

# Two Higgs doublets, effective interactions and a strong first-order electroweak phase transition

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**ABSTRACT:** It is well-known that type II two Higgs doublet models (2HDMs) can struggle to facilitate a strong first-order electroweak phase transition in the early universe whilst remaining theoretically appealing scenarios for many reasons. We analyse this apparent shortfall from the perspective of additional new physics. Starting from a consistent dimension-6 effective field theory Higgs potential extension, we identify the Higgs potential extensions that provide the necessary additional contributions required to achieve a strong first-order electroweak phase transition and trace their phenomenological implications for the Large Hadron Collider. In passing, we critically assess the reliability of the dimension-6 approximation depending on the expected 2HDM phenomenology. In particular, we focus on the role of Higgs pair production (resonant and non-resonant) and interference effects expected in top final states, which are the prime candidates of 2HDM exotics discoveries.

**KEYWORDS:** Higgs Properties, Multi-Higgs Models, Phase Transitions in the Early Universe, Specific BSM Phenomenology

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## 1 Introduction

The null results of searches for new physics beyond the Standard Model (BSM) chiefly performed at the Large Hadron Collider (LHC) have left particle physics in a delicate status quo: the standard paradigms that have shaped BSM model building for the past decades stand challenged, and the role of the TeV scale in nature alongside its microscopic origin are profoundly unclear. This observation is accompanied by insurmountable evidence that new physics is required to reconcile physics at the smallest distances with astrophysical and cosmological observations. The Sakharov criteria [1] provide a strong motivation to incorporate additional sources for CP violation and dynamics responsible for a strong first-order electroweak phase transition to our particle physics picture for efficient baryogenesis. There are various ways to achieve the latter which venture away from minimal SM extensions [2–4]. Nonetheless electroweak baryogenesis remains an attractive avenue and the potential implications for TeV-scale LHC measurements are phenomenologically relevant [5–8].

Along these lines, two Higgs doublet models (2HDMs) remain attractive theories; they have seen continued scrutiny in the literature [9–18]. On the one hand, currently available experimental results are not sensitive enough to move exotic scalar bosons beyond the kinematic reach of the LHC [19]. On the other hand, electroweak precision constraints are avoided similar as in the Standard Model.<sup>1</sup> In particular the 2HDM type II that provides a tangible link to supersymmetric UV completions of the SM is being cornered to provide a strong first-order electroweak phase transition (EWPT) for existing parameter constraints [15, 17, 21, 22]. As 2HDM dynamics alone could not be quite enough to furnish

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<sup>1</sup>It should be noted that the observed tension in the Kaon sector [17] and the recently experimentally hardened  $(g - 2)_\mu$  observation [17, 20] remain difficult to explain in the minimal implementations of the 2HDM.

$O_6^{111111}$	$(\Phi_1^\dagger \Phi_1)^3$	$O_6^{222222}$	$(\Phi_2^\dagger \Phi_2)^3$
$O_6^{111122}$	$(\Phi_1^\dagger \Phi_1)^2 (\Phi_2^\dagger \Phi_2)$	$O_6^{112222}$	$(\Phi_1^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2)^2$
$O_6^{122111}$	$(\Phi_1^\dagger \Phi_2) (\Phi_2^\dagger \Phi_1) (\Phi_1^\dagger \Phi_1)$	$O_6^{122122}$	$(\Phi_1^\dagger \Phi_2) (\Phi_2^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2)$
$O_6^{121211}$	$(\Phi_1^\dagger \Phi_2)^2 (\Phi_1^\dagger \Phi_1) + \text{h.c.}$	$O_6^{121222}$	$(\Phi_1^\dagger \Phi_2)^2 (\Phi_2^\dagger \Phi_2) + \text{h.c.}$

**Table 1.** Dimension-6 operators of class  $\Phi^6$  involving  $\Phi_1$  and  $\Phi_2$ .

a strong first-order EWPT, it is the purpose of this paper to clarify the extra dynamics that are required for the 2HDM to provide a sufficiently large EWPT for electroweak baryogenesis. Concretely, we approach this by means of effective field theory (see also [23–25]) and focus in this work on extensions of the scalar potential of the softly broken  $\mathbb{Z}_2$ -symmetric and CP-conserving 2HDM as a well-motivated sector to facilitate a strong first-order EWPT [26–31]. We will focus on the 2HDM type II in this work, but as we will focus mostly on the implications for multi-Higgs production and phenomenological prospects for multi-top final states, our findings generalise to the 2HDM type I straightforwardly.

We organise this work as follows: in section 2 we review the basics of the 2HDM alongside the effective field theory (EFT) modifications we consider in this work. Section 3 provides a short overview of our computational methods. Section 4 is devoted to our results: we provide scans of operators to achieve a strong first-order EWPT and clarify the correlated phenomenological implications relevant for the LHC in multi-Higgs and multi-top final states. We summarise and conclude in section 5.

## 2 2HDMs and dimension-6 Higgs potential extensions

The tree-level dimension-4 potential of the 2HDM is given by [32, 33]

$$\begin{aligned}
 V_{\text{tree}}(\Phi_1, \Phi_2) = & m_{11}^2 (\Phi_1^\dagger \Phi_1) + m_{22}^2 (\Phi_2^\dagger \Phi_2) - m_{12}^2 (\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1) + \lambda_1 (\Phi_1^\dagger \Phi_1)^2 + \lambda_2 (\Phi_2^\dagger \Phi_2)^2 \\
 & + \lambda_3 (\Phi_1^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2) + \lambda_4 (\Phi_1^\dagger \Phi_2) (\Phi_2^\dagger \Phi_1) + \frac{1}{2} \lambda_5 \left[ (\Phi_1^\dagger \Phi_2)^2 + (\Phi_2^\dagger \Phi_1)^2 \right] \\
 & + \left( \lambda_6 (\Phi_1^\dagger \Phi_1) + \lambda_7 (\Phi_2^\dagger \Phi_2) \right) (\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1)
 \end{aligned} \tag{2.1}$$

where  $\Phi_{1,2}$  are  $SU(2)_L$  doublets with hypercharge  $Y = 1$ . The absence of tree-level flavour-changing neutral interactions can be guaranteed by imposing a  $\mathbb{Z}_2$  symmetry [34], which is softly broken by the term proportional  $m_{12}^2$ . In the following, we will assume only such soft breaking and choose the couplings  $\lambda_{6,7} = 0$ , which induces a hard breaking of  $\mathbb{Z}_2$ , to be zero. We furthermore take the values of the remaining coupling and mass parameters,  $\lambda_i$  ( $i = 1, \dots, 5$ ) and  $m_{ab}^2$  ( $a, b = 1, 2$ ), to be real.

