

Version 0.1 DRAFT

# ATLAS+CMS DARK MATTER FORUM RECOMMENDATIONS

Author/contributor list to be added as document is finalized.

April 13, 2015



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## *Introduction*

This is a citation test Harris and Kousouris [2011].



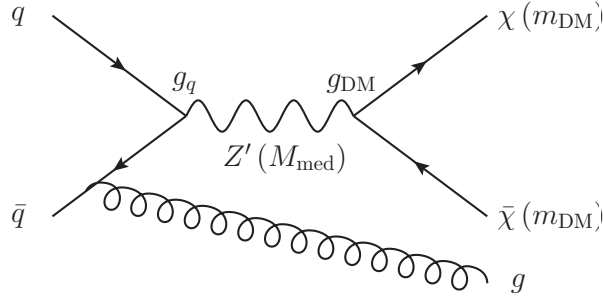


Figure 2.1: The diagram shows the pair production of dark matter particles in association with a parton from the initial state via an s-channel vector or axial-vector mediator. The process is specified by  $(M_{\text{med}}, m_{\text{DM}}, g_{\text{DM}}, g_q)$ , the mediator and dark matter masses, and the mediator couplings to dark matter and quarks respectively.

## 2

### *List of simplified models: choices and implementation*

General topics:

- choice of Dark Matter type: Dirac (unless specified otherwise) and what we might be missing
- MFV and what we might be missing

#### 2.1 *Generic models for mono-jet signatures*

*Vector and axial vector mediator, s-channel exchange*

- Matrix Element implementations (with references)
  - Production mechanism
  - Lagrangian We consider the case of a dark matter particle that is a Dirac fermion and where the production proceeds via the exchange of a spin-1 s-channel mediator. We consider the following interactions between the DM and SM fields including a vector mediator with:
    - (a) vector couplings to DM and SM.
    - (b) axial-vector couplings to DM and SM.

$$\mathcal{L}_{\text{vector}} = \sum_q g_q Z'_\mu \bar{q} \gamma^\mu q + g_{\text{DM}} Z'_\mu \bar{\chi} \gamma^\mu \chi \quad (2.1)$$

$$\mathcal{L}_{\text{axial}} = \sum_q g_q Z'_\mu \bar{q} \gamma^\mu \gamma^5 q + g_{\text{DM}} Z'_\mu \bar{\chi} \gamma^\mu \gamma^5 \chi \quad (2.2)$$

where the coupling extends over all the quarks and universal couplings are assumed for all the quarks. It is also possible to consider another model in which mixed vector and axial-vector couplings are considered, for instance the couplings to the quarks are vector whereas those to DM are axial-vector. As a starting point, we consider only the models with the vector couplings only and axial vector couplings only. Studies have been performed to see if the case of a mixed coupling can be simply extracted from the other models by some reweighting procedure to take account of the difference in cross section. This would assume that the difference between the pure and mixed couplings case does not affect the kinematics of the event.

- Definition of minimal width We assume that no additional visible or invisible decays contribute to the width of the mediator, this is referred to as the minimal width and it is defined as follows for the vector and axial-vector models.

$$\Gamma_{\text{min}} = \Gamma_{\bar{\chi}\chi} + \sum_q N_c \Gamma_{\bar{q}q} \quad (2.3)$$

where the individual contributions to this from the partial width are from,

$$\Gamma_{\bar{\chi}\chi}^{\text{vector}} = \frac{g_{\text{DM}}^2 M_{\text{med}}}{12\pi} \left( 1 + \frac{2m_{\text{DM}}^2}{M_{\text{med}}^2} \right) \sqrt{1 - \frac{4m_{\text{DM}}^2}{M_{\text{med}}^2}} \quad (2.4)$$

$$\Gamma_{\bar{q}q}^{\text{vector}} = \frac{3g_q^2 M_{\text{med}}}{12\pi} \left( 1 + \frac{2m_q^2}{M_{\text{med}}^2} \right) \sqrt{1 - \frac{4m_q^2}{M_{\text{med}}^2}} \quad (2.5)$$

$$\Gamma_{\bar{\chi}\chi}^{\text{axial}} = \frac{g_{\text{DM}}^2 M_{\text{med}}}{12\pi} \left( 1 - \frac{4m_{\text{DM}}^2}{M_{\text{med}}^2} \right)^{3/2} \quad (2.6)$$

$$\Gamma_{\bar{q}q}^{\text{axial}} = \frac{3g_q^2 M_{\text{med}}}{12\pi} \left( 1 - \frac{4m_q^2}{M_{\text{med}}^2} \right)^{3/2}. \quad (2.7)$$

- Couplings
- Parameter choices (for scan) Vary mediator mass and DM mass
- Generator implementation There are several matrix element implementations of the s-channel vector mediated DM production. This is available in POWHEG, MADGRAPH and also MCFM. The

implementation in POWHEG generates DM pair production with 1 parton at Next-to-Leading-Order, whilst Madgraph and MCFM are at leading order. As shown in POWHEG paper Haisch:2013ata, including NLO corrections result in an enhancement in the cross section as compared to leading-order (LO) and though this is not significant, it does lead to a substantial reduction in the dependence on the choice of the renormalisation and factorisation scale and hence the theoretical uncertainty on the signal prediction. Since NLO calculations are available for the process in POWHEG, we recommend to proceed with POWHEG as the generator of choice. In addition to this, studies conducted within the DM forum have shown that POWHEG is more efficient for the generation of events all the way out to the tails of the kinematic distributions (<https://indico.cern.ch/event/374678/session/o/material/3/1.pdf>). The input configuration in POWHEG allows you to set parameters to not generate events below a given  $kT$  cut ('bornktmin') and an additional parameter that ensures sufficient statistics at high transverse momentum ('bornsupfact'). With these flags set to appropriate variables, it is then possible to use a single POWHEG sample to generate the Monte Carlo for all signal regions, whereas with Madgraph more individual samples to be stitched together to achieve the required statistics out to the tails of the kinematic distributions. The POWHEG and Madgraph implementations were compared and the yields obtained from both were found to be compatible.

