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Directional detection of dark matter with two-dimensional targets

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ABSTRACT

We propose two-dimensional materials as targets for direct detection of dark matter. Using graphene as an example, we focus on the case where dark matter scattering deposits sufficient energy on a valence-band electron to eject it from the target. We show that the sensitivity of graphene to dark matter of MeV to GeV mass can be comparable, for similar exposure and background levels, to that of semiconductor targets such as silicon and germanium. Moreover, a two-dimensional target is an excellent directional detector, as the ejected electron retains information about the angular dependence of the incident dark matter particle. This proposal can be implemented by the PTOLEMY experiment, presenting for the first time an opportunity for directional detection of sub-GeV dark matter.

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1. Introduction

The Weakly Interacting Massive Particle (WIMP) is currently the dominant theoretical paradigm for dark matter (DM), and has guided experimental search efforts in recent decades. Direct detection experiments, which search for DM-nucleus collisions, are currently targeting the WIMP parameter space [1–7]. However, null results from these searches motivate renewed consideration for a broader range of DM models. One possibility involves DM particles below the ∼GeV scale, which arise in a variety of theory scenarios [8–21]. Current direct detection experiments lose sensitivity to sub-GeV DM because the nuclear recoil energy is too small to be detected. However, DM with mass below a target nucleus deposits a greater fraction of its kinetic energy on an electron than a nucleus, making electrons a favorable target for light DM detection.

Consider the case of MeV-scale DM, which carries about an eV of kinetic energy, enough to excite atomic electrons after scattering [22]. The first limits on such processes have been set using data from the Xenon10 experiment [23], with recent work extending this analysis to Xenon100 [24]. The energy gap for electronic excitations in noble gases is ∼10 eV, which places a lower bound on the DM mass that can be probed with these methods. However, a smaller energy deposit can up-scatter valence electrons in semiconductors with band gaps ∼1 eV [22,25]. As a result, semiconductor targets are more sensitive to DM in the 1–10 MeV mass range [26,27]. Superconducting targets with ∼meV energy gaps are capable of reaching ∼keV masses [28,29].

This Letter proposes an alternative approach using two-dimensional (2D) materials as targets. In this setup, an incident DM particle can deposit sufficient energy on a valence electron to eject it from the target. The energy and direction of the recoiling electron is then directly measured with a combination of position measurements, time-of-flight, and energy deposition in a calorimeter. This is in contrast to scattering in bulk targets, where the scattered particle (nucleus or electron) produces secondary excitations before measurement [2,30,26,27], erasing the initial directional information in the scattering. Using 2D targets, DM masses down to the MeV scale can be probed if the energy required to eject the electron is a few eV.

Most importantly, 2D targets allow one to measure the direction of the incoming DM because the differential cross section for the outgoing electron is peaked in the forward direction. The lattice structure of the target can even yield diffraction patterns in the electron angular distribution for certain kinematics. Directional detection has long been recognized as a powerful tool in the study of DM, both as a discriminator against background sources and also because it leads to a daily modulation of the signal rate [31].

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There are currently no feasible proposals for directional detection of sub-GeV DM [32], making the use of 2D targets a powerful tool in pushing sensitivities to lower DM masses. We will describe a potential experimental realization using the PTOLEMY experiment [33].

2. Dark matter scattering in graphene

As a concrete example of a 2D target, we focus on monolayer graphene, which is especially convenient because analytic solutions for the electron wavefunctions in the tight-binding approximation are tractable due to the symmetries of the lattice [34]. We compute the DM scattering rate here, and in Section 3 we show that the direction of the scattered electron retains a strong directional correlation with the DM direction.

Monolayer graphene consists of carbon atoms arranged in a two-dimensional honeycomb lattice. The distance between neighboring carbon atoms is $a = 0.142$ nm. The lattice is built from two distinct triangular sub-lattices. Four out of the six electrons of a carbon atom are valence electrons, occupying $(2s)(2p)^3$ orbitals.[1] The 2s orbital becomes 'hybridized' with the in-plane p_x and p_y orbitals, such that the energy eigenstates (called σ bonds) are linear combinations of 2s, $2p_x$, and $2p_y$. The out-of-plane p_z orbitals remain unhybridized and form covalent bonds, called π. We outline the important features of the unhybridized π electron wavefunction here, relegating further details and a discussion of the σ electrons to the Supplementary Material.