Including higher-dimensional EFT contributions to the Higgs potential leads to the operators of table 1. The  $\Phi^6$  operators<sup>2</sup> that we focus on in this first investigation give rise

<sup>2</sup>The dimension-6 operators are classified following the Warsaw basis convention [35].

to a dimension-6 extension of the 2HDM potential [36–39]

$$\mathcal{L}_{\text{EFT}} = \mathcal{L}_{\text{2HDM}} + \sum_i \frac{C_6^i}{\Lambda^2} O_6^i \quad \Rightarrow \quad V_{\text{dim-6}} = - \sum_i \frac{C_6^i}{\Lambda^2} O_6^i. \quad (2.2)$$

Here,  $O_6^i$  are the dimension-6 operators given in table 1 and  $C_6^i$  are the corresponding Wilson Coefficients (WCs). As in the unperturbed 2HDM we can solve the tadpole equations to relate  $m_{11}^2$  and  $m_{22}^2$  to the remaining potential parameters in the vacuum. Diagonalisation of the charged and CP-odd Higgs mass mixing matrices is described by the characteristic angle

$$\tan \beta = \frac{v_2}{v_1}, \quad (2.3)$$

which means that even in the presence of our  $\Phi^6$  interactions  $\tan \beta$  is directly linked to the ratio of vacuum expectation values, while  $(246 \text{ GeV})^2 \approx v^2 = v_1^2 + v_2^2$  is fixed by the  $W$  boson mass (or equivalently the Fermi constant). Explicit expressions for the masses of the neutral and charged Higgs bosons can be obtained similar as in [36], however, it is convenient for us to choose masses and mixing angles as input parameters. Taking inspiration from the on-shell renormalisation scheme introduced in [15], we perform shifts of the renormalisable part of the Lagrangian,  $\lambda_i \rightarrow \lambda_i + \delta\lambda_i$  and  $m_{12}^2 \rightarrow m_{12}^2 + \delta m_{12}^2$ , with

$$\begin{aligned} \delta\lambda_1^{\text{d6}} &= \frac{1}{4\Lambda^2 v_1^2} \left[ 6C_6^{111111} v_1^4 + (2C_6^{121211} + C_6^{122111}) v_1^2 v_2^2 \right. \\ &\quad \left. - \left\{ 2(C_6^{112222} + C_6^{121222}) + C_6^{122122} \right\} v_2^4 \right], \\ \delta\lambda_2^{\text{d6}} &= - \frac{1}{4\Lambda^2 v_2^2} \left[ \left\{ 2(C_6^{111122} + C_6^{121211}) + C_6^{122111} \right\} v_1^4 \right. \\ &\quad \left. - (2C_6^{121222} + C_6^{122122}) v_1^2 v_2^2 - 6C_6^{222222} v_2^4 \right], \\ \delta\lambda_4^{\text{d6}} &= \frac{v_1^2}{\Lambda^2} (C_6^{111122} + C_6^{121211} + C_6^{122111}) \\ &\quad + \frac{v_2^2}{\Lambda^2} (C_6^{112222} + C_6^{121222} + C_6^{122122}), \\ \delta\lambda_5^{\text{d6}} &= \frac{1}{2\Lambda^2} \left[ (2C_6^{111122} + 4C_6^{121211} + C_6^{122111}) v_1^2 \right. \\ &\quad \left. + (2C_6^{112222} + 4C_6^{121222} + C_6^{122122}) v_2^2 \right], \\ \delta m_{12}^{2\text{d6}} &= \frac{v_1 v_2}{2\Lambda^2} \left[ \left\{ 2(C_6^{111122} + C_6^{121211}) + C_6^{122111} \right\} v_1^2 \right. \\ &\quad \left. + \left\{ 2(C_6^{112222} + C_6^{121222}) + C_6^{122122} \right\} v_2^2 \right]. \end{aligned} \quad (2.4)$$

They directly yield mass eigenvalues and mixing angles as in the  $d = 4$  2HDM.<sup>3</sup> In the following, we will refer to  $h$  and  $H$  as the lighter and the heavier CP-even Higgs boson,

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<sup>3</sup>A similar consistent choice could be obtained by a different subset of the potential parameters in the dimension-4 Lagrangian.

Model	I	II
$\xi_h^u$	$\cos \alpha / \sin \beta$	$\cos \alpha / \sin \beta$
$\xi_h^d$	$\cos \alpha / \sin \beta$	$-\sin \alpha / \cos \beta$
$\xi_H^u$	$\sin \alpha / \sin \beta$	$\sin \alpha / \sin \beta$
$\xi_H^d$	$\sin \alpha / \sin \beta$	$\cos \alpha / \cos \beta$
$\xi_A^u$	$\cot \beta$	$\cot \beta$
$\xi_A^d$	$\cot \beta$	$-\tan \beta$

**Table 2.** Coupling modifiers  $\xi$  for 2HDM type I and II and up- and down-type quarks relevant for this study.

respectively.  $A$  denotes the CP-odd scalar and  $H^+$  is the charged scalar degree of freedom. This means that for any choice of Wilson coefficients we obtain the same mass spectrum and neutral mixing angle in the vacuum as for eq. (2.1). Any direct coupling of the Higgs bosons to SM matter is therefore insensitive to the Wilson coefficients and the choice of eq. (2.4) shifts correlations into Higgs self-couplings and multi-Higgs final states, given single Higgs measurements as the transparent input that is typically provided by experimental collaborations and checked for in parameter scans with BSM tools such as **ScannerS** [40–42], **HiggsBounds** [43–46], or **HiggsSignals** [47, 48].<sup>4</sup>

Consequently, the usual classification of the 2HDM according to the  $\mathbb{Z}_2$  assignments applies to this work as well. The Higgs boson couplings to fermions  $f$  in the mass basis are given by

$$\mathcal{L}_{\text{Yuk}} = - \sum_{f=u,d,\ell} \frac{m_f}{v} \left( \xi_h^f \bar{f} f h + \xi_H^f \bar{f} f H - i \xi_A^f \bar{f} \gamma_5 f A \right) - \left[ \frac{\sqrt{2} V_{ud}}{v} \bar{u} \left( m_d \xi_A^d P_R - m_u \xi_A^u P_L \right) d H^+ + \frac{\sqrt{2}}{v} m_\ell \xi_A^\ell \left( \bar{\nu} P_R \ell \right) H^+ + \text{h.c.} \right], \quad (2.5)$$

where  $P_{L,R}$  are the left and right chirality projectors and the coupling modifiers  $\xi$  are given in table 2.

