

Dark-Matter–Electron Detectors for Dark-Matter–Nucleon Interactions

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In a seminal paper now a decade old, it was shown that dark-matter detectors geared at probing interactions with nucleons could also be used to probe dark-matter interactions with electrons. In this Letter, we show that new detector concepts designed to probe dark-matter–electron interactions at low masses can similarly be used to probe new parameter space for dark-matter–nucleon interactions. We demonstrate the power of this approach by using existing data from superconducting detectors to place new limits on the interactions of nuclei with MeV-scale dark matter. Further, we show that advances in detector technology that have been anticipated for electronic interactions will automatically extend sensitivity deep into uncharted territory for nuclear interactions. This doubles the effective science output of future low-threshold experiments.

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Introduction—The nature of dark matter (DM) remains one of the largest open problems in modern physics. For decades, most laboratory DM searches operated under the assumption that DM particles would be associated with the weak scale of the standard model, and thus targeted heavy DM particles with masses well above the GeV scale [1–5]. In the wake of null results, the community has shifted focus to a broader set of candidates and has developed searches sensitive to different mass scales and interactions. One notable step of this transition was the realization by Essig *et al.* [6] that existing experiments designed to search for DM-nucleon scattering could also be used to probe DM-electron scattering at masses well below the GeV scale. This set off a new exploratory phase of DM searches, and sensitivity to DM-electron interactions was quickly extended to DM masses well below those probed via nuclear recoils.

Since then, advancements in technology have yielded many new detection modalities based on a variety of different electronic systems, offering sensitivity to extremely low energy deposits and thus to low DM masses [6–24]. Among the proposed platforms, superconductors [9–13] have been quickly advancing in both threshold and scale. The small gap of the superconducting phase, $\mathcal{O}(\text{meV})$ in common

materials, allows in principle for the detection of DM down to the keV scale, where cosmological bounds become relevant. Several new experiments leveraging such new ideas are already underway [25–28], and prototype superconducting detectors have already produced new limits on DM-electron interactions [11,13,29].

However, a complete and cohesive low-mass DM detection program must include detection strategies for DM-nucleon scattering as well as DM-electron scattering. Many interesting models feature only DM-nucleon interactions, and these would be invisible to an experiment that is only sensitive to DM-electron interactions. It is especially desirable to identify single systems that can probe nuclear and electronic interactions simultaneously. A large class of electron-recoil detectors have been proposed and designed with the specific aim of detecting DM interactions with electrons themselves. This is sensible from a historical perspective since their reach into low DM masses was originally driven by kinematics: when the DM mass is below the GeV scale and the relevant process for detection is two-body elastic scattering, interactions with electrons can transfer energy much more efficiently than interactions with heavier nuclei.

But as the DM mass decreases further, the interaction rate becomes dominated by many-body processes rather than two-body elastic scattering. In this regime, detectors are primarily sensitive to the excitation of collective modes, such as the production of quasiparticles and phonons, rather than the recoil of individual electrons and nuclei. Here, a

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central challenge for DM searches is to detect extremely low energy deposits, regardless of the mode that sourced them. This is exactly where recent generations of electron-recoil detectors truly excel. Moreover, since the different classes of collective excitations are coupled to one another and to single-particle excitations, a system that specializes in detecting excitations among the electrons should have automatic sensitivity to excitations among the nuclei, and vice versa.

Accordingly, in this Letter, rather than asking whether a nuclear-recoil detector can automatically detect low-energy electronic recoils, as was done in the seminal work of Ref. [6], we ask whether an electron-recoil detector can automatically detect low-energy nuclear recoils. We show that experiments that were designed to detect DM-electron interactions can in fact already probe DM-nucleon interactions at low DM masses. Using existing data from superconducting detectors and focusing on interactions with a scalar mediator, we place new limits on DM-nucleon scattering in parameter space not otherwise constrained by elastic scattering experiments. We thus demonstrate the potential for low-threshold DM experiments to double their science output.

