A Path to the Direct Detection of sub-GeV Dark Matter Using Calorimetric Readout of a Superfluid $^4$He Target

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A promising technology concept for sub-GeV dark matter detection is described, in which low-temperature microcalorimeters serve as the sensors and superfluid $^4$He serves as the target material. A superfluid helium target has several advantageous properties, including a light nuclear mass for better kinematic matching with light dark matter particles, copious production of scintillation light, extremely good intrinsic radiopurity, a high impedance to external vibration noise, and a unique mechanism for observing phonon-like modes via liberation of $^4$He atoms into a vacuum (‘quantum evaporation’). In this concept, both scintillation photons and triplet exciters are detected using calorimeters, including calorimeters immersed in the superfluid. Kinetic excitations of the superfluid medium (rotons and phonons) are detected using quantum evaporation and subsequent atomic adsorption onto a microcalorimeter suspended in vacuum above the target helium. The energy of adsorption amplifies the phonon/roton signal before calorimetric sensing, producing a gain mechanism that can reduce the technology’s recoil energy threshold below the calorimeter energy threshold. We describe signal production and signal sensing probabilities, and estimate electron recoil discrimination. We then simulate radioactive backgrounds from gamma rays and neutrons. Dark matter - nucleon elastic scattering cross-section sensitivities are projected, demonstrating that even very small (sub-kg) target masses can probe wide regions of as-yet untested dark matter parameter space.

I. INTRODUCTION

In the $\Lambda$CDM model of cosmology, dark matter makes up 26.8% of the mass-energy density of the universe [1]. Gravitational effects of this dark matter are evident at many distance and time scales, and the dark matter strongly affects the evolution of the universe. The nature of this dark matter is a mystery, and resolution of this mystery would have a profound impact on the fields of astrophysics, cosmology, and particle physics. In particular, the existence of dark matter is strong evidence for physics beyond the Standard Model, and measurement of the mass of the dark matter particle and its interaction modes with ordinary matter would open new vistas in particle physics.

For the past several decades, experimental efforts to directly detect such dark matter interactions have focused on axions [2] and Weakly Interacting Massive Particles (WIMPs) [3]. In the latter case, the dark matter particle must have a mass above the Lee-Weinberg limit [4] of $\sim 2$ GeV; a lower WIMP mass results in an annihilation cross-section that is too small and produces too much dark matter in the early universe. However, if a new force carrier exists, then dark matter particles which interact through this new mediator are viable below the Lee-Weinberg scale. Models of sub-GeV dark matter include freeze-out dark matter [5–12], asymmetric dark matter [13, 14], and freeze-in dark matter [15]. A summary of the physics motivation for sub-GeV dark matter may be found in the recent reviews [16, 17].

Direct detection of sub-GeV dark matter through nuclear recoils is a particularly difficult challenge because the transfer of kinetic energy from the dark matter particle is very inefficient if its mass is much less than that of the target nucleus (generically true for the sub-GeV case). Some approaches designed to avoid this limitation are described in [18]. One can also consider dark matter scattering with electrons, and there are several experimental approaches in development [16, 17]. Methods of detecting sub-GeV dark matter interactions with electrons include use of charge signal in noble liquid experiments [19–21], charge-coupled devices [22], and electron-hole pair detection in semiconductors employing Luke-Neganov gain [20]. Many dark matter models do not predict much dark matter interaction with electrons, so experiments designed to detect dark matter-nucleus interactions must be included in a broad experimental program and are naturally complementary to those searching for dark matter-electron interactions. Two recently proposed approaches are to detect color centers [23] or spin avalanches [24] produced through dark matter-nucleus interactions. Light dark matter may also be detected through its coupling to optical phonons in polar materials [25].

Superfluid $^4$He has been previously considered for WIMP detection in [26] as part of the HERON project [27, 28], and has recently gained attention in the context of low-mass dark matter detection [29–31]. Advantages of superfluid $^4$He include a) Low target...
mass, allowing relatively good kinematic matching to low-mass dark matter particles; b) Multiple observable and distinguishable signal channels summing to the total recoil energy, including phonons and rotons (commonly referred to collectively as ‘quasiparticles’), substantial scintillation light, and triplet helium excimers; c) Inhibited vibrational coupling of the target mass to the environment (the container walls), due to the distinct superfluid phonon/roton dispersion relation; d) High radiopurity, as helium has no long-lived isotopes, may be purified using getters or cold traps, and impurities freeze out of the bulk; d) A large band gap energy of 19.77 eV (the energy needed to excite atomic helium to an n = 2 state), inhibiting all electronic excitation backgrounds below this energy; e) Quasiparticle excitations which are long-lived and ballistic, thereby preserving recoil information encoded in their production; and f) A liquid state down to zero K, enabling mK-temperature calorimetric readout of an easily-scalable liquid target mass. Superfluid $^4$He is being used as an ultracold neutron production, storage, and detection material for measurements of the neutron lifetime [32] and the neutron electric dipole moment [33]. Superfluid $^3$He has also been proposed as a dark matter detection material, using oscillating wires immersed the superfluid to detect dark matter particles [34, 35]. Here we elaborate on the possibility of using superfluid $^4$He for direct detection of sub-GeV mass dark matter particles, with an approach relying entirely on calorimetric measurement of multiple signal carriers: scintillation light, triplet excimers, and quasiparticles. This method offers the possibility of discrimination between dark matter particles and backgrounds using the ratios of different signal channels, while enabling extremely low energy threshold.

II. DETECTOR LAYOUT

A general detector geometry is described in Figure 1. The $^4$He target mass is contained within a passive surrounding vessel. In the vacuum above the liquid surface a large-area calorimeter is suspended, serving as the primary detector for quasiparticles, via quantum evaporation of $^4$He atoms into the vacuum. Other large-area calorimeters are immersed within the target liquid, approximately covering the vessel surface and providing nearly complete area coverage.

Interfaces between superfluid $^4$He and solid materials exhibit an exceptionally large Kapitza resistance, which would inhibit the transmission of quasiparticle states into the vessel or immersed calorimeters. After multiple internal reflections, a dominant fraction of quasiparticle energy can escape the liquid as atomic evaporation. The suspended calorimeter senses the arrival of these evaporated atoms, with the dominant energy per atom being the adhesion energy of the atom to the calorimeter surface. This ‘adhesion gain’ requires a ‘dry’ calorimeter, free of the $^4$He that typically coats all available surfaces at these temperatures. Various technologies have been demonstrated which can prevent film flow to the suspended calorimeter, including a film burner as used in the HERON project [36], a knife edge of atomic sharpness [37–39], and a clean surface of non-wettable material such as Rb or Cs [40, 41].