### 3 Finite temperature and phase transitions

In order to calculate the strength of the EWPT in our model, we determine the global minimum of the one-loop corrected effective potential at finite temperature. The derivation of the effective potential is reviewed in section 3.1, our scan is described in section 3.2.

#### 3.1 Review of calculational methods

The effective potential describes the exact vacuum state [49] of a theory including finite temperature effects [50–52]. It can be derived through a perturbative expansion of the

<sup>4</sup>The dimension-4 parameters can then be obtained by inverting  $\delta\lambda_i, \delta m_{12}^2$  in the  $\Lambda^{-1}$  expansion.

generating functional of one-particle irreducible Green's functions. Consequently, the one-loop contribution includes the inverse propagator, independent of the structure of the underlying Lagrangian. In terms of the static field configurations described by  $\vec{\omega}$ , and the temperature  $T$ , the one-loop contribution to the effective potential at finite temperature has the general form

$$V_{\text{eff}}^{(1)}(\vec{\omega}, T) = \sum_{X=S,G,F} (-1)^{2s_X} (1 + 2s_X) I^X, \quad (3.1)$$

with scalar ( $S$ ), gauge-boson ( $G$ ) and fermion ( $F$ ) contributions that have the following form

$$I^S = \frac{T}{2} \sum_n^{\text{Bos}} \int \frac{d^3k}{(2\pi)^3} \sum_i \left[ \log \det \left( -\mathcal{D}_{S,i}^{-1} \right) \right], \quad (3.2a)$$

$$I^G = \frac{T}{2} \sum_n^{\text{Bos}} \int \frac{d^3k}{(2\pi)^3} \sum_i \left[ \log \det \left( -\mathcal{D}_{GB,i}^{-1} \right) \right], \quad (3.2b)$$

$$I^F = -T \sum_n^{\text{Ferm}} \int \frac{d^3k}{(2\pi)^3} \sum_i \left[ \log \det \left( -\mathcal{D}_{F,i}^{-1} \right) \right], \quad (3.2c)$$

with inverse propagators  $\mathcal{D}^{-1}$  calculated in finite temperature field theory. Within the imaginary time formalism, the propagators receive temperature-dependent corrections, e.g. the inverse scalar propagator in momentum space has the form  $\mathcal{D}_S^{-1} = \omega_n^2 + \omega_k^2$  in terms of the discrete Matsubara modes

$$\omega_n^2 = (2n\pi T)^2, \quad n \in \mathbb{N}_0 \quad (3.3)$$

and

$$\omega_k^2 = \mathbf{k}^2 + m^2, \quad (3.4)$$

where in the case of eq. (3.4), the mass term receives corrections from the EFT operators. At  $T = 0$  this gives rise to the vacuum expectation value  $v \simeq 246$  GeV as described above.

The integrals in eq. (3.2) split into a UV-divergent temperature-independent part and a UV-finite, but IR-divergent temperature-dependent part. In the  $\overline{\text{MS}}$ -scheme, the one-loop contributions read

$$I_{\overline{\text{MS}}}^X = \underbrace{\frac{m_X^4}{64\pi^2} \left[ \log \left( \frac{m_X^2}{\mu^2} \right) - k_X \right]}_{\equiv V_{\text{CW}}(\vec{\omega})} + \underbrace{\frac{T^4}{2\pi^2} J_{\pm} \left( \frac{m_X^2}{T^2} \right)}_{\equiv V_T(\vec{\omega}, T)} \quad (3.5)$$

with  $X = \{(S), (G), (F)\}$  and the renormalisation constant  $k_X$

$$k_X = \begin{cases} \frac{5}{6}, & \text{for gauge bosons} \\ \frac{3}{2}, & \text{otherwise} \end{cases} \quad (3.6)$$

and the *thermal* fermionic (+) and bosonic (−) function  $J_{\pm}$  [50, 52, 53]

$$J_{\pm}(x^2) = \int_0^{\infty} dk k^2 \log(1 \pm e^{-\sqrt{k^2+x^2}}). \quad (3.7)$$

The bosonic Matsubara zero modes,  $n = 0$ , lead to IR divergences that are cancelled by resumming the thermal masses  $\Pi$  [51, 53–57]. The scalar thermal masses are calculated as thermal scalar self-energy corrections in the soft-momentum limit,

$$\begin{aligned} \Pi_{ij}^{(1)}(\mathbf{p} \rightarrow 0, \omega_n \rightarrow 0) &\equiv \Pi_{ij}^{(1)}(0) \\ &= \sum_k \kappa_{ij}^k T \sum_n \int \frac{d^3p}{(2\pi)^3} \mathcal{D}_{kk}(\omega_n, \omega_p) \\ &\quad + \sum_{k,l} \kappa_{ij}^{kl} T^2 \sum_{n,m} \int \frac{d^3p_1}{(2\pi)^3} \mathcal{D}_{kk}(\omega_n, \omega_{p_1}) \int \frac{d^3p_2}{(2\pi)^3} \mathcal{D}_{ll}(\omega_m, \omega_{p_2}). \end{aligned} \quad (3.8)$$

Here, quartic scalar couplings are labelled with  $\kappa_{ij}^k$ , while couplings between six scalars are encoded in  $\kappa_{ij}^{kl}$ . Note, that the latter is already a two-loop correction to the scalar self-energy. The finite temperature field theory integral is evaluated in the high-temperature limit  $m/T \ll 1$  as

$$\begin{aligned} &T \sum_n \int \frac{d^3p}{(2\pi)^3} \mathcal{D}_{ij}(\omega_n, \omega_p) \\ &= T \sum_n \int \frac{d^3p}{(2\pi)^3} \frac{1}{\omega_n^2 + \omega_p^2} \\ &= T \sum_n \int \frac{d^3p}{(2\pi)^3} \frac{1}{\omega_n^2 + \mathbf{p}^2 + m^2} \simeq T \sum_n \int \frac{d^3p}{(2\pi)^3} \frac{1}{\omega_n^2 + \mathbf{p}^2} = \dots = \frac{T^2}{12} \left[ 1 + \mathcal{O}\left(\frac{m}{T}\right) \right]. \end{aligned} \quad (3.9)$$

The dimension-6 operators generate 2-loop contributions  $\sim T^4$  (see e.g. [28]) which we include throughout our calculation. These derive straightforwardly from the six-point interaction vertices of table 1. Applying the Arnold-Espinosa method [57],<sup>5</sup> the thermal potential  $V_T(\vec{\omega}, T)$  is replaced as

$$V_T(\vec{\omega}, T) \rightarrow V_T(\vec{\omega}, T) + V_{\text{daisy}}(\vec{\omega}, T), \quad (3.10)$$

where

$$V_{\text{daisy}}(\vec{\omega}, T) = -\frac{T}{12\pi} \left[ \sum_{i=1}^{n_{\text{Higgs}}} \left( (\bar{m}_i^2)^{3/2} - (m_i^2)^{3/2} \right) + \sum_a^{n_{\text{gauge}}} \left( (\bar{m}_a^2)^{3/2} - (m_a^2)^{3/2} \right) \right], \quad (3.11)$$

with  $n_{\text{Higgs}}$  denoting the number of real Higgs fields and  $n_{\text{gauge}}$  the number of gauge bosons in the adjoint representation of the gauge group. The  $\bar{m}$  denote the thermal masses that include the thermal corrections  $\Pi^{(1)}$ .

