Simplified dark matter models with a spin-2 mediator at the LHC

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Abstract. We consider simplified dark matter models where a dark matter candidate couples to the standard model (SM) particles via an s-channel spin-2 mediator, and study constraints on the model parameter space from the current LHC data. Our focus lies on the complementarity among different searches, in particular monojet and multijet plus missing energy searches and resonance searches. For universal couplings of the mediator to SM particles, missing-energy searches can give stronger constraints than WW, ZZ, dijet, dihiggs, t ¯t, b ¯b resonance searches in the low-mass region and/or when the coupling of the mediator to dark matter is much larger than its couplings to SM particles. The strongest constraints however come from diphoton and dilepton resonance searches. Only if these modes are suppressed, missing-energy searches can be competitive in constraining dark matter models with a spin-2 mediator.

1 Introduction

Convincing astrophysical and cosmological observations for the existence of dark matter (DM) provide us one of the strong motivations to consider physics beyond the standard model (SM). The search for DM is thus one of the main pillars of the LHC physics program.

As the nature of DM is known so little, a so-called simplified-model approach [1] has been widely adopted, and concrete simplified DM models have recently been proposed by the LHC DM working group to conduct the systematic DM searches at the LHC Run-II [2]. Following the proposal, the Run-I data as well as the early Run-II data have already been analysed to constrain simplified DM models with s-channel spin-1 and spin-0 mediators, see e.g. [3][12]. On the other hand, the model with a spin-2 mediator [13][14] has not been fully explored for the LHC yet—it is one of the next-generation simplified DM models [15].

In this article, we consider simplified DM models where a DM candidate couples to the SM particles via an s-channel spin-2 mediator, and study constraints on the model parameter space from searches in final states with and without missing energy in the current LHC data. This work follows the DMStMP framework [16][18], which provides the DM model files for event generators such as MadGraph5 AMC@NLO [19] as well as for DM tools such as MicrOMEGAs [20][22] and MadDM [23][24]. The same framework was used previously to study the cases of s-channel spin-1 and spin-0 mediators.

We note that, to keep the analysis of the LHC constraints fully general, we do not impose any astrophysical constraints like relic density or (in)direct detection limits on the DM candidate, as these partly depend on astrophysical assumptions. Moreover, in a full model, the DM may couple to other new particles that are irrelevant for the collider phenomenology discussed here. We refer readers to [13][14] for the astrophysical constraints, and to [25] for a discussion of spectral features in indirect detection.

The article is organised as follows. The simplified model is presented in Section 2 and the production and decays of the spin-2 mediator in Section 3. The re-interpretation of the LHC results is discussed in Section 4. Section 5 contains a summary and conclusions. Supplemental material for recasting is provided in the Appendix.

2 Model

Gravity-mediated DM was proposed in [13][14], where the dark sector communicates with the SM sector through a new spin-0 particle (radion) and spin-2 particles (Kaluza-Klein (KK) gravitons) in warped extra-dimension models as well as in the dual composite picture.

In this work, following the approach of simplified DM models, we consider DM particles which interact with the
SM particles via an s-channel spin-2 mediator. The interaction Lagrangian of a spin-2 mediator ($Y_2$) with DM ($X$) is given by

$$\mathcal{L}_X^{Y_2} = -\frac{1}{4} g_{X} T_{\mu\nu} Y_{2}^{\mu\nu},$$

(1)

where $A$ is the scale parameter of the theory, $g_{X}^T$ is the coupling parameter, and $T_{\mu\nu}^X$ is the energy–momentum tensor of a DM field. Here, we consider three types of DM independently; a real scalar ($X_R$), a Dirac fermion ($X_D$), and a vector ($X_V$). The interaction with SM particles is obtained by

$$\mathcal{L}_{SM}^{Y_2} = -\frac{1}{A} \sum_i g_i^T T_{\mu\nu} Y_{2}^{\mu\nu},$$

(2)

where $i$ denotes each SM field, i.e. the Higgs doublet ($H$), quarks ($q$), leptons ($\ell$), and $SU(3)_C$, $SU(2)_L$ and $U(1)_Y$ gauge bosons ($g, W, B$). Following [26,27] we introduce the phenomenological coupling parameters

$$g_i^T = \{g_H, g_q^T, g_\ell^T, g_g, g_W, g_B\}$$

(3)

without assuming any UV model[1]. The energy–momentum tensors of the DM are

$$T_{\mu\nu}^{X_R} = -\frac{1}{2} g_{X} \rho \partial_\mu X_R \partial_\nu X_R - m_X^2 X_R^{\mu\nu},$$

(4)

$$T_{\mu\nu}^{X_D} = -g_{X} \partial_\mu X_D \partial_\nu X_D + \frac{1}{2} g_{X} \rho \partial_\mu (\gamma_\nu \partial_\rho X_D) + \frac{1}{2} X_D \partial_\mu (\gamma_\nu \partial_\rho X_D) - \frac{1}{4} \partial_\mu (X_D \gamma_\nu \partial_\rho X_D) - \frac{1}{4} \partial_\nu (X_D \gamma_\mu \partial_\rho X_D),$$

(5)

$$T_{\mu\nu}^{X_V} = -g_{X} \left( -\frac{1}{4} F_{\rho\sigma} F^{\rho\sigma} + \frac{m_X^2}{2} X_V^{\mu\nu} \right) + F_{\mu\nu} F_\nu^T + m_X^2 X_V^{\mu\nu},$$

(6)

where $F_{\mu\nu}$ is the field strength tensor. Those of the SM fields are similar; see e.g. [28] for the explicit formulae.

Complying with the simplified-model idea, it is instructive to consider universal couplings between the spin-2 mediator and the SM particles:

$$g_{SM} \equiv g_H^T = g_q^T = g_\ell^T = g_g^T = g_W^T = g_B^T.$$  

(7)

With this simplification, the model has only four independent parameters, two masses and two couplings:

$$\{m_X, m_Y, g_X/A, g_{SM}/A\},$$

(8)

where we dropped the superscript $T$ for simplicity. Such a universal coupling to SM particles is realised, e.g., in the original Randall–Sundrum (RS) model of localised gravity [29]. The parameters are related as

$$m_Y/A = x_1 k/M_{Pl},$$

(9)

where $x_1 = 3.83$ is the first root of the Bessel function of the first kind, $k$ is the curvature of the warped extra dimension, and $M_{Pl} = 2.4 \times 10^{18}$ GeV is the reduced four-dimensional Planck scale. On the other hand, in the so-called bulk RS model [30,31], where the SM particles also propagate in the extra dimension, $g_i^T$ can take different values depending on the setup.

