

Deconstructing resonant Higgs boson pair production at the LHC: Effects of colored and neutral scalars in the NMSSM test case

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We study resonant production of pairs of Standard Model (SM)-like Higgs bosons, in the presence of new neutral Higgs states together with new colored scalars (stops or sbottoms) in loops within the Next-to-Minimal Supersymmetric Standard Model. This is used as a test case to prove that the Large Hadron Collider has sensitivity to a variety of effects stemming from interferences between resonant (heavy) Higgs diagrams and/or among these and nonresonant topologies involving loops of both tops and stops. These effects can alter significantly the naive description of individual s -channel Breit-Wigner resonances, leading to distortions of the latter, which, on the one hand, may mask their presence but, on the other hand, could enable one to extract features of the underlying new physics scenario. This last aspect is made possible through a decomposition of the $gg \rightarrow hh$ signal process into all its amplitude components, each of which has a well-defined coupling structure. Ultimately, such effects can be traced back to the relevant Feynman diagrams and can enable a detailed interpretation of this process. To illustrate this, we introduce various benchmark points that exhibit potentially observable features during the current and/or upcoming runs of the LHC in one or more of the three customary di-Higgs decay channels: $b\bar{b}b\bar{b}$, $b\bar{b}\tau^+\tau^-$, and $b\bar{b}\gamma\gamma$.

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I. INTRODUCTION

Within the Standard Model (SM), Higgs pair, or di-Higgs, (hh) production is at leading order (LO), a one-loop process with destructive interference between the box and the triangle diagrams [1,2], which thus makes it particularly sensitive to Beyond the Standard Model (BSM) physics effects, which enter at the same perturbative order. The largest of such effects are modifications to the trilinear Higgs self-coupling [3–8], the possibility of new colored particles in the loops [9–12] or the existence of new Higgs bosons, which can give an s -channel resonant contribution to di-Higgs production [13–16]. In principle, the cross section can also be modified by a deviation of the top quark Yukawa coupling from its Standard Model (SM) value, but,

in practice, such effects are already somewhat constrained as the $t\bar{t}h$ production cross section does not show a large deviation from the SM prediction [17–19].

The phenomenological importance of SM Higgs pair production at the LHC is primarily related to the ability of accessing the aforementioned trilinear Higgs self-coupling, hhh , and modifications thereof (or indeed additional effects emerging in BSM scenarios). This is of paramount importance, as it gives insights into the shape of the Higgs potential, thereby enabling one to potentially pinpoint the underlying Higgs mechanism of spontaneous symmetry breaking. The latter ought to be embedded in some viable theory of the electroweak (EW) sector, by which one intends a theoretical construct able to remedy the so-called hierarchy problem of the SM. A prevalent theory in this connection is supersymmetry (SUSY).

In a previous work [11], we looked at nonresonant SM-like di-Higgs production in the context of SUSY. In this work, we concentrate instead on the aforementioned s -channel resonant contributions and their interferences with other diagrams that contribute to di-Higgs production in the same theoretical framework. To gain better understanding of all effects, we use the deconstruction approach we previously used to describe squark effects in nonresonant SM-like di-Higgs production [11], now extended to

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include the corresponding resonant case. Interference effects are known to be potentially important in various processes, with recent studies also including di-Higgs production [20–22]. Our approach allows us to show the interference patterns clearly, and, in addition, we shall discuss what kind of interpretations can be done within specific models if an interference pattern is observed.

We study specifically the Nonminimal Supersymmetric SM (NMSSM) [23] (see Ref. [24] for a review of alternative realizations of it) for a number of reasons. Concerning the aforementioned hierarchy (or fine-tuning) problem of the SM, SUSY is a well-motivated framework for BSM physics in relation to the Higgs sector as its symmetry between fermions and bosons produces a cancellation of loop contributions to the Higgs mass making a light Higgs boson natural, thereby elegantly solving this flaw of the SM. In fact, another appealing feature of SUSY is that the Higgs potential is therein generated dynamically (instead, in the SM, it is put in by hand). Concerning model realizations of SUSY, it should be noted that the parameter space of the MSSM, the simplest realization of it, is constrained when it comes to resonant di-Higgs production as a 125 GeV Higgs state requires a large value of $\tan\beta$. At large $\tan\beta$, the constraints from single Higgs boson searches force the heavier Higgs states to be beyond a TeV [25,26], including the heavy CP -even one, H , which will then suppress the corresponding signal $gg \rightarrow H \rightarrow hh$. The NMSSM has both a singlet Higgs field, for which the constraints are weaker, and the possibility of having a 125 GeV Higgs at lower values of $\tan\beta$ [27], which is also a less constrained region. Finally, a feature of SUSY models is that their Higgs self-couplings have less freedom than in the non-SUSY scenario, which will help in the interpretation of di-Higgs results, as discussed below.

Higgs pair production in the NMSSM has been studied before. References [28,29] tackled the case where the Higgs pairs H_1H_1 , H_1H_2 , and H_2H_2 were produced, when H_2 is the SM-like Higgs boson (the aforementioned h state) and H_1 is dominantly a singlet state. In Refs. [30–32], the authors studied Higgs pair production with both mass orderings but only at the total cross section level without specifying resonant and nonresonant contributions. The nonresonant, resonant, and interference contributions were addressed in Ref. [33], but the interference patterns were not studied therein. Here, we will concentrate on the case where the 125 GeV Higgs boson pairs will be produced and our key addition to existing studies is to deconstruct the BSM effects to the differential cross section and to show what are the interference patterns between the different contributions.

We stress, however, that the NMSSM scenarios addressed here are only some of the possible examples which can be studied with the deconstruction method. With the same dataset we have produced, it is possible to study scenarios involving *any* number of colored scalars [of any representation of $SU(3)$] and neutral scalars. A subset of the Monte Carlo (MC) event samples we used for this

analysis corresponds to those used in Ref. [11], and, since the samples involving the neutral scalars have been modularly added to the dataset, it is possible to select only these ones and study scenarios involving only neutral scalars, again, in any number, thereby allowing for full flexibility within this particle content. (New colored fermions will also be included in due course, e.g., to allow for the effects of vectorlike quarks in compositeness [12].)

The plan of the paper is as follows. In the next section, we explain our deconstruction of the $gg \rightarrow hh$ signal. We then proceed to review triple-Higgs couplings in the NMSSM, following which we will discuss our results. Conclusions will then follow.

II. DECONSTRUCTION OF THE SIGNAL CROSS SECTION

The di-Higgs signal can be deconstructed in a finite number of kinematically independent contributions, which can be simulated individually and recombined *a posteriori* to obtain the full signal. This strategy, adopted also in Ref. [11], allows one to assess the role of each contribution in the determination of the kinematical properties of the final state. At the same time, it shows how specific choices of coupling values weight differently the various amplitude contributions for a given set of kinematically relevant parameters (masses, widths, and spins). In this analysis, we consider new colored and new neutral scalars propagating in the loops at the same time as SM objects.

The starting point is to write the simplified Lagrangians, Eq. (1), which minimally extend the SM with modifications of the trilinear hhh coupling and of the Yukawa couplings of the Higgs with SM top and bottom quarks, alongside the contribution of any number of colored scalars $\tilde{s}_i$ and of new neutral scalars S_I^0 ,

$$\mathcal{L}_M = -(\lambda^{\text{SM}} + \kappa_{hhh})vh^3 - \frac{1}{\sqrt{2}}(y_f^{\text{SM}} + \kappa_{hff})h\bar{f}f, \quad (1a)$$

$$\begin{aligned} \mathcal{L}_{\tilde{s}} = & \sum_i \kappa_{h\tilde{s}\tilde{s}}^{ii} v h \tilde{s}_i^* \tilde{s}_i + \kappa_{hh\tilde{s}\tilde{s}}^{ii} h h \tilde{s}_i^* \tilde{s}_i \\ & + \left(\sum_{i>j} \kappa_{h\tilde{s}\tilde{s}}^{ij} v h \tilde{s}_i^* \tilde{s}_j + \text{H.c.} \right), \end{aligned} \quad (1b)$$

$$\mathcal{L}_S = \sum_I \kappa_{Shh}^I v S_I^0 h h + \kappa_{Sff}^I S_I^0 \bar{f} f, \quad (1c)$$

$$\mathcal{L}_{S\tilde{s}} = \sum_{I,i} \kappa_{S\tilde{s}\tilde{s}}^{Ii} v S_I^0 \tilde{s}_i^* \tilde{s}_i, \quad (1d)$$

where all couplings κ are independent dimensionless parameters [a Higgs vacuum expectation value (VEV) has been factorized for all dimension-3 operators] and we have not included couplings between a neutral scalar and two different colored scalars because this vertex does

TABLE I. Complete list of topologies describing gluon-initiated di-Higgs production with modified SM couplings hhh and htt and propagation of any number of colored scalars ($\tilde{s}_i$) and/or neutral scalars (S_I^0). The topologies are classified according to the different products of new couplings, defined in the Lagrangians of Eq. (1), to which the amplitudes are proportional.

Topology type	Feynman diagrams	Amplitude
1 Modified hhh coupling		$A_i \propto \kappa_{hhh}$
2 One modified hff coupling		$A_i \propto \kappa_{hff}$
3 Modified hhh coupling and modified hff coupling		$A_i \propto \kappa_{hhh}\kappa_{hff}$
4 Two modified hff couplings		$A_i \propto \kappa_{hff}^2$
5 Scalar bubble and triangle with $h\tilde{s}\tilde{s}$ couplings		$A_i \propto \kappa_{h\tilde{s}\tilde{s}}^{ii}$
6 Modified hhh coupling + Scalar bubble and triangle with $h\tilde{s}\tilde{s}$ coupling		$A_i \propto \kappa_{hhh}\kappa_{h\tilde{s}\tilde{s}}^{ii}$
7 Scalar triangle and box with two $h\tilde{s}\tilde{s}$ couplings		$A_i \propto \kappa_{h\tilde{s}\tilde{s}}^{ij} ^2$
8 Scalar bubble and triangle with $hh\tilde{s}\tilde{s}$ coupling		$A_i \propto \kappa_{hh\tilde{s}\tilde{s}}^{ii}$
9 Neutral scalar		$A_i \propto \kappa_{Shh}^I \kappa_{Sff}^I$
10 Neutral scalar + colored scalar		$A_i \propto \kappa_{Shh}^I \kappa_{S\tilde{s}\tilde{s}}^{Ii}$

not appear in the di-Higgs process at leading order. It is important to notice that in this Lagrangian we are working with mass eigenstates and that the colored scalars do not necessarily need to be in the fundamental representation of $SU(3)$. Our analysis will indeed consider squarks, but the method we use can be straightforwardly generalized to other cases, such as scalar sextets, octets, and so on.

The signal and interferences contributions to the cross section σ can be parameterized in terms of a sum of different terms, proportional to unique functions of the couplings and to *reduced* cross sections $\hat{\sigma}$ depending exclusively on the masses of all scalars and on the width of the (potentially resonant) neutral scalar. The topologies into which the signal can be deconstructed are represented

in Table I, where the proportionality of the amplitude with respect to the couplings is indicated. The explicit expressions of all unique contributions to the cross section are provided in Appendix A.

From a practical point of view, we have performed MC simulations corresponding to each term of Eq. (A1) in a grid of masses and, only for the neutral scalar, total widths, such that the total cross sections and differential distributions can be built by summing each individual result with weights given as external numerical inputs. There are multiple advantages of this procedure:

- (1) The event samples can be recycled to account for multiple colored or neutral scalars appearing in different models of new physics. In fact, the numerical

model used for the simulations, available as auxiliary data, contains only the *minimal* number of particles needed to simulate all combinations. In particular, the largest amount of neutral scalars which can appear in any amplitude squared for the di-Higgs final states is 2 (corresponding to interference terms between topologies involving two different scalars): if more neutral scalars are present in a specific theory, the same grid of samples can be used, one just needs to pick the closest masses and total widths and rescale by the appropriate (only coupling-related) weights. For what concerns the squarks, one can immediately realize that the electric charge does not play any role in the determination of the cross sections and that the minimal amount of squarks needed to determine the cross section is 4, corresponding to interferences between amplitudes of topology type 7 in Table I where all squark indices are different.

