

Contributions of inert electroweak multiplets to Higgs properties

Hugues Beauchesne^a and Cheng-Wei Chiang^{b,a}

^a*Physics Division, National Center for Theoretical Sciences,
Taipei 10617, Taiwan*

^b*Department of Physics and Center for Theoretical Physics, National Taiwan University,
Taipei 10617, Taiwan*

E-mail: beauchesneh@phys.ncts.ntu.edu.tw, chengwei@phys.ntu.edu.tw

ABSTRACT: New physics could manifest itself in the form of electroweak multiplets that interact at tree level with the Higgs boson but do not mix with Standard Model fields or acquire expectation values. In this paper, we study the potential contributions of such inert multiplets to several crucial Higgs properties, namely, the branching ratio of the Higgs to a Z boson and a photon (or massless dark photon) and the triple Higgs coupling. Constraints from the Higgs signal strengths, oblique parameters and unitarity are taken into account.

KEYWORDS: Anomalous Higgs Couplings, Higgs Properties, New Gauge Interactions, Specific BSM Phenomenology

ARXIV EPRINT: [2308.06901](https://arxiv.org/abs/2308.06901)

Contents

1	Introduction	1
2	Possible interaction terms	2
3	Constraints and observables	6
3.1	Lagrangian	6
3.2	Relevant Higgs decays and Higgs signal strengths	8
3.2.1	Higgs to (dark) photons	8
3.2.2	Higgs decay to Z and a (dark) photon	9
3.2.3	Constraints on the Higgs signal strengths	10
3.3	Oblique parameters	10
3.4	Unitarity	11
3.5	Triple Higgs coupling	13
4	Results	14
4.1	Scanning procedure	14
4.2	One interaction term	15
4.3	Multiple interaction terms	20
5	Conclusion	21

1 Introduction

Electroweak multiplets that neither mix with Standard Model (SM) fields nor acquire vacuum expectation values (VEVs) are ubiquitous in physics beyond the Standard Model (BSM). Such multiplets, which we will refer to as inert, can provide potential dark matter candidates [1–4], be responsible for baryogenesis [5, 6], or explain certain flavour anomalies [7].

Inert multiplets can in principle interact at tree level with the Higgs boson h and therefore affect some of its properties. Interestingly, specifying these interaction terms and the masses of the particles involved is sufficient to compute their leading contributions to many Higgs properties. Since these interaction terms can only take a limited number of forms, this allows for a very broad analysis of the impact of inert multiplets on specific Higgs properties that are sensitive to new physics contributions. In this paper, we will study three such properties, all affected by the new particles at the one-loop level.

The first one is the branching ratio of the Higgs boson to a photon and a Z boson. Since this is a loop process in the SM, new particles could potentially lead to sizable deviations from the SM prediction. Furthermore, the current measurement of this branching ratio

is 2.1 ± 0.7 times larger than its SM value, corresponding to a deviation of 1.9σ [8, 9]. Previous works on possible deviations of this branching ratio include refs. [10, 11].

The second Higgs property of interest is the branching ratio of the Higgs to a Z boson and a massless dark photon A' [12]. An extensive literature on the decay $h \rightarrow AA'$ already exists [13–21], but it was recently demonstrated in refs. [19–21] that this branching ratio is considerably more constrained than the previous analysis ref. [13]. Some of the constraints on $h \rightarrow AA'$ simply do not apply to $h \rightarrow ZA'$ and no bound on this branching ratio currently exists in the literature.

The third Higgs property considered is the triple Higgs coupling. This coupling provides additional information on the form of the Higgs potential. In addition, a larger value of this coupling could lead to a sufficiently strong first-order electroweak phase transition in the early Universe and thus explain matter abundance via electroweak baryogenesis [22, 23].

More precisely, the goal of this paper is to study the potential contributions of inert multiplets to the Higgs branching ratios $\text{BR}(h \rightarrow ZA^{(\prime)})$ and the triple Higgs coupling. To do so, we will consider all possible tree-level interactions between the Higgs and inert multiplets and then compute their contributions to these Higgs properties. Constraints from the Higgs signal strengths, the electroweak precision tests and perturbative unitarity will be taken into account. In many ways, this paper is an extension of the formalism of refs. [19, 20] and its application to new Higgs properties.

We find the following results. The branching ratio $\text{BR}(h \rightarrow ZA)$ can easily be enhanced by $\mathcal{O}(20\%)$ for simple models. Considerably larger enhancements are possible, but require more complicated models and careful fine-tuning. The branching ratio $\text{BR}(h \rightarrow ZA')$ could in principle be above 1%. It would however require new neutral particles barely above half the mass of the Higgs boson. A branching ratio of $\mathcal{O}(0.1\%)$ is however relatively easy to obtain. The triple Higgs coupling can easily be enhanced by a factor of several.

This paper is organized as follows. We introduce in section 2 the inert multiplets and all their relevant interactions with the Higgs doublet. Section 3 explains how we apply constraints and compute the above-mentioned properties. The results are shown in section 4. Some concluding remarks are presented in section 5.

2 Possible interaction terms

If inert multiplets are to have any sizable impact on Higgs properties, they need to interact with the Higgs boson at tree level. In this section, we present all renormalizable interaction terms relevant to our observables. Most of these terms were already included in refs. [19, 20], from which we borrow heavily. Contrary to these papers, we will not in general require the inert multiplets to be charged under a new $U(1)'$ symmetry. This will allow for a few additional terms as well as real multiplets. We will require the interaction terms to be able to contribute to our observables at one loop. In combination with the requirement of the fields not acquiring VEVs, this will force the interaction terms to always contain exactly two (distinct or not) inert multiplets. We will write quantum numbers of fields as $(SU(2)_L, U(1)_Y, U(1)')$. The $SU(2)_L$ indices are labeled by Latin letters and range from 1 to the size of the representation of the corresponding field. Considering the constraints on

new coloured particles, all fields will be assumed to be $SU(3)_C$ singlets. Terms that are equivalent to those below up to a field redefinition are not included in the list, but could easily be included when the field redefinition cannot be performed simultaneously for all terms in the Lagrangian. See ref. [24] for a similar exercise but with different assumptions and motivations.

Fermion case. Consider the fermions

$$\psi_1 : (\mathbf{p}, Y^p, Q'), \quad \psi_2 : (\mathbf{n}, Y^n, Q'), \quad (2.1)$$

with $Y^p = Y^n + 1/2$ and $p = n \pm 1$. Throughout the paper, we use the boldfaced \mathbf{p} and \mathbf{n} to denote the $SU(2)_L$ representation of a multiplet and italic p and n for their corresponding dimensions. The following term is allowed¹

$$F(\psi_1, \psi_2) = -\hat{d}_{abc}^{pn} \bar{\psi}_1^a (A_L P_L + A_R P_R) \psi_2^b H^c + \text{H.c.} \quad (2.2)$$

The $SU(2)_L$ tensor is given by the Clebsch-Gordan coefficient

$$\hat{d}_{abc}^{pn} = C_{j_1 m_1 j_2 m_2}^{JM} \equiv \langle j_1 j_2 m_1 m_2 | JM \rangle, \quad (2.3)$$

where

$$\begin{aligned} j_1 &= \frac{n-1}{2}, & j_2 &= \frac{1}{2}, & J &= \frac{p-1}{2}, \\ m_1 &= \frac{n+1-2b}{2}, & m_2 &= \frac{3-2c}{2}, & M &= \frac{p+1-2a}{2}. \end{aligned} \quad (2.4)$$

Depending on their quantum numbers, either ψ_1 or ψ_2 can be a real multiplet.

Scalar case I. Consider the scalars

$$\phi_1 : (\mathbf{p}, Y^p, Q'), \quad \phi_2 : (\mathbf{n}, Y^n, Q'), \quad (2.5)$$

with $Y^p = Y^n + 1/2$ and $p = n \pm 1$. The following term is allowed

$$S_1(\phi_1, \phi_2) = -\mu \hat{d}_{abc}^{pn} \phi_1^{a\dagger} \phi_2^b H^c + \text{H.c.} \quad (2.6)$$

The $SU(2)_L$ tensor is given by

$$\hat{d}_{abc}^{pn} = C_{j_1 m_1 j_2 m_2}^{JM}, \quad (2.7)$$

where

$$\begin{aligned} j_1 &= \frac{n-1}{2}, & j_2 &= \frac{1}{2}, & J &= \frac{p-1}{2}, \\ m_1 &= \frac{n+1-2b}{2}, & m_2 &= \frac{3-2c}{2}, & M &= \frac{p+1-2a}{2}. \end{aligned} \quad (2.8)$$

Depending on their quantum numbers, either ϕ_1 or ϕ_2 can be a real multiplet. This case includes the electroweak (EW) case of refs. [13–17] when their neutral scalar S is replaced by its expectation value.

¹Summation over repeated symbols is implicit when trivial but will be written explicitly otherwise.

Scalar case II. Consider the scalar

$$\phi : (\mathbf{n}, Y^n, Q'). \quad (2.9)$$

The following term is allowed

$$S_2(\phi) = - \sum_{r \in \{n-1, n+1\}} \lambda^r \hat{d}_{abcd}^{nr} H^{a\dagger} H^b \phi^{c\dagger} \phi^d. \quad (2.10)$$

The $SU(2)_L$ tensor is given by

$$\hat{d}_{abcd}^{nr} = \sum_M C_{j_1 m_1 j_2 m_2}^{JM} C_{j_3 m_3 j_4 m_4}^{JM}, \quad (2.11)$$

where M is summed over $\{-J, -J+1, -J+2, \dots, +J\}$ and

$$\begin{aligned} j_1 &= \frac{1}{2}, & j_2 &= \frac{n-1}{2}, & j_3 &= \frac{1}{2}, & j_4 &= \frac{n-1}{2}, & J &= \frac{r-1}{2}, \\ m_1 &= \frac{3-2a}{2}, & m_2 &= \frac{n+1-2c}{2}, & m_3 &= \frac{3-2b}{2}, & m_4 &= \frac{n+1-2d}{2}. \end{aligned} \quad (2.12)$$

Unless ϕ is a singlet, there are in general two possible contractions of the $SU(2)_L$ indices and therefore two coefficients. Depending on its quantum numbers, ϕ can be either a real or complex multiplet.