The primary role of the immersed calorimetry, on the other hand, is to detect signal carriers resulting from atomic excitation, each of higher (eV-scale) energies. The immersed calorimetry may require slight modification to adapt for the immersed environment. The superfluid’s high Kapitza resistance mitigates but does not completely eliminate the leakage of phonon energy out of the calorimeter. A straightforward solution to the leakage could be to cover a large fraction of the calorimeter surface with Al, in which incident energy is efficiently converted to Bogoliubov quasiparticles (‘broken cooper pairs’), which should exhibit negligible leakage.

A secondary concept might also be of similar interest, in which the vessel is transparent to optical photons, the inside vessel surface is coated in a wavelength shifter such as tetraphenyl butadiene (TPB), and the previously-immersed calorimetry is instead suspended in the vacuum outside the transparent vessel. The relative merits of the two layouts are not yet known quantitatively.

The microcalorimeters (in both primary and secondary concepts) pair a microscopic energy sensor with a large-area (few-cm scale) thin (<mm) absorber. While other sensor technologies may also provide the necessary sensitivity, the sensor technology which has received the most study in this large area application is the Transition Edge Sensor (TES). Pyle et al. [42] point out that historically the timescale of energy diffusion from absorber to TES has been severely mismatched with the TES response timescales, leading to significant

![FIG. 1. Simplified detector layout. Here, superfluid $^4$He is blue, large-area microcalorimetry is red, and the passive containing vessel is grey.](image-url)
and avoidable degradation in threshold and resolution. With this and other recent conceptual advances, there has been significant recent laboratory progress towards larger areas and lower thresholds. A recent R&D device has demonstrated a 3.5 eV baseline resolution ($\sigma$) on a 45.6 cm$^2$ collection area [43]. There are no obstacles expected in further refinement, we expect the thresholds in coming years to advance into the sub-eV regime, while retaining areas. Such sensors are capable of counting individual eV-scale deposits (for example, from scintillation photons) with no relevant dark count rate.

III. ENERGY PARTITIONING IN SUPERFLUID $^4$He

The energy of a particle recoil in liquid $^4$He is partitioned among several channels: ionization, electronic excitation, and quasiparticle excitations (phonons and rotons).

At low applied fields, geminate recombination converts nearly all ionization into neutral (but excited) atoms. Electronic excitations decay via IR emission to either a singlet or triplet state of the first excited state. Such atomic excited states appear in the liquid as dimer excimers, the singlet as $A^1\Sigma_u^+$ and the triplet as $a^3\Sigma_u^+$.

The singlet excimer decays on a ns-scale via UV photon emission ($\sim$16 eV). The triplet excimer decays with the exceptionally long half life of 13 s [44]. Despite the triplet excimer’s long lifetime, it typically releases its energy on shorter timescales after propagation to material boundaries. This triplet excimer propagation is ballistic thanks to the low density of phonon excitations in the superfluid medium (assuming a temperature <100 mK), and the ballistic velocity has been measured to be $O(m/s)$ [45], with variation dependent on the phonon environment in which the recombination process occurs. The quenching of single-atom He triplet states on material surfaces is now a standard technique for probing surface electron states in a vacuum environment [46–48]. In this atomic case, the triplet quenching process is observed to occur via charge exchange between the excited atom and the surface, injectig some fraction of excitation energy into the surface (as excited electrons), and ejeclting some fraction away from the surface (as Auger electrons or x-rays). Similar triplet quenching processes occur in the excimer case and in the superfluid environment. Observation of the injected energy via microcalorimetry has recently been demonstrated [49].

To estimate the energy partitioning fractions for nuclear recoils, we first use the Lindhard nuclear and electronic energy partition [50]

$$E = \nu + \eta, \quad (1)$$

where $\nu$ is energy deposited via nuclear stopping and $\eta$ is the energy deposited by electronic stopping through ionizations and excitations of helium atoms. The Lindhard $\nu$ is given by

$$\nu(\epsilon) = \frac{\epsilon}{1 + k g}, \quad (2)$$

where $\epsilon = 11.5 E / Z^{7/3}$ is a reduced energy (with $E$ in keV and $Z$ being the atomic number), $k = 0.133 Z^{2/3} A^{-1/2}$ (with $A$ being the atomic mass) and $g$ is well-approximated by $g = 3e^{-0.15} + 0.6e^{0.7} + \epsilon$ [51]. The relative fraction of energy appearing in each Lindhard channel is the ratio of $\nu$ to $\eta$, and we assume the fraction of energy deposited in the detector through nuclear stopping can be detected as quasiparticle excitations. The ionization and excitation production ratio, which has been described and modeled for helium-helium collisions by Guo & McKinsey [29] and Ito & Seidel [30], can be derived from measured cross sections. As in Guo & McKinsey, we have neglected secondary electron effects in our calculation of the electronic stopping power since we are concerned with recoils of energies well below 100 keV.

Following Ito & Seidel, we assume a recombination ratio of singlets to triplets of 1 to 3 for nuclear recoils, and we estimate the excitation ratio of singlets to triplets to be 0.86:0.14 with a total cross section for excitations equal to 1.4 times the $2^1P$ excitation cross section for helium projectiles. With these assumptions, we estimate the ratio of the ionization, singlet excitation, and triplet
excitation cross sections as a function of recoil energy. Based on this ratio, we compute the energy appearing in each signal channel by assigning an average energy of 15.5 eV to each singlet excitation, 18 eV to each triplet, 4 eV for each ionization and 0.5 eV for each excitation to the IR channel, and we add to the quasiparticle channel 8 eV for the average secondary electron contribution, 2 eV for dimerization and 4 eV for dissociation of ground state excimers [52]. We also estimate the effect of the Penning quenching process on excimers in the recoil region with the model presented by Ito & Seidel, which treats the density of excited atoms at a recoil site by the differential equation

$$\frac{dn}{dt} = -\gamma n^2 - \frac{r n}{\tau},$$

where $\gamma$ is a bimolecular rate taken to be the same for all species, $r = 0.4$ is determined by the fraction of singlet excitations, and $\tau$ is the singlet lifetime. The Penning quenching factor

$$f = \frac{1}{n_0} \int_0^\infty \frac{r n}{\tau} dt = \frac{\ln(1 + \xi)}{\xi},$$

with $\xi = n_0 \gamma \tau / r$, is the fraction of excimers that decay radiatively, while rest of the energy is quenched and appears in the quasiparticle channel. Ito & Seidel fix the bimolecular rate with the calorimetric observation that $f = 0.5$ for 5.5-MeV $\alpha$ particles [30, 53, 54]. Since secondary electrons have a non-negligible effect on the ionization and excitation stopping powers at this energy, we use Ito & Seidel’s track density calculations for a rough estimate of $\gamma = 13$ cm$^{-1}$ s$^{-1}$. Since the ionization and excitation cross section data does not extend below 100 eV, we extrapolate while ensuring that all of the energy of a recoil to goes into the quasiparticle channel below the excitation threshold of 19.77 eV.