- (2) Reinterpreting the results for different scenarios does not require new simulations: since event samples are combined according to Eq. (A1), different benchmark points (BPs) characterized by different masses, total widths and/or couplings can be analysed using the same dataset. However, while couplings can be continuously modified as they enter as multiplicative parameters of the linear combination, masses and widths are linked to the grid of values used for the simulations. Interpolation for intermediate values is not trivial, as the kinematical features associated to resonances and threshold effects result in nonlinear behaviors of the differential distributions, but results are reliable for points sufficiently close to the grid. (Interpolation methods are currently being studied to overcome this limitation.)
- (3) Nonlinear features of differential distributions related to the interplay of interferences or threshold effects can be inspected with unprecedented detail by isolating their sources and straightforwardly relating these to the underlying physics. This is especially relevant when the mass scales of such nonlinear effects are similar for different new particles, so determining the shape of distributions and their dependence upon variations of the input parameters would be a difficult task using dedicated simulations, especially if such new particles are many. We stress here that the individual deconstructed samples are not physical *per se*; they become physical only once they have been all combined through Eq. (A1) for a given set of input parameters. However, they provide essential information about their role in the determination of the final result.

The key element to achieve the deconstruction of the signal is to build a numerical model where each new interaction is labeled by a specific string, so that, when

doing the simulations, interactions can be counted at the level of amplitudes squared. A dedicated model was thus built in Universal FeynRules Output (UFO) format [34], and individual coupling orders have been assigned to each BSM interaction of the Lagrangians in Eq. (1) (see Appendix B for more details). MG5_aMC [35,36] was then used, as it allows us to generate process by specifying a given number of couplings with a given coupling order in the amplitude squared. The MC simulations were performed at 13 TeV, using the LO set of the NNPDF3.0 PDFs [37]. Pythia8 [38] was used to take care of Higgs decays, hadronization/fragmentation, and parton shower. The reconstruction of objects in the final state was done using the anti- k_T algorithm [39] with jet cone radius $\Delta R = 0.4$, implemented in FastJet [40]. MadAnalysis 5 [41] was used to generate the numerical distributions for the deconstructed samples, combined *a posteriori* with a dedicated script.

A. Width corrections

As previously described, the deconstruction method relies on combining event datasets on a grid of masses and widths of the new particles. As far as masses are concerned, if the studied benchmark points are close enough to the simulated grid, it is possible identify key features related to the interplay of contributions for different physical scenarios in a reasonable approximation. For the total widths, however, if the narrow-width approximation holds, it is possible to apply corrections for points which are not on the grid, allowing us to accurately determine the size of peaks and the total cross section for more flexible range of input parameters.

We start by the consideration that the *partial* width describing the resonant decay of a neutral scalar to the final di-Higgs state is proportional to the trilinear Shh coupling, used to compute the weights of the corresponding contribution in Eq. (A1): $\Gamma_{S \rightarrow hh} \propto \kappa_{Shh}^2$.

The resonant contribution to the di-Higgs production process can be written as

$$\sigma_{pp \rightarrow S \rightarrow hh} = \sigma_{pp \rightarrow S} \times BR_{S \rightarrow hh} = \sigma_{pp \rightarrow S} \times \frac{\Gamma_{S \rightarrow hh}}{\Gamma_S^{\text{tot}}}, \quad (2)$$

The cross section to produce the neutral scalar, $\sigma_{pp \rightarrow S}$, does not depend on the scalar width, and therefore it must be the same whether we consider the closest point on the widths grid or the BP under study. Under the assumption that the BP masses are close enough to the simulation grid, it follows that

$$\begin{aligned} \sigma_{pp \rightarrow S \rightarrow hh}^{\text{BP}} &= \sigma_{pp \rightarrow S \rightarrow hh}^{\text{sim}} \times \frac{BR_{S \rightarrow hh}^{\text{BP}}}{BR_{S \rightarrow hh}^{\text{sim}}} \simeq \sigma_{pp \rightarrow S \rightarrow hh}^{\text{sim}} \\ &\times \left(\frac{\kappa_{Shh}^{\text{BP}}}{\kappa_{Shh}^{\text{sim}}} \right)^2 \times \frac{\Gamma_S^{\text{tot, sim}}}{\Gamma_S^{\text{tot, BP}}}. \end{aligned} \quad (3)$$

Therefore, to take into account widths which are not on the simulation grid, it is sufficient to do the approximate substitution

$$\kappa_{Shh}^{\text{BP}} \rightarrow \kappa_{Shh}^{\text{BP}} \sqrt{\frac{\Gamma_S^{\text{tot, sim}}}{\Gamma_S^{\text{tot, BP}}}} \quad (4)$$

for each neutral scalar with mass larger than $2m_h$ in the BP, using the closest width on the simulation grid. This substitution becomes less and less accurate as the total widths of the scalar increases. Notice also that this rescaling only affects the *height* of the resonant peak. The *shape* of the peak in our results is still determined by the simulated width of the scalar, so this rescaling will lead to the correct contribution to the cross section, but the shape of the peak is only approximate. In the narrow-width approximation, however, this should be sufficient in most practical situations, considering the limited experimental resolution after reconstructions and selections are applied.

III. NMSSM CASE

The particle content we have at hand, colored and neutral BSM scalars, exists naturally in SUSY models. The MSSM allows only heavy new Higgs bosons, so resonant di-Higgs production will be suppressed. Therefore, we look at the NMSSM, where light singlets can evade the experimental constraints and the parameter space for the heavy doublet states also allows lighter Higgs states.

A. Triple-Higgs couplings in the NMSSM

Resonant di-Higgs production depends on the various triple Higgs couplings, so we shall have a look at them first. The MSSM is defined through the superpotential

$$W_{\text{MSSM}} = y_u \hat{Q} \hat{H}_u \hat{U}^c + y_d \hat{Q} \hat{H}_d \hat{D}^c + y_e \hat{L} \hat{H}_d \hat{E}^c + \mu \hat{H}_u \hat{H}_d. \quad (5)$$

The NMSSM extends the field content with a singlet superfield $\hat{S}$ and generates the μ -term dynamically through

$$V(H_u^0, H_d^0, S) = \frac{1}{2} m_{H_u}^2 (H_u^0)^2 + \frac{1}{2} m_{H_d}^2 (H_d^0)^2 + \frac{1}{2} m_S^2 S^2 + \frac{A_\lambda}{2\sqrt{2}} S H_u^0 H_d^0 + \frac{A_\kappa}{6\sqrt{2}} S^3 + \frac{1}{32} (g^2 + g'^2) ((H_d^0)^2 - (H_u^0)^2)^2 + \frac{1}{8} g^2 (H_d^0 H_u^0)^2 + \frac{\lambda^2}{4} (|H_u^0 H_d^0|^2 + |H_u^0 S|^2 + |H_d^0 S|^2) + \frac{\kappa\lambda}{4} H_u^0 H_d^0 S^2 + \frac{\kappa^2}{4} S^4, \quad (10)$$

S being here the CP -even component of the singlet scalar field having a VEV $\langle S \rangle = v_S$. In addition, there will be loop corrections from the top-stop sector, which give a significant correction to the H_u^4 term, which we denote as

$$V_{\text{loop}} = \delta |H_u^0|^4. \quad (11)$$

the VEV of the scalar component of the singlet field. The superpotential of the NMSSM reads

$$W_{\text{NMSSM}} = W_{\text{MSSM}}(\mu = 0) + \lambda \hat{S} \hat{H}_u \hat{H}_d + \frac{\kappa}{3} \hat{S}^3. \quad (6)$$

The SUSY part of the tree-level scalar potential is CP conserving, and we assume here that also the soft SUSY breaking terms respect CP invariance. In such a case, only the CP -even states contribute to resonant di-Higgs production.

We write the scalar parts of the doublet Higgs superfields as

$$H_u = \begin{pmatrix} H_u^+ \\ \frac{1}{\sqrt{2}} (H_u^0 + iA_u^0) \end{pmatrix}, \quad H_d = \begin{pmatrix} \frac{1}{\sqrt{2}} (H_d^0 + iA_d^0) \\ H_d^- \end{pmatrix}, \quad (7)$$

and we define the VEVs such that $\langle H_{u,d}^0 \rangle = v_{u,d}$.

The scalar potential gets three types of contributions. The superpotential gives the F-term contribution

$$V_F = \sum_i \left| \frac{\partial W}{\partial \phi_i} \right|^2, \quad (8)$$

where the sum is over all superfields and we retain only the scalar part. The coupling of chiral superfields to vector superfields give the D-term contributions, which have the same form as in the MSSM:

$$V_D = \frac{1}{2} g^2 \left(H_d^\dagger \frac{\sigma^\mu}{2} H_d + H_u^\dagger \frac{\sigma^\mu}{2} H_u \right)^2 + \frac{1}{2} \left(\frac{g'}{2} \right)^2 (H_u^\dagger H_u - H_d^\dagger H_d)^2. \quad (9)$$

The third contribution is the soft SUSY breaking terms, which include mass terms for the scalar fields and trilinear scalar interactions, which are proportional to superpotential terms. For a superpotential coupling α , the corresponding trilinear SUSY breaking coupling will be denoted as A_α .

The scalar potential of the CP -even Higgs sector at tree level will be $(H_u^0, H_d^0, \text{ and } S \text{ being real scalar fields})$

The Higgs coupling data are compatible with the SM predictions [42,43]. This suggests that we are close to the so-called alignment limit [44], in which one of the mass eigenstates is aligned with the VEV of the doublet states. Therefore, we shall derive the triple Higgs couplings under the assumption that we are in the alignment limit, which means that the CP -even mass matrix is diagonalized by the angle β for which $\tan\beta = v_u/v_d$. Hence, we can represent the doublet states as the linear combinations

$$h = \sin\beta H_u^0 + \cos\beta H_d^0, \quad (12)$$

$$H = \cos\beta H_u^0 - \sin\beta H_d^0. \quad (13)$$

We shall first neglect the mixing with the singlet state. If the latter is negligible, the trilinear Higgs couplings will be determined solely by the quartic couplings, i.e., gauge couplings and λ . The equations for h and H can be inverted to give

$$H_u^0 = \sin\beta h + \cos\beta H, \quad (14)$$

$$H_d^0 = \cos\beta h - \sin\beta H. \quad (15)$$

When we insert this into the scalar potential, we get the following quartic interactions between the mass eigenstates in the alignment limit:

$$\begin{aligned} V(h, H) = & \frac{\cos^2 2\beta}{32} (g^2 + g'^2) (hhhh - 2hhHH + HHHH) + \frac{g^2 + 2\lambda^2}{32} [\sin^2 2\beta hhhh + 4 \sin 2\beta \cos 2\beta hhhhH \\ & + (4\cos^2 2\beta - 2\sin^2 2\beta) hhHH - 4 \sin 2\beta \cos 2\beta hHHH + \sin^2 2\beta HHHH] \\ & + \delta (\sin^4 \beta hhhh + 4 \sin^3 \beta \cos \beta hhhhH \dots). \end{aligned} \quad (16)$$

In the alignment limit, $\langle h \rangle = v$ and $\langle H \rangle = 0$, so the triple Higgs couplings relevant for di-Higgs production become

$$\lambda_{hhh} = \left(\frac{g^2 + g'^2 \cos^2 2\beta}{8} + \frac{\lambda^2 \sin^2 2\beta}{4} + 4\delta \sin^4 \beta \right) v, \quad (17)$$

$$\lambda_{Hhh} = \left(\frac{3(g^2 + 2\lambda^2)}{16} \sin 4\beta + 12\delta \sin^3 \beta \cos \beta \right) v. \quad (18)$$

Both of these couplings are subject to the constraint that $m_h = 125$ GeV. The first of these couplings governs nonresonant di-Higgs production, whereas the second one gives the strength of the resonant di-Higgs production via the heavy doublet Higgs.

If $\tan\beta$ is large, the term proportional to λ^2 in λ_{hhh} gets suppressed, and the remaining terms lead to λ_{hhh} being close to its SM value as the part of the Higgs potential relevant for Higgs mass generation is of the form $\lambda_{\text{eff}} |H_u^0|^4$, like in the SM. At low $\tan\beta$, the λ^2 term becomes relevant, and, if it is greater than the SM gauge couplings, it can enhance λ_{hhh} above its SM value, a feature well known in the literature [6].