Scalar case III. Consider the scalars

$$\phi_1 : (\mathbf{p}, Y^p, Q'), \quad \phi_2 : (\mathbf{n}, Y^n, Q'), \quad (2.13)$$

with $p \in \{n-2, n, n+2\}$ and $Y^p = Y^n$. The following term is allowed

$$S_3(\phi_1, \phi_2) = - \sum_{r \in \mathcal{R}} \lambda^r \hat{d}_{abcd}^{pnr} H^{a\dagger} H^b \phi_1^{c\dagger} \phi_2^d + \text{H.c.}, \quad (2.14)$$

where $\mathcal{R} = \{n-1, n+1\} \cap \{p-1, p+1\}$. The $SU(2)_L$ tensor is given by

$$\hat{d}_{abcd}^{pnr} = \sum_M C_{j_1 m_1 j_2 m_2}^{JM} C_{j_3 m_3 j_4 m_4}^{JM}, \quad (2.15)$$

where M is summed over $\{-J, -J+1, -J+2, \dots, +J\}$ and

$$\begin{aligned} j_1 &= \frac{1}{2}, & j_2 &= \frac{p-1}{2}, & j_3 &= \frac{1}{2}, & j_4 &= \frac{n-1}{2}, & J &= \frac{r-1}{2}, \\ m_1 &= \frac{3-2a}{2}, & m_2 &= \frac{p+1-2c}{2}, & m_3 &= \frac{3-2b}{2}, & m_4 &= \frac{n+1-2d}{2}. \end{aligned} \quad (2.16)$$

If p and n differ by two, there is only one possible contraction. If $p = n$, there are two possible contractions unless $p = n = 1$. Depending on their quantum numbers, either ϕ_1 and ϕ_2 can be a real multiplet.

Scalar case IV. Consider the scalars

$$\phi_1 : (\mathbf{p}, Y^p, Q'), \quad \phi_2 : (\mathbf{n}, Y^n, Q'), \quad (2.17)$$

with $p \in \{n-2, n, n+2\}$ and $Y^p = Y^n + 1$. The following term is allowed

$$S_4(\phi_1, \phi_2) = -\lambda \hat{d}_{abcd}^{pn} H^a H^b \phi_1^{c\dagger} \phi_2^d + \text{H.c.} \quad (2.18)$$

The $SU(2)_L$ tensor is given by

$$\hat{d}_{abcd}^{pn} = \sum_{M_1} C_{j_1 m_1 j_2 m_2}^{J_1 M_1} C_{J_1 M_1 j_3 m_3}^{J_2 M_2}, \quad (2.19)$$

where M_1 is summed over $\{-1, 0, 1\}$ and

$$\begin{aligned} j_1 &= \frac{1}{2}, & j_2 &= \frac{1}{2}, & j_3 &= \frac{n-1}{2}, & J_1 &= 1, & J_2 &= \frac{p-1}{2}, \\ m_1 &= \frac{3-2a}{2}, & m_2 &= \frac{3-2b}{2}, & m_3 &= \frac{n+1-2d}{2}, & M_2 &= \frac{p+1-2c}{2}. \end{aligned} \quad (2.20)$$

Only one contraction of the $SU(2)_L$ indices is allowed. Depending on their quantum numbers, either ϕ_1 or ϕ_2 can be a real multiplet.

Scalar case V. Consider the scalar

$$\phi : (\mathbf{n}, 1/2, 0). \quad (2.21)$$

The following term is allowed

$$S_5(\phi) = -\lambda \hat{d}_{abcd}^n H^a H^b \phi^{c\dagger} \phi^{d\dagger} + \text{H.c.} \quad (2.22)$$

The $SU(2)_L$ tensor is given by

$$\hat{d}_{abcd}^n = \sum_M C_{j_1 m_1 j_2 m_2}^{JM} C_{J_3 m_3 j_4 m_4}^{JM}, \quad (2.23)$$

where M is summed over $\{-1, 0, 1\}$ and

$$\begin{aligned} j_1 &= \frac{1}{2}, & j_2 &= \frac{1}{2}, & j_3 &= \frac{n-1}{2}, & j_4 &= \frac{n-1}{2}, & J &= 1, \\ m_1 &= \frac{3-2a}{2}, & m_2 &= \frac{3-2b}{2}, & m_3 &= \frac{n+1-2c}{2}, & m_4 &= \frac{n+1-2d}{2}. \end{aligned} \quad (2.24)$$

Only one non-zero contraction of the $SU(2)_L$ indices is possible. The scalar ϕ must be a complex multiplet. It must also be of even dimension to be non-zero, which can easily be verified from a table of Clebsch-Gordan coefficients. This interaction term does not allow ϕ to be charged under $U(1)'$ and thus did not appear in refs. [19, 20].

Scalar case VI. Consider the scalar

$$\phi : (\mathbf{n}, 0, 0). \tag{2.25}$$

The following term is allowed

$$S_6(\phi) = - \sum_{r \in \{1,3\}} \lambda^r \hat{d}_{abcd}^{nr} H^{a\dagger} H^b \phi^c \phi^d + \text{H.c.} \tag{2.26}$$

The $SU(2)_L$ tensor is given by

$$\hat{d}_{abcd}^{nr} = \sum_M C_{j_1 m_1 j_2 m_2}^{J_1 M} C_{J_1 M j_3 m_3}^{J_2 M_2}, \tag{2.27}$$

where M is summed over $\{-J_1, -J_1 + 1, -J_1 + 2, \dots, +J_1\}$ and

$$\begin{aligned} j_1 &= \frac{n-1}{2}, & j_2 &= \frac{n-1}{2}, & j_3 &= \frac{1}{2}, & J_1 &= \frac{r-1}{2}, & J_2 &= \frac{1}{2}, \\ m_1 &= \frac{n+1-2c}{2}, & m_2 &= \frac{n+1-2d}{2}, & m_3 &= \frac{3-2b}{2}, & M_2 &= \frac{3-2a}{2}. \end{aligned} \tag{2.28}$$

Two contractions of the $SU(2)_L$ indices are possible, though in practice only $r = 1$ contributes if n is odd and only $r = 3$ if n is even. The scalar ϕ can be either real or complex. This term was not permitted by the requirements of refs. [19, 20].

3 Constraints and observables

In this section, we introduce useful notation, explain how constraints are applied and present the computations of relevant Higgs properties. Many results from this section can be found in refs. [19, 20], though we modified the notation and expanded them to apply to additional scenarios.

3.1 Lagrangian

The only parts of the Lagrangian necessary to compute our observables are the kinematic terms and the interaction terms of the last section. In this section, we will write down the most general form of the relevant Lagrangian terms that can result from them. All new particles considered will be either complex scalars ϕ_i^C , real scalars ϕ_i^R , Dirac fermions ψ_i^D or Majorana fermions ψ_i^M , the only exception being a potential dark photon A' . The Lagrangian is expressed in terms of gauge eigenstates. Many parameters will be introduced for convenience, though some are related either by construction or gauge symmetries.

Mass terms

$$\begin{aligned} \mathcal{L}^{\text{Mass}} &= - (m_C^2)_{ij} \phi_i^{C\dagger} \phi_j^C - \frac{1}{2} (m_R^2)_{ij} \phi_i^R \phi_j^R \\ &\quad - \bar{\psi}_i^D \left((m_D^L)_{ij} P_L + (m_D^R)_{ij} P_R \right) \psi_j^D - \frac{1}{2} \bar{\psi}_i^M \left((m_M^L)_{ij} P_L + (m_M^R)_{ij} P_R \right) \psi_j^M. \end{aligned} \tag{3.1}$$

Yukawa interactions

$$\begin{aligned} \mathcal{L}^{\text{Yukawa}} = & -h(\Omega_C)_{ij}\phi_i^{C\dagger}\phi_j^C - \frac{h}{2}(\Omega_R)_{ij}\phi_i^R\phi_j^R \\ & -h\bar{\psi}_i^D\left((\Omega_D^L)_{ij}P_L + (\Omega_D^R)_{ij}P_R\right)\psi_j^D - \frac{h}{2}\bar{\psi}_i^M\left((\Omega_M^L)_{ij}P_L + (\Omega_M^R)_{ij}P_R\right)\psi_j^M. \end{aligned} \quad (3.2)$$

Gauge interactions of complex scalars

$$\begin{aligned} \mathcal{L}_C^{\text{Gauge}} = & \left(-i(A_C)_{ii}A_\mu\phi_i^{C\dagger}\partial^\mu\phi_i^C + \text{H.c.}\right) + (B_C)_{ii}A_\mu A^\mu\phi_i^{C\dagger}\phi_i^C \\ & + \left(-i(C_C)_{ij}Z_\mu\phi_i^{C\dagger}\partial^\mu\phi_j^C + \text{H.c.}\right) + (D_C)_{ij}Z_\mu Z^\mu\phi_i^{C\dagger}\phi_j^C + 2(E_C)_{ij}A_\mu Z^\mu\phi_i^{C\dagger}\phi_j^C \\ & + \left(-i(F_C)_{ij}W_\mu^+\left(\phi_i^{C\dagger}\partial^\mu\phi_j^C - \partial^\mu\phi_i^{C\dagger}\phi_j^C\right) + \text{H.c.}\right) + 2(G_C)_{ij}W_\mu^+W^{-\mu}\phi_i^{C\dagger}\phi_j^C \\ & + \left(-i\frac{(H_C^1)_{ij}}{\sqrt{2}}W_\mu^+\left(\phi_i^{C\dagger}\partial^\mu\phi_j^{C\dagger} - \phi_i^{C\dagger}\partial^\mu\phi_j^{C\dagger}\right) + \text{H.c.}\right) \\ & + \left(-i\frac{(H_C^2)_{ij}}{\sqrt{2}}W_\mu^+\left(\phi_i^C\partial^\mu\phi_j^C - \phi_i^C\partial^\mu\phi_j^C\right) + \text{H.c.}\right) \\ & + \left(-i(A'_C)_{ii}A'_\mu\phi_i^{C\dagger}\partial^\mu\phi_i^C + \text{H.c.}\right) + 2(B'_C)_{ii}A'_\mu A'^\mu\phi_i^{C\dagger}\phi_i^C + (B''_C)_{ii}A'_\mu A'^\mu\phi_i^{C\dagger}\phi_i^C \\ & + 2(E'_C)_{ij}A'_\mu Z^\mu\phi_i^{C\dagger}\phi_j^C. \end{aligned} \quad (3.3)$$

Gauge interactions of real scalars

$$\mathcal{L}_R^{\text{Gauge}} = \frac{(C_R)_{ij}}{2}Z_\mu\left(\phi_i^R\partial^\mu\phi_j^R - \partial^\mu\phi_i^R\phi_j^R\right) + \frac{(D_R)_{ij}}{2}Z_\mu Z^\mu\phi_i^R\phi_j^R + (G_R)_{ij}W_\mu^+W^{-\mu}\phi_i^R\phi_j^R. \quad (3.4)$$

Gauge interactions of both complex and real scalars

$$\mathcal{L}_{RC}^{\text{Gauge}} = -i(F_{RC}^1)_{ij}W_\mu^+\left(\phi_i^R\partial^\mu\phi_j^C - \partial^\mu\phi_i^R\phi_j^C\right) - i(F_{RC}^2)_{ij}W_\mu^+\left(\phi_i^{C\dagger}\partial^\mu\phi_j^R - \partial^\mu\phi_i^{C\dagger}\phi_j^R\right) + \text{H.c.} \quad (3.5)$$

