Complementary Constraints on Light Dark Matter from Heavy Quarkonium Decays

Nicolas Fernandez, Jason Kumar, Ilsoo Seong, and Patrick Stengel

Department of Physics and Astronomy,
University of Hawaii, Honolulu, HI 96822 USA

Abstract

We investigate constraints on the properties of light dark matter which can be obtained from analysis of invisible quarkonium decays at high intensity electron-positron colliders in the framework of a low energy effective field theory. A matrix element analysis of all contact operators pertinent for these meson decays allows for a model-independent calculation of associated dark matter-nucleon scattering and dark matter annihilation cross sections. Assuming dark matter couples universally to all quark flavors, we then obtain bounds on nucleon scattering which complement direct dark matter detection searches. In contrast to similar analyses of monojet searches at high energy colliders, B and charm factories are more suitable probes of light dark matter interactions with less massive mediators. Relevant bounds on dark matter annihilation arising from gamma ray searches of dwarf spheroidal galaxies are also presented.
I. INTRODUCTION

Despite strong observational evidence for non-baryonic dark matter (DM) which interacts gravitationally \[1\], the detection of dark matter interactions with the Standard Model (SM) remains elusive. Many extensions of the SM predict dark matter candidates which should leave signatures in the direct and indirect dark matter detection experiments, and at hadron colliders. If the particles mediating dark matter-Standard Model interactions are much heavier than the energy scales involved, then the constraints on dark matter interactions arising from these disparate detection strategies can be related to each other in a model independent fashion via a generalized effective field theory (EFT) framework, in which the details of the ultraviolet (UV) physics have been integrated out of the Lagrangian \[2\]–\[15\], and dark matter-Standard Model interactions occur through contact operators.

Weakly interacting massive particles (WIMPs) are stable dark matter candidates predicted by many models of physics beyond the SM \[1\]. Although recent hints of possible WIMP signals may be encouraging, the lack of clear and convincing evidence for the discovery of WIMP dark matter motivates consideration of dark matter candidates which deviate from the expectations of the WIMP paradigm. A well-motivated example is light dark matter (LDM), a class of dark matter candidates with masses typically \(\sim 10 \text{ MeV} - 10 \text{ GeV}\). LDM would elastically scatter at direct detection experiments, with nuclear recoil energies which are relatively small and may be below the experimental threshold, rendering them undetectable. In this case, other experimental means aside from direct detection would be required to probe DM-SM interactions. For example, complimentary bounds on LDM scattering can be inferred from collider monojet searches \[16\]–\[27\]. Nonresonant LDM production at low energy \(e^+/e^-\) colliders can also be used to set model independent limits on electron scattering and, if the LDM couples universally, nucleon scattering \[28\].

In this work, we consider the prospects for probing LDM-quark interactions through bounds on invisible decays of heavy quarkonium states at colliders \[29\]–\[33\]. Such bounds have already been considered in a variety of contexts, including \(B\) and \(D\) meson decays \[34\], and \(\Upsilon\) decays into scalar LDM \[35\]–\[36\]. However, the constraints which one can obtain on dark matter-quark interactions depend in detail on the quantum numbers of the heavy meson, as well on the choice of final state (i.e., \(\rightarrow \text{invisible}\) or \(\rightarrow \gamma + \text{invisible}\)). The angular momentum and \(C/P\) transformation properties of the initial state (as well as the presence or absence of a photon in the final state) together determine which of the possible dark matter-quark interaction structures can participate in the decay process (and can thus be bounded by constraints on invisible decays). We consider invisible decays of the heavy quarkonium states \(\Upsilon(1S)\) and \(J/\Psi\), mesons with \(J^{PC} = 1^{-+}\). As the quark constituents annihilate in an \(s\)-wave,
the dependence of the meson decay matrix element on the associated nonrelativistic bound state wavefunction is very simple and can be determined experimentally, with relatively little uncertainty. Moreover, since the mesons which we consider each have a quark and anti-quark of the same flavor, the DM-SM interactions which we introduce are not constrained by bounds on flavor-violation and, given universal quark coupling, can contribute to nucleon scattering.

At quark level, the matrix element relevant for meson decay ($\bar{q}q \rightarrow \bar{X}X$) is also relevant for monojet/photon/$W,Z$ searches at the LHC \cite{39-43}. As a result, these searches will share many features, and event rates will have the same dependence on the energy of the process. Moreover, bounds arising from both invisible meson decay rates and LHC monojet searches do not weaken as the dark matter mass decreases, in notable contrast to direct detection searches. However, if the particle mediating LDM interactions is light, then the interaction is poorly approximated by a contact operator for the purposes of LHC mono-anything searches, and model independent bounds on the interaction strength can no longer be obtained. Even given a mediating particle massive enough to warrant the use of the contact interaction approximation, the masses of LDM particles would be difficult to resolve at LHC searches due to the large center-of-mass energy of the beam. The bounds obtained from meson decays thus provide a unique handle on some dark matter interaction models, which is complementary to the information provided by other search strategies.