Nuclear scattering events in electronic detectors—For the purposes of this Letter, we specialize to superconducting detectors, although similar statements hold for other classes of low-threshold electron-recoil detectors. Superconducting detectors have been used as readout sensors in a variety of DM experiments. Here, we are interested in the experiments probing the lowest DM masses, where the DM interaction takes place inside the superconductor. Such detectors are ultimately sensitive not to single electron recoils but to changes in the properties of the superconducting phase in the detector. For example, this might be an increase in the phonon density, as in a transition-edge sensor (TES) [26–28]; an increase in the quasiparticle density, as in a kinetic inductance detector (KID) [29]; or a complete transition of a portion of the detector from the superconducting phase to the normal metal phase, as in a superconducting nanowire single photon detector (SNSPD) [11,13].

In all searches for DM interactions within superconductors to date, only DM-electron interactions have been considered as possible triggers of such phase changes, but each of these effects can also be caused by DM-nucleon interactions. Low-energy DM-nucleon scattering events lead directly to the production of phonons. Phonons can, in turn, down-convert by breaking Cooper pairs and producing free quasiparticles. If the initial deposit is large enough, the excess phonons and quasiparticles can prompt a transition to the normal metal state, just as with electronic recoils. Each of the detectable phenomena in a superconducting detector is thus sensitive to DM-nucleon scattering.

In general, the dependence of the event rate on the properties of the target system is determined by the dynamic structure factor $S(\mathbf{q}, \omega)$, which measures the response of the

system to a deposited momentum $\mathbf{q}$ and energy ω . Working in natural units with $c = \hbar = 1$, the DM scattering rate then takes the form

$$R = \frac{\pi \bar{\sigma}_n \rho_\chi}{\mu_{\chi n}^2 \rho_T m_\chi} \int \frac{d^3 \mathbf{q}}{(2\pi)^3} d^3 \mathbf{v} d\omega [f_\chi(\mathbf{v}) \mathcal{F}_{\text{med}}^2(\mathbf{q}) \times S(\mathbf{q}, \omega) \delta(\omega - \omega_{\mathbf{q}})], \quad (1)$$

where $\rho_\chi = 0.4 \text{ GeV/cm}^3$ is the DM density; ρ_T is the target density; $\mathcal{F}_{\text{med}}^2(\mathbf{q})$ is the square of the form factor determined by the structure of the interaction, given by 1 for a heavy mediator and $(q_0/q)^2$ for a light mediator, where $q_0 = m_\chi v_0$ for v_0 the mean DM velocity and $q \equiv |\mathbf{q}|$; $f_\chi(\mathbf{v})$ is the DM velocity distribution; and $\omega_{\mathbf{q}} \equiv \mathbf{v} \cdot \mathbf{q} - \mathbf{q}^2/2m_\chi$ is the energy deposited. Given a matrix element $\mathcal{M}_{\chi n}(q) = \mathcal{M}_0 \mathcal{F}_{\text{med}}(q)$, we define a reference cross section $\bar{\sigma}_n = (\mu_{\chi n}^2/\pi) |\mathcal{M}_{\chi n}(q_0)|^2$, where $\mu_{\chi n}$ is the reduced mass. We take f_χ to be a shifted Maxwell-Boltzmann distribution with $v_0 = 230 \text{ km/s}$, Earth velocity $v_e = 240 \text{ km/s}$, and an escape velocity $v_{\text{esc}} = 600 \text{ km/s}$.

As we are considering multiple interaction channels for the DM, it is important to consider the differential sensitivity of the experimental system to each channel. In particular, TESs [26–28] and SNSPDs [11,13] are nominally sensitive to heat, while KIDs [29] are nominally sensitive to quasiparticles in the final state. Fortunately, for energy deposits ω in a superconductor that are much larger than the superconducting gap 2Δ , the partitioning of energy between phonons and quasiparticles in the final state is independent of ω and q [12,30]. The detector threshold is typically derived from calibration experiments using photon absorption, and, for all of the detectors we consider in this Letter, these calibration experiments are performed with $\omega \gg 2\Delta$. Thus, the determination of the threshold is subject to the same partitioning of energy that occurs for an electronic or nuclear scattering event. For example, in aluminum, an energetic event transferring more than $\sim 5\Delta$ in the superconductor deposits about 60% of this energy into ballistic quasiparticles, and the rest into phonons [12,31]. This means that there is no need to treat DM interactions differently from photon absorption for the purposes of determining the threshold: since both processes are well above the superconducting gap, the final state is insensitive to the channel and depends only on the total energy.

