Constraints on Sub-GeV Dark-Matter–Electron Scattering from the DarkSide-50 Experiment
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The nature of dark matter (DM) remains unknown despite several decades of increasingly compelling gravitational evidence [1–5]. While the most favored candidate in a particle physics interpretation is the Weakly Interacting Massive Particle (WIMP) [6, 7], which obtains its relic abundance by thermal freeze-out through weak interactions, there is as yet no unambiguous evidence of WIMP direct detection, warranting searches for other possible DM paradigms.

Another well-motivated class of DM candidates is sub-GeV particles interacting through a vector mediator with couplings smaller than the weak-scale. These light DM candidates arise in a variety of models [8–12], and there are a number of proposed mechanisms that naturally obtain the expected relic abundance for light DM [13–27]. Light DM may have couplings to electrons, and because the energy transferred by the DM particle to the target depends on the reduced mass of the system, electron targets provide more efficient absorption of the kinetic energy of sub-GeV-scale light DM than a nuclear target [28].

There is currently a substantial experimental effort to search for light DM through multiple techniques, see Refs. [29, 30] and references therein. In particular, dual-phase time projection chambers (TPCs) are an excellent probe of light DM, which can ionize atoms to create an electroluminescence signal (S2) even when the corresponding prompt scintillation signal (S1), typically used to identify
nuclear recoils, is below the detector threshold [31, 174.
In this letter, we present the first limits on light DM-electron scattering from the DarkSide-50 experiment[176
(DS-50). This analysis closely follows Ref. [32, 177
which contains additional details about the detect-
tor, data selection, detector response, and cut-effi-
ciences.

DS-50 is a dual-phase time projection chamber[188
with a (46.4 ± 0.7) kg target of low-radioactivity un-
derground argon (UAr) [33–36] outfitted with 38 3” PMTs, 19 above the anode and 19 below the cathode [37]. Particle interactions within the target[189
volume create primary UAr scintillation (S1) and ion-
ized electrons. These electrons are drifted tow-
ards the anode of the TPC and extracted into a gas layer where they create gas-proportional scin-
tillation (S2). The electron extraction efficiency is no
ter than 99.9% [38]. While the trigger efficiency[190
for S1 signals drops to zero below approximately 0.6 keVee, the S2 trigger efficiency remains 100% [191
above 0.05 keVee due to the high S2 photon yield[192
per electron, (23 ± 1) PE/e− in the central PMT as[193
measured by single-electron events caused by impu-
rities within the argon that trap and release single
des. S2 signals are identified offline using a soft-
ware pulse finding algorithm that is effectively 100%
efficient above 0.05 keVee, and a set of basic cuts[195
are applied to the data to reject spurious events. A fiducial cut is then applied that only accepts events[196
whose maximum signal occurs within one of the cen-
tral seven PMTs in the top PMT array. After all cuts[197
are applied, the detector acceptance is (0.43 ± 0.01)% due[198
almost entirely to fiducialization. A correction is[199
applied to events that occur under the six PMTs[200
surrounding the central one to correct for a radial
variation in photon yield observed in 83Kr source
data.

A DM particle may scatter off a bound electron
within the DS-50 detector, ionizing an argon atom.[201
We evaluate the dark matter recoil spectra for ar-
gon following the calculation of Refs. [28, 39]. The
velocity averaged differential ionization cross sec-
tion for bound electrons in the (n, l) shell is given by[202
\[
\frac{d\langle \sigma v \rangle_{\text{ion}}}{d \ln E_{\text{cr}}} = \frac{\sigma_e}{8 \mu_e^2}
\times \int dq |f_{\text{ion}}^{nl}(k', q)|^2 |F_{\text{DM}}(q)|^2 \eta(v_{\text{min}}),
\]
where the reference cross section, \(\sigma_e\), parameterizes[204
the strength of the interaction and is equivalent to[205
the cross section for elastic scattering on free elec-
trons; \(\mu_e\) is the DM-electron reduced mass; \(q\) is the[206
3-momentum transfer; \(f_{\text{ion}}^{nl}(k', q)\) is the ionization[207
form-factor, which models the effects of the bound[208
electron initial state and the outgoing final state per-
turbed by the potential of the ion from which the
electron escaped; \(k'\) is the electron recoil momen-
tum; \(F_{\text{DM}}(q)\) is the DM form factor; and the DM
velocity profile is encoded in the inverse mean speed
function, \(\eta(v_{\text{min}}) = \frac{1}{2} \Theta(v - v_{\text{min}})\), where \(v_{\text{min}}\) is
the minimum velocity required to eject an electron
with kinetic energy \(E_{\text{cr}}\) given the momentum trans-
fer \(q\) and \(\Theta\) is the Heaviside step function.