In this paper, we use the limits on bound state decay widths to constrain the coupling of scalar, fermion or vector LDM to Standard Model quarks through all contact operators of dimension six or lower. In section II, we review the relevant effective contact interactions and calculate the resulting meson decay rates and dark matter annihilation and scattering cross sections. In section III we present constraints on all of the relevant dark matter-quark interaction structures arising from $\Upsilon(1S)$ decay and Fermi gamma-ray searches of dwarf spheroidal galaxies, and relate these constraints to those arising from direct detection experiments and LHC searches. We conclude in section IV with a discussion of our results.

II. FRAMEWORK AND CONSTRAINTS

We will consider a framework in which dark matter-quark interactions can be parametrized by a four-point contact effective operator. Such a structure can be written as an appropriate Lorentz contraction of a Standard Model quark bilinear and a dark matter bilinear. We are only interested in effective operators which can yield a non-zero matrix element when acting on a $1^-\bar{\psi}$ meson state, such as the $\Upsilon(1S)$ or $J/\psi$; for such operators, the quark bilinear must be either $\bar{q}\gamma^i q$ or $\bar{q}\sigma^{0i} q$, where $i$ is a spatial index \cite{7}. The angular momentum quantum numbers of the quark/anti-quark bound states of interest are $S = 1$, $L = 0$, $J = 1$. In
general, the final state need not have the same $C$ and $P$ transformation properties as the initial state (and thus need not have the same $S$ and $L$ quantum numbers), but must have the same total angular momentum $J$. The effective operators of dimension 6 or less which can have a non-zero matrix element with either an $Υ(1S)$ or $J/ψ$ initial state are listed in Table I [7] (the operators are labeled using the conventions of [7] and [20]). Note that for Majorana fermion dark matter, $F6$ is the only non-vanishing contact operator.

| Name | Interaction Structure | Annihilation | Scattering |
|------|-----------------------|--------------|------------|
| $F5$ | $(1/Λ^2)Xγ^μXqγ_μq$ | Yes          | SI         |
| $F6$ | $(1/Λ^2)Xγ^μγ^5Xqγ_μq$ | No           | No         |
| $F9$ | $(1/Λ^2)Xσ^{μν}Xqσ_μ^νq$ | Yes          | SD         |
| $F10$ | $(1/Λ^2)Xσμνγ^5Xqσ_μ^νq$ | Yes          | No         |
| $S3$ | $(1/Λ^2)iIm(φ^†∂_μφ)qγ^μq$ | No           | SI         |
| $V3$ | $(1/Λ^2)iIm(B^†_μB_ν)qγ^μq$ | No           | SI         |
| $V5$ | $(1/Λ)(B^†_μB_ν − B^†_νB_μ)qσμνq$ | Yes          | SD         |
| $V6$ | $(1/Λ)(B^†_μB_ν − B^†_νB_μ)qσμνγ^5q$ | Yes          | No         |
| $V7$ | $(1/Λ^2)B^†_μB_ν(1/2)γ^μq$ | No           | No         |
| $V9$ | $(1/Λ^2)ε^{μνχσ}B^†_μ(1/2)γ^νq$ | No           | No         |

TABLE I. Effective contact operators which can mediate the decay of a $J^{PC} = 1^{−−}$ quarkonium bound state. We also indicate if the operator can permit an $s$-wave dark matter initial state to annihilate to a quark/anti-quark pair; if so, then a bound can also be set by indirect observations of photons originating from dwarf spheroidal galaxies. Lastly, we indicate if the effective operator can mediate velocity-independent nucleon scattering which is either spin-independent (SI) or spin-dependent (SD).

The measured limits on the invisible decay branching fractions could possibly be contaminated by decays to an invisible final state as well as a soft photon. However, contributions to the processes $Υ(1S), J/ψ → Xγ$ would not arise from the operators we consider. A contribution could arise from the other dark matter-quark four-point operators which we do not consider, but the rate would be suppressed by an additional factor of $α$; to constrain such operators, it would be more fruitful to search directly for $γ + invisible$ decays. The contribution of soft photon processes to fake decays was also considered in [44], where the contribution was also small.

The meson decay rate matrix element can be determined by convolving the quark/antiquark annihilation ($qq → Xγ$) matrix element with the meson bound state wavefunction. Using crossing symmetry, one can relate the quark annihilation matrix element to the matrix

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1 Henceforth, we refer to spin-0, spin-1/2 and spin-1 dark matter fields with $φ$, $X$ and $B^μ$, respectively. When describing a dark matter particle of arbitrary spin, we will use $X$. 