For the electron recoil partitioning fractions, we use cross section data for electron-impact ionizations and excitations to the lowest-lying singlet and triplet states in helium [55]. We estimate the ratio of ionizations, singlet excitations, and triplet excitations by the ratio of these cross sections, this time assuming the geminate recombination fraction of 50% singlets, and assign energies for each ionization and excitation as described above for the nuclear recoil partitioning [52]. The average partitioning of energy for nuclear and electron recoils is shown in the upper row of Figure 2.

IV. DETECTOR RESPONSE AND ELECTRON RECOIL DISCRIMINATION

Given a model for energy partitioning, we can simulate statistical distributions of observed excitation counts using a simple Monte Carlo method. The three partition fractions for atomic excitation quanta (singlet, triplet, IR) are taken as in the upper panel of Figure 2. Then, a poissonian and uncorrelated production mechanism is assumed for these three excitation modes, resulting in atomic excitation counts. The energies of the resulting atomic excitations are summed, subtracted from the recoil energy, and the remainder is assumed to appear as quasiparticle excitations (for now, assuming each quasiparticle is of energy 0.8 meV).

Once excitation counts in all four signal channels have been produced, these counts flow through a similarly simple detection Monte Carlo, in which it is assumed that each type of excitation is observed according to a binomial process with observation efficiencies specific to each signal type. Given the primary detector layout described in Figure 1, we assume detection efficiencies of 0.95 for singlet excitations (16 eV photons) and IR photons. We assume a triplet excitation detection efficiency of 5/6, accounting for some uncertainty in the behavior of excimers incident on the liquid-vacuum interface (though we expect the surface facilitates a quenching process which converts their energy into an observable signal). As discussed in subsequent sections, the probability of observing quasiparticle excitations through the evaporation channel is uncertain; for the purposes of Figure 2 we assume a low evaporation efficiency of 0.05 (though this evaporation fraction is largely irrelevant within the illustrated energy range or to the discrimination metric we soon calculate).

For the atomic excitations, the calorimeter observes an amount of energy equal to the excitation energy. Quasiparticle observations, on the other hand, appear not as the quasiparticle energy but as the energy of adsorption of a $^4$He atom on the upper calorimeter. Assuming a quasiparticle energy of 0.8 meV and an atomic binding energy at the liquid surface of 0.62 meV, the kinetic energy of the liberated atom is quite low, only 0.18 meV. Comparing 0.8 meV to the typical adhesion energy of He to solid materials ($\sim$10 eV for most materials, see a compendium in [59]), we see that the adhesion energy dominates the calorimetric measurement, and that it can provide a gain in energy, before sensing, of $>10\times$. This ‘adhesion gain’ factor can be enhanced by coating the calorimeter with a layer, perhaps only atomically thin, of a material of particularly high He adhesion energy. We assume a per-atom adhesion energy of 42.9 meV, appropriate for fluoro-graphene [57]. Given these energies, each meV-scale quasiparticle excitation, if observed via the quantum evaporation process, would appear as an energy deposit in the calorimeter of $E_{\text{dep}} = E_{\text{atom}} + E_{\text{binding(cal)}} = (E_{\text{gp}} - E_{\text{binding(LHe)}}) + E_{\text{binding(cal)}} = (0.8 - 0.62) + 42.9 = 43.1$ meV.

Finally, for all signals (photon, excimer, and evaporation), calorimeter noise is modeled by adding a gaussian broadening of $\sigma_{\text{cal}} = 0.5$ eV, only incrementally better than existing large-area microcalorimetry [43].

The lower panels of Figure 2 are the result of putting many individual energy deposits through the simple partitioning and detection monte carlo. The energy scale
is a reconstructed energy from observations of all signal quanta, constructed by multiplying the observed signal counts by their respective energies, and dividing by each signal type’s (assumed known) detection efficiencies. The y-axis is similarly an observed energy fraction, using the reconstructed energy to determine a reconstructed energy fraction in each signal channel. Below 19.77 eV, all signal appears in the quasiparticle channel, with broadening entirely from efficiency and calorimetry resolution. At \( \sim 100 \) eV, bands appear representing production and detection of small numbers of singlet photons and triplet excimers.

It can be seen in both the upper and lower panels of Figure 3 that nuclear and electron recoils exhibit distinct ratios of various atomic excitation modes. Because signal quanta can be distinctly identifiable by energy and arrival time, one can imagine constructing and applying multiple discrimination quantities, including a powerful singlet:triplet excitation ratio as in LAr-based experiments. Here we estimate only one simple discrimination quantity, using the ratio of observed energy in the quasiparticle modes vs the atomic excitation modes (adding singlet, triplet, and IR photons together).

The electron recoil leakage fraction at 50% nuclear recoil acceptance is a common measure of electron recoil discrimination power. Given this definition and monte carlo simulations as in the lower panels of Figure 2 we arrive at a discrimination estimate illustrated in the upper panel of Figure 3. The high production efficiency and high detection efficiency of the singlet and triplet excitations leads to extreme electron recoil rejection abilities at rather moderate energies (e.g., a leakage fraction of \( \sim 10^{-6} \) at 1 keV). Some amount of electron recoil discrimination should exist all the way down to the threshold of exciting electrons into the \( n=2 \) state at 19.77 eV.

This 19.77 eV energy marks an important transition not only in rejection power but more fundamentally in detector response. Below this energy, electronic excitation in \(^4\)He is impossible, given the possible electronic states of a helium atom. How this transition affects background expectations is described in Section VI.

V. THE PHONON AND ROTON SIGNAL CHANNEL

We briefly summarize the physics of the phonon and roton ‘quasiparticle’ excitations, as relevant to their production by a particle recoil, their propagation through the medium, and their interaction with boundaries of the medium. We then fold that literature into a simple simulation to better understand essential properties of quasiparticle-induced evaporation signals.

Kinetic excitations of the superfluid \(^4\)He medium, often termed ‘quasiparticles’, exhibit a distinctive dispersion relation as shown in the top panel of Figure 3. While the phonon branch is not dissimilar from normal sound, the higher-momentum portions (termed ‘rotons’, despite their lack of angular momentum) give the superfluid much of its interesting thermal properties. We describe this dispersion relation using a high-order polynomial fit to the data of [58].

Quasiparticles produced by a particle recoil are expected to undergo some level of thermalization at the recoil site, via quasiparticle-quasiparticle interactions. This thermalization process should be nearly complete at high recoil energies, leading to a distribution proportional to \( p^4 \). This assumption of thermalization becomes less motivated at lower recoil energies, where the initial quasiparticle density may be low enough so as to not encourage quasiparticle-quasiparticle scattering.