<sup>5</sup>For further remarks on this approach and how it compares to the Parwani approach, cf. [54, 57]. Further discussions and comparisons are given in [58, 59].

As stated above, we introduce shifts to the parameters of the Higgs potential to absorb the effect of the dimension-6 operators such that the masses and mixing angles that we use as input parameters remain unchanged. We also require the one-loop corrections to leave the masses and mixing angles at their tree-level values. For this we introduce additional finite counterterms summarised in the counterterm potential  $V^{\text{CT}}$  and apply the renormalisation conditions

$$0 = \partial_{\phi_i} \left( V^{\text{CW}} + V^{\text{CT}} \Big|_{\vec{\omega}=\vec{\omega}_{\text{tree}}} \right) \tag{3.12}$$

$$0 = \partial_{\phi_i} \partial_{\phi_j} \left( V^{\text{CW}} + V^{\text{CT}} \Big|_{\vec{\omega}=\vec{\omega}_{\text{tree}}} \right), \tag{3.13}$$

where  $\phi_i$  denote the scalar Higgs doublet fields developing a non-zero VEV  $\bar{\omega}_i$ , with  $\bar{\omega}_{\text{tree},i}$  being the corresponding tree-level VEV. Note, that for the construction of our counterterm potential, we choose the free parameters arising in the derivation of the potential, *cf.* [15], such that the dimension-6 terms do not introduce new coupling structures. The dimension-6 operator contributions enter through the partial derivatives of  $V_{\text{CW}}$  which, however, are modified w.r.t. the 2HDM through the dimension-6 contributions to the mass terms. Note, that the dimension-6 operators do not introduce additional new counterterm structures in the above two conditions. For further details on the renormalisation procedure, we refer to [60, 61].

With the dimension-6 extended tree-level potential  $V_{\text{tree,dim-6}}$ ,

$$V_{\text{tree,dim-6}} \equiv V_{\text{tree}} + V_{\text{dim-6}} \tag{3.14}$$

in terms of the tree-level potential  $V_{\text{tree}}$  of eq. (2.1) and the dimension-6 potential  $V_{\text{dim-6}}$  of eq. (2.2), we hence have for the loop-corrected effective potential at finite temperature as function of the classical field configuration  $\vec{\omega}$ ,

$$V(\vec{\omega}, T) = V_{\text{tree,dim-6}}(\vec{\omega}) + V_{\text{CW}}(\vec{\omega}) + V_{\text{CT}}(\vec{\omega}) + V_T(\vec{\omega}, T), \tag{3.15}$$

which we have implemented in the C++ code `BSMPT` [60, 61].

For an electroweak phase transition to be of strong first order, the ratio of the critical VEV  $v_c$  at the critical temperature  $T_c$  has to fulfil the (conventionally chosen) criterion  $\xi_c \equiv v_c/T_c > 1$  to avoid too large baryon washout. The critical temperature is defined as the temperature where there exist two degenerate global minima, one at  $v = 0$  and the other at the critical VEV  $v_c \neq 0$ . The values of  $T_c$ ,  $v_c$  and hence  $\xi_c$  are obtained from `BSMPT` which computes the vacuum expectation value  $v(T)$  at a given temperature  $T$  through the minimisation of the effective potential  $V(\vec{\omega}, T)$ . The value  $v$  is obtained as

$$v(T) = \left( \sum_{k=1}^{n_H} \bar{\omega}_k^2 \right)^{\frac{1}{2}}, \tag{3.16}$$

where  $n_H$  means that the sum is performed over all field directions in which we allow for the development of a non-zero electroweak VEV, which are given by the fields that couple to the electroweak gauge bosons. The  $\bar{\omega}_k$  denote the field configurations that minimise the loop-corrected effective potential  $V(\vec{\omega}, T)$ .



$m_h$ [GeV]	$m_H$ [GeV]	$m_A$ [GeV]	$m_{H^\pm}$ [GeV]	$\tan \beta$	$c_{HVV}$	$m_{12}^2$ [GeV <sup>2</sup> ]
2HDM I/II (light)						
125.09	130...3000	30...3000	85/800...3000	0.8...30	-0.3...1.0	$10^{-5}...10^7$

**Table 3.** Scan ranges of the 2HDM input parameters, where light refers to the set-up where the lighter of the two CP-even neutral Higgs bosons is the SM-like Higgs  $H_{\text{SM}}$ , i.e.  $h \equiv H_{\text{SM}}$ .

### 3.2 Scanning methodology

For our numerical analysis, we use parameter points that are compatible with all relevant theoretical and experimental constraints. We resort here to a parameter sample that has been generated recently for the 2HDM [62]. The scan was performed with the help of the program `ScannerS` [40–42]. `ScannerS` chooses as scan parameters the masses of the 2HDM Higgs bosons,  $\tan \beta$ , the soft breaking mass term  $m_{12}^2$  and the coupling  $c_{HVV}$  of the heavier Higgs boson to massive gauge bosons  $V \equiv W^\pm, Z$  (instead of the mixing angle  $\alpha$  that diagonalizes the neutral CP-even mass matrix). We performed an additional scan to select parameter points with significant branching ratios  $\text{BR}(H \rightarrow hh)$  of  $H$  into a pair of lighter Higgs bosons  $h$ . The total scan ranges of the two merged parameter samples are listed in table 3 for the scenario where the lighter of the two neutral CP-even Higgs bosons,  $h$ , takes the role of the SM-like Higgs, denoted as  $H_{\text{SM}}$  in the following.<sup>6</sup> We restrict ourselves to the type I and II models. `ScannerS` checks for the theoretical constraints, requiring that the potential is bounded from below, that perturbative unitarity holds and that the electroweak vacuum is the global minimum. For the latter it uses the discriminant from [63].