In [28], the SM sector of the above model was implemented in FeynRules/NloCT [32,33] (based on [34–36]), and the $Y_2$ production and decay rates at next-to-leading order (NLO) QCD accuracy were presented. In this work, we include the three DM species ($X_R$, $X_D$, $X_V$) with the corresponding interactions, and add the model into the DMStimp framework [37] as the simplified DM model with a spin-2 mediator.
3 Phenomenology at the LHC

3.1 Decay of the spin-2 mediator

Regarding LHC phenomenology, let us begin by discussing the spin-2 mediator decays. The partial widths for the decays into a pair of spin-0 \((S = X_R, h)\), spin-1/2 \((F = X_D, q, \ell)\) and spin-1 \((V = X_V, g, \gamma, Z, W)\) DM or SM particles are given by

\[
\Gamma_S = \frac{g_X^2 m_Y^2}{160 \pi A^2} \beta_S^3, \tag{10}
\]

\[
\Gamma_F = \frac{g_X^2 N_F N_Y m_Y^3}{160 \pi A^2} \beta_F^3 \left(1 + \frac{8}{3} r_F\right), \tag{11}
\]

\[
\Gamma_V = \frac{g_X^2 N_F N_Y m_Y^3}{40 \pi A^2} \beta_V f(r_V), \tag{12}
\]

where \(\beta_i = \sqrt{1 - 4r_i}\) with \(r_i = m_Y^2/m_i^2, \quad \beta_r = g_B \cos^2 \theta_W + g_W \sin^2 \theta_W\) and \(g_Z = g_B \sin^2 \theta_W + g_W \cos^2 \theta_W\) with the weak-mixing angle \(\theta_W\), and \(f(r_V) = 1 + \frac{3}{2} \beta^2_H - r_V (3 - \frac{20}{5} \beta^2_H - \frac{20}{3} \beta^2_H + \frac{14}{5} \beta^2_H)\) with \(\beta_H = g_H/g_V\).

For gluons and photons, \(\beta_g = 0\), while \(\beta_H = 1\) for vector DM. The factors \(N_F = 1/2\) for neutrinos and \(N_F = 1/2\) for two identical particles, and are unity otherwise; \(N_C^{F,V}\) is the number of colours. We note that \(B(Y_2 \rightarrow Z\gamma) = 0\) for \(g_W = g_B\) as the decay rate is proportional to \(g_Z^2 = [g_W - g_B \cos \theta_W \sin \theta_W]^2\). We see that, due to the different overall prefactors, the partial widths become larger in order of scalar, fermion, and vector DM. Moreover, the different powers \((5, 3, 1)\) of the velocity factor \(\beta_i\) indicate that the decay proceeds mainly via a D, P, and S wave for the scalar, fermion, and vector case, respectively.

Figure 1 shows the \(Y_2\) total width scaled by the mass, \(\Gamma_Y/m_Y\), and the decay branching ratios for the case that only decays into SM particles are allowed. \textsc{MadWidth} [38] provides the partial decay rates numerically for each parameter point. In Table 1 we provide the explicit values for a few representative mass points. We see that, for a universal coupling \(g_{SM}\), decays into gluons and light quarks, leading to a dijet signature, are completely dom-
nant (≥80% depending on \(m_Y\)). The diphoton channel has 4–5% branching ratio; other diboson channels (WW and ZZ) as well as \(tt\) are important as well when kinematically allowed. Finally, it is important to note that decays into neutrinos have 3–4% branching ratio, leading to missing energy signatures independent of decays to DM.\(^2\) The width is proportional to \(m_Y^4\), and from the upper panel in Fig. 1 we see that for \(g_{SM}/A (3 \text{ TeV})^{-1}\), the resonance is always very narrow (\(\Gamma_Y/m_Y < 1\%\)) up to \(m_Y \sim 1 \text{ TeV}\). Note here, that \(A\) is simply a scale parameter, not a physical cut-off of the theory.

When decays into DM are allowed, their relative importance depends on \(g_X\) and the type of DM (scalar, Dirac or vector) as illustrated in Fig. 2; see also Eqs. (10)–(12). Two mass scales are considered: \(m_Y = 100 \text{ GeV}\) and 1 TeV, with \(m_X = 10 \text{ GeV}\) and \(g_{SM} = 1\).\(^3\) We see that decays into DM can be important and even dominant, but the resonance remains narrow for any choice of \(A \geq 3 \text{ TeV}\) for \(m_Y \lesssim 1 \text{ TeV}\). Another important observation is that for scalar DM (\(X_R\)), for \(g_X \approx g_{SM}\) the decay into \(Y_2 \to X_R X_R\) is practically irrelevant; one needs \(g_X/g_{SM} \approx 3\) for the decay into DM to exceed the one into neutrinos, and \(g_X/g_{SM} \approx 5–6\) to reach the 10% level. For Dirac (\(X_D\)) and vector (\(X_V\)) DM, the decays into DM and into neutrinos are of comparable magnitude at \(g_X \sim g_{SM}\), both contributing to missing-energy signatures. For \(g_X/g_{SM} = 2\), the branching ratio of \(Y_2 \to X_D X_D (X_V X_V)\) attains about 10% (20%). These differences depending on the type of DM will be important later for the collider limits.

### 3.2 Production of the spin-2 mediator

Turning to the production modes, the potentially interesting channels are inclusive \(Y_2\) production (\(pp \to Y_2\)), as well as the production with an extra hard tagging jet (\(pp \to Y_2 j\)) or an electroweak boson (e.g., \(pp \to Y_2 \gamma\)). With the \(Y_2\) decaying into SM particles, the former gives resonant peak signatures (without missing energy). On the other hand, the latter two give the typical monojet or monophoton signatures when the mediator decays invisibly. Moreover, the latter two play a role in the low-mass resonance search in dijet events with initial-state radiation (ISR) as seen later.