Gauge interactions of Dirac fermions

$$\begin{aligned} \mathcal{L}_D^{\text{Gauge}} = & -(A_D)_{ii}A_\mu\bar{\psi}_i^D\gamma^\mu\psi_i^D - Z_\mu\bar{\psi}_i^D\gamma^\mu\left((B_D^L)_{ij}P_L + (B_D^R)_{ij}P_R\right)\psi_j^D \\ & - W_\mu^+\bar{\psi}_i^D\gamma^\mu\left((F_D^L)_{ij}P_L + (F_D^R)_{ij}P_R\right)\psi_j^D + \text{H.c.} \\ & - (A'_D)_{ii}A'_\mu\bar{\psi}_i^D\gamma^\mu\psi_i^D. \end{aligned} \quad (3.6)$$

Gauge interactions of Majorana fermions

$$\mathcal{L}_M^{\text{Gauge}} = -\frac{Z_\mu}{2}\bar{\psi}_i^M\gamma^\mu\left((F_M^L)_{ij}P_L + (F_M^R)_{ij}P_R\right)\psi_j^M. \quad (3.7)$$

Gauge interactions of both Dirac and Majorana fermions

$$\begin{aligned} \mathcal{L}_{MD}^{\text{Gauge}} = & -W_\mu^+\bar{\psi}_i^M\gamma^\mu\left((F_{MD}^{1L})_{ij}P_L + (F_{MD}^{1R})_{ij}P_R\right)\psi_j^D + \text{H.c.} \\ & - W_\mu^+\bar{\psi}_i^D\gamma^\mu\left((F_{MD}^{2L})_{ij}P_L + (F_{MD}^{2R})_{ij}P_R\right)\psi_j^M + \text{H.c.} \end{aligned} \quad (3.8)$$

Quartic couplings

$$\mathcal{L}^{\text{Quartic}} = -\frac{\lambda_{ij}^C}{2} h^2 \phi_i^{C\dagger} \phi_j^C - \frac{\lambda_{ij}^R}{4} h^2 \phi_i^R \phi_j^R. \quad (3.9)$$

The different powers of 2 that appear in these equations are chosen to simplify the results of the following sections. In practice, we compute analytically the different coefficients of the Lagrangian via a simple program. The mass matrices of eq. (3.1) are then diagonalized numerically. The resulting mass eigenstates are labelled with a hat and their corresponding masses are referred to as \hat{m}_i^α , with α being either C , R , D or M as in eq. (3.1). The coefficients in the mass eigenstates basis are also labeled with a hat. The H_C^1 and H_C^2 terms only appear in very exotic cases.

3.2 Relevant Higgs decays and Higgs signal strengths

We present in this section all relevant decay widths of the Higgs boson. We also discuss our implementation of the constraints on the Higgs signal strengths.

3.2.1 Higgs to (dark) photons

Only complex scalars and Dirac fermions can contribute at one loop to the Higgs decay to photons or dark photons. The amplitudes take the general form

$$M^{h \rightarrow A^{(\prime)} A^{(\prime)}} = S^{h \rightarrow A^{(\prime)} A^{(\prime)}} (p_1 \cdot p_2 g_{\mu\nu} - p_{1\mu} p_{2\nu}) \epsilon_{p_1}^\nu \epsilon_{p_2}^\mu + i \tilde{S}^{h \rightarrow A^{(\prime)} A^{(\prime)}} \epsilon_{\mu\nu\alpha\beta} p_1^\alpha p_2^\beta \epsilon_{p_1}^\nu \epsilon_{p_2}^\mu, \quad (3.10)$$

with p_i and ϵ_i being respectively the momentum and polarization of the gauge bosons. The new physics contributions to the coefficients are

$$\begin{aligned} 4\pi^2 S^{h \rightarrow A^{(\prime)} A^{(\prime)}} &= \sum_i (\hat{A}_C^{(\prime)})_{ii} (\hat{A}_C^{(\prime)})_{ii} (\hat{\Omega}_C)_{ii} f_1(\hat{m}_i^C) + \sum_i (\hat{A}_D^{(\prime)})_{ii} (\hat{A}_D^{(\prime)})_{ii} \text{Re}((\hat{\Omega}_D^L)_{ii}) f_2(\hat{m}_i^D), \\ 4\pi^2 \tilde{S}^{h \rightarrow A^{(\prime)} A^{(\prime)}} &= \sum_i (\hat{A}_D^{(\prime)})_{ii} (\hat{A}_D^{(\prime)})_{ii} \text{Im}((\hat{\Omega}_D^L)_{ii}) f_3(\hat{m}_i^D), \end{aligned} \quad (3.11)$$

where

$$\begin{aligned} f_1(m_i) &= \frac{1 + 2m_i^2 C_0(0, 0, m_h^2; m_i, m_i, m_i)}{m_h^2}, \\ f_2(m_i) &= -2 \frac{m_i}{m_h^2} \left[2 + (4m_i^2 - m_h^2) C_0(0, 0, m_h^2; m_i, m_i, m_i) \right], \\ f_3(m_i) &= -2im_i C_0(0, 0, m_h^2; m_i, m_i, m_i), \end{aligned} \quad (3.12)$$

with $C_0(s_1, s_{12}, s_2; m_0, m_1, m_2)$ being the scalar three-point Passarino-Veltman function [25].² In obtaining these results, we have used the fact that certain coefficients are related by gauge symmetries (e.g., $(\hat{B}_C)_{ii} = (\hat{A}_C)_{ii}^2$) and also $(\hat{\Omega}_D^R)_{ii}^* = (\hat{\Omega}_D^L)_{ii}$. The decay widths are then

$$\Gamma^{h \rightarrow A^{(\prime)} A^{(\prime)}} = \frac{|S^{h \rightarrow A^{(\prime)} A^{(\prime)}}|^2 + |\tilde{S}^{h \rightarrow A^{(\prime)} A^{(\prime)}}|^2}{32\pi n_S} m_h^3, \quad (3.13)$$

where n_S is 2 for AA and $A'A'$ and 1 for AA' .

²All loop computations were performed with the assistance of Package-X [26].

3.2.2 Higgs decay to Z and a (dark) photon

Once again, only complex scalars and Dirac fermions can contribute to the Higgs decay to a Z boson and a photon or dark photon. The amplitudes take the general form

$$M^{h \rightarrow A^{(\prime)} Z} = S^{h \rightarrow A^{(\prime)} Z} (p_1 \cdot p_2 g_{\mu\nu} - p_{1\mu} p_{2\nu}) \epsilon_{p_1}^\nu \epsilon_{p_2}^\mu + i \tilde{S}^{h \rightarrow A^{(\prime)} Z} \epsilon_{\mu\nu\alpha\beta} p_1^\alpha p_2^\beta \epsilon_{p_1}^\nu \epsilon_{p_2}^\mu. \quad (3.14)$$

The new physics contributions to the coefficients are

$$\begin{aligned} 4\pi^2 S^{h \rightarrow A^{(\prime)} Z} &= \sum_{i,j} (\hat{A}_C^{(\prime)})_{ii} (C_C)_{ji} (\hat{\Omega}_C)_{ij} f_4(\hat{m}_i^C, \hat{m}_j^C) \\ &\quad + \sum_{i,j} (\hat{A}_D^{(\prime)})_{ii} \left[(\hat{B}_D^L)_{ji} (\hat{\Omega}_D^L)_{ij} + (\hat{B}_D^R)_{ji} (\hat{\Omega}_D^R)_{ij} \right] f_5(\hat{m}_i^D, \hat{m}_j^D) \\ &\quad + \sum_{i,j} (\hat{A}_D^{(\prime)})_{ii} \left[(\hat{B}_D^R)_{ji} (\hat{\Omega}_D^L)_{ij} + (\hat{B}_D^L)_{ji} (\hat{\Omega}_D^R)_{ij} \right] f_5(\hat{m}_j^D, \hat{m}_i^D), \\ 4\pi^2 \tilde{S}^{h \rightarrow A^{(\prime)} Z} &= \sum_{i,j} (\hat{A}_D^{(\prime)})_{ii} \left[(\hat{B}_D^L)_{ji} (\hat{\Omega}_D^L)_{ij} - (\hat{B}_D^R)_{ji} (\hat{\Omega}_D^R)_{ij} \right] f_6(\hat{m}_i^D, \hat{m}_j^D) \\ &\quad + \sum_{i,j} (\hat{A}_D^{(\prime)})_{ii} \left[(\hat{B}_D^R)_{ji} (\hat{\Omega}_D^L)_{ij} - (\hat{B}_D^L)_{ji} (\hat{\Omega}_D^R)_{ij} \right] f_6(\hat{m}_j^D, \hat{m}_i^D), \end{aligned} \quad (3.15)$$

where

$$\begin{aligned} f_4(m_i, m_j) &= \frac{1}{m_h^2 - m_Z^2} \left[1 + \frac{m_i^2 - m_j^2}{2m_h^2} \ln \frac{m_i^2}{m_j^2} + \frac{m_Z^2 (\Lambda(m_h^2, m_i, m_j) - \Lambda(m_Z^2, m_i, m_j))}{m_h^2 - m_Z^2} \right. \\ &\quad \left. + m_i^2 C_0(0, m_h^2, m_Z^2; m_i, m_i, m_j) + m_j^2 C_0(0, m_h^2, m_Z^2; m_j, m_j, m_i) \right] \\ f_5(m_i, m_j) &= -m_i f_4(m_i, m_j) + \frac{m_i}{2} C_0(0, m_h^2, m_Z^2; m_i, m_i, m_j) \\ f_6(m_i, m_j) &= -\frac{m_i}{2} C_0(0, m_h^2, m_Z^2; m_i, m_i, m_j), \end{aligned} \quad (3.16)$$

and $\Lambda(s; m_i, m_j)$ is defined as

$$\Lambda(s; m_i, m_j) = \lim_{\epsilon \rightarrow 0^+} \frac{\sqrt{\lambda(s, m_i^2, m_j^2)}}{s} \ln \left(\frac{m_i^2 + m_j^2 - s + \sqrt{\lambda(s, m_i^2, m_j^2)}}{2m_i m_j} + i\epsilon \right), \quad (3.17)$$

with $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz$ being the Källén function. The decay widths are then

$$\Gamma^{h \rightarrow A^{(\prime)} Z} = \frac{|S^{h \rightarrow A^{(\prime)} Z}|^2 + |\tilde{S}^{h \rightarrow A^{(\prime)} Z}|^2}{32\pi} \left(1 - \frac{m_Z^2}{m_h^2} \right)^3 m_h^3. \quad (3.18)$$

Note that one major qualitative difference between these decay widths and $\Gamma^{h \rightarrow A^{(\prime)} A^{(\prime)}}$ is the suppression factor of $(1 - m_Z^2/m_h^2)^3 \simeq 0.104$ from the phase space.

