Exotic colored scalars at the LHC

Kfir Blum, Aielet Efrati, Claudia Frugiuele and Yosef Nir

Department of Particle Physics and Astrophysics, Weizmann Institute of Science, Rehovot, 7610001 Israel
E-mail: kfir.blum@weizmann.ac.il, aielet.efrati@weizmann.ac.il, claudia.frugiuele@weizmann.ac.il, yosef.nir@weizmann.ac.il

ABSTRACT: We study the phenomenology of exotic color-triplet scalar particles $X$ with charge $|Q| = 2/3, 4/3, 5/3, 7/3, 8/3$ and $10/3$. If $X$ is an SU(2)$_W$-non-singlet, mass splitting within the multiplet allows for cascade decays of the members into the lightest state. We study examples where the lightest state, in turn, decays into a three-body $W^\pm jj$ final state, and show that in such case the entire multiplet is compatible with indirect precision tests and with direct collider searches for continuum pair production of $X$ down to $m_X \approx 250$ GeV. However, bound states $S$, made of $XX^\dagger$ pairs at $m_S \approx 2m_X$, form under rather generic conditions and their decay to diphoton can be the first discovery channel of the model. Furthermore, for SU(2)$_W$-non-singlets, the mode $S \to W^+W^-$ may be observable and the width of $S \to \gamma\gamma$ and $S \to jj$ may appear large as a consequence of mass splittings within the $X$-multiplet. As an example we study in detail the case of an SU(2)$_W$-quartet, finding that $m_X \approx 450$ GeV is allowed by all current searches.

KEYWORDS: Exotics, Hadron-Hadron scattering (experiments), Particle and resonance production

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1 Introduction

The large hadron collider (LHC) search for new physics at or below the TeV scale is far from complete, even for strongly interacting particles. As concerns the commonly studied Standard Model (SM) extensions [1–3], the dedicated searches by CMS and ATLAS for new strongly interacting light degrees of freedom are covering a large part of the parameter space. However, new colored particles beyond these standard scenarios could still have unexpected phenomenology and, in this case, traditional LHC searches often lose much of their power. In this work we consider colored scalar states with exotic EM charges, with a focus on SU(2)$_W$-non-singlets. Such particles, while being copiously produced at the LHC, could still be hiding undiscovered amidst the large QCD background. Three different paths can be pursued in the experimental search for these particles:

1. Direct collider searches for QCD continuum pair production of $X_Q$, a colored particle with EM charge $Q$. Such searches are potentially effective, but depend on the decay modes of $X_Q$ and hence are model dependent.

2. Precision measurements of electroweak (EW) processes, constituting an indirect search for $X_Q$.

3. Direct collider searches for $S_Q$, the bound state formed out of $X_Q X_Q^\dagger$ through Coulomb gluon exchange, with mass $m_{S_Q} \simeq 2m_{X_Q}$. $S_Q$ decays into diboson final states, with branching ratios that are determined to a large extent by the quantum numbers of $X_Q$. For exotic states the consequent constraints are often less model dependent than continuum pair production searches (see e.g. [4, 5]).

We pursue all three avenues in this work.

Color-triplet fields with exotic charges and/or in high SU(2)$_W$ representations are a rather generic outcome of unified models (see, e.g., ref. [6]). Within SU(5) models, SU(2)-triplets are parts of the 35, 45 and 70 representations, and an SU(2)$_W$-quartet is contained in the 70$^\prime$ representation [7]. Some of these representations are embedded in, for example, the 126 and 320 representations of the SO(10) group. These large representations are often invoked in GUT scenarios that address the issues of neutrino masses and of doublet-triplet splitting. Of course, the masses of these multiplets are not necessarily light. On a more phenomenological level, color-triplets of exotic charges have been introduced to explain various anomalies, such as the $B \rightarrow D^{(*)}\tau\nu$ anomaly [8], and the forward-backward asymmetry in $t\bar{t}$ events [9]. In these cases, these new degrees of freedom must be at the electroweak scale. Finally, we believe that, given the current status of experimental search for new physics, it is appropriate and timely to consider scenarios which are not necessarily related to the fine-tuning problem. In particular, special attention should be drawn to particles which can be produced abundantly at the LHC but would evade detection due to their distinct signature. Looking for novel signatures which were not the main focus of the experimental searches in recent years might encode new and interesting surprises.

The paper is organized as follows. In section 2 we present our theoretical framework and the relevant representations for our study. Section 3 details the experimental bounds.
from direct searches for continuum QCD pair production of $X_Q$. In section 4 we discuss mass splittings within $SU(2)_W$ multiplets and the implications for cascade decays. In section 5 we present a benchmark model. Section 6 deals with the unique phenomenology of $SU(2)_W$ multiplets, and the footprint it might leave in indirect probes such as electroweak precision measurements (EWPM), Higgs couplings and the renormalisation group evolution of various couplings. In section 7 we study the QCD bound states formed out of $X_QX_Q^\dagger$ pairs, and the possible signatures at the LHC. We conclude in section 8. Various technical details are presented in the appendices.

2 Theoretical framework

Consider a scalar $X$ in the $(R; n)_Y$ representation of the $SU(3)_C \times SU(2)_W \times U(1)_Y$ gauge group. The Lagrangian is given by

$$\mathcal{L} = \mathcal{L}_{SM} + |D_\mu X|^2 + \mathcal{L}_{YX} - V(H, X),$$

(2.1)

where $D_\mu$ is the covariant derivative, determined by the quantum numbers (QN) of $X$, and $H$ is the SM Higgs doublet, $H \sim (1, 2)_{+1/2}$. The scalar potential $V(H, X)$ has the form

$$V(H, X) = V^{SM}(H) + m_X^2 X^\dagger X + \frac{\lambda_X}{2} (X^\dagger X)^2 + \lambda_{XH} X^\dagger X H^\dagger H + \lambda'_{XH} (X^\dagger T^a_n X)(H^\dagger T^2_a H),$$

(2.2)

where $T^a_n$ are the $SU(2)_W$ generators in the $n$ representation. As we explore below, $\lambda_{XH} \neq 0$ generates mass splitting between the various states $X_Q$. Both $\lambda_{XH} \neq 0$ and $\lambda'_{XH} \neq 0$ modify Higgs couplings to SM fermions and gauge bosons.

We comment that eq. (2.2) is not the most general form possible for $V(H, X)$. Additional $X^4$ couplings may arise e.g. for color triplets in a non-singlet $SU(2)_W$ representation. As long as these couplings are small compared to $g_s^2 \sim 1$, they are not essential in most of our analysis and we omit them here.

As concerns the $SU(3)_C$ representation of $X$, we focus on color-triplets. This is a common starting point in many analyses, often considering quantum numbers similar to those of the SM quarks as occurs in supersymmetric models. The common lore is that first and second generation squarks are ruled out below 1.4 TeV while stops should be heavier than 900 GeV [10]. We study how this discussion is affected once exotic $SU(2)_W \times U(1)_Y$ representations are considered.

Other $SU(3)_C$ assignments have been studied in various contexts. For instance, supersymmetric models with Dirac gauginos introduce a color-octet scalar as the superpartner of the fermion which marries the gluino to form a Dirac fermion [11]. Color-sextets have been introduced in some models of grand unification [12, 13]. (See also [14] for a relevant discussion.) Despite this interest, we keep our focus on $R = 3$ for concreteness, though we include generic representations $R$ in some parts of the analysis where it does not introduce excess clutter.

The terms in $\mathcal{L}_{YX}$ break $X$ number and thus control the decay of $X$ to SM final states. With some abuse of notation, we refer to the terms in $\mathcal{L}_{YX}$ as Yukawa interactions.
maintain this terminology also to nonrenormalizable operators which, when the Higgs fields are replaced by their vacuum expectation values, lead to effective Yukawa couplings of $X$ with SM fermions. A doublet or a triplet of SU(2)$_W$ can couple to a fermion pair in a renormalizable operator, while other representations of SU(2)$_W$ require higher dimensional operators for the decay of their members. The inclusion of effective operators truncates the validity of our model at some cut-off scale $\Lambda$. To avoid the need for low cut-off scale, we restrict our discussion to effective operators with mass dimension $\leq 6$. This, in turn, leads us to consider $n \leq 5$, and limits the possible hypercharge assignments for $X$.

In table 1 we list all possible representations of $X$, for which we can find $X$-decay operators compatible with the restriction $d \leq 6$ for $\mathcal{L}_X$. We also list the corresponding diquark and/or leptoquark $X$-number violating operators. We denote the SM left-handed doublets as $Q$ and $L$, and the right-handed singlets as $U$, $D$ and $E$. Throughout the analysis we will assume that, when several operators are available in table 1, only one of them exists while the others are absent or negligible. For brevity, we omit $d \leq 6$ operators which include derivative interactions, as they introduce no new representations for $X$.