One may immediately observe that λ_{hhh} is positive whereas the sign of λ_{Hhh} is not definite and depends on $\tan\beta$ as well as the relative size of λ and δ . Also, it is easy to observe that at large $\tan\beta$ both terms in λ_{Hhh} go to zero, so the resonant contribution becomes small. While this result was derived in the context of the NMSSM, we notice that

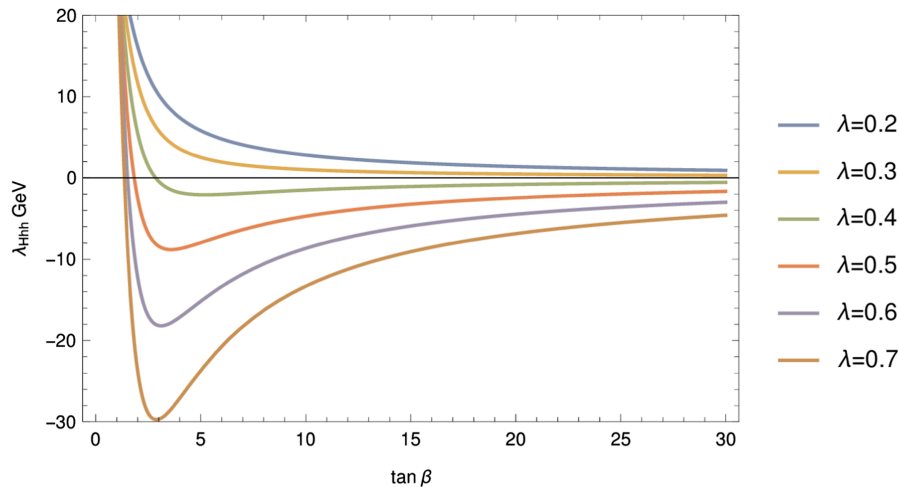


FIG. 1. The schematic dependence of the Hhh triple Higgs coupling on $\tan\beta$ for values of λ between 0.2 (top) and 0.7 (bottom) with a fixed value of $\delta = 0.04$. Large and small values of λ lead to different interference patterns in di-Higgs production.

we have no other Higgs coupling structures in the MSSM. Hence, the general conclusion of small λ_{Hhh} at large $\tan\beta$ applies also to the MSSM. However, since a large value of $\tan\beta$ is required to obtain a 125 GeV Higgs mass in the MSSM (unless the SUSY breaking scale is extremely large), resonant effects are expected to be small, which is a further reason to use the NMSSM as our example instead of the MSSM.

We show schematically the dependence of λ_{Hhh} on $\tan\beta$ in Fig. 1 using $\lambda = 0.2 \dots 0.7$ and $\delta = 0.04$. If we limit ourselves to $\tan\beta > 1$, for large λ , the coupling is negative except for a small range of values close to $\tan\beta = 1$. Close to $\tan\beta = 3$, the coupling gets its most negative value, and at higher values of $\tan\beta$, it tends to zero (as mentioned above). If λ is small, the second term will be larger, and the coupling is positive at all $\tan\beta$, which can be seen from Fig. 1. Therefore, the sign of this coupling and consequently the interference pattern of resonant di-Higgs production with the di-Higgs continuum will depend on whether the 125 GeV Higgs mass is achieved through tree-level superpotential contributions or stop loops.

We then look at the CP -even singlet state in resonant di-Higgs production. In practice, there will be a small mixing between the singlet and doublet states through the λ term. Even in the absence of this mixing, there is a chance for resonant di-Higgs production via the singlet as the scalar potential contains a term $\frac{1}{2} y_t \lambda \tilde{t}_L \tilde{t}_R S H_d^0$, which will induce a non-negligible coupling between stops and the singlet. Unfortunately, even if there is a relatively light stop and stop mixing is large, the resonant cross section is very small, as we show below through some examples. Therefore, the singlet states are in practice produced through top loops resulting from the mixing between the singlet and doublet states.

The coupling between the singlet and the SM-like Higgs states in the perfect alignment limit is

$$\lambda_{Shh} = \frac{\lambda^2}{2} v_S + \left(\frac{A_\lambda}{4\sqrt{2}} + \frac{\kappa\lambda}{4} v_S \right) \sin 2\beta, \quad (19)$$

where v_S is the VEV of the singlet field and A_λ is the soft SUSY-breaking trilinear coupling between S , H_u , and H_d . If λ is large, the first term is large and positive, while the second term may have either sign. The second term has its largest absolute value at $\tan\beta = 1$ and becomes small at large $\tan\beta$. If there is to be a large resonant signal from the singlet state, λ_{Shh} must be positive, and A_λ should be of the same sign as the other terms. The nonzero mixing with the doublet states lead to some corrections, but if the expression in Eq. (19) is not small, the correction is subleading.

Both singlet and doublet states can give a large contribution to the resonant di-Higgs production cross section in the so-called λ SUSY limit, where λ is large and $\tan\beta$ not much larger than 1. In that case, the signs of the λ_{Hhh} and

λ_{Shh} couplings will be opposite, which has implications for the interference patterns, as we will show below.

B. Resonant di-Higgs signals

We shall now show representative examples of the impact of new Higgs bosons onto the di-Higgs differential cross section. There are overall three classes of deformation of the cross section from the SM: the modification of the Higgs trilinear coupling or top Yukawa (denoted by M in the figures), squark effects (denoted by s in the figures), and effects of additional scalars (denoted by S in the figures). Interferences between classes A and B are denoted by A|B and the class *mix* arises from contributions where more than one type of classes ends up with the same coupling structure. The class Ss denotes diagrams where BSM Higgs states are produced through squark loops. The grouping is done by the combination of couplings, so sometimes contributions are not where they would seem to belong by the naming of the groups.¹ The detailed constraints of each category are given in Appendix A.

In all figures, the black line denotes the nonresonant di-Higgs production in the SM, the red curve (*signal*) the sum of all BSM effects and the blue curve the total differential cross section. We stress that *only the black and blue curves represent physical events* and must therefore be positive definite in each bin. The red curve, representing the contribution of the sum of all signal and interference terms, can have bins with negative event rates; this contribution is unphysical if taken alone, but it is correct as it only makes sense when summed to the irreducible background (black curve) to give the (physical) total number of events of the blue curve. Any curve with negative events indicates destructive interference between the diagrams involved, again unphysical if taken alone, but leads to the correct differential event rate when all contributions are summed. Prominent features of all plots will be the rise of the cross section from $\approx 2m_h$ (where the phase space begins) to slightly above the $\approx 2m_t$ threshold (produced by the top loops), followed by a nearly exponential decrease of the SM cross section, upon which the various BSM effects will install (as described in the forthcoming subsections), the latter having the common feature of a second threshold at $\approx 2m_{\tilde{t}_1}$ (produced by the lightest stop), which may be visible or not depending on the actual value of $m_{\tilde{t}_1}$ and other parameters specific to a given BP.

We show the values of some of the essential model parameters for our BPs in Table II. The spectrum of masses, widths, and couplings were computed with SPheno4.0.4 [45,46], where we modified the SPheno output to include

¹For instance, the diagram 2 squared (which one would think belongs to the category M) is proportional to κ_{hff}^2 as is the interference of diagram 4 with the SM (which one would think belongs to the category M|B). Now, they are both in the category “mix” instead.

TABLE II. The summary of the relevant parameters for resonant di-Higgs production for our BPs. Here, m_H denotes the mass of the heavy (mostly) doublet Higgs, and m_S is the mass of the (mostly) singlet Higgs.

	$\tan\beta$	λ	κ	A_λ/GeV	v_S/GeV	m_H/GeV	m_S/GeV	$m_{\tilde{t}_1}/\text{GeV}$	Peculiar features
BP1	30	0.043	0.04	150	13150	2000	350	1400	Light singlet on the SM top threshold
BP2	1.38	0.69	0.43	-340	1250	800	500	600	Light squark and scalars (all sub-TeV), doublet heavier than singlet
BP3	2.5	0.7	0.54	-345	1210	1200	800	600	Doublet scalar on the $\tilde{t}_1$ threshold
BP4	2.31	0.65	0.68	220	1280	800	1200	600	Light squark and doublet, singlet heavier than doublet
BP5	7.0	0.21	0.16	-550	943	800	100	1400	100 GeV singlet, sub-TeV doublet

the various couplings needed by the UFO model, normalized these, and fixed their signs according to Eq. (1). The SPheno model was generated through Sarah4.14.5 [47].

1. BP1: Light singlet scalar

Our first BP introduces a singlet scalar at $m_S = 350$ GeV. Since the mixing between the singlet and the SM-like Higgs pushes the SM-like Higgs mass down, one needs a relatively small value for λ , a large value for $\tan\beta$, and heavy squarks to achieve a 125 GeV Higgs mass. A large value for $\tan\beta$ forces the Higgs self-coupling to be SM-like. This case is shown in Fig. 2.

The production of the singlet scalar relies here almost completely on the mixing with the SM-like Higgs, i.e., the inherited coupling with the top quark. Other contributions to resonant production are at a per mille level. The sign of the triple Higgs coupling is positive [in the convention of

Eq. (19)], which leads to an interference pattern producing an excess for $m_{hh} < m_S$ and a deficit for $m_{hh} > m_S$ with respect to the SM. The interference effects near the resonant peak are $\sim 30\%$ – 40% of the SM cross section at parton level. The singlet states always have a very small width, with $\Gamma/m \lesssim 0.001$, which limits the size of the interference effects.

2. BP2: Intermediate scale scalars with light squarks

This BP is actually our second NMSSM one from Ref. [11], where we had a 600 GeV squark combined with a large modification of the Higgs trilinear self-coupling. This is achieved through a low value for $\tan\beta$ and a large value for λ . This BP had also a singlet scalar with a mass close to 500 GeV and a doublet scalar with mass close to 800 GeV. Here, we show what happens when you include the effects of the new Higgs states. The

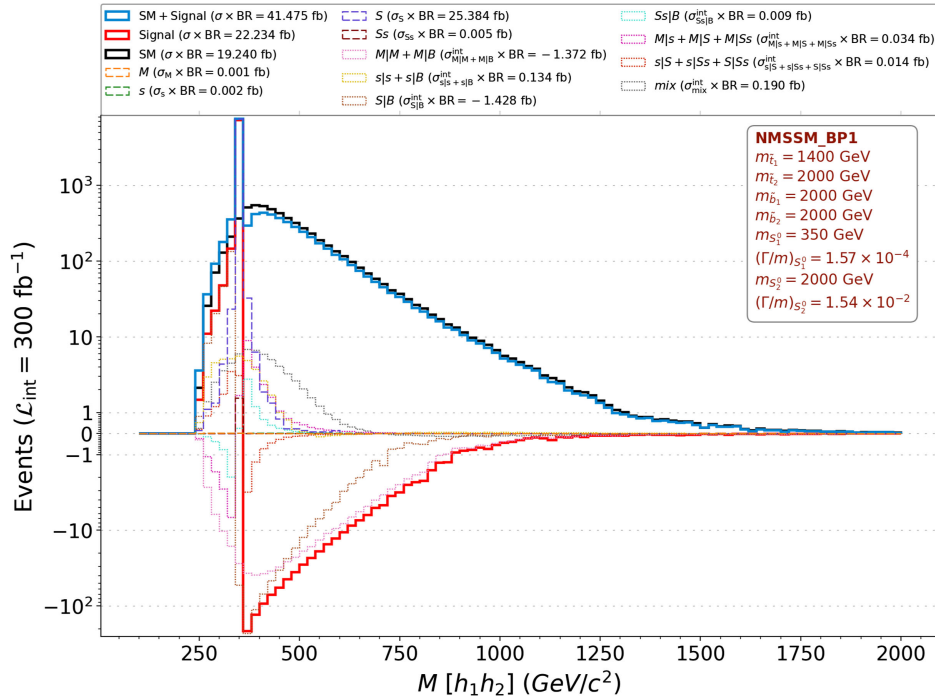


FIG. 2. Differential distribution of the di-Higgs invariant mass for BP1, which features a singlet Higgs with $m_S = 350$ GeV. In this and following figures of this section, the y axis represents the number of events per bin, associated to an integrated luminosity of 300/fb.