The \(Y_2\) production cross sections at NLO QCD accuracy for \(pp\) collisions at 13 TeV are depicted in Fig. 3 as a function of the mediator mass.\(^4\) We employ MADGraph5_AMC@NLO\(^5\) to calculate the cross sections and generate events with the LO/NLO NNPDF2.3.\(^6\) The factorisation and renormalisation scales are taken at the sum of the transverse masses of the final states as a dynamical scale choice. In our simplified model, the cross sections depend solely on \(g_{SM}/A\) and scale with \((g_{SM}/A)^2\).

The dashed lines showing \(g_{SM}/A = (10 \text{ TeV})^{-1}\) are therefore an order of magnitude below the corresponding solid lines for \(g_{SM}/A = (3 \text{ TeV})^{-1}\). Also noteworthy is the fact that \(pp \to Y_2\) is mostly gluon-initiated for the low-mass case \((39\%\), 97%, and 83%\) of the LO total rate for \(m_Y = 100 \text{ GeV}, 1 \text{ TeV},\) and 5 TeV, respectively, stem from \(gg\) fusion. Since the radiation of an initial-state photon (\(Z/W\)) can only occur in the quark-initiated process, \(Y_2 + \text{photon} (Z/W)\) production is much suppressed as compared to \(Y_2 + \text{jet}\) production. This is also the reason that the process has a huge \(K\) factor especially in the low-mass region \((28\%)\).

In the context of DM searches, the monojet signature is expected to give important constraints on the model. The fiducial cross sections for \(pp \to Y_2 j\) with \(p_{Tj} > 200 \text{ GeV}\) and \(|\eta_j| < 5\) are shown in Fig. 4, where one can estimate the monojet cross section by taking into account the \(Y_2\) branching ratio into DM particles (and/or neutrinos) when \(m_Y > 2m_X\). In Fig. 4 we also plot the fiducial cross sections for \(pp \to j + \nu\eta\) as a function of the DM mass, \(|\eta_j| < 2.5\). \(K\) factors are also shown in the lower panel as a reference.

\(^2\) These decay branching ratios were already presented in Eq. (39) for the case of the RS graviton. We repeat them here for the sake of completeness. Our numbers agree with Eq. (39) apart from a factor 1/2 for decays into neutrinos.

\(^3\) As can be deduced from Fig. 1 above the WW threshold up to high masses the picture does not change much apart from the \(tt\) and/or \(hh\) channels being open or not.

\(^4\) See also Fig. 12 (bottom) for \(\sigma(pp \to Y_2)\) at \(\sqrt{s} = 8 \text{ TeV}\).
separating the contributions from neutrinos (black lines) and DM (red lines) produced through the spin-2 mediator. For definiteness, we take \( m_Y = 500 \text{ GeV} \), \( \Lambda = 3 \text{ TeV} \), \( g_{\text{SM}} = 1 \) and compare \( g_X = 1, 2 \) and 10 for Dirac DM. As already seen in Fig. 3, their relative importance depends on \( g_X \). For \( m_Y < 2m_X \), a pair of DM is produced via the off-shell mediator and the cross section is strongly suppressed. Therefore, the neutrino contribution always dominates the monojet signature for the \( m_Y < 2m_X \) region even if \( g_X / g_{\text{SM}} = 10 \). For the other DM types, scalar and vector, the picture is similar, but the relative importance to the neutrino channel is different; see Fig. 3. This is one of the characteristic features of the spin-2 mediator DM model with universal couplings, as compared to the \( s \)-channel spin-1 and spin-0 models, whose mediators do not couple to charged leptons and neutrinos in the minimal setup [8].

4 Constraints from current LHC data

4.1 Searches with missing energy

The ATLAS and CMS experiments have been searching for new physics in a large variety of final states. As mentioned above, in the context of DM searches, the monojet signature is regarded as particularly interesting. In practice, at 13 TeV, the monojet analyses require one hard jet recoiling against \( E_T \), but allow for additional jets from QCD radiation. Therefore one can expect that multijet+\( E_T \) searches are also relevant [11, 42].

To work out the current constraints on the spin-2 mediator DM model from these searches, we consider the following early Run-II analyses:

- ATLAS monojet with 3.2 fb\(^{-1} \)[5],
- ATLAS 2–6 jets + \( E_T \) with 3.2 fb\(^{-1} \)[43].

In the monojet analysis [5], a simplified DM model with an \( s \)-channel spin-1 mediator is considered. Events are required to have at least one hard jet with \( p_T > 250 \text{ GeV} \) and \( |\eta| < 2.4 \), and a maximum of four jets with \( p_T > 30 \text{ GeV} \) and \( |\eta| < 2.8 \) are allowed. Several inclusive and exclusive signal regions (SRs) are considered with increasing \( E_T \) requirements from 250 GeV to 700 GeV. The multijet+\( E_T \) analysis [43] is designed to search for squarks and gluinos in supersymmetric models, where neutralinos lead to missing energy. Several SRs are characterised by minimum jet multiplicity from two to six; \( E_T > 200 \text{ GeV} \) is required for all SRs, while different thresholds are applied on jet momenta and on the azimuthal separation between jets and \( E_T \).

To reinterpret the above analyses in the context of our spin-2 mediator simplified DM model, we use CheckMATE2 [44], which is a public recasting tool providing confidence limits from simulated signal events and includes a number of 13 TeV analyses. We generate hadron-level signal samples by using the tree-level matrix-element plus parton-shower (ME+PS) merging procedure. In practice, we make use of the shower-\( k_T \) scheme [45], implemented in MadGraph5_\textsc{amc@nlo} [19] with Pythia6 [46], and generate signal events with parton multiplicity from one to two partons. We impose \( E_T > 200 \text{ GeV} \) and set \( Q_{\text{cut}} = 200 \text{ GeV} \) for the merging separation parameter at the parton level; these values are chosen for an efficient event generation without affecting the final results. The event rate is normalised to the \( pp \rightarrow Y_{2j} \) NLO cross sections shown in Fig. 2 (Note, however, that NLO corrections may also affect the shapes of the kinematic distributions, as shown for the spin-1 and spin-0 cases in [17]; a detailed study of this aspect will be reported elsewhere.)