### Table 1

| $(R, n)_Y$ | $\lvert Q \rvert_{\text{high}}$ | Hadronic operators  | Lepto-quark operators |
|------------|-------------------------------|----------------------|-----------------------|
| (3, 1)$_Y$ | $2/3$                         | $XDUH^\dagger, XD_lD_jH, XD_lQ_jH$ | $XDL, XUEH^\dagger, XQLH^\dagger, XULH, XQEH^\dagger$ |
| (3, 2)$_{+1/6}$ | $4/3$                         | $XQ_lQ_jH, Xu_lU_jH^\dagger, XUDH$ | $XQLH^\dagger, XUEH^\dagger, XDEH^\dagger, XDLH$ |
| (3, 3)$_{-1/3}$ | $4/3$                         | $XQ_lQ_jH, Xu_lU_jH^\dagger, XD_lD_jHH$ | $XQL, XDLH, XDEH^\dagger H$ |
| (3, 2)$_{+7/6}$ | $5/3$                         | $XD_lD_jH^\dagger$ | $XQE, XUL, XDLH^\dagger H$ |
| (3, 3)$_{+2/3}$ | $5/3$                         | $XD_lD_jH^\dagger, XUDH^\dagger H^\dagger$ | $XQE, XULH, XDLH^\dagger$ |
| (3, 4)$_{-1/3}$ | $5/3$                         | $XQ_lQ_jH^\dagger$ | $XQLH^\dagger, XQEH, XDLH^\dagger H, XULH$ |
| (3, 2)$_{-11/6}$ | $7/3$                         | $XU_lU_jH$ | $XDEH$ |
| (3, 3)$_{-4/3}$ | $7/3$                         | $XU_lU_jH^\dagger H, XUDH, XQ_lQ_jH^\dagger H$ | $XUEH, XDEH^\dagger H$ |
| (3, 4)$_{-5/6}$ | $7/3$                         | $XQ_lQ_jH^\dagger$ | $XDLH, XQEH$ |
| (3, 5)$_{-1/3}$ | $7/3$                         | $XQ_lQ_jH^\dagger H^\dagger$ | $XQLH^\dagger H$ |
| (3, 2)$_{+13/6}$ | $8/3$                         | $XD_lD_jH^\dagger H^\dagger$ | $XQEH^\dagger H^\dagger, XQEH^\dagger H^\dagger$ |
| (3, 3)$_{+5/3}$ | $8/3$                         | $XD_lD_jH^\dagger H^\dagger$ | $XULH^\dagger, XQEH^\dagger$ |
| (3, 4)$_{+7/6}$ | $8/3$                         | $XQ_lQ_jH^\dagger H^\dagger$ | $XQEH^\dagger H^\dagger, XDLH^\dagger H^\dagger$ |
| (3, 5)$_{-2/3}$ | $8/3$                         | $XQ_lQ_jH^\dagger H^\dagger$ | $XQEH^\dagger H^\dagger$ |
| (3, 3)$_{-7/3}$ | $10/3$                        | $XU_lU_jH$ | $XDEH$ |
| (3, 5)$_{-4/3}$ | $10/3$                        | $XQ_lQ_jH$ | $XQLH$ |

### 3 Direct searches for continuum pair production

Colored particles are pair-produced at the LHC via initial state gluons. In this section we study the direct searches for continuum pair production of color triplet $X_Q$. The EM charge $Q$ dictates the possible decay modes and, subsequently, the experimental signatures. The SU(2)$_W$ quantum numbers are provisionally left out of the discussion.
Table 2 summarizes the possible decay final states of $X_Q$ for a given charge. We distinguish between two different decay topologies: 1) fully hadronic, in which $X_Q$ decays to two jets and possibly also $W$ bosons (we omit potential $jjh$ and $jjZ$ decay modes, as these are subdominant to an allowed $jj$ decay), and, 2) lepto-quark signature, in which $X_Q$ decays to a lepton (possibly a neutrino) and a jet.

Let us first analyze prompt signatures, highlighting the mass range $250 \text{ GeV} \leq m_{X_Q} \leq 1000 \text{ GeV}$. For some $X_Q$ decay topologies, dedicated searches were carried out by ATLAS, CMS, or the Tevatron collaborations. These decay modes, along with the relevant searches, are summarized in table 3. However, some of the signatures we study have no dedicated experimental analysis. We identify relevant searches which are sensitive to these topologies and estimate the corresponding efficiencies for our signal. For this purpose we implement our model in FeynRules [15] and simulate the signal in MadGraph5 [16] using Pythia 8 [17, 18] for showering and hadronization. Detector effects are simulated in Delphes [19] using the standard configuration. We stress that, for the recasted channels, our results should be taken as an estimation only. A detailed description of our recast procedure can be found in appendices A, B and C.

Our findings are presented in figure 1(a) for the dijet decays, figure 1(b) for the jet and charged lepton signals, and figure 1(c) for the neutrino-jet topology. We also consider the case where a jet is replaced by heavy flavor quark. In each figure we show the current limit on the pair-production cross section times BR$^2$, normalized to the NLO+NLL cross section of a scalar colored triplet taken from [20–22]. Presented this way, when a single mode dominates the decay (namely BR = 1), the $y$ axis corresponds to the number of copies of the $X$ representation that are experimentally allowed.

An important ingredient for collider phenomenology is the lifetime of $X_Q$. Stable particles and non-prompt decays are studied by the experimental collaborations in dedicated searches, leading to bounds in the ballpark of $m_{X_Q} \gtrsim 700–900 \text{ GeV}$ for color-triplet scalars. Refs. [23, 24] analyzed displaced signatures in the context of RPV SUSY models. They find that $X_Q$ in the mass range of 100–1000 GeV, decaying to dijet, or to a jet and a charged lepton, or to a jet and a neutrino, would not be captured by the displaced-track searches if its mean-free path is less than 0.3–10 mm. While the exact number depends on the particle mass and decay mode, we conservatively use in the following 0.3 mm as an upper bound.
| Diquark       | jj          | bj          | tj, tb                  |
|--------------|-------------|-------------|-------------------------|
| Refs.        | [25–29]     | [25, 30, 31]| [32, 33]                |
| Comments     | RPV SUSY    | RPV SUSY    | RPV SUSY search for     |
|              |             |             | m ≤ 600 GeV, with j ≠ b, c. |
|              |             |             | t* → tg search for higher masses. |
|              |             |             | We assume similar efficiencies for higher pr gluon and quark jets. |
| Wqq          |             |             |                         |
| Refs.        | [34–43]     | [32, 44, 45]|                         |
| Comments     | Recast Wj, Wb and Wt | Recast of multi-lepton searches. |
|              |             |             | reduction. See appendix A. |
| Wt           |             |             |                         |
| Comments     |             |             |                         |
|              |             |             |                         |
| Lepto-quark  | ℓj, ℓb     | ℓt          | τj                      |
| Refs.        | [46–52]     | [45]        | [55] for m_X ≥ 600 GeV, |
| Comments     | 1st and 2nd generation lepto-quark searches | Recast of multi-lepton searches. |
|              | w/o b veto  |             | recasting νj searches using 30%–50% efficiency reduction |
| τt           |             |             |                         |
| Comments     | [56]        | [55, 57, 58]|                         |
| νj           |             |             |                         |
| Comments     | [53, 54]    | [47, 59]    | [47, 59]                |
| νb           |             |             |                         |
| Comments     | Standard SUSY searches | lepto-quark searches |
|              | for squark pair searches | lepto-quark searches |

| Table 3. Direct searches for X_Q used in our analysis. The resulting bounds are shown in figure 1. We use the notation q for all six quark flavors, while j = u, d, s, c, and ℓ = e, μ. |

on a two-body decay length. We are not aware of any dedicated analysis for displaced signature of a three- (or four-) body final state. We estimate that the larger multiplicity of the final objects would increase the efficiency of these searches at high m_X, while the low m_X regime will suffer from the typically lower energy carried by each final object. Over all, we expect that the sensitivity is comparable to the other topologies, and so we consider cτ ≤ 1 mm for three-body decay. We then apply the following ‘promptness’ requirements on X_Q decay rates:

\[
\Gamma_{2\text{-body}} \gtrsim 7 \times 10^{-13} \text{ GeV}, \quad \Gamma_{3\text{-body}} \gtrsim 2 \times 10^{-13} \text{ GeV},
\]

(3.1)

which translate into lower bounds on the Yukawa coupling of X_Q to SM states. If X decays via non-renormalizable operators, the exact cut-off of its Yukawa operators depends on the details of the UV interactions, the mass of X, and its representation under the SM gauge group. Generically, for m_X ≈ 500 GeV, dimension-five operators would require \( \Lambda \lesssim 10^6 \text{ TeV} \), while dimension-six operators would require \( \Lambda \lesssim 100–1000 \text{ TeV} \).
Figure 1. Bounds from direct searches for $X_Q$ continuum pair production in the various final state topologies. The $y$ axes give $\sigma \times \text{BR}^2$ normalized to the pair production cross-section of a color-triplet scalar. Sharp features are caused by considering multiple searches for each channel; see the text for more details.

Concluding this section, we learn the following:

- The lepto-quark topology is strongly constrained by direct searches. As can be seen in figure 1(b), none of the decay modes in this category allows for more than two states below $m_{X_Q} \simeq 750$ GeV.

- The neutrino-quark topology is subject to the standard SUSY searches for jet and missing energy. As can be seen in figure 1(c), the corresponding bounds on $m_{X_Q}$ are even stronger than in the $jl$ category.

- The hadronic decay modes are significantly less constrained by direct searches. This is expected given the large QCD backgrounds at the LHC.