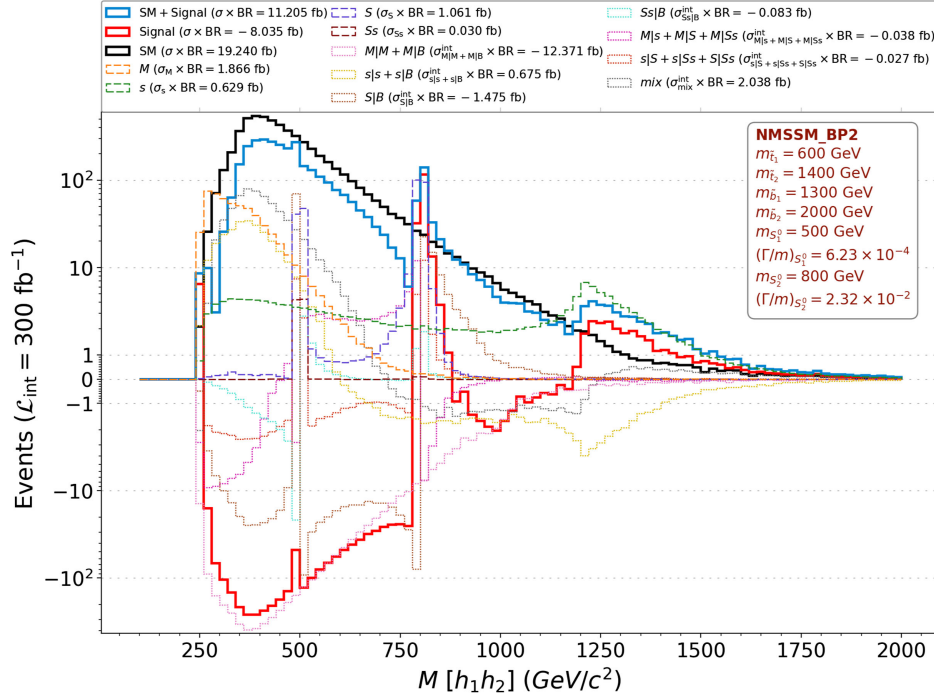


FIG. 3. Differential distribution of the di-Higgs invariant mass for BP2, which features a singlet Higgs with $m_S = 500$ GeV and a doublet Higgs with $m_H = 800$ GeV.

differential cross section for this benchmark is given in Fig. 3.

The singlet Higgs is nearly invisible due to its very small couplings to the SM quarks, but even for this BP, the quark contribution is larger than that from squark loops. The main effect of the singlet state is to assist the doublet Higgs in creating a negative interference with the SM diagrams over the range 500–800 GeV. This negative interference together with the resonant peak at 800 GeV are the main modifications given by the two scalars. For doublet states, we typically have $\Gamma/m \sim 0.01$, which means that the interference effects are larger than for singlet states. The destructive interference for BP2 just below the doublet peak is $\sim 75\%$ of the SM cross section. Interference effects between the squarks and new scalars are small but negative, slightly reducing the squark threshold effect.

There is a local minimum in the parton level cross section at low m_{hh} , which occurs if the triple Higgs coupling λ_{hhh} is well above its SM value. The low event rate in that bin combined with the finite energy/momentum resolution (which leads to migration of events between adjacent bins with larger event rates) make this feature impossible to be observed experimentally.

3. BP3: Doublet scalar on top of the squark threshold

Here, we take a BP with two quite heavy scalars, $m_S = 800$ GeV and $m_H = 1200$ GeV, together with a 600 GeV stop. In principle, the heavy Higgs is right on top of the

squark threshold ($2m_{\tilde{t}_1}$). With these kinds of BPs, there is a complementarity between resonant effects and squark effects. Resonant effects are large if we have a large tree-level contribution to the SM-like Higgs mass (large λ and low $\tan\beta$), while squark effects are large if we have a large loop-level contribution to the Higgs mass (small λ and large $\tan\beta$, including large stop mixing). Our benchmark configuration is an example of the former.

The differential cross section at parton level is given in Fig. 4. In this BP, the tree-level effect to the SM-like Higgs mass dominates, and therefore the behavior at 1200 GeV is dictated by a resonance structure. We see clearly the typical doublet interference pattern for large λ , i.e., destructive interference when $m_{hh} < m_H$ and constructive interference when $m_{hh} > m_H$. Interferences with the squarks and new scalars come with an opposite sign with respect to the quarks but are subdominant.

There is a singlet Higgs at 800 GeV, but in this case, it is completely invisible. In this BP, the production through the lightest stop dominates, but the interference terms with the SM box diagram are larger than the resonant diagram squared. The interferences with the squark and quark diagrams come with opposite signs, which further suppresses the singlet contribution to become completely invisible.

4. BP4: Doublet lighter than singlet

Now, we reverse the roles of the singlet and doublet compared to the previous BP; i.e., we set the doublet state to have $m_H = 800$ GeV and the singlet state to have

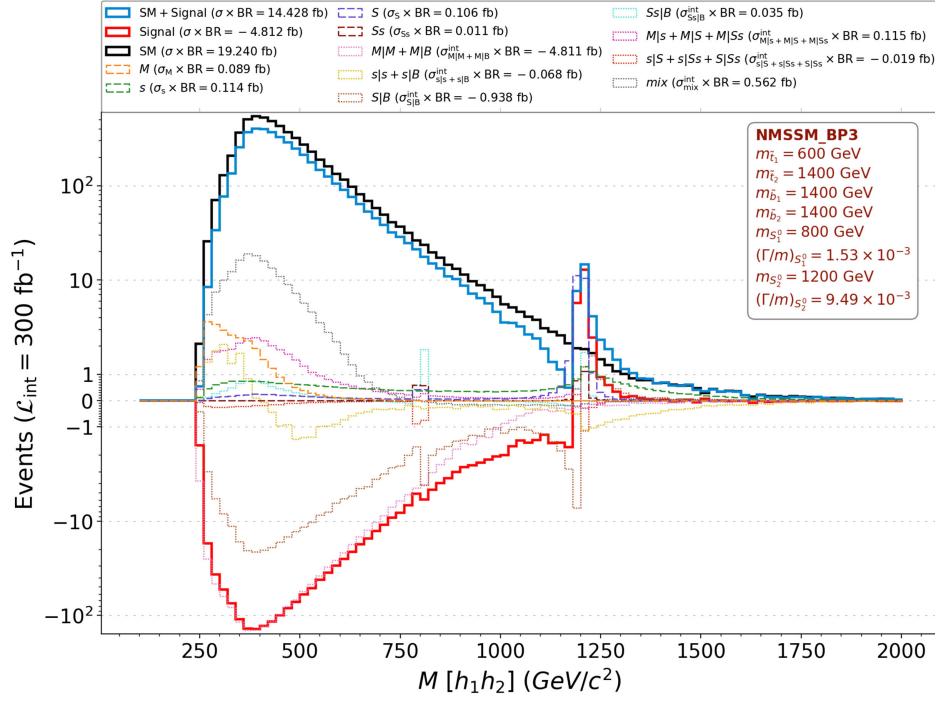


FIG. 4. Differential distribution of the di-Higgs invariant mass for BP3, which features a singlet Higgs with $m_S = 800$ GeV and a doublet Higgs with $m_H = 1200$ GeV.

$m_S = 1200$ GeV. We also have a large value of λ and a small one of $\tan\beta$, which lead to an enhanced value of λ_{hhh} compared to the SM. As in the previous BP, the squark contribution becomes rather small.

In Fig. 5 you can see that both scalars give a resonant peak with their typical interference structures at large values of λ . There is constructive interference in the interval $m_H < m_{hh} < m_S$ and destructive interference over the range

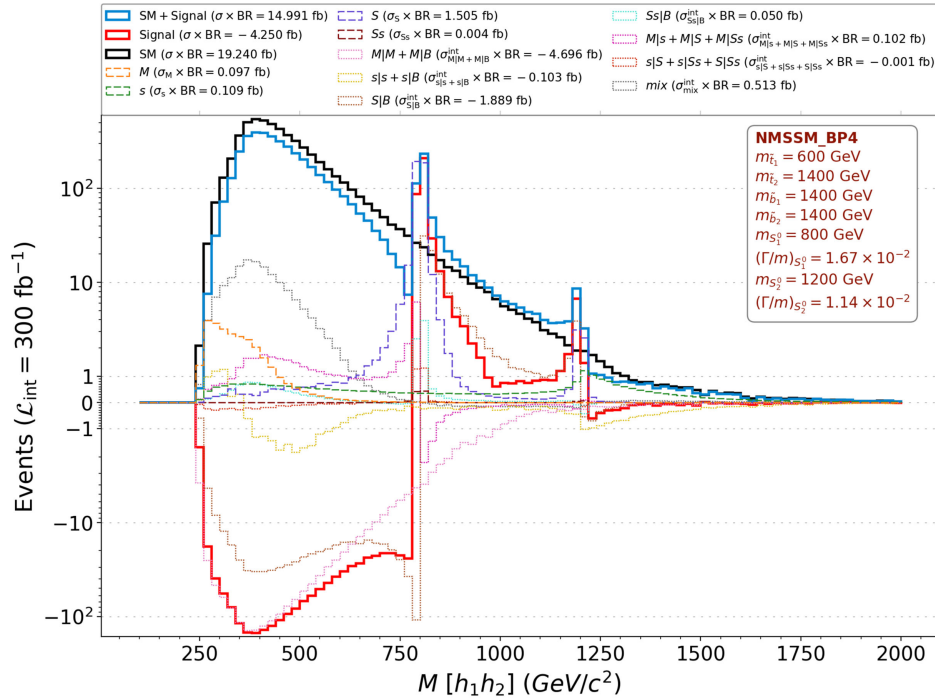


FIG. 5. Differential distribution of the di-Higgs invariant mass for BP4, which features a doublet Higgs with $m_H = 800$ GeV and a singlet Higgs with $m_S = 1200$ GeV.

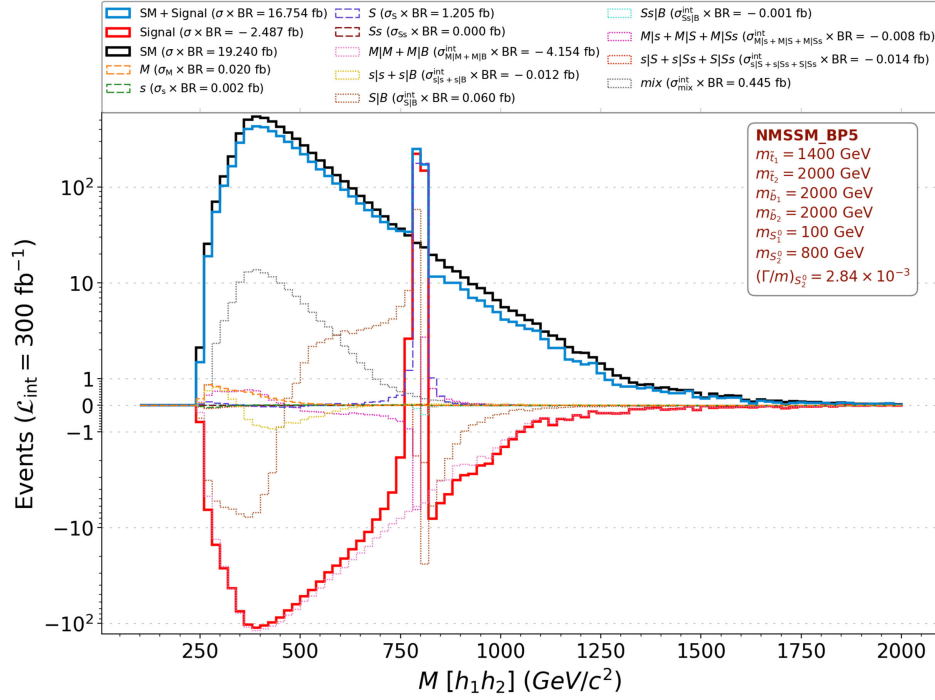


FIG. 6. Differential distribution of the di-Higgs invariant mass for BP5, which features a singlet Higgs state with $m_S = 100$ GeV and a doublet Higgs state with $m_H = 800$ GeV.