It turns out that, for an on-shell mediator of given mass, the selection efficiencies are independent of the mass and spin of the invisible decay products. Moreover, contributions from off-shell production are negligible for the scenarios considered here. The efficiencies can thus be evaluated as a function of the mediator mass only; see also Appendix A. 1 In the following, we normalise the number of events with NLO cross sections, shown in Fig. 2 and the total branching ratio into invisible final states (DM and neutrino). We note that for a given mediator mass the leading jet for the spin-2 mediator case is harder and more forward than that for the spin-1 case. This is partly because the spin-2 mediator with a parton is produced not only through the \( q\bar{q} \) and \( g\bar{g} \) initial states but also dominantly through the \( gg \) initial state, and partly because the spin-2 mediator is also emitted from a gluon as well as from the \( gggY_2 \) and \( qgY_2 \) four-point vertices.
Figure 5 shows the ratio of signal events over the number of events excluded at 95% confidence level (CL), $S/S_{95}$, as a function of the mediator mass, for the three types of DM (taking $g_X = 1$ or 2 with $\Lambda = 3$ TeV, $g_{\text{SM}} = 1$ and $m_X = 10$ GeV as a benchmark case). As expected from the discussion in the previous section, the scalar DM case is the least constrained, with the $E_T$ coming dominantly from the neutrino channel; for $g_X = 1$ (2), we find the limit $m_Y \gtrsim 600$ (750) GeV from the monojet analysis and $m_Y \gtrsim 750$ (850) GeV from the multijet+$E_T$ analysis. For Dirac DM the limit increases to $m_Y \gtrsim 950$ (1300) GeV owing to the contribution from $Y_2 \rightarrow X_D X_P$. Finally, for vector DM we have $m_Y \gtrsim 1100$ (1550) GeV. For the monojet analysis, the inclusive SR with the $E_T$ cut of 500, 600, and 700 GeV (denoted as IM5, IM6, and IM7 in [4]) gives the limit for the low ($100 \sim 300$ GeV), middle ($300 \sim 450$ GeV), and high ($\gtrsim 450$ GeV) mass region, respectively. For the multijet+$E_T$ analysis, the 2-jet loose (2jl) SR gives the limit for the mass range of $100 \sim 300$ GeV, while the 2-jet medium (2jm) SR does for $\gtrsim 300$ GeV. See [43] for the detailed selection criteria.

As the production rate scales as $1/\Lambda^2$, the upper limit of $\Lambda$ can be estimated from the plots. For instance, for vector DM with $m_Y = 100$ GeV, $\Lambda$ should be larger than around 10 TeV for $g_{\text{SM}} = g_X = 1$. It should be noted that, due to the $K$ factors of $1.7 \sim 1.2$ for $m_Y = 100 \sim 2000$ GeV (see Fig. 3), these limits are slightly stronger than what would be obtained with LO production rates.

The 95% CL exclusion in the $m_X$ vs. $m_Y$ plane is shown in Fig. 6. Due to the different threshold behaviours, as seen in Eqs. (10) and (12), the excluded region near $m_Y = 2m_X$ strongly depends on the type of DM.

We note that we compared the CHECKMATE results with those obtained by the equivalent analysis implementations in MADANALYSIS 5 [47, 48] (recast codes [49, 50]) and RIVET 2.5 [51] for a couple of representative mass

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**Fig. 5.** Ratio of signal events over the number of events excluded at 95% CL as a function of the mediator mass, for $g_X = 1$ or 2 with $\Lambda = 3$ TeV, $g_{\text{SM}} = 1$ and $m_X = 10$ GeV, where the ATLAS 13 TeV (3.2 fb$^{-1}$) monojet [5] and multijet+$E_T$ [43] analyses are considered. From left to right: scalar, Dirac and vector DM.

**Fig. 6.** 95% CL exclusion from the ATLAS 13 TeV (3.2 fb$^{-1}$) monojet [5] and multijet+$E_T$ [43] analyses in the plane of the DM vs. mediator masses, for $g_X = 1$ or 2 with $\Lambda = 3$ TeV and $g_{\text{SM}} = 1$. From left to right: scalar, Dirac and vector DM.

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6 While both analyses have very similar sensitivity, i.e. their expected limits are basically the same, the monojet results have over- and under-fluctuations in some SRs. Therefore the expected and observed limits slightly differ from each other for the monojet analysis. Overall, the multijet+$E_T$ analysis tends to give the stronger limit.
choices and found agreement at the level of 20% within all three tools.

The monophoton (as well as mono-Z/W) signature could also be interesting to explore the spin-2 model. However, as seen in Sec. 4.2, the production rate for a pair of DM with a photon is strongly suppressed. We checked that there is no constraint for the above benchmark points from the CMS 13 TeV monophoton analysis ($12.9 \text{ fb}^{-1}$).

An interesting alternative to the universal coupling $g_{SM}$ is a leptophobic scenario with

$$g_T^L \ll g_{SM} = g_H^T = g_q^T = g_W^T = g_B^T. \quad (13)$$

In this case, the $E_T$ signatures come exclusively from decays into DM, because $Y_2$ decays into neutrinos are switched off. Moreover, constraints from dilepton resonance searches, which as we will see in the next subsection are quite severe, are evaded. The results for the leptophobic scenario are presented in Figs. 7 and 8 in analogy to Figs. 5 and 6. As expected, the $m_Y < 2 m_X$ region is no longer constrained. Also, for $g_X = 1$, the exclusion becomes considerably weaker for all the DM types; in particular there is no more constraint for scalar DM. For $g_X = 2$, except scalar DM, the mediator decays into DM dominates the neutrino decay mode even for the universal coupling scenario (see Fig. 2), and hence the $m_Y$ limits are very similar.

### 4.2 Resonance searches

Direct resonance searches can also be used to explore s-channel mediator DM models, see e.g. $[55, 66]$ for the spin-1 and spin-0 mediator models, respectively. Results from Run-II data are already available for a large variety of final states (dijet, dilepton, diphoton, WW, ZZ, bb, tt, hh) from ATLAS $[52, 53, 57, 59, 60]$ and CMS $[9, 56, 58, 67, 69]$, and give powerful constraints for mediator masses of a few hundred GeV up to several TeV. Lower masses are partly covered by Run-I results.$^7$

Table 2 lists the current resonance search results which we use to constrain our spin-2 simplified model. The RS massive graviton is considered in the analyses for pairs of electroweak gauge or Higgs bosons $[54, 57, 58, 60, 63, 64]$. We thank the referee for pointing us to the ATLAS analysis $[61]$, which looked for narrow scalar resonances in the diphoton invariant mass spectrum down to 65 GeV.
as one of the new physics hypotheses. For the fermionic and jet final states in $(jj)$, $(j+j)$, and model-independent Gaussian-shaped resonance have been studied. Except the dijet and di-jet analyses at 13 TeV and the low-mass diphoton analysis at 8 TeV from ATLAS, the limits are provided directly on the cross section in the given channel, and hence we obtain the model constraints by simply using the $Y_2$ production cross section and the branching ratio discussed in Section 3. For the analyses with different hypotheses from the spin-2 resonance, we assume that the acceptance and efficiency are similar. When limits are given on the fiducial cross section, $\sigma \times B \times A$, we generate LO events normalised by the NLO cross section and apply the fiducial cuts at the parton level by using MadAnalysis5 [62].