3.2.3 Constraints on the Higgs signal strengths

Constraints on the Higgs signal strengths are imposed by using the κ formalism [27]. Given a production mechanism i with cross section σ_i or a decay process i with width Γ_i , the parameter κ_i is defined such that

$$\kappa_i^2 = \frac{\sigma_i}{\sigma_i^{\text{SM}}} \quad \text{or} \quad \kappa_i^2 = \frac{\Gamma_i}{\Gamma_i^{\text{SM}}}, \quad (3.19)$$

where σ_i^{SM} and Γ_i^{SM} are the corresponding SM quantities. The only two Higgs couplings affected at leading order are those to AA and AZ . The corresponding κ 's are

$$\kappa_{AA}^2 = \frac{|S^{h \rightarrow AA}|^2 + |\tilde{S}^{h \rightarrow AA}|^2}{|S_{\text{SM}}^{h \rightarrow AA}|^2 + |\tilde{S}_{\text{SM}}^{h \rightarrow AA}|^2}, \quad \kappa_{ZA}^2 = \frac{|S^{h \rightarrow ZA}|^2 + |\tilde{S}^{h \rightarrow ZA}|^2}{|S_{\text{SM}}^{h \rightarrow ZA}|^2 + |\tilde{S}_{\text{SM}}^{h \rightarrow ZA}|^2}, \quad (3.20)$$

where $\tilde{S}_{\text{SM}}^{h \rightarrow AA}$ and $\tilde{S}_{\text{SM}}^{h \rightarrow ZA}$ are both zero at leading order. The invisible ($A'A'$) and semi-invisible ($A'A$ and $A'Z$) decays of the Higgs boson are accounted for by rescaling the signal strengths. The experimental input is the Higgs signal strength measurements of ref. [28] by CMS and ref. [29] by ATLAS. These studies provide the measurements, uncertainties and correlations necessary to produce our own χ^2 fit. The two searches are assumed to be uncorrelated.

3.3 Oblique parameters

We present in this section all contributions to the oblique parameters [30]. Define

$$\begin{aligned} \alpha T_\phi(m, C, D, F, G, H^1, H^2) = & \frac{1}{32\pi^2 m_W^2} \left[\sum_i (G_{ii} - c_W^2 D_{ii}) F_1(m_i) + \sum_{i,j} (|F_{ij}|^2 + |H_{ij}^1|^2 + |H_{ij}^2|^2 - c_W^2 |C_{ij}|^2) F_2(m_i, m_j) \right], \\ \alpha T_\psi(m, B^L, B^R, F^L, F^R) = & \frac{1}{32\pi^2 m_W^2} \left[(|F_{ij}^L|^2 + |F_{ij}^R|^2) F_3(m_i, m_j) + 2\text{Re}(F_{ij}^L F_{ij}^{R*}) F_4(m_i, m_j) \right. \\ & \left. - c_W^2 [(|B_{ij}^L|^2 + |B_{ij}^R|^2) F_3(m_i, m_j) + 2\text{Re}(B_{ij}^L B_{ij}^{R*}) F_4(m_i, m_j)] \right], \\ \alpha S_\phi(m, A, C) = & \frac{s_W^2 c_W^2}{8\pi^2} \sum_{i,j} \left(|C_{ij}|^2 - \frac{c_W^2 - s_W^2}{c_W s_W} C_{ij} A_{ji} - A_{ij}^2 \right) F_5(m_i, m_j), \\ \alpha S_\psi(m, A, B^L, B^R) = & \frac{s_W^2 c_W^2}{8\pi^2} \sum_{i,j} \left[(|B_{ij}^L|^2 + |B_{ij}^R|^2) F_6(m_i, m_j) + 2\text{Re}(B_{ij}^L B_{ij}^{R*}) F_7(m_i, m_j) \right. \\ & \left. - \left(\frac{c_W^2 - s_W^2}{c_W s_W} (B_{ij}^L + B_{ij}^R) A_{ij} + 2|A_{ij}|^2 \right) [F_6(m_i, m_j) + F_7(m_i, m_j)] \right], \end{aligned} \quad (3.21)$$

where

$$\begin{aligned} F_1(m_i) &= -4m_i^2(1 - \ln m_i^2), \\ F_2(m_i, m_j) &= (m_i^2 + m_j^2) \left(3 - \frac{2(m_i^4 \ln m_i^2 - m_j^4 \ln m_j^2)}{m_i^4 - m_j^4} \right), \end{aligned}$$

$$\begin{aligned}
 F_3(m_i, m_j) &= (m_i^2 + m_j^2) \left(1 - \frac{2(m_i^4 \ln m_i^2 - m_j^4 \ln m_j^2)}{m_i^4 - m_j^4} \right), \\
 F_4(m_i, m_j) &= -4m_i m_j \left(1 - \frac{(m_i^2 \ln m_i^2 - m_j^2 \ln m_j^2)}{m_i^2 - m_j^2} \right), \\
 F_5(m_i, m_j) &= -\frac{5m_i^4 - 22m_i^2 m_j^2 + 5m_j^4}{9(m_i^2 - m_j^2)^2} + \frac{2 \left[m_i^4(m_i^2 - 3m_j^2) \ln m_i^2 - m_j^4(m_j^2 - 3m_i^2) \ln m_j^2 \right]}{3(m_i^2 - m_j^2)^3}, \\
 F_6(m_i, m_j) &= -\frac{4(m_i^4 - 8m_i^2 m_j^2 + m_j^4)}{9(m_i^2 - m_j^2)^2} + \frac{4 \left[m_i^4(m_i^2 - 3m_j^2) \ln m_i^2 - m_j^4(m_j^2 - 3m_i^2) \ln m_j^2 \right]}{3(m_i^2 - m_j^2)^3}, \\
 F_7(m_i, m_j) &= -2m_i m_j \frac{(m_i^2 + m_j^2)}{(m_i^2 - m_j^2)^2} \left(1 - \frac{2m_i^2 m_j^2}{m_i^4 - m_j^4} \ln \frac{m_i^2}{m_j^2} \right). \tag{3.22}
 \end{aligned}$$

The oblique parameters are then conveniently given by

$$\begin{aligned}
 T &= T_\phi(\hat{m}_C, \hat{C}_C, \hat{D}_C, \hat{F}_C, \hat{G}_C, \hat{H}_C^1, \hat{H}_C^2) + \frac{1}{2} T_\phi(\hat{m}_R, \hat{C}_R, \hat{D}_R, 0, \hat{G}_R, 0, 0) \\
 &\quad + T_\phi(\hat{m}_{RC}, 0, 0, \hat{F}_{RC}, 0, 0, 0) + T_\psi(\hat{m}_D, \hat{B}_D^L, \hat{B}_D^R, \hat{F}_D^L, \hat{F}_D^R) + \frac{1}{2} T_\psi(\hat{m}_M, \hat{B}_M^L, \hat{B}_M^R, 0, 0) \\
 &\quad + T_\psi(\hat{m}_{MD}, 0, 0, \hat{F}_{MD}^L, \hat{F}_{MD}^R), \\
 S &= S_\phi(\hat{m}_C, \hat{A}_C, \hat{C}_C) + \frac{1}{2} S_\phi(\hat{m}_R, 0, \hat{C}_R) + S_\psi(\hat{m}_D, \hat{A}_D, \hat{B}_D^L, \hat{B}_D^R) + \frac{1}{2} S_\psi(\hat{m}_M, 0, \hat{B}_M^L, \hat{B}_M^R), \tag{3.23}
 \end{aligned}$$

where

$$\begin{aligned}
 \hat{m}_{RC} &= (\hat{m}_R \ \hat{m}_C), & \hat{m}_{MD} &= (\hat{m}_M \ \hat{m}_D), \\
 \hat{F}_{RC} &= \begin{pmatrix} 0 & \hat{F}_{RC}^1 \\ \hat{F}_{RC}^2 & 0 \end{pmatrix}, & \hat{F}_{MD}^{L/R} &= \begin{pmatrix} 0 & \hat{F}_{MD}^{1L/R} \\ \hat{F}_{MD}^{2L/R} & 0 \end{pmatrix}. \tag{3.24}
 \end{aligned}$$

A χ^2 fit is performed using [31]

$$S = -0.01 \pm 0.07, \quad T = 0.04 \pm 0.06, \tag{3.25}$$

with a correlation of 0.92.

3.4 Unitarity

Consider the amplitude \mathcal{M} of a $2 \rightarrow 2$ scattering process. It can be decomposed in terms of the Legendre polynomials $P_\ell(\cos \theta)$ as

$$\mathcal{M} = 16\pi \sum_\ell (2\ell + 1) a_\ell P_\ell(\cos \theta). \tag{3.26}$$

For complex scalars, real scalars, Dirac spinors and Majorana spinors, we will define matrices of a_ℓ coefficients called a_ℓ^{mat} . Unitarity will then impose

$$a_\ell^{\text{max}} \equiv \max \left(\left| \text{Re} \left(a_\ell^{\text{eig}} \right) \right| \right) \leq \frac{1}{2}, \tag{3.27}$$

where a_ℓ^{eig} is the list of eigenvalues of a_ℓ^{mat} . In the presence of identical particles in the incoming or outgoing state, a_0 is multiplied by a factor of $1/\sqrt{2}$. The limit of very high energy is assumed, in which case we can work with gauge eigenstates, though the basis invariance of the final results will be apparent. We will consider scattering from two components of an inert multiplet to two Higgs bosons and take inspiration from ref. [32].

As a side note, we mention that the treatment of the fermion case differs from that of refs. [19, 20]. Although the treatment of these references is very convenient for simpler cases, its generalization to more general cases can prove cumbersome.

Complex scalars. Consider the basis of field pairs

$$hh, \phi_1^{C\dagger} \phi_1^C, \phi_1^{C\dagger} \phi_2^C, \dots, \phi_1^{C\dagger} \phi_n^C, \phi_2^{C\dagger} \phi_1^C, \phi_2^{C\dagger} \phi_2^C, \dots, \phi_2^{C\dagger} \phi_n^C, \dots, \phi_n^{C\dagger} \phi_1^C, \phi_n^{C\dagger} \phi_2^C, \dots, \phi_n^{C\dagger} \phi_n^C, \quad (3.28)$$

where n is the number of complex scalars. The a_0^{mat} matrix for scattering from one pair to another is

$$a_0^{\text{mat}} = -\frac{1}{16\sqrt{2}\pi} \begin{pmatrix} 0 & v^C \\ v^{C\dagger} & 0 \end{pmatrix}, \quad (3.29)$$

where

$$v^C = (\lambda_{11}^C \ \lambda_{12}^C \ \dots \ \lambda_{1n}^C \ \lambda_{21}^C \ \lambda_{22}^C \ \dots \ \lambda_{2n}^C \ \dots \ \lambda_{n1}^C \ \lambda_{n2}^C \ \dots \ \lambda_{nn}^C). \quad (3.30)$$

In practice, this simply gives

$$a_0^{\text{max}} = \frac{\sqrt{\text{Tr}[(\lambda^C)^2]}}{16\sqrt{2}\pi}. \quad (3.31)$$