$m_{hh} < m_H$ and $m_{hh} > m_S$. The negative interference above $m_{hh} = 1200$ GeV is larger than the squark contribution, so the squark threshold becomes invisible. The destructive interference near the doublet state is the largest one among our BPs, reaching in excess of 75% of the SM cross section at parton level.

Even though the singlet has a mass of $2m_{\tilde{t}}$, the squark diagrams do not contribute significantly to its production; in fact, only about 1% of the resonant singlet contribution is due to squarks.

5. BP5: 100 GeV singlet

In singlet extensions of the Higgs sector, it is possible that the 125 GeV state is not the lightest Higgs boson. In fact, there are small excesses in various channels at LEP [48], ATLAS [49], and CMS [26,50], which could be explained through a new scalar state with mass in the range 95–100 GeV [51–56]. We show what could be the typical impact of such a scalar onto SM-like di-Higgs production.

We present the differential cross section for such a BP in Fig. 6. We have added a 100 GeV mostly singlet scalar together with a 800 GeV doublet. To get a 125 GeV Higgs and a light singlet, we must choose small values for λ and κ , a moderate value for $\tan \beta$, and heavy stops. The small value of λ together with $\tan \beta = 7$ leads to a SM-like hhh coupling.

The light singlet state leads a suppression of the di-Higgs signal at low m_{hh} . This is due to some suppression in the SM-like Higgs coupling to top quarks, which could also be measured in $t\bar{t}h$ production. The singlet state itself has a negative interference with the SM diagrams, which has an impact at low m_{hh} , but this is subdominant compared to the modification of the top Yukawa. These effects together give a reduction in the low m_{hh} region, similar to what we could have if λ_{hhh} were above its SM value. This is again an example of where a simple coupling modifier approach could lead to an incorrect interpretation.

The doublet state at 800 GeV has now an interference pattern with an excess for $m_{hh} < m_H$ and a deficit for

TABLE III. Selection and kinematic cuts for the three final states considered in the analysis.

$b\bar{b}\gamma\gamma$	$b\bar{b}\tau\tau$	$b\bar{b}b\bar{b}$
$N(b) > 1$	$N(b) > 1$	$N(b) > 3$
$N(\gamma) > 1$	$N(\tau) > 1$	...
$p_T(b) > 45(20)$ GeV	$p_T(b) > 45(20)$ GeV	$p_T(b) > 40$ GeV
$ \eta(b) < 2.5$	$ \eta(b) < 2.5$	$ \eta(b) < 2.5$
$ \eta(\gamma) < 2.5$	$ \eta(\tau) < 2.5$	...
$120 \text{ GeV} < M(\gamma\gamma) < 130 \text{ GeV}$	...	...

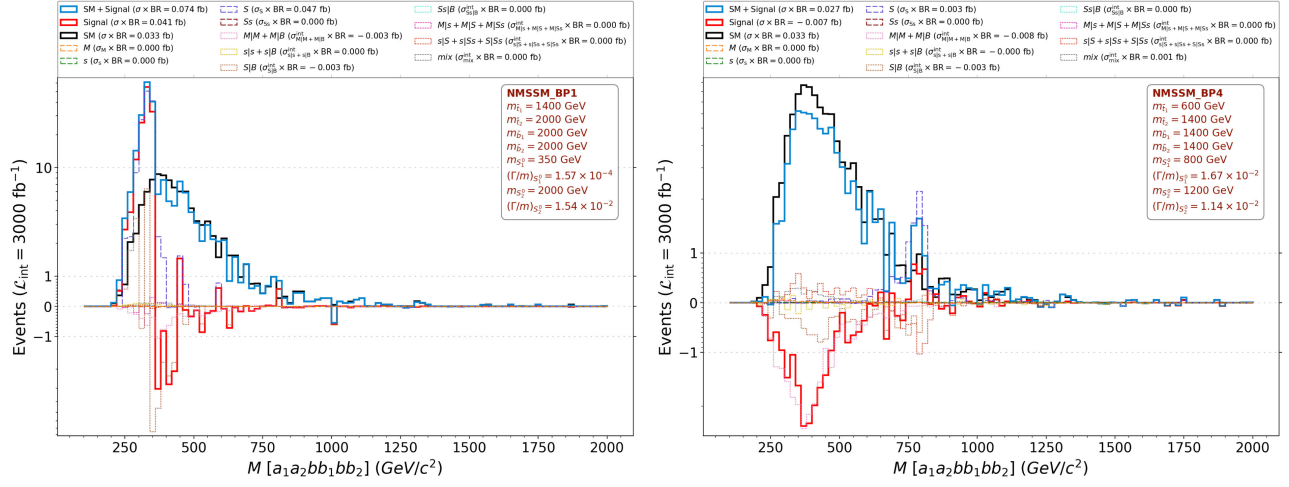


FIG. 7. Differential distribution of the di-Higgs invariant mass in the $bb\gamma\gamma$ channel for BP1 and BP4 at 3000 fb^{-1} .

$m_{hh} > m_H$. This is an example of the Hhh coupling behavior for small λ shown in Fig. 1 and indicates that the SM-like Higgs mass relies largely on loop corrections and not the tree-level enhancement due to a large λ .

C. Prospects in different decay channels

We look at some of the benchmarks in the three most sensitive decay channels $bb\gamma\gamma$, $b\bar{b}\tau\tau$, and $b\bar{b}b\bar{b}$. In Table III, we show the preselection criteria applied before plotting each of the final states. However, note that the preselection only includes kinematical and geometrical criteria and does not include the expected reconstruction efficiency. Hence, the numbers shown in the plots should be multiplied by the reconstruction efficiency. We think it can be useful to remind the reader here that, like in the plots of the previous section, the only physical curves are the solid black and blue ones, representing the irreducible background and the sum of signal and background and their interference, respectively. All the other curves, and noticeably the solid red one, which represents the sum of all signal and interference contributions, do not represent physical observables and can therefore contain bins with a negative number of events: these bins represent a dominant contribution by negative interference. It is also worth noticing that in the tails of some distribution, a negative number of events can be seen in the overall sums of the blue curves, but this clearly is an effect of statistical fluctuations in regions with very few MC events and has no physical meaning.

1. Two b -jets and two photons

The channel with two b -jets and two photons offers a low background and also a good mass resolution. The downside of this channel is the low event rate due to a $\text{BR}(hh \rightarrow b\bar{b}\gamma\gamma) = 0.26\%$ and a reconstruction efficiency of around 10% after preselection. Therefore, there is limited sensitivity to most of our benchmarks in this channel at nominal

High Luminosity phase of the Large Hadron Collider (HL-LHC) integrated luminosity (3000 fb^{-1}), which is used to scale our results. The only exception is BP1 which has a large rate due to the low mass of the heavy Higgs and likely is at the edge of being significant.

An interesting feature for many of the benchmarks presented is that they show destructive interference at low mass, where the $bb\gamma\gamma$ channel is a significant contributor to a potential detection of a LHC SM signal. This means that a first di-Higgs discovery might be delayed if alternative scenarios of this type are realized. This feature arises from an enhanced Higgs self-coupling, which would be needed for a first order electroweak phase transition [57–59].

The distribution of the $bb\gamma\gamma$ events after the preselection described in Table III for benchmarks BP1 and BP4 are shown in Fig. 7. Around a hundred excess events will survive the preselection for BP1, but destructive interference at $m_{hh} > m_S$ is not as prominent. For an example of the other benchmarks, we show the distribution for BP4. For BP4, the resonant signals or their interference patterns are likely not resolvable even at HL-LHC.

2. Two b -jets and two taus

This final state has a larger rate than $bb\gamma\gamma$, but the mass resolution is not as good due to the neutrinos in the final state. However, the presence of the taus helps in reducing the QCD background. Usually, this final state has the best sensitivity to intermediate invariant masses.

In Fig. 8, we show the invariant mass distributions after the preselection of Table III in the $b\bar{b}\tau\tau$ channel for BP4 and BP5. Both of these have a heavy Higgs at 800 GeV, but their interference patterns are opposite. There is an overall migration toward lower invariant masses in this channel due to the neutrinos in the final state. Therefore, the constructive interference at $m_{hh} < m_H$ for BP5 is more visible than the destructive interference of BP4.

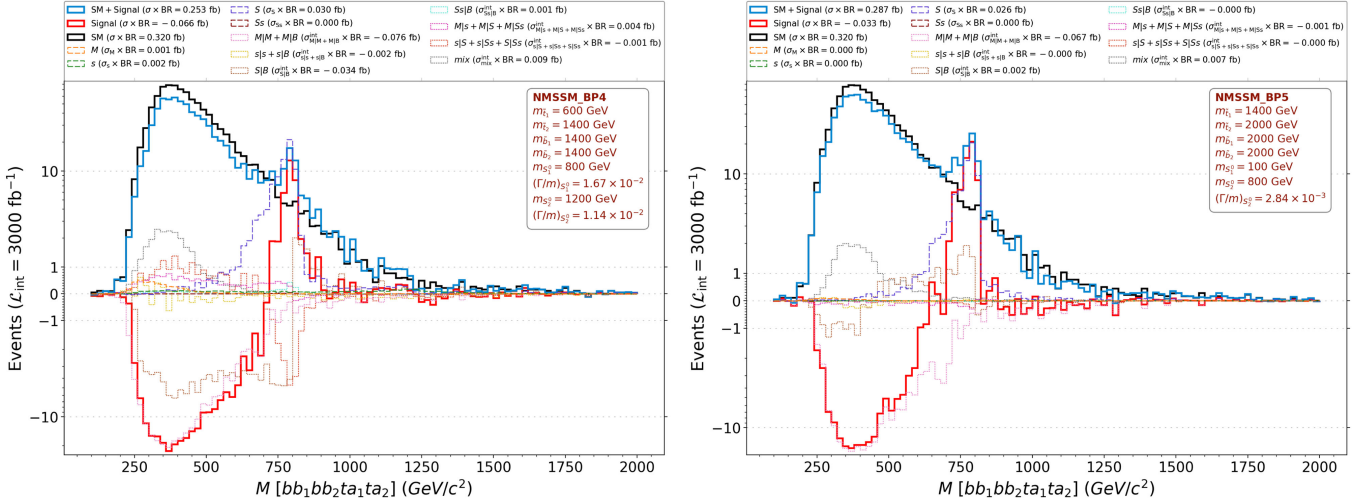


FIG. 8. Differential distribution of the di-Higgs invariant mass in the $bb\tau\tau$ channel for BP4 and BP5 at 3000 fb^{-1} .

Nevertheless, in comparison, there is a difference between the interference patterns, but the event rate at $m_H = 800 \text{ GeV}$ is so low that only HL-LHC can possibly distinguish the two cases.

3. Four b -jets

This final state has the largest event rate and the best sensitivity for high-mass resonances. However, due to low event yields in the high-mass tail, it will be more difficult to significantly disentangle any positive or negative interference compared to similar effects at lower mass.

In Fig. 9, we show the invariant mass distribution in the $4b$ final state for BP4 and BP5. The reconstruction efficiency for $b\bar{b}b\bar{b}$ is slightly below 10% in the relevant mass range. The interference patterns are even less clear than in the $bb\tau\tau$ final state, although there is some difference in shapes. For constructive interference at $m_{hh} < m_H$

and destructive interference at $m_{hh} > m_H$, the distribution is nearly flat and drops sharply at m_H , while the distribution for the opposite case is smoother.

We may also notice that, even though BP4 has a heavy Higgs at 1200 GeV , which is visible in the parton-level plot, the resolution of the $4b$ channel in the resolved jet analysis mode without any kinematical fit is not sufficient to produce any visible signal at that mass.

D. Discussion

The previous examples in NMSSM have shown the complex interplay between resonant and nonresonant contributions in Higgs pair production at the LHC, leading to different and potentially observable effects in the total cross sections and differential distributions.