- A $Wjj$ signature is poorly constrained by the LHC. As we show below, such topology could be the signature of multiple states which undergo cascade SU(2)$_W$ decays. This is an important gap in the LHC coverage for colored new particles which motivates dedicated searches for this decay topology.
4 Mass splitting and cascade decays

In general, two members of an SU(2)_W multiplet with EM charges Q and Q' are split in mass. Tree level mass splittings are induced by the $\lambda'_{XH}$ term:

$$ (m_Q^2 - m_{Q'}^2)_{\text{tree}} = -\frac{\lambda'_{XH} v^2}{4} (Q - Q') $$

$$ \Rightarrow (m_Q - m_{Q'})_{\text{tree}} \simeq 1.7 \text{GeV} \times (Q' - Q) \left(\frac{\lambda'_{XH}}{0.1}\right) \left(\frac{m_X}{450 \text{ GeV}}\right)^{-1}. \quad (4.1) $$

Mass splittings also arise through electroweak gauge boson loops from the kinetic term $(D_\mu X)^\dagger (D^\mu X)$ [60]:

$$ (m_Q - m_{Q'})_{\text{1-loop}} = \frac{\alpha m_Z}{2} \left\{ (Q^2 - Q'^2) \tilde{f}(x_Z) + (Q - Q')(Q + Q' - 2Y) \frac{1}{s_W^2} [c_W \tilde{f}(x_W) - \tilde{f}(x_Z)] \right\} $$

$$ \Rightarrow (m_Q - m_{Q'})_{\text{1-loop}} \simeq -0.15 \text{GeV} \times (Q' - Q)(Q' + Q + 2.3Y), \quad (4.2) $$

where $s_W^2 \equiv \sin^2 \theta_W$, $c_W \equiv \cos \theta_W$, $x_V = m_V/m_X$, and $\tilde{f}(x) = -\frac{1}{8\pi} (2x^3 \log x + (x^2 - 4)^{3/2}) \log[(x^2 - 2 - x\sqrt{x^2 - 4})/2] = 1 - \frac{x}{3} + O(x^2)$.

Assuming no fine-tuned cancelations between the tree and loop contributions, a mass splitting of at least $O(100 \text{ MeV})$ between adjacent members of the multiplet $(Q = Q' + 1)$ is unavoidable. Much larger splittings are possible, depending on the value of $\lambda'_{XH}$. If the tree contribution dominates, the splitting can be of either sign, and the lightest colored scalar is the one with either the highest or the lowest Q.

The mass splitting between the members of an SU(2)_W multiplet leads to W-mediated decays within the multiplet, $X_m \rightarrow X_{m \pm 1} W^{\mp(*)}$. (Note that we change notations in this section from $X_Q$ to $X_m$, with $Q = m + Y$.) For the three-body decay, $X_m \rightarrow X_{m+1} f \overline{f'}$, with massless fermions, we obtain

$$ \Gamma(X_m \rightarrow X_{m+1} f \overline{f'}) = \frac{G_F^2}{15\pi^3} (j - m)(j + m + 1)(\Delta M)^5 $$

$$ \simeq 3 \times 10^{-13} \text{GeV} \left(\frac{\Delta M}{1 \text{ GeV}}\right)^5 (j - m)(j + m + 1). \quad (4.3) $$

If $\Delta M > m_\pi$, we have the two body decay $X_m \rightarrow X_{m+1} \pi^-$, in which case

$$ \Gamma(X_m \rightarrow X_{m+1} \pi^-) = \frac{G_F^2}{\pi} (j - m)(j + m + 1)(\Delta M)^3 f^2 \sqrt{1 - \frac{m_\pi^2}{(\Delta M)^2}} $$

$$ \simeq 7 \times 10^{-13} \text{GeV} \left(\frac{\Delta M}{1 \text{ GeV}}\right)^3 (j - m)(j + m + 1). \quad (4.4) $$

For $m = -1$ we recover the results of ref. [60]. We do not consider $\Delta M > m_W$.

To determine the phenomenological significance of these decays (for all but the lightest member of the multiplet), we need to compare their rate to those of the Yukawa mediated decays. We will do so in the next section.
5 A model example: $X(3,4)_{_{+1/6}}$

In the following we discuss the model example $X \sim (3,4)_{_{+1/6}}$, containing a state with $Q = +5/3$ as the highest charge state. We assign $X$ zero lepton number which, given our assumptions in section 2, forces $X_{+5/3}$ to decay into the hadronic three body state $\bar{d}_i \bar{d}_j W^+$ via the operator

$$\mathcal{L}_{YX} = \frac{Y^Q_{ij}}{2} Q^i Q^j H^+ X + \text{h.c.},$$

with $Y^Q_{ij}$ antisymmetric in the flavor indices $i, j$, and of dimension mass$^{-1}$.

We consider two specific scenarios:

- Case A: degenerate $X_Q$ states.
- Case B: non-degenerate $X_Q$ states.

We now show that these two cases exhibit distinct phenomenology.

5.1 Degenerate SU(2)$_W$-quartet

The Lagrangian (5.1) gives the following component interaction terms (to leading order in CKM rotation) for the four multiplet members (with $Q = +5/3, +2/3, -1/3, -4/3$):

$$\mathcal{L} = Y^Q_{ij} \left[ X_{+5/3} d_i d_j h^- + X_{-4/3} \left( \frac{v}{\sqrt{2}} u_i u_j + \frac{1}{\sqrt{2}} u_i u_j (h + i\rho) \right) + X_{+2/3} \left( \frac{v}{\sqrt{6}} d_i d_j + \frac{1}{\sqrt{6}} d_i d_j (h + i\rho) - \frac{1}{\sqrt{3}} d_i u_j h^- - \frac{1}{\sqrt{3}} u_i d_j h^- \right) + X_{-1/3} \left( \frac{v}{\sqrt{6}} d_i u_j + \frac{v}{\sqrt{6}} u_i d_j + \frac{1}{\sqrt{6}} d_i u_j h + \frac{1}{\sqrt{6}} u_i d_j h - \frac{1}{\sqrt{3}} u_i u_j h^- \right) \right] + \text{h.c.}.$$  

(5.2)

(We work in the Feynman-'t Hooft gauge $\xi = 1$ as to straightforwardly keep track of longitudinal $W^+, Z$ modes.) These terms allow two body decays of $X_{+2/3}$, $X_{-1/3}$ and $X_{-4/3}$:

$$\Gamma(X_Q \to \bar{q}_L\bar{q}_L) = c_Q |Y^Q_{ij}|^2 \frac{3v^2}{16\pi} m_X \lambda[m_i^2, m_j^2, m_X^2]^{1/2} \beta[m_i + m_j, m_X],$$

(5.3)

where $c_Q = \frac{1}{3}, \frac{1}{3}, 1$ for $Q = +2/3, -1/3, -4/3$, respectively, $\lambda[x, y, z] = (1 - x/z - y/z)^2 - 4xy/z^2$, and $\beta[x, y] = (1 - x^2/y^2)$. They also allow three body decays of all four members:

$$\Gamma(X \to \phi\bar{q}_L\bar{q}_L) = \tilde{c}_Q |Y^Q_{ij}|^2 \frac{m_X^2}{512\pi^3} m_X \left[ 1 + \frac{2\phi}{m_X^2} \left( 9 + 6 \log \frac{m_X^2}{m_\phi^2} - 6 \frac{m_\phi^2}{m_X^2} + O \left( \frac{m_\phi}{m_X} \right)^4 \right) \right],$$

(5.4)

Here $\tilde{c}_Q = 1, \frac{1}{3}, \frac{1}{3}, \frac{1}{2}$ for $Q = +5/3, +2/3, -1/3, -4/3$, respectively, and we take $m_j = 0$. The boson $\phi$ is a neutral Higgs or a longitudinal gauge boson.

For $m_X \lesssim 8$ TeV, the two-body decays of eq. (5.3), where available, dominate over the three-body decays of eq. (5.4). If the $Y^Q_{ij}$ term dominates the decay rates of all members
of the quartet, then

\[
\begin{align*}
\text{BR}[X_{+5/3} \rightarrow \bar{d}_i \bar{d}_j W^+] & \simeq 1, \\
\text{BR}[X_{-4/3} \rightarrow \bar{u}_i \bar{u}_j] & \simeq 1, \\
\text{BR}[X_{+2/3} \rightarrow \bar{d}_i \bar{d}_j] & \simeq 1, \\
\text{BR}[X_{-1/3} \rightarrow \bar{u}_i \bar{d}_j] = \text{BR}[X_{-1/3} \rightarrow \bar{d}_i \bar{u}_j] & \simeq \frac{1}{2}.
\end{align*}
\]

(5.5)

For \(i, j = 1, 2\), we have three states decaying into a \(jj\) final state, and one state decaying into a \(Wjj\) topology. This is allowed for \(m_X \gtrsim 630 \text{ GeV}\). For \(i = 3\), we have effectively 1.5 members decaying into \(jb\) and \(jt\) each. Looking at \(N \times \text{BR}^2 = 1.25\) in figure 1(a) we conclude that \(m_{X_Q} = 520 \text{ GeV}\) is a viable possibility. We use this mass as our benchmark point in the following. To guarantee prompt decay of \(X_{+5/3}\) we impose

\[
Y_{QQ}^{X_{+5}} \gtrsim 7 \times 10^{-9} \text{ GeV}^{-1}(520 \text{ GeV}/m_X)^{3/2}.
\]

(5.6)

5.2 Non-degenerate SU(2)_W-quartet

Mass splitting between the members of the quartet allow for fast cascade decays of the three heavier ones. In order to establish their phenomenological relevance one needs to compare the rate of these weak decays with the rate of the Yukawa mediated decay modes, which depend on the dimensional coupling \(Y_{QQ}^{X_{+5}}\), eqs. (5.3) and (5.4). The dominant terms need to induce prompt decays for all the members of the \(X\) multiplet. We distinguish between two cases:

1. \(X_{-4/3}\) is the lightest: In this case, either all states decay dominantly via their Yukawa coupling, or the \(Q = +5/3\) state (and possibly also the \(Q = +2/3\) and \(Q = -1/3\) states) decay via \(W\)-mediated cascade decays. In either case, we have at least three color-triplet states decaying into two jets. The mass of the lightest state should then be similar to the mass considered in the degenerate quartet scenario.