FIG. 3. Several fundamental characteristics of superfluid \(^4\)He quasiparticles are here illustrated. TOP: the dispersion relation. MIDDLE: the group velocity. BOTTOM: transmission probabilities at normal incidence in two cases, incident on a \(^4\)He-solid interface with solid phonon outgoing state (red dashed) and incident on a \(^4\)He-vacuum interface with outgoing state a \(^4\)He atom (blue solid). At both high and low momentum quasiparticles are of finite lifetime, and unlikely to reach an interface before decay.
When kinematically allowed, phonons spontaneously decay to two lower-energy phonons at a rate of \( \Gamma = A \epsilon^2 \), where \( \epsilon \) is the energy in K, and \( A = 7.12 \times 10^5 \text{ sec}^{-1} \). Over a significant portion of the dispersion curve, however, spontaneous decay is prohibited due to fundamental laws of energy and momentum conservation [61]. Quasiparticle decay into two quasiparticles is prohibited for \( p > 0.83 \text{ keV/c} \), and decay to any number of quasiparticles is prohibited for \( p > 1.10 \text{ keV/c} \). Experiments have confirmed the existence and approximate placement of this boundary between instability at low momentum and long-term stability at higher momentum [62]. At even higher momentum, there exists an upper bound to the stable momentum window, often taken to be \( p = 4.54 \text{ keV/c} \). This momentum window of infinite quasiparticle lifetime is highlighted in Figure 3.

Quasiparticle propagation can be taken to be ballistic within the \(^4\text{He}\) itself, assuming low (but readily achieved) number densities for scattering sites of three different types. Thermal phonons are no longer a relevant scattering mechanism at temperatures of \( \lesssim 100 \text{ mK} \). (The temperature dependence is steep, because such scattering is dominated by multi-phonon processes). \(^3\text{He}\) isotopic impurities can also serve as quasiparticle scattering sites (with scattering cross section \( \sim 10^{14} \text{ cm}^2 \)). To reach macroscopic mean free paths, the \(^3\text{He}\) concentration must be \( \lesssim 10^{-9} \). While this is lower than the natural concentration (\( \sim 10^{-7} \), \(^3\text{He}\) removal is straightforward using a heatflash method in the superfluid state [63]. \(^4\text{He}\) with isotopic contamination lower than \( 10^{-12} \) is commercially available [64]. It should also be noted that at \( \lesssim 100 \text{ mK} \) nearly all \(^3\text{He}\) will be forced the boundaries of the superfluid material. Last, the third possible quasiparticle scattering site population is quantum vortices, which are long-lived bulk excitations of angular momentum surrounding a hollow (or normal-fluid-filled) core. While the scattering cross section off such vortex cores is large, the concentration of quantum vortices can be assumed low.

Assuming quasiparticles which are both ballistic and stable against spontaneous decay, all complexity is limited to their initial production by the recoil and their subsequent interactions with the boundaries of the \(^4\text{He}\) superfluid. There exist two types of material boundaries: \(^4\text{He}-\text{to-solid}\) (immersed calorimetry or passive structures) and \(^4\text{He}-\text{to-vacuum}\) (the top surface, which in our case is instrumented with calorimetry above the vacuum gap). The ideal solid interface is one that is highly reflective to quasiparticles, such that a quasiparticle may reflect several times efficiently before finally escaping as quantum evaporation. Similarly, quasiparticle transmission into solid surfaces, through 1-to-1 transmission into phonons, can be considered a signal loss mechanism due to our primary reliance on the evaporation channel for signal gain and threshold suppression. A second loss process is quasiparticle downconversion (1-to-n processes), which may play a key role in quasiparticle signal loss, by degrading the quasiparticle energy even only slightly, to below the quantum evaporation threshold of 0.62 meV. For each surface interaction, the probability of each outgoing state (reflection, transmission, downconversion) is dependent on the incoming quasiparticle momentum and incident angle. Unfortunately, all such interface interaction probabilities are poorly constrained by the experimental literature.

**Solid Interfaces** In the thermal regime, the extreme Kapitza resistance observed at this interface can be used to imply an upper limit on the roton contribution to thermal conductivity. From these thermal models, one expects the phonon transmission probability to be \( \sim 10^{-4} \) and the roton transmission probability to be at least one order of magnitude lower [67]. While the acoustic mismatch theory of Khalatnikov [68] predicts very little thermal transport across the boundary between a solid and superfluid helium, most thermal transport experiments show much greater transmission. In a model proposed by Adamenko and Fuks [69], such thermal transport is increased for rough surfaces, with resonant transmission occurring when the quasiparticle wavelength is comparable to the length scale of surface roughness. This model has received recent experimental support [70] for phonons passing from a silicon single crystal into superfluid helium, indicating that controlling surface roughness on atomic scales might be used to enhance or suppress phonon and roton reflectivity.

There is strong tension between this thermal case (both model and observation) and observations of athermal pulses of quasiparticle excitation. For example, laboratory studies related to the HERON project showed that, for quasiparticle populations produced by alpha recoil, the quasiparticle reflection probability was only \( \sim 30\% \), with only very subtle variation with material characteristics and preparation [71]. At the same time, measurements by Brown and Wyatt [72] show the probability of a R\(^+\) roton transmitting its energy into the surface of an immersed bolometer to be only \( 2.8 \times 10^{-3} \), implying a very high reflection probability.

It is evident that quasiparticle interactions at liquid-solid interfaces are poorly understood. In our simulations here, we use as a starting point the purely theoretical work of I.V. Tanatarov et al. [65], and then suppress the reflection probability by a scaling factor, varied to account for the large uncertainty in the literature. We treat all solid reflection as diffuse rather than specular, consistent again with the experiments of HERON [71] and others.

**Vacuum Interfaces** Experiments observing quasiparticle interactions with the vacuum interface exhibit much greater agreement with theoretical expectation (perhaps due to the lack of surface roughness). The evaporation probability is quite high for high-momentum phonons at all angles, with some non-zero probability for high-momentum R\(^+\) rotons. The evaporation probability is near-zero for R\(^-\) rotons due to such excitations’ anti-
parallel propagation and momentum vectors. Several theoretical descriptions of quantum evaporation have been given (see [73–77]), we copy that of M.B. Sobnack et al. [66] (as in Figure 4 upper right), with the exception that in the simulation work here we scale this theoretical expectation down by a factor of two to be consistent with experiment (as in [78]). Reflection probabilities at this interface are taken again from the work of I.V. Tanatarov et al. [65], scaled where necessary to accommodate the Sobnack evaporation probability.

