Uncovering the relation of a scalar resonance to the Higgs boson

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We consider the associated production of a scalar resonance with the standard model Higgs boson. We demonstrate via a realistic phenomenological analysis that couplings of such a resonance to the Higgs boson can be constrained in a meaningful way in future runs of the LHC, providing insights on its origin and its relation to the electroweak symmetry breaking sector. Moreover, the final state can provide a direct way to determine whether the new resonance is produced predominantly in gluon fusion or quark-anti-quark annihilation. The analysis focuses on a resonance coming from a scalar field with vanishing vacuum expectation value and its decay to a photon pair. It can however be straightforwardly generalised to other scenarios.

I. NEW SCALAR RESONANCES AT COLLIDERS.

Models with an additional (pseudo-)scalar singlet with a mass of several hundred GeV represent a well motivated class of extensions of the Standard Model (SM) of particle physics, including composite Higgs scenarios, supersymmetry, Coleman-Weinberg models, models addressing the strong CP problem, models of flavor, as well as generic Higgs portal setups (see e.g. [1–9]). A particularly promising channel to search for and analyze such a particle is its decay to two photons. Beyond being possibly sizable in certain scenarios, it offers a robust and clean way to detect a signal, emerging over a steeply falling background [10, 11].

After its discovery, an important aspect of scrutinising any new resonance is in fact to measure its couplings, and hence determine its relation, to the known particle content of the Standard Model (SM). A crucial component of this task is to uncover its role in the arena of electroweak symmetry breaking (EWSB). As a first step in this direction, determination of the couplings of the new scalar to the SM-like Higgs boson is mandatory, which is the main focus of this article, employing its di-photon (γγ) decay channel.

For a γγ resonance originating from a scalar field \( \mathcal{S} \), neutral under the SM gauge group, the relevant effective Lagrangian for our study – augmenting the SM at dimension \( D \leq 5 \) – is

\[
\mathcal{L}_{\text{eff}} \supset \frac{1}{2} \partial_{\mu} \mathcal{S} \partial^{\mu} \mathcal{S} - \frac{1}{2} \mu_{\mathcal{S}}^2 \mathcal{S}^2 - \frac{1}{2} \mu_{\mathcal{B}}^2 \mathcal{B}^2 - \frac{1}{2} \mu_{\mathcal{W}}^2 \mathcal{W}^2 + \lambda_{\mathcal{B}} \mathcal{B}^2 \mathcal{S}^2 + \lambda_{\mathcal{W}} \mathcal{W}^2 \mathcal{S}^2 + \lambda \mathcal{S}^4 + \frac{1}{4} \frac{\lambda_{\mathcal{H} \mathcal{S}}}{\Lambda^2} \mathcal{H}^2 \mathcal{S}^2 - \frac{1}{4} \frac{\lambda_{\mathcal{H} \mathcal{S} \mathcal{H}}}{\Lambda^4} \mathcal{H}^4 - \frac{1}{2} \mu_{\mathcal{H}}^2 \mathcal{H}^2 + \text{linear terms} + \text{higher dimensional operators},
\]

(1)

Here, \( Q_{L}^i \) is the \( i \)-th generation left-handed \( SU(2)_L \) fermion doublet, \( d_R^j \) and \( u_R^j \) are the right-handed \( SU(2)_L \) fermion singlets for generation \( j \), \( (y_{S}^q)^{ij} \) are the corresponding Yukawa-like couplings, \( c_{B} \), \( c_{W} \) are the couplings of \( \mathcal{S} \) to the \( U(1)_Y \) and \( SU(2)_L \) gauge fields \( B \) and \( W \), \( c_{G}^{S} \) is the coupling to the gluon fields, \( H \) is the Higgs boson doublet, \( \lambda_{\mathcal{H} \mathcal{S}} \) is the Higgs-Scalar portal coupling and \( \lambda_{\mathcal{S}} \) the new scalar quartic. Moreover, \( \Lambda \) denotes the scale of heavy new physics (NP), mediating the contact interactions of \( \mathcal{S} \) with SM gauge bosons and fermions (the latter involving \( H \) to generate a gauge singlet). Note that we do not include terms with an odd number of \( \mathcal{S} \) fields containing only scalars (as well as lepton fields). The corresponding interaction vertices will turn out irrelevant in general for the process we will consider, see below.† Beyond that, terms linear in \( \mathcal{S} \) could also lead to the singlet mixing with the Higgs boson after EWSB, which would in fact affect its phenomenology. Although such effects could still be present at a non-negligible level, they are expected to be sub-leading and we neglect them for simplicity, see Appendix [3]. Furthermore, the analysis that follows is independent of the CP properties of \( \mathcal{S} \) and its interactions and henceforth, for simplicity, we assume it to be CP-even with CP-conserving interactions.

With the potential for the Higgs doublet \( H \) taking the conventional form:

\[
\mathcal{V} = \lambda_{H} |H|^4 - \mu_{H}^2 |H|^2 ,
\]

(2)

†The full list of potential \( D \leq 5 \) operators is listed in Appendix A.

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A specific motivation for the first version of this manuscript was provided by the apparent \( \gamma \gamma \) resonance at \( M_{\gamma \gamma} \sim 750 \text{ GeV} \) in ATLAS [12, 13] and CMS [14, 15] data. This turned out not to be present in the 2016 data [16, 17]. Consequently the article was generalised to other mass scales of a potential scalar resonance, which remains well motivated, taking into account new limits on its cross section - see below. Comprehensive analyses studying constraints on (other) possible couplings of a di-photon resonance as well as detailed examinations of indirect footprints of new (high multiplicity) sectors, linked to its productions or decay appeared e.g. in [18–45].
and assuming that $H$ is the only scalar that gets a vacuum expectation value (vev), $|\langle H \rangle| = v/\sqrt{2}$, triggering EWSB, we obtain the condition\(^\dagger\)

$$\lambda_{HS}^2 - \lambda_H \mu_S^2 < 0, \quad \mu_H^2 < 0,$$  \hspace{1cm} (3)

The physical mass of the singlet thus reads $M = \sqrt{\mu_S^2 + \lambda_{HS}^2}$. The resulting trilinear interactions between the physical scalar resonances after EWSB are described by

$$\mathcal{L}_\text{scalar}^4 = -\frac{M^2}{2v} h^3 - \lambda_{HS} v h S^2,$$  \hspace{1cm} (4)

where $h$ is the Higgs boson, which (due to the case of negligible scalar mixing) is basically fully embedded in $H$ and describes excitations around its vev, such that in unitary gauge $H \approx \sqrt{2}(0, v + h)\dagger$, and $S$ is a new scalar resonance, which can also be (approximately) identified as $S = S$. Moreover, $v \approx 246$ GeV is the Higgs vev, $M_h \approx 125$ GeV is the measured Higgs boson mass and the portal coupling $\lambda_{HS}$ is to be determined, being a main scope of this paper. This coupling would be basically unconstrained by direct observation of a di-photon resonance. Nevertheless, loose indirect constraints can be derived, requiring vacuum stability not to be spoiled. They read

$$\lambda_S > 0, \quad \lambda_H > 0, \quad \lambda_{HS}^2 < \lambda_H \lambda_S,$$  \hspace{1cm} (5)

and need to be imposed at least at the scale where the NP enters, i.e., the TeV scale (see, e.g., \cite{11, 16}).\(^\ddagger\) Requiring $\lambda_H \sim 0.13$, to fit the observed Higgs mass, as well as $\lambda_S < (4\pi)^2$, we thus obtain $-4.5 \lesssim \lambda_{HS} \lesssim 4.5$. We will see below that in general our analysis can put stronger bounds than these on $\lambda_{HS}$.\(^\ddagger\)

In the present study we consider measurement of the coupling $\lambda_{HS}$ at the LHC, where it can be probed via associated production of the new resonance with the SM-like Higgs boson: $pp \rightarrow hS$. For this process, the interactions neglected in \cite{11} play no role to good approximation: they would either not enter at leading order (LO), or, as is the case for the $|H|^2 S$ interactions, contribute at most to a diagram with a (strongly suppressed) off-shell Higgs boson propagator, for details see the appendices \cite{A} and \cite{B}.