#### *Scalar and pseudoscalar mediator, s-channel exchange*

*Colored scalar mediator, t-channel exchange* An alternative set of simplified models exist where the mediator is exchanged in the  $t$ -channel, thereby coupling the quark and dark matter particle directly. Under the assumption that  $\chi$  is a Standard Model (SM) singlet, the mediating particle, labeled  $\phi$ , is necessarily charged and coloured. This model is parallel to, and partially motivated by, the squark of the MSSM, but in this case the  $\chi$  is chosen to be Dirac. Following the example of Ref. Papucci et al. [2014], the interaction Lagrangian is written as

$$\mathcal{L}_{\text{int}} = g \sum_{i=1,2,3} (\phi_L^i \bar{Q}_L^i + \phi_{uR}^i \bar{u}_R^i + \phi_{dR}^i \bar{d}_R^i) \chi \quad (2.8)$$

(Note: Papucci et al. [2014] uses only  $i = 1, 2$ , but I think it's fine to extend this to 3 here.) where  $Q_L^i$ ,  $u_R^i$  and  $d_R^i$  are the SM quarks and  $\phi_L^i$ ,  $\phi_{uR}^i$  and  $\phi_{dR}^i$  are the corresponding mediators, which (unlike the  $s$ -channel mediators) must be heavier than  $\chi$ . These mediators have SM gauge representations under  $(SU(3), SU(2))_Y$  of  $(3, 2)_{-1/6}$ ,  $(3, 1)_{2/3}$

and  $(3, 1)_{-1/3}$  respectively. Variations of the model previously studied include coupling to the left-handed quarks only Chang et al. [2014], Busoni et al. [2014a], to the  $\phi_{uR}^i$  or  $\phi_{dR}^i$  Papucci et al. [2014], Abdallah et al. [2014], or some combination Bai and Berger [2013], An et al. [2014].

Minimal Flavour Violation (MFV) requires that the mediator masses for each flavour be equal; the same logic also applies to the couplings  $g$ . The available parameters are then

$$\{m_\chi, M_\phi, g\}. \quad (2.9)$$

In practice, the third mediator mass and coupling could be separated from the other two, if higher order corrections to the MFV prediction arise due to the large top Yukawa coupling – a common variation is then to define this split between the first two generations and the third, so the parameters are extended to

$$\{m_\chi, M_{\phi_{1,2}}, M_{\phi_3}, g_{1,2}, g_3\}. \quad (2.10)$$

The width of each mediator is expressed, using the example of decay to an up quark, as

$$\begin{aligned} \Gamma(\phi_i \rightarrow \bar{u}_i \chi) &= \frac{g_i^2}{16\pi M_{\phi_i}^3} (M_{\phi_i}^2 - m_{u_i}^2 - m_\chi^2) \\ &\times \sqrt{M_{\phi_i}^4 + m_{u_i}^4 + m_\chi^4 - 2M_{\phi_i}^2 m_{u_i}^2 - 2M_{\phi_i}^2 m_\chi^2 - 2m_{u_i}^2 m_\chi^2}, \end{aligned} \quad (2.11)$$

this reduces to

$$\frac{g_i^2 M_{\phi_i}}{16\pi} \left(1 - \frac{m_\chi^2}{M_{\phi_i}^2}\right)^2 \quad (2.12)$$

in the limit  $M_{\phi_i}, m_\chi \gg m_{u_i}$ .

An interesting point of difference with the  $s$ -channel simplified models is that the mediator can radiate a SM object, such as a jet or gauge boson, thus providing three separate mono-X diagrams which must be considered together in calculations. This model can also give a signal in the di-jet + MET channel when, for example, the  $\chi$  is exchanged in the  $t$ -channel and the resulting  $\phi$  pair each decay to a jet +  $\chi$ .



## 2.2 Scalar models

## 2.3 Spin-0 Mediators

One of the most straightforward Simplified Models to contemplate is connecting dark matter to the visible sector through a spin-0 mediator, either a scalar or a pseudoscalar. Such models have intriguing connections with Higgs physics, and can be viewed as generalizations of the Higgs Portal to dark matter. The most general scalar mediator models will of course have renormalizable interactions between the Standard Model Higgs and the new scalar  $\phi$  or pseudoscalar  $A$ , as well as  $\phi/A$  interactions with electroweak gauge bosons. Such interactions are model-dependent, often subject to constraints from electroweak-precision tests, and would suggest specialized searches which cannot be generalized to a broad class of models (unlike the  $\cancel{E}_T$  plus jets searches, for example). As a result, for this class of simplified models with spin-0 mediators, we suggest focusing exclusively on the couplings to fermions, and induced couplings to gluons, leaving the possibilities opened up by couplings to the electroweak sector to the discussion of Higgs Portal dark matter.

In our benchmark models, we will consider two possibilities for the CP assignment of the mediator (scalar and pseudoscalar), and two spin-assignments for the dark matter itself (scalar and fermionic). We provide here the Simplified Model for the interactions of the mediator, the relevant equations for scattering with nucleons and self-annihilation, and a discussion of some important issues in simulating events at the LHC. Throughout, we will assume Minimal Flavor Violation (MFV) for the couplings of the mediators to Standard Model fermions.