2. \(X_{+5/3}\) is the lightest: In this case, \(X_{+5/3}\) decays to a \(\bar{d}_i \bar{d}_j W^+\) final state. As concerns the three heavier states, they can either decay into two jets, or cascade into the \(X_{+5/3}\) state. The latter would lead to effectively four states decaying to \(Wqq\) in the final state, assuming the other cascade products are too soft to be detected (this is the case for a few GeV splitting). As far as direct searches for continuum pair production are concerned, we estimate the sensitivity of top-partner searches at the Tevatron [34] and find that, in this case, \(X_{+5/3}\) can be as light as 250 GeV. As we will see next, the direct searches for an \(X_Q X_Q^\dagger\) bound state place a stronger limit, of \(m_X \gtrsim 450 \text{ GeV}\), with a corresponding lower bound on the Yukawa coupling,

\[
Y_{\text{min}}^{QQ} \simeq 9.1 \times 10^{-9} \text{ GeV}^{-1}(450 \text{ GeV}/m_X)^{3/2},
\]

(5.7)

to ensure its prompt decay. Using \(Y_{\text{min}}^{QQ}\) as a convenient reference, and recalling that the two-body decay rate is faster than the three-body one for \(m_X \lesssim 8\text{ TeV}\), a mass splitting of

\[
\Delta M \gtrsim 2.7 \text{ GeV} \left(\frac{450 \text{ GeV}}{m_X}\right)^{2/5} \left(\frac{Y_{QQ}^{X_Q}}{Y_{\text{min}}^{QQ}}\right)^{2/5}
\]

(5.8)
between two ‘adjacent’ members of the multiplet would effectively cause the four members of $X(3,4)_{+1/6}$ decay to $Wqq$ final states. The precise coefficient varies a little between the different SU(2)$_W$ members.

We therefore consider, for our second scenario, the following spectrum:

\[ m_{X_{+5/3}} = 450 \text{ GeV}, \quad m_{X_{+2/3}} = 452.8 \text{ GeV}, \]
\[ m_{X_{-1/3}} = 455.7 \text{ GeV}, \quad m_{X_{-4/3}} = 458.5 \text{ GeV}, \]

which is the result of $\lambda_{XH} = 0.17$.

6 SU(2)$_W$ phenomenology

In this section we explore the distinct phenomenology of colored SU(2)$_W$ non-singlet scalars.

6.1 Electroweak precision measurements (EWPM)

Large mass splitting within an SU(2)$_W$ multiplet is constrained by EWPM. Specifically, it modifies the oblique $T$ and $S$ parameters [61], where the leading effect comes from generating the dimension six operators $O_T$ and $O_{WB}$ (see appendix E for the definition of these operators). For an $(R;n)_{Y}$ representation, we have

\[
T = \frac{v^2}{\alpha} c_T = \left( \frac{v^2}{4608 \pi^2 \alpha} \right) \left( \frac{\lambda_{XH}^2}{m_X^2} \right) R n (n^2 - 1) \\
= 4.4 \times 10^{-3} \left( \frac{\lambda_{XH}}{0.17} \right)^2 \times \left( \frac{R n (n^2 - 1)}{180} \right) \left( \frac{450 \text{ GeV}}{m_X} \right)^2 ,
\]
\[
S = 16 \pi v^2 c_{WB} = \left( \frac{v^2}{144 \pi} \right) \left( \frac{\lambda_{XH}^2}{m_X^2} \right) R n (n^2 - 1) Y \\
= 3.4 \times 10^{-3} \left( \frac{\lambda_{XH}}{0.17} \right) \times \left( \frac{R n (n^2 - 1) Y}{30} \right) \left( \frac{450 \text{ GeV}}{m_X} \right)^2 ,
\]

where in the second equation of each line we normalize to the quantum numbers of $X(3,4)_{+1/6}$ and to the value of $\lambda_{XH}^2$ which we use for case B in section 5. The EWPM constraints read (for $U = 0$) [62]

\[
T = 0.10 \pm 0.07 , \quad S = 0.06 \pm 0.09 ,
\]

with correlation of $\rho = 0.91$. Using one dimensional $\chi^2(\lambda')$ function we find that $|m_{X_Q} - m_{X_{Q+1}}| \lesssim 13$–16 GeV is allowed around 450 GeV, where a positive (negative) $\lambda_{XH}^2$ implies that $X_{+5/3}$ ($X_{-4/3}$) is the lightest member of the multiplet. Clearly, EWPM allow the mass splitting we consider in case B.

In the limit of custodial symmetry, modifications to the EW vacuum polarization amplitude alter the oblique $Y$ and $W$ parameters [63]. These are primarily encoded in the
dimension six operators $O_{2B}, O_{2W}$:

\[
Y = 2m_W^2 c_{2B} = \frac{g^2}{240\pi^2} \frac{m_W^2}{m_X^2} Rn Y^2 \simeq 5.7 \times 10^{-7} \times \left( \frac{RnY^2}{1/3} \right)^2 \left( \frac{450 \text{ GeV}}{m_X} \right)^2,
\]

\[
W = 2m_W^2 c_{2W} = \frac{g^2}{2880\pi^2} \frac{m_W^2}{m_X^2} Rn(n^2 - 1) \simeq 8.9 \times 10^{-5} \times \left( \frac{Rn(n^2 - 1)}{180} \right)^2 \left( \frac{450 \text{ GeV}}{m_X} \right)^2.
\]

(6.3)

These contributions to $Y$ and $W$ are below the current sensitivity of LEP (see, e.g., table 4 of [64]) and the LHC [65–67]. The values we take for the various coupling constants are listed in appendix D.

6.2 Gauge coupling running

The presence of $X \sim (R, n)_Y$ modifies the running of the gauge coupling constants. We describe this effect, at one-loop level, in appendix D. In particular, high SU(2)$_W$ representations change significantly the running of $\alpha_2$. For instance, color-triplet in the quartet (or higher) representation of SU(2)$_W$ flips the sign of the SU(2)$_W$ beta function. In particular, for $X(3,5)_Y$, $\alpha_2$ becomes non-perturbative at $\mu \simeq 10^{15}$ GeV. Since the decay of $X$ already requires some cut-off at a lower scale, this is insignificant to our study.

Additional probe for the running of EW gauge coupling is the differential distribution of Drell-Yan processes at various energies, as was proposed in ref. [64]. Refs. [65, 66] find that for $m_\psi = 520$ GeV, $N_\psi Q^2 \geq 46$ is excluded at the 2σ level, where $N_\psi$ is the number of copies of vector-like fermions transforming as $\psi \sim (3, 1)_Q$. This scenario would generate a 23% (50%) relative increase in the Drell-Yan rate at $m_{t\bar{t}} = 1 (1.5)$ TeV, which excludes $b_2^X \leq -46$. In our model example of section 5, $b_2^X = -Rn(n^2 - 1)/36 = -5$, clearly within bounds. A more recent analysis done in ref. [67] yields the same conclusion.

6.3 Additional constraints

SM Higgs couplings. Integrating out $X$ generates dimension six effective operators involving the Higgs field. These, in turn, modify the Higgs couplings to fermions and gauge-bosons with respect to their SM values. LHC Higgs data constrain these modifications, resulting in bounds on the quartic couplings $\lambda_{XH}$ and $\lambda_{XH}'$. At present, EWPM induce stronger constraints on $\lambda_{XH}$. The Higgs data do constrain $\lambda_{XH}$, but this coupling is not directly relevant to our analysis. We present our numerical results of the Higgs data for $X(3,4)_{+1/6}$ in appendix G, and the resulting minor effects on the various $S \to VV$ decays in appendix I.

Scalar quartic coupling running. In addition to modifying the SM Higgs couplings to fermions and gauge bosons, the presence of $X$ changes the running of the SM Higgs couplings. We calculate these effects in appendix H. We find that no dangerous runaway behavior is generated. The same conclusion holds for the $X$ quartic coupling, and the mixed couplings $\lambda_{XH}$ and $\lambda_{XH}'$. 

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7 QCD bound state

In the previous sections we obtained constraints from both direct and indirect probes on the existence of exotic colored scalars. The interesting result is that these constraints can be quite mild, allowing rather light colored scalars. For example, as demonstrated by the non-degenerate quartet scenario (case B in section 5), the data still allow four colored states with \( m_X \approx 250 \text{ GeV} \). In this section we study another way to discover light colored scalars, which might go first through the observation of their QCD bound state \([4, 5, 68]\). Moreover, constraints derived from bound state searches are less model dependent, in the sense that they do not depend on the decay mode of \( X \).

A pair of \( X_Q X_Q^I \) near threshold can form a QCD bound state, which we denote by \( S_Q \). If the decay rates of its constituents are smaller than \( \Gamma_{S_Q} \), and its width is smaller than the respective binding energy, \( S_Q \) can be seen as a resonance as it annihilates into SM particles. For a review we refer the reader to ref. \([69]\) and references therein. Heavy constituents exhibit Coulomb-like potential with a binding energy

\[
E_b = -\frac{1}{4 n_F^E} [C_2(R)]^2 \bar{\alpha}^2 s m_X ,
\]

where \( n_F \) is the excitation index (\( n = 1 \) is the ground state), \( \bar{\alpha} \) is the strong coupling evaluated at the bound-state typical scale (for which we use, following ref. \([69]\), the Bohr radius) and \( C_2(R) \) is the quadratic SU(3) Casimir of representation \( R \), with \( C_2(3) = 4/3 \). We assume that the resulting bound state is an SU(3) singlet. The mass of \( S \) is \( m_S = 2 m_X + E_b \).