The details of the argon atom’s electronic struc-
ture and the outgoing state of the recoil electron are
contained in \(f_{\text{ion}}^{nl}(k', q)\), which is a property of the
argon target and independent of the DM physics.
Computing \(f_{\text{ion}}^{nl}(k', q)\) requires one to model both
the initial bound states and the final continuum
outgoing states of the electron. The target elec-
trons are modeled as single-particle states of an
isolated argon atom described by the Roothaan-Hartree-Fock wavefunctions. This conservatively
neglects the band structure of liquid argon which, if
included, should enhance the total electron yield
due to the decreased ionization energy in the liq-
uid state [40]. The recoil electron is modeled as the
full positive-energy wavefunction obtained by solv-
ing the Schrödinger equation with a hydrogenic po-
tential of some effective screened charge \(Z_{\text{eff}}\) [41].
We choose a \(Z_{\text{eff}}\) that reproduces the energy levels
of the argon atom assuming a pure Coulomb poten-
tial. Further details on the computation of \(f_{\text{ion}}^{nl}(k', q)\)
are provided in the Appendix.

The DM form factor, \(F_{\text{DM}}(q)\), parametrizes the
fundamental momentum transfer dependence of the
DM-electron interaction and has the following lim-
iting values:

\[
F_{\text{DM}}(q) = \frac{m_A^2 + \alpha^2 m_e^2}{m_A^2 + q^2} \approx \begin{cases} 1, & m_{A'} \gg \alpha m_e \\ \frac{\alpha^2 m_e^2}{q^2}, & m_{A'} \ll \alpha m_e \end{cases},
\]
where \(m_{A'}\) is the mass of the vector mediator, \(m_e\)
is the electron mass, and \(\alpha\) is the fine-structure con-
stant. Because \(F_{\text{DM}}(q)\) is dimensionless by defini-
tion, the form factor needs to be defined with re-
spect to a reference momentum scale. The conven-
tional choice is \(q_0 = \alpha m_e = 1/a_0\), where \(a_0\) is the
Bohr radius, because this is typical of atomic mo-
menta. The case where \(F_{\text{DM}}(q) = 1\) corresponds to
the “heavy mediator” regime, where \(m_{A'}\) is much
larger than the typical momentum scale. The case
where \(F_{\text{DM}}(q) \times 1/q^2\) corresponds to the “light me-
diator” regime.

The inverse mean speed, \(\eta(v_{\text{min}})\), is defined
through the DM velocity distribution in the same way
as for GeV-scale WIMPs and nuclear scattering.
We have assumed the Standard Halo Model
with escape velocity $v_{\text{esc}} = 544 \text{ km/s}$ [42], circular velocity $v_0 = 220 \text{ km/s}$, and the Earth velocity as specified in [43] and evaluated at $t = 199$ days ($v_E \approx 244$ km/s), the median run live-time for DarkSide-50. Note that the definition of $v_{\text{min}}$ is different for electron scattering from a bound initial state than for elastic nuclear recoils. The relation $E_R = q^2/2m_N$, which is valid in two-body elastic scattering, no longer holds. For a bound electron with principal quantum number $n$ and angular momentum $l$ [39]

$$v_{\text{min}}(q, E_{nl}^b, E_{\text{er}}) = \frac{|E_{nl}^b| + E_{\text{er}}}{q} + \frac{q}{2m_X}, \quad (3)$$

where $|E_{nl}^b| + E_{\text{er}}$ is the total energy transferred to the ionized electron, which is a sum of the energy needed to overcome the binding energy, $E_{nl}^b$, and the recoil energy of the outgoing electron, $E_{\text{er}}$.