\(^\dagger\)The fact that $\langle S \rangle = 0$ guarantees the full absence of scalar mixing, that could otherwise occur even without linear terms in $S$.

\(^\ddagger\)While the first two conditions need to hold at all scales, for $\lambda_{HS} > 0$ the last condition might be violated at higher scales, while still the electroweak vacuum remains stable.\(^\dag\) Moreover, for the latter condition odd terms in $S$ are assumed to vanish.

\(^\dag\)Note that requiring a more conservative limit, such as $\lambda_S < O(10)$ (corresponding to, e.g., $d\lambda_S/\lambda_S < 1$ \cite{11}), restricts $\lambda_{HS}$ to be not much larger than 1 and would remove a considerable portion of the parameter space where our analysis exhibits sensitivity. However, in any case, the limits presented here are complementary to such considerations.

In principle several decay modes of $S$ can be considered. Here we focus on the process $pp \rightarrow hS \rightarrow h\gamma\gamma$, where the new particle decays to a pair of photons. Given that the tentative cross section of the resonant di-photon production, $pp \rightarrow S \rightarrow \gamma\gamma$, will be known and will be well-measured in the case of discovery, this allows constraints on the coupling $\lambda_{HS}$ to be imposed almost independently of the couplings to the initial-state partons and final-state photons, given that only a single production mode is relevant. We consider production via gluon fusion and quark-anti-quark annihilation, mediated through non-vanishing coefficients $c_{G}^S$, $(y_d^S)^{\pi\pi}$ or $(y_d^S)^{\pi\pi}$, respectively, and show how these modes could be disentangled by appropriate measurements. We will implicitly assume not too large values of $\Gamma(S \rightarrow \gamma\gamma)$, in such a way that photo-production is always subdominant.

Although we will focus on three specific benchmark masses of $M = 600, 750, 900$ GeV, our analysis could be applied to the general case of associated production of a Higgs boson with a scalar di-photon resonance of any mass. Moreover, several features of the final state studied here, such as the invariant mass of the final-state scalar or the total invariant mass of the process, will exhibit similar features when considering other decay modes.

The article is organised as follows: in Section \textbf{II} we examine the process of associated production of a scalar resonance and a Higgs boson, in Section \textbf{III} we describe the event generation and detector simulation setup and in Section \textbf{IV} we provide details of the analysis and results. Finally, we conclude in Section \textbf{V}.

\section{Associated Production of a Scalar Resonance and a Higgs Boson.}

\subsection{Production through gluon fusion.}

The dominant diagrams at LO contributing to the production of the $hS$ final state in gluon fusion and subsequent decay of the resonance $S$ to a pair of photons via the interactions

$$\mathcal{L}_\text{eff} \supset - S \frac{1}{16\pi^2} \left[ e^2 (c_{G}^S + c_{W}^S) F_{\mu\nu} F^{\mu\nu} + g_s c_{G} G^{\mu\nu} G_{\mu\nu} \right]$$

$$= - S \frac{1}{16\pi^2} \left[ e^2 c_{G}^S F_{\mu\nu} F^{\mu\nu} + g_s c_{G} G^{\mu\nu} G_{\mu\nu} \right],$$

are shown in Fig. \ref{fig:diagram}. In the analysis of the present article, we will consider the Higgs boson decaying to a bottom-quark pair, since this maximizes the expected number of events, which would be modest in general. There exist both the $s$-channel $S$ exchange, involving the portal coupling $\lambda_{HS}$ and depicted in the upper panel (a), as well as the ‘direct’ $hS$ production, via $t$-channel gluon exchange, depicted in the lower panel (b).
B. Production through quark-anti-quark annihilation.

For the case of quark anti-quark annihilation, a new diagram arises from the contact interaction $q\bar{q}hS$.\footnote{The $t$-channel diagram with the $q\bar{q}b$ interaction is suppressed due to a small Yukawa coupling.} Both contributing graphs are shown in Fig. 2. The new diagram (b) distinguishes the $q\bar{q}$ annihilation from the gluon fusion case. An important fact is that now the $hS$ process is non-vanishing and significant even in the absence of the portal coupling $\lambda_{HS}$. This indicates that one can employ this final state to exclude $q\bar{q}$ annihilation as the dominant production process (in the absence of a signal).

In the following we will focus on the cases of $q = b, q = s$ or $q = c$.

It will be useful in both scenarios to construct the ratio of the associated production process $pp \rightarrow hS \rightarrow h\gamma\gamma$, through all possible intermediate states, with the subsequent decay of the new resonance to a $\gamma\gamma$ final state, to that of the single production $pp \rightarrow S \rightarrow \gamma\gamma$,

$$\rho(xx') = \frac{\sigma(xx' \rightarrow hS \rightarrow h\gamma\gamma)}{\sigma(xx' \rightarrow S \rightarrow \gamma\gamma)},$$

where we consider $xx' = \{gg, b\bar{b}, s\bar{s}, c\bar{c}\}$.\footnote{In the most general setup, the analysis of this article can constrain the sum of the squares of the couplings of $S$ to all quark generations (for a given $\lambda_{HS}$), appropriately weighted by the parton density functions.} The ratio is a useful quantity since it removes the dependence on the product of couplings of the new resonance to the initial-state partons and final-state photons. Moreover, it can be used to absorb, at least approximately, theoretical and experimental systematic uncertainties.\footnote{For a similar idea investigated in the context of Higgs boson pair production, see \cite{47}.}

![Fig. 1: The diagrams contributing to the process $gg \rightarrow hS \rightarrow (b\bar{b})(\gamma\gamma)$ at the LHC at LO.](image1)

![Fig. 2: The diagrams contributing to the process $q\bar{q} \rightarrow hS \rightarrow (b\bar{b})(\gamma\gamma)$ at the LHC at LO.](image2)

![Fig. 3: The ratios $\rho(gg)$, $\rho(b\bar{b})$, $\rho(s\bar{s})$ and $\rho(c\bar{c})$ for $gg$, $b\bar{b}$, $s\bar{s}$ and $c\bar{c}$ initial states respectively, defined between the associated production $pp \rightarrow hS \rightarrow h\gamma\gamma$ and the single production $pp \rightarrow S \rightarrow \gamma\gamma$, as functions of the portal coupling $\lambda_{HS}$. The mass of the scalar resonance was taken to be $M = 750$ GeV and the width $\Gamma = 1$ GeV.](image3)
obtained for example, if $c_S^S \sim \mathcal{O}(10)$ and $c_G \sim \mathcal{O}(1)$ or $(g_G^2)^{33} \sim \mathcal{O}(1)$, for $\Lambda = 1$ TeV, with a cross section in $pp \to hS \to h\gamma\gamma$ compatible with current constraints. Similar behaviour of the ratio $\rho$ is obtained for different scalar $S$ masses and widths.