On the experimental side, we impose compatibility with the electroweak precision data and demand the computed  $S$ ,  $T$  and  $U$  values to be within  $2\sigma$  of the SM fit [64], taking into account the full correlation among the three parameters. One of the neutral CP-even Higgs bosons, in our study chosen to be  $h$ , is required to have a mass of [65]

$$m_{H_{\text{SM}}} = 125.09 \text{ GeV}, \tag{3.17}$$

and to behave SM-like. `ScannerS` checks for compatibility with the Higgs signal data through the link to `HiggsSignals` version 2.6.1 [47]. Scenarios with interfering Higgs signals are excluded by forcing the non-SM-like Higgs bosons to lie outside an interval of 5 GeV around 125 GeV in order to avoid interference effects that require a dedicated thorough study beyond the focus of this work. We require 95% C.L. exclusion limits on non-observed scalar states by using `HiggsBounds` version 5.9.0 [43–46]. The sample is also checked with respect to the recent ATLAS analyses in the  $ZZ$  [66] and  $\gamma\gamma$  [67] final states that were not yet included in `HiggsBounds`. Flavour constraints are taken into account by testing for compatibility with  $\mathcal{R}_b$  [68, 69] and  $B \rightarrow X_s \gamma$  [69–74] in the  $m_{H^\pm} - \tan \beta$  plane. We imposed in the 2HDM type II the latest bound on the charged Higgs mass given in [74],  $m_{H^\pm} \geq 800 \text{ GeV}$ , for essentially all values of  $\tan \beta$ . In the type I models, this bound is much weaker and is strongly correlated with  $\tan \beta$ .

<sup>6</sup>While also scenarios with  $H \equiv H_{\text{SM}}$  still comply with all applied constraints, the plots shown in the following have  $h \equiv H_{\text{SM}}$ .

$m_h$ [GeV]	$m_H$ [GeV]	$m_A$ [GeV]	$m_{H^\pm}$ [GeV]	$\tan \beta$	$c_{HVV}$	$m_{12}^2$ [GeV <sup>2</sup> ]
125.09	681	855	884	1.362	-0.00459	220945
		$T_c^{\text{d4}}$ [GeV]	$v(T_c)^{\text{d4}}$ [GeV]	$\xi_c^{\text{d4}}$		
		250.55	226.76	0.91		

**Table 4.** Input parameters of the benchmark point used for figure 1.

The inclusion of the Wilson coefficients in the 2HDM potential modifies the 2HDM mass values and mixing angles. By applying the shifts given in eqs. (2.4) we ensure, however, that they remain unchanged also after inclusion of the dimension-6 contributions. This allows us to use the `ScannerS` sample of parameter points and work with parameter sets that are compatible with all relevant theoretical and experimental constraints. In practice, we apply within `BSMPT` on the parameters  $\lambda_{1,2,4,5}$  and  $m_{12}^2$  of the respective `ScannerS` sample point the shifts of eqs. (2.4). The values of  $m_{11}^2$  and  $m_{22}^2$  are obtained from the minimisation conditions taking into account the dimension-6 contributions to the Higgs potential. With these parameters and  $\lambda_3$  for the given parameter point under investigation, `BSMPT` then computes the EWPT for the 2HDM potential including the dimension-6 operators.

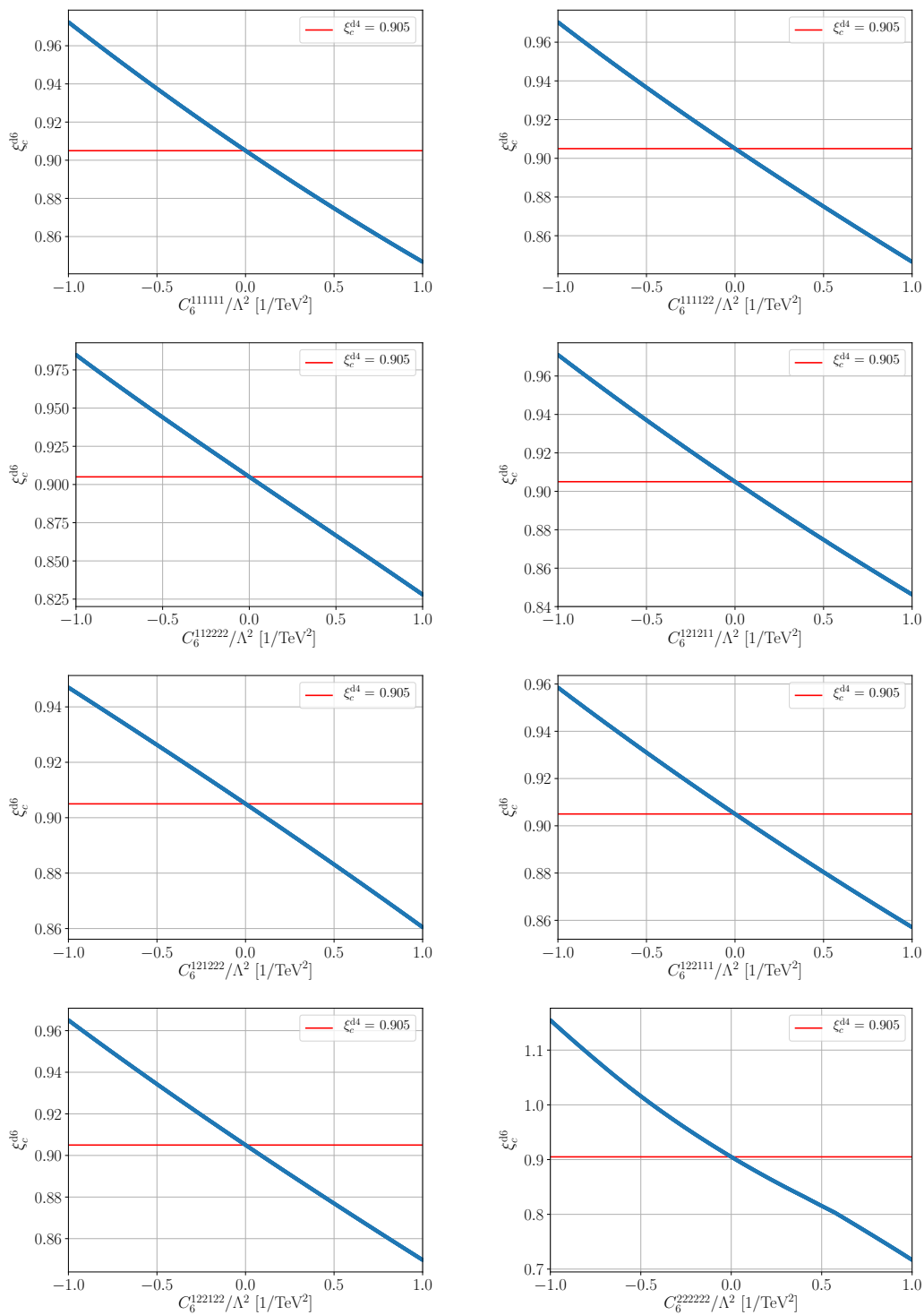
#### 4 Phenomenological aspects of effective 2HDM phase transitions

The general requirements for EFT methods to provide appropriate approximations for momentum-independent Wilson coefficients are twofold. Firstly, the heavy degrees of freedom that are integrated out need to be sufficiently heavy compared to the characteristic energy scale that is probed at the respective laboratory. And secondly, perturbativity, which constitutes the overarching concept of the field-theoretic aspects investigated in this work imposes the additional requirement of dimension-6 terms to be a small correction in relation to the renormalisable dimension-4 result. This provides confidence in neglecting higher order terms in the  $\Lambda^{-1}$  expansion.