The velocity averaged differential ionization cross section, Eq. 1, is used to calculate the DM-electron differential ionization rate,

$$\frac{dR}{d\ln E_{\text{er}}} = N_T \frac{\rho_X}{m_X} \sum_{nl} \frac{d\langle\sigma v_{\text{ion}}^l\rangle}{d\ln E_{\text{er}}}, \quad (4)$$

where $N_T$ is the number of target atoms per unit mass, $\rho_X = 0.4 \text{ GeV/cm}^3$ is the local DM density used in Ref. [39], and $m_X$ is the DM mass. The sum is over the outer-shell 3p (16.08 eV binding energy) and 3s (34.76 eV binding energy) electrons. Fig. 1 shows the contributions of the individual atomic shells to the total DM-electron scattering rate. For low electron recoil energies, the outer shell contribution (3p) dominates, while at higher energy, the contribution from the 3s shell increases. This behavior becomes more pronounced as the DM mass increases. The same behavior is observed for the differential ionization rate, section, Eq. 1, is used to calculate the DM-electron recoil rate assuming a WIMP-electron cross section of $10^{-40} \text{ cm}^2$ and $F_{\text{DM}} = 1$ for a 100 MeV/c$^2$ DM particle (solid) and a 1000 MeV/c$^2$ DM particle (dashed).

The calculated DM-electron recoil spectra are converted to the ionization spectra measured in DS-50 using a scale conversion based on a fit to low energy peaks of known energy, as shown in Fig. 2 and described in [32]. The resulting ionization spectra are then smeared assuming the ionization yield and recombination processes follow a binomial distribution and convolved with the detector response, measured from single-electron events [32]. This procedure correctly reconstructs the measured width of the $^{35}$Ar K-shell (2.82 keV) and L-shell (0.27 keV) peaks. The expected DM-electron scattering ionization spectra in the case of a heavy mediator, $F_{\text{DM}} = 1$, and in the case of a light mediator, $F_{\text{DM}} \propto 1/q^2$, are shown in Fig. 3.

We use a 500 day dataset collected between April 30, 2015, and April 25, 2017, corresponding to a 6786.0 kg d exposure, to place limits on DM with masses below 1 GeV/c$^2$. The 500 day ionization spectrum used for the search is shown in Fig. 3.
FIG. 3. The 500 day DarkSide-50 ionization spectrum compared with predicted spectra from the G4DS background simulation [46]. These are the same data and background spectra shown in Ref. [32]. Also shown are calculated DM-electron scattering spectra for DM particles with masses $m_{\chi}$ of 10, 100, and 1000 MeV/c$^2$, reference cross section $\sigma_e = 10^{-36}$ cm$^2$/stop (top) and $\sigma_e = 10^{-33}$ cm$^2$/stop (bottom), and $F_{DM}(q) = 1/30$ (top) and $F_{DM}(q) \propto 1/q^2$ (bottom). The vertical dashed line indicates the $N_{e^-} = 3$ analysis threshold.

Limits are calculated using a binned profile likelihood method implemented in RooStats [47–49]. We use an analysis threshold of $N_{e^-} = 3$, approximately equivalent to 0.05 keVee, lower than the threshold used in [32]. This increases the signal acceptance at the expense of a larger background rate from coincident single-electron events, which are not included in the background model and contribute as signal during the limit calculation. The background model used in the analysis is determined by a detailed Monte Carlo simulation of the DarkSide-50 apparatus. Spectral features at high energy are used to constrain the simulated radiological activity within detector components to predict the background spectrum in the region of interest [50]. The predicted spectrum is plotted alongside the data in Fig. 3 and described in greater detail in [32]. During the analysis, the overall normalization of the background model is constrained near its predicted value by a Gaussian nuisance term in the likelihood function. Additional gaussian constraints on the backgrounds and signal spectral shape are included based on the uncertainty of the fit in Fig. 2 and the uncertainty in the S2 to $N_{e^-}$ conversion factor, extracted from single-electron data.

The resulting 90% C.L. limits are shown in Fig. 4 for two assumptions of DM form-factors, $F_{DM}(q) = 1$ and $F_{DM}(q) \propto 1/q^2$. In the case of a light mediator, $F_{DM}(q) \propto 1/q^2$, the constraints from DS-50 are not as stringent as the XENON10 experiment due to the higher ($N_{e^-} = 3$) analysis threshold adopted in this work but better than the XENON10 limit due the lower background rate. For a heavy mediator, $F_{DM}(q) = 1$, we improve the existing limits from XENON10 and XENON100 [39] for dark matter masses between 30 MeV/c$^2$ to 100 MeV/c$^2$, seeing a factor of 3 improvement at 50 MeV/c$^2$.