### 2.3.1 Fermionic Dark Matter

Assuming dark matter is a fermion  $\chi$  who's interactions with the Standard Model proceed only through a scalar  $\phi$  or pseudoscalar  $a$ , the most general Lagrangian at tree-level are

$$\begin{aligned}\mathcal{L}_{\text{fermion},\phi} &= \mathcal{L}_{\text{SM}} + i\bar{\chi}\not{\partial}\chi + m_\chi\bar{\chi}\chi + |\partial_\mu\phi|^2 + \frac{1}{2}m_\phi^2\phi^2 + g_\chi\phi\bar{\chi}\chi + \sum_f \frac{g_v y_f}{\sqrt{2}}\phi\bar{f}f \\ \mathcal{L}_{\text{fermion},a} &= \mathcal{L}_{\text{SM}} + i\bar{\chi}\not{\partial}\chi + m_\chi\bar{\chi}\chi + |\partial_\mu a|^2 + \frac{1}{2}m_a^2a^2 + ig_\chi a\bar{\chi}\gamma^5\chi + \sum_f i\frac{g_v y_f}{\sqrt{2}}a\bar{f}\gamma^5 f\end{aligned}\quad (2.13)$$

Here, we have made several simplifying assumptions. First, we assume that the coupling to visible-sector fermions is MFV, and proportional to a single universal coupling  $g_v$ . Thus, the coupling of  $\phi$  or  $a$  to any flavor of fermion  $f$  is set by  $g_v \times y_f$ , where  $y_f$  is the Standard Model

yukawa  $y_f = \sqrt{2}m_f/v$ .<sup>1</sup> It is not hard to imagine scenarios that still possess the positive qualities of the MFV assumption but have non-universal  $g_v$ ; for example, couplings only to up-type quarks, or only to leptons. We do not single out any of these options here in our benchmark models, but remind the reader that it is desirable to have experimental constraints sensitive to couplings to different flavors of fermions. Similarly, since there is no “MFV” motivation for the structure of dark matter-mediator couplings in the dark sector, and it is of course not known whether the dark matter mass  $m_\chi$  is set by only by the Higgs vev  $v$  (indeed this would seem to be somewhat unlikely, given the direct detection constraints on dark matter as containing a pure  $SU(2)_L$  doublet) we parametrize the dark matter-mediator coupling by  $g_\chi$ , rather than by some number times a yukawa coupling proportional to  $m_\chi$ .

<sup>1</sup> This normalization for the yukawas assumes  $v = 246$  GeV and  $y_t \sim 1$ .

Finally, the most general Lagrangians including new scalars or pseudoscalars should have a potential  $V$  containing possible interactions with the Higgs  $h$ . As stated in the introduction, we choose to take a more minimal set of possible interactions, and leave discussions of the Higgs interactions to the sections on the Higgs Portal.

Given the Lagrangians in Eqs. (2.13) and 2.14, the low-energy Lagrangian will develop loop-level couplings between the mediator  $\phi/a$  and gluons and photons. This proceeds through loops exactly analogous to the loop-level couplings between the Higgs and gluons and photons. Due to our MFV assumption, as with Higgs physics, the top quark loop will dominate, with the bottom quark playing a minor role. Ignoring all couplings except those to the top, the induced couplings to *on-shell* external gluons and photons are

$$\begin{aligned}\mathcal{L}_{\text{loop},\phi} &= \frac{\alpha_S}{8\pi} \frac{g_v y_t}{v} f_\phi \left( \frac{4m_t^2}{m_\phi^2} \right) \phi G^{a,\mu\nu} G_{\mu\nu}^a + \frac{\alpha}{8\pi} (N_c Q_t^2) \frac{g_v y_t}{v} f_\phi \left( \frac{4m_t^2}{m_\phi^2} \right) \phi F^{\mu\nu} F_{\mu\nu} \\ \mathcal{L}_{\text{loop},a} &= \frac{\alpha_S}{4\pi} \frac{g_v y_t}{v} f_a \left( \frac{4m_t^2}{m_\phi^2} \right) a G^{a,\mu\nu} \tilde{G}_{\mu\nu}^a + \frac{\alpha}{4\pi} (N_c Q_t^2) \frac{g_v y_t}{v} f_a \left( \frac{4m_t^2}{m_\phi^2} \right) a F^{\mu\nu} \tilde{F}_{\mu\nu}\end{aligned}$$

where  $\alpha_S$  and  $\alpha$  are the QCD and QED fine-structure constants,  $N_c = 3$  is the number of quark colors,  $Q_t = 2/3$  is the top-quark charge, and the loop integrals are

$$\begin{aligned}f_\phi(\tau) &= \begin{cases} \tau \left( 1 + (1 - \tau) \left[ \arcsin \frac{1}{\sqrt{\tau}} \right]^2 \right), & \tau < 1, \\ \tau \left( 1 + (1 - \tau) \left( -\frac{1}{4} \right) \left[ \ln \left( \frac{1 + \sqrt{1 - \tau}}{1 - \sqrt{1 - \tau}} \right) - i\pi \right]^2 \right), & \tau > 1, \end{cases} \quad (2.17) \\ f_a(\tau) &= \begin{cases} \tau \left[ \arcsin \frac{1}{\sqrt{\tau}} \right]^2, & \tau < 1, \\ \tau \left( -\frac{1}{4} \right) \left[ \ln \left( \frac{1 + \sqrt{1 - \tau}}{1 - \sqrt{1 - \tau}} \right) - i\pi \right]^2, & \tau > 1. \end{cases} \quad (2.18)\end{aligned}$$

*It is important to remember the values of the loop-induced couplings shown here are correct only in the limit of on-shell external gauge bosons, and*

where the internal momenta in the loops are small compared to the top mass. They should not be used in Monte Carlo even-generation with gluon jets that have large  $p_T$  when compared to  $m_t$ , or when the mediators themselves have high  $p_T$ . The tree-level couplings of scalar and pseudoscalar mediators to quarks can be used in event-generation through programs like MadGraph5 as with any other new physics model, and model files are available for these purposes. Correct event generation of scalars/pseudoscalar mediators being produced primarily through the gluon couplings must use more specialized event generation routines, capable of resolving the loop. Such codes include MCFM and Sherpa, but are not as yet ready for out-of-the-box use in the same manner as MadGraph.