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elements for either dark matter annihilation ($\bar{XX} \to \bar{q}q$) or dark matter-nucleon scattering ($XN \to XN$).

The signals which can be observed at various experiments depend on which quarks appear in the effective operator. The $b$- and $c$-quark couplings will be relevant for the decay of $\Upsilon$ or $J/\psi$ states, respectively, whereas $u$-, $d$- and $s$-quark couplings are the most relevant for either dark matter annihilation, dark matter-nucleon scattering, or dark matter production at the LHC.

A. Bound State Decays

The matrix element for the decay of a bound state is given by the convolution of the nonrelativistic bound state wavefunction with the annihilation matrix element for a free quark/anti-quark pair. Since we consider $s$-wave meson bound states in the nonrelativistic approximation, this convolution depends only on the value of the spatial wavefunction at the origin, $\psi(0)$. The wavefunction at the origin can then be determined from the well-measured decay branching fraction to $e^+e^-$, yielding

$$B(\Upsilon(1S) \to e^+e^-) = 16\pi\alpha^2 Q_b^2 |\psi_\Upsilon(0)|^2 \frac{\Gamma_\Upsilon M_\Upsilon^2}{M_\Upsilon^2} = 0.0238 \pm 0.0011,$$

$$B(J/\Psi \to e^+e^-) = 16\pi\alpha^2 Q_c^2 |\psi_{J/\psi}(0)|^2 \frac{\Gamma_{J/\psi} M_{J/\psi}^2}{M_{J/\psi}^2} = 0.0594 \pm 0.0006,$$  \quad (1)

with $M_\Upsilon = 9460.30 \pm 0.26$ MeV, $\Gamma_\Upsilon = 54.02 \pm 1.25$ keV, $M_{J/\psi} = 3096.916 \pm 0.011$ MeV and $\Gamma_{J/\psi} = 92.9 \pm 2.8$ keV [45]. Note, we have ignored the contribution from $Z,h$-exchange; this contribution is smaller than the uncertainties in the measured branching fraction. Searches for $\Upsilon(1S)$ invisible decays have been performed by Belle [46] and BaBar [47] operating at the $\Upsilon(3S)$ resonance. They use the transition $\Upsilon(3S) \to \pi^+\pi^-\Upsilon(1S)$ to detect invisible $\Upsilon(1S)$ decays and reconstruct the presence of the $\Upsilon(1S)$ from the $\Upsilon(1S)$ peak in the recoil mass distribution, $M_{\text{rec}}$, by tagging $\pi^+\pi^-$ pairs with kinematics

$$M_{\text{rec}}^2 \equiv s + M_{\pi\pi}^2 - 2\sqrt{s}E_{\pi\pi}^*, \quad (2)$$

where $M_{\pi\pi}$ is the invariant mass of the dipion system, $E_{\pi\pi}^*$ is the energy of the dipion system in the center-of-mass (CM) frame of the $\Upsilon(3S)$, and $\sqrt{s} = 10.3552$ GeV is $\Upsilon(3S)$ resonance energy. Similar searches for invisible decays of $J/\Psi$ are based on the transition $\Psi(2S) \to \pi^+\pi^-J/\Psi$. The 90% CL constraints on branching fractions for invisible decays of $\Upsilon(1S)$ and $J/\Psi$, as measured by BaBar and BES [48], respectively, are

$$B(\Upsilon(1S) \to \text{invisible}) < 3.0 \times 10^{-4},$$

$$B(J/\Psi \to \text{invisible}) < 7.2 \times 10^{-4}. \quad (3)$$
There is a Standard Model contribution to invisible bound state decays, namely, the decay of a meson to \( \bar{\nu}\nu \) via a Z boson. But these partial widths have been calculated and are negligible [49]:

\[
B(\Upsilon(1S) \to \nu\bar{\nu}) = 9.85 \times 10^{-6},
B(J/\Psi \to \nu\bar{\nu}) = 2.70 \times 10^{-8}.
\]

For each contact operator, one can calculate the branching fraction for the bound state to decay to dark matter in terms of the bound state mass, the mediation scale, the dark matter mass, and the branching fraction to \( e^+e^- \) (assuming \( q = b \) for \( \Upsilon(1S) \) decay, or \( q = c \) for \( J/\psi \) decay):