Real scalars. Consider the basis of field pairs

$$hh, \phi_1^R \phi_1^R, \phi_1^R \phi_2^R, \dots, \phi_1^R \phi_n^R, \phi_2^R \phi_2^R, \phi_2^R \phi_3^R, \dots, \phi_2^R \phi_n^R, \dots, \phi_n^R \phi_n^R, \quad (3.32)$$

where n is the number of real scalars. The a_0^{mat} matrix for scattering from one pair to another is

$$a_0^{\text{mat}} = -\frac{1}{16\sqrt{2}\pi} \begin{pmatrix} 0 & v^R \\ v^{R\dagger} & 0 \end{pmatrix}, \quad (3.33)$$

where

$$v^R = \left(\frac{\lambda_{11}^R}{\sqrt{2}} \ \lambda_{12}^R \ \dots \ \lambda_{1n}^R \ \frac{\lambda_{22}^R}{\sqrt{2}} \ \lambda_{23}^R \ \dots \ \lambda_{2n}^R \ \dots \ \frac{\lambda_{nn}^R}{\sqrt{2}} \right). \quad (3.34)$$

In practice, this simply gives

$$a_0^{\text{max}} = \frac{\sqrt{\text{Tr}[(\lambda^R)^2]}}{32\pi}. \quad (3.35)$$

Dirac fermions. Consider the basis of field pairs

$$hh, \bar{\psi}_1^D \psi_1^D, \bar{\psi}_1^D \psi_2^D, \dots, \bar{\psi}_1^D \psi_n^D, \bar{\psi}_2^D \psi_1^D, \bar{\psi}_2^D \psi_2^D, \dots, \bar{\psi}_2^D \psi_n^D, \dots, \bar{\psi}_n^D \psi_1^D, \bar{\psi}_n^D \psi_2^D, \dots, \bar{\psi}_n^D \psi_n^D, \quad (3.36)$$

where n is the number of Dirac fermions. The a_0 coefficients for $\bar{\psi}_i^D \psi_j^D \rightarrow hh$ are null. The a_1^{mat} matrix for scattering from one pair to another is

$$a_1^{\text{mat}} = -\frac{1}{16\sqrt{2}} \begin{pmatrix} 0 & v^D \\ v^{D\dagger} & 0 \end{pmatrix}, \quad (3.37)$$

where

$$v^D = \left(A_{11}^D \ A_{12}^D \ \dots \ A_{1n}^D \ A_{21}^D \ A_{22}^D \ \dots \ A_{2n}^D \ A_{n1}^D \ A_{n2}^D \ \dots \ A_{nn}^D \right), \quad (3.38)$$

with

$$A_{ij}^D = \left(\sum_k (\Omega_D^R)_{ik} (\Omega_D^L)_{kj} \ \sum_k (\Omega_D^L)_{ik} (\Omega_D^R)_{kj} \right), \quad (3.39)$$

in the helicity basis ($\uparrow\downarrow, \downarrow\uparrow$). In practice, this simply gives

$$a_1^{\max} = \frac{\sqrt{\text{Tr}[(\Omega_D^L \Omega_D^R)^2]}}{16}. \quad (3.40)$$

Majorana fermions. Consider the basis of field pairs

$$hh, \psi_1^M \psi_1^M, \psi_1^M \psi_2^M, \dots, \psi_1^M \psi_n^M, \psi_2^M \psi_2^M, \psi_2^M \psi_3^M, \dots, \psi_2^M \psi_n^M, \dots, \psi_n^M \psi_n^M, \quad (3.41)$$

where n is the number of Majorana fermions. The a_0 coefficients for $\psi_i^M \psi_j^M \rightarrow hh$ are null. The a_1^{mat} matrix for scattering from one pair to another is

$$a_1^{\text{mat}} = -\frac{1}{16\sqrt{2}} \begin{pmatrix} 0 & v^M \\ v^{M\dagger} & 0 \end{pmatrix}, \quad (3.42)$$

where

$$v^M = \left(\frac{A_{11}^M}{\sqrt{2}} \ A_{12}^M \ \dots \ A_{1n}^M \ \frac{A_{22}^M}{\sqrt{2}} \ A_{23}^M \ \dots \ A_{2n}^M \ \dots \ \frac{A_{nn}^M}{\sqrt{2}} \right), \quad (3.43)$$

with

$$A_{ij}^M = \left(\sum_k (\Omega_M^R)_{ik} (\Omega_M^L)_{kj} \ \sum_k (\Omega_M^L)_{ik} (\Omega_M^R)_{kj} \right), \quad (3.44)$$

in the helicity basis ($\uparrow\downarrow, \downarrow\uparrow$). In practice, this simply gives

$$a_1^{\max} = \frac{\sqrt{\text{Tr}[(\Omega_M^L \Omega_M^R)^2]}}{16\sqrt{2}}, \quad (3.45)$$

where we have assumed that Ω_M^R and Ω_M^L are constructed such that they are symmetric, which can always be performed.

3.5 Triple Higgs coupling

With our formalism, the triple Higgs coupling can easily be computed using the Coleman-Weinberg potential [33]. In the \overline{MS} renormalization scheme, it is given by

$$V_{CW} = \frac{1}{64\pi^2} \left[\sum_i (\hat{m}_R^4)_{ii} \left(\ln \left(\frac{(\hat{m}_R^2)_{ii}}{Q^2} \right) - \frac{3}{2} \right) + 2 \sum_i (\hat{m}_C^4)_{ii} \left(\ln \left(\frac{(\hat{m}_C^2)_{ii}}{Q^2} \right) - \frac{3}{2} \right) \right. \\ \left. - 2 \sum_i (\hat{m}_M^4)_{ii} \left(\ln \left(\frac{(\hat{m}_M^2)_{ii}}{Q^2} \right) - \frac{3}{2} \right) - 4 \sum_i (\hat{m}_D^4)_{ii} \left(\ln \left(\frac{(\hat{m}_D^2)_{ii}}{Q^2} \right) - \frac{3}{2} \right) \right], \quad (3.46)$$

where the different masses are understood to be the field-dependent masses and Q is some scale that cancels in the final result. Requesting to reproduce the correct mass and VEV of the Higgs, the triple Higgs coupling λ_{hhh} is given by

$$\frac{\lambda_{hhh}}{3!} = \frac{m_h^2}{2v} + \left[\left(\frac{1}{2v^2} \frac{\partial}{\partial h} - \frac{1}{2v} \frac{\partial^2}{\partial h^2} + \frac{1}{3!} \frac{\partial^3}{\partial h^3} \right) V_{CW} \right] \Big|_{h=0}, \quad (3.47)$$

where we use the convention that $v \approx 246$ GeV.

4 Results

We present in this section the potential contributions of inert multiplets to $\text{BR}(h \rightarrow ZA^{(\prime)})$ and the triple Higgs coupling. We first discuss our scanning procedure and then present results in the presence of either one or multiple interaction terms.

4.1 Scanning procedure

For each benchmark model considered, the entire parameter space is scanned using a Markov chain with the Metropolis-Hasting algorithm. To increase the density of points toward the limits and therefore obtain faster convergence, a non-flat prior is chosen proportional to

$$\left(\frac{\delta\lambda_{hhh}}{a}\right)^2 + \left(\frac{\delta\text{BR}(h \rightarrow AZ)}{b}\right)^2 + \left(\frac{\text{BR}(h \rightarrow A'Z)}{c}\right)^2 \quad (4.1)$$

where we set $a = b = 0.1$ and $c = 0.001$. The relative deviation of the triple Higgs coupling from its SM value is defined as

$$\delta\lambda_{hhh} = \frac{\lambda_{hhh} - \lambda_{hhh}^{\text{SM}}}{\lambda_{hhh}^{\text{SM}}}. \quad (4.2)$$

The relative deviation of $\text{BR}(h \rightarrow AZ)$ from its SM value is defined as

$$\delta\text{BR}(h \rightarrow AZ) = \frac{\text{BR}(h \rightarrow AZ) - \text{BR}(h \rightarrow AZ)_{\text{SM}}}{\text{BR}(h \rightarrow AZ)_{\text{SM}}}. \quad (4.3)$$

We have verified that the Markov chain converges to the same results irrespective of the choice of prior. To keep the number of figures reasonable, the results are finely binned and presented as histograms. When multiple contractions of $\text{SU}(2)_L$ indices are possible, all contractions are considered. For scalar multiplets, points are dismissed in the presence of tachyons, as the scalars acquiring VEVs fall outside the scope of this work and can even potentially break electromagnetism.³ Points that contain charged particles below 100 GeV are also rejected, as such particles are ruled out by LEP [34, 35]. As discussed in refs. [19, 20], constraints from the Higgs signal strengths can be partially avoided by having purely imaginary $\hat{\Omega}_D^{L/R}$ couplings. This however leads to a strongly excluded contribution to the electron electric dipole moment, which forces $\hat{\Omega}_D^{L/R}$ to be almost purely real. In practice, the results are equivalent to taking $A_{L/R}$ to be real. For the sake of simplicity and to speed up computations, we will therefore take $A_{L/R}$ real. For complex scalars, unitarity imposes a constraint on the couplings to dark photons of [32]

$$\sum_i (A'_C)_{ii}^4 < 16\pi^2. \quad (4.4)$$

³We mention that Higgs vacuum instability could also be a problem, for example because the fermionic multiplets tend to make the Higgs quartic run more negative. However, bounds on vacuum stability would imply the assumption that no new physics exists between the electroweak scale and the instability scale. Since we cannot guarantee this assumption is respected, we do not impose such bounds. In the same vein, the existence of additional minima that break electromagnetism could potentially be problematic. However, such bounds would depend on the full potential. The additional terms could address the stability issues while leaving our observables unchanged at leading order. Because of this, we do not consider this issue any further.

For Dirac fermions, we take the less restrictive $\max(|A'_D|_{ii}) < \sqrt{4\pi}$. As a validation procedure, we have verified that all divergences in the loop decay widths and oblique parameters cancel for all points of every benchmark. Finally, as pointed out in refs. [19, 20], the Higgs signal strengths are considerably insensitive to the scenario that the BSM contributions to $S^{h \rightarrow AA}$ ($S^{h \rightarrow AZ}$) are about $-2S_{\text{SM}}^{h \rightarrow AA}$ ($-2S_{\text{SM}}^{h \rightarrow AZ}$), since this simply changes the sign of $S^{h \rightarrow AA}$ ($S^{h \rightarrow AZ}$) which the Higgs signal strengths are insensitive to. This would however require tremendous tuning, if even possible, and we will ignore such points.

4.2 One interaction term

We first present results in the presence of a single interaction term of the types presented in section 2. The chosen benchmarks are presented in tables 1 and 2. Each entry contains the relevant fields and their interaction Lagrangian as defined in section 2. All bounds are at 95% confidence level.

First, figure 1 shows the relative deviation of $\text{BR}(h \rightarrow AZ)$ from its SM value as a function of the mass of the lightest new particles m_{min} . As can be seen, deviations of $\pm \mathcal{O}(20\%)$ are possible. The scalar cases IV and V are especially constrained and do not lead to any sizable contribution to this observable. This will be the case for all other observables. The only exception is if the decay of the Higgs boson to two new scalars is allowed, in which case $\text{BR}(h \rightarrow AZ)$ can be reduced. Many other benchmark models with a single interaction term have been analyzed, but none allowed for a significantly larger range of $\delta\text{BR}(h \rightarrow AZ)$.