These Simplified Models have four free parameters: the universal coupling  $g_v$ , the dark coupling  $g_\chi$ , the dark matter mass  $m_\chi$ , and the mediator mass  $m_\phi$  or  $m_a$ . From this, all phenomenology can be calculated. However, one of the critical derived quantities, the mediator width, deserves special discussion. Under the minimal model, the widths for the mediators are given by:

$$\Gamma_\phi = \sum_f N_C \frac{y_f^2 g_v^2 m_\phi}{16\pi} \left(1 - \frac{4m_f^2}{m_\phi^2}\right)^{3/2} + \frac{g_\chi^2 m_\phi}{8\pi} \left(1 - \frac{4m_\chi^2}{m_\phi^2}\right)^{3/2} \quad (2.19)$$

$$\begin{aligned} & + \frac{\alpha_s^2 y_t^2 g_v^2 m_\phi^3}{32\pi^3 v^2} \left|f_\phi \left(\frac{4m_t^2}{m_\phi^2}\right)\right|^2 + \frac{\alpha^2 y_t^2 g_v^2 m_\phi^3}{16 \times 9\pi^3 v^2} \left|f_\phi \left(\frac{4m_t^2}{m_\phi^2}\right)\right|^2 \\ \Gamma_a = & \sum_f N_C \frac{y_f^2 g_v^2 m_a}{16\pi} \left(1 - \frac{4m_f^2}{m_a^2}\right)^{1/2} + \frac{g_\chi^2 m_a}{8\pi} \left(1 - \frac{4m_\chi^2}{m_a^2}\right)^{1/2} \quad (2.20) \\ & + \frac{\alpha_s^2 y_t^2 g_v^2 m_a^3}{8\pi^3 v^2} \left|f_a \left(\frac{4m_t^2}{m_\phi^2}\right)\right|^2 + \frac{\alpha^2 y_t^2 g_v^2 m_a^3}{4 \times 9\pi^3 v^2} \left|f_a \left(\frac{4m_t^2}{m_a^2}\right)\right|^2 \end{aligned}$$

Here, the first term in each width corresponds to the decay into Standard Model fermions (the sum runs over all kinematically available fermions,  $N_C = 3$  for quarks and  $N_C = 1$  for leptons). The second term is the decay into dark matter (assuming that this decay is kinematically allowed), and the last two terms correspond to decay into gluons and photon pairs. The factor of 2 between the decay into Standard Model fermions and into dark matter is a result of our choice of normalization of the yukawa couplings.

In colliders, if the mediator is produced on-shell, as is the primary mode of dark matter production when  $m_\chi < m_{\phi/a}/2$  and  $m_{\phi/a} \ll \sqrt{\hat{s}}$  (where  $\sqrt{\hat{s}}$  is some characteristic c.o.m. at the collider in question), the cross section of dark matter production will be proportional to the branching ratio into dark matter. The total production of the mediator will go as  $g_v^2$ , and the decay into invisible dark matter will be  $\propto g_\chi^2/\Gamma_{\phi/a}$ , with the appropriate kinematic factors. This confounds the easy factorization of limits on the four-dimensional parameter space,

since while the total cross section will have an overall dependence on the product  $g_\chi^2 \times g_v^2$ , it will also depend on  $\Gamma_{\phi/a}$ , which even in the minimal model depends on the sum of  $g_\chi^2$  and  $g_v^2$  (with kinematic factors inserted).

In addition, as this is a Simplified Model, it is possible that the mediator can decay into additional states present in a full theory that we have neglected. For example, the mediator could decay into additional new charged particles which themselves eventually decay into dark matter, but with additional visible particles that would move the event out of the selection criteria of monojets or similar missing energy searches. Thus, the widths calculated in Eqs. (2.19) and (2.20) are lower bounds on the total width.

As a result of these issues, the width of the mediator is often treated as an independent variable in Simplified Models with  $s$ -channel production of dark matter. Fortunately, for on-shell production, the effect of changing the width is only a rescaling of the total event rate assuming that  $\Gamma_{\phi/a} < m_{\phi/a}$ , which is a necessary condition for a valid weakly coupled theory. As a result, changing the width just rescales the total event rate at colliders. In the case when the dark matter is produced through an off-shell mediator, the width is not relevant. We therefore recommend that experimental and theoretical bounds on  $s$ -channel models make the assumption that the width is minimal, set by Eqs. (2.19) or (2.20), and treat it as a dependent quantity, rather than an additional free parameter.

Furthermore, as we are dealing with a 4D parameter space, it is necessary to consider how to reduce the parameter space to display results which can be compared between experiments and with theory. Clearly, the dependence of the constraints from experiments on the masses  $m_\chi$  and  $m_{\phi/a}$  cannot be suppressed, as the bounds will have non-trivial kinematic dependences on the masses. This leaves the couplings  $g_\chi$  and  $g_v$ . The most straightforward assumption to make is to set  $g_\chi = g_v$ . This has the somewhat unfortunate result of making certain types of experimental results appear more effective than others. For example, with  $g_\chi = g_v$ , we expect the mediator branching ratio to dark matter to completely dominate over the visible channels, unless the top channel is kinematically allowed. This would make resonance searches in visible channels for the mediator appear to be uncompetitive. However, some choice must be made, and setting the two couplings equal to each other is perhaps the simplest.

We now turn to the constraints on these models from non-collider experiments: thermal relic abundances, indirect detection, and direct detection. The first two results can be considered together, as they depend on the same set of annihilation cross sections.