\[
\begin{align*}
B_{F5}(\bar{X}X) &= \frac{B(e^+e^-)M^4}{16\pi^2\alpha^2Q^4\Lambda^4} \left(1 - \frac{4m_X^2}{M^2}\right)^{1/2} \left(1 + \frac{2m_X^2}{M^2}\right), \\
B_{F6}(\bar{X}X) &= \frac{B(e^+e^-)M^4}{16\pi^2\alpha^2Q^2\Lambda^4} \left(1 - \frac{4m_X^2}{M^2}\right)^{3/2}, \\
B_{F9}(\bar{X}X) &= \frac{B(e^+e^-)M^4}{8\pi^2\alpha^2Q^2\Lambda^4} \left(1 - \frac{4m_X^2}{M^2}\right)^{1/2} \left(1 + \frac{8m_X^2}{M^2}\right), \\
B_{F10}(\bar{X}X) &= \frac{B(e^+e^-)M^4}{8\pi^2\alpha^2Q^2\Lambda^4} \left(1 - \frac{4m_X^2}{M^2}\right)^{3/2}, \\
B_{S3}(\bar{X}X) &= \frac{B(e^+e^-)M^4}{256\pi^2\alpha^2Q^2\Lambda^4} \left(1 - \frac{4m_X^2}{M^2}\right)^{3/2}, \\
B_{V3}(\bar{X}X) &= \frac{B(e^+e^-)M^4}{128\pi\alpha^2Q^2\Lambda^4} \left(1 - \frac{4m_X^2}{M^2}\right)^{3/2} \left(1 + \frac{M^4}{8m_X^2} \left(1 - \frac{2m_X^2}{M^2}\right)^2\right), \\
B_{V5}(\bar{X}X) &= \frac{B(e^+e^-)M^2}{16\pi^2\alpha^2Q^2\Lambda^2} \left(1 - \frac{4m_X^2}{M^2}\right)^{3/2} \frac{M^2}{m_X^2} \left(1 + \frac{M^2}{4m_X^2}\right), \\
B_{V6}(\bar{X}X) &= \frac{B(e^+e^-)M^2}{16\pi^2\alpha^2Q^2\Lambda^2} \left(1 - \frac{4m_X^2}{M^2}\right)^{1/2} \frac{M^2}{m_X^2} \left(1 + \frac{2m_X^2}{M^2}\right), \\
B_{V7}(\bar{X}X) &= \frac{B(e^+e^-)M^4}{64\pi^2\alpha^2Q^2\Lambda^4} \left(1 - \frac{4m_X^2}{M^2}\right)^{3/2} \frac{M^2}{m_X^2}, \\
B_{V9}(\bar{X}X) &= \frac{B(e^+e^-)M^4}{256\pi^2\alpha^2Q^2\Lambda^4} \left(1 - \frac{4m_X^2}{M^2}\right)^{5/2} \frac{M^2}{m_X^2}.
\end{align*}
\]

Note that for operator F6 we have written the branching fraction assuming that dark matter is a Dirac fermion. If the dark matter were instead a Majorana fermion, the branching fraction would be larger by a factor of 2 (with a factor of 4 arising from the squared matrix element, and a factor of 1/2 from the reduced phase space of the final state particles). The result for operator S3, and the corresponding constraints, match that found in [35].
Note that if dark matter is spin-1, the decay rates have terms which scale as $m_X^{-2}$ or $m_X^{-4}$. These terms arise from final states if either one or both of the dark matter particles is longitudinally polarized [7].

In general, unitarity constrains the magnitude of the squared matrix element for both elastic and inelastic scattering. One may worry that, for small $m_X$, unitarity would require the presence of corrections which invalidate the use of the contact operator approximation at tree-level; for example, if the dark matter is a gauge boson which has become massive due to spontaneous breaking of a dark sector symmetry, then one may also need to include additional diagrams involving the fields responsible for the spontaneous symmetry breaking. However, the constraints from unitarity are trivial when initial particles are at rest, because the elastic scattering cross-section is at threshold (for example, see [50]). As a result, the contact operator approximation is consistent with unitarity for the meson decay process in the non-relativistic bound state limit. This approximation becomes more problematic as one departs from the non-relativistic limit, but that analysis is beyond the scope of this work.

B. Dark Matter Annihilation

Applying crossing symmetry to the quark annihilation matrix elements yields the matrix element for dark matter annihilation to quarks. If the dark matter can annihilate from an $s$-wave initial state, then gamma ray observations can be used to set limits on $\Lambda$. But one should note that, although we have restricted ourselves to effective operators which have a non-zero matrix element with an $s$-wave meson state, these operators need not necessarily have a non-zero matrix element with an $s$-wave dark matter initial state. Only five of these operators which we consider can permit unsuppressed dark matter annihilation. The corresponding annihilation cross-sections, at tree-level, are given by

$$
\langle \sigma^F_5 v \rangle = \frac{3}{2\pi \Lambda^4} \left( 1 - \frac{m_q^2}{m_X^2} \right)^{1/2} \left( 2m_X^2 + m_q^2 \right),
$$