We recall that, for a given mediator mass, the $Y_2$ production cross section depends solely on $g_{SM}/A$, while the branching ratio depends also on the parameters related to DM, i.e. $g_X$ and $m_X$, as well as on the type of DM. In the decoupling limit of the dark sector, the constraints on $A/g_{SM}$ are the most stringent. When decays to DM are relevant, the branching ratios to SM particles become smaller and hence the constraints are weakened.

Figure 9 shows the constraints on $A/g_{SM}$ from observed 95% CL upper limits of resonance searches at the 13 TeV (solid) and 8 TeV (dashed) LHC as a function of the spin-2 mediator mass. We assume a negligible branching ratio to DM, except for a dotted line, where the vector DM coupling $g_X/g_{SM} = 10$ with $m_X = 10$ GeV is taken into account as a reference. Regions below each line are excluded. Information on the mediator width-to-mass ratio is given by the grey dotted lines.
and $WW/ZZ$ resonance searches lead to a constraint of a few tens of TeV on $\Lambda/g_{SM}$ for around 1 TeV mediator mass. We note again that the limits are obtained based on the NLO production rates which are larger than the LO ones, especially for $pp \rightarrow (Y_2 \rightarrow jj)\gamma$; see Fig. 3. We also note that, as indicated by grey dotted lines in Fig. 3, the mediator width can be very large at high mass and low $\Lambda/g_{SM}$; as the experimental analyses often assume a narrow width, this region has to be regarded with caution.

The weakening of the constraints when $Y_2$ decays into DM is shown in Fig. 9, depicted by a dotted line, where we assume vector DM and take $g_X = 10$ and $m_X = 10$ GeV. For instance, at $m_Y = 1$ TeV, the dilepton (electron and muon) branching ratio becomes 0.8%, i.e. the dilepton production rate becomes smaller by a factor of five, reducing the limit on $\Lambda/g_{SM}$ by $1/\sqrt{5}$. As seen in Fig. 3, the above assumption gives the largest DM branching ratio within the scenarios we consider.

Therefore, the diphoton resonance searches, and for $m_Y > 200$ GeV also the dilepton resonance searches, provide stronger constraints on the universal coupling scenario than the searches with missing energy.

To avoid such severe constraints from resonance searches, it is interesting to consider scenarios beyond the universal coupling case. The dilepton constraints could be avoided, for example, in the leptophobic scenario, $g_T = 0$, as already discussed in the previous subsection. To avoid the diphoton constraints is somewhat more complicated. One possibility would be the gravity-mediated DM model, where the KK graviton mainly couples to massive particles —DM, Higgs, massive gauge bosons and top quarks— while the couplings to photons, gluons and light quarks are highly suppressed. In such scenarios, the branching ratios and the production cross sections of the spin-2 resonance strongly depend on the setup and can be very different from those in the universal coupling case. In fact, associated production of the mediator with a $W$ or $Z$ boson, or mediator production in vector boson fusion may be more relevant than $s$-channel production in $q\bar{q}$ or $gg$ fusion. While such setups can in principle be studied easily in the simplified model framework by appropriately choosing the free parameters $g_T^q$ and $g_T^T$ in Eq. (3), such an analysis is beyond the scope of this paper. A final caveat is that non-universal couplings to gluons and quarks, $g_T^q \neq g_T^T$, give rise to a unitarity violating behaviour at higher order in QCD. We therefore only consider phenomenological scenarios with $g_T^q = g_T^T$.

5 Summary

We considered a simplified DM model where the DM candidate couples to the SM particles via an $s$-channel spin-2 mediator, $Y_2$, and studied the constraints from the current LHC data. In particular, we compared the constraints from searches with and without missing energy.

For universal couplings of the mediator to SM particles, we found that diphoton resonance searches provide the strongest constraints, $\Lambda/g_{SM} \gtrsim 100$ TeV for $Y_2$ masses up to $\sim 1$ TeV. For $\Lambda/g_{SM} = 10$ (3) TeV, the exclusion extends up to 4 (beyond 5) TeV in $m_Y$. The dilepton channel provides a similarly strong constraint for mediator masses above 200 GeV. Monojet and multijet+$E_T$ searches are competitive only if the mediator decays into photons and leptons are heavily suppressed; in this case they could provide complementary constraints to the other
resonance searches in particular in the low-mass region below 0.5–1 TeV, depending on $g_X/g_{SM}$.

For $m_Y < 2m_X$, $E_T$ signatures arise solely from $Y_2$ decays into neutrinos, leading to $m_Y \gtrsim 700$ GeV for $g_X/A = g_{SM}/A = (3 \text{ TeV})^{-1}$, based on $3.2 \text{ fb}^{-1}$ of data at $\sqrt{s} = 13 \text{ TeV}$. For $m_Y > 2m_X$, the limit crucially depends on $g_X$ and the type of dark matter. The dependence on the DM mass is less pronounced unless one approaches the threshold region. For $m_X = 10$ GeV and $g_X/A = g_{SM}/A = (3 \text{ TeV})^{-1}$, we found $m_Y \gtrsim 750, 950,$ and $1100$ GeV for scalar, Dirac, and vector DM, respectively. This increases to 850, 1300, and 1550 GeV when doubling $g_X$. We note that the obtained limits are based on the NLO-QCD predictions, which give a larger production rate than at the LO. The $K$ factor depends on the mediator mass and the production channel.

The complementarity among the different searches is illustrated in Fig. 10, where we have rescaled the reach of the jets + $E_T$ searches from 3.2 to 15 fb$^{-1}$ in order to make a fair comparison. We see that, for the same amount of data, in case of $g_X \simeq g_{SM}$ the missing energy searches are roughly competitive with the dijet and heavy diboson $(WW, ZZ)$ searches, pushing $A/g_{SM}$ beyond 20 TeV. (As mentioned, when the dilepton and diphoton constraints hold, they give even stronger limits.)