The condition that pair annihilation dominates the decay of \( S_Q \) reads

\[
2 \Gamma_X < \Gamma_{S_X} = 2 \times 10^{-5} m_X .
\]

The r.h.s. is well above the lower bounds in eq. (3.1). In fact, (7.2) is fulfilled quite generally by the exotic states on which we focus the analysis. The argument goes as follows. Suppose that \( X \) decays into a two fermion final state, with effective coupling \( y \). The condition (7.2) translates into \( y < 10^{-2} \). If the effective coupling comes from a dimension \( d \) operator, we have \( y = \hat{y}(v/\Lambda)^{d-4} \), where \( \hat{y} \) is dimensionless and \( \Lambda \) is the scale of new physics. We assume perturbativity (\( \hat{y} \lesssim 1 \)), and a NP scale that is not very low (\( \Lambda \gtrsim 10 \text{ TeV} \)). Then, for \( d = 6 \) operators, the condition is always fulfilled. For \( d = 5 \) operators it is not fulfilled only in a small region of parameter space where \( \Lambda \lesssim 25 \text{ TeV} \) and \( \hat{y} > 0.4 \). Fully hadronic decays via renormalizable operators (\( d = 4 \)) are possible only in a single case of SU(2)_W non-singlet, that is \( X(3, 3)_{-1/3} \), and even then the condition is fulfilled for \( \hat{y} < 0.01 \). The condition (7.2) applies in all cases of dominant three body final state. We conclude that the search for bound states is truly a generic tool to look for exotic colored scalars \([4]\).

The quantum numbers of \( X \) determine the gluon fusion (ggF) production cross section of \( S \) as well as its decay rates into pairs of vector bosons: \( gg, \gamma\gamma, ZZ, Z\gamma \) and \( WW \). Assuming that the \( X + X^I \) production is dominated by ggF, and that there are no additional decay modes that give a significant contribution to the total width of \( S \), then
\( \sigma(pp \to S) \times \text{BR}(S \to V_1 V_2) \) is predicted. The ggF partonic production cross section is given by

\[
\hat{\sigma}_{gg \to S} = \frac{\pi^2 \Gamma(S \to gg)}{8 m_S} \delta(\hat{s} - m_S^2). \tag{7.3}
\]

We convolute \( \hat{\sigma} \) with the partonic luminosity function

\[
\sigma = \frac{\hat{\sigma}}{m_S^2} \frac{\tau dL}{d\tau} \tag{7.4}
\]

evaluated at \( \tau = m_S^2/s \), where \( \sqrt{s} \) is the CoM energy. For the various two-body decay rates, we use (see [4] and references therein)

\[
\Gamma(S_Q \to V_1 V_2) = \frac{R}{8\pi(1 + \delta_{V_1 V_2})} \frac{|\psi(0)|^2}{m_S^2} [\mathcal{M}_{V_1 V_2}]^2 \lambda^{1/2}(m_S^2, m_{V_1}^2, m_{V_2}^2), \tag{7.5}
\]

where \( \lambda[x, y, z] \) is defined below eq. (5.3), and \( |\psi(0)| \) is the joint wave function of \( X_Q X_Q^1 \) at the origin, which controls the probability to form a bound state, and is given by

\[
|\psi(0)|^2 = \frac{[C_2(R)]^2 m_X^3}{8\pi n^3}. \tag{7.6}
\]

The full expressions for \( [\mathcal{M}_{V_1 V_2}]^2 \) can be found in appendix I. We provide here the ratios between the different decay rates of \( S_Q \) (with \( Q = m + Y \)), denoting \( R_{X/Y}^Q = \Gamma(S_Q \to X)/\Gamma(S_Q \to Y) \), and neglecting contributions proportional to \( \lambda_{X H}^m = \lambda_{X H} - (m/2)\lambda_{X H}' \) and phase space suppressions:

\[
\begin{align*}
R_{gg/\gamma\gamma}^Q &= \frac{C_2(R)\alpha_s^2}{8Q^4\alpha_{EM}^2}, & R_{ZZ/\gamma\gamma}^Q &= \frac{[m - Qs_W^2]^4}{s_W^4 c_W^4 Q^4}, \\
R_{Z/\gamma\gamma}^Q &= \frac{2[m - Qs_W^2]^2}{s_W^4 c_W^4 Q^2}, & R_{WW/\gamma\gamma}^Q &= \frac{(n^2 - 1 - 4m^2)^2}{32s_W^4 c_W^4 Q^4}.
\end{align*} \tag{7.7}
\]

In the limit of small mass splitting, the various \( V_1 V_2 \) signals depend on the sum of the branching ratio of each member, rather than on the sum of \( R_Q \). They are the same if the total width of all the \( S_Q \) members is equal, which is the case if the digluon mode dominates the total width. In table 4 we calculate the ratios between the different \( V_1 V_2 \) signals, summing over all \( S_Q \)’s. Note that the running of the gauge coupling slightly modifies the numerical values of these ratios for various bound state masses. For concreteness, we quote these values at \( m_S = 800 \text{ GeV} \), and denote

\[
R_{X/Y} = \frac{\sum_Q \text{Br}[S_Q \to X]}{\sum_Q \text{Br}[S_Q \to Y]}. \tag{7.8}
\]

We further specify, in table 4, \( \sigma_{\gamma\gamma}^{13} \), the expected diphoton signal at the 13 TeV LHC for the various representations we consider, taking \( m_S = 800 \text{ GeV} \). Bound state composed of \( \text{SU}(2)_W \) non-singlet exhibit several interesting features, which we discuss next.
Leptoquark involving the lower bound on $m_X$. Consider, for example, the quartet $X$ ones from direct continuum searches, but have the advantage of being less model dependent. It is around 800 GeV for decays into a lepton pair, but can be very weak for fully hadronic decays and depends on the SU(2)$_W$. These bounds are effective: in fact, for SU(2)$_W$ singlets the bound is stronger than the bound from LHC direct continuum pair production searches in a large region of the parameter space. For instance, as discussed in the previous section, there are only very week bounds for an $X$ singlet to the one obtained from the highest SU(2)$_W$ representation listed in table 1. For the same charge $Q$ we notice a dependence on the SU(2)$_W$ representation.

The experimental upper bounds on $\sigma_{\gamma\gamma}$ at 13 TeV translate into a lower bound on $m_S$ and, consequently, on $m_X$. These bounds are effective: in fact, for SU(2)$_W$ singlets the bound is stronger than the bound from LHC direct continuum pair production searches in a large region of the parameter space. For instance, as discussed in the previous section, there are only very week bounds for an $X$ singlet from direct continuum pair production searches, while the search for diphoton resonance gives $m_{X_{\pm3}} \gtrsim 600$ GeV.

For higher SU(2)$_W$ representations, the bound state limits can be weaker than the ones from direct continuum searches, but have the advantage of being less model dependent. Consider, for example, the quartet $X$ $3(4,4)_{+1/2}$. As discussed in the previous section, the lower bound on $m_X$ is very model dependent. It is around 800 GeV for decays into a leptoquark involving $e$ or $\mu$, but can be very weak for fully hadronic decays and reason-

\[
\begin{array}{ccccccccc}
(R, n)_Y & |\mathcal{Q}|_{\text{high}} & \sigma_{\gamma\gamma}^{13} [\text{fb}] & R_{WW/\gamma\gamma} & R_{Z\gamma/\gamma\gamma} & R_{ZZ/\gamma\gamma} & R_{gg/\gamma\gamma} \\
(3, 1)_Y & |Y| & 0.48 Y^{-4} & 0 & 0.09 & 0.6 & 30 Y^{-4} \\
(3, 2)_{+1/6} & 2/3 & 0.09 & 22 & 6.8 & 3.8 & 286 \\
(3, 2)_{-5/6} & 4/3 & 1.3 & 1.4 & 1.1 & 0.3 & 19 \\
(3, 3)_{-1/3} & 4/3 & 0.9 & 15 & 6.8 & 3.9 & 26 \\
(3, 2)_{+7/6} & 5/3 & 3.0 & 0.6 & 0.7 & 0.3 & 7.6 \\
(3, 3)_{+2/3} & 5/3 & 1.9 & 6.7 & 3.4 & 1.6 & 11 \\
(3, 4)_{+1/6} & 5/3 & 1.1 & 25 & 10.7 & 6.1 & 11 \\
(3, 2)_{-11/6} & 7/3 & 8.0 & 0.1 & 0.4 & 0.4 & 1.8 \\
(3, 3)_{-4/3} & 7/3 & 6.1 & 1.8 & 1.3 & 0.5 & 2.8 \\
(3, 4)_{-5/6} & 7/3 & 3.1 & 8.8 & 4.0 & 2.0 & 3.8 \\
(3, 5)_{-1/3} & 7/3 & 1.7 & 25 & 9.7 & 5.6 & 4.0 \\
(3, 2)_{+13/6} & 8/3 & 10 & 0.08 & 0.3 & 0.4 & 1.0 \\
(3, 3)_{+5/3} & 8/3 & 9.1 & 1.1 & 0.9 & 0.4 & 1.6 \\
(3, 4)_{+7/6} & 8/3 & 5.0 & 5.2 & 2.5 & 1.1 & 2.2 \\
(3, 5)_{+2/3} & 8/3 & 2.5 & 17 & 6.6 & 3.6 & 2.7 \\
(3, 3)_{-7/3} & 10/3 & 15 & 0.4 & 0.6 & 0.4 & 0.6 \\
(3, 5)_{-4/3} & 10/3 & 6.0 & 7.0 & 3.0 & 1.4 & 1.1 \\
\end{array}
\]

Table 4. The $\sigma_{\gamma\gamma}^{13}$ cross section for $m_S = 800$ GeV and the ratios between $S \rightarrow V_1 V_2$ and the diphoton signals for the various SU(3)$_C \times SU(2)_W \times U(1)_Y$ representations. The singlet values are valid for small hypercharge, assuming the digluon decay mode dominates the total width of $S$.