*Thermal Cross Sections* The thermally-average annihilation of dark matter through the spin-0 mediators can be calculated from the Simplified Model Eqs. (2.13) and (2.14). The resulting cross sections for annihilation into Standard Model fermions, as a function of the dark matter temperature  $T$  are

$$\begin{aligned}\langle\sigma v\rangle(\chi\bar{\chi}\rightarrow\phi^*\rightarrow f\bar{f}) &= N_c\frac{3g_\chi^2g_v^2y_f^2(m_\chi^2-m_f^2)^{3/2}}{8\pi m_\chi^2\left[(m_\phi^2-4m_\chi^2)^2+m_\phi^2\Gamma_\phi^2\right]}T, \\ \langle\sigma v\rangle(\chi\bar{\chi}\rightarrow a^*\rightarrow f\bar{f}) &= N_c\frac{g_\chi^2g_v^2y_f^2}{4\pi\left[(m_a^2-4m_\chi^2)^2+m_a^2\Gamma_a^2\right]}\left[m_\chi^2\sqrt{1-\frac{m_f^2}{m_\chi^2}}+\frac{3m_f^2}{4m_\chi\sqrt{1-\frac{m_f^2}{m_\chi^2}}}T\right].\end{aligned}$$

Notably, the scalar mediators do not have a temperature-independent contribution their annihilation cross section, while pseudoscalars do. As  $T \propto v^2$ , where  $v$  is the dark matter velocity, there is no velocity-independent annihilation through scalars. As, in the Universe today,  $v \lesssim 10^{-3}$ , this means there are no non-trivial constraints on dark matter annihilation from indirect detection in the scalar mediator model.

The pseudoscalar model, on the other hand, does have relevant constraints from indirect detection. These can be obtained from Eq. 2.21 by setting  $T \rightarrow 0$ , and considering annihilation into the relevant Standard Model channel(s). Most constraints from indirect detection are written in terms of a single annihilation channel, and so the constraints for the full Simplified Model (with multiple annihilation channels open) require some modification of the available results. Good estimates can be obtained by considering the most massive fermion into which the dark matter can annihilate (typically the  $b$  or  $t$  quark), as this will tend to dominate the annihilation cross section. Note that, outside of resonance, the width is relatively unimportant to the indirect detection constraints.

The thermal relic calculation requires the same input cross sections as the indirect detection. Here, the cross sections are summed over all kinematically available final states, and can be written parametrically as

$$\langle\sigma v\rangle = a + bT.$$

The thermal relic abundance of dark matter is then

$$\Omega_\chi h^2 = \frac{1.04 \times 10^9 \text{ GeV}}{m_{\text{Planck}}} \frac{x_f}{\sqrt{g_*}} \frac{1}{a + bm_\chi x_f^{-1}}, \quad (2.21)$$

where  $x_f \sim 25$  is  $m_\chi$  over the freeze-out temperature, and  $g_{\text{star}}$  is the number of degrees of freedom active at the time of freeze-out. For

reasonable early Universe parameters, the correct relic abundance occurs when

$$3 \times 10^{-26} \text{ cm}^3/\text{s} = 2.57 \times 10^{-9} \text{ GeV}^{-2} = a + \frac{bm_\chi}{x_f}. \quad (2.22)$$

Keep in mind that these equations require some modification when the dark matter-mediator system is on resonance. Further, recall that we do not know dark matter is a thermal relic, or that the only annihilation process in play in the early Universe is through the mediator. Therefore, while it is appropriate to compare the sensitivity of experimental results to the thermal cross section, this is not the only range of parameters of theoretical interest.

*Direct Detection* As noted previously, the scalar mediator model has no indirect detection constraints, while the pseudoscalar does. The situation is reversed in direct detection: the pseudoscalar mediator has no velocity- or momentum-unsuppressed interactions with nucleons, while the scalar mediator induces a spin-independent scattering cross section. These constraints are very powerful compared to the present collider bounds, especially when  $m_\chi > 10 \text{ GeV}$ . As with indirect detection, for weakly coupled theories, the direct detection bounds are relatively independent of the width.

For the scalar mediator model, the interaction with direct detection nuclear targets is (to good approximation) isospin-conserving. Therefore, the bound on scattering with nucleons can be applied to either the dark matter-proton or dark matter-neutron cross section, given by

$$\begin{aligned} \sigma_{\chi-p,n} &= \frac{\mu^2}{\pi} f_{p,n}^2 \quad (2.23) \\ f_{p,n} &= \sum_{q=u,d,s} f_q^{p,n} \frac{m_{p,n}}{m_q} \left( \frac{g_\chi g_v y_q}{\sqrt{2} m_\phi^2} \right) + \frac{2}{27} f_{\text{TG}}^{p,n} \sum_{q=c,b,t} \frac{m_{p,n}}{m_q} \left( \frac{g_\chi g_v y_q}{\sqrt{2} m_\phi^2} \right) \end{aligned}$$

where  $\mu$  is the dark matter-nucleon reduced mass  $\mu = (m_\chi m_{p,n}) / (m_\chi + m_{p,n})$ , and the nuclear matrix elements  $f_q^{p,n}$  and  $f_{\text{TG}}^{p,n}$  must be extracted from lattice QCD. For our purposes, the neutron and proton matrix elements are essentially identical. Using the values from Ref. ? gives (for both protons and neutrons)

$$f_u = 0.02 \quad (2.25)$$

$$f_d = 0.026 \quad (2.26)$$

$$f_s = 0.118 \quad (2.27)$$

$$f_{\text{TG}} = 0.84 \quad (2.28)$$

## 2.4 Specific models for signatures with EW bosons

In this Section, we consider models with a photon, a W boson, a Z boson or a Higgs boson in the final state, accompanied by Dark Matter particles that either couple directly to the boson or are mediated by a new particle. The experimental signature is identified as  $V+MET$ .