The modifications of the 2HDM introduced in section 2 provide a rich landscape for phenomenological deviations. This is particularly interesting for the 2HDM type II that we will mostly focus on in this section. However, as we are mainly considering interactions of the extra Higgs bosons with the top sector, our findings generalise to the 2HDM type I, where a strong first-order electroweak phase transition (SFOEWPT) already at dimension-4 level can be found more easily than in type II.

On the one hand, correlations of masses and couplings are modified away from the dimension-4 expectation when considering effective interactions. On the other hand, such deviations are still allowed to be significant as the LHC has so far only shed limited light on the structure of the Higgs self-interactions. The choice of input parameters enables us to directly choose  $\alpha$ ,<sup>7</sup>  $\tan \beta$  and the Higgs boson masses as relevant input parameters of this study. This has the benefit that electroweak precision constraints and Higgs signal strengths

<sup>7</sup>In terms of  $c_{HVV}$  in the `ScannerS` scan.



**Figure 1.** Representative behaviour of  $\xi_c^{d6}$  for a representative parameter point of the 2HDM type II of table 4 with  $\xi_c^{d4} \simeq 0.9$  when the impact of the individual Wilson coefficients is considered. .

are largely unchanged from the dimension-4 result;<sup>8</sup> the modifications of the dimension-6 interactions of section 2 are then primarily visible in modified Higgs self-interactions (e.g. Higgs pair production). Consistency of single Higgs collider physics observables with the SM follows closely the usual type I or II paradigms (see [17, 19, 62] for recent analyses).

We first turn to the impact of individual  $\Phi^6$  operators of table 1 and how they can contribute to a first-order phase transition.<sup>9</sup> Reflecting the fact that these additional interactions should be small compared to the dimension-4 theory, in first instance, we consider points that show a relatively strong phase transition  $\xi_c^{d4} \simeq 0.9$  with the measurements as reflected in `HiggsBounds`, `HiggsSignals`, via `ScannerS`. As can be seen in figure 1, the effect of additional contributions to the potential creates to good approximation a linear dependence  $\sim C_6^i$ , which demonstrates the robustness of the approach, i.e. the inclusion of non-linear parts in the  $\Lambda^{-1}$  expansion via the Debye masses and the BSMPT approach is numerically insignificant. Furthermore, figure 1 clearly shows that an SFOEWPT can be achieved in the 2HDM type II when considering new effective contributions to the Higgs potential at moderate Wilson coefficient sizes in agreement with current experimental constraints. A short remark concerning the perturbativity and unitarity of the EFT extension is in order. The operators considered in this work do not lead to a kinematic enhancement of scalar amplitudes as they are manifestly momentum-independent. This means that the additional contribution to the dominant zeroth partial wave analysis of the scalar amplitudes is parametrically given by

$$\delta a_{d6}^0 \sim \frac{C_6^i}{32\pi} v^2 / \Lambda^2$$

for high energies compared to the particle thresholds (the  $\mathcal{O}(1)$  factor is determined by  $\tan\beta$  and  $\alpha$  and identical particle factors contributing to the amplitude, see e.g. [75, 76]). This leaves only loose constraints on the Wilson coefficients  $|C^i|/\Lambda^2 < \mathcal{O}(1) \times 32\pi/\text{TeV}^2$ .

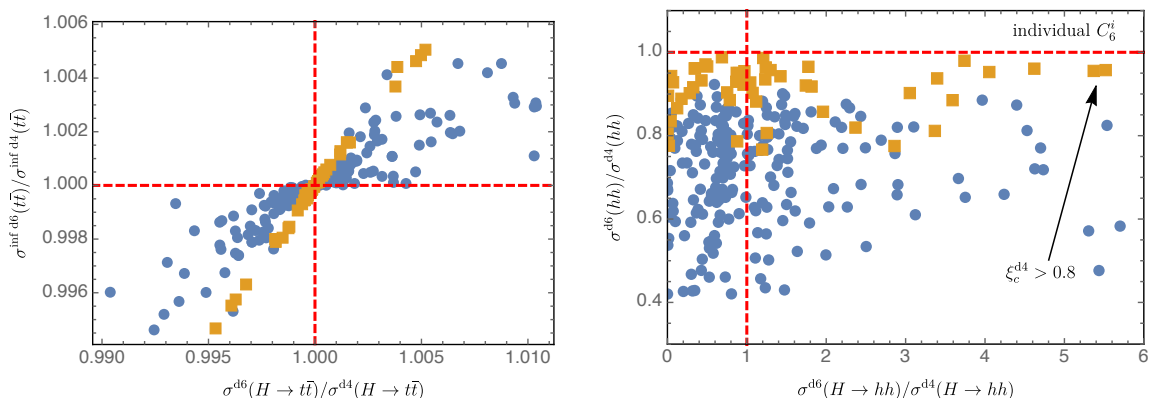
The arguably most interesting question now becomes whether the presence of such operators has collider-relevant implications. The latter come in a range of different guises; the phenomenology of the extra Higgs bosons is dominated by top quark final states, in particular when we are relatively close to the alignment limit which is favoured by current collider observations [19]. It is well-known that these are subject to large interference effects that can render the narrow-width approximation unreliable and could even lead to a vanishing direct sensitivity for the naively best-motivated LHC signatures [77] (see also [78–85]). The LHC experiments include these interference effects to their 2HDM searches, e.g. [86, 87]. In  $gg \rightarrow H \rightarrow t\bar{t}$  the exotic Higgs width is the crucial parameter<sup>10</sup> and it is therefore worthwhile to understand the correlation of  $\xi_c$  with the Higgs width feeding into  $H \rightarrow t\bar{t}$  searches. Along these lines, the potential lineshape analysis of a future discovery  $H \rightarrow t\bar{t}$  could serve as an indirect measurement of the phase transition in the

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<sup>8</sup>We explicitly check that the modified charged Higgs contributions do not impact the SM-like Higgs decay into  $\gamma\gamma$ .