7.1 Diphoton signature

Interestingly, if $X$ transforms in a large SU(2)$_W$ representation, its total width can be much larger than its partial width into $gg$. This can deplete the various $S$ signals, in particular the $S \rightarrow \gamma \gamma$ one. We demonstrate this effect in figure 2, where we show, for a given charge, the differences between the diphoton signal of an SU(2)$_W$ singlet to the one obtained from the highest SU(2)$_W$ representation listed in table 1. For the same charge $Q$ we notice a dependence on the SU(2)$_W$ representation.
able mass splitting. Diphoton resonance searches set a solid bound of 450 GeV which is independent of these details of the model. Similar statements can be made for other high SU(2)_{W} representations.

### 7.2 Distinct features of a bound state composed of SU(2)_{W}-non-singlet constituents

If an X-onium \( S \) involves \( X \) that is an SU(2)_{W}-non-singlet, then it might exhibit two features that would clearly distinguish it from the SU(2)_{W}-singlet case: a large branching ratio into \( W^+W^- \) and an apparent large width. In this subsection we explain these two features.

**Large BR(\( S \to W^+W^- \)).** Observation of any diboson decay mode of \( S \to \gamma\gamma, W^+W^-, ZZ, Z\gamma \) — will help to close in on the representation of \( X \). Our main focus is on cases where the \( S \to W^+W^- \) decay rate is large. For the sake of concreteness, we examine whether \( R_{WW/\gamma\gamma} \geq 10 \) is possible. Table 4 shows five candidates. We list them by the order of the lower bound on their mass from diphoton searches:

| \( |Q|_{\text{high}} \) | \( \sigma_{\gamma\gamma} \) ratio |
|------------------|------------------|
| 2/3              | 1.0              |
| 4/3              | 1.6              |
| 5/3              | 2.4              |
| 7/3              | 3.4              |
| 8/3              | 2.7              |
| 10/3             | 1.4              |

Table 5. The ratio between the diphoton signals in the singlet and high-representation cases. The depletion of the signal for high representations is clearly seen for the various charges. The numbers are given for \( m_S = 800 \) GeV and vary only little due to RGE effects.
\begin{itemize}
  \item $(3, 2)_{+1/6}$, with $m_S \gtrsim 500$ GeV.
  \item $(3, 3)_{-1/3}$, with $m_S \gtrsim 850$ GeV.
  \item $(3, 4)_{+1/6}$, with $m_S \gtrsim 900$ GeV.
  \item $(3, 5)_{-1/3}$, with $m_S \gtrsim 1.1$ TeV.
  \item $(3, 5)_{+2/3}$, with $m_S \gtrsim 1.2$ TeV.
\end{itemize}

We assume that all members of the $X$-multiplet are close enough in mass that they are observed as a single $X$-onium resonance. Another option would be separated signatures, in which, for example, a diphoton signal would come mainly from the $|Q|_{\text{high}}$ state, while the $W^+W^-$ signature arises mainly from the $|m|_{\text{low}}$ state/s, possibly at different mass. For half-integer $SU(2)_W$ representations, the $m = \pm 1/2$ states would not appear wide in the $W^+W^-$ channel, given its resolution, even for $\lambda_{XH} \simeq 1$.

We note again that $X \sim (3, 4)_{+1/6}$ can be as light as 450 GeV only if $X_{+5/3}$ is the lightest state and the mass splitting is large enough to let all the other states decay to it via three body decay. We further discuss this possibility in the next section, in the context of the second scenario we study.

**Large apparent $\Gamma_S$.** The mass splitting between members of an $SU(2)_W$ multiplet may cause an apparent large width in the $X$-onium diphoton signal. To this end, it is important that the contribution to the diphoton events is not completely dominated by a single member of the multiplet. However, since the contribution of a particle of charge $Q$ to the diphoton signal is proportional to $Q^4$, a single member dominance is the case more often than not. For example, for the $(3, 2)_{-5/6}$ multiplet, the contribution of the $Q = -4/3$ particle is 256 times larger than that of the $Q = -1/3$ particle. From the representations in table 1, only two could result in an apparent large diphoton width:

\begin{itemize}
  \item $(3, 4)_{+1/6}$, with $\sigma_{-4/3}^{\gamma\gamma}/\sigma_{+5/3}^{\gamma\gamma} \simeq 0.40$.
  \item $(3, 5)_{-1/3}$, with $\sigma_{+5/3}^{\gamma\gamma}/\sigma_{-7/3}^{\gamma\gamma} \simeq 0.26$.
\end{itemize}

The mass splitting between two extreme bound states of an $SU(2)_W$ $n$-tuple is $\Delta m_S \simeq -\lambda_{XH}^2(n - 1)v^2/(2m_S)$. Therefore, a quartic coupling of size

$$|\lambda_{XH}| = \frac{1}{50(n - 1)} \frac{m_S^2}{v^2},$$

would saturate an estimated 1\% mass resolution of the diphoton signal (see, e.g. [70]). Such a small quartic coupling is allowed by EWPM and has no observed impact on Higgs couplings. Note that in order to understand whether the whole multiplet contributes to the resonance, or just the lightest member, one needs to make sure that the $W$-mediated decays within the multiplet, $X_m \to X_{m \pm 1}W^\pm$ (eqs. (4.3) and (4.4)), are not faster than the decay rate of $S$. This condition is generally satisfied below the $m_W$ threshold.
7.3 Back to our model examples

Let us now describe the phenomenology of the QCD bound state for our two benchmark scenarios of section 5.

7.3.1 Degenerate SU(2)$_W$-quartet

In this scenario with $m_X = 520$ GeV, the bound state has a mass $m_S = 1036$ GeV, with possible small splitting between the various $S_Q$ states. It exhibits the following features:

- $\gamma\gamma$: possible large apparent width in diphoton signals, with $\sigma_{\gamma\gamma}^{13} \simeq 0.25$ fb.
- $gg$: possible large apparent width in dijet signals, with $R_{gg/\gamma\gamma} \simeq 11$.
- $W^+W^-$: large $W^+W^-$ signal, with $R_{WW/\gamma\gamma} \simeq 25$.
- $ZZ$: enhanced $ZZ$ signal, with $R_{ZZ/\gamma\gamma} \simeq 11$.

In particular, a discovery of $S$ with $m_S$ slightly above TeV is, in this case, within the reach of upcoming diphoton searches.

7.3.2 Non-degenerate SU(2)$_W$-quartet

This is an example in which the bound state search is more powerful than the direct searches of $X_Q$ due to the lack of sensitivity for the three body final state $Wjj$ which would allow quartet as light as 250 GeV. Diphoton searches for $S_Q$ exclude $m_S \leq 900$ GeV, which corresponds to $m_X \lesssim 450$ GeV. At the 13 TeV with increased luminosity we expect a resonance which exhibits the following features:

- $\gamma\gamma$: possibly two resolved diphoton resonances, with a total diphoton signal $\sigma_{\gamma\gamma}^{13} \simeq 0.58$ fb.
- $gg$: wide dijet signal, with $R_{gg/\gamma\gamma} \simeq 11$.
- $W^+W^-$: large $W^+W^-$ signal, with $R_{WW/\gamma\gamma} \simeq 25$.
- $ZZ$: enhanced $ZZ$ signal, with $R_{ZZ/\gamma\gamma} \simeq 11$.

8 Summary and conclusions

The LHC search for new physics at or below the TeV scale is far from complete, even for strongly interacting particles. New particles might have surprising features, different from those predicted by the commonly studied extensions of the standard model. We studied the phenomenology of color-triplet scalar particles transforming in non-trivial representation of SU(2)$_W$ and potentially carrying exotic EM charges. Our main results are as follows.

- Color-triplet scalars ($X$), transforming in exotic representations of SU(2)$_W$ with masses at a few hundred GeV, are far from being experimentally excluded.
Depending on the electromagnetic charges of such colored scalars, their dominant decay modes could be into three or four body final states. Some of these decay topologies, in particular the $W^\pm jj$ one, are essentially unexplored by current analyses.

In large parts of the parameter space, $XX^\dagger$ for exotic $X$ would form a QCD-bound state ($S$). It is easy to find examples where the observation of di-electroweak boson (e.g. diphoton) resonance at $m_S$ will precede the direct discovery of $X$.

If $X$ is an SU(2)$_W$-non-singlet, the phenomenology of $S$ might involve intriguing features, such as $WW$ resonance at the same invariant mass as the diphoton resonance or somewhat removed from it, and a large apparent width for $S$.

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A ($W^+ jj)(W^- jj)$ final state

A dedicated search for the three body decay topology $Wjj$ has not been performed by the experimental collaborations. There are, however, a few analyses which are potentially sensitive to this final state. As detailed in section 3, we simulated our signal in MC simulation and compared between the efficiencies of these analyses for our ($W^+ jj)(W^- jj)$ signal and for the topologies that originally served as benchmark models. We stress that the limits obtained in this way should be taken as indicative of the sensitivity of certain searches to our final state, rather than as a complete recast of the analyses.