The resonant signatures arising from a singletlike or a doubletlike Higgs have peculiar features. The doublet state

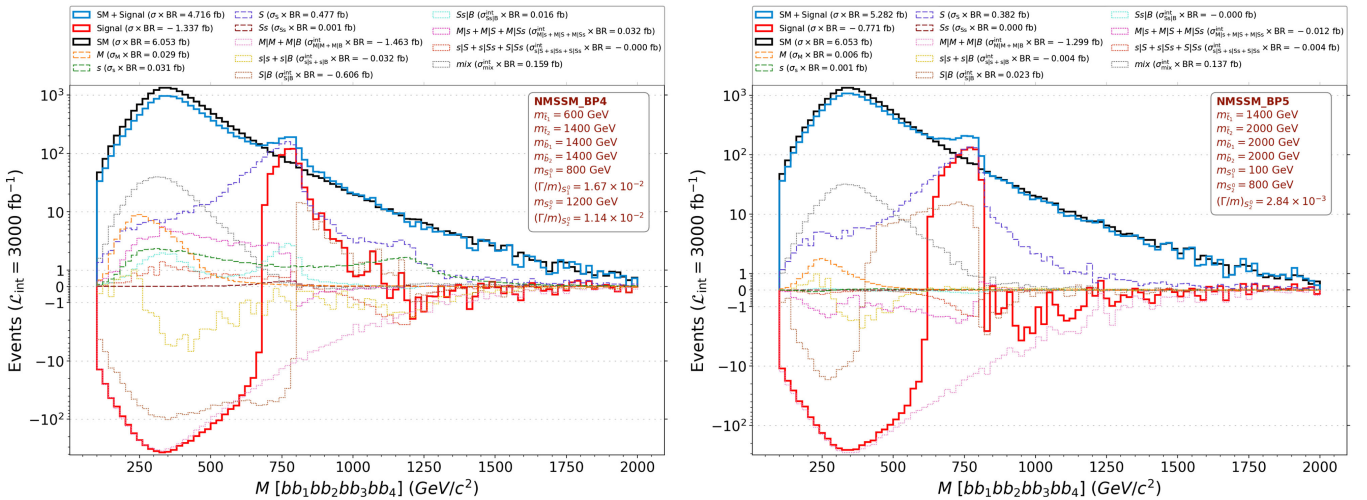


FIG. 9. Differential distribution of the di-Higgs invariant mass in the $4b$ channel for BP4 and BP5 at 3000 fb^{-1} .

has a stronger coupling to quarks and hence a larger overall cross section. It also has a larger width, which makes the interference effects stronger. The singlet has an interference pattern with an excess at $m_{hh} < m_S$ and a deficit at $m_{hh} > m_S$ whenever the singlet is visible [i.e., when A_λ is positive in the conventions of Eq. (10)]. The doublet interference pattern tells us qualitatively about the contributions to the 125 GeV mass. There is destructive interference at $m_{hh} < m_H$ if the mass is due to the tree-level superpotential contributions (the λ -term) and constructive interference at $m_{hh} < m_H$ if the mass is mainly due to loop-level top-stop contributions.

The presence of light squarks can modify these patterns, especially in the presence of neutral scalars with mass close to twice the mass of the squarks, where threshold effects play a significant role.

If a qualitative understanding of the interference patterns can be achieved, one can gain understanding of the nature of the Higgs sector. The deconstruction of the signal provides precious insight in this respect, allowing us to investigate the role of individual terms, cancellations, and enhancements and relate these to the relevant Lagrangian parameters, making it easy to connect the underlying physics to the corresponding features in the m_{hh} distribution. Such an analysis can obviously be extended beyond our NMSSM example, though the restrictions that supersymmetry imposes on the quartic scalar couplings make the interpretation somewhat easier.

At a practical level, the mass resolution of the most sensitive final states is, however, limited in current experimental setups, so interference effects get partially smeared out due to migration of events between nearby bins. Constructive interference survives better to detector level than destructive interference, so much so that we verified that, even at detector level, the distributions exhibit an asymmetric shape around the peak. Any definite conclusions of interference patterns will require the full integrated luminosity of the HL-LHC. Correlating more observables in the presence of a larger statistics associated to higher luminosities will also clearly help.

At the moment, our approach is limited to LO (one-loop) only. It is well known that di-Higgs production in the SM has large higher-order corrections [60–64]. In the heavy top limit, this leads to a relatively constant K-factor close to 1.8, but finite mass effects lead to a substantial dependence on the top mass scheme [65], and in different schemes, the variation of the K-factor as a function of m_{hh} varies. This can lead to shape errors of 20%–30%, though local variation is milder. For the BSM part, the next-to-leading-order (NLO) corrections are not always known. For squarks the NLO-QCD corrections have been computed in the limit of vanishing external momenta [66], and the K-factor corrections are close to the SM ones. For the resonant part, K-factors of heavy Higgs production in the MSSM are in the same ballpark as the ones for di-Higgs production in the SM [67,68], so our qualitative conclusions should hold.

IV. CONCLUSIONS

We have shown that customary approaches to resonant SM-like di-Higgs production at the LHC that treat s -channel heavy resonances in the narrow width approximation, or even as Breit-Wigner distributions, neglecting the contributions to the signal due to interference effects can miss relevant effects. Crucially, for our purposes, we have shown that not only can interferences between the heavy Higgs state graph and SM diagrams be relevant, but also those between such a graph and BSM diagrams carrying in the loop companion states to the top quark can be relevant. We have proven this to be the case for a prevalent theory of the EW scale, chiefly, SUSY, in its NMSSM incarnation (wherein such top quark companion states are top squarks). In fact, in the latter SUSY construct, one could even have two heavy Higgs states as resonant objects in s -channel, which an even more complicated pattern of interferences will further blur the validity of the aforementioned naive approaches.

This has been illustrated in the case of a variety of NMSSM BPs, wherein heavy Higgs states have been spaced in mass in such a way as to enter the entire di-Higgs invariant mass spectrum from 100 GeV, to $2m_t$, to 500 GeV, to 800 GeV, and to $2m_{\tilde{t}}$, thereby capturing all relevant possible phenomenological features of the differential di-Higgs cross section in this BSM framework. We have also shown that the qualitative features of interferences and relative strengths of squark and heavy Higgs modifications have a connection to the SM-like Higgs mass generation.

Furthermore, we have shown that these interferences effects, originally studied at the parton level for on-shell SM-like Higgs states, survive to a tangible extent upon allowing for the decays of the latter into $b\bar{b}\gamma\gamma$, $b\bar{b}\tau^+\tau^-$, and $b\bar{b}b\bar{b}$ final states in the presence of parton shower, fragmentation/hadronization, and realistic detector effects. However, this largely depends upon the achievable di-Higgs mass resolution in the aforementioned decay channels, in the sense that, while not all individual interference effects may be visible in actual data, the new signal definition (including a globally interfering contribution to it) induces the generic feature of a visible distortion of the BW shape which needs to be accounted for in experimental analysis.

This work is a continuation of a previous one of ours, which tackled the nonresonant di-Higgs process, wherein we had introduced a dedicated framework to deconstruct the signal in a finite set of independent terms to be recombined using specific coupling combinations as weights. This deconstruction technique is able to capture the individual contributions of interferences, finite widths, or threshold effects and has allowed us to map the full results of the di-Higgs process in terms of simplified model structures, providing a prompt interpretation of the underlying BSM dynamics. The modularity of this technique allows

straightforward generalizations, such as the present case where resonant di-Higgs production has been included. Indeed, such a framework can readily be extended to also accommodate other types of other colored particles, such as fermionic top quark partners like, e.g., vectorlike quarks in compositeness, so as to generalize our approach to other theories of the EW scale with top companion states.

ACKNOWLEDGMENTS

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DATA AVAILABILITY

Some of the data that support the findings of this article are openly available [69]. Benchmark points used in the study are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: DECONSTRUCTED CROSS SECTIONS

In this appendix, we provide the complete set of contributions to the di-Higgs cross section, Eq. (A1). The conventions adopted throughout the equations are the following:

- (i) The reduced cross sections $\hat{\sigma}$ have been labeled according to the numbering scheme of Table I, and

their arguments are the masses (and widths for neutral scalars) of the particles propagating in the loops; factors of 2 from interference terms have been included in the $\hat{\sigma}$'s.

- (ii) The SM irreducible background is labeled as “B,” the contribution of the topologies containing only modified SM couplings (1 to 4 in Table I) are labeled as “M,” the contribution of topologies associated with the propagation of colored scalars (5 to 8 in Table I) is labeled as “ $\tilde{s}$,” those associated with the neutral scalars (9 in Table I) are labeled as S , and the contributions with both neutral and colored scalar (10 in Table I) are labeled as $S\tilde{s}$.
- (iii) Summations over SM quarks have index f, f' , over squarks have index i, j , and over neutral scalars have index I, J ; when summing over multiple indices, if the cross section does not change by exchange of the indices, the summation is only for larger values of the second index ($\sum_i \sum_{j>i}$), and otherwise the summation is just for different values ($\sum_i \sum_{j \neq i}$).
- (iv) $\hat{\sigma}_{x|y}^{\text{int}}$ and $\hat{\sigma}_{x|B}^{\text{int}}$ represent, respectively, the interference terms between the x and y classes of topologies and between x and the SM background, while $\hat{\sigma}_{x|y-j|k}^{\text{int}}$ and $\hat{\sigma}_{x|y-j|B}^{\text{int}}$ are signal combinations proportional to the same function of couplings.
- (v) For reduced cross sections involving squarks or scalars, the summation index associated with the new particles is reported next to the topology class in the subscript of the $\hat{\sigma}$; e.g., $\hat{\sigma}_{7i|7jk}^{\text{int}}$ represent interference between two topologies of type 7 in which only the squark i circulates in one and the squarks j and k circulate in the other.

The complete list of contributions is

$$\sigma_M = \kappa_{hhh}^2 \hat{\sigma}_1 + \sum_{f=t,b} [(\kappa_{hhh} \kappa_{hff})^2 \hat{\sigma}_{3f} + \kappa_{hff}^4 \hat{\sigma}_{4f}], \quad (\text{A1a})$$

$$\begin{aligned} \sigma_{\tilde{s}} = & \sum_i \left[\kappa_{h\tilde{s}\tilde{s}}^{ii} \sum_{j>i} \kappa_{h\tilde{s}\tilde{s}}^{jj} \hat{\sigma}_{5i|5j}^{\text{int}}(m_{\tilde{s}_{i,j}}) + \kappa_{hhh}^2 (\kappa_{h\tilde{s}\tilde{s}}^{ii})^2 \hat{\sigma}_{6i}(m_{\tilde{s}_i}) + \kappa_{hhh}^2 \kappa_{h\tilde{s}\tilde{s}}^{ii} \sum_{j>i} \kappa_{h\tilde{s}\tilde{s}}^{jj} \hat{\sigma}_{6i|6j}^{\text{int}}(m_{\tilde{s}_{i,j}}) \right. \\ & + (\kappa_{h\tilde{s}\tilde{s}}^{ii})^4 \hat{\sigma}_{7i}(m_{\tilde{s}_i}) + \sum_{j>i} (\kappa_{h\tilde{s}\tilde{s}}^{ij})^4 \hat{\sigma}_{7ij}(m_{\tilde{s}_{i,j}}) + (\kappa_{h\tilde{s}\tilde{s}}^{ii})^2 \sum_{j>i} (\kappa_{h\tilde{s}\tilde{s}}^{jj})^2 \hat{\sigma}_{7i|7j}^{\text{int}}(m_{\tilde{s}_{i,j}}) + (\kappa_{h\tilde{s}\tilde{s}}^{ii})^2 \sum_{j \neq i} (\kappa_{h\tilde{s}\tilde{s}}^{ij})^2 \hat{\sigma}_{7i|7ij}^{\text{int}}(m_{\tilde{s}_{i,j}}) \\ & + (\kappa_{h\tilde{s}\tilde{s}}^{ii})^2 \sum_{j \neq i} \sum_{\substack{k>j \\ k \neq i}} (\kappa_{h\tilde{s}\tilde{s}}^{jk})^2 \hat{\sigma}_{7i|7jk}^{\text{int}}(m_{\tilde{s}_{i,j,k}}) + \sum_{j>i} \sum_{\substack{k>j \\ k \neq i}} \kappa_{h\tilde{s}\tilde{s}}^{ij} \kappa_{h\tilde{s}\tilde{s}}^{ik} \hat{\sigma}_{7ij|7ik}^{\text{int}}(m_{\tilde{s}_{i,j,k}}) + \sum_{j>i} \sum_{\substack{k \neq i \\ k \neq j}} \sum_{\substack{l>k \\ l \neq i \\ l \neq j}} \kappa_{h\tilde{s}\tilde{s}}^{ij} \kappa_{h\tilde{s}\tilde{s}}^{kl} \hat{\sigma}_{7ij|7kl}^{\text{int}}(m_{\tilde{s}_{i,j,k,l}}) \\ & \left. + (\kappa_{h\tilde{s}\tilde{s}}^{ii})^2 \hat{\sigma}_{8i}(m_{\tilde{s}_i}) + \kappa_{h\tilde{s}\tilde{s}}^{ii} \sum_{j>i} \kappa_{h\tilde{s}\tilde{s}}^{jj} \hat{\sigma}_{8i|8j}^{\text{int}}(m_{\tilde{s}_{i,j}}) \right], \quad (\text{A1b}) \end{aligned}$$