Second, figure 2 displays the allowed range of $\text{BR}(h \rightarrow A'Z)$. As can be seen, this branching ratio could in principle reach values above 1%. However, doing so would require new neutral particles just above half the mass of the Higgs boson, which requires careful fine-tuning between the different parameters and is not possible for every model. Otherwise, a branching ratio of $\mathcal{O}(0.1\%)$ is relatively easy to obtain. We mention that the coefficient $S^{h \rightarrow A'Z}$ is less constrained than $S^{h \rightarrow A'A}$. As alluded to before, the reason that the constraints on $\text{BR}(h \rightarrow A'Z)$ are generally stronger than those on $\text{BR}(h \rightarrow A'A)$ of refs. [19, 20] is the phase-space suppression in eq. (3.18). The branching ratio $S^{h \rightarrow A'Z}$ is generally optimized for couplings of the dark photon A'_D or A'_C of $\mathcal{O}(1)$, with the exact value being model dependent.

Third, figure 3 shows examples of the relative deviation of the triple Higgs coupling $\delta\lambda_{hhh}$ as a function of m_{min} . As can be seen, certain inert multiplets can easily give large contributions to the triple Higgs couplings, including both positive and negative ones. Even larger corrections are possible, though current constraints on the Higgs self-interactions render such numbers less relevant [36].

We mention that constraints on direct production of the new particles would generally apply to a fully defined model. However, doing so in our case would require both specifying the Lagrangian beyond the interaction terms and searches that do not presently exist, making it practically unfeasible. If this were possible, constraints could be enhanced.

Name	Fields	Gauge numbers	Type	Lagrangian
FA	ψ_1	$(\mathbf{2}, 1/2, 1)$	Complex	$F(\psi_1, \psi_2)$
	ψ_2	$(\mathbf{1}, 0, 1)$	Complex	
FB	ψ_1	$(\mathbf{1}, 0, 1)$	Complex	$F(\psi_1, \psi_2)$
	ψ_2	$(\mathbf{2}, -1/2, 1)$	Complex	
FC	ψ_1	$(\mathbf{3}, 0, 1)$	Complex	$F(\psi_1, \psi_2)$
	ψ_2	$(\mathbf{2}, -1/2, 1)$	Complex	
FD	ψ_1	$(\mathbf{2}, 1/2, 1)$	Complex	$F(\psi_1, \psi_2)$
	ψ_2	$(\mathbf{3}, 0, 1)$	Complex	
FE	ψ_1	$(\mathbf{2}, 1/2, 0)$	Complex	$F(\psi_1, \psi_2)$
	ψ_2	$(\mathbf{1}, 0, 0)$	Real	
FF	ψ_1	$(\mathbf{1}, 0, 0)$	Real	$F(\psi_1, \psi_2)$
	ψ_2	$(\mathbf{2}, -1/2, 0)$	Complex	
FG	ψ_1	$(\mathbf{3}, 0, 0)$	Real	$F(\psi_1, \psi_2)$
	ψ_2	$(\mathbf{2}, -1/2, 0)$	Complex	
FH	ψ_1	$(\mathbf{2}, 1/2, 0)$	Complex	$F(\psi_1, \psi_2)$
	ψ_2	$(\mathbf{3}, 0, 0)$	Real	
S1A	ϕ_1	$(\mathbf{2}, 1/2, 1)$	Complex	$S_1(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{1}, 0, 1)$	Complex	
S1B	ϕ_1	$(\mathbf{1}, 0, 1)$	Complex	$S_1(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{2}, -1/2, 1)$	Complex	
S1C	ϕ_1	$(\mathbf{3}, 0, 1)$	Complex	$S_1(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{2}, -1/2, 1)$	Complex	
S1D	ϕ_1	$(\mathbf{2}, 1/2, 1)$	Complex	$S_1(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{3}, 0, 1)$	Complex	
S1E	ϕ_1	$(\mathbf{2}, 1/2, 0)$	Complex	$S_1(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{1}, 0, 0)$	Real	
S1F	ϕ_1	$(\mathbf{1}, 0, 0)$	Real	$S_1(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{2}, -1/2, 0)$	Complex	
S1G	ϕ_1	$(\mathbf{3}, 0, 0)$	Real	$S_1(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{2}, -1/2, 0)$	Complex	
S1H	ϕ_1	$(\mathbf{2}, 1/2, 0)$	Complex	$S_1(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{3}, 0, 0)$	Real	
S2A	ϕ	$(\mathbf{1}, 1, 1)$	Complex	$S_2(\phi)$
S2B	ϕ	$(\mathbf{1}, 2, 1)$	Complex	$S_2(\phi)$
S2C	ϕ	$(\mathbf{2}, 1/2, 1)$	Complex	$S_2(\phi)$
S2D	ϕ	$(\mathbf{2}, 3/2, 1)$	Complex	$S_2(\phi)$

Table 1. Different benchmarks for the single interaction term case. The type represents whether the multiplet is real or complex. The gauge numbers are in the format $(\text{SU}(2)_L, \text{U}(1)_Y, \text{U}(1)')$.

Name	Fields	Gauge numbers	Type	Lagrangian
S2E	ϕ	$(\mathbf{3}, 0, 1)$	Complex	$S_2(\phi)$
S2F	ϕ	$(\mathbf{3}, 1, 1)$	Complex	$S_2(\phi)$
S2G	ϕ	$(\mathbf{1}, 0, 0)$	Real	$S_2(\phi)$
S2H	ϕ	$(\mathbf{3}, 0, 0)$	Real	$S_2(\phi)$
S3A	ϕ_1	$(\mathbf{1}, 1, 1)$	Complex	$S_3(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{1}, 1, 1)$	Complex	
S3B	ϕ_1	$(\mathbf{1}, 2, 1)$	Complex	$S_3(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{1}, 2, 1)$	Complex	
S3C	ϕ_1	$(\mathbf{2}, 1/2, 1)$	Complex	$S_3(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{2}, 1/2, 1)$	Complex	
S3D	ϕ_1	$(\mathbf{2}, 3/2, 1)$	Complex	$S_3(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{2}, 3/2, 1)$	Complex	
S3E	ϕ_1	$(\mathbf{3}, 0, 1)$	Complex	$S_3(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{3}, 0, 1)$	Complex	
S3F	ϕ_1	$(\mathbf{3}, 1, 1)$	Complex	$S_3(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{3}, 1, 1)$	Complex	
S3G	ϕ_1	$(\mathbf{1}, 0, 0)$	Real	$S_3(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{1}, 0, 0)$	Real	
S3H	ϕ_1	$(\mathbf{3}, 0, 0)$	Complex	$S_3(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{3}, 0, 0)$	Real	
S3I	ϕ_1	$(\mathbf{3}, 0, 0)$	Real	$S_3(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{3}, 0, 0)$	Real	
S4A	ϕ_1	$(\mathbf{2}, 1, 1)$	Complex	$S_4(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{2}, 0, 1)$	Complex	
S4B	ϕ_1	$(\mathbf{3}, 1, 1)$	Complex	$S_4(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{1}, 0, 1)$	Complex	
S4C	ϕ_1	$(\mathbf{3}, 1, 1)$	Complex	$S_4(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{3}, 0, 1)$	Complex	
S4D	ϕ_1	$(\mathbf{3}, 1, 0)$	Complex	$S_4(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{1}, 0, 0)$	Real	
S4E	ϕ_1	$(\mathbf{3}, 1, 0)$	Complex	$S_4(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{3}, 0, 0)$	Real	
S5A	ϕ	$(\mathbf{2}, 1/2, 0)$	Complex	$S_5(\phi)$
S5B	ϕ	$(\mathbf{4}, 1/2, 0)$	Complex	$S_5(\phi)$
S6A	ϕ	$(\mathbf{2}, 0, 0)$	Complex	$S_6(\phi)$
S6B	ϕ	$(\mathbf{3}, 0, 0)$	Complex	$S_6(\phi)$
S6C	ϕ	$(\mathbf{3}, 0, 0)$	Real	$S_6(\phi)$

Table 2. Table 1 continued.

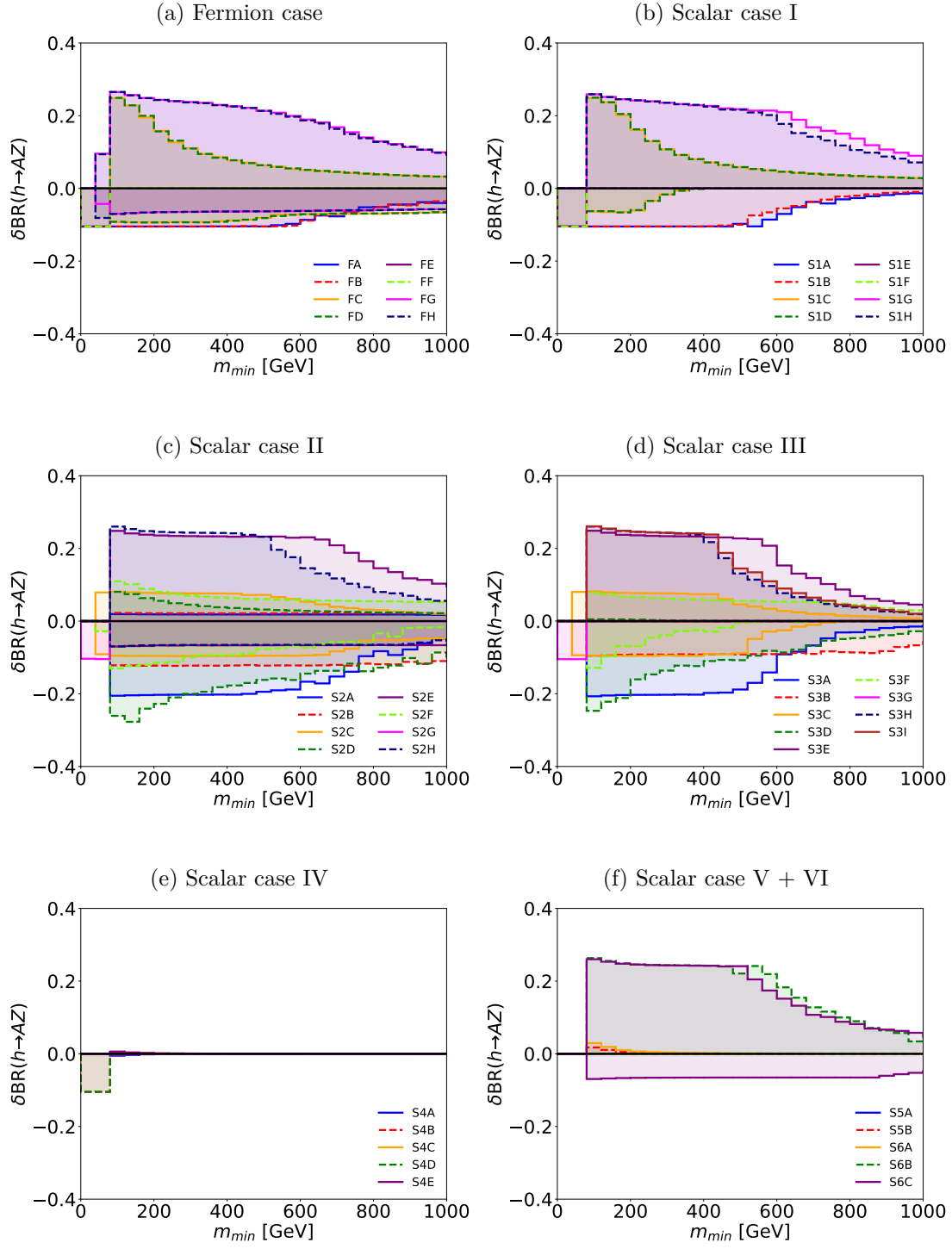


Figure 1. Allowed range of $\delta\text{BR}(h \rightarrow AZ)$ for different cases of a single interaction term. See tables 1 and 2 for a description of the benchmarks.