Determining the sensitivity to DM-nucleon interactions is now reduced to the computation of the dynamic structure factor for each interaction process of interest. Here, the processes we consider are elastic nuclear recoils, single phonon production, and multiple phonon production. We now calculate $S(\mathbf{q}, \omega)$ for each of these channels.

Elastic nuclear recoils: In general, the interaction between a DM particle and a nucleon in the target material is complicated by the response of the nuclear lattice.

However, if the deposited energy ω is much larger than the maximal phonon energy $\omega_{\text{ph}}^{\text{max}}$, the DM-nucleon interaction is decoupled from the lattice structure of the target. If the target is composed of several types of atoms, then the dynamic structure factor for a hadrophilic interaction is given by a weighted average of the contributions of each atom as [32]

$$S(\mathbf{q}, \omega) = \frac{2\pi\rho_T}{\sum_N A_N} \sum_N \frac{A_N^3}{m_N} F_N^2(\mathbf{q}) \delta\left(\omega - \frac{\mathbf{q}^2}{2m_N}\right). \quad (2)$$

Here, N indexes the nuclei in a unit cell, m_N is the atomic mass, $A_N = m_N/u$ is the mass number, and $F_N(\mathbf{q})$ is the nuclear form factor. Throughout, we use the Helm form factor [33], $F_N(q) = [3j_1(qr_N)/(qr_N)]e^{-(qs)^2/2}$, where j_1 is the spherical Bessel function; $r_N \approx A_N^{1/3} \times 1.14$ fm is the nuclear radius; and s is the nuclear skin thickness. We use $\{A_W, A_{Si}, A_{Ti}, A_N, A_{Al}\} \approx \{183.85, 28.09, 47.87, 14.01, 26.98\}$, and we set $s = 0.9$ fm.

Single phonons: On the other hand, for the smallest deposits, $\omega \ll \omega_{\text{ph}}^{\text{max}}$, the relevant excitations are single phonons. The dynamic structure factor for the production of a single phonon is given by [32]

$$S(\mathbf{q}, \omega) = \frac{\pi}{\Omega} \sum_{\nu} \frac{\delta(\omega - \omega_{\nu, \mathbf{k}})}{\omega_{\nu, \mathbf{k}}} \left| \sum_j \frac{e^{-W_j(\mathbf{q})}}{\sqrt{m_j}} e^{i\mathbf{G} \cdot \mathbf{x}_j^0} \times A_{N_j} F_{N_j}(q) \mathbf{q} \cdot \boldsymbol{\epsilon}_{\nu, \mathbf{k}, j}^* \right|^2. \quad (3)$$

Here, Ω is the volume of the primitive cell; ν indexes the phonon branches and j indexes atoms in the primitive cell; $\mathbf{q} = \mathbf{k} + \mathbf{G}$, with $\mathbf{k}$ lying in the first Brillouin zone and with $\mathbf{G}$ a reciprocal lattice vector; $\omega_{\nu, \mathbf{k}}$ is the phonon frequency; m_j is the atomic mass; $\boldsymbol{\epsilon}_{\nu, \mathbf{k}, j}$ is the polarization vector; $F_{N_j}(q)$ is the nuclear form factor; and $W_j(\mathbf{q})$ is the Debye-Waller factor, which describes the physics of the creation and annihilation of phonons from the vacuum. The latter is given by

$$W_j(\mathbf{q}) = \frac{1}{4N_L m_j} \sum_{\nu} \sum_{\mathbf{k} \in \text{1BZ}} \frac{|\mathbf{q} \cdot \boldsymbol{\epsilon}_{\nu, \mathbf{k}, j}|^2}{\omega_{\nu, \mathbf{k}}}, \quad (4)$$

where N_L is the number of primitive unit cells in the lattice, which give rise to N_L discrete values of $\mathbf{k}$. In practice, we take the continuum limit $N_L \rightarrow \infty$. The Debye-Waller factor suppresses the contribution of higher momentum transfers to the phonon production rate.

