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Mono-everything: Combined limits on dark matter production at colliders from multiple final states

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Searches for dark matter production at particle colliders are complementary to direct-detection and indirect-detection experiments and especially powerful for small masses, \( m_{\chi} < 100 \text{ GeV} \). An important collider dark matter signature is due to the production of a pair of these invisible particles with the initial-state radiation of a standard model particle. Currently, collider searches use individual and nearly orthogonal final states to search for initial-state jets, photons or massive gauge bosons. We combine these results across final states and across experiments to give the strongest current collider-based limits in the context of effective field theories and map these to limits on dark matter interactions with nuclei and to dark matter self-annihilation.

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Although the presence of dark matter in the Universe has been well-established, little is known of its particle nature or its nongravitational interactions. A vibrant experimental program is searching for a weakly interacting massive particle (WIMP), denoted as \( \chi \), and interactions with standard model particles via some as-yet-unknown mediator. If the mediator is too heavy to be resolved, the interaction can be modeled as an effective field theory with a four-point interaction.

One critical component of this program is the search for pair production of WIMPs at particle colliders, specifically \( pp \to \chi \bar{\chi} \) at the LHC via some unknown intermediate state. As the final-state WIMPs are invisible to the detectors, the events can only be seen if there is associated initial-state radiation of a standard model particle \([1–3]\), see Fig. 1, recoiling against the dark matter pair.

The LHC collaborations have reported limits on the cross section of \( pp \to \chi \bar{\chi} + X \) where \( X \) is a gluon or quark \([4,5]\), photon \([6,7]\), and other searches have been repurposed to study the cases where \( X \) is a \( W \) \([8]\) or \( Z \) boson \([9,10]\). In each case, limits are reported in terms of the mass scale \( M_* \) of the unknown interaction expressed in an effective field theory \([1–3,11–19]\). These various initial-state tags probe the same effective theory but are largely statistically independent due to their nearly orthogonal event selection requirements. As the relative rates of radiation of gluons (quarks), photons, \( W \) or \( Z \) bosons from the incoming quark (gluon) legs are determined by the standard model, the various probes may be combined to give the strongest limits without any loss of generality or additional theoretical assumptions.

Recently, an analysis of multijet final states was shown to add some sensitivity to the monojet analyses \([20]\); that sample is not statistically independent from the monojet results used here and is not included. An earlier global analysis of indirect and direct constraints with Tevatron data and monojet data from ATLAS provided an initial set of combined constraints \([21]\) using the approximations of a \( \chi^2 \) technique.

In this paper, we perform a full statistical combination of the limits from all available channels (monojet, monophoton and mono-\( Z \)) from both ATLAS and CMS at \( \sqrt{s} = 7 \text{ TeV} \), accounting for the dominant correlations and providing the most powerful current collider constraints. While the limits reported by the experimental collaborations are typically given for a few select effective operators, we calculate the efficiencies of their selections and reinterpret their searches for the complete set of operators relevant for Dirac fermion or complex scalar WIMPs.

I. MODELS

The effective theories of dark matter considered here consider the possibility that the final-state WIMPs are a Dirac fermion (operators D1-D14 in Ref. \([14]\)) or a complex scalar (operators C1-C6 in Ref. \([14]\)). These four-point effective operators assume that the unknown intermediate particles have a heavy mass scale; we use a suppression scale, \( M_* \). Cross sections at leading order for production in \( pp \) collisions at \( \sqrt{s} = 7 \text{ TeV} \) are shown in Fig. 2 for select operators with \( M_* = 1 \text{ TeV} \) for illustration. Recently, next-to-leading-order calculations have been performed for monojet and monophoton processes \([23]\) showing ratios of \( \sigma_{\text{NLO}}/\sigma_{\text{LO}} \approx 1.2–1.5 \); our monojet results partially include this effect by generating and matching multiple-parton emission.

For some operators, cross sections of dark matter production at the LHC can be transformed into cross sections for WIMP-nucleon interaction, \( \sigma(\chi - n) \) \([3]\), or WIMP...
annihilations [2]. Therefore, the effective field theories allow us to map measurements performed at the LHC to the quantities relevant for direct-detection and indirect-detection dark matter search experiments.

The effective-field-theory approach is valid as long as the unknown new mediator particles that couple the dark-matter particles to standard model quarks or gluons are too heavy to be resolved: 

\[ q < M^* \]

where \( q \) is the momentum transfer. The breakdown of the effective approach depends ultimately on the details of the new and unknown physics, specifically on the number of new mediator particles and the new couplings. Therefore, these theories cannot be treated generically and must be interpreted with some care. To guide the interpretation, we indicate the range of validity as lower bounds on the mass suppression scale \( M^\ast \) following Ref. [3]. We note that any range of validity of the effective field theory involves assumptions about the unknown physics; see Refs. [20,24] for additional unitarity arguments and more stringent validity ranges.