Since the dominant matrix-element contribution to the $gg$-initiated $hS$ process is proportional to the portal coupling, the process approximately vanishes as $\lambda_S \to 0$, and hence $\rho(gg) \approx \rho_{2gg}\lambda_S^2$, where $\rho_{2gg} \approx 0.00133$, for $M = 750$ GeV and $\Gamma = 1$ GeV, obtained by performing a quadratic fit of the $gg$ curve in Fig. 3. As already discussed, this does not hold for the $q\bar{q}$-initiated process due to the contact interaction diagram. This results in a non-negligible minimum for $\rho(q\bar{q})$. A fit to the cross section, again for $M = 750$ GeV and $\Gamma = 1$ GeV, yields $\rho(q\bar{q}) \approx \rho_{2qq}\lambda_S^2 + \rho_{1qq}\lambda_S + \rho_{0qq}$ with $\rho_{0bb} \approx 0.00828$, $\rho_{1bb} \approx 0.00309$, $\rho_{2bb} \approx 0.00086$, corresponding to $bb$ initial states, and $\rho_{0ss} \approx 0.01025$, $\rho_{1ss} \approx 0.00371$, $\rho_{2ss} \approx 0.00091$, corresponding to $ss$ initial states. The case of $c\bar{c}$ is similar to the $ss$ case and therefore in the rest of the article we focus on the cases $q = b$ and $q = s$. The positivity of the coefficient $\rho_{1qq}$ indicates constructive interference between the contact interaction and resonant diagrams (for $\lambda_S > 0$). For an extended fit of the ratio $\rho$, including additional diagrams with the production of an intermediate Higgs boson due to $S|H|^{2,4}$ interactions, that turn out to be sub-dominant, see Appendix B.

The fact that for quark-anti-quark annihilation the $hS$ process is non-vanishing for all values of the portal coupling $\lambda_S$ indicates that one could employ this final state to exclude $bb$, $ss$ or $c\bar{c}$ annihilation as the dominant production process. The analysis that will follow in the present article suggests however, that the di-photon decay of the $S$ alone may not be sufficient for that purpose for the benchmark points that we consider.

We show in Fig. 4 the variation of the ratio $\rho$ with the mass of the resonance, $M$, for the $gg$-initiated process and $\lambda_S = 1$, and for the $bb$-initiated process for $\lambda_S = 0$ (no portal) and $\lambda_S = 1$. We have fixed the width to $\Gamma = 1$ GeV. Interestingly, the pure $qq$-induced processes exhibit an increase of the ratio $\rho$ with increasing mass – related to the new $q\bar{q}hS$ interaction growing with momentum – whereas the pure $gg$-induced process exhibits a slight decrease.

If the di-photon resonance is wide, the analyses performed for the $hS$ final state will differ in the details due to changes in the kinematics. We show in Fig. 5 the variation of the ratio $\rho$ with the width over the mass, $\Gamma/M$, at a fixed mass $M = 750$ GeV, for $\lambda_S = 1$, and for the $bb$-initiated process for $\lambda_S = 0$ (no portal) and $\lambda_S = 1$. One can observe that the central value of the ratio remains approximately constant in all cases, with only a slight decrease with increasing width.

In both Figs. 4 and 5 we also provide, as coloured bands, the parton density function uncertainty for the MMHT14nlo68cl set combined in quadrature with the scale variation between 0.5 and 2.0 times the default central dynamical scale implemented in MadGraph5_aMC@NLO. For a mass of 750 GeV, the total theoretical uncertainties due to scale and PDF variations
are $\pm 10\%$ for the $gg$-induced process, $\sim 10\%$ for the $bb$-induced and $\sim 30\%$ for the $ss$-induced cases (the latter not shown in the figure for simplicity).

Assuming a total cross section for the production of a $\gamma\gamma$ resonance of mass $M = 750$ GeV of, say, $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 5$ fb (see below), one would expect a total of $\mathcal{O}(20)$ $hS \rightarrow h\gamma\gamma$ events at the high-luminosity LHC (HL-LHC, assuming 3000 fb$^{-1}$ of integrated luminosity) if the process is gluon-fusion initiated and $\mathcal{O}(200)$ events for $bb$-initiated production, for a portal coupling $\lambda_{HS} = 1$. Moreover, the minimum expected number of events for the $bb$-initiated process is $\mathcal{O}(80)$, arising for $\lambda_{HS} \approx -1.8$ and for the $ss$-initiated process one expects a minimum of $\mathcal{O}(100)$ events for $\lambda_{HS} \approx -2.0$. We note here that the positions of the minima for the $qq$-initiated process will change after cuts due to the varying effect of the analysis on the different pieces contributing to the cross section.

The kinematic structure of the $pp \rightarrow hS \rightarrow h\gamma\gamma$ process can be well-described by examining the distribution of the invariant mass of the $\gamma\gamma$ state, $M_{\gamma\gamma}$, or the distribution of the invariant mass of the Higgs boson and di-photon combination, $M_{h\gamma\gamma}$. In Figs. 6 and 7 we show, respectively, these distributions for the gluon-fusion-initiated process, for two widths, $\Gamma = 1$ GeV, $\Gamma = 45$ GeV. For the sake of clarity, here we only show distributions for a scalar mass of $M = 750$ GeV, but the main features remain unaltered as long as the scalar is heavier than the Higgs boson, $M > M_h$. The distributions clearly show the existence of two regions: a region in which the intermediate $s$-channel propagator for the $S$ scalar is on-shell and the final-state $S(\gamma\gamma)$ is off-shell, and a region in which the $s$-channel internal propagator is instead off-shell and the final-state $S(\gamma\gamma)$ is on-shell. The existence of the former region, $M_{\gamma\gamma} \lesssim M - M_h$, $M_{h\gamma\gamma} \sim M$ which henceforth we will call “three-body decay” since the intermediate $S$ is decaying approximately on-shell, is made possible by the fact that the mass of the particle produced in association with $S$, the Higgs boson, is smaller than the masses of $S$ we are considering, $M > M_h = 125$ GeV. The other region, $M_{\gamma\gamma} \sim M$, $M_{h\gamma\gamma} \gtrsim M + M_h$ which we will refer to as “on-shell di-photon”, exists irrespective of the mass of $S$. Note that both the three-body decay and on-shell di-photon regions exist even for $\Gamma/M \ll 1$. The normalised distributions look identical for all values of the portal coupling, $\lambda_{HS}$ ($\neq 0$), since the dominant contribution stems by far from the diagram shown in Fig. 1 (a).
\[ \frac{\Gamma}{\sigma} \left( \frac{d\sigma}{dM_{h\gamma\gamma}} \right) \]

\[ \Gamma = 1 \text{ GeV}, \lambda_{HS} = 1.0 \]

\[ \Gamma = 45 \text{ GeV}, \lambda_{HS} = 1.0 \]

\[ \Gamma = 1 \text{ GeV}, \lambda_{HS} = 0.0 \]

\[ \Gamma = 45 \text{ GeV}, \lambda_{HS} = 0.0 \]

FIG. 9: The matrix-element level distribution of the combined Higgs boson and di-photon invariant mass, \( M_{h\gamma\gamma} \), in the \( b\bar{b} \rightarrow hS \rightarrow h\gamma\gamma \) process, normalised to unity, for two width scenarios, \( \Gamma = 1, \) 45 GeV and two values of the portal coupling, \( \lambda_{HS} = 0, \) 1, for \( M = 750 \) GeV.

\( \Gamma = 1 \text{ GeV}, \) with the on-shell di-photon region dominating. For \( \lambda_{HS} = 0 \), the “three-body decay” region disappears completely since the resonant s-channel diagram of Fig. 2 (a) vanishes. For large width the two regions merge into one and the effect of the vanishing three-body decay region for \( \lambda_{HS} = 0 \) is not as evident as in the case of small width. The distributions for the \( ss \)-initiated process exhibit similar features, with different “mixtures” between the two regions arising from the differences between the strange and bottom quark parton density functions. We omit them for the sake of simplicity.

FIG. 10: The matrix-element level distribution of the combined Higgs boson and di-photon invariant mass, \( M_{h\gamma\gamma} \), in the \( gg \rightarrow hS \rightarrow h\gamma\gamma \) process, normalised to unity, for three benchmark mass scenarios, \( M = 600, 750, 900 \) GeV, and \( \Gamma = 1 \) GeV, for \( \lambda_{HS} = 1.0 \).