For $g_X/g_{SM} = 10$ (or $g_X/g_{SM} = 10$), also the resonance constraints strongly depend on the type of DM. Therefore, in the right plot in Fig. 10 only the vector DM case is shown. We see that the jets+$E_T$ searches give stronger constraints than the dijet and heavy diboson searches up to mediator masses of about 1.2 TeV. The dilepton and diphoton constraints are weakened by about a factor of two but still give the strongest constraints.

We hope our work will be useful to find reasonable benchmark scenarios for spin-2 mediated DM searches at the LHC as well as to construct viable UV-completed models which can give predictions for those parameters. We also note that our study on resonance searches in Sec. 4.2 can be applied not only for spin-2 mediated DM models but also for usual RS-type graviton searches; see also, e.g. [71]. As a final remark we like to point out that in a full model the presence of KK excitations might alter the LHC phenomenology as compared to the simplified model scenarios discussed here. Examples are limits on gauge KK modes providing additional constraints on light gravitons, or KK excitations of the DM fields contributing to $E_T$ signatures. While this goes well beyond the simplified model picture, it is certainly an interesting topic for future studies.

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A Supplemental material for recasting

A.1 Searches with missing energy

As mentioned in the main part of the paper, in case of the monojet and the 2–6 jets + $E_T$ searches, the signal comes solely from on-shell mediator production with the $Y_2$ decaying into neutrinos and/or DM. The signal selection efficiency (more precisely acceptance times efficiency, $A \times \epsilon$) depends only on the properties of the mediator, but not on those of the invisible decay products. Figure 11 shows $A \times \epsilon$ for those SRs which, depending on $m_Y$, can be the most sensitive ones in each of the two ATLAS analyses considered in this paper. As a service to the reader and potential user of our work, the complete $A \times \epsilon$ tables for all SRs are available in numerical form at [72].

A.2 Resonance searches

In Fig. 12 we show observed 95% CL upper limits on resonant production cross section times branching ratio (times
Fig. 12. Observed 95% CL upper limits on resonant production cross section times branching ratio (times acceptance) as a function of the resonance mass from each experimental paper; see Table 2 for more detailed information. Dashed lines denote limits including cut acceptance. For reference, NLO production cross sections of the spin-2 mediator are shown by dotted lines for different values of $g_{SM}/\Lambda$.

As indicated in Table 2, the dijet (+ ISR jet/photon) and $t\bar{t}$ analyses at 13 TeV as well as the ATLAS 8 TeV diphoton analysis provide tables with the numbers corresponding to the lines in the exclusion plots, which is very convenient for our purpose. The other analyses do not provide explicit values, and hence we have to extract these data from the exclusion plots ‘by hand’, e.g. using WebPlotDigitizer [73], a public software. To avoid that other people have to redo this exercise, our digitised data files are available at [72] and on the new PHENODATA database [74]. We encourage the experimental collaborations to provide digitised data together with their plots, in order to make it easier to use their results.

Finally, we notice a caveat regarding the re-interpretation of the low-mass resonance search in dijet plus ISR final states [53]. We found that final state radiation (FSR) may be also important and give rise to a non-trivial structure in the dijet invariant mass spectrum. Technically, simulated event shapes can differ by including FSR or not in the matrix elements, which may affect the parameter fitting procedure for a bump search.
1. CMS New Physics Working Group Collaboration, D. Alves, Simplified Models for LHC New Physics Searches, J. Phys. G39 (2012) 105005, arXiv:1105.2838.

2. D. Abercrombie et al., Dark Matter Benchmark Models for Early LHC Run-2 Searches: Report of the ATLAS/CMS Dark Matter Forum, arXiv:1507.00966.

3. CMS Collaboration, V. Khachatryan et al., Search for dark matter in proton-proton collisions at √s = 13 TeV with the ATLAS detector, JHEP 06 (2016) 059, arXiv:1604.01306.

4. CMS Collaboration, M. Aaboud et al., Search for new phenomena in events with a photon and missing transverse momentum in pp collisions at √s = 13 TeV with the ATLAS detector, JHEP 06 (2016) 059, arXiv:1604.01306.

5. ATLAS Collaboration, M. Aaboud et al., Search for new phenomena in final states with an energetic jet and large missing transverse momentum in pp collisions at √s = 13 TeV using the ATLAS detector, Phys. Rev. D94 (2016), no. 3 032005, arXiv:1604.07773.

6. ATLAS Collaboration, M. Aaboud et al., Search for dark matter produced in association with a hadronically decaying vector boson in pp collisions at √s = 13 TeV with the ATLAS detector, Phys. Lett. B763 (2016) 251–268, arXiv:1606.02379.

7. ATLAS Collaboration, M. Aaboud et al., Search for dark matter in association with a Higgs boson decaying to b-quarks in pp collisions at √s = 13 TeV with the ATLAS detector, Phys. Lett. B765 (2017) 11–31, arXiv:1609.04572.

8. ATLAS Collaboration, Search for Dark Matter production associated with bottom quarks with 13.3 fb⁻¹ of pp collisions at √s = 13 TeV with the ATLAS detector at the LHC, Tech. Rep. ATLAS-CONF-2016-086, CERN, Geneva, Aug. 2016.

9. CMS Collaboration, A. M. Sirunyan et al., Search for dijet resonances in proton-proton collisions at √s = 13 TeV and constraints on dark matter and other models, Submitted to: Phys. Lett. B (2016) arXiv:1611.03589.

10. CMS Collaboration, Search for dark matter in final states with an energetic jet, or a hadronically decaying W or Z boson using 12.9 fb⁻¹ of data at √s = 13 TeV, Tech. Rep. CMS-PAS-EXO-16-037, CERN, Geneva, 2016.

11. CMS Collaboration, Search for dark matter in Z+E_{T}^{miss} events using 12.9 fb⁻¹ of 2016 data, Tech. Rep. CMS-PAS-EXO-16-038, CERN, Geneva, 2016.

12. CMS Collaboration, Search for dark matter and graviton produced in association with a photon in pp collisions at √s = 13 TeV with an integrated luminosity of 12.9 fb⁻¹, Tech. Rep. CMS-PAS-EXO-16-039, CERN, Geneva, 2016.