Within the tight-binding model, we approximate the wavefunction by a sum over nearest neighbors, corresponding to four lattice sites. The Bloch function for a π electron is given by

$$\Psi_\pi(\boldsymbol{\ell}, \mathbf{r}) \approx \mathcal{N}_{\boldsymbol{\ell}} \left(\phi_{2p_z}(\mathbf{r}) + e^{i\varphi_{\boldsymbol{\ell}}} \sum_{j=1}^{3} e^{i\boldsymbol{\ell}\cdot\mathbf{R}_j} \phi_{2p_z}(\mathbf{r} - \mathbf{R}_j) \right) \tag{1}$$

for lattice momentum $\boldsymbol{\ell} = (\ell_x, \ell_y) \in \mathrm{BZ}$ in the Brillouin zone. Here, $\mathcal{N}_{\boldsymbol{\ell}}$ is a normalization constant, $\mathbf{R}_j$ are the nearest-neighbor vectors, and $\varphi_{\boldsymbol{\ell}}$ is an $\boldsymbol{\ell}$-dependent phase. We take a hydrogenic orbital for the $2p_z$ wavefunction of carbon,

$$\phi_{2p_z}(\mathbf{r}) = \mathcal{N} a_0^{-3/2} \frac{r}{a_0} e^{-Z_{\mathrm{eff}} r/2a_0} \cos\theta \,, \tag{2}$$

where a_0 is the Bohr radius and $\mathcal{N}$ is the normalization. The effective nuclear charge $Z_{\mathrm{eff}} \simeq 4.03$ is chosen to fit the numerical solution for the overlap between adjacent $2p_z$ orbitals. The Fourier transform of Eq. (1) is

$$\widetilde{\Psi}_\pi(\boldsymbol{\ell}, \mathbf{k}) = \mathcal{N}_{\boldsymbol{\ell}} \left(1 + e^{i\varphi_{\boldsymbol{\ell}}} f(\boldsymbol{\ell} + \mathbf{k})\right) \widetilde{\phi}_{2p_z}(\mathbf{k}), \tag{3}$$

where $\mathbf{k}$ is the momentum conjugate to $\mathbf{r}$, $f(\boldsymbol{\ell} + \mathbf{k}) = \sum_{j=1}^{3} e^{i(\boldsymbol{\ell}+\mathbf{k})\cdot\mathbf{R}_j}$ is a sum of phase factors, and the Fourier transform of the atomic orbital is well-approximated by

$$\widetilde{\phi}_{2p_z}(\mathbf{k}) \approx \widetilde{\mathcal{N}} a_0^{3/2} \frac{a_0 k_z}{\left(a_0^2 |\mathbf{k}|^2 + (Z_{\mathrm{eff}}/2)^2\right)^3} \tag{4}$$

with normalization $\widetilde{\mathcal{N}}$.

Analytic forms for the σ electron wavefunctions are also possible to derive, but are more complicated than their π counterparts because the coefficients of the basis orbitals must be computed by diagonalizing a 6×6 Hamiltonian. The π (σ_1) electrons have binding energies ∼0–6 (13–18) eV.

If the scattered electron is ejected from the material after scattering, then its final-state wavefunction is well-modeled by a plane wave [35]. The initial-state wavefunction corresponds to an electron in any of graphene's four valence bands. The cross section for a DM particle of mass m_χ and initial velocity $\mathbf{v}$ to scatter off an electron in band $i = \pi, \sigma_1, \sigma_2, \sigma_3$ with lattice momentum $\boldsymbol{\ell}$ is then

$$v\,\sigma_i(\boldsymbol{\ell}) = \frac{\bar{\sigma}_e}{\mu_{e\chi}^2} \int \frac{d^3k_f}{(2\pi)^3} \frac{d^3q}{4\pi} |F_{\mathrm{DM}}(q)|^2 \left|\widetilde{\Psi}_i(\boldsymbol{\ell}, \mathbf{q} - \mathbf{k}_f)\right|^2 \times \delta\left(\frac{k_f^2}{2m_e} + E_i(\boldsymbol{\ell}) + \Phi + \frac{q^2}{2m_\chi} - \mathbf{q}\cdot\mathbf{v}\right), \tag{5}$$