FIG. 11: The matrix-element level distribution of the combined Higgs boson and di-photon invariant mass, \( M_{h\gamma\gamma} \), in the \( b\bar{b} \rightarrow hS \rightarrow h\gamma\gamma \) process, normalised to unity, for three benchmark mass scenarios, \( M = 600, 750, 900 \) GeV, and \( \Gamma = 1 \) GeV, for \( \lambda_{HS} = 1 \).

that we will consider as “benchmark” scenarios in our analysis, \( M = 600, 750, 900 \) GeV (see below) and \( \Gamma = 1 \) GeV. For the \( b\bar{b} \) case we only show the \( \lambda_{HS} = 1 \) distributions for simplicity. They all clearly demonstrate the existence of the main features described for the \( M = 750 \) GeV case, i.e. the “three-body decay” and “on-shell di-photon” regions.

III. EVENT GENERATION AND DETECTOR SIMULATION.

A. Event generation.

The signal model was generated via an implementation of the Lagrangian of Eq. 1 in FeynRules \([49, 50]\). Via the UFO interface \([51]\) this was used to generate parton-level events employing MadGraph5_aMC@NLO \([52, 53]\). The background processes were also generated using MadGraph5_aMC@NLO, with appropriate generation-level cuts to reduce the initial cross sections to a manageable level. All the events were passed through the HERWIG 7 \([54, 55]\) Monte Carlo for simulation of the parton shower, the underlying event and hadronization. As before, the MMHT14nlo68cl PDF set was employed. To remain conservative, we consider collisions at the LHC at a centre-of-mass energy of 13 TeV. The possible increase of energy to 14 TeV will increase rates in the considered processes by \( \mathcal{O}(10\%) \).

Since we expect to impose cuts on the di-photon mass window, \( M_{\gamma\gamma} \), that are sufficiently far away from the Higgs boson resonance, we can immediately exclude any background processes containing \( h \rightarrow \gamma\gamma \) from the analysis. For this reason we do not include associated Higgs boson production with a vector boson or Higgs boson pair production, \( t\bar{t}h \) production, and so on. This implies that
the relevant backgrounds are those with non-resonant $\gamma\gamma$ production, other processes that involve $S \to \gamma\gamma$, and reducible backgrounds. We thus consider the following processes: $\gamma+$-jets, $\gamma\gamma+$-jets, events with at least one true $b$-quark at parton-level ($b+$-jets), $Z\gamma\gamma$ with $Z \to bb$, $t\bar{t}\gamma\gamma$ including all the decay modes of the top quarks and the production of the resonance $S$ in association with a non-resonant $bb$ pair. All the multi-jet processes are generated without merging to the parton shower, in the five-flavour scheme, with four outgoing partons at the matrix-element level.

The calculation of higher-order QCD corrections to these multi-leg processes, particularly when restricting the phase space with cuts, is numerically challenging at present. To remain conservative, we will assume that the corrections are large and apply $K$-factors of $K = 2$ to all the background processes. For the signal and the $bbS$ associated production we do not apply any $K$-factors since the corrections are approximately absorbed into the ratio with the single inclusive production of the $S$ resonance, see below. Throughout this article we assume that $\sigma(pp \to S \to \gamma\gamma) = 10$, 5, 1 fb, corresponding to the benchmark masses $M = 600$, 750, 900 GeV, fixing the product $c_\rho^Sc_\rho^S$ (or $(y_d^S)^2c_\rho^S$), which drops out in the ratio $\rho$. The values of the cross sections are motivated by the current ATLAS [16] and CMS [17] limits on di-photon resonances.

Note that it turns out that the non-resonant $bbS$ process is only relevant for gluon-fusion production of $S$, and we only report numbers for that in what follows.

### B. Detector simulation.

In the hadron-level analysis that follows, performed without using any dedicated detector simulation software, we consider all particles within a pseudo-rapidity of $|\eta| < 5$ and $p_T > 100$ MeV. We smear the momenta of all reconstructed objects (i.e. jets, electrons, muons and photons) according to HL-LHC projections [60–62]. We also apply the relevant reconstruction efficiencies. We simulate $b$-jet tagging by looking for jets containing $B$-hadrons, that we have set to stable in the simulation, and considering them as the $b$-jet candidates. The mis-tagging of $c$-jets to $b$-jets is performed by choosing $c$-jet candidates (after hadronization) as those jets that lie within a distance $\Delta R < 0.4$ from $c$-quarks (after the parton shower), with transverse momentum $p_T > 1$ GeV. We apply a flat $b$-tagging efficiency of 70% and a mis-tag rate of 1% for light-flavour jets and 10% for charm-quark-initiated jets.

We reconstruct jets using the anti-$k_t$ algorithm available in the FastJet package [62–63], with a radius parameter of $R = 0.4$. We only consider jets, photons and leptons with $p_T > 30$ GeV within pseudo-rapidity $|\eta| < 2.5$ in our analysis. The jet-to-lepton mis-identification probability is taken to be $P_{\gamma \to \ell} = 0.0048 \times e^{-0.035p_T/\text{GeV}}$ and the jet-to-photon mis-identification probability was taken to be $P_{\gamma \to \gamma} = 0.0093 \times e^{-0.036p_T/\text{GeV}} [63–65]$, both flat in pseudo-rapidity. We demand all leptons and photons to be isolated, where an isolated object is defined to have $\sum_{j} p_{T,j} < 15\%$ of its transverse momentum in a cone of $\Delta R = 0.2$ around it.

### IV. Detailed Analysis.

We consider events with two reconstructed $b$-jets and two isolated photons as defined in Section [11]. Note that this final state has been previously considered in the context of searches for Higgs boson pair production, e.g. in [67–68]. We impose the following ‘acceptance’ cuts to all samples:

- $b$-jets: transverse momenta $p_{T,1} > 30$ GeV, $p_{T,2} > 30$ GeV, all $b$-jets within $|\eta| < 2.5$.
- photons: transverse momenta $p_{T,1} > 30$ GeV, $p_{T,2} > 30$ GeV, all photons within $|\eta| < 2.5$.
- invariant mass of the two $b$-jets $M_{bb} \in [90, 160]$ GeV.
- invariant mass of the two photons $M_{\gamma\gamma} > M - 300$ GeV.
- veto events with leptons of $p_T > 25$ GeV within $|\eta| < 2.5$,

for each of the considered di-photon resonance masses, $M$.

The cross sections after application of the acceptance cuts are given in Table I for two values of the widths $\Gamma = 1$ GeV and $\Gamma = 45$ GeV and for $M = 750$ GeV. For the case of $q\bar{q}$ we consider as examples $\lambda_{HS} = 1$ and $\lambda_{HS} = 0$. Throughout this analysis, the total signal cross section was calculated by using the ratio $\rho$ (derived in Section [11]) as

$$\sigma(pp \to hS \to h\gamma\gamma) = \rho \times \sigma(pp \to S \to \gamma\gamma),$$

where $\sigma(pp \to S \to \gamma\gamma) = 10$, 5, 1 fb for $M = 600$, 750, 900 GeV, and including the decay $h \to bb$. This cross section was employed as the normalisation of the signal event samples (before analysis cuts). The expected number of signal events, for $\lambda_{HS} \sim O(1)$, after acceptance cuts is $O(10)$ at $3000$ fb$^{-1}$ of integrated luminosity. However, as already discussed, one should keep in mind that the cross section grows with $\lambda_{HS}^2$ in both $gg$- and $q\bar{q}$-initiated production.

The resulting di-photon invariant mass after acceptance cuts is shown in Fig. [12] for the example of $M = 750$ GeV. The $M_{\gamma\gamma}$ observable can be used to separate the analysis into the two regions described in Section [11].

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55 We also considered the $h\gamma\gamma$ process, including the loop-induced pieces [59], but found that it possess a negligible cross section.

56 Fixing the product $c_\rho^Sc_\rho^S$ (or $(y_d^S)^2c_\rho^S$), which drops out in the ratio $\rho$. The values of the cross sections are motivated by the current ATLAS [16] and CMS [17] limits on di-photon resonances.