These models are interesting both as extensions of models where the gluon provides the experimentally detectable signature, and as stand-alone models with final states that cannot be generated by the models in Section 2.1.

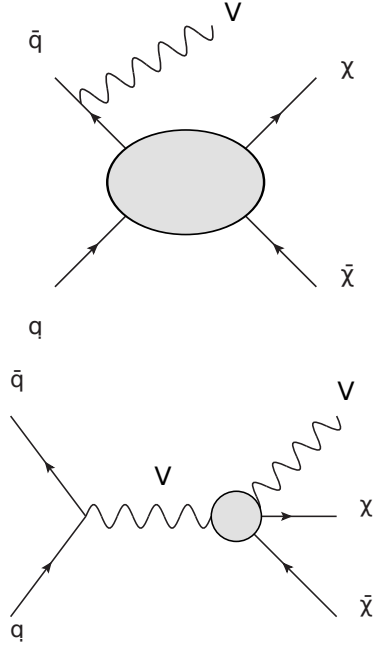


Figure 2.2: Sketch of EFT models for  $V+MET$  searches, adapted from [Nelson et al., 2014].

The models considered can be divided in four categories:

*EFT models where the boson is radiated from the initial state* As depicted in the top diagram of Figure 2.2, these models follow the nomenclature and theory for the EFT benchmarks commonly used by MET+X searches [Goodman et al., 2010]. These models have been used in past experimental searches [Kha, 2014, Aad, 2014a, Khachatryan et al., 2014, Aad, 2014a, Aad et al., 2014, Aad, 2014b], and they will not be shown here.

*EFT models where the boson is directly coupled to DM* Shown in the bottom of Figure 2.2, these models allow for an EFT vertex that directly couples the boson to Dark Matter.

*Simplified models where the boson is radiated from the initial state* These models follow those already described in Section 2.1, replacing the initial state gluon with a boson.

*V-specific simplified models* These models postulate direct couplings of new mediators to bosons, e.g. they couple the Higgs boson to a new scalar [Carpenter et al., 2014].

The following Sections describe the models within these categories, the parameters for each of the benchmark models chosen, the studies towards the choices of the parameters to be scanned, and finally point to the location of their Matrix Element implementation.

*Simplified models with ISR boson radiation* Searches in the jet+MET final state are generally more sensitive with respect to final states including bosons, due to the much larger rates of signal events featuring quark or gluon radiation with respect to radiation of bosons [Zhou et al., 2013], in combination with the low branching ratios if leptons from boson decays are required in the final state. The rates for the Higgs boson radiation is too low for these models to be considered a viable benchmark [Carpenter et al., 2014].

However, the presence of photons leptons from W and Z decays and W or Z bosons decaying hadronically allows to reject the background more effectively, making Z/gamma/W+MET searches still worth comparing with searches in the jet+MET final state.

The case for searches with W bosons in the final state is strengthened by the presence of particular choices of couplings between the WIMP and the up and down quarks which enhance W radiation [Bai and Tait, 2013], in the case of the exchange of a vector mediator in the s-channel. We consider three sample cases for the product of up and down quark couplings to the mediator  $\xi$ :

- No couplings between mediator and either up or down quarks ( $\xi = 0$ );
- Same coupling between mediator and each of the quark types ( $\xi = 1$ );
- Coupling of opposite sign between mediator and each of the quark types ( $\xi = -1$ );

The  $\xi = -1$  case produces constructive interference between the two diagrams in which a W is produced from an initial state of an up and a down quark. This in turn increases the cross-section of the process and the hardness of the spectrum of missing transverse energy or transverse mass used for the searches. The sensitivity of the W+MET search for this benchmark in the case of constructive interference surpasses that of the jet+MET search.

The above considerations lead to the choice of a vector mediator exchanged in the s-channel as the main benchmark model, in the case of a W/Z boson or a photon are radiated in the initial state.



[CD: Plots I would like for this following section: 1. mT for Wlep, 2. MET for gamma, 3. MET for Z, 4. (wishful thinking) fat jet mass for Whad]

The parameters for the definition and characterization of this model are:

- the mass of the DM particle ( $m_{DM}$ );
- the mediator mass ( $m_{Med}$ );
- the mediator width ( $\Gamma_{Med}$ );
- the couplings between the DM and the mediator ( $g_{DM}$ ), and between the mediator and the initial state quarks ( $g_{SM}$ ); together with the product of the couplings between the mediator and up/down quarks ( $\xi$ );
- the chirality of the couplings between DM and mediator, and between mediator and initial state quarks (vector-vector, axial-vector, axial-axial, vector-axial).

The DM and mediator mass ( $m_{DM}, m_{Med}$ ) are scanned independently. Different kinematic distributions are obtained with a scan of  $N$  points in the DM/mediator mass plane, as shown in Figure ?? . The spacing of the points is chosen because [CD: we need to justify this.]

The couplings between the DM and the mediator  $g_{DM}$ , and the couplings between mediator and initial state quarks  $g_{SM}$  are scanned independently following the recommendation for the jet+MET models [CD: link to monojet parameter scan]. [CD: do we have a plot of the dependence of the kinematics on the couplings?]. The products of the couplings of the mediator to up and down quarks  $\xi$  needs to be considered for the three cases separately in the case of the W+MET signature. Figure ?? shows that the constructive and destructive interference options have a similar kinematics. It is therefore recommended to generate two out of the three coupling ratio options ( $\xi=1$  and  $\xi=0$ ) and rescale the third. [CD: Kerstin's plot of mT for different values of  $\xi$  - how to make it more clear that we can rescale? Or: can we show analytically that having two out of three is sufficient?].