$$\sigma_S = \sum_I \left[\sum_{f=t,b} (\kappa_{Shh}^I \kappa_{Sff}^I)^2 \hat{\sigma}_{9If}(m_{S_I^0}, \Gamma_{S_I^0}) + (\kappa_{Shh}^I)^2 \kappa_{Stt}^I \kappa_{Sbb}^I \hat{\sigma}_{9It|9Ib}^{\text{int}}(m_{S_I^0}, \Gamma_{S_I^0}) \right. \\ \left. + \sum_{J>I} \sum_{f=t,b} \kappa_{Shh}^J \kappa_{Shh}^I \kappa_{Sff}^J \kappa_{Sff}^I \hat{\sigma}_{9If|9Jf}^{\text{int}}(m_{S_{I,J}^0}, \Gamma_{S_{I,J}^0}) + \sum_{J \neq I} \kappa_{Shh}^J \kappa_{Shh}^I \kappa_{Stt}^J \kappa_{Sbb}^I \hat{\sigma}_{9It|9Jb}^{\text{int}}(m_{S_{I,J}^0}, \Gamma_{S_{I,J}^0}) \right], \quad (\text{A1c})$$

$$\sigma_{S\bar{S}} = \sum_{i,l} \left[(\kappa_{Shh}^l \kappa_{S\bar{S}\bar{S}}^{li})^2 \hat{\sigma}_{10Ii}(m_{\bar{S}_i}, m_{S_I^0}, \Gamma_{S_I^0}) + \sum_{J>I} \kappa_{Shh}^J \kappa_{Shh}^I \kappa_{S\bar{S}\bar{S}}^{Ji} \kappa_{S\bar{S}\bar{S}}^{Ii} \hat{\sigma}_{10Ii|10Ji}^{\text{int}}(m_{\bar{S}_i}; m_{S_{I,J}^0}, \Gamma_{S_{I,J}^0}) \right. \\ \left. + \sum_{J>I} (\kappa_{Shh}^J)^2 \kappa_{S\bar{S}\bar{S}}^{Ji} \kappa_{S\bar{S}\bar{S}}^{Ij} \hat{\sigma}_{10Ii|10Jj}^{\text{int}}(m_{\bar{S}_{i,j}}; m_{S_I^0}, \Gamma_{S_I^0}) + \sum_{J \neq I} \sum_{j \neq i} \kappa_{Shh}^J \kappa_{Shh}^I \kappa_{S\bar{S}\bar{S}}^{Ji} \kappa_{S\bar{S}\bar{S}}^{Jj} \hat{\sigma}_{10Ii|10Jj}^{\text{int}}(m_{\bar{S}_{i,j}}; m_{S_{I,J}^0}, \Gamma_{S_{I,J}^0}) \right], \quad (\text{A1d})$$

$$\sigma_{M|B}^{\text{int}} = \kappa_{hhh} \hat{\sigma}_{1|B} + \sum_{f=t,b} \kappa_{hff} \hat{\sigma}_{2f|B}, \quad (\text{A1e})$$

$$\sigma_{\bar{S}|B}^{\text{int}} = \sum_i \left[\kappa_{h\bar{S}\bar{S}}^{ii} \hat{\sigma}_{5i|B}^{\text{int}}(m_{\bar{S}_i}) + \sum_{j>i} (\kappa_{h\bar{S}\bar{S}}^{ij})^2 \hat{\sigma}_{7ij|B}^{\text{int}}(m_{\bar{S}_{i,j}}) + \kappa_{h\bar{S}\bar{S}}^{ii} \hat{\sigma}_{8i|B}^{\text{int}}(m_{\bar{S}_i}) \right], \quad (\text{A1f})$$

$$\sigma_{S|B}^{\text{int}} = \sum_I \sum_{f=t,b} \kappa_{Shh}^I \kappa_{Sff}^I \hat{\sigma}_{9If|B}^{\text{int}}(m_{S_I^0}, \Gamma_{S_I^0}), \quad (\text{A1g})$$

$$\sigma_{S\bar{S}|B}^{\text{int}} = \sum_{i,l} \kappa_{Shh}^l \kappa_{S\bar{S}\bar{S}}^{li} \hat{\sigma}_{10Ii|B}^{\text{int}}(m_{\bar{S}_i}; m_{S_I^0}, \Gamma_{S_I^0}), \quad (\text{A1h})$$

$$\sigma_{M|M}^{\text{int}} = \sum_{f=t,b} \left[\kappa_{hhh}^2 \kappa_{hff} \hat{\sigma}_{1|3f}^{\text{int}} + \kappa_{hhh} \kappa_{hff}^2 \hat{\sigma}_{1|4f-2f|3f}^{\text{int}} + \kappa_{hff}^3 \hat{\sigma}_{2f|4f}^{\text{int}} + \kappa_{hhh} \kappa_{hff}^3 \hat{\sigma}_{3f|4f}^{\text{int}} \right. \\ \left. + \sum_{f' \neq f} (\kappa_{hhh} \kappa_{hff} \kappa_{hff'} \hat{\sigma}_{2f|3f'}^{\text{int}} + \kappa_{hff} \kappa_{hff'}^2 \hat{\sigma}_{2f|4f'}^{\text{int}} + \kappa_{hhh} \kappa_{hff} \kappa_{hff'}^2 \hat{\sigma}_{3f|4f'}^{\text{int}}) \right], \quad (\text{A1i})$$

$$\sigma_{\bar{S}|\bar{S}}^{\text{int}} = \sum_i \left[\kappa_{hhh} \kappa_{h\bar{S}\bar{S}}^{ii} \sum_{j>i} \kappa_{h\bar{S}\bar{S}}^{jj} \hat{\sigma}_{5i|6j}^{\text{int}}(m_{\bar{S}_{i,j}}) + (\kappa_{h\bar{S}\bar{S}}^{ii})^3 \hat{\sigma}_{5i|7i}^{\text{int}}(m_{\bar{S}_i}) + \kappa_{h\bar{S}\bar{S}}^{ii} \sum_{j \neq i} (\kappa_{h\bar{S}\bar{S}}^{jj})^2 \hat{\sigma}_{5i|7j}^{\text{int}}(m_{\bar{S}_{i,j}}) \right. \\ + \kappa_{h\bar{S}\bar{S}}^{ii} \sum_{j \neq i} (\kappa_{h\bar{S}\bar{S}}^{ij})^2 \hat{\sigma}_{5i|7ij}^{\text{int}}(m_{\bar{S}_{i,j}}) + \kappa_{h\bar{S}\bar{S}}^{ii} \sum_{j \neq i} \sum_{k>j} (\kappa_{h\bar{S}\bar{S}}^{jk})^2 \hat{\sigma}_{5i|7jk}^{\text{int}}(m_{\bar{S}_{i,j,k}}) + \kappa_{h\bar{S}\bar{S}}^{ii} \kappa_{h\bar{S}\bar{S}}^{ii} \hat{\sigma}_{5i|8i}^{\text{int}}(m_{\bar{S}_i}) \\ + \kappa_{h\bar{S}\bar{S}}^{ii} \sum_{j \neq i} \kappa_{h\bar{S}\bar{S}}^{jj} \hat{\sigma}_{5i|8j}^{\text{int}}(m_{\bar{S}_{i,j}}) + \kappa_{hhh} (\kappa_{h\bar{S}\bar{S}}^{ii})^3 \hat{\sigma}_{6i|7i}^{\text{int}}(m_{\bar{S}_i}) + \kappa_{hhh} \kappa_{h\bar{S}\bar{S}}^{ii} \sum_{j \neq i} (\kappa_{h\bar{S}\bar{S}}^{jj})^2 \hat{\sigma}_{6i|7j}^{\text{int}}(m_{\bar{S}_{i,j}}) \\ + \kappa_{hhh} \kappa_{h\bar{S}\bar{S}}^{ii} \sum_{j \neq i} (\kappa_{h\bar{S}\bar{S}}^{ij})^2 \hat{\sigma}_{6i|7ij}^{\text{int}}(m_{\bar{S}_{i,j}}) + \kappa_{hhh} \kappa_{h\bar{S}\bar{S}}^{ii} \sum_{j \neq i} \sum_{k>j} (\kappa_{h\bar{S}\bar{S}}^{jk})^2 \hat{\sigma}_{6i|7jk}^{\text{int}}(m_{\bar{S}_{i,j,k}}) + \kappa_{hhh} \kappa_{h\bar{S}\bar{S}}^{ii} \kappa_{h\bar{S}\bar{S}}^{ii} \hat{\sigma}_{6i|8i}^{\text{int}}(m_{\bar{S}_i}) \\ + \kappa_{hhh} \kappa_{h\bar{S}\bar{S}}^{ii} \sum_{j \neq i} \kappa_{h\bar{S}\bar{S}}^{jj} \hat{\sigma}_{6i|8j}^{\text{int}}(m_{\bar{S}_{i,j}}) + (\kappa_{h\bar{S}\bar{S}}^{ii})^2 \kappa_{h\bar{S}\bar{S}}^{ii} \hat{\sigma}_{7i|8i}^{\text{int}}(m_{\bar{S}_i}) + (\kappa_{h\bar{S}\bar{S}}^{ii})^2 \sum_{j \neq i} \kappa_{h\bar{S}\bar{S}}^{jj} \hat{\sigma}_{7i|8j}^{\text{int}}(m_{\bar{S}_{i,j}}) \\ + \sum_{j \neq i} (\kappa_{h\bar{S}\bar{S}}^{ij})^2 \kappa_{h\bar{S}\bar{S}}^{ii} \hat{\sigma}_{7ij|8i}^{\text{int}}(m_{\bar{S}_{i,j}}) + \sum_{j>i} \sum_{k \neq i} (\kappa_{h\bar{S}\bar{S}}^{ij})^2 \kappa_{h\bar{S}\bar{S}}^{kk} \hat{\sigma}_{7ij|8k}^{\text{int}}(m_{\bar{S}_{i,j,k}}) \left. \right], \quad (\text{A1j})$$