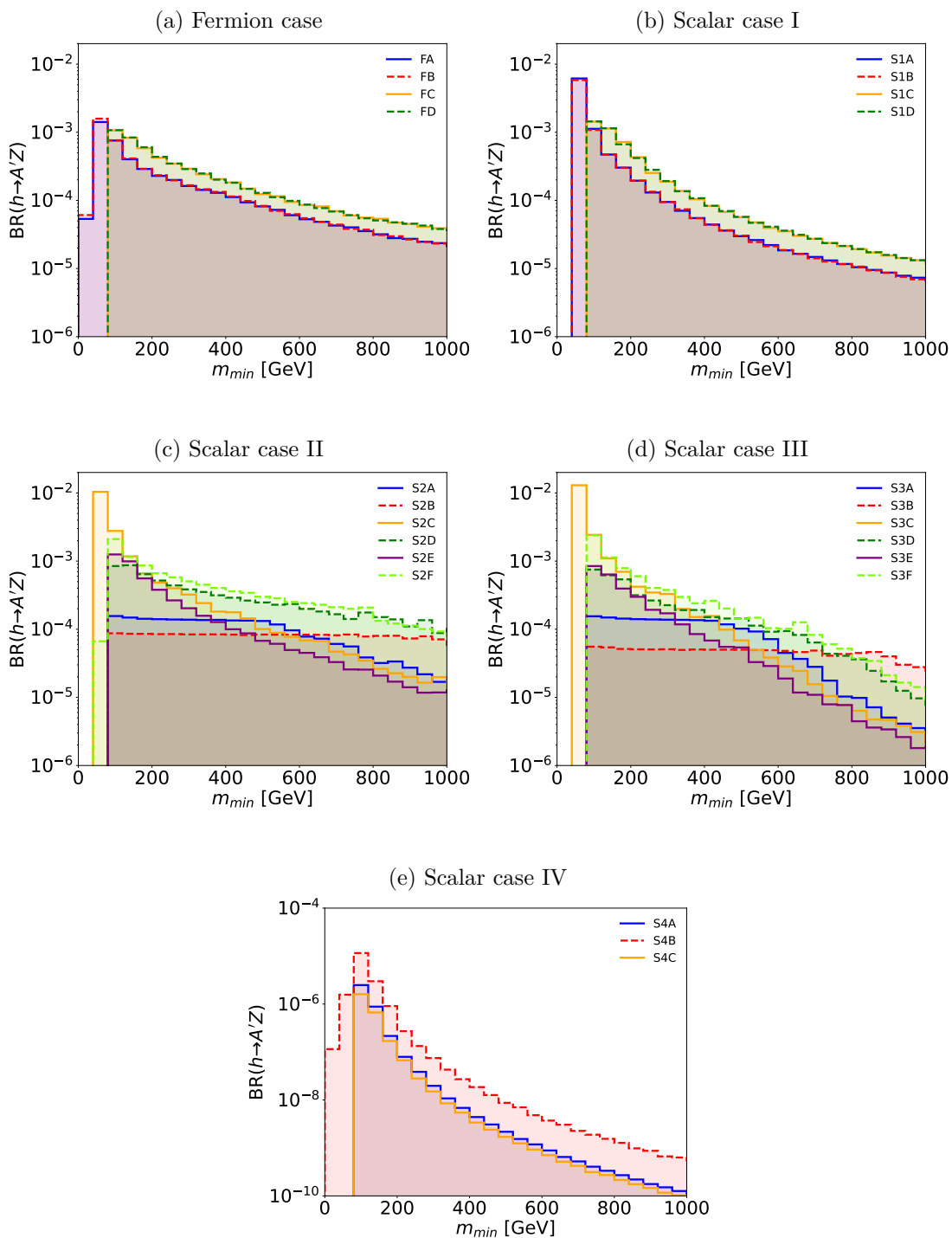


Figure 2. Allowed range of $BR(h \rightarrow A'Z)$ for different cases of a single interaction term. See tables 1 and 2 for a description of the benchmarks. The missing benchmarks do not contribute to this decay.

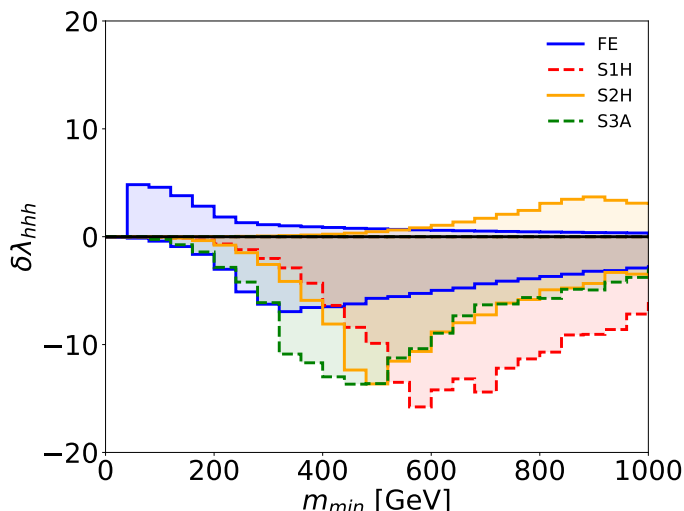


Figure 3. Examples of allowed $\delta\lambda_{hhh}$ range. See tables 1 and 2 for a description of the benchmarks.

Name	Fields	Gauge numbers	Type	Lagrangian
C1	ϕ_1	$(\mathbf{3}, 0, 1)$	Complex	$S_2(\phi_1) + S_2(\phi_2)$
	ϕ_2	$(\mathbf{1}, 1, 1)$	Complex	
C2	ϕ_1	$(\mathbf{3}, 0, 1)$	Complex	$S_2(\phi_1) + S_2(\phi_2) + S_1(\phi_2, \phi_1)$
	ϕ_2	$(\mathbf{2}, 1/2, 1)$	Complex	
C3	ψ_1	$(\mathbf{2}, 1/2, 1)$	Complex	$F(\psi_1, \psi_2) + F(\psi_1, \psi_3)$
	ψ_2	$(\mathbf{3}, 0, 1)$	Complex	
	ψ_3	$(\mathbf{1}, 0, 1)$	Complex	
C4	ϕ_1	$(\mathbf{3}, 0, 1)$	Complex	$S_2(\phi_1) + S_2(\phi_2) + S_3(\phi_1, \phi_2)$
	ϕ_2	$(\mathbf{3}, 0, 1)$	Complex	
C5	ϕ_1	$(\mathbf{3}, 0, 1)$	Complex	$S_2(\phi_1) + S_2(\phi_2) + S_2(\phi_3) + S_1(\phi_2, \phi_1) + S_1(\phi_1, \phi_3)$
	ϕ_2	$(\mathbf{2}, 1/2, 1)$	Complex	
	ϕ_3	$(\mathbf{2}, -1/2, 1)$	Complex	

Table 3. Benchmarks with multiple interaction terms.

4.3 Multiple interaction terms

We now present results in the presence of multiple interactions of the types presented in section 2. The chosen benchmarks are presented in table 3 and the results in figure 4.

As can be seen, the range of allowed $\delta\text{BR}(h \rightarrow AZ)$ can be considerably extended. However, this requires careful tuning between different interaction terms to ensure that their contributions to constrained observables interfere destructively and constructively for $\delta\text{BR}(h \rightarrow AZ)$. This is why much larger values of $\delta\text{BR}(h \rightarrow AZ)$ are allowed for combinations of interactions than with a single interaction term. This is reflected in practice by the numerical difficulty in obtaining such large values and is also why the contours are less smooth in these plots. Additionally, such careful tuning is not even possible for all models, as certain combinations of interaction terms do not allow for a larger range than

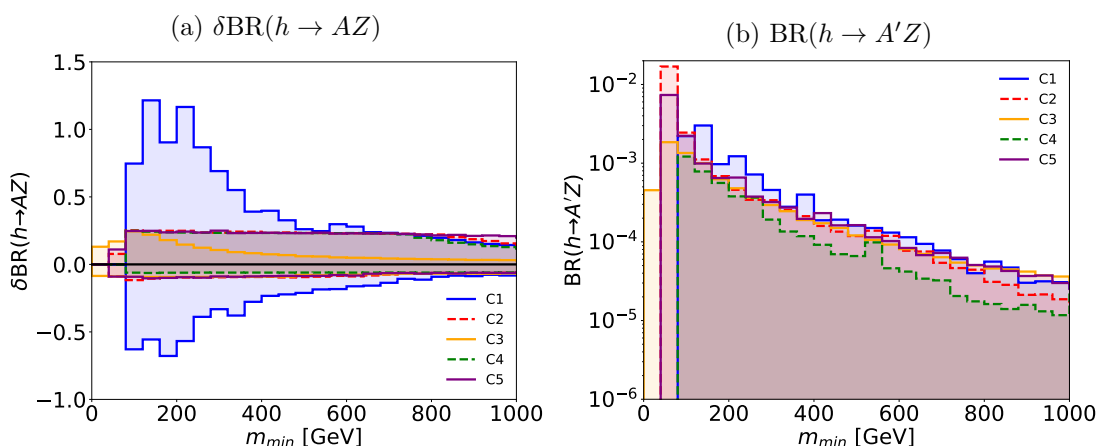


Figure 4. Allowed ranges of (a) $\delta\text{BR}(h \rightarrow AZ)$ and (b) $\text{BR}(h \rightarrow A'Z)$ for the benchmarks with multiple interactions. See table 3 for a description of the benchmarks.

the individual terms do. We do not observe any qualitative enhancement in the allowed range of $\text{BR}(h \rightarrow A'Z)$.