$$
\langle \sigma^F_9 v \rangle = \frac{6}{\pi \Lambda^4} \left( 1 - \frac{m_q^2}{m_X^2} \right)^{1/2} \left( m_X^2 + 2m_q^2 \right),
$$

$$
\langle \sigma^{F10} v \rangle = \frac{6}{\pi \Lambda^4} \left( 1 - \frac{m_q^2}{m_X^2} \right)^{3/2} m_X^2,
$$

$$
\langle \sigma^V_5 v \rangle = \frac{2}{3\pi \Lambda^2} \left( 1 - \frac{m_q^2}{m_X^2} \right)^{3/2},
$$

$$
\langle \sigma^V_6 v \rangle = \frac{2}{3\pi \Lambda^2} \left( 1 - \frac{m_q^2}{m_X^2} \right)^{1/2} \left( 1 + \frac{2m_q^2}{m_X^2} \right).$$

(6)
These cross sections can be bounded by a stacked analysis of the number of photons arriving from dwarf spheroidal galaxies [51, 55]. The number of photons expected to result from dark matter annihilation is the product of an astrophysics-dependent factor and a particle-physics dependent factor. The astrophysics-dependent factor depends on the density profile of the dwarf spheroidal galaxies, and can be estimated from the rotation curves of visible matter. The particle-physics dependent factor can be expressed as

$$\Phi_{PP} = \frac{\langle \sigma_A v \rangle}{8\pi m_X^2} \int_{E_{thr}}^{m_X} \sum_f B_f \frac{dN_f}{dE} dE,$$  

(7)

where $B_f$ is the branching ratio for dark matter to annihilate to a channel $f$ and $E_{thr} = 1\, \text{GeV}$ is the photon energy analysis threshold. $dN_f/dE$ is the photon spectrum for a given dark matter annihilation channel. Note that our bounds from dark matter annihilation assume universal quark coupling, with the strongest contribution to the photon spectrum coming from the $u$- and $d$- quark channels.

The 95% CL limit on $\Phi_{PP}$ arising from observations by Fermi-LAT of dwarf spheroidal galaxies is given by [52]

$$\Phi_{PP} < 5.0^{+4.3}_{-4.5} \times 10^{-30} \, \text{cm}^3 \, \text{s}^{-1} \, \text{GeV}^{-2},$$  

(8)

where the asymmetric uncertainties are 95% CL systematic errors [53], resulting from the uncertainty in the mass density profiles of the various satellites. This bound assumes that dark matter is its own anti-particle (if the dark matter field is complex, as in the cases we consider, this number increases by a factor of 2). The associated bounds on the annihilation cross sections are plotted in Figure 1. For each annihilation channel, the photon spectrum was produced by Pythia 6.403 [56] and accounts for collinear photons generated during the quark final state parton shower evolution.

The general shape of the bound on the annihilation cross-section is easily understood; for larger dark matter mass, the bound strengthens with decreasing mass because of the resulting increase in the dark matter number density. At lower masses the bound begins to sharply weaken because it is no longer possible to create photons above the analysis threshold. We do not plot bounds on dark matter annihilation for $m_X \lesssim 4\, \text{GeV}$ since the energy of the final state quarks becomes close to the hadronization scale. Annihilation constraints for lighter dark matter would also be cut off by either, for the $\bar{c}c$ and $\bar{s}s$ channels, the threshold to create quark/antiquark final states or, for the $\bar{u}u$ and $\bar{d}d$ channels, the analysis threshold.

Given any particle physics model for the dark matter annihilation cross section and final state branching fractions, one can then determine if the model is consistent with data from Fermi-LAT [58]. Note, however, that since all of the relevant operators are only non-vanishing when dark matter is complex, it is consistent with this analysis for dark matter to
be asymmetric. In this case, only the particle would be abundant at the current epoch, not the anti-particle; the bounds on dark matter annihilation in dwarf spheroidal galaxies would thus necessarily be unconstraining.

![Graph showing bounds on the annihilation cross section](image)

**FIG. 1.** Bounds on the annihilation cross section, $\langle \sigma_A v \rangle$, for dark matter of mass $m_X$ annihilating to quarks in dwarf spheroidal galaxies. Note that the results for annihilation to the $u\bar{u}$ and $d\bar{d}$ channels are visually identical.

It should be noted that, for $m_X \lesssim \mathcal{O}(1)$ GeV, $s$-wave dark matter annihilation can be strongly constrained by anisotropies in the cosmic microwave background (CMB) arising from the energy deposition caused by annihilation during recombination [57]. But since we do not assume any coupling between dark matter and leptons, the only relevant annihilation channels are $\bar{X}X \rightarrow \bar{q}q$. For these final states, if $m_X \lesssim 1$ GeV, it is more difficult to determine CMB constraints on LDM because of the effects of hadronization. A full treatment of these constraints is thus beyond the scope of this work. But one can note that, in general, the constraints on the scale $\Lambda$ arising from CMB studies (as with gamma-ray searches of dwarf spheroidal galaxies) either tend to become weaker as $m_X$ decreases, for fermionic dark matter, or strengthen less, relative to collider-based constraints, for vector dark matter.