13. H. M. Lee, M. Park, and V. Sanz, Gravity-mediated (or Composite) Dark Matter, Eur. Phys. J. C74 (2014) 2715, arXiv:1306.4107.

14. H. M. Lee, M. Park, and V. Sanz, Gravity-mediated (or Composite) Dark Matter Confronts Astrophysical Data, JHEP 05 (2014) 063, arXiv:1401.5301.

15. M. Bauer et al., Towards the next generation of simplified Dark Matter models, arXiv:1607.06680.

16. O. Mattelaer and E. Vryonidou, Dark matter production through loop-induced processes at the LHC: the s-channel mediator case, Eur. Phys. J. C75 (2015), no. 9 436, arXiv:1508.00564.

17. M. Backovic, M. Krämer, F. Maltoni, A. Martini, K. Mawatari, and M. Pellen, Higher-order QCD predictions for dark matter production at the LHC in simplified models with s-channel mediators, Eur. Phys. J. C75 (2015), no. 10 482, arXiv:1508.05327.

18. M. Neubert, J. Wang, and C. Zhang, Higher-Order QCD Predictions for Dark Matter Production in Mono-Z Searches at the LHC, JHEP 02 (2016) 082, arXiv:1509.05785.

19. J. Alwall, R. Frederix, S. Frixione, V. Hirschi, F. Maltoni, O. Mattelaer, H. S. Shao, T. Stelzer, P. Torrielli, and M. Zaro, The automated computation of tree-level and next-to-leading order differential cross sections, and their matching to parton shower simulations, JHEP 07 (2014) 079, arXiv:1405.0301.

20. G. Belanger, P. Boudjema, A. Pukhov, and A. Semenov, MicrOMEGAs 2.0: A Program to calculate the relic density of dark matter in a generic model, Comput. Phys. Commun. 176 (2007) 367–382, hep-ph/0607059.

21. G. Belanger, F. Boudjema, A. Pukhov, and A. Semenov, Dark matter direct detection rate in a generic model with micrOMEGAs 2.2, Comput. Phys. Commun. 180 (2009) 747–767, arXiv:0803.2360.

22. D. Budarici, G. Belanger, J. Bernon, F. Boudjema, J. Da Silva, S. Kraml, U. Laa, and A. Pukhov, Collider limits on new physics within micrOMEGAs, arXiv:1605.08334.

23. M. Backovic, K. Kong, and M. McCaskey, MadDM v.1.0: Computation of Dark Matter Relic Abundance Using MadGraph5, Physics of the Dark Universe 5–6 (2014) 18–28, arXiv:1308.4955.

24. M. Backovic, A. Martini, O. Mattelaer, K. Kong, and M. Mohlabeng, Direct Detection of Dark Matter with MadDM v.2.0, Phys. Dark Univ. 9-10 (2015) 37–50, arXiv:1505.04190.

25. C. Garcia-Cely and J. Heeck, Indirect searches of dark matter via polynomial spectral features, JCAP 1608 (2016) 023, arXiv:1605.08049.

26. J. Ellis, R. Fok, D. S. Hwang, V. Sanz, and T. You, Distinguishing ‘Higgs’ spin hypotheses using γγ and WW decays, Eur. Phys. J. C73 (2013) 2488, arXiv:1210.5229.

27. C. Englert, D. Goucalves-Netto, K. Mawatari, and T. Plehn, Higgs Quantum Numbers in Weak Boson Fusion, JHEP 1301 (2013) 148, arXiv:1212.0843.

28. G. Das, C. Degrande, V. Hirschi, F. Maltoni, and H.-S. Shao, NLO predictions for the production of a (750 GeV) spin-two particle at the LHC, arXiv:1605.09359.

29. L. Randall and R. Sundrum, A Large mass hierarchy from a small extra dimension, Phys. Rev. Lett. 83 (1999) 3370–3373, hep-ph/9905221.

30. K. Agashe, H. Davoudiasl, G. Perez, and A. Soni, Warped Gravitations in the LHC and Beyond, Phys. Rev. D76 (2007) 036006, hep-ph/0701180.

31. A. L. Fitzpatrick, J. Kaplan, L. Randall, and L.-T. Wang, Searching for the Kaluza-Klein Graviton in Bulk RS Models, JHEP 09 (2007) 013, hep-ph/0701150.