A. Simulation of evaporation channel signal characteristics

Although the quasiparticle interactions at surfaces are poorly understood, it is instructive to construct a quasiparticle propagation simulation to gain some expectation of quasiparticle signal characteristics. A detector geometry of 20 cm diameter and 20 cm liquid height serves as the baseline (~1 kg $^4$He), with a calorimeter for the evaporation sensor immediately above the liquid-vacuum interface. In this simulation, quasiparticles are released isotropically from an origin point on the central axis. Quasiparticles of momentum outside the momentum window of stability (as in Figure 3) are immediately removed from the population; quasiparticles in the stable window are simulated until they are either transmitted across a solid surface as a phonon, transmitted across a vacuum surface as an atom, or are ‘lost’ via a general scaling factor applied at each solid interaction. This loss probability is meant to capture the general uncertainty in physics at this interface.

Some basic results of this propagation monte carlo are shown in Figure 5 and Figure 6 including variation in the loss probability per surface interaction. Figure 5 indicates that the falltime of quasiparticle evaporation signal is expected to be order 10 to 100 ms, dependent

![Figure 4](image1.png)

**FIG. 4.** Quasiparticle transmission and reflection probabilities as a function of incoming state momentum (x-axis) and incidence angle (y-axis). We combine the quasiparticle reflection description of I.V. Tanatarov et al. [65] with the evaporation description of M.B. Sobnack et al. [66]. The solid transmission probability (upper left panel) has been multiplied by a factor of 20 for visibility. The vacuum transmission probability (upper right panel) has not yet been reduced by a factor of 2 to better match experiment. Solid white lines indicate the boundaries between phonon, $R^-$, and $R^+$ regions. Dashed white lines indicate the boundaries of the region for which the dispersion relation is multi-valued in energy.

![Figure 5](image2.png)

**FIG. 5.** Simulated pulses from a helium evaporation channel, showing the pulse decay timescales given several distinct quasiparticle loss probabilities per interaction with solid surface (blue: 0.7, green: 0.1, red: 0.01). The pulse cannot be described using a single exponential falltime; the exponential time constant increases at late times in each case due to the survival of slow-moving excitations.
on loss probability per interaction. This is reassuring, in that these timescales are short enough to avoid a pileup background.

In Figure 6, quasiparticle evaporation probability (allowing for many reflections) is plotted vs. initial quasiparticle momentum. The figure shows that the overall quasiparticle detection efficiency are strongly dependent on surface quasiparticle reflectivity. This is particularly true for those momenta with low evaporation probabilities, such as the R− case (mid momenta). The blue curve in Figure 6 illustrates a loss per surface interaction of 0.7, similar to what was observed in the HERON project R&D (resulting in a few-percent evaporation efficiency after allowing multiple reflections). The red curve represents a loss probability per surface interaction of 0.01, somewhat closer to what is expected theoretically in the case of a perfectly smooth surface. It is clear that understanding and increasing quasiparticle reflectivity at solid interfaces is an important aspect of future R&D.

Further results of this quasiparticle propagation simulation appear in an appendix.

![Figure 6](image-url)

**FIG. 6.** Probabilities for eventual evaporation after many reflections (including times up to one second after recoil), as a function of initial quasiparticle momentum. As in Figure 3, color represents different probabilities of loss per solid surface interaction (blue: 0.7, green: 0.1, red: 0.01). It can be seen that the loss probability has a strong effect on eventual evaporation probability, given the many expected reflections. The geometry of the simulated liquid volume is also varied: Solid lines are for a cylindrical volume of equal height and diameter (20 cm), while dashed lines are for a ‘pancake’ geometry in which the depth is much less than the diameter (20 cm depth, infinite diameter). It can be seen that a pancake geometry enhances the evaporation probability by boosting the fraction of boundary interactions occurring on the vacuum interface.

### B. The quasiparticle-only energy regime

As mentioned previously, atomic excitations (IR photons, 15.5 eV singlet photons, and triplet excimers) require a minimum threshold energy of 19.77 eV. For a search focused on dark matter masses below several hundred MeV/c², the nuclear recoil signal spectrum would lie entirely below the atomic excitation energy scales, entirely within a quasiparticle-only energy scale.

For a search focused on the lowest energies, any atomic excitations could be treated as a veto flag for efficiently tagging and rejecting background recoils at higher energies.

Although electron recoils are inhibited in this window, discrimination information is still useful for the exclusion of any possible ‘dark rate’ of detector-induced false signals. At higher energies multiple ‘flavors’ of atomic excitation enable background rejection, due to their distinguishable energies and arrival times. At these lower quasiparticle-only energies, multiple ‘flavors’ of quasiparticles (phonon, R−, R+) can serve a similar purpose, contributing to the observation a rough description of the initial quasiparticle momentum distribution. There are two methods for observing momentum-distribution information: 1) Quasiparticles populations of differing momentum exhibit significantly differing velocity. The first arrival of these populations at the liquid surface can appear as several distinct peaks in evaporation rate, separated by ms timescales, with amplitude proportional to spectral density (see appendix), 2) Given an energy threshold for evaporation, some fraction of quasiparticle energy can not evaporate and will instead be absorbed by the immersed calorimeter. Although not benefiting from the ‘adhesion gain’ principle, the immersed calorimetry might be able to provide information on what fraction of quasiparticle energy fell below the evaporation threshold of 0.62 meV. One could even imagine distinguishing between phonon populations, by purposefully varying the immersed calorimeter surface roughness such that different momentum populations were transmitted into different calorimeters, perhaps at different rates.

Given the 0.62 meV evaporation threshold and the 4 GeV nuclear mass, dark matter sensitivity in the simple nuclear recoil channel is ultimately limited to a mass region of ≥1 MeV. The nuclear recoil endpoint for O(keV) dark matter masses lies below the energy of a single evaporation quantum, given the requirement of momentum conservation. However, O(keV) dark matter can recoil instead off the bulk superfluid material, in effect bypassing the constraints imposed by the nuclear mass. Such a recoil can directly produce an off-shell quasiparticle (of comparatively high energy and low momentum) which then decays in multiple on-shell observable quasiparticles. While this off-shell process is significantly suppressed relative to the nuclear recoil case, it allows for up to 100% of the dark matter kinetic energy to be transmitted to the target material in the form of observable excitations. The amplitudes for multi-
excitation production are known from the ultracold-neutron field, and the associate dark matter signal sensitivities have recently been calculated, extending the reach of low-threshold $^4$He targets into the keV range\cite{13, 79}. An additional advantage of this detection approach is that the multi-excitations are produced in a distinct back-to-back (momentum-canceling) relative orientation, potentially allowing for background rejection through coincidence requirements.

VI. BACKGROUND SIMULATIONS

Since the number density of dark matter varies inversely with its mass $M_{\chi}$, the requirements on active mass and backgrounds are significantly relaxed for searches in the keV-GeV DM mass range. Additionally, the small recoil energies of these interactions mean that there is little overlap with the Compton scattering and beta backgrounds with characteristic energies of $O(100$ keV), so underground operation and use of radiopure materials developed for high mass WIMP searches should be sufficient to guarantee sub-dominant radioactive backgrounds.