**C. Nuclear Scattering**

Several of the effective operators which we consider will also yield velocity independent terms in the associated dark matter-nucleon scattering matrix element.
Operators with vector quark bilinears will yield spin-independent scattering, with associated cross sections given by

\[ \sigma_{SI}^{p,n} = \frac{\mu_p^2}{32\pi(2J_X + 1)} \sum_{\text{spins}} \left| \sum_{q} \frac{B_q^{p,n}}{m_{Xq} m_q} \mathcal{M}_{Xq \to Xq} \right|^2, \]  

(9)

where \( \mu_p \) is the reduced mass of the dark matter-nucleon system and \( J_X \) is the dark matter spin. The nucleon form factors associated with the vector quark bilinear are \( B_u^p = B_d^n = 2 \), \( B_u^n = B_d^p = 1 \), and \( B_{s,c,b,t}^{p,n} = 0 \) [59].

Similarly, operators with tensor quark bilinears will yield spin-dependent scattering, with associated cross sections given by

\[ \sigma_{SD}^{p,n} = \frac{\mu_p^2}{32\pi(2J_X + 1)} \sum_{\text{spins}} \left| \sum_{q} \frac{\delta_q^{p,n}}{m_{Xq} m_q} \mathcal{M}_{Xq \to Xq} \right|^2, \]  

(10)

The nucleon spin form factors \( \delta_q^{p,n} \) can be extracted from data and are roughly given by \( \delta_u^p = 0.54^{+0.09}_{-0.22} \) and \( \delta_d^p = -0.23^{+0.09}_{-0.16} \) [60, 61]. These form factors are in slight disagreement with lattice calculations [62, 63], although using alternative form factors would only change our scattering cross sections by a factor of order unity.

Note that for contact operators which we consider, only the coupling to first-generation quarks is relevant for scattering. In general, dark matter can exhibit velocity-independent spin-independent scattering from nucleons through interactions with heavy quarks, but only if the dark matter bilinear is a scalar. We do not consider such operators, however, because they cannot contribute to the decay of a \( 1^- \) meson. Thus, the bounds on dark matter-quark interactions arising from \( \Upsilon(1S) \) or \( J/\psi \) decay can only be related to bounds on dark matter-nucleon scattering if one makes a particular choice for the relative strength of dark matter coupling to light and heavy quarks.

III. RESULTS

We assume that dark matter couples to quarks only through a single effective operator, but with equal coupling to all quark flavors. An example of a model which would yield this effective operator realization would be the case where the mediating particle was a massive vector boson for a new \( U(1) \) symmetry under which all quarks have equal charge and under which the dark matter is also charged (for example, this \( U(1) \) could be a linear combination of \( U(1)_{\text{baryon}} \) and another \( U(1) \) symmetry under which the dark matter is charged). But, of course, other choices are possible and can be well-motivated by other UV completions. We focus on the case of equal couplings simply as a benchmark. We note also that the coefficient of the effective contact operator can exhibit RG-running between the energy scale relevant
for meson decay and the scale relevant for nuclear scattering; however, this is expected to be a relatively small effect [65, 66].

A. Mediator Scale

In Figure 2 we plot bounds on $\Lambda$, as a function of $m_X$, arising from limits on invisible $\Upsilon(1S) \rightarrow \text{nothing}$ decays and from dark matter annihilation in dwarf spheroidal galaxies. Note that the bounds arising from $\Upsilon(1S)$ decays are sensitive only to dark matter coupling to $b$-quarks. But for values of $m_X$ for which $\Upsilon(1S)$ decays are kinematically allowed, dark matter annihilation to $\bar{b}b$ is kinematically forbidden\(^2\) instead, the bounds on dark matter annihilation arise from dark matter couplings to the lighter quarks. Note that the mediator scale, $\Lambda$, is always larger than $\sim 10$ GeV, implying that the contact operator approximation is valid at all of the colliders which are relevant for meson decay bounds.

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\(^2\) For the case where $m_X > m_b$, bounds arising from radiative $\Upsilon$ decays at the LHC have been discussed in [32, 33].
always $p$-wave if mediated by a contact operator which can also mediate the decay of a $1^{--}$ meson, and it is thus unconstrained by gamma ray observations. Bounds on dark matter annihilation in dwarf spheroidal galaxies are not constraining for $m_X \leq 1$ GeV, because very light dark matter cannot produce any photons above the analysis threshold. Bounds on the dark matter annihilation cross section have systematic uncertainties related to the dark matter density profile; these uncertainties can weaken these bounds by up to a factor of 2 or strengthen them by up to a factor of 10.