57 We assume that the non-resonant $bbS$ process is only relevant for gluon-fusion production of $S$, and we only report numbers for that in what follows.

58 Note that it turns out that the non-resonant $bbS$ process is only relevant for gluon-fusion production of $S$, and we only report numbers for that in what follows.

59 This procedure of associating jets to $c$-quarks is expected to be conservative.
the “three-body decay” region (“TBD”) and the “on-shell di-photon” region (“OS$$\gamma\gamma$$”). The separation is identical in both $$gg$$- and $$qq$$-initiated processes. We choose: $$M_{\gamma\gamma} < M - 50$$ GeV for the “TBD” region and $$M_{\gamma\gamma} > M - 50$$ GeV for the “OS$$\gamma\gamma$$” region for a di-photon resonance mass, $$M$$. We also show the distribution of the combined invariant mass of the two b-jet candidates and the di-photon system, $$M_{bb\gamma\gamma}$$ in Fig. 13 which also clearly demonstrates the existence of the two regions.

We apply further cuts to improve signal and background discrimination. As we did not attempt to fully optimize the cuts in the present analysis, we apply a common set of cuts along with invariant mass cuts on the observables $$M_{\gamma\gamma}$$ and $$M_{bb\gamma\gamma}$$ that provide the main distinction between the two regions. The common cuts applied in each region are shown in Table II and the specific invariant mass cuts are shown in Table III. Effectively, the cuts aim to exploit the fact that the photons in the signal are harder than in the backgrounds and also feature tighter di-photon and $$bb$$ mass windows, particularly in the “OS$$\gamma\gamma$$” region for the former.

![FIG. 12: The distribution of the di-photon invariant mass, $$M_{\gamma\gamma}$$, for the $$gg \rightarrow hS \rightarrow bb\gamma\gamma$$ process after acceptance cuts, normalised to unity, for two width scenarios, $$\Gamma = 1$$ GeV and $$\Gamma = 45$$ GeV. All branching ratios, acceptances and tagging rates have been applied. We have assumed that the single production cross section for a di-photon scalar resonance of $$M = 750$$ GeV is $$\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 5$$ fb.](image)

![FIG. 13: The distribution of the combined di-photon + $$bb$$ invariant mass, $$M_{bb\gamma\gamma}$$, for the $$gg \rightarrow hS \rightarrow bb\gamma\gamma$$ process after acceptance cuts, normalised to unity, for two width scenarios, $$\Gamma = 1$$ GeV and $$\Gamma = 45$$ GeV, while $$M = 750$$ GeV.](image)

![TABLE II: The additional common cuts applied along with the acceptance cuts in both the “three-body decay” and the “on-shell di-photon” region for the $$gg$$- and $$qq$$-initiated processes. The labels “1” and “2” correspond to the hardest and second hardest reconstructed objects (photons or b-jets), respectively.](table)

| observable | cut |
|------------|-----|
| $$P_T^{\gamma_1}$$ | > 200 GeV |
| $$P_T^{\gamma_2}$$ | > 120 GeV |
| $$\Delta R(\gamma, \gamma)$$ | $$\in [2.0, 4.0]$$ |
| $$M_{bb}$$ | $$\in [100, 150]$$ GeV |
| $$\Delta R(b, b)$$ | $$\in [0.8, 3.0]$$ |
| $$\Delta R(\gamma\gamma, h_2)$$ | < 3.0 |

We show the resulting cross sections after the application of these further cuts in Table IV for the case of $$M = 750$$ GeV. A high efficiency is maintained for the signal, with high rejection factors for the background processes. We note again that the $$bbS$$ associated production process is relevant only for the gluon-fusion scenario.

To obtain the 95% confidence-level exclusion regions for $$\lambda_{HS}$$ we use Poissonian statistics to calculate the probabilities. Since we have assumed that the production of a scalar di-photon resonance will have been observed, we have to construct a null hypothesis compatible with such an observation providing the expected number of events.
mass windows were made according to the width of the resonance, \(\Gamma\), of the photon resonance of mass \(gg\) single production cross section for a di-photon scalar resonance of mass \(M\). The different choices for the mass windows were made according to the width of the resonance, \(\Gamma\).

| \(\Gamma = 1\) GeV | “TBD” | “OS-\(\gamma\gamma\)” |
|----------------------|--------|----------------------|
| \(M_{h}\) | \(\in M \sim 10^{10}\) GeV | \(\in M \pm 5\) GeV |
| \(M_{b\bar{b}(\gamma\gamma)}\) | \(\in M \pm 30\) GeV | - |

| \(\Gamma = 45\) GeV | “TBD” | “OS-\(\gamma\gamma\)” |
|----------------------|--------|----------------------|
| \(M_{h}\) | \(\in M \sim 10^{10}\) GeV | \(\in M \pm 40\) GeV |
| \(M_{b\bar{b}(\gamma\gamma)}\) | \(\in M \pm 40\) GeV | - |

TABLE III: The additional invariant mass cuts applied along with the further cuts of Table II in the “three-body decay” region (“TBD”) and the “on-shell di-photon” region (“OS-\(\gamma\gamma\)”), for both the \(gg\)- and \(qq\)-initiated processes for a scalar di-photon resonance of mass \(M\). The different choices for the mass windows were made according to the width of the resonance, \(\Gamma\).

| process | “TBD” | “OS-\(\gamma\gamma\)” |
|---------|--------|----------------------|
| \(gg \rightarrow h(bb)S(\gamma\gamma), \Gamma = 1\) GeV, \(\lambda_{HS} = 1\) | \(0.00007\) | \(0.00020\) |
| \(gg \rightarrow h(bb)S(\gamma\gamma), \Gamma = 45\) GeV, \(\lambda_{HS} = 1\) | \(0.00006\) | \(0.00014\) |
| \(bb \rightarrow h(bb)S(\gamma\gamma), \Gamma = 1\) GeV, \(\lambda_{HS} = 1\) | \(0.00007\) | \(0.000105\) |
| \(bb \rightarrow h(bb)S(\gamma\gamma), \Gamma = 45\) GeV, \(\lambda_{HS} = 1\) | \(0.000038\) | \(0.000074\) |
| \(ss \rightarrow h(bb)S(\gamma\gamma), \Gamma = 1\) GeV, \(\lambda_{HS} = 1\) | \(0.000007\) | \(0.00000005\) |
| \(ss \rightarrow h(bb)S(\gamma\gamma), \Gamma = 45\) GeV, \(\lambda_{HS} = 1\) | \(0.000005\) | \(0.000107\) |

at least one b-quark + jets \(\mathcal{O}(10^{-6})\) \(\mathcal{O}(10^{-6})\)

| process | “TBD” | “OS-\(\gamma\gamma\)” |
|---------|--------|----------------------|
| \(gg \rightarrow b\bar{b}S(\gamma\gamma), \Gamma = 1\) GeV | \(0.00025\) | \(0.00075\) |
| \(gg \rightarrow b\bar{b}S(\gamma\gamma), \Gamma = 45\) GeV | \(0.00003\) | \(0.00018\) |
| \(gg \rightarrow b\bar{b}S(\gamma\gamma), \Gamma = 75\) GeV | \(0.000025\) | \(0.000075\) |

TABLE IV: The expected cross sections at 13 TeV \(pp\) collision energy in fb for all the considered processes after the application of further cuts as in Tables II and III for \(M = 750\) GeV and \(\Gamma = 1, 45\) GeV. All branching ratios, acceptances and tagging rates have been applied. We have assumed that the single production cross section for a di-photon scalar resonance of \(M = 750\) GeV is \(\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 5\) fb.