Multiple phonons: For intermediate deposits which are neither much smaller nor much larger than $\omega_{\text{ph}}^{\text{max}}$, an interaction can produce several phonons in the material. The analysis in this case is more involved. In this Letter, we use the incoherent approximation as described in Ref. [34], where the dynamic structure factor is written as

$$S(\mathbf{q}, \omega) \approx \frac{2\pi}{\Omega} \sum_j A_{N_j}^2 e^{-2W_j(\mathbf{q})} \sum_n \frac{1}{n!} \left(\frac{q^2}{2m_j} \right)^n \times \left(\prod_{i=1}^n \int d\omega_i \frac{D_j(\omega_i)}{\omega_i} \right) \delta\left(\sum_i \omega_i - \omega\right), \quad (5)$$

where the n th term is the contribution of the n -phonon final state and we sum the terms up to $n = 10$. Here, D_j is the partial density of states for the j th atom in the unit cell given by

$$D_j(\omega) = \frac{1}{3N_L} \sum_{\nu} \sum_{\mathbf{k} \in \text{1BZ}} |\boldsymbol{\epsilon}_{\nu, \mathbf{k}, j}|^2 \delta(\omega - \omega_{\nu, \mathbf{k}}). \quad (6)$$

In Eq. (5), W_j denotes the partial Debye-Waller factor, which is given by $W_j(\mathbf{q}) = (q^2/4m_j) \int d\omega' D_j(\omega')/\omega'$.

We note that the incoherent approximation is not expected to be valid for momentum transfers smaller than the size of the Brillouin zone, q_{BZ} . The length scale corresponding to $q < q_{\text{BZ}}$ extends over more than one unit cell, so coherent contributions from atoms in different unit cells become important. In this region, the scattering rate is again dominated by single phonon excitations. We thus omit the $n = 1$ term in the sum of Eq. (5) and add in its place the accurate single phonon contribution of Eq. (3). At fixed ω , the sum of contributions from n -phonon final states in Eq. (5) converges more slowly as q grows due to the factor of $q^2/2m_j$. For this reason, we follow Ref. [34] and use the impulse approximation for the dynamic structure factor when $q > 2\sqrt{2m_j\bar{\omega}_j}$, where $\bar{\omega}_j$ is the mean frequency weighted by the partial density of states D_j . The dynamic structure factor in the impulse approximation is then given by

$$S(\mathbf{q}, \omega) = \frac{2\pi}{\Omega} \sum_j \frac{A_{N_j}^2}{\sqrt{2\pi\Delta_j^2}} \exp\left(-\frac{(\omega - q^2/2m_j)^2}{2\Delta_j^2}\right), \quad (7)$$

where the Gaussian width is $\Delta_j^2 = q^2\bar{\omega}_j/(2m_j)$.

Results—We have now set the stage to compute interaction rates in various targets and to determine the sensitivity of real experiments to DM parameter space. We use existing experimental data collected from two prototype superconducting detectors operated above ground: an SNSPD [11,13] and a KID [29]. (The experimental setups are summarized in the End Matter.) These two detectors have already demonstrated extraordinarily low energy thresholds and dark count rates. Moreover, each of these devices has a dynamic range wider than the range of DM masses that we constrain, so we need not impose a maximum energy transfer. We derive constraints by comparing the measured count rate in these experiments to the event rate implied by Eq. (1), with $S(\mathbf{q}, \omega)$ determined by Eqs. (2), (3), (5), and (7) in the appropriate regimes.