Assuming the simplest possible structure of new physics (mediation via exactly one new heavy mediator of mass \( M^\ast = M/\sqrt{g_1 g_2} \), \( g_1 \) and \( g_2 \) being coupling constants), bounds on the suppression scale can be placed by requiring \( M > 2m_\chi \) and that the new physics be as strongly coupled as possible for it to be still perturbative (\( \sqrt{g_1 g_2} < 4\pi \)):

\[
M^\ast > \frac{m_\chi}{2\pi} \quad (D5 \text{ to } D14 \text{ and C3 to C6}), \\
\left( \frac{M^2}{m_q^2} > \frac{m_\chi}{2\pi} \right) \quad (D1 \text{ to } D4), \\
\left( \frac{M^2}{m_q^2} > \frac{m_\chi}{2\pi} \right) \quad (C1 \text{ and C2}).
\]

Note that we are accounting for additional factors of \( m_q \) in the definitions of operators D1 to D4 and C1, C2 of Ref. [3].

**II. EXPERIMENTAL SEARCHES**

The experimental searches typically require one or more high-\( p_T \) object and missing transverse momentum; see FIG. 1. Pair production of WIMPs (\( \chi \bar{\chi} \)) in proton-proton collisions at the LHC via an unknown intermediate state, with initial-state radiation of a standard model particle.

![FIG. 1](image)

**FIG. 1.** Pair production of WIMPs (\( \chi \bar{\chi} \)) in proton-proton collisions at the LHC via an unknown intermediate state, with initial-state radiation of a standard model particle.

**FIG. 2 (color online).** Cross sections for \( pp \to \chi \bar{\chi} + X \) production where \( X \) is the initial-state radiation of a jet, photon or Z boson. Jet and photon final states include a \( p_T > 80 \text{ GeV} \) cut at the parton level. Each pane shows the cross section for a different effective operator: the top is D5, the center is D8, and the bottom is D9. See Ref. [3] for operator definitions.

Table I for a summary and comparison of the monophoton and monojet selections.

The mono-Z analysis [10] uses the ATLAS \( ZZ \to \ell \ell \nu\nu \) cross-section measurement [9], which requires:

(i) two same-flavor opposite-sign electrons or muons, each with \( p_T > 20 \text{ GeV} \), \( |\eta| < 2.5 \); 
(ii) dilepton invariant mass close to the Z-boson mass: 
\[
m_{\ell\ell} \in [m_Z - 15, m_Z + 15] \text{ GeV};
\]


(iii) no particle-level jet with $p_T^j > 25$ GeV and $|\eta^j| < 4.5$;
(iv) $(|p_T^{\ell\nu} - p_T^{\ell\ell}|)/p_T^{\ell\ell} < 0.6$;
(v) $-p_T^{\ell\nu} \times \cos(\Delta \phi(p_T^{\ell\nu}, p_T^{\ell\ell})) > 80$ GeV.

The selection efficiency of each selection for each operator is given in Table II and was estimated in the following way. References [4–7] provide signal efficiency for several select operators; this efficiency is the product of geometric and kinematic acceptance for the selection criteria and object reconstruction efficiency. The object reconstruction efficiency depends on the details of the detector performance but is largely independent of the operator. The geometric and kinematic acceptances can be reliably estimated using parton-level simulated event samples [25]. We measure the geometric and kinematic efficiency for each operator and use the quoted total efficiencies to deduce the object reconstruction efficiencies. This allows us to estimate the total efficiency for each operator.

### III. COMBINATION

The separate analyses, each of which are single-bin counting experiments, are combined into a multibin counting experiment. This allows for a coherent signal rate to be tested across channels but preserves their distinct signal-to-background ratios.

The background estimates are taken directly from the experimental publications, see a summary in Table III, and are assumed to be uncorrelated across channels, as they are typically dominated by channel-specific or detector-specific uncertainties. For example, in some cases, the background estimates are data-driven, and the dominant uncertainties are in the finite statistics of independent control samples. Inclusion of correlations up to 20% does not qualitatively impact the results of the combination.
The backgrounds, their uncertainties and the observed yield can be used to calculate a 90% C.L. upper limit on the number of signal events $N$ in the sample, see Tables III and IV, using the C.L.s method \cite{26,27}. This value is almost completely model-independent. Translating it into a limit on the cross section for the $pp \rightarrow \chi \bar{\chi} + X$ signal requires the efficiency of the signal in each selection; see Table III. These individual limits reproduce well the results reported by the experiments.