$$\begin{aligned}
\sigma_{M|\bar{s}}^{\text{int}} = & \sum_i \left[\kappa_{hhh}^2 \kappa_{h\bar{s}\bar{s}}^{ii} \hat{\sigma}_{1|6i}^{\text{int}}(m_{\bar{s}_i}) + \kappa_{hhh} \sum_{j>i} (\kappa_{h\bar{s}\bar{s}}^{ij})^2 \hat{\sigma}_{1|7ij}^{\text{int}}(m_{\bar{s}_{i,j}}) + \kappa_{hhh} \kappa_{h\bar{s}\bar{s}}^{ii} \hat{\sigma}_{1|8i}^{\text{int}}(m_{\bar{s}_i}) \right. \\
& + \sum_{f=t,b} \left(\kappa_{hff} \kappa_{h\bar{s}\bar{s}}^{ii} \hat{\sigma}_{2f|5i}^{\text{int}}(m_{\bar{s}_i}) + \kappa_{hhh} \kappa_{hff} \kappa_{h\bar{s}\bar{s}}^{ii} \hat{\sigma}_{2f|6i-3f|5i}^{\text{int}}(m_{\bar{s}_i}) + \kappa_{hff} (\kappa_{h\bar{s}\bar{s}}^{ii})^2 \hat{\sigma}_{2f|7o}^{\text{int}}(m_{\bar{s}_i}) \right. \\
& + \kappa_{hff} \sum_{j>i} (\kappa_{h\bar{s}\bar{s}}^{ij})^2 \hat{\sigma}_{2f|7ij}^{\text{int}}(m_{\bar{s}_{i,j}}) + \kappa_{hff} \kappa_{hh\bar{s}\bar{s}}^{ii} \hat{\sigma}_{2f|8i}^{\text{int}}(m_{\bar{s}_i}) + \kappa_{hhh}^2 \kappa_{hff} \kappa_{h\bar{s}\bar{s}}^{ii} \hat{\sigma}_{3f|6i}^{\text{int}}(m_{\bar{s}_i}) + \kappa_{hhh} \kappa_{hff} (\kappa_{h\bar{s}\bar{s}}^{ii})^2 \hat{\sigma}_{3f|7i}^{\text{int}}(m_{\bar{s}_i}) \\
& + \kappa_{hhh} \kappa_{hff} \sum_{j>i} (\kappa_{h\bar{s}\bar{s}}^{ij})^2 \hat{\sigma}_{3f|7ij}^{\text{int}}(m_{\bar{s}_{i,j}}) + \kappa_{hhh} \kappa_{hff} \kappa_{hh\bar{s}\bar{s}}^{ii} \hat{\sigma}_{3f|8i}^{\text{int}}(m_{\bar{s}_i}) + \kappa_{hff}^2 \kappa_{h\bar{s}\bar{s}}^{ii} \hat{\sigma}_{4f|5i}^{\text{int}}(m_{\bar{s}_i}) + \kappa_{hhh} \kappa_{hff}^2 \kappa_{h\bar{s}\bar{s}}^{ii} \hat{\sigma}_{4f|6i}^{\text{int}}(m_{\bar{s}_i}) \\
& \left. \left. + \kappa_{hff}^2 (\kappa_{h\bar{s}\bar{s}}^{ii})^2 \hat{\sigma}_{4f|7i}^{\text{int}}(m_{\bar{s}_i}) + \kappa_{hff}^2 \sum_{j>i} (\kappa_{h\bar{s}\bar{s}}^{ij})^2 \hat{\sigma}_{4f|7ij}^{\text{int}}(m_{\bar{s}_{i,j}}) + \kappa_{hff}^2 \kappa_{hh\bar{s}\bar{s}}^{ii} \hat{\sigma}_{4f|8i}^{\text{int}}(m_{\bar{s}_i}) \right) \right], \quad (\text{A1k})
\end{aligned}$$

$$\begin{aligned}
\sigma_{M|S}^{\text{int}} = & \sum_I \sum_{f=t,b} \left[\kappa_{hhh} \kappa_{Shh}^L \kappa_{Sff}^L \hat{\sigma}_{1|9If}^{\text{int}}(m_{S_I^0}, \Gamma_{S_I^0}) + \kappa_{hff} \kappa_{Shh}^L \kappa_{Sff}^L \hat{\sigma}_{2f|9If}^{\text{int}}(m_{S_I^0}, \Gamma_{S_I^0}) \right. \\
& + \kappa_{hhh} \kappa_{hff} \kappa_{Shh}^L \kappa_{Sff}^L \hat{\sigma}_{3f|9If}^{\text{int}}(m_{S_I^0}, \Gamma_{S_I^0}) + \kappa_{hff}^2 \kappa_{Shh}^L \kappa_{Sff}^L \hat{\sigma}_{4f|9If}^{\text{int}}(m_{S_I^0}, \Gamma_{S_I^0}) \\
& + \sum_{f' \neq f} (\kappa_{hff} \kappa_{Shh}^L \kappa_{Sff'}^L \hat{\sigma}_{2f|9If'}^{\text{int}}(m_{S_I^0}, \Gamma_{S_I^0}) + \kappa_{hhh} \kappa_{hff} \kappa_{Shh}^L \kappa_{Sff'}^L \hat{\sigma}_{3f|9If'}^{\text{int}}(m_{S_I^0}, \Gamma_{S_I^0}) \\
& \left. + \kappa_{hff}^2 \kappa_{Shh}^L \kappa_{Sff'}^L \hat{\sigma}_{4f|9If'}^{\text{int}}(m_{S_I^0}, \Gamma_{S_I^0}) \right) \quad (\text{A1l})
\end{aligned}$$

$$\begin{aligned}
\sigma_{M|\bar{S}\bar{S}}^{\text{int}} = & \sum_{I,i} \left[\kappa_{hhh} \kappa_{Shh}^L \kappa_{S\bar{s}\bar{s}}^{Li} \hat{\sigma}_{1|10Ii}^{\text{int}}(m_{\bar{s}_i}; m_{S_I^0}, \Gamma_{S_I^0}) + \sum_{f=t,b} (\kappa_{hff} \kappa_{Shh}^L \kappa_{S\bar{s}\bar{s}}^{Li} \hat{\sigma}_{2f|10Ii}^{\text{int}}(m_{\bar{s}_i}; m_{S_I^0}, \Gamma_{S_I^0}) \right. \\
& \left. + \kappa_{hhh} \kappa_{hff} \kappa_{Shh}^L \kappa_{S\bar{s}\bar{s}}^{Li} \hat{\sigma}_{3f|10Ii}^{\text{int}}(m_{\bar{s}_i}; m_{S_I^0}, \Gamma_{S_I^0}) + \kappa_{hff}^2 \kappa_{Shh}^L \kappa_{S\bar{s}\bar{s}}^{Li} \hat{\sigma}_{4f|10Ii}^{\text{int}}(m_{\bar{s}_i}; m_{S_I^0}, \Gamma_{S_I^0}) \right] \quad (\text{A1m})
\end{aligned}$$

$$\begin{aligned}
\sigma_{\bar{s}|S}^{\text{int}} = & \sum_{I,i} \sum_{f=t,b} \left[\kappa_{h\bar{s}\bar{s}}^i \kappa_{Shh}^L \kappa_{Sff}^L \hat{\sigma}_{5i|9If}^{\text{int}}(m_{\bar{s}_i}; m_{S_I^0}, \Gamma_{S_I^0}) + \kappa_{hhh} \kappa_{h\bar{s}\bar{s}}^i \kappa_{Shh}^L \kappa_{Sff}^L \hat{\sigma}_{6i|9If}^{\text{int}}(m_{\bar{s}_i}; m_{S_I^0}, \Gamma_{S_I^0}) \right. \\
& + (\kappa_{h\bar{s}\bar{s}}^i)^2 \kappa_{Shh}^L \kappa_{Sff}^L \hat{\sigma}_{7i|9If}^{\text{int}}(m_{\bar{s}_i}; m_{S_I^0}, \Gamma_{S_I^0}) + \sum_{j>i} \kappa_{h\bar{s}\bar{s}}^{ij} \kappa_{Shh}^L \kappa_{Sff}^L \hat{\sigma}_{7ij|9If}^{\text{int}}(m_{\bar{s}_{i,j}}; m_{S_I^0}, \Gamma_{S_I^0}) \\
& \left. + \kappa_{h\bar{s}\bar{s}}^i \kappa_{Shh}^L \kappa_{Sff}^L \hat{\sigma}_{8i|9If}^{\text{int}}(m_{\bar{s}_i}; m_{S_I^0}, \Gamma_{S_I^0}) \right] \quad (\text{A1n})
\end{aligned}$$

$$\begin{aligned}
\sigma_{\bar{s}|\bar{S}\bar{S}}^{\text{int}} = & \sum_{I,i,j} \left[\kappa_{h\bar{s}\bar{s}}^i \kappa_{Shh}^L \kappa_{S\bar{s}\bar{s}}^{Lj} \hat{\sigma}_{5i|10Ij}^{\text{int}}(m_{\bar{s}_{i,j}}; m_{S_I^0}, \Gamma_{S_I^0}) + \kappa_{hhh} \kappa_{h\bar{s}\bar{s}}^i \kappa_{Shh}^L \kappa_{S\bar{s}\bar{s}}^{Lj} \hat{\sigma}_{6i|10Ij}^{\text{int}}(m_{\bar{s}_{i,j}}; m_{S_I^0}, \Gamma_{S_I^0}) \right. \\
& \left. + \sum_{k \geq i} (\kappa_{h\bar{s}\bar{s}}^{ik})^2 \kappa_{Shh}^L \kappa_{S\bar{s}\bar{s}}^{Lj} \hat{\sigma}_{7ik|10Ij}^{\text{int}}(m_{\bar{s}_{i,j,k}}; m_{S_I^0}, \Gamma_{S_I^0}) + \kappa_{h\bar{s}\bar{s}}^i \kappa_{Shh}^L \kappa_{S\bar{s}\bar{s}}^{Lj} \hat{\sigma}_{8i|10Ij}^{\text{int}}(m_{\bar{s}_{i,j}}; m_{S_I^0}, \Gamma_{S_I^0}) \right] \quad (\text{A1o})
\end{aligned}$$

$$\sigma_{S|\bar{S}\bar{S}}^{\text{int}} = \sum_{I,J,i} \sum_{f=t,b} \kappa_{Shh}^L \kappa_{Sff}^L \kappa_{Shh}^L \kappa_{S\bar{s}\bar{s}}^{Lj} \hat{\sigma}_{9If|10Ji}^{\text{int}}(m_{\bar{s}_i}; m_{S_{I,J}^0}, \Gamma_{S_{I,J}^0}) \quad (\text{A1p})$$

$$\begin{aligned}
\sigma_{\text{mix}}^{\text{int}} = & \sum_{f=t,b} [\kappa_{hhh} \kappa_{hff} \hat{\sigma}_{1|2f-3f|B}^{\text{int}} + \kappa_{hff}^2 \hat{\sigma}_{2f-4f|B}^{\text{int}}] \\
& + \sum_i [\kappa_{hhh} \kappa_{h\bar{s}\bar{s}}^{ii} \hat{\sigma}_{1|5i-6i|B}^{\text{int}}(m_{\bar{s}_i}) + \kappa_{hhh} (\kappa_{h\bar{s}\bar{s}}^{ii})^2 \hat{\sigma}_{1|7i-5i|6i}^{\text{int}}(m_{\bar{s}_i}) + (\kappa_{h\bar{s}\bar{s}}^{ii})^2 \hat{\sigma}_{5i-7i|B}^{\text{int}}(m_{\bar{s}_i})]. \quad (\text{A1q})
\end{aligned}$$

APPENDIX B: UFO MODEL CONTENT

The UFO model provided with this analysis can be downloaded at [69]. It contains the minimal particle content to generate any signal and interference term for di-Higgs studies involving the propagation of new colored or neutral particles, at LO, using MG5_aMC.

As mentioned in Sec. II, this amount in four colored scalars of unspecified charge, labeled as s_{q1} , s_{q2} , s_{q3} , and s_{q4} (already present in the UFO model used for our previous analysis [11]), two neutral scalars labeled as s_{01} and s_{02} , four vectorlike quarks with charge $2/3$ (t_{p1} , t_{p2} , t_{p3} , and t_{p4}), and three vectorlike quarks with charge $-1/3$ (b_{p1} , b_{p2} , and b_{p3}). The new fermions have not been used in this analysis but will be included in a future complete treatment of the process. For fermions, the charge is important, as they can mix with the SM third-generation

fermions, and the different number of top and bottom partners is again related to the minimal number of particles necessary to describe all combinations appearing when squaring amplitudes in the deconstruction. The couplings, coupling orders associated to the new particles, and simulation syntax for deconstructed elements follow the same rationale described in Appendix A of Ref. [11], to which we refer for further details.

The model contains three restriction cards. They are for the two Higgs doublet model, the MSSM, and NMSSM particle contents. Naturally, these can be used for any other model with the same relevant particle content; e.g., the 2HDM restriction card would be appropriate also for the SM extended with a singlet Higgs. To use the model, one will need to calculate the couplings using the conventions of Eq. (1).

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