The search for exotic vector-like fermions decaying into $Wj$ [36]. The bounds from this search are presented in figure 1 as they are found to be the most sensitive ones. We find that the selection efficiencies of our signal and the targeted topology ($W^+ j)(W^- j)$ are comparable. Yet, the binned analysis performed by the collaborations relies on the mass reconstruction of the parent fermion. Therefore our signal, originating from a three-body decay, suffers from a broadening of the $m_{Wj}$ distribution. We conservatively estimate this reduction to be between 30% and 50%. Under this assumption, this search does not rule out the existence of a light quartet with all components cascading down to $Wjj$, except at a small mass window between 375–440 GeV. At the Tevatron both CDF and D0 [34, 35] performed a similar search for fourth generation quarks in the mass range 200–500 GeV.
Assuming similar reduction in the efficiency, these searches give the best sensitivity at the low mass ranges. They exclude, for example, an SU(2)$_W$ quartet below 250 GeV.

**The search for supersymmetric multi-jet with 0/1/2 leptons.**

- **Fully hadronic:** almost half of the events of our signal are purely hadronic. However, multi-jet searches suffer from the large QCD background and do not exclude a scalar color-triplet, even when taking into account the high multiplicity of a quartet. (These searches are more effective for gluinos, which have a significantly higher cross section [22].)

- **One lepton:** the final state contains a single lepton, jets and missing energy, which is common to many SUSY scenarios. In particular, the ATLAS search of ref. [72] targets, among others, the double production of first and second generation squarks $\tilde{q}$ decaying into $W^\pm j\chi_0$ via on-shell chargino. An important parameter for this signal is $x = \Delta m(\chi^+, \chi^0)/\Delta m(\tilde{q}, \chi^0)$. For $x = 1$ the squarks and chargino are degenerate, while for $x = 0$ the chargino and neutralino are degenerate. Originating from a three body decay, the kinematics of the $W$’s in our signal resemble more the low $x$ case. Taking $x = 0.2$ as representative of the low $x$ region, we find that the efficiency of our signal is lower by 50% than the one of the targeted signal. This reduction originates from the lower missing energy which, however, is partially compensated by the enlarged jet activity. Taking $x = 0.8$ as representative of the high $x$ region, we find that the efficiency reduction becomes less than 10%. The similar analysis at 13 TeV has lower sensitivity as it typically targets higher masses [73, 74]. The search reported in ref. [75] might also have some sensitivity to this signal, but it relies on the assumption that there is no significant signal contribution to events with five or six jets, which is not the case for our ($W^+jj$)($W^-jj$) final state.

- **Two leptons:** searches for a final state containing two leptons, missing energy and jets have a potentially similar reach, but pay a higher price in the leptonic branching ratio of the $W$ bosons. Therefore, they do not provide the best limits on our signal.

**The search for first or second generation leptoquarks.** The LQ searches typically suffers from a 25% reduction in the efficiency for our signal. This, together with the small leptonic $W$ branching ratios, yield bounds that are insignificant. We note that a mixed $(e^\pm j)(\mu^\mp j)$ search, which is currently not done by the collaborations, may have better sensitivity due to lower expected background.

**Searches for various states containing $b$ jets.**

- The CMS 7 and 8 TeV analyses [39, 76] search for heavy top-like quark ($t'$) decaying to $Wb$ final state. These searches might be sensitive to a $Wbj$ topology. Yet, as previously discussed, the $t'$ mass reconstruction weakens the reach of this search to the $Wbj$ topology. We again estimate this reduction to be between 30% and 50% and show the resulting bounds in figure 1. The same is done for the heavy bottom-like quark searches [32, 44, 45].
The CMS RPV-SUSY search \cite{32} for $\tilde{b} \rightarrow tj$, where $\tilde{b}$ is the bottom squark, could have some sensitivity to $Wbj$ topology. However, it requires the reconstruction of $t$ quarks which reduces significantly the sensitivity to our signal.

SUSY stop searches, e.g. \cite{77}, look for a single lepton, missing energy and $b$-jets final state. We find these searches to be less sensitive than the heavy quark searches, as in the SUSY multi-jet searches with 1 lepton.

We conclude that the $Wjj$ decay mode is presently poorly constrained, irrespective of the flavor of the jets in the final state.

Precision cross-section measurements. Precision measurement of the $tt$ and $W^+W^-$ cross sections might probe best the low mass region of a $(W^+jj)(W^-jj)$ signal. However, for $m_X \geq 250$ GeV we find that these are not sensitive even at multiplicity as high as $n = 5$; the argument goes as follows. We consider the NNLO-NNLL $tt$ production cross section (see \cite{78} and references therein), with $m_t = 172.5$ GeV, and combine scale uncertainty and the uncertainty associated with variations of the PDF and $\alpha_s$ (see \cite{79-82}). At $m_X = 250$ GeV, the production cross section for a quintuplet is below the theoretical uncertainty, assuming the efficiency of the $tt$ search to be 50\% smaller than the efficiency for the $tt$ sample itself. This is a plausible estimate in the case of the $Wbj$ topology, and a conservative one for the $Wjj$ topology, even if we allow a large mistagging rate. Therefore, a quintuplet at 250 GeV is not constrained by the $tt$ measurements. As for the $W^+W^-$ cross section measurements, the relevant analyses veto on $N_j \geq 1$. Since our signal contains many jets in the final state, it would not contribute significantly to these measurements.

B $(W^+W^+jj)(W^-W^-jj)$ final state

There are no dedicated searches for the four body $WWjj$ decay mode, but other searches are potentially sensitive to it. For the fully hadronic final states and for the ones containing only one or two leptons, conclusions similar to those made for the $Wjj$ decay mode hold. However, for this topology, the most promising search strategy is to look for multilepton final states. The low SM background compensates for the branching ratio suppression of four $W$’s decaying leptonically.

We analyze the RPV multilepton CMS search \cite{32,45} which does not rely on any missing energy cut. This analysis contains many exclusive signal regions, depending on the number of leptons, the presence of hadronically decaying $\tau$, the presence of $b$ jets, and the number of opposite-sign-same-flavor (OSSF) lepton pairs. We consider the low background regions, with four leptons, zero hadronic $\tau$’s and 1 pair of OSSF leptons, summing over all $S_T$ bins. To be conservative, we allow the number of background events to fluctuate up by 95\% C.L. and the number of signal events to fluctuate down by 95\% C.L., assuming Poisson statistics. We take $N_{\text{sig}} = L\sigma\text{BR}_{4W\rightarrow4\ell,\text{OSSF}}$ with very high efficiency $\epsilon = 80\%-90\%$.

A somewhat stronger bound comes from the ATLAS analyses of ref. \cite{44}. For this, we consider the two overlapping signal regions, SR3L1 and SR0b1, with the corresponding bounds of $\sigma_{\text{SR3L1}} \leq 0.59$ fb and $\sigma_{\text{SR0b1}} \leq 0.37$ fb, set at 95\% C.L.. (For the exact description
of these signal regions we refer the reader to ref. [44].) Since this search was specifically designed to be applicable to any SUSY RPV scenario, we assume the efficiency for our signal to be similar to the one quoted. We therefore use $\epsilon = 2\%$—$5\%$. The resulting limits are presented in figure 1.

C \((\ell^{\pm}t)(\ell^{\mp}t)\) final state

Similar to the four-body decays, $X \to t\ell^{\pm}$ decay would be captured by the multi-lepton searches aiming at RPV SUSY signals. For this signature we estimate the reach of the CMS 8 TeV search [45] in the signal region with four leptons, zero hadronic taus, one pair of OSSF leptons and one tagged $b$-jet. As before, we allow upward fluctuation of the background and downward fluctuation of the signal, both within 95\% C.L. Assuming efficiency of 60\%–80\%, we find an excluded cross section of $\sigma_{XX}^X \leq 5$–6.6 fb. The resulting bounds as a function of $m_X$ are presented in figure 1.

We note that the 13 TeV analysis of CMS [32] veto $b$-jets, while the ATLAS 13 TeV analysis [44] uses large jet multiplicity ($N_j \geq 6$) and relatively large missing energy ($E_T^{\text{miss}} \geq 200$) GeV. Both of these searches are therefore less sensitive to our signal in this case.

D Running of gauge coupling constants

At one loop,

$$\alpha(\mu)^{-1} = \alpha(\mu_0)^{-1} + \frac{(b^\text{SM} + b^X)}{2\pi} \log\left(\frac{\mu}{\mu_0}\right) \tag{D.1}$$

with (we use the common GUT inspired definition $g_1 = \sqrt{3/5}g'$)

$$b_1^\text{SM} = -\frac{41}{10}, \quad b_2^\text{SM} = \frac{19}{6}, \quad b_3^\text{SM} = 7,$$

$$b_1^X = -\frac{1}{5}RnY^2,$$

$$b_2^X = -\frac{1}{3}RC(n) = -\frac{1}{36}Rn(n^2 - 1),$$

$$b_3^X = -\frac{1}{3}nC(R) = \begin{cases} -\frac{6}{5} & \text{for } R = 3 \\ -\frac{5}{6} & \text{for } R = 6 \end{cases}, \quad (D.2)$$

where $C(n)$ [$C(R)$] is the Casimir of the $n$ [$R$] representation of SU(2) [SU(3)]. For numerical evaluation, we use, at $m_Z = 91.1876$ GeV [83]

$$\alpha_{\text{EM}}(m_Z) = 127.94^{-1}, \quad s_W^2(m_Z) = 0.22333, \quad \alpha_3(m_Z) = 0.1185$$

$$\alpha_1 = \frac{5}{3s_W^2} \alpha_{\text{EM}}, \quad \alpha_2 = \frac{1}{s_W^2} \alpha_{\text{EM}}, \quad (D.3)$$

and $m_W = 80.385$ GeV.
Table 6. Dimension six operators generated by integrating out a scalar \( X(R, n)_Y \).

