A Proposal to Detect Dark Matter Using Axionic Topological Antiferromagnets

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Astrophysical and cosmological observations of the last 40 years provide strong evidence for the existence of non-baryonic dark matter (DM) \cite{1,2}. Among possible candidates are dark axions (DA) \cite{5,6,7,8,9,10,11}, hypothetical particles \cite{12,13,14,15,16} suggested to solve the $C\bar{P}$ problem in quantum chromodynamics (QCD) \cite{17}. Searching for the DA is challenging due to its weak coupling to ordinary matter (e.g. photons). For DA masses $m_a \lesssim 0.2$ eV the local DA field, $\theta_D$, can be described as a classical coherent state. If the local DM is indeed a DA, then $\rho_{DM} = |\theta_D(t)|^2 m_a f_a^2/2$, where $m_a$ and $f_a$ are the unknown axion mass and “decay constant”, and the local DM density is $\rho_{DM} \approx 0.4$ GeV cm$^{-3}$ \cite{3}. The DA field oscillates in time, with a frequency dominated by the rest energy, $m_a c^2$, and an intrinsic width set by the galactic velocity dispersion, $\sigma_v \approx 230$ km s$^{-1}$ $\Rightarrow \Delta \omega_a/\omega_a = \sigma_v^2/c^2 \approx 10^{-6}$. The central frequency is $\nu = 0.25(m_a/\text{meV})$ THz. The QCD axion mass can be computed in chiral perturbation theory or on the lattice, and is given by $m_a = 0.6$ meV (10$^{10}$ GeV/$f_a$) \cite{13,14,15} (we use units $h = c = 1$ if not stated otherwise).

Only one DM search, the Axion Dark Matter eXperiment (ADMX) \cite{17,18}, has made a significant constraint on the QCD axion parameter space predicted by the favoured Kim-Shifman-Vainshtein-Zhadarov (KSVZ) \cite{19,20} and Dine-Fischler-Srednicki-Zhitnitsky (DFSZ) \cite{21,22} axion models. The QCD axion mass can span roughly $10^{-12}$ $\lesssim m_a \lesssim 10^{-2}$ eV, satisfying astrophysical constraints on the couplings \cite{23,24}, and with $f_a$ less than the Planck scale. There are hints, however, pointing to the meV range, particularly for DFSZ-type models \cite{25,26}. Furthermore, constraints from the search for solar axions by the Cern Axion Solar Telescope (CAST) \cite{27}, combined with the prediction of KSVZ and DFSZ models provide a target range for the axion-photon coupling $10^{-10}$ GeV$^{-1} \lesssim g_\gamma \lesssim 10^{-13}$ GeV$^{-1}$ for $m_a \approx 1$ meV.

The problem of detecting DAs can be seen from the cycle average power output from the axion-induced electric field in the loop diagram in Fig. 1(a):

$$P_0 = \frac{1}{2} E_0^2 V_{\text{eff}} \omega_a = g_\gamma^2 B_0^2 \rho_{DM} m_a^2 V_{\text{eff}} \omega_a.$$

(1)

Taking the effective volume $V_{\text{eff}} \approx (2\pi/m_a)^3$ from the vacuum dispersion relation, $m_a = 1$ meV, $g_\gamma = 10^{-10}$ GeV$^{-1}$, and $B_0 = 1$ T, gives $P_0 = 10^{-27}$ W. DA searches must amplify this tiny signal, for example by resonance in a microwave cavity \cite{17,18}, ferromagnetic resonance \cite{28}, or coherent enhancement \cite{29}. These methods, however, and many others \cite{28,30,31,32,33,34}, are only effective up to around 0.5 meV.

Axionic degrees of freedom were predicted to materialise as quasiparticles in magnetically doped topological insulators (TIs) \cite{35}, Cr$_2$O$_3$ ($\theta_Q = \pi/36$) \cite{36,37}, $\alpha$-Fe$_2$O$_3$ \cite{38} with a corundum structure, spinels \cite{39} and magnetic TI heterostructures \cite{40}. The signatures of the topological magnetoelectric effect, a.k.a. static axion electrodynamics, were recently reported as quantized magneto-optical effects in TIs \cite{41,42,43}, and quantized magneto and electrical resistance changes in artificial antiferromagnetic heterostructures of magnetically doped TIs \cite{44,45,46}. Finally, dynamical axion quasiparticles (AQ) in the form of magnetic fluctuations were predicted in magnetically doped TIs (MTI) \cite{45}, spin-orbit coupled Mott insulators \cite{47}, in MTI superlattices \cite{48}, and also an inverse chiral magnetoelectric effect was suggested \cite{49}.

Here we propose to use AQs in antiferromagnetically doped TIs (A-TI) as a detector of DAs. The conversion process of DAs to visible photons is shown in Fig. 1(b). Antiferromagnets provide the correct ball-park THz frequency owing to the resonance frequency exchange enhancement $\omega \sim \sqrt{(2H_E + H_A)H_A}$ ($H_E, H_A$ are exchange and anisotropy fields respectively). Inside
FIG. 1. (a) The chiral anomaly [50, 51]. \( \theta \) is a pseudoscalar chirally coupled to charged Dirac fermions, \( \psi \). For the AQ, \( \theta_Q \) is the pseudoscalar part of the spin wave. With applied \( B_0 \), \( \theta_Q \) mixes with \( E \) leading to the