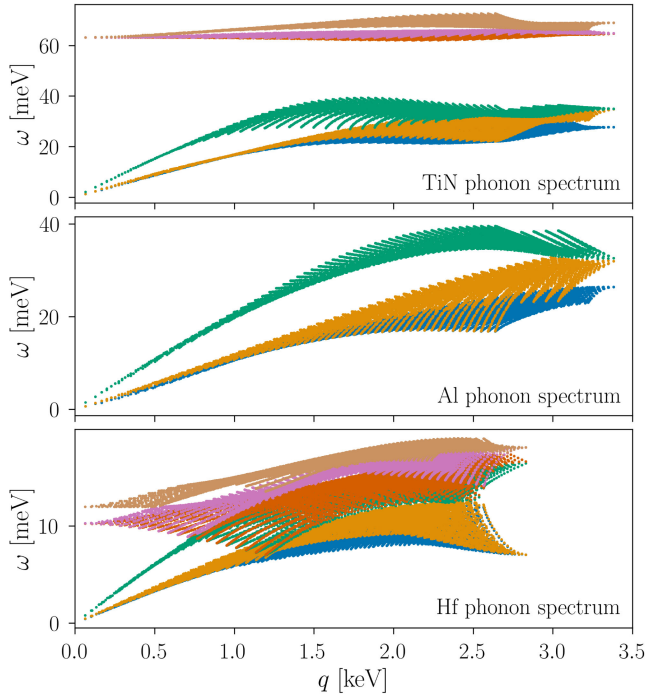


FIG. 1. Phonon spectra. The computed spectrum of phonons in TiN (top panel), Al (middle panel), and Hf (bottom panel), with bands differentiated by color. TiN and Hf exhibit both acoustic phonons, which have a linear dispersion at small q , and optical phonons, which have a nonzero gap with a nearly flat dispersion. The spectra are computed in the normal metal phase, as corrections from superconductivity are negligible at the deposits we consider.

The SNSPD [11,13] considered here is composed of WSi with a mass of 4.3 ng and a threshold of 0.73 eV. A count rate of $6 \times 10^{-6} \text{ s}^{-1}$ was measured over 180 h. The threshold is large compared to the scale of single phonon energies in the material, so we consider only elastic nuclear recoils. Hence, our new bound is based on the interaction rate of a DM particle with the W and Si target nuclei within this prototype detector. The KID experiment, which is fully described in Ref. [29], used a detector composed of TiN with a volume of $2 \times 10^{-12} \text{ cm}^3$. A count rate of $2 \times 10^{-3} \text{ s}^{-1}$ was measured over 26 h.

We obtain the phonon spectrum in TiN via the finite displacement method using PHONOPY [35,36] with density functional theory (DFT) calculations carried out in VASP [37–40] (see the End Matter for details). The resulting population of phonon modes is shown in the top panel of Fig. 1. The effective threshold of the device is as low as $\sim 200 \text{ meV}$, whereas $\omega_{\text{ph}}^{\text{max}} \approx 70 \text{ meV}$. (The threshold of the KID is determined by the characteristics of the noise in the device at low energies and is subject to modeling assumptions. See Ref. [29] for details.) Our new constraints receive contributions from multiphonon production and lie mainly in the many-phonon impulse regime.

Figure 2 summarizes our new limits on the DM-nucleon cross section placed by the SNSPD (shaded green areas) and KID data (shaded magenta areas), as well as projections for the sensitivity of future experiments with lower thresholds and larger exposures. Our bounds are set at the 95% confidence level for scattering of DM via a light (left panel of Fig. 2) or heavy mediator (right panel of Fig. 2), and we incorporate the measured count rates using the Feldman-Cousins procedure [41]. In principle, the KID constraints continue to lower masses and fall off only as a power law for $1 \text{ meV} \lesssim m_\chi \lesssim 20 \text{ meV}$, as multiphonon excitations remain kinematically available below the threshold for elastic nuclear recoils. However, we do not show constraints on cross sections above 10^{-25} cm^2 due to significant effects from atmospheric scattering en route to the detector. A full treatment of overburden is complicated, particularly for light mediators. The cross sections we show are not typically subject to significant atmospheric scattering in surface experiments at comparable DM masses [42], but our upper limit of 10^{-25} cm^2 is approximate, and constraints should be interpreted with caution near the top of the figures. For comparison, the shaded gray regions of Fig. 2 depict the best existing constraints from terrestrial experiments [43–53]. The darker gray region in the right panel of Fig. 2 indicates direct constraints from DM-nucleon scattering, while the lighter shading indicates constraints inferred from electronic excitations via the Migdal effect [54]. Model-dependent cosmological bounds may also apply in portions of this plane. (See Ref. [55] for an overview.)