“Weisser method [70], where the combined significance, \(\Omega\), is given, in terms of the individual significances \(\Omega_i\), as:

\[
\Omega = \frac{1}{\sqrt{k}} \sum_{i=1}^{k} \Omega_i .
\]

We show the resulting expected limits (assuming our null hypothesis is true) as a function of the integrated luminosity for the different benchmark scenarios that we consider in Figs. [14, 25]. For the case \(M = 750\) GeV we show results for \(\Gamma = 45\) GeV as well. For \(\Gamma = 1\) GeV, we obtain more stringent constraints, limiting, for \(M = 600, 750, 900\) GeV respectively, \(|\lambda_{HS}| \lesssim 2, 4, 5\) for the \(gg\)-initiated process and \(\lambda_{HS} \in [\sim 4, \sim 1, [\sim 7, \sim 3, [\sim 8, \sim 4]\), both for the \(bb\) and \(ss\)-initiated processes, at the end of the HL-LHC run (3000 fb\(^{-1}\)). The variation between the results for \(bb\) and \(ss\)-initiated processes is visible in the plots – it is very small and can be attributed to the differences between the parton density functions for the strange and bottom quarks, as already mentioned.

The scenario with the larger width, \(M = 750\) GeV, \(\Gamma = 45\) GeV, clearly exhibits weaker constraints, with the \(gg\)-initiated processes yielding \(|\lambda_{HS}| \lesssim 5\), the \(bb\)-initiated process \(\lambda_{HS} \in [\sim 5]\) and the \(ss\)-initiated process \(\lambda_{HS} \in [\sim 4]\). It is conceivable that if further decay channels of the resonance \(S\) are discovered, the remaining unconstrained regions in the \(qq\) cases can be covered (in particular for narrow width), allowing determination of the initial state partons that produce the resonance.

The lower bound on \(\lambda_{HS}\) for the \(qq\)-initiated cases is driven by the “TBD” region. This is understood by the fact that the “TBD” region always vanishes near \(\lambda_{HS} \sim 0\), as it is dominated by the diagram with an on-shell s-channel \(S\), making the exclusion region resulting from it symmetric with respect to \(\lambda_{HS} \sim 0\), whereas the “OS-\(\gamma\gamma\)” region possesses a symmetry point somewhere in \(\lambda_{HS} < 0\).

Mixed production of the di-photon resonance

So far we have investigated production of the scalar resonance initiated purely either via gluon fusion or \(qq\) annihilation. We can generalize this to “mixed” production via \(gg\) and \(qq\) initial states simultaneously. We concentrate on the scenario of \(gg\) and \(bb\) for simplicity, with the extension to additional quark flavours being straightforward. In that case, the ratio of cross sections, \(\rho_{mixed}\), defined between the associated production and single production modes can be written as:

\[
\rho_{mixed} = \frac{\sigma(bb \rightarrow hS \rightarrow h\gamma\gamma) + \sigma(gg \rightarrow hS \rightarrow h\gamma\gamma)}{\sigma(bb \rightarrow S \rightarrow \gamma\gamma) + \sigma(gg \rightarrow S \rightarrow \gamma\gamma)} = \frac{B_2(\lambda_{HS})(|y_2^S|^3)^2 + G_2(\lambda_{HS})(c_G^S)^2}{B_1(|y_1^S|^3)^2 + G_1(c_G^S)^2},
\]

where \(B_2(\lambda_{HS}), G_2(\lambda_{HS})\) are functions of the portal coupling \(\lambda_{HS}\) and \(B_1\), \(G_1\) are constants with respect
FIG. 14: The grey-shaded area shows the 95% confidence-level exclusion region for the portal coupling \( \lambda_{HS} \) for the \( gg \)-induced case and \( M = 600 \) GeV, \( \Gamma = 1 \) GeV, for the combination of the “TBD” (red) and “OS\( \gamma\gamma \)” (green) regions as defined in the main text. We have assumed that the single production cross section \( \sigma(pp \to S \to \gamma\gamma) = 10 \) fb.

FIG. 15: The grey-shaded area shows the 95% confidence-level exclusion region for the portal coupling \( \lambda_{HS} \) for the \( b\bar{b} \)-induced case and \( M = 600 \) GeV, \( \Gamma = 1 \) GeV, for the combination of the “TBD” (red) and “OS\( \gamma\gamma \)” (green) regions as defined in the main text. We have assumed that the single production cross section \( \sigma(pp \to S \to \gamma\gamma) = 10 \) fb.

FIG. 16: The grey-shaded area shows the 95% confidence-level exclusion region for the portal coupling \( \lambda_{HS} \) for the \( s\bar{s} \)-induced case and \( M = 600 \) GeV, \( \Gamma = 1 \) GeV, for the combination of the “TBD” (red) and “OS\( \gamma\gamma \)” (green) regions as defined in the main text. We have assumed that the single production cross section \( \sigma(pp \to S \to \gamma\gamma) = 10 \) fb.

FIG. 17: The grey-shaded area shows the 95% confidence-level exclusion region for the portal coupling \( \lambda_{HS} \) for the \( gg \)-induced case and \( M = 750 \) GeV, \( \Gamma = 1 \) GeV, for the combination of the “TBD” (red) and “OS\( \gamma\gamma \)” (green) regions as defined in the main text. We have assumed that the single production cross section \( \sigma(pp \to S \to \gamma\gamma) = 5 \) fb.

To the portal coupling, all to be determined. Defining \( \theta \equiv |(y^S_{d3})^3|/c_S^G \), the above expression can be re-written as:

\[
\rho_{\text{mixed}} = \frac{B_2(\lambda_{HS})\theta^2 + G_2(\lambda_{HS})}{B_1\theta^2 + G_1}.
\]

(10)

By considering the limits \( \theta \to 0 \) and \( \theta \to \infty \), and dividing the numerator and denominator by \( G_1 \), we can deduce that

\[
\rho_{\text{mixed}} = \frac{\rho(b\bar{b}, \lambda_{HS})r_{bg}\theta^2 + \rho(gg, \lambda_{HS})}{r_{bg}\theta^2 + 1},
\]

(11)

where \( \rho(b\bar{b}, \lambda_{HS}) \) and \( \rho(gg, \lambda_{HS}) \) are the ratios of cross sections as functions of the portal coupling as defined in Eq. [7] for the cases of pure production initiated either via \( b\bar{b} \) or \( gg \), respectively, and \( r_{bg} \) is the ratio of pure single production cross sections for \( \theta = 1 \): \( r_{bg} = B_1/G_1 \). The former two functions have already been determined in the analysis of the pure cases. The constant \( r_{bg} \) was calculated explicitly via Monte Carlo to be \( r_{bg} \approx 1.02, 0.69, 0.49 \) for \( M = 600, 750, 900 \) GeV respectively. The values are approximately equal for both the narrow width case (\( \Gamma = 1 \) GeV) and the larger width case (\( \Gamma = 45 \) GeV).
FIG. 18: The grey-shaded area shows the 95% confidence-level exclusion region for the portal coupling $\lambda_{HS}$ for the $gg$-induced case and $M = 750$ GeV, $\Gamma = 45$ GeV, for the combination of the “TBD” (red) and “OS$\gamma\gamma$” (green) regions as defined in the main text. We have assumed that the single production cross section $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 5$ fb.

Using Eq. [11] we can deduce an expression for the predicted number of signal events after the application of analysis cuts in the mixed production case:

$$N_{\text{mixed}} = \frac{r_{bg}\theta^2}{r_{bg}\theta^2 + 1}N(bb, \lambda_{HS}) + \frac{1}{r_{bg}\theta^2 + 1}N(gg, \lambda_{HS}),$$

(12)

where $N(bb, \lambda_{HS})$ and $N(gg, \lambda_{HS})$ are the expected signal events for either pure $bb$ or pure $gg$ production for a given value of $\lambda_{HS}$, at a specific integrated luminosity.