In order to quantify the magnitude of radiogenic gamma backgrounds, GEANT4 simulations \cite{81–83} were performed on a simplified detector with shielding geometry modeled on the proposed design of the SuperCDMS SNOLAB experiment \cite{80}. Two types of backgrounds were investigated: electron recoil events caused by Compton scattering and photoabsorption, and coherent gamma scattering events in the form of Rayleigh, nuclear Thomson, and Delbrück scattering, which become significant sources of nuclear recoils at low recoil energies \cite{84}. The simulation geometry, depicted in Figure 7 consisted of a cube of active helium volume in a copper cryostat surrounded by layers of high density polyethylene (HDPE), lead, and water, with thicknesses derived from the SuperCDMS SNOLAB design. Simulations were performed on geometries with 1 kg, 10 kg, and 100 kg active helium masses.

Gamma rays produced in significant amounts in the $^{238}$U, $^{232}$Th, $^{40}$K, $^{60}$Co, and $^{137}$Cs decay chains, assuming secular equilibrium where relevant, were generated uniformly and isotropically in each component of the simulation geometry. Events were tracked to produce either the spectrum of gamma energies entering the active helium volume or the spectrum of electron recoil events. Simulations used a modified version of the SHIELDING physics list to account for atomic shell effects in Compton scattering and neglected any detector response effects in producing either the gamma flux into the helium volume or the electron recoil spectrum. Spectra were produced by assuming the same impurity concentrations as SuperCDMS SNOLAB \cite{80}. To

![Simulation geometry used to obtain the Compton recoil energy spectrum and incident photon energy spectrum on an active helium volume. Each component is cylindrically symmetric except for the helium volume, which is cubic. The thicknesses of the shielding components, listed in parentheses, were adapted from SuperCDMS \cite{80}.](image)

FIG. 7. Simulation geometry used to obtain the Compton recoil energy spectrum and incident photon energy spectrum on an active helium volume. Each component is cylindrically symmetric except for the helium volume, which is cubic. The thicknesses of the shielding components, listed in parentheses, were adapted from SuperCDMS \cite{80}.

![Graph showing leakage fraction versus recoil energy and predicted recoil energy spectrum.](image)

FIG. 8. (a) The leakage fraction from the energy partitioning shown in Figure 2. Discrimination based on light and heat signals are not possible in the shaded region, but it may be possible to extend roton-only discrimination to lower recoil energies. (b) The predicted recoil energy spectrum contributing to electron recoil (ER) and nuclear recoil (NR) backgrounds in the 1 kg liquid helium detector mass. Total gamma ER backgrounds (dashed yellow) were simulated directly and combined with the leakage fraction for discriminated gamma ER background rates (yellow). Gamma NR backgrounds (blue), the sum of Rayleigh (dashed blue), nuclear Thomson (dash-dot blue), and Delbrück (dotted blue) scattering, were calculated from cross sections and the simulated incident spectrum. Other backgrounds considered were NR events from astrophysical neutrinos (green) and radiogenic neutrons (red). See text for details.
calculate the coherent scattering recoil spectrum, the radiogenic gamma flux into the detector was combined with the coherent sum of cross sections contributing to elastic gamma scattering, assuming only single scattering events. These cross sections were obtained from non-relativistic form factors for Rayleigh scattering [85], direct calculation for nuclear Thomson scattering, and interpolated amplitudes for Delbrück scattering [84, 86]. Structure effects, which may become significant at recoil interpolated amplitudes for Delbrück scattering [84, 86]. relativistic form factors for Rayleigh scattering [85], events. These cross sections were obtained from non-elastic gamma scattering, assuming only single scattering with the coherent sum of cross sections contributing to radiogenic gamma flux into the detector was combined calculate the coherent scattering recoil spectrum, the gamma flux, and d
R/dEγ = D \cdot \int dN/dEγ \cdot dσ/dEγ dEγ, \quad (5)

where dR/dEγ is the differential recoil rate, D is the number of helium nuclei per unit mass, dN/dEγ is the gamma flux, and dσ/dEγ is the differential cross section at a particular recoil energy. While it may be possible to discriminate multiple and single scattering events by examining TES hit patterns, doing so does not substantially alter background rates for recoil energies of interest since the mean free path for Compton scattering in helium is relatively long. In the analysis of simulated events, we do not consider discrimination between single and multiple scattering. Figure 8 shows the gamma background rates.

Coherent neutrino scattering from solar, atmospheric, and supernova neutrinos was modeled as a background using the method in [87]. The neutron background was assumed to be flat in recoil energy, with a constant value of $2.5 \times 10^{-5}$ DRU, which is on par with the neutron background projected by SuperCDMS [80].

VII. SENSITIVITY PROJECTIONS

Projected sensitivities to DM interactions are calculated using a profile likelihood ratio (PLR) analysis [88]. The PLR likelihood function is

$$
\mathcal{L}(\sigma_{\chi-n}) = \frac{e^{-\left(\mu_x + \sum_i \mu_j \right)}}{N!} \times \prod_{i=1}^{N} \left[ \mu_x \phi_i(E_{R_i}) + \sum_j \mu_j f_j(E_{R_i}) \right],
$$

where $i$ iterates over observed events, $j$ iterates over different background species, $\mu_x$ ($\mu_j$) is the expected number of signal (background) events, $N$ is the total number of observed events, and $f_x$ ($f_j$) is the signal (background) recoil energy PDF. At each point in parameter space, we quantify the degree to which a typical background-only simulation can reject the signal hypothesis, where "typical" is defined to be the median value of the PLR test statistic. The projected sensitivity is defined to be the curve in parameter space on which the signal hypothesis can be rejected with 90% confidence.

The detector is simulated as a liquid He-4 detector with recoil energy as the only observable and 100% efficiency at all recoil energies above threshold. We consider four generations of experiments with threshold-mass-runtime combinations of (40 eV, 10 g, 100 days), (10 eV, 1 kg, 1 year), (0.1 eV, 10 kg, 1 year), and (1 meV, 100 kg, 1 year). The first generation experiment we describe is ‘shovel ready’ in that it combines several already-demonstrated technologies with no required new R&D: a calorimeter of 3.5 eV baseline resolution (sigma) [49], an efficiency of converting recoil energy to evaporation of ~5% [71], and a 9× ‘adhesion gain’ of He atoms on a Si surface [71]. Significant future advancement in threshold appears plausible, given the three routes towards threshold reductions: 1) improving the quasiparticle evaporation efficiency, perhaps by reducing surface roughness on the solid surfaces, 2) the addition of high-adhesion-gain coatings on the calorimeter surface, and 3) the continued advancement of TES-based large area calorimetry, which has yet to hit any fundamental limit.