<sup>9</sup>As stated above, for all results that we present we take parameter scenarios where the lighter of the two CP-even Higgs bosons,  $h$ , is the SM-like one.

<sup>10</sup>Theoretically this is apparent from the requirement to evaluate the signal component of the process at the complex Higgs pole [88, 89] to guarantee gauge-independence as a consequence of the Nielsen identities [90].



(a) Modification of  $gg \rightarrow H \rightarrow t\bar{t}$  and interference effects with continuum  $gg \rightarrow t\bar{t}$ . (b) Modification of Higgs pair production  $gg \rightarrow hh$  and its correlation with the resonance contribution  $gg \rightarrow H \rightarrow hh$ .

**Figure 2.** Correlation between EFT-extended cross sections and their dimension-4 2HDM counterparts. We scan over parameters that are allowed by the constraints detailed in section 3.2, and identify individual Wilson coefficients to achieve  $\xi_c^{d6} = 1$ . We consider points with  $\xi_c^{d4} > 0.3$  (blue points) and highlight  $\xi_c^{d4} > 0.8$  for comparison (orange squares). As can be seen the dimension-6 modifications are much smaller for the  $\xi_c^{d4} > 0.8$  points as the phenomenologically relevant  $\sigma^{d6}(hh)/\sigma^{d4}(hh)$  approaches unity.

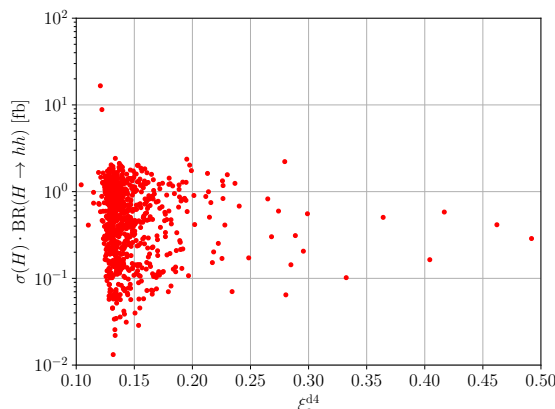
favoured  $t\bar{t}$  channel. To this end, we define the interference cross section between gluon fusion  $gg \rightarrow H \rightarrow t\bar{t}$  signal and QCD  $gg \rightarrow t\bar{t}$  continuum as

$$d\sigma^{\text{inf}} \sim 2\text{Re} \{ \mathcal{M}(gg \rightarrow H \rightarrow t\bar{t}) \mathcal{M}^*(gg \rightarrow t\bar{t}) \} \tag{4.1}$$

where  $\mathcal{M}$  denotes the amplitude (we have suppressed identical phase space and parton density factors and work to leading order accuracy in the following).

In figure 2 we show a parameter sample of points in agreement with the constraints described in section 3.2, for  $\xi_c^{d4} > 0.3$  with  $\xi_c^{d6} \simeq 1$  through choices for single Wilson coefficients (we will study the effect of combined Wilson coefficients below). We do not distinguish between the individual Wilson coefficients as the phenomenological outcome is qualitatively similar. The scan also includes relatively large Wilson coefficient choices which are necessary to achieve  $\xi_c^{d6} \simeq 1$  starting from  $\xi_c^{d4} \simeq 0.3$ ; for illustration purposes we highlight smaller dimension-6 couplings resulting from  $\xi_c^{d4} \geq 0.8$  in figure 2. For these latter points the modification of the electroweak phenomenology is smaller than for points far away from  $\xi_c \sim 1$ . Consequently, the di-Higgs production cross section receives a smaller modification than for points with  $0.3 \leq \xi_c^{d4} \leq 0.8$  (highlighted as blue dots in figure 2).

The phenomenological baseline of the  $d = 4$  points shown in figure 2 is a top-philic one;  $t\bar{t}$  final states are the preferred decay channels of the exotic Higgs bosons with typically  $\text{BR}(H \rightarrow t\bar{t}) \gtrsim 0.8$ . The changes that are introduced by the dimension-6 interactions do not (and to be perturbatively robust must not) change this behaviour dramatically. In fact, neither the  $t\bar{t}$  final states, nor their width-sensitive interference effects show phenomenologically observable modifications, Figure 2(a). There is a trend that reflects



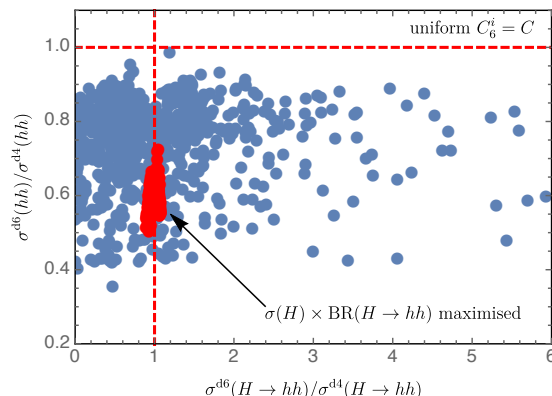
**Figure 3.** Correlation of the 2HDM value of  $\xi_c^{d4}$  and the exotic Higgs production cross section in the double Higgs decay channel.

the overall  $\xi_c$  behaviour, i.e. the closer  $\xi_c^{d4}$  gets to unity, the smaller the  $gg \rightarrow H \rightarrow t\bar{t}$  modification becomes as a result of a smaller modification of the total  $H$  decay width. In any case for the generic top-dominated final states, such per mille level effects are well beyond the sensitivity that can be obtained at hadron colliders.

This leaves multi-Higgs final states as motivated signatures as shown in figure 2 (b). The resonant  $H \rightarrow hh$  contribution is small as  $H \rightarrow t\bar{t}$  is preferred, but there can be a modification of the resonance signal  $gg \rightarrow H \rightarrow hh$ , which is correlated with a modified trilinear  $Hhh$  coupling. However, the overall  $gg \rightarrow hh$  rate is decreased. For instance we find deviations of 125 GeV Higgs boson pair production of  $\sigma^{d6}(hh)/\sigma^{d4}(hh) \simeq 0.4$  (0.8) for  $\xi_c^{d4} = 0.3$  (0.9) when sampling individual Wilson coefficient directions. For large distances  $1 - \xi_c^{d4}$  it is clear that the EFT contribution needs to overcome the 2HDM contribution alone, which eventually will put pressure on the dimension-6 EFT assumption, highlighted through non-linear dependencies of  $\xi_c^{d6}(\{C_6^i\})$ . The individual Wilson coefficient scans that we have focussed so far remain in their linear regime and hence robust when viewed according to this criterion (this quickly changes for correlated Wilson coefficient choices, see below).