The mediator width is chosen to be the minimal possible width for decays in quarks and Dark Matter assuming the couplings chosen for the scan point, as in Equation ?? . The width of this mediator as a function of the mediator mass [CD: or some other plot, just to show it does not explode and that we are treating  $\xi$  correctly] is shown in Figure ?? . [CD: is there any plot of the dependence of the MET from the mediator width? This should be the same as the monojet, small difference in what happens to the ISR and MET, so we can just leave the rescaling to larger widths to theorists and trust their

**width is not too large so that they incur in PDF effects etc. Does this also hold for  $M_{Wjet}$ ?**

Similarly to the considerations made for the same model in the jet+MET case, the **chiral structure of the couplings** does not change the kinematics of the main variables used for the search **[CD: this will need a plot to be confirmed]**, so we recommend to restrict to one of the cases, e.g. the axial-vector coupling case for SM and DM particles to the mediator respectively. **[CD: using axial-vector and justifying it with suppression at DD may not be correct, see D'Eramo et al on RGE, will add paper link]**

These models are generated at leading order with MadGraph 2.2.2, and parameter cards can be found on SVN **[CD: how to put this in write-up?]**. The parton shower is done using Pythia 8, with a matching scale of... **[CD: need to figure this out.]**

*EFT models with direct DM-boson couplings* A complete list of effective operators with direct DM/boson couplings for Dirac DM, up to dimension 7, can be found in [Cotta et al., 2013].

Following the notation of [Carpenter et al., 2013], the dimension 5 benchmark models from this category have a Lagrangian that includes terms such as:

$$\frac{m_W^2}{\Lambda_5^3} \bar{\chi} \chi W^{+\mu} W_\mu^- + \frac{m_Z^2}{2\Lambda_5^3} \bar{\chi} \chi Z^\mu Z_\mu . \quad (2.29)$$

where  $m_Z$  and  $m_W$  are the masses of the Z and W boson,  $W^\mu$  and  $Z^\mu$  are the fields of the gauge bosons,  $\chi$  denote the Dark Matter fields and  $\Lambda_5$  is the effective field theory scale. This operator induces signatures with MET in conjunction with Z and W bosons at tree level, while at loop level it induces couplings to photon pairs and  $Z\gamma$  through W loops. **[CD: I believe the paper but I couldn't write this down. How much detail do we want here?]**

The dimension 7 benchmark models include couplings to the kinetic terms of the EW bosons ( $F_i^{\mu\nu}$ , with  $F_i = 1, 2, 3$  being the field strengths of the SM  $U(1)$  and  $SU(2)$  gauge groups and  $\tilde{F}_i^{\mu\nu}$  their dual tensors). The Lagrangian for the scalar coupling of DM and bosons include terms such as the following:

$$\frac{1}{\Lambda_{7,S}^3} \bar{\chi} \chi \sum_i k_i F_i^{\mu\nu} F_{\mu\nu}^i + \frac{1}{\Lambda_{7,S}^3} \bar{\chi} \chi \sum_i k_i F_i^{\mu\nu} \tilde{F}_{\mu\nu}^i \quad (2.30)$$

The Lagrangian with pseudoscalar coupling includes the following term:

$$\frac{1}{\Lambda_{7,PS}^3} \bar{\chi} \gamma^5 \chi \sum_i k_i F_i^{\mu\nu} F_{\mu\nu}^i + \frac{1}{\Lambda_{7,PS}^3} \bar{\chi} \gamma^5 \chi \sum_i k_i F_i^{\mu\nu} \tilde{F}_{\mu\nu}^i \quad (2.31)$$

The cut-off scales  $\Lambda$  for the separate terms can be related to operators with different Lorentz structure from Ref. [Cotta et al., 2013]. Given that they do not lead to substantial differences for collider searches as shown in Figure 2 of Ref. [Carpenter et al., 2013], they have been denoted as  $\Lambda_{7,S}$  for the scalar case and  $\Lambda_{7,PS}$  for the pseudoscalar case. When no distinction is made, they will be simply called  $\Lambda_7$ .

The  $k_i$  coefficients are related to the couplings of DM to pairs of gauge bosons by gauge invariance:

$$g_{WW} = \frac{2k_2}{s_w^2 \Lambda_7^3} \quad (2.32)$$

$$g_{ZZ} = \frac{1}{4s_w^2 \Lambda_7^3} \left( \frac{k_1 s_w^2}{c_w^2} + \frac{k_2 c_w^2}{s_w^2} \right) \quad (2.33)$$

$$g_{\gamma\gamma} = \frac{1}{4c_w^2} \frac{k_1 + k_2}{\Lambda_7^3} \quad (2.34)$$

$$g_{Z\gamma} = \frac{1}{2s_w c_w \Lambda_7^3} \left( \frac{k_2}{s_w^2} - \frac{k_1}{c_w^2} \right) \quad (2.35)$$

where  $s_w$  and  $c_w$  are respectively the sine and cosine of the weak mixing angle. **[CD: add here correspondence with Uli's coefficients - it may be useful.]**

The coefficients  $k_i$  determine the relative importance of each of the boson channels, and their correlations. For example, for what concerns searches with W, Z and photons:

- $k_2$  alone controls the rate of the coupling to W boson pairs;
- If  $k_1 = k_2$  contributions from both Z and  $\gamma$  exchange appear;
- If  $k_1 = c_w^2/s_w^2 k_2$  the  $\gamma$  exchange is negligible.

**[CD: discussion on the grid scan, need better internet for that]**

*Specific simplified models*

## 2.5 Specific models for signatures with heavy flavor quarks

## 2.6 SUSY-inspired simplified models



### 3

## *Validity of EFT approach*

Effective Field Theories (EFTs) are an extremely useful tool for DM searches at the LHC. Given the current lack of indications about the nature of the DM particle and its interactions, a model independent interpretation of the collider bounds appears mandatory, especially in complementarity with the reinterpretation of the exclusion limits within a choice of simplified models, which cannot exhaust the set of possible completions of an effective Lagrangian. However EFTs must be used with caution at LHC energies, where the energy scale of the interaction is at a scale where the EFT approximation can no longer be assumed to be valid. Here we summarise some methods that can be used to ensure the validity of the EFT approximation. These methods are described in detail in Refs. Busoni et al. [2014b,a,a], Aad et al. [2015], Racco et al. [2015].