The nuclear recoil energy spectrum from dark matter-nucleus elastic scattering is modeled as in [87] [89], with a Maxwell-Boltzmann velocity distribution for the dark matter halo and a Helm form factor. We further model the backgrounds as described in Section IV. We reject electronic recoil backgrounds, namely Compton scattering, using the discrimination model described in Section III.C. Our projected sensitivity to spin-independent dark matter-nucleus scattering is shown in Figure 9 along with selected experimental constraints [87] [90] [99].

The "neutrino floor" in Figure 9 represents the curve in parameter space at which coherent elastic neutrino-nucleus scattering (CEvNS) becomes a limiting background for helium-target detectors. It is calculated using a technique similar to [87]. The coherent elastic scattering of solar neutrinos on helium nuclei is considered to be the only background in a hypothetical detector, and we define a new PLR by extending the likelihood function to include uncertainties in neutrino flux as nuisance parameters:

$$
\mathcal{L}(\sigma_{\chi-n}, \phi_j) = \frac{e^{-\left(\mu_x + \sum_i \mu_j \right)}}{N!} \times \prod_{i=1}^{N} \left[ \mu_x \phi_i(E_{R_i}) + \sum_j \mu_j f_j(E_{R_i}) \right] \times \prod_j \mathcal{G}_j(\phi_j, \Delta \phi_j).
$$

Here, $\mathcal{G}_j$ is a Gaussian distribution centered on the mean value of the flux of neutrino species $j$, and with standard deviation given by the uncertainty in that flux.

\[ (3.5 \text{ eV} \times 5\sigma) / (9\times \text{gain}) / (0.05 \text{eff.}) = 39 \text{ eV}. \]
The neutrino species we consider and their associated uncertainties are: \( pp \) (1%), \( pep \) (1.7%), \( ^7Be \) (10.5%), \( ^8B \) (8.8%), \( hep \) (15.5%), and \( CNO \) (30%). These values are a combination of theoretical [100, 101], experimental [102], and estimated [87] uncertainties. We do not consider atmospheric and diffuse supernova background neutrinos because it is kinematically unlikely that these recoils could ever mimic a WIMP signal in helium.

CEvNS becomes a limiting, irreducible background for dark matter experiments when the exposure is high enough that flux uncertainties rival Poisson fluctuations. For recoil energies less than about 1 keV, the dominant neutrino species are \( pp \) and \( ^7Be \). The geometric mean of their uncertainties is 3.2%, indicating that the solar neutrino background becomes significant at an exposure corresponding to about 980 expected recoils. We thus define the neutrino floor as the projected sensitivity of a He-4 detector with 1.6 tonne-yr exposure, for which the corresponding to about 980 expected recoils. We thus set arbitrarily low.

In the case of a heavy dark photon mediator \( F_{DM} = 1 \), the DM-nucleon and DM-electron scattering cross-sections are related:

\[
\frac{\sigma_e}{\sigma_{\chi-n}} = \left( \frac{A}{Z} \right)^2 \left( \frac{\mu_{\chi-e}}{\mu_{\chi-n}} \right)^2,
\]

where \( A \) and \( Z \) are the atomic mass number and atomic number of the target nucleus, respectively, and \( \mu_{\chi-e} \) (\( \mu_{\chi-n} \)) is the reduced mass between the DM particle and an electron (nucleon). Thus, we can translate our projected sensitivities into DM-electron space and compare to existing constraints on dark photon interactions. This is done in Figure 10 with current NR constraints translated into the \( \sigma_e \) plane using Equation 8. Note that we have not translated the sensitivities for our third and fourth generation experiments into \( \sigma_e \) parameter space. In the case of a dark photon-mediated nuclear recoil in helium, the photon propagator is modified by in-medium effects. These effects are negligible for recoil energies greater than 10 eV [113], legitimizing the translation of the first and second generation sensitivities. However, if they are significant for lower recoil energies, Equation 8 will not hold. Further work needs to be done to determine whether the other sensitivities can be similarly translated, based on a detailed calculation of the in-helium photon propagator.

The flux and kinetic energy of the dark matter can degrade because of the earth shielding, if dark matter has relatively high cross-sections with normal matter so that the mean free path of dark matter in the earth is comparable to or lower than the diameter of the earth. For underground experiments, dark matter could lose so much energy in interacting with the earth that the remaining kinetic energy of dark matter is below the energy threshold of the detector \( E_{thr} \) when it reaches the depth of the detector, thus indetectable by the...

![FIG. 9. Projected sensitivity of the helium detector to DM-nucleon SI interaction with 90% confidence level, through detection of elastic scattering and Bremsstrahlung emission [103]. Four combinations of exposure and energy threshold have been investigated: 1 kg-day with 40 eV (solid red), 1 kg-yr with 10 eV (dashed red), 10 kg-yr with 0.1 eV (dotted red), and 100 kg-yr with 1 meV (dashed dotted red). The label 100 m and 1478 m refers to the experiment site depth, related to the limit curves due to the earth shielding effect discussed in the main text. The dashed dotted dotted red curve corresponds to limit taking off-shell phonon sensitivity into account, assuming a massive mediator with 100 kg-yr exposure and 1 meV energy threshold, extrapolated from Knapen et al. [15]. The dashed dotted dotted black curve corresponds to the neutrino floor calculated for helium, as discussed in the main text. Some other limits are also plotted for comparison: the neutrino floor for xenon (solid black) [87], cosmic Microwave Background anisotropy (dotted black) [90], galaxy gas cooling (dashed black) [93], XQC experiment (dashed dotted black) [96], CRESST surface (dashed dotted green) [87], CRESST-II (dashed dotted green) [93], CDMS-Lite (dotted green) [92], XENON-1T (solid green) [93], and LUX (dashed green) [94]. The earth shielding limits for CRESST surface and CRESST-II are also shown [90]. The green and grey shaded region correspond to parameter space that has been excluded by direct detection experiments and astronomy, respectively.]
From the surface of the Earth, and the initial velocity is lower than a cut off speed, for computational reasons; ii) the DM particle reaches the depth of the experiment. The speed of the DM particles are recorded when they reach the depth of the experiment. The energy deposit available to the dark matter detector is conservatively assumed to be the kinetic energy of the DM particles. Consequently, the minimal velocity of DM needed to deposit energy above the energy threshold of the detector is determined when the average speed of the DM particles reaching the depth of the detector \(v\) and the standard deviation of that speed distribution \(\Delta v\) are related to the minimum velocity by \(v + 5\Delta v < v_{\text{min}}\). For more simulation details see Emken et al. [114]. As shown in Fig. 9, the cross-section lower limits constrained by the earth scattering are well above the cross-section upper limits. Also, the result shows the lower limits on cross-sections depend on detector overburden as well as the detector energy threshold.

