ \zeta_i^2 = \frac{\Gamma(H_2 \to i)}{\Gamma(h \to i)_{SM}} = s_\alpha^2. \]  

(6.8)

Total decay width of \( H_2 \) is then also approximately

\[ \zeta_H^2 = \frac{\Gamma(H_2)_{tot}}{\Gamma(h)_{tot}}_{SM} \sim s_\alpha^2. \]  

(6.9)

Similarly to previous case, LEP bound overlapped with favored region of global fit results in eq.(5.13) is shown in the right panel of Fig. 7. If future data would predict the same central value with smaller error, we might obtain a band of favored region which can narrow down the allowed mass range. Diphoton resonance search [123] and \( ZZ^* \to 4\ell \) resonance search at LHC [114] constrain following signal strength parametrizations

\[ R(\sigma(pp \to H_2) \times BR(H_2 \to \gamma\gamma)) = \zeta_\gamma^2 \]  

(6.10)

\[ R(\sigma(pp \to H_2) \times BR(H_2 \to ZZ^*)) = \zeta_Z^2 \]  

(6.11)

where \( \zeta_i^2 = \zeta_i^2/\zeta_H \) as before. Bounds on the parameter space from these resonance searches and global fits are shown in Fig. 9. As search bounds vary with the mass of a resonance, we pick four representative results in the plot. Contours of constant \( \zeta_i^2 \) are also shown for reference. SM is consistent with both global fit and resonance searches. Nonetheless, some region preferred by global fits is excluded by resonance searches.

### 6.3 Addendum with 136.5GeV diphoton accumulation

We also carry out a quick and interesting aside. CMS has recently observed the slight excess of diphoton resonance data at around 136.5GeV [124]. ATLAS has not so far observed corresponding excess. No detailed CMS data is available yet. But we assume following data \( \mu_{\gamma\gamma}(136.5 \text{ GeV}) = 2.3 \pm 0.45 \) and illustrate how our method can be applied to interpret this data.

We use \( \{c_\alpha, c_\gamma\} \) ansatz to fit both 125GeV and 136.5GeV CMS data simultaneously. ATLAS data is not included. From previous subsection, 136.5GeV diphoton signal strength (compared to that of 136.5GeV SM Higgs rate) is given by eq.(6.6) and eq.(6.11). Here, \( b_\gamma(136.5 \text{ GeV}) = 1.069 \) and \( c_\gamma(136.5 \text{ GeV}) = c_\gamma(125 \text{ GeV}) = c_\gamma \) is constant. Using these, our best fit yields

\[ c_\gamma = 1.54, \quad c_\alpha = 0.997, \quad \chi^2/\nu = 1.65/5 = 0.33 \]  

(6.12)
Figure 9. In addition to 1, 2, 3σ favored region of global fit of \( \{\alpha, c_\gamma\} \) as black lines, contours of \( \zeta_2^2 = R(\sigma(pp \rightarrow H_2) \times BR(H_2 \rightarrow \gamma\gamma)) \) are shown as red-dashed. From two central lines, \( \zeta_2^2 = 0.1, 0.5, 1 \). Regions between two thick red lines are allowed from 95%CL bound of diphoton resonance search at LHC [123]. Shaded region in the last plot is excluded from \( ZZ^* \rightarrow 4\ell \) resonance search at LHC by 95%CL [114]. This bound is too weak to be shown in other panels.

which can be compared with best-fit results without 136.5GeV data: \( c_\gamma = 0.173, c_\alpha = 0.942, \chi^2/\nu = 1.13/4 = 0.28 \) (only CMS) and \( c_\alpha = 0.98, c_\gamma = -0.55, \chi^2/\nu = 11.1/8 = 1.39 \) (both CMS and ATLAS as reported in eq.(5.13)). Notably, sizable \( c_\gamma \) is preferred compared to CMS-alone results while positive \( c_\gamma \) is preferred as opposed to CMS+ATLAS results. The latter may imply that ATLAS 125GeV diphoton enhancement and CMS 136.5GeV diphoton enhancement may not consistently coexist in the model of singlet extension (see Sec.4.3 for discussions on general trends of ATLAS and CMS datasets). We comment that we restrict \( s_\alpha > 0 \) throughout this paper which doesn’t matter if only 125GeV data is used in the global fit, but changing the sign of \( s_\alpha \) yields somewhat worse fit here.
7 Conclusions

In this paper, we performed comprehensive analysis of the current LHC data on the 125 GeV Higgs-like resonance assuming that the SM Higgs properties can be modified either by new physics in the loop or by a mixing with a nearby singlet scalar boson. Our approach is conceptually more general than other approaches based on the effective Lagrangian with higher dimensional operators with the SM Higgs boson $h$. We imposed the SM gauge symmetry $SU(3)_c \times SU(2)_L \times U(1)_Y$ instead of $SU(3)_c \times U(1)_{em}$, and distinguish the nature of the SM Higgs $h$ and a singlet scalar $s$. In practice, our fitting procedure is identical to others, but the interpretation of the results could be different. We could obtain more detailed informations on possible new physics impacts on the SM Higgs properties within our approach.

We presented a number of interesting class of BSM’s where new singlet scalar bosons appear with couplings to the SM fields as well as to some new fields such as new charged vector bosons, vectorlike fermions or hidden sector dark matters, etc.. The singlet scalar boson(s) mix with the SM Higgs boson, and thus modify the Higgs properties.

To parametrize these general modification due to a single singlet scalar, we introduce separate coupling constants to the Higgs and singlet interaction eigenstates. Their mixing is parameterized by a single mixing angle. We later have also discussed how this parametrization can be further used to constrain the singlet scalar interactions from other resonance searches at collider.

We carried out global fits assuming various ansatzs representing BSM’s in literature. Interestingly enough, the LHC data after Moriond 2013, according to our global fit, implies that the SM gives the best $\chi^2/\nu = 1.20$ when both ATLAS and CMS data are used for the fit. All other models we considered, such as models with extra vector-like fermions or charged vector bosons that were considered for the enhancement of $H \rightarrow \gamma\gamma$ or new sequential fermions or 2 HDM’s, or models with extra singlet scalar bosons, yield larger $\chi^2/\nu$ compared with the SM (see Tables 5 and 7). If we consider the ATLAS and the CMS data separately and perform the $\chi^2$ minimization fit, the CMS data seems to prefer a singlet scalar boson with some suppression of the signal strengths. If the SM is assumed, $\chi^2/\nu = 0.466$ is obtained for the CMS-only which is smaller than ATLAS + CMS SM fit ($\chi^2/\nu = 1.20$). On the other hand, the ATLAS data can be better fitted with BSM in terms of $b_i \neq 1$, but if the SM is assumed, ATLAS’s $\chi^2/\nu = 1.94$ becomes worse than the ATLAS + CMS SM fit. Of course, the number of degrees of freedom in our fits maybe too small to make a strong conclusion at the moment, but this observation may be suggestive.

In order to make a stronger conclusion, we have to await more data accumulation from the next run of the LHC at 14 TeV, with better statistics and systematics, as well as better determination of the signal strengths for different Higgs production channels. In many BSM’s we presented in this paper, there appears a second neutral Higgs-like scalar boson, a mixture of the SM Higgs boson $h$ and a singlet scalar $s$. It would not be easy to search for this second scalar, because it should be mostly a singlet scalar, considering the current data on the 125 GeV Higgs-like resonance. Still, search for this second scalar is clearly warranted because it may be the only way to probe dark sector made of the SM singlet.
fields at colliders.

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