We also plot bounds on the scale $\Lambda$ for each of the operators (for spin-1/2 dark matter) arising from searches for monojet and monophoton production at the LHC [39–42] (similar bounds can also be found for mono-$W,Z$ production [43]). These searches place bounds on the cross section for the process $pp \to XX + jet, \gamma$, where dark matter interacts with Standard Model quarks through a contact operator. As such, these bounds are very similar in spirit to bounds on dark matter-quark couplings arising from meson decay, and in particular these bounds do not dramatically worsen for very light dark matter. However, it is important to note one significant difference. These LHC monojet bounds are only valid if the contact operator approximation is valid even at the energies of LHC processes, which requires the mediator mass to be $\mathcal{O}(\text{TeV})$. For lighter mediators, one cannot perform a model-independent operator analysis; although stringent bounds may be possible [37, 38], they depend on the details of dark matter-quark interaction. Of course, the bounds arising from meson decay are also only valid if the contact operator approximation is valid, but in this case the relevant energy scale is the meson mass, which is $\mathcal{O}(1 – 10$ GeV). For mediators much heavier than $\sim 10$ GeV, the bounds which arise from meson decay will be valid.

Dark matter with $m_X \gtrsim 10$ MeV is heavy enough to be consistent with the cold dark matter paradigm. Note, however, that our analysis does not require that dark matter is a thermal relic or that dark matter annihilation be predominantly to Standard Model final states. Instead, our focus is on constraints on dark matter interactions arising from astrophysical observations of the annihilation products in the current epoch, not on constraints on the mechanism of dark matter generation. Even if the the cross section for dark matter annihilation to visible matter is small, there could be a large branching fraction for annihilation to the dark sector, evading indirect detection constraints but allowing for a thermal relic density which can satisfy observational constraints.

B. Complementary Dark Matter Scattering

Of the contact operators which we consider, the ones which permit velocity-independent scattering are F5 (SI), F9 (SD), S3 (SI), V3 (SI), and V5 (SD). The corresponding total dark
matter-proton scattering cross sections are given by

\[ \sigma_{SI}^{F5} = \frac{\mu_p^2}{4\pi\Lambda^4} (B_u^p + B_d^p)^2, \]

\[ \sigma_{SI}^{S3} = \sigma_{SI}^{V3} = \frac{\mu_p^2}{4\pi\Lambda^4} (B_u^p + B_d^p)^2, \]

\[ (11) \]

\[ \sigma_{SD}^{F9} = \frac{12\mu_p^2}{\pi\Lambda^4} (\delta_{up} + \delta_{dp})^2, \]

\[ \sigma_{SD}^{V5} = \frac{2\mu_p^2}{\pi\Lambda^2 m_X^2} (\delta_{up} + \delta_{dp})^2. \]

In Figure 3, we plot the bounds on spin-independent (left panel) and spin-dependent (right panel) scattering mediated by each of the relevant operators. We also plot 95% CL bounds arising from Fermi-LAT constraints on dark matter annihilation in dwarf spheroidal galaxies and 90% CL bounds arising from monojet and monophoton searches (CMS [39, 40] and ATLAS [41, 42]). The DAMA/LIBRA [67], CRESST II (95% CL) [68], CoGeNT [69], and CDMS II (Silicon) [70] 90% CL signal regions are also shown, as are the 90% CL exclusion contours from SuperCDMS [71], LUX [72], SIMPLE [73], PICASSO [74], and COUPP [75].

Various experiments are able to set bounds on DM-electron scattering [28, 76], but some assumption of universal dark matter coupling to quarks and leptons is required to, in turn, bound dark matter-nucleon scattering.

Similar bounds on dark matter interactions can be obtained from bounds on \( J/\Psi \) invisible decays, and are presented in the appendix. While these bounds are weaker than those obtained from invisible \( Y(1S) \) decays, they are valid for a larger range of mediator masses \( (\gtrsim M_{J/\Psi}) \) and directly probe the couple of dark matter to \( c \)-quarks, thus providing non-trivial complementary constraints. Note that the contact operator approximation will begin to break down for mediator scales smaller than \( \sim 10 \) GeV.

IV. CONCLUSIONS

We have presented bounds on dark matter-quark contact interactions which can be obtained from high luminosity \( B/charm \)-factories by constraining decays of the form \( Y(1S), J/\psi \to \bar{X}X \). These bounds on low mass dark matter probe a mass range significantly below the threshold of direct dark matter detection experiments and complement bounds on dark matter interactions obtained from gamma ray searches of dwarf spheroidal galaxies and from monojet/monophoton/mono-\( W,Z \) searches at hadron colliders.