We also consider the prospects for future experiments. Here, we evaluate Hf, which has proven useful in recent KID detectors [57], and Al, which is a common benchmark in this parameter space [12,58]. The blue and orange curves in Fig. 2 show the projected reach for future experiments consisting of Al and Hf targets, respectively, with thresholds of 10, 30, and 50 meV through multiple phonon excitations in the target. The phonon populations we use for Al and Hf are shown in the middle and bottom panels of Fig. 1. These projections are made for a 1 kg yr exposure, assuming no background events at the 95% confidence level, for comparison with the previous literature. For comparison, we show in dotted gray the expected reach of a proposed experiment utilizing superfluid He [56] with a 1 meV threshold and a 1 kg yr exposure. The pronounced features in the low-threshold Al and Hf projections arise from single phonon excitations: this channel is available only for thresholds $\omega_{\text{min}} < \omega_{\text{ph}}^{\text{max}}$. The acoustic phonons can only be produced by DM scattering for $m_\chi \gtrsim 180 \text{ keV} \times (10^{-5}/c_s) \times (\omega_{\text{min}}/10 \text{ meV})$, where $c_{s,\text{Al}} \approx 2 \times 10^{-5}$ and $c_{s,\text{Hf}} \approx 1 \times 10^{-5}$. The Hf spectrum includes optical phonons with a gap $\omega_O \approx 10 \text{ meV}$, which can be excited by $m_\chi \gtrsim 3 \text{ keV} \times (\omega_O/10 \text{ meV})$. For small DM masses, single phonons dominate the rate when kinematically allowed.