5 Conclusion

Inert multiplets can only interact with the Higgs boson via a finite number of interaction terms and the forms of these terms are enough to determine their leading contributions to many Higgs properties. In this paper, we have studied the contributions of inert multiplets to the branching ratios of the Higgs to a Z boson and either a photon or a dark photon and to the triple Higgs coupling.

We reach the following conclusions. The branching ratio $\text{BR}(h \rightarrow AZ)$ can deviate from its SM value by $\mathcal{O}(20\%)$ even for simple models of inert multiplets. Larger deviations are possible, but require complicated models and precise fine-tuning. Therefore, the current measurement of $\text{BR}(h \rightarrow AZ)$ being 2.1 ± 0.7 larger than the Standard Model value might be explainable by inert multiplets, but any such model would have to be contrived and very fine-tuned.

The branching ratio $\text{BR}(h \rightarrow A'Z)$ could in principle reach above 1%. However, this would again require considerable fine-tuning, and values of $\mathcal{O}(0.1\%)$ are more realistic. Such small branching ratios are difficult to probe at the LHC for two reasons. First, hadronic decays of the Z boson would be difficult to reconstruct and its branching ratio to leptons is small. Second, the fact that the mass of the Z boson is not that far off from that of the Higgs boson would lead to much less missing transverse momentum than for $h \rightarrow AA'$. It therefore seems unlikely that this decay channel would be observable at the LHC.

Large deviations on the triple Higgs coupling are however perfectly possible and could be an ideal observable to probe the inert multiplet scenarios considered in this work. Though a full study is beyond the scope of this paper, it seems to indicate that a first-order electroweak phase transition could be explained by inert multiplets.

Acknowledgments

This work was supported by the National Science and Technology Council under Grant No. NSTC-111-2112-M-002-018-MY3, the Ministry of Education (Higher Education Sprout Project NTU-112L104022), and the National Center for Theoretical Sciences of Taiwan.

Open Access. This article is distributed under the terms of the Creative Commons Attribution License ([CC-BY 4.0](https://creativecommons.org/licenses/by/4.0/)), which permits any use, distribution and reproduction in any medium, provided the original author(s) and source are credited.

References

- [1] L. Lopez Honorez, E. Nezri, J.F. Oliver and M.H.G. Tytgat, *The Inert Doublet Model: An Archetype for Dark Matter*, *JCAP* **02** (2007) 028 [[hep-ph/0612275](#)] [[INSPIRE](#)].
- [2] M. Gustafsson, E. Lundstrom, L. Bergstrom and J. Edsjo, *Significant Gamma Lines from Inert Higgs Dark Matter*, *Phys. Rev. Lett.* **99** (2007) 041301 [[astro-ph/0703512](#)] [[INSPIRE](#)].
- [3] F.P. Huang and J.-H. Yu, *Exploring inert dark matter blind spots with gravitational wave signatures*, *Phys. Rev. D* **98** (2018) 095022 [[arXiv:1704.04201](#)] [[INSPIRE](#)].
- [4] A. Betancur, G. Palacio and A. Rivera, *Inert doublet as multicomponent dark matter*, *Nucl. Phys. B* **962** (2021) 115276 [[arXiv:2002.02036](#)] [[INSPIRE](#)].
- [5] T. Hambye and M.H.G. Tytgat, *Electroweak symmetry breaking induced by dark matter*, *Phys. Lett. B* **659** (2008) 651 [[arXiv:0707.0633](#)] [[INSPIRE](#)].
- [6] G. Gil, P. Chankowski and M. Krawczyk, *Inert Dark Matter and Strong Electroweak Phase Transition*, *Phys. Lett. B* **717** (2012) 396 [[arXiv:1207.0084](#)] [[INSPIRE](#)].
- [7] C.-H. Chen, C.-W. Chiang and C.-W. Su, *Top-quark FCNC decays, LFVs, lepton $g - 2$, and W mass anomaly with inert charged Higgses*, [arXiv:2301.07070](#) [[INSPIRE](#)].
- [8] ATLAS collaboration, *Evidence for the Higgs boson decay to a Z boson and a photon at the LHC*, ATLAS-CONF-2023-025 (2023) [[INSPIRE](#)].
- [9] ATLAS collaboration, *Evidence for the Higgs boson decay to a Z boson and a photon at the LHC*, ATLAS-CONF-2023-025 (2023) [[INSPIRE](#)].
- [10] P. Archer-Smith, D. Stolarski and R. Vega-Morales, *On new physics contributions to the Higgs decay to $Z\gamma$* , *JHEP* **10** (2021) 247 [[arXiv:2012.01440](#)] [[INSPIRE](#)].
- [11] R. Benbrik et al., *Higgs-like particle decays into γZ and $\gamma\gamma$: Fingerprints of some non-supersymmetric models*, *Nucl. Phys. B* **990** (2023) 116154 [[arXiv:2211.12546](#)] [[INSPIRE](#)].
- [12] B. Holdom, *Two U(1)'s and Epsilon Charge Shifts*, *Phys. Lett. B* **166** (1986) 196 [[INSPIRE](#)].
- [13] E. Gabrielli, M. Heikinheimo, B. Mele and M. Raidal, *Dark photons and resonant monophoton signatures in Higgs boson decays at the LHC*, *Phys. Rev. D* **90** (2014) 055032 [[arXiv:1405.5196](#)] [[INSPIRE](#)].
- [14] S. Biswas, E. Gabrielli, M. Heikinheimo and B. Mele, *Higgs-boson production in association with a dark photon in e^+e^- collisions*, *JHEP* **06** (2015) 102 [[arXiv:1503.05836](#)] [[INSPIRE](#)].

- [15] S. Biswas, E. Gabrielli, M. Heikinheimo and B. Mele, *Dark-Photon searches via Higgs-boson production at the LHC*, *Phys. Rev. D* **93** (2016) 093011 [[arXiv:1603.01377](#)] [[INSPIRE](#)].
- [16] S. Biswas, E. Gabrielli, M. Heikinheimo and B. Mele, *Dark-photon searches via ZH production at e^+e^- colliders*, *Phys. Rev. D* **96** (2017) 055012 [[arXiv:1703.00402](#)] [[INSPIRE](#)].
- [17] S. Biswas, E. Gabrielli, M. Heikinheimo and B. Mele, *Searching for massless Dark Photons at the LHC via Higgs boson production*, *PoS EPS-HEP2017* (2017) 315 [[INSPIRE](#)].
- [18] H. Beauchesne and C.-W. Chiang, *Measuring properties of a dark photon from semi-invisible decay of the Higgs boson*, *JHEP* **04** (2022) 127 [[arXiv:2201.04658](#)] [[INSPIRE](#)].
- [19] H. Beauchesne and C.-W. Chiang, *Is the Decay of the Higgs Boson to a Photon and a Dark Photon Currently Observable at the LHC?*, *Phys. Rev. Lett.* **130** (2023) 141801 [[arXiv:2205.10976](#)] [[INSPIRE](#)].
- [20] H. Beauchesne and C.-W. Chiang, *Observability of the Higgs boson decay to a photon and a dark photon*, *Phys. Rev. D* **108** (2023) 015018 [[arXiv:2304.04165](#)] [[INSPIRE](#)].
- [21] S. Biswas, E. Gabrielli and B. Mele, *Dark Photon Searches via Higgs Boson Production at the LHC and Beyond*, *Symmetry* **14** (2022) 1522 [[arXiv:2206.05297](#)] [[INSPIRE](#)].
- [22] V.A. Kuzmin, V.A. Rubakov and M.E. Shaposhnikov, *On the Anomalous Electroweak Baryon Number Nonconservation in the Early Universe*, *Phys. Lett. B* **155** (1985) 36 [[INSPIRE](#)].
- [23] A.G. Cohen, D.B. Kaplan and A.E. Nelson, *Baryogenesis at the weak phase transition*, *Nucl. Phys. B* **349** (1991) 727 [[INSPIRE](#)].
- [24] I. Banta et al., *Non-decoupling new particles*, *JHEP* **02** (2022) 029 [[arXiv:2110.02967](#)] [[INSPIRE](#)].
- [25] G. Passarino and M.J.G. Veltman, *One Loop Corrections for e^+e^- Annihilation Into $\mu^+\mu^-$ in the Weinberg Model*, *Nucl. Phys. B* **160** (1979) 151 [[INSPIRE](#)].
- [26] H.H. Patel, *Package-X: A Mathematica package for the analytic calculation of one-loop integrals*, *Comput. Phys. Commun.* **197** (2015) 276 [[arXiv:1503.01469](#)] [[INSPIRE](#)].
- [27] LHC HIGGS CROSS SECTION WORKING GROUP collaboration, *Handbook of LHC Higgs Cross Sections: 3. Higgs Properties*, [arXiv:1307.1347](#) [[DOI:10.5170/CERN-2013-004](#)] [[INSPIRE](#)].
- [28] CMS collaboration, *A portrait of the Higgs boson by the CMS experiment ten years after the discovery*, *Nature* **607** (2022) 60 [[arXiv:2207.00043](#)] [[INSPIRE](#)].
- [29] ATLAS collaboration, *A detailed map of Higgs boson interactions by the ATLAS experiment ten years after the discovery*, *Nature* **607** (2022) 52 [Erratum *ibid.* **612** (2022) E24] [[arXiv:2207.00092](#)] [[INSPIRE](#)].
- [30] M.E. Peskin and T. Takeuchi, *A new constraint on a strongly interacting Higgs sector*, *Phys. Rev. Lett.* **65** (1990) 964 [[INSPIRE](#)].
- [31] PARTICLE DATA GROUP collaboration, *Review of Particle Physics*, *PTEP* **2022** (2022) 083C01 [[INSPIRE](#)].
- [32] K. Hally, H.E. Logan and T. Pilkington, *Constraints on large scalar multiplets from perturbative unitarity*, *Phys. Rev. D* **85** (2012) 095017 [[arXiv:1202.5073](#)] [[INSPIRE](#)].
- [33] S.R. Coleman and E.J. Weinberg, *Radiative Corrections as the Origin of Spontaneous Symmetry Breaking*, *Phys. Rev. D* **7** (1973) 1888 [[INSPIRE](#)].

- [34] LEPSUSYWG, ALEPH, DELPHI, L3 and OPAL collaborations, *Combined lep chargino results, up to 208 gev for large m_0* ,
http://lepsusy.web.cern.ch/lepsusy/www/inos_moriond01/charginos_pub.html.
- [35] LEPSUSYWG, ALEPH, DELPHI, L3 and OPAL collaborations, *Combined LEP Chargino Results, up to 208 GeV for low DM* ,
http://lepsusy.web.cern.ch/lepsusy/www/inoslowdmsummer02/charginolowdm_pub.html.
- [36] ATLAS collaboration, *Constraints on the Higgs boson self-coupling from single- and double-Higgs production with the ATLAS detector using pp collisions at $s = 13$ TeV*, *Phys. Lett. B* **843** (2023) 137745 [[arXiv:2211.01216](https://arxiv.org/abs/2211.01216)] [[INSPIRE](#)].