The predicted number of background events after a given set of analysis cuts is constant with $\theta$, apart from the associated production of the di-photon resonance with a $bb$ pair, which was found to be significant only in the $gg$-initiated scenario. This background scales as $N_{\text{assoc.}} = N_{\text{assoc.},0}/(r_{bg}\theta^2 + 1)$, where $N_{\text{assoc.},0}$ is the expected number of events for the $bbS$ associated production after analysis cuts in the case of pure $gg$-initiated production.

Using the expression of Eq. [12] and the event numbers for the two analysis regions “TBD” and “OS$\gamma\gamma$” obtained for the pure production modes, we can derive constraints on the $(\theta, \lambda_{HS})$-plane. These are shown in Figs. 26-28 and 29 for masses $M = 600$, 750, 900 GeV and $\Gamma = 1$ GeV and for $M = 750$ GeV and $\Gamma = 45$ GeV,
FIG. 22: The grey-shaded area shows the 95% confidence-level exclusion region for the portal coupling $\lambda_{HS}$ for the $s\bar{s}$-induced case and $M = 750$ GeV, $\Gamma = 45$ GeV, for the combination of the “TBD” (red) and “OS$\gamma\gamma$” (green) regions as defined in the main text. We have assumed that the single production cross section $\sigma(pp \to S \to \gamma\gamma) = 5$ fb.

FIG. 23: The grey-shaded area shows the 95% confidence-level exclusion region for the portal coupling $\lambda_{HS}$ for the $gg$-induced case and $M = 900$ GeV, $\Gamma = 1$ GeV, for the combination of the “TBD” (red) and “OS$\gamma\gamma$” (green) regions as defined in the main text. We have assumed that the single production cross section $\sigma(pp \to S \to \gamma\gamma) = 1$ fb.

FIG. 24: The grey-shaded area shows the 95% confidence-level exclusion region for the portal coupling $\lambda_{HS}$ for the $b\bar{b}$-induced case and $M = 900$ GeV, $\Gamma = 1$ GeV, for the combination of the “TBD” (red) and “OS$\gamma\gamma$” (green) regions as defined in the main text. We have assumed that the single production cross section $\sigma(pp \to S \to \gamma\gamma) = 1$ fb.

FIG. 25: The grey-shaded area shows the 95% confidence-level exclusion region for the portal coupling $\lambda_{HS}$ for the $s\bar{s}$-induced case and $M = 900$ GeV, $\Gamma = 1$ GeV, for the combination of the “TBD” (red) and “OS$\gamma\gamma$” (green) regions as defined in the main text. We have assumed that the single production cross section $\sigma(pp \to S \to \gamma\gamma) = 1$ fb.

respectively, at an integrated luminosity of 3000 fb$^{-1}$ at 13 TeV. It can be seen that in the limits $\theta \to 0$ and $\theta \gg 1$, corresponding to $gg$- or $b\bar{b}$-dominated production respectively, one can recover the pure production constraints of Figs. [14][25] obtained at 3000 fb$^{-1}$.

V. DISCUSSION AND CONCLUSIONS.

We have investigated the associated production of a diphoton scalar resonance with a Higgs boson and have employed the $pp \to hS \to (b\bar{b})(\gamma\gamma)$ final state to obtain constraints on the portal coupling with the SM Higgs boson $\lambda_{HS}$, at the LHC. We have considered three benchmark scalar masses, $M = 600$, 750, 900 GeV, and we have assumed that the inclusive single production cross section is $\sigma(pp \to S \to \gamma\gamma) = 10, 5, 1$ fb, respectively, compatible with current experimental constraints. To construct expected constraints we considered the null hypothesis (i.e. the supposed ‘true’ underlying theory) to correspond to gluon-fusion-initiated production with vanishing portal coupling, $\lambda_{HS} = 0$. We then first analysed the case of either pure gluon-fusion-induced production or production via quark-anti-quark annihilation. For gluon-fusion
production one can impose constraints on the portal coupling at the end of the HL-LHC run of $|\lambda_{HS}| \lesssim 2, 4, 5$ for $M = 600, 750, 900$ GeV, respectively, and $\Gamma = 1$ GeV, while $|\lambda_{HS}| \lesssim 5$ in the case of large width ($\Gamma = 45$ GeV) and $M = 750$ GeV. For quark-anti-quark annihilation, the production of an on-shell scalar and the Higgs boson is enhanced by a contact interaction $\sim q\bar{q}hS$, originating from the same $D = 5$ operator that mediates single $S$ production. This implies that the cross section is non-negligible even for vanishing portal coupling $\lambda_{HS} = 0$. This fact can be exploited to exclude the whole plane of $\lambda_{HS}$, thus excluding the hypothesised production via $q\bar{q}$ annihilation. For the case of $b\bar{b}$ we find that a region inside $\lambda_{HS} \in [\sim -4, \sim 1], [\sim -7, \sim 3], [\sim -8, \sim 4]$ for $M = 600, 750, 900$ GeV, respectively, could remain unconstrained at the end of the HL-LHC in the narrow width scenario, while $\lambda_{HS} \in [\sim -8, \sim 5]$ for large width and $M = 750$ GeV. For the case of $s\bar{s}$ annihilation we find the same unconstrained regions to good approximation for narrow width, while we obtain $\lambda_{HS} \in [\sim -8, \sim 4]$ for
in the case of large width, where \( M = 750 \text{ GeV} \). We have also considered ‘mixed’ production via gluon fusion and \( bb \) annihilation and derived constraints on the plane of ratio of the corresponding couplings versus \( \lambda_{HS} \) for an integrated luminosity of 30000 \( \text{fb}^{-1} \).

Note that since a measurement of \( \lambda_{HS} \) can straightforwardly be translated into a measurement of \( \mu_S \), these numbers suggest that it is possible to exclude, for example, the scenario where the mass of \( S \) stems from EWSB (\( \mu_S = 0 \)), which corresponds to \( \lambda_{HS} = M^2/v^2 \approx 5.9, 9.3, 13.4 \) (still assuming linear terms in \( S \) to vanish).\(^{*}^{*}^{*} \) Furthermore, if additional decay modes of the resonance \( S \) are discovered beyond \( \gamma \gamma \), it is conceivable that a combined analysis in various channels would be able to exclude all possible values of \( \lambda_{HS} \) - for the case of \( qq \) production - thus providing an independent method of determining the production mode. Conversely, the analysis of the present article demonstrates that if the production mechanism is constrained via alternative means, it will be possible to obtain meaningful constraints on the interaction of a new scalar resonance and the Higgs boson, allowing determination of its relation to electroweak symmetry breaking.

Acknowledgments

We would like to thank Peter Richardson, Jernej Kamnik, José Zurita, Tania Robens and Eleni Vryonidou for useful comments and discussion. AC acknowledges support by a Marie Sklodowska-Curie Individual Fellowship of the European Community’s Horizon 2020 Framework Programme for Research and Innovation under contract number 659239 (NP4theLHC14). The research of F.G. is supported by a Marie Curie Intra European Fellowship within the 7th European Community Framework Programme (grant no. PIEF-GA-2013-628224). AP acknowledges support by the MCnetITN FP7 Marie Curie Initial Training Network PITN-GA-2012-315877 and a Marie Curie Intra European Fellowship within the 7th European Community Framework Programme (grant no. PIEF-GA-2013-622071).

Appendix A: The full set of operators.

In this appendix, we provide the full set of operators for the SM in the presence of an additional dynamical scalar singlet \( S \), up to \( D = 5 \). The fact that \( S \) does not transform under the SM gauge group makes the construction straightforward: in the case that it is CP even (with CP-conserving interactions), each additional operator will consist of a gauge invariant SM operator, multiplied by powers of \( S \) (and potentially derivatives).\(^{†††} \) This means that in particular that, schematically, \( \cal{L} \supset \cal{L}_{SM} + \cal{L}_{SM} \cdot S \).