*Outline of the procedure described in Refs. Aad et al. [2015]*

For a tree-level interaction between DM and the Standard Model (SM) via some mediator with mass  $M$ , the EFT approximation corresponds to expanding the propagator in powers of  $Q_{\text{tr}}^2/M^2$ , truncating at lowest order, and combining the remaining parameters into a single parameter  $M_*$  (also called  $\Lambda$ ). For an example scenario with a  $Z'$ -type mediator (leading to some combination of operators D5 to D8 in the EFT limit) this corresponds to setting

$$\frac{g_{\text{DM}}g_q}{Q_{\text{tr}}^2 - M^2} = -\frac{g_{\text{DM}}g_q}{M^2} \left( 1 + \frac{Q_{\text{tr}}^2}{M^2} + \mathcal{O}\left(\frac{Q_{\text{tr}}^4}{M^4}\right) \right) \simeq -\frac{1}{M_*^2}, \quad (3.1)$$

where  $Q_{\text{tr}}$  is the momentum carried by the mediator, and  $g_{\text{DM}}, g_q$  are the DM-mediator and quark-mediator couplings respectively. Similar expressions exist for other operators. Clearly the condition that must be satisfied for this approximation to be valid is that  $Q_{\text{tr}}^2 < M^2 = g_{\text{DM}}g_q M_*^2$ .

We can use this condition to enforce the validity of the EFT approximation by restricting the signal (after the imposition of the cuts of

the analysis) to events for which  $Q_{\text{tr}}^2 < M^2$ . This truncated signal can then be used to derive the new, truncated limit on  $M_*$  as a function of  $(m_{\text{DM}}, g_{\text{DM}} g_q)$ .

For the example D5-like operator,  $\sigma \propto M_*^{-4}$ , and so there is a simple rule for converting a rescaled cross section into a rescaled constraint on  $M_*$  if the original limit is based on a simple cut-and-count procedure. Defining  $\sigma_{\text{EFT}}^{\text{cut}}$  as the cross section truncated such that all events pass the condition  $\sqrt{g_{\text{DM}} g_q} M_*^{\text{rescaled}} > Q_{\text{tr}}$ , we have

$$M_*^{\text{rescaled}} = \left( \frac{\sigma_{\text{EFT}}^{\text{cut}}}{\sigma_{\text{EFT}}^{\text{original}}} \right)^{1/4} M_*^{\text{original}}, \quad (3.2)$$

which can be solved for  $M_*^{\text{rescaled}}$  via either iteration or a scan (note that  $M_*^{\text{rescaled}}$  appears on both the LHS and RHS of the equation). Similar relations exist for a given UV completion of each operator. The details and application of this procedure to ATLAS results can be found in Ref. Aad et al. [2015] for a range of operators. Since this method uses the physical couplings and energy scale  $Q_{\text{tr}}$ , it gives the strongest possible constraints in the EFT limit while remaining robust by ensuring the validity of the EFT approximation.

*Outline of the procedure described in Ref. Racco et al. [2015]*

In Racco et al. [2015] a procedure to extract model independent and consistent bounds within the EFT is described. This procedure can be applied to any effective Lagrangian describing the interactions between the DM and the SM, and provides limits that can be directly reinterpreted in any completion of the EFT.

The range of applicability of the EFT is defined by a mass scale  $M_{\text{cut}}$ , a parameter which marks the upper limit of the range of energy scales at which the EFT can be used reliably, independently of the particular completion of the model. Regardless of the details of the full theory, the energy scale probing the validity of the EFT is less than or equal to the centre-of-mass energy  $E_{\text{cm}}$ , the total invariant mass of the hard final states of the reaction. Therefore, the condition ensuring the validity of the EFT is, by definition of  $M_{\text{cut}}$ ,

$$E_{\text{cm}} < M_{\text{cut}}. \quad (3.3)$$

For example, in the specific case of a tree level mediation with a single mediator,  $M_{\text{cut}}$  can be interpreted as the mass of that mediator.

There are then at least three free parameters describing an EFT: the DM mass  $m_{\text{DM}}$ , the scale  $M_*$  of the interaction, and the cutoff scale  $M_{\text{cut}}$ .

We can use the same technique as above to restrict the signal to the events for which  $E_{\text{cm}} < M_{\text{cut}}$ , using only these events to derive the exclusion limits on  $M_*$  as a function of  $(m_{\text{DM}}, M_{\text{cut}})$ . We can also define an *effective coupling strength*  $M_{\text{cut}} = g_* M_*$ , where  $g_*$  is a free parameter that substitutes the parameter  $M_{\text{cut}}$ , and therefore derive exclusions on  $M_*$  as a function of  $(m_{\text{DM}}, g_*)$ . This allows us to see how much of the theoretically allowed parameter space has been actually tested and how much is still unexplored; For example, in the  $Z'$ -type model considered above,  $g_*$  is equal to  $\sqrt{g_{\text{DM}} g_q}$ . The resulting plots are shown in Racco et al. [2015] for a particular effective operator.

The advantage of this procedure is that the obtained bounds can be directly and easily recast in any completion of the EFT, by computing the parameters  $M_*$ ,  $M_{\text{cut}}$  in the full model as functions of the parameters of the complete theory. On the other hand, the resulting limits will be weaker than those obtained using  $Q_{\text{tr}}$  and a specific UV completion.





4

*Recommendations for expressing collider constraints*



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