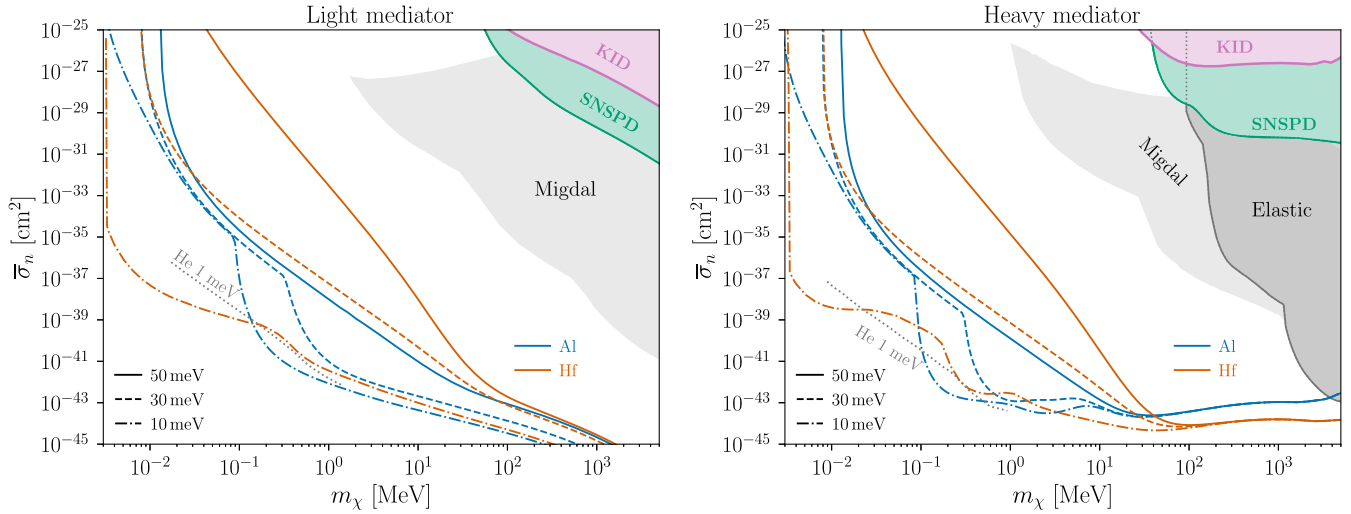


FIG. 2. Dark-matter scattering. Results for hadrophilic dark-matter scattering with nucleons mediated by a light (left panel) or heavy (right panel) scalar mediator at the 95% confidence level. The green and magenta shaded regions indicate the new bounds derived in this Letter using the SNSPD and KID data of Refs. [11,13,29]. The shaded dark and light gray regions indicate the strongest existing terrestrial constraints to date based on elastic interactions and the Migdal effect, respectively. These include constraints from EDELWEISS [43,44], DarkSide-50 [45,46], SuperCDMS [47,48], CRESST-III [49], SENSEI [50], PandaX-4T [51], LUX [52], and XENON1T [53] data. Constraints may be relaxed for some models due to atmospheric scattering, particularly for cross sections above 10^{-26} cm². The blue and orange curves show the projected sensitivity of experiments with superconducting Al and Hf targets, respectively, via excitations of single or multiple phonons in the target, for a kg-yr exposure. Solid, dashed, and dot-dashed curves indicate thresholds of 50, 30, and 10 meV, respectively. The dotted gray curve indicates the projected reach of a superfluid He experiment [56] with a threshold of 1 meV and kg-yr exposure.

Discussion—We have shown how detectors geared at probing DM-electron interactions can be simultaneously sensitive to DM-nucleon interactions. This is possible due to the electron-phonon coupling in materials, which allows deposits in one type of degrees of freedom to be transferred to the other. (The use of a coupling between different channels is reminiscent of the Migdal effect at higher masses [53,54,59,60], in which a nuclear coupling is used to produce an electronic ionization signal.) We have demonstrated the power of this approach using published data from superconducting detectors—previously used to set world-leading limits on DM scattering with electrons—to place new constraints on DM-nucleon interactions. We have computed the reach of future Al and Hf targets, laying the groundwork to double the science that will be extracted from existing and future detectors.

This has particularly significant implications for the future of superconducting detectors. Such devices were originally proposed as a light DM scattering target almost ten years ago by Ref. [9]. At that time, they were introduced strictly as electron-recoil detectors, intended to be complementary to other technologies probing nuclear recoils. Our results from prototype experiments illustrate that there is no need for a distinction between electron-recoil and nuclear-recoil experiments: the same systems probe both interactions. Several existing and planned low-threshold electron-recoil experiments are simultaneously capable of probing DM-nucleon couplings.

We stress that the prototype superconducting targets we used to place new limits on DM parameter space in this Letter are significantly smaller than other existing detectors. Nonetheless, the low thresholds of these detectors already enable us to probe DM masses lower than any previously probed by elastic DM-nucleon scattering. Efforts are already underway to achieve low thresholds with larger target masses [61,62]. With the future scaling of such experiments, already motivated by their sensitivity to DM-electron interactions, superconducting detectors promise to probe deep into uncharted parameter space for DM-nucleon interactions as well.

Our results from existing experiments are sensitive mostly to parameter space that is nominally probed by larger, higher-threshold semiconductor and time projection chamber experiments using the Migdal effect. In this Letter, we have focused on the relatively simple processes that arise directly from DM-nucleon interactions: single phonon production, multiphonon production, and elastic nuclear recoils. As Fig. 2 demonstrates, even these processes will eventually probe much lower masses than is possible with the Migdal effect. These are, however, not the only processes that can lead to detectable events. (See, for example, Ref. [63], in which hadronic loops dominate for scattering via a vector mediator.) These detectors are also sensitive to the production of prompt quasiparticle pairs via an off-shell phonon. Quasiparticle pairs can offer better kinematical matching to low-mass DM, potentially

extending the reach of superconducting detectors further into unconstrained parameter space. This will be the subject of future work [64].

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Data availability—The data that support the findings of this Letter are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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End Matter

Here, we provide additional details regarding the implementation of the experiments whose data was used in this Letter, as well as the DFT computations performed to obtain phonon spectra.

Experimental details—The details of the SNSPD experiment considered here are fully described in Refs. [11,13]. The detector is composed of a wire of WSi 140 nm in width with a mass of 4.3 ng, deposited in a 7 nm layer atop a SiO₂ substrate. The device was operated with a bias current of 4.5 μ A at a temperature of 0.3 K. Deposited energy prompting a transition to the normal metal phase increases the resistance of the device, leading to a voltage pulse. Under these conditions, the SNSPD was determined to have high detection efficiency for photons with energy of at least 0.73 eV. Sensitivity to photons at lower energies cannot be ruled out, but we conservatively neglect that

possibility in this Letter. Over the 180-h science run, four counts were observed, corresponding to a dark count rate of $6 \times 10^{-6} \text{ s}^{-1}$. This SNSPD is a count detector with very low energy resolving power. Thus, no spectral information can be used for the DM constraint. Note that when this device was used to set limits on electronic scattering by Ref. [13], the effects of the thin-layer geometry were carefully evaluated and found to be negligible. The typical momentum transfers we consider in this Letter are larger than those encountered in electronic scattering, so thin-layer corrections are even less significant.