If all operators in the SM would be marginal, \( i.e. \), feature \( D = 4 \), this would already correspond to the full Lagrangian at \( D \leq 5 \), connecting the SM with the new resonance. In turn, the only operators that are missing (up to pure NP terms) are those containing the single relevant \( i.e. \), \( D < 4 \) operator of the SM, \( |H|^2 \). This can be multiplied by up to three powers of \( S \) as well as by \( \partial^2 S \) (the latter being equivalent via integration by parts to \( S \partial^2 |H|^2 \)). This finally leads to the Lagrangian \( ^{†††} \)

\[
\begin{align*}
\cal{L}_{\text{eff}} &= \cal{L}_{\text{SM}} + \frac{\lambda_S}{2\sqrt{2}} v S^3 - \frac{\lambda_{HS} v}{4} |H|^2 S - \lambda_{HS} |H|^2 S^2 \\
&\quad - \frac{\lambda_S}{2} \partial^2 S - \frac{\lambda_{HS} v}{4} |H|^2 S - \lambda_{HS} |H|^2 S^2 \\
&\quad - \frac{S}{\Lambda} \left[ c_{\lambda S} S^4 + c_{\lambda HS} |H|^2 S^2 + c_{\lambda H} |H|^4 \right] \\
&\quad - (g_d^{S,i})^2 \frac{S}{\Lambda} Q_L^i H d_R^i - (g_u^{S,i})^2 \frac{S}{\Lambda} Q_L^i \bar{H} q_R^i + \text{h.c.} \\
&\quad - \frac{S}{\Lambda} \frac{1}{16\pi^2} \left[ g^2 c_G^B B_{\mu \nu} B^{\mu \nu} + g^2 c_G^W W^{\mu \nu} W^{\mu \nu} \right. \\
&\quad \left. + g_s^2 c_G^S G^{\mu \nu} G_{\mu \nu} \right]. \quad (A1)
\end{align*}
\]

Note that we have used equations of motion/field redefinitions to eliminate the operators \( S/\Lambda|\partial^2 S|^2, S/\Lambda |D_{\mu} H|^2, S/\Lambda \partial^2 |H|^2 \) and those of type \( S/\Lambda q D_{\mu} \gamma^\mu q \). The couplings in the third line of (A1) and \( \lambda_{S,HS} \), which are those that have been neglected in (1), do not contribute to the process under consideration to good approximation: They either do not enter at leading order, or (in the case of the \( |H|^2 S \) interactions) at most contribute to a diagram with a (strongly suppressed) off-shell Higgs boson propagator, discussed in Appendix B.

The same holds true for the \( S^4 \) term, that was kept for the discussion in Section II.

Appendix B: Fits for the ratio \( \rho \) including a linear term.

In addition to the diagrams of Figs. I and II, there can be contributions to the final state under consideration originating from \( hhS \) interactions (for example, an intermediate \( s \)-channel Higgs boson to \( hS \) or an \( S \rightarrow hh \) intermediate state), provided the following triple scalar

\[ \rho \]

\[ ^{†††} \] We neglect the only \( D = 5 \) operator consisting only of SM building blocks – the Weinberg operator – since the tiny neutrino masses suggest that it is suppressed by a very large scale \( \Lambda \sim M_{\text{GUT}} \).

\[ ^{†††} \] See also [49-72]. We do not include leptons, as they do not enter the examined processes to the order considered. Moreover, note that (induced) tadpoles in \( S \) can always be removed via a field redefinition.
interaction, linear in S, is non-vanishing:
\[-\kappa H_S v^2 S \subset -\lambda'_{H S} v |H|^2 S - c_{\lambda H} \frac{S}{\Lambda} |H|^4, \quad (B1)\]
with \(\kappa_{HS} = 1/2\lambda'_{HS} + 3/(2\Lambda)c_{\lambda H} H\).

In this scenario, terms will appear in the cross section proportional to \(\kappa_{HS}\) and \(\kappa'_{HS}\). Due to interference with the new \(hH\) diagrams (note that also the ‘non-signal’ \(S \to hh\) interferes with the signal diagrams after an off-shell \(h \to \gamma\gamma\)), some dependence on the coupling of the resonance \(S\) to photons, \(c_S^2\), as well as the production couplings, \((y^S_d)_{22}, (y^S_d)_{33}\) and \(c^S_G\), is introduced in the denominator of the cross section ratio. Due to this, smaller values of \(|c^S_H|\) and the production couplings enhance the effect of the new contributions proportional to \(\kappa_{HS}\) and \(\kappa'_{HS}\). To take these effects into account and at the same time remain conservative, we set \(|c^S_H|\) to the possible minimal value that produces a \(\sigma(pp \to S \to \gamma\gamma) \sim 5\) fb for \(M = 750\) GeV, as derived in Ref. \[24\]. This leads to \(|c^S_H| \geq \{60, 100, 10\}\) for \(ss, bb\) and \(gg\) production, respectively, where here and in the following we assume a normalisation of \(\Lambda = 1\) TeV. Moreover, one can derive conservative lower bounds on \(|(y^S_d)_{22}|, |(y^S_d)_{33}|\) and \(|c^S_G|\), by demanding \(\sigma(pp \to S) \geq 5\) fb, leading to \(|(y^S_d)_{22}| \gtrsim 0.15, |(y^S_d)_{33}| \gtrsim 0.2\) and \(|c^S_G| \gtrsim 0.25\). We set these to their minimal values as well in what follows.

We can then parametrize the ratio \(\rho(xx') = \sigma(xx' \to h\gamma\gamma)/\sigma(xx' \to S \to \gamma\gamma)\), \(xx' = bb, ss, gg\), by\(^{95}^{96}\)
\[\rho(xx') = \delta_1^S + \delta_2^S \kappa_{HS} \lambda_{HS} + \delta_3^S \kappa_{HS}^2 \lambda_{HS} + \delta_4^S \kappa_{HS}^2 \lambda_{HS}^2, \quad (B2)\]
where \(\delta_i^{b,s,g}\) are coefficients to be determined.

\[
\begin{array}{ccc}
\hline
i & \delta_i^b & \delta_i^s & \delta_i^g \\
\hline
1 & 8.27 \times 10^{-3} & 1.05 \times 10^{-4} & < O(10^{-5}) \\
2 & 1.93 \times 10^{-4} & -2.30 \times 10^{-4} & O(10^{-6}) \\
3 & 8.26 \times 10^{-4} & 2.54 \times 10^{-5} & O(10^{-6}) \\
4 & 3.08 \times 10^{-3} & 3.71 \times 10^{-4} & O(10^{-5}) \\
5 & 2.40 \times 10^{-5} & 6.00 \times 10^{-4} & O(10^{-5}) \\
6 & 8.60 \times 10^{-4} & 9.27 \times 10^{-4} & 1.33 \times 10^{-3} \\
\hline
\end{array}
\]

TABLE V: The \(\rho\) fit coefficients for the general case of including an interaction term proportional to \(\mathcal{L} \supset -\kappa H_S v^2 S\) for width \(\Gamma = 1\) GeV and \(M = 750\) GeV.

An estimate of the coefficients in this conservative scenario, obtained numerically, is given in Table\[\]\[\] for width \(\Gamma = 1\) GeV and \(M = 750\) GeV. The contributions from diagrams with an off-shell Higgs boson are a priori not fully negligible for the \(bb\)-induced case, simply because of interference of the \(s\)-channel \(h \to hS\) with the large matrix elements with the on-shell \(S\) in the final state.

\(^{95}^{96}\)For the case of vanishing \(\kappa_{HS}\), this more general definition coincides with Eq.\[7\] to good approximation.

\(^{95}^{96}\)Here, we have neglected the impact of the second operator in \(\mathcal{B1}\), \(S|H|^4\), which breaks the correlation between Higgs-scalar mixing and the \(h^2S\) interaction.
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