The signal regions are nearly orthogonal but not exactly. For example, the monojet analyses do not veto events with a photon, and the monophoton analyses allow the presence of one jet. From our parton-level simulated event samples, we estimated the overlaps among different channels and found that the overlap fraction is less than 1%.

The individual analyses include signal uncertainties of up to 20% on the cross section, mostly due to uncertainties in jet energy calibration and levels of initial-state radiation. These uncertainties do not affect the cross-section limits but can be simply applied to limits on $M_*$. In each case, we quote the limit using the central value.

To summarize, the assumptions made in this combination are valid.

### Table III
90% C.L. limits on $N_{\text{events}}$, efficiencies for $m_{\chi} = 10$ GeV and limits on $\sigma(pp \rightarrow \chi \bar{\chi} + X)$ using the D5 operator.

<table>
<thead>
<tr>
<th>Channel</th>
<th>Background</th>
<th>Observed</th>
<th>Limit $N$</th>
<th>Efficiency</th>
<th>Luminosity (fb$^{-1}$)</th>
<th>Limit $\sigma$ (fb)</th>
</tr>
</thead>
<tbody>
<tr>
<td>ATLAS jet + $E_T$</td>
<td>$750 \pm 60$</td>
<td>$785$</td>
<td>$139.3$</td>
<td>$1.7%$</td>
<td>$4.8$</td>
<td>$1,700$</td>
</tr>
<tr>
<td>CMS jet + $E_T$</td>
<td>$1225 \pm 101$</td>
<td>$1142$</td>
<td>$125.2$</td>
<td>$2.2%$</td>
<td>$5.0$</td>
<td>$1,140$</td>
</tr>
<tr>
<td>ATLAS $\gamma + E_T$</td>
<td>$137 \pm 20$</td>
<td>$116$</td>
<td>$27.4$</td>
<td>$18%$</td>
<td>$4.6$</td>
<td>$33$</td>
</tr>
<tr>
<td>CMS $\gamma + E_T$</td>
<td>$75.1 \pm 9.4$</td>
<td>$73$</td>
<td>$19.3$</td>
<td>$11%$</td>
<td>$5.0$</td>
<td>$35$</td>
</tr>
<tr>
<td>ATLAS $Z + E_T$</td>
<td>$86.2 \pm 7.2$</td>
<td>$87$</td>
<td>$21.7$</td>
<td>$13%$</td>
<td>$4.6$</td>
<td>$36$</td>
</tr>
</tbody>
</table>

### Table IV
90% C.L. limits on $N_{\text{events}}$, efficiencies for $m_{\chi} = 10$ GeV and limits on $\sigma(pp \rightarrow \chi \bar{\chi} + X)$ using the D9 operator.

<table>
<thead>
<tr>
<th>Channel</th>
<th>Background</th>
<th>Observed</th>
<th>Limit $N$</th>
<th>Efficiency</th>
<th>Luminosity (fb$^{-1}$)</th>
<th>Limit $\sigma$ (fb)</th>
</tr>
</thead>
<tbody>
<tr>
<td>ATLAS jet + $E_T$</td>
<td>$83 \pm 14$</td>
<td>$77$</td>
<td>$25.5$</td>
<td>$0.9%$</td>
<td>$4.8$</td>
<td>$590$</td>
</tr>
<tr>
<td>CMS jet + $E_T$</td>
<td>$1225 \pm 101$</td>
<td>$1142$</td>
<td>$125.2$</td>
<td>$4.1%$</td>
<td>$5.0$</td>
<td>$610$</td>
</tr>
</tbody>
</table>

### Table V
90% C.L. limits on $\sigma(pp \rightarrow \chi \bar{\chi} + X)$ for $m_{\chi} = 10$ GeV, theory prediction for $M_* = 1$ TeV, and limits on $M_*$ using the D5 operator. In the case of the $Z + E_T$ final state, the predictions include the $Z \rightarrow \ell \ell$ branching fraction.