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APPENDIX

Here we provide additional details on the DM-electron scattering rate calculation described in the text. The explicit forms of the radial part of the wavefunction used to compute the atomic form factor, $|f_{ion}^{nl}(k',q)|^2$, are given by the Roothaan-Hartree-Fock (RHF) wavefunctions [53], which are linear combinations of Slater-type orbitals:

$$R_{nl}(r) = a_0^{-3/2} \sum_j C_{jln} (2Z_j)^{n_j'+1/2} \sqrt{(2n_j')!} \times \left( \frac{r}{a_0} \right)^{n_j'-1} e^{-Z_j r/a_0},$$

where the coefficients $C_{jln}$, $Z_j$, and $n_j'$ are given in Ref. [53].

In the literature, different procedures have been used to approximate the outgoing electron wavefunction in such scattering scenarios. One common approximation is to treat the final state as a pure plane-wave corrected by a Fermi factor,

$$F(k', Z_{eff}) = \frac{2\pi Z_{eff}}{k' a_0} \frac{1}{1 - e^{-2\pi Z_{eff}/(k'a_0)}},$$

which parameterizes the distortion of the outgoing electron wavefunction by the effective screened Coulomb potential of the nucleus. While the approximate shape of the ionization form factors, $f_{ion}^{nl}$, are consistent between the plane-wave solution and the continuum-state solution used in this work, the detailed structure does vary between the two. At large momentum transfers, the plane-wave and continuum solutions approach each other, but they diverge at lower momentum transfers where the form factor is dominated by the overlap between the bound and continuum wavefunctions near the origin. This is because the Fermi factor reproduces the behavior of the full wavefunction at the origin, but outer-shell orbitals have most of their support away from the origin, such that the overlap with the outgoing wavefunction is maximized away from the origin. Thus, smaller atoms and inner shells have better agreement. For this reason, the discrepancy between using continuum versus plane-wave final states is smaller for argon than for xenon. We however choose to use the full-continuum solutions for the presentation of all final results.

The continuum-state solutions to the Schrödinger equation with potential $-Z_{eff}/r$ have radial wavefunctions indexed by $l$ and $k$, given by [41]

$$\tilde{R}_{kl}(r) = (2\pi)^{3/2}(2k)^l \sqrt{\frac{2}{\pi}} \Gamma \left( l + 1 - \frac{iZ_{eff}}{k a_0} \right) \frac{e^{-ikr}}{(2l + 1)!} \times e^{-ikr} \frac{1}{F_1} \left( l + 1 + \frac{iZ_{eff}}{k a_0}, 2l + 2, 2ikr \right). \quad (A.7)$$

The ratio of the wavefunction at the origin to the wavefunction at infinity gives the Fermi factor:

$$\left| \frac{\tilde{R}_{kl}(r = 0)}{\tilde{R}_{kl}(r = \infty)} \right|^2 = F(k, Z_{eff}).$$

The normalization for these unbound wavefunctions is

$$\int dr r^2 \tilde{R}_{kl}^*(r) \tilde{R}_{kl'}(r) = (2\pi)^3 \frac{1}{k^6} \delta_{ll'} \delta(k - k'), \quad (A.9)$$

so that $\tilde{R}_{kl}(r)$ itself is dimensionless. In terms of these wavefunctions, the ionization form factor is given by

$$|f_{ion}^{nl}(k', q)|^2 = \frac{4k'^3}{(2\pi)^3} \sum_{l'} \sum_{L=|l'-l|} (2l + 1)(2l' + 1)(2L + 1) \times \left[ \frac{l'}{0} \frac{l'}{0} \frac{L}{0} \right]^2 \int dr r^2 \tilde{R}_{kl'}(r) R_{nl}(r) j_L(qr) \left| j_{l'}(qr) \right|^2 \quad (A.10)$$

The term in brackets is the Wigner-3$j$ symbol evaluated at $m_1 = m_2 = m_3 = 0$, and $j_L$ is the spherical Bessel function of order $L$.

Following [31, 39], the procedure used to determine $Z_{eff}$ is:

1. Treat the bound-state orbital $R_{nl}$ as a bound state of a pure Coulomb potential $-Z_{eff}^2/r$, rather than the self-consistent potential giving rise to the RHF wavefunctions.

2. Determine $Z_{eff}^n$ by matching the energy eigenvalue to the RHF eigenvalue.

3. Use this $Z_{eff}^n$ to construct all $\tilde{R}_{kl'}(r)$ in the sum in Eq. (A.10).

For example, for the 3p shell of argon, $E_{3p} = 16.08$ eV, so we solve

$$13.6 \text{ eV} \times \left( \frac{Z_{eff}^{3p}}{3.26} \right)^2 = 16.08 \text{ eV} \quad \Rightarrow \quad Z_{eff}^{3p} = 3.26.$$