In particular, the effective interactions which permit decay of a \( 1^{--} \) meson state can also permit velocity-independent dark matter-nucleon scattering (either spin-independent or
FIG. 3. Bounds on the dark matter-proton spin-independent (left panel) and spin-dependent (right panel) scattering cross section for dark matter of mass $m_X$ coupling universally to quarks through the indicated effective contact operator. The labeled exclusion contours indicate 90% CL bounds arising from limits on invisible decays of $\Upsilon(1S)$, 95% CL bounds arising from Fermi-LAT constraints on dark matter annihilation in dwarf spheroidal galaxies, and 90% CL bounds arising from monojet searches (CMS [39, 40] and ATLAS [41]). The DAMA/LIBRA [67], CRESST II (95% CL) [68], CoGeNT [69], and CDMS II(Silicon) [70] 90% CL signal regions are also shown, as are the 90% CL exclusion contours from SuperCDMS [71], LUX [72], SIMPLE [73], PICASSO [74], and COUPP [75].

spin-dependent). For $m_X \sim 1 - 5$ GeV, the bounds obtained from meson decay can thus potentially complement those obtained from direct detection experiments. For the case of spin-independent scattering, direct detection experiments already place bounds which well exceed those obtained from meson decay. However, for spin-dependent scattering, bounds arising from $\Upsilon(1S)$ decay via the F9 and V5 operators are comparable to those obtained from direct detection experiments. This is not surprising, as direct detection experiments typically have much weaker sensitivity to spin-dependent scattering, due to the lack of constructive interference in coherent scattering.

Moreover, bounds on spin-1 dark matter interactions improve dramatically as $m_X$ decreases, because of the enhancement in the matrix element which arises when the dark matter particles are longitudinally polarized. Relating the suppression scale $\Lambda$ to the mediator mass scale $m_{\text{med}}$, and coupling $g$ by $\Lambda \sim m_{\text{med}}/g$, this implies that interactions between spin-1 dark matter and quarks can be constrained even if the coupling is very weak.

We have seen that invisible quarkonium decays probe the same parton-level process ($\bar{q}q \rightarrow \bar{X}X$) as monojet/photon/$W,Z$ searches at hadron colliders. However, quarkonium decays
provide complementary information, allowing robust probes of models with relatively light mediators (\(\gtrsim 10\ \text{GeV}\)) for which the contact operator approximation would fail at the LHC. It is also worth noting that, since the heavy quarkonium bound states are non-relativistic, searches based on invisible heavy quarkonium decays readily distinguish between DM-SM interactions which vanish in the limit of non-relativistic quarks and those which do not. This probe thus nicely complements LHC searches, in which the partons are highly relativistic.

A similar analysis can be performed of heavy quarkonium decays to a photon and missing energy; although the set of relevant contact operators would be different for such an analysis, those branching fractions are much more tightly constrained. Improved bounds on the \(\Upsilon(1S) \rightarrow \text{nothing}\) decay rate from Belle II and a factor of \(\sim 14\) enhancement in sensitivity to \(J/\Psi \rightarrow \text{nothing}\) decay rate from BESIII [77] will allow for even tighter constraints on the interactions of low mass dark matter with Standard Model particles.

Acknowledgements

We are grateful to Y. G. Aditya, T. Browder, F. Harris, D. Marfatia, A. Petrov, A. Rajaraman, X. Tata, S. Vahsen, D. Walker and A. Wijangco for useful discussions. This work is supported in part by Department of Energy grant de-sc0010504.

Appendix A: Constraints from \(J/\Psi\) Decay

In Figure 4, we plot bounds on \(\Lambda\), as a function of \(m_X\) arising from limits on invisible \(J/\Psi\) decays and from dark matter annihilation in dwarf spheroidal galaxies. We also plot bounds on the scale \(\Lambda\) for each of the operators (for spin-1/2 dark matter) arising from searches for monojet and monophoton production at the LHC [39–42] (similar bounds can also be found for monophoton or mono-W,Z production [43]). Note that the contact operator approximation will begin to break down for mediator scales smaller than \(\sim 10\ \text{GeV}\).

In Figure 5, we plot the bounds on spin-independent (left panel) and spin-dependent (right panel) scattering mediated by each of the relevant operators. We also plot 95% CL bounds arising from Fermi-LAT constraints on dark matter annihilation in dwarf spheroidal galaxies and 90% CL bounds arising from monojet and monophoton searches (CMS [39, 40] and ATLAS [41, 42]). The DAMA/LIBRA [67], CRESST II (95% CL) [68], CoGeNT [69], and CDMS II(Silicon) [70] 90% CL signal regions are also shown, as are the 90% CL exclusion contours from SuperCDMS [71], LUX [72], SIMPLE [73], PICASSO [74], and COUPP [75].

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