32. A. Alloul, N. D. Christensen, C. De Grande, C. Duhur, and B. Fuls, FeynRules 2.0 - A complete toolbox for tree-level phenomenology, Comput. Phys. Commun. 185 (2014) 2250–2300, arXiv:1310.1921.
33. C. Degrande, Automatic evaluation of UV and R2 terms for beyond the Standard Model Lagrangians: a proof-of-principle, Comput. Phys. Commun. 197 (2015) 239–262. arXiv:1406.3030.
34. K. Hagiwara, J. Kanzaki, Q. Li, and K. Kawatari, HEP and MadGraph/MadEvent with spin-2 particles, Eur. Phys. J. C56 (2008) 435–447, arXiv:0805.2554.
35. P. de Aquino, K. Hagiwara, Q. Li, and F. Maltoni, Simulating graviton production at hadron colliders, JHEP 06 (2011) 132, arXiv:1101.5499.
36. P. Artoisenet, P. de Aquino, F. Demartin, R. Frederix, S. Friizione, et al., A framework for Higgs characterisation, JHEP 1311 (2013) 043, arXiv:1306.6464.
37. http://feynrules.irmp.ucl.ac.be/wiki/DMsimp
38. J. Alwall, C. Duhr, B. Fuks, O. Mattelaer, D. G. Öztürk, and C.-H. Shen, A framework for Higgs characterisation, JHEP 1311 (2013) 043, arXiv:1306.6464.
39. J. Alwall, S. de Visscher, and F. Maltoni, QCD radiation in the production of heavy colored particles at the LHC, JHEP 0902 (2009) 017, arXiv:0810.5350.
40. T. Sjostrand, S. Mrenna, and P.Z. Skands, PYTHIA 6.4 Physics and Manual, JHEP 0605 (2006) 026, hep-ph/0603175.
41. E. Conte, B. Dumont, B. Fuks, and C. Wyman, Designing and recasting LHC analyses with MadAnalysis 5, Eur. Phys. J. C74 (2014), no. 10 3103, arXiv:1405.3892.
42. B. Dumont, B. Fuks, S. Kraml, S. Bein, G. Chalons, E. Conte, S. Kulkarni, D. Sengupta, and C. Wyman, Toward a public analysis database for LHC new physics searches using MADANALYSIS 5, Eur. Phys. J. C75 (2015), no. 2 56, arXiv:1407.3278.
43. D. Sengupta, “MadAnalysis implementation of the ATLAS monojet and missing transverse momentum search documented in arXiv:1604.07773.” http://doi.org/10.7484/INSPIREHEP.DATA.GTH3.RN26. 2016-07-19.
44. B. Fuks, S. Banerjee, and B. Zaidi, “MadAnalysis5 implementation of the multijet analysis of ATLAS (arXiv:1605.03814).” http://doi.org/10.7484/INSPIREHEP.DATA.GTF5.RN03. 2017-01-25.
45. A. Buckley, J. Butterworth, L. Lonnblad, D. Grellscheid, H. Hoeth, J. Monk, H. Schulz, and F. Siegert, Rivet user manual, Comput. Phys. Commun. 184 (2013) 2803–2819, arXiv:1003.0694.
46. ATLAS Collaboration, Search for New Phenomena in Dijet Events with the ATLAS Detector at $\sqrt{s}=13$ TeV with 2015 and 2016 data, Tech. Rep. ATLAS-CONF-2016-009, CERN, Geneva, Aug, 2016.
47. ATLAS Collaboration, Search for new light resonances decaying to jet pairs and produced in association with a photon or a jet in proton-proton collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, Tech. Rep. ATLAS-CONF-2016-070, CERN, Geneva, Aug, 2016.
48. ATLAS Collaboration, Search for diboson resonance production in the $t\bar{t}qq$ final state using pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector at the LHC, Tech. Rep. ATLAS-CONF-2016-062, CERN, Geneva, Aug, 2016.
49. ATLAS Collaboration, Search for resonances in the mass distribution of jet pairs with one or two jets identified as b-jets with the ATLAS detector with 2015 and 2016 data, Tech. Rep. ATLAS-CONF-2016-060, CERN, Geneva, Aug, 2016.
50. CMS Collaboration, Search for tt resonances in boosted semileptonic final states in pp collisions at $\sqrt{s} = 13$ TeV, Tech. Rep. CMS-PAS-B2G-15-002, CERN, Geneva, 2016.
51. ATLAS Collaboration, Searches for heavy ZZ and ZW resonances in the $t\ell_qq$ and $\nu\bar{q}qq$ final states in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, Tech. Rep. ATLAS-CONF-2016-082, CERN, Geneva, Aug, 2016.
52. CMS Collaboration, V. Khachatryan et al., Search for high-mass diphoton resonances in proton-proton collisions at 13 TeV and combination with 8 TeV search, Submitted to: Phys. Lett. B (2016) arXiv:1609.02507.
53. ATLAS Collaboration, Search for new high-mass resonances in the dilepton final state using proton-proton collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, Tech. Rep. ATLAS-CONF-2016-045, CERN, Geneva, Aug, 2016.
54. ATLAS Collaboration, Search for pair production of Higgs bosons in the bb bbb final state using proton–proton collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, Tech. Rep. ATLAS-CONF-2016-049, CERN, Geneva, Aug, 2016.
55. ATLAS Collaboration, G. Aad et al., Search for Scalar Diphoton Resonances in the Mass Range 65 – 600 GeV with the ATLAS Detector in pp Collision Data at $\sqrt{s} = 8$ TeV, Phys. Rev. Lett. 113 (2014), no. 17 171801, arXiv:1407.6683.
56. https://hepdata.net/record/ins1307756.
57. CMS Collaboration, V. Khachatryan et al., Search for diphoton resonances in the mass range from 150 to 850 GeV in pp collisions at $\sqrt{s} = 8$ TeV, Phys. Lett. B750 (2015) 494–519, arXiv:1506.02301.
58. ATLAS Collaboration, G. Aad et al., Combination of searches for WW, WZ, and ZZ resonances in pp collisions at $\sqrt{s} = 8$ TeV with the ATLAS detector, Phys. Lett. B755 (2016) 285–305, arXiv:1512.05099.
65. M. Chala, F. Kahlhoefer, M. McCullough, G. Nardini, and K. Schmidt-Hoberg, *Constraining Dark Sectors with Monojets and Dijets*, JHEP 07 (2015) 089, [arXiv:1503.05916](http://arxiv.org/abs/1503.05916).

66. C. Arina et al., *A comprehensive approach to dark matter studies: exploration of simplified top-philic models*, JHEP 11 (2016) 111, [arXiv:1605.09242](http://arxiv.org/abs/1605.09242).

67. CMS Collaboration, *Search for massive resonances decaying into pairs of boosted $W$ and $Z$ bosons at $\sqrt{s} = 13$ TeV*, Tech. Rep. CMS-PAS-EXO-15-002, CERN, Geneva, 2015.

68. CMS Collaboration, *Search for a high-mass resonance decaying into a dilepton final state in 13 fb$^{-1}$ of pp collisions at $\sqrt{s} = 13$ TeV*, Tech. Rep. CMS-PAS-EXO-16-031, CERN, Geneva, 2016.

69. CMS Collaboration, V. Khachatryan et al., *Search for heavy resonances decaying to tau lepton pairs in proton-proton collisions at $\sqrt{s} = 13$ TeV*, JHEP 02 (2017) 048, [arXiv:1611.06594](http://arxiv.org/abs/1611.06594).

70. E. Conte, B. Fuks, and G. Serret, *MadAnalysis 5, A User-Friendly Framework for Collider Phenomenology*, Comput. Phys. Commun. 184 (2013) 222–256, [arXiv:1206.1599](http://arxiv.org/abs/1206.1599).

71. E. Alvarez, L. Da Rold, J. Mazzitelli, and A. Szynkman, *Graviton resonance phenomenology and a pNGB Higgs at the LHC*, [arXiv:1610.08451](http://arxiv.org/abs/1610.08451).

72. [http://lpsc.in2p3.fr/projects-th/recasting/spin2mediator/](http://lpsc.in2p3.fr/projects-th/recasting/spin2mediator/)

73. [http://arohatgi.info/WebPlotDigitizer/](http://arohatgi.info/WebPlotDigitizer/)

74. [https://hepmdb.soton.ac.uk/phenodata/](https://hepmdb.soton.ac.uk/phenodata/)