The KID experiment and the construction of the detector is described at length in Ref. [29]. The experiment uses a $2 \times 10^{-12} \text{ cm}^3$ detector composed of TiN, with a thickness of 22 nm. The superconducting TiN sample has a gap of 0.18 meV, corresponding to a critical temperature of 1.2 K. The superconductor is placed in series with a capacitor and

an inductor, forming an RLC resonator with a sharply defined resonance frequency $\bar{f}$. Energy deposited in the superconductor increases the quasiparticle density, leading to an increase in the kinetic inductance and a transient shift in the resonant frequency. The sensor is thus read out via the frequency quadrature $\Delta f(t)$, which is matched to a template pulse via an optimal filter. The template pulse shape is determined from the same experimental configuration with 0.8 eV calibration photons. This device has good energy resolution compared to SNSPDs, with the resolving power $R \approx 2.4$ at $\omega = 0.8$ eV, and the energy distribution of the dark counts was measured using the procedure detailed in Ref. [29]. The dark count rate is thus strongly energy dependent. We use this distribution in order to place constraints in the manner described in Ref. [29]. Note that a number of new KID-based detectors are already being constructed using Hf targets (see, e.g., Ref. [57]), so we show projections for Hf in Fig. 2 in addition to our TiN results.

For each of the devices considered here, the superconducting layer of the detector is in direct contact with a substrate with a substantially larger volume. It is possible that DM scattering in the substrate will produce an excitation in the superconductor, leading to a detection event, particularly if the deposited energy is large. However, since our interest lies in constraining the lowest DM masses with the smallest energy deposits and since the efficiency of detection for events in the substrate is difficult to reliably compute, we neglect all DM interactions in the substrate and place conservative constraints. Future study of excitations in the substrate, along with calibration of the efficiency for substrate events, might lead to substantially stronger limits at masses that are well above threshold.

There are promising prospects for improving and scaling experiments of the type used in this Letter: for example, SNSPDs have already been multiplexed into arrays with

hundreds of thousands of sensors [65], and individual sensors have already demonstrated thresholds of just tens of meV [66]. Thus, the technology already exists to implement an experiment that is intermediate between our current results and the more ambitious goal of a kg-yr low-threshold exposure.

As experiments scale in exposure, background count rates may well become a limiting factor. More than simple counts will be needed in order to claim a DM discovery. Electron-recoil searches in crystalline detectors can very easily accommodate directional sensitivity due to anisotropy in the excitation spectrum [67]. This is an essential tool for background rejection [68] since the velocity distribution of DM in the laboratory is anisotropic. This may also be viable for nuclear-recoil searches, as may the approach described in Ref. [12] for electron recoils in superconductors. We defer the consideration of directional sensitivity to future work [69].

Density functional theory calculations—We compute the phonon spectrum in TiN using the Perdew-Burke-Ernzerhof exchange-correlation functional [70] and projector augmented wave pseudopotentials [71,72] including 12 and 5 electrons as valences for Ti and N, respectively. We use a plane-wave energy cutoff of 800 eV, Γ -centered k -point grids with k -point spacing less than $0.17 \text{ \AA}^{-1}$ and convergence criteria for an electronic self-consistent loop set to 10^{-8} eV. Structure optimization using a maximum force criterion of $10^{-3} \text{ eV \AA}^{-1}$ yielded the lattice constant $4.249 \text{ \AA}$. We calculate the phonon spectrum with a 216 atom supercell on a $40 \times 40 \times 40$ q -point mesh. For Al, we use a 256 atom supercell and three electrons as the valence, with all other calculation settings the same as for TiN, resulting in a lattice constant of $4.038 \text{ \AA}$. For Hf, we use a 150 atom supercell and ten electrons as the valence, resulting in lattice constants $\{a, c\} = \{3.202, 5.055\} \text{ \AA}$.