E Effective operators

Consider a scalar \( X(R, n)_Y \) with the Lagrangian given in eqs. (2.1) and (2.2). The impact of \( X \) on the SM fields is mainly captured by the dimension six operators, generated at one-loop order upon integration out of \( X \). Refs. [84, 85] compute the Wilson coefficient of these effective interactions for a general scalar. We present their results in table 6. Note that even though this list is not completely independent when the Higgs and gauge bosons equations of motion are considered, we find it convenient for our purposes to determine the oblique parameters and the Higgs couplings, as long as no redundancy is used when considering physical parameters.

F Oblique parameters

A scalar \( X(R, n)_Y \) alters the vacuum polarization amplitudes of the EW gauge fields. These effects are conveniently parameterized by the oblique parameters \( S, T \) and \( U \) [61] and \( V, X, Y \) and \( W \) (for a review see [63]). The leading contributions to the oblique parameters read

\[
T = \frac{v^2}{\alpha} c_T = \left( \frac{v^2}{4608\pi^2\alpha} \right) \left( \frac{\lambda_{XH}^2}{m_X^2} \right) R n(n^2 - 1),
\]

\[
S = 16\pi v^2 c_{WB} = \left( \frac{v^2}{144\pi} \right) \left( \frac{\lambda_{XH}^2}{m_X^2} \right) R n(n^2 - 1) Y,
\]

\[
Y = 2m_W^2 c_{2B} = \frac{g^2}{240\pi^2} \frac{m_W^2}{m_X^2} R n Y^2,
\]

\[
W = 2m_W^2 c_{2W} = \frac{g^2}{2880\pi^2} \frac{m_W^2}{m_X^2} R n(n^2 - 1),
\]

(F.1)
The quartic scalar couplings $\lambda_{XH}$ and $\lambda'_{XH}$ modify the light Higgs couplings from their SM values. For an $X(R, n)_Y$ representation, these modifications read

\[ \delta c_\gamma = 4\pi^2 v^2 \left( c_{BB} + c_{WW} - c_{WB} \right) = \frac{v^2}{144m_X^2} nR \left( \frac{n^2 - 1}{2} + 3Y^2 \right) \lambda - \frac{(n^2 - 1)Y\lambda'}{2}, \]

\[ \delta c_g = 48\pi^2 v^2 c_{GG} = \frac{v^2 \lambda_{XH}}{4m_X^2} nC(R), \]

\[ \delta c_W = -\frac{c_H v^2}{2} + \frac{2c_W^2 v^2}{c_W - s_W} c_T - \frac{16\pi v^2}{c_W - s_W} c_{WB} \]

\[ = \frac{v^2}{12m_X^2} nR \left[ -\frac{\lambda_{XH}^2}{16\pi^2} + \frac{c_W^2 (n^2 - 1)\lambda_{XH}^2}{96(c_W^2 - s_W^2)} + \frac{\alpha Y(n^2 - 1)\lambda_{XH}'}{6\pi(c_W^2 - s_W^2)} \right], \]

\[ \delta c_Z = -\frac{c_H v^2}{2} - \frac{v^2 \lambda_{XH}^2}{192\pi^2 m_X^2} nR, \]

\[ \delta c_f = -\frac{c_H v^2}{2} = -\frac{v^2 \lambda_{XH}^2}{192\pi^2 m_X^2} nR. \]

The $hgg$ and $h\gamma\gamma$ couplings are computed using the Higgs effective low energy theory [86]:

\[ \delta c_\gamma = \frac{R}{24} \sum_Q Q^2 v^2 \frac{\partial \log M^2_Q}{\partial v}, \]

\[ \delta c_g = \frac{C(R)}{2} \sum_Q v^2 \frac{\partial \log M^2_Q}{\partial v}, \]

where

\[ M^2_Q = m_X^2 + \left( \lambda_{XH} - \frac{\lambda'_{XH} Q}{2} \right) v^2. \]

Other couplings are computed by their definition in terms of the Wilson coefficients, for which we use the results of refs. [83, 87].

For our numerical results we use table 14 of [88] with $B_{BSM} = 0$. We take as a concrete example the case of $X \sim (3, 4)_{+1/6}$. The exact results, including EWPM constraints, are shown in figure 3. The constraints on $\lambda_{XH}$ and $\lambda'_{XH}$ are rather mild and do not affect our conclusions.

### H Quartic coupling running

In this appendix we obtain the one loop $\beta$ function for the four quartic couplings of the scalar potential, where

\[ \frac{d\lambda}{d \log \mu} = \beta_\lambda. \]

We use the normalization $\lambda_H = m_H^2/v^2$ for the SM Higgs quartic coupling. Other quartic couplings are defined in eq. (2.2). EW corrections of the order $g^2, g^2$ are neglected. The $\lambda_H$ RGE is given by

\[ \beta_\lambda = \frac{1}{16\pi^2} \left[ 12\lambda_H^2 + 2Rn\lambda_H^2 + \frac{Rn(n^2 - 1)}{24}\lambda_H^2 + \cdots \right]. \]
Figure 3. Higgs decay (blue) and EWPM (gray) constraints on the quartic couplings $\lambda_{XH}$ and $\lambda'_{XH}$ at 95% C.L. for $X \sim (3, 4, +1/6)$ at $m_X = 450$ GeV. Blue point is the best fit value from the Higgs data. Black line is the best fit value from EWPM.

where $\ldots$ indicates other SM contributions (coming from, e.g., the top-quark). The $\lambda_X$ RGE is given by

$$\beta_{\lambda_X} = \frac{1}{16\pi^2} \left[ (2Rn + 8)\lambda_X^2 + 4\lambda_{XH}^2 + \frac{13}{3} g_4^2 - 16g_3^2 \lambda_X \right].$$

(H.3)

As for the mixed terms, we find:

$$\begin{align*}
\beta_{\lambda_{XH}} &= \frac{1}{16\pi^2} \left[ 2(Rn + 1)\lambda_X\lambda_{XH} + 6\lambda_H\lambda_{XH} + 4\lambda_{XH}^2 + \left( \frac{n^2 - 1}{4} \right) \lambda_X^2 - 8g_3^2 \lambda_{XH} \right], \\
\beta_{\lambda'_{XH}} &= \frac{1}{16\pi^2} \left[ 2\lambda_X\lambda'_{XH} + 2\lambda_H\lambda'_{XH} + 8\lambda_{XH}\lambda'_{XH} - 8g_3^2 \lambda'_{XH} \right].
\end{align*}$$

(H.4)

I $S \rightarrow VV$ decays

For the $S \rightarrow VV$ decays, we use (see [4] and references therein)

$$\Gamma(S_Q \rightarrow V_1 V_2) = \frac{R}{8\pi(1 + \delta_{V_1 V_2})} \lambda^{1/2}(M_S^2, m_{V_1}^2, m_{V_2}^2) \left| \psi(0) \right|^2 \left| \overline{\mathcal{M}}_{V_1 V_2} \right|^2,$$

(I.1)

where $\lambda[x, y, z]$ is defined below eq. (5.3), and the squared amplitudes are given by [89, 90]

$$\left| \overline{\mathcal{M}}_{gg} \right|^2 = C_2(R)^2(16\pi^2 \alpha_s^2),$$

$$\left| \overline{\mathcal{M}}_{gg} \right|^2 = 8\pi^4 Q^2,$$

$$\left| \overline{\mathcal{M}}_{ZZ} \right|^2 = \frac{8e^4(m - Q s_W^2)^4}{s_W^2 c_W^2} + (\lambda_{XH}^m)^2 + O\left( m_Z^2, m_H^2 / m_X^2 \right),$$

$$\left| \overline{\mathcal{M}}_{Z^*} \right|^2 = 8\left( \frac{m - Q s_W^2}{s_W c_W} \right)^2 Q^2 e^4,$$

$$\left| \overline{\mathcal{M}}_{WW} \right|^2 = e^4 (n^2 - 1 - 4m^2)^2 + (\lambda_{XH}^m)^2 + O\left( m_W^2, m_H^2 / m_X^2 \right),$$

$$\left| \overline{\mathcal{M}}_{hh} \right|^2 = \left[ -\lambda_{XH}^m - 3\lambda_{XH}^m \left( \frac{m_h^2}{4m_X^2 - m_h^2} \right) + (\lambda_{XH}^m)^2 \left( \frac{2v^2}{m_X^2 - m_h^2} \right) \right]^2$$

$$= (\lambda_{XH}^m)^2 + O\left( m_h^2 / m_X^2, v^2 / m_X^2 \right).$$

(I.2)
Here $C_2(R)$ is the quadratic Casimir, with $C_2(3) = 4/3$. The quartic couplings $\lambda_{\chi H}^m = \lambda_{\chi H}^m - (m/2)\lambda_{\chi H}^0$ (with $Q = m + Y$) and $\lambda_{\chi H}^0$ change the decay rates of $S_Q$ into $WW$ and $ZZ$ final states in a mild way, and generate $S_Q \rightarrow hh$ decays.

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