<table>
<thead>
<tr>
<th>Channel</th>
<th>Limit $\sigma$ (fb)</th>
<th>Predicted (fb)</th>
<th>Limit $M_*$ (GeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>ATLAS jet + $E_T$</td>
<td>$1,700$</td>
<td>$370$</td>
<td>$685$</td>
</tr>
<tr>
<td>CMS jet + $E_T$</td>
<td>$1,140$</td>
<td>$370$</td>
<td>$750$</td>
</tr>
<tr>
<td>ATLAS $\gamma + E_T$</td>
<td>$33$</td>
<td>$3.7$</td>
<td>$580$</td>
</tr>
<tr>
<td>CMS $\gamma + E_T$</td>
<td>$35$</td>
<td>$3.7$</td>
<td>$570$</td>
</tr>
<tr>
<td>ATLAS $Z + E_T$</td>
<td>$36$</td>
<td>$0.5$</td>
<td>$340$</td>
</tr>
</tbody>
</table>

**FIG. 3 (color online).** Limits at 90% C.L. in $M_*$ (top) and in the spin-independent WIMP-nucleon cross section (bottom) for individual and combined limits using the D5 operator as a function of $m_\chi$. 
(i) the background uncertainties are monolithic and uncorrelated;
(ii) the signal selections are orthogonal.

Combining channels is then straightforward, although the intermediate step of a model-independent limit on the number of events \( N \) is no longer possible, as the limits depend on the relative distribution of signal events across channels, which is model-specific. Instead, cross-section limits are obtained directly. These limits are then converted into limits on \( M^2 \), using the relationships from Ref. [14].

The individual-channel limits, combination across experiments and the grand combination of all channels are shown in Table V for the D5 operator and one choice of \( m/C31 \).

Clearly the monojet analyses are the most powerful, and the greatest gain in combination is from combining the ATLAS and CMS monojet analyses, although the addition of the monophoton and mono-Z gives a non-negligible improvement in the combined result.

Limits on \( M^2 \) for the D5 and D8 operators are shown in Figs. 3 and 4 as well as limits on \( \sigma(\chi - n) \). Where the \( M^2 \) limits exceed the thermal relic values taken from Ref. [3], assuming that dark matter is entirely composed of thermal relics, the resulting dark matter density of the Universe would contradict WMAP measurements; therefore, WIMPs cannot couple to quarks or gluons exclusively via the given operator and account entirely for the relic density. This \( m/C31 \) region is either excluded or requires that annihilation channels to leptons must exist or participation of different operators which interfere negatively, thereby reducing the limits on \( M^2 \).
IV. APPLICATION TO OTHER MODELS

While the experimental results are usually quoted for a small selection of the effective operator models, the analyses are clearly relevant for all of them.

We reinterpret the experimental analyses in the context of each operator and perform the grand combination across all channels. Figure 5 and Table VI show the limits on $M_*$, translated to the WIMP-nucleon cross section where possible.

V. CONCLUSIONS

We have presented the first combination of collider-based searches for dark matter pair production, using
final states involving jets, photons and leptonically decaying Z bosons in the context of effective field theories. The most powerful results are from the monojet analyses, and the greatest gains come from the combination of the independent analyses from ATLAS and CMS, although the other final states make a non-negligible improvement. The results are the strongest limits to date from collider searches in the effective field theory context.

In addition, we have reinterpreted the experimental results, quoted by ATLAS and CMS only for a few effective operators, across a broad range of operators, providing a comprehensive view of the power of these searches to constrain the weak-level or weaker

FIG. 7 (color online). Combined limits on $M_\chi$ vs dark matter mass $m_\chi$ for operators D1, D2, D3 and D4. The $M_\chi$ values at which dark matter particles of a given mass would result in the required relic abundance are shown as green dashed lines [3], assuming annihilation in the early Universe proceeded exclusively via the given operator.

FIG. 8 (color online). Combined limits on $M_\chi$ vs dark matter mass $m_\chi$ for operators D5, D6 and D7. The $M_\chi$ values at which dark matter particles of a given mass would result in the required relic abundance are shown as green dashed lines [3], assuming annihilation in the early Universe proceeded exclusively via the given operator.
interactions between dark matter and standard model particles.

We have made use of the effective field theory framework to convert the ATLAS and CMS results to quantities relevant for direct-detection and indirect-detection dark matter searches. Under the assumptions made for the effective operators, LHC limits can be very competitive, in particular, for low-mass dark matter particles $m_\chi \leq 10$ GeV.

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**APPENDIX: INDIVIDUAL OPERATORS**

In Figs. 7–12, we show the combined limits for each operator, compared to the thermal relic values. Where the limits exceed the thermal relic values, assuming that dark matter is entirely composed of thermal relics, the dark matter density of the Universe would contradict measurements and hence cannot couple to quarks or gluons exclusively via the given operator. This region is either excluded, else other annihilation channels to leptons must exist, or finally different operators may interfere negatively thereby reducing the limits on $M_\chi$.

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