Bremssstrahlung photons from the dark matter–nucleus inelastic interaction can be used to extend the sensitivity reach of a dark matter experiment to lower mass dark matter [103]. As mentioned before, because of kinematic mismatch, low-mass dark matter only deposits a very small amount of energy into the nucleus through the elastic scattering channel. In contrast, the energy transfer through the inelastic scattering channel with emission of Bremssstrahlung photon can be much larger. For light dark matter, the hierarchy for the maximum energy of a Bremssstrahlung photon \(\omega_{\text{max}}\) and the maximum energy deposited into nucleus \(E_{R,\text{max}}\) is [103]

\[
E_{R,\text{max}} = 4 \left( \frac{m_X}{m_N} \right) \omega_{\text{max}} \ll \omega_{\text{max}} \quad (m_X \ll m_N)
\]

where \(\omega_{\text{max}} = \mu_{X-N} \cdot v^2 / 2\).

Fig. 10 shows that the Bremssstrahlung signal could extend the experimental sensitivity to lower dark matter mass than the reach from elastic scattering.
The use of superfluid helium with calorimetric readout offers a unique avenue for carving out a vast swath of dark matter parameter space. Current technology will allow probing of dark matter masses as low as 60 MeV/c². With further advancements in calorimeter threshold and helium quasiparticle reflectivity, the technology can probe dark matter masses as low as 600 keV/c² (via simple elastic nuclear recoils). We have also shown how the neutrino floor behaves for low dark matter masses, which can be useful for many direct detection approaches.

In the upcoming years, when experiments like LZ and SuperCDMS begin to reach the neutrino floor, it will be crucial to explore lower-threshold and smaller-scale technologies including the concept described in this letter.

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Appendix A: Pulse shape information in the quasiparticle-only regime

In this appendix we apply the quasiparticle propagation monte carlo of Section V to ask what types of quasiparticle information might be contained (and readable) in the evaporation pulse shape. Specifically, we ask two questions: does the evaporation pulse shape contain significant information on the phonon momentum distribution? And, does the evaporation pulse shape contain significant information on the recoil depth below the liquid surface? We tentatively find that the evaporation channel conveys significant information in both respects.

First, we inspect the upper panel of Figure 11 in which evaporation time is plotted on the x-axis and initial quasiparticle momentum is plotted on the y-axis. Dashed lines bound the momentum window of long-lived quasiparticles. Coloring is by number of reflections before evaporation: black is 0, red is 1, green is 2, blue is ≥ 3.

This particular population of quasiparticles is released isotropically from the center of a 20 cm diameter 20 cm height (≈ 1 kg) $^4$He cylinder, with vacuum interface at top. In can be seen in this plot that the high-momentum $R^+$ states arrive at the liquid surface first, and with a small range evaporation times. This restricted range of $R^+$ evaporation times is partly due to a small range in $R^+$ velocities and partly due to the restricted range of incident angles that allow $R^+$-induced evaporation (see Figure 4, upper right panel). After a first burst of $R^+$ evaporation, a slightly slower pulse of phonon-induced evaporation begins, smeared over a ms timescale due to a range of phonon velocities and an unrestricted range of phonon incident angles that allow evaporation. Last, at ≈ 1.9 ms an ‘echo’ appears off the bottom surface. This echo is the first opportunity $R^-$ rotons have had to induce evaporation, as they typically cannot induce evaporation in the $R^-$ state. These $R^+$ rotons convert to other more evaporable momenta upon reflection on the bottom surface (see again Figure 4 $R^-$ reflection probabilities into phonon and $R^+$ modes).

Given the three peaks in time, representing in order $R^+$, then phonon, then $R^-$ populations, we expect the relative scale of these initial momentum populations to be observable in the relative scale of these pulse features. This is what we see in the middle panel of Figure 11. We vary the initial distribution from the naive thermal expectation $p^2$ in blue, to a less thermal $p^4$ in green, to a flat ‘white’ $p^0$ distribution in red. The relative height of the $R^+$ peak decreases while the phonon peak increases, as expected. The short timing gap between the $R^+$ pulse and the phonon pulse is key to this measurement, and would perhaps be a design driver in a future device. In the case of a merging of these two components, a ‘prompt fraction’ shape quantity might be used as a rough metric for the $R^+$:phonon ratio.

The bottom panel of Figure 11 again varies the initial momentum distribution, and now plots the spatial distribution of evaporation from the surface in the first ms. The thermal $p^2$ initial population exhibits a sharp boundary marking the maximum angle of $R^+$ evaporation (again see Figure 4, upper right panel). Some estimate of the $R^+$:phonon ratio can again be gleaned here, by constructing a ‘low radius fraction’ quantity.

Figure 12 is a similar study of pulse shape and spatial distribution of evaporation, but here we vary recoil depth below the liquid surface (red: 1 cm from bottom, green: the center, blue: 1 cm from top). Several features are evident. Separation time between the $R^+$ peak and the phonon pulse depends on distance to the surface (with the two features ultimately merging for recoils near the top, blue). Separation time between the $R^+$ peak and the $R^-$ ‘echo’ is even more dramatically dependent on depth, but in the opposite direction (with the $R^-$ echo ultimately merging with the phonon population for recoils near the bottom, red).

The lower panel Figure 12 shows a strong spatial dependence in the evaporation signal (again only showing the first ms), depending on the depth of the recoil. For a recoil very near the liquid surface, the evaporation is tightly restricted in radius, with this spatial pattern loosing as the depth increases.
FIG. 11. An illustration of evaporation pulses communicating information of a recoil’s initial quasiparticle momentum distribution. For details and interpretation, see the text. In all panels, the quasiparticle population is released from the center of a 20×20 cm cylindrical volume. TOP: The initial momentum is plotted vs the evaporation time (if any). Dashed horizontal lines indicate the window of stability in momentum; dotted horizontal lines indicate the boundaries between phonon, R⁺, and R⁺⁺ momentum regions. Coloring indicates the number of boundary reflections before evaporation (black:0, red:1, green:2, blue: ≥3). MIDDLE: the evaporation pulse shapes of three initial momentum distributions are compared: p⁰ (red), p¹ (green), and p² (blue). BOTTOM: the radial distribution of evaporation in the first 1 ms is shown for these same three momentum distributions.
FIG. 12. An illustration of evaporation pulses communicating information of a recoil’s depth in the $^4$He target. Here the vertical position of the recoil is varied in a 20cm diameter 20 cm height cylindrical vessel (red: 1 cm from bottom, green: center, blue: 1 cm from top). The initial momentum distribution is assumed isotropic and following the thermal distribution as $p^2$. TOP: Evaporation pulse shape for these three depths, showing increasing and decreasing timescales between features. BOTTOM: Evaporation spatial distribution for these three depths, showing enhancement or lack of enhancement at central radii.