The behaviour of the Higgs pair production cross sections in figure 2 (b) is mostly due to the fact that additional potential contributions lead to an enhancement of the trilinear SM-like Higgs self-coupling  $\sim \lambda_{hhh}$  at the order of  $\sim 50\%$ . For these coupling deviations the dominant  $gg \rightarrow hh$  contribution shows a decreasing behaviour with a rescaling of the light Higgs self-interaction  $\kappa_\lambda > 1$  [91–98]. In the light of existing projections of Higgs boson pair production [99, 100] this can be a manageable albeit challenging task at the LHC.<sup>11</sup> The analysis of the separated resonant  $H \rightarrow hh$  signal in direct comparison to the  $hh$  continuum production can therefore serve as an indirect constraint on  $\xi_c \sim 1$ . Given that discovery of the  $H$  state should become possible in  $H \rightarrow t\bar{t}$  first, there is significant scope of data-driven methods in separating continuum from on-shell  $H$  production.

<sup>11</sup>The  $t\bar{t}hh$  final state showing an increasing cross section for  $\kappa_\lambda > 1$  could provide additional sensitivity [101].



**Figure 4.** Similar to figure 2(b) but with uniformly chosen Wilson coefficients  $C_6^i = C$  and scanned over  $C$  to achieve  $\xi_c^{d6} \simeq 1$ . We highlight the Higgs-philic scan result points of figure 3 in red. The large deviation of  $\sigma^{d6}(hh)/\sigma^{d4}(hh)$  signifies that large Higgs potential modifications are required to achieve  $\xi_c^{d6} \simeq 1$  for such points.

So far all of our results have been dominated by the general top-philic nature of the exotic Higgs bosons in the 2HDM. Moving to larger dimension-4 couplings in the Higgs sector we can furnish situations where the branching ratio of  $H \rightarrow hh$  is significant whilst maintaining reasonable production rates via virtual top quarks. In such an instance the correlation of on-shell production and di-Higgs continuum is less statistically limited and therefore experimentally more feasible. As can be seen in figure 3, such parameter points in the 2HDM type II are typically characterised by a larger distance  $|1 - \xi_c^{d4}|$ . Achieving  $\xi_c > 1$  in a controlled way therefore relies on the interplay of different effective operators as can indeed be expected in concrete UV scenarios (e.g. in singlet extensions of the scalar sector [22, 30, 102–105]). In figure 4, we show the results of a scan of uniform Wilson coefficients  $C_6^i = C$  to achieve  $\xi_c^{d6} \simeq 1$ , again for  $\xi_c^{d4} \geq 0.3$ . Furthermore, we also include the Higgs-philic points of figure 3 to figure 4. These are typically characterised by relatively low  $\xi_c^{d4}$  — the price of Higgs-philic  $H$  phenomenology. As can be seen, in general our findings are similar to the individual Wilson coefficient scan, with significant enhancements possible in the  $H \rightarrow hh$  rate. As this starts from a relatively low cross section rate for top-philic  $H$  decays, the largest enhancements  $\sigma^{d6}(H \rightarrow hh)/\sigma^{d4}(H \rightarrow hh) > 3$  arise from small  $H \rightarrow hh$  dimension-4 cross sections. In this instance, a large enhancement is not directly phenomenologically relevant as the cross section still remains small when including  $d = 6$  contributions. Yet, enhancements of factors of  $\sim 2.5$  are possible for cross sections in the fb range and we can therefore anticipate some LHC sensitivity here in the  $b\bar{b}b\bar{b}$  [106–110] and  $b\bar{b}\tau\tau$  channels [94, 111–116]. When turning to points that have a larger  $H \rightarrow hh$  probability (highlighted in figure 4 in red), the resonance contribution is modified at the 5–10% level, while the continuum receives a 50% modification. We note that for uniform Wilson coefficients squared dimension-6 terms  $\sim C_6^i C_6^j / \Lambda^4$  will induce a non-linear behaviour thus highlighting the importance of a full (matching) calculation to obtain more realistic estimates. This is particularly relevant for parameter points with sizeable Higgs-philic branching ratios given in figure 3.

## 5 Summary and conclusions

The SM alone is known to provide an insufficiently strong first-order phase transition for electroweak baryogenesis. On the one hand, this could mean that baryogenesis proceeds through mechanisms not associated with the TeV scale. On the other hand, this could indicate additional BSM physics close to the TeV scale with LHC-relevant implications. Well-motivated SM extensions such as the 2HDM become increasingly under pressure to provide a sufficiently large SFOEWPT which we take as motivation to understand and address  $|1 - \xi_c^{d4}|$  in terms of additional dynamics that facilitate a strong first-order electroweak phase transition as a minimal modification of the 2HDM. Reverting to effective field theory techniques for the 2HDM, we consider modifications of the Higgs potential at dimension-6 level for the  $\mathbb{Z}_2$ -symmetric, CP-conserving 2HDM. While the 2HDM type II typically falls short of an SFOEWPT, the distance  $|1 - \xi_c^{d4}|$  can be overcome by EFT contributions to the Higgs sector.

Additional dimension-6 dynamics that push the 2HDM over the SFOEWPT finishing line (according to the  $\xi_c \simeq 1$  criterion) can then lead to phenomenological consequences for LHC physics. Interference effects of heavy Higgs production in the top final state are width-dependent and therefore sensitive to EFT modifications. The overall effect, however, for the top-philic final states currently preferred by experimental data through the alignment limit renders these effects too small to be measurable at the LHC.

Higgs pair production is an important tool for fingerprinting an SFOEWPT, and the distance  $|1 - \xi_c^{d4}|$  is directly correlated with expected Higgs pair production deviation. Current extrapolations of Higgs pair production to the 3/ab high-luminosity frontier indicate that the LHC should become sensitive enough to partially explore this region (see also the recent [117]), potentially assisted by discoveries in the  $t\bar{t}$  channels. For more strongly-coupled Higgs interactions of the renormalisable 2HDM, the  $H \rightarrow hh$  signature is more motivated as a signature for 2HDM discovery. In such an instance, the 2HDM type II is not capable of producing an SFOEWPT and a significant modification of the 2HDM Higgs potential is required. While this stretches the reliability of the dimension-6 approximation, there are phenomenologically relevant implications, predominantly the reduction of the  $gg \rightarrow hh$  rate and modifications of  $gg \rightarrow H \rightarrow hh$ . The LHC is capable of exploring both phenomenological arenas to some extent and the discovery of an additional Higgs boson that follows a 2HDM paradigm could therefore be analysed from an SFOEWPT dimension-6 Higgs-EFT angle.

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