where $-E_i(\boldsymbol{\ell})$ is the band energy, m_e is the electron mass, $\mathbf{k}_f$ is the final electron momentum, $\mathbf{q}$ is the momentum transfer (*i.e.*, the outgoing DM has momentum $m_\chi \mathbf{v} - \mathbf{q}$), and $\mu_{e\chi}$ is the DM-electron reduced mass. Note that $\mathbf{q}$ and $\mathbf{k}_f$ are independent scattering variables since the initial bound-state wavefunction is an energy eigenstate but not a momentum eigenstate, with Fourier components at all $\mathbf{k}$ values. For the same reason, Eq. (5) contains only a single delta function enforcing energy conservation, with no corresponding delta function for momentum conservation. $\Phi \simeq 4.3$ eV is the work function of graphene [35], defined as the energy difference between the Fermi surface and the vacuum.[2] Following Ref. [22], we define

$$\bar{\sigma}_e \equiv \frac{\mu_{e\chi}^2}{16\pi m_\chi^2 m_e^2} \overline{|\mathcal{M}_{e\chi}(q)|^2}\Big|_{q^2 = \alpha^2 m_e^2} \,, \tag{6}$$

with $\overline{|\mathcal{M}_{e\chi}(q)|^2}$ the spin-averaged amplitude, to be the scattering cross section for DM off a free electron with $q = \alpha m_e$. The momentum dependence of the matrix element is then absorbed into the DM form factor $F_{\mathrm{DM}}(q) = \overline{|\mathcal{M}_{e\chi}(q)|}/\overline{|\mathcal{M}_{e\chi}(\alpha m_e)|}$. We do not include the so-called Fermi factor, which enhances the rate at low recoil energies due to the distortion of the outgoing electron wavefunction by the Coulomb field of the nucleus. This factor is significant for bulk materials, but negligible for a 2D material for two reasons: the ionized electron energy must be high enough to overcome the work function, and the ionized electron travels single-atom distances and thus spends little time in the vicinity of the nucleus.

To obtain the total rate per unit time and detector mass, we must integrate Eq. (5) over all $\boldsymbol{\ell} \in \mathrm{BZ}$ and all incoming DM velocities, then sum the contributions from the four valence bands:

$$R = 2 \sum_{i=\pi,\sigma_{1,2,3}} \frac{\rho_\chi}{m_\chi} N_C A_{\mathrm{uc}} \int \frac{d^2\ell}{(2\pi)^2} d^3v\, g(\mathbf{v})\, v\, \sigma_i(\boldsymbol{\ell}) \,, \tag{7}$$

where $g(\mathbf{v})$ is the lab-frame DM velocity distribution, $A_{\mathrm{uc}} = 3\sqrt{3}a^2/2$ is the area of the unit cell, $N_C \simeq 5 \times 10^{25}$ kg^{-1} is the density of carbon atoms in graphene, and $\rho_\chi \simeq 0.4$ GeV/cm^3 is the local DM density [37]. The factor of two in Eq. (7) accounts for the degenerate spin states in each band.

The kinematics of the scattering process dictate that there is a minimal DM velocity required to eject an electron of momentum k_f from the target via a momentum transfer q:

$$v_{\mathrm{min}}^i(\boldsymbol{\ell}, k_f, q) = \frac{E_{\mathrm{er}} + E_i(\boldsymbol{\ell}) + \Phi}{q} + \frac{q}{2m_\chi} \,, \tag{8}$$

where $E_{\mathrm{er}} \equiv k_f^2/2m_e$. We assume the Standard Halo Model (SHM) [38] for $g(\mathbf{v})$, with $v_0 \simeq 220$ km/s [39] and $v_{\mathrm{esc}} \simeq 550$ km/s [40]

[1] The core 1s electrons have binding energies of several hundred eV and contribute negligibly to the scattering rate.

[2] The work function is not an intrinsic property of graphene, and can be manipulated with a suitable choice of substrate; see *e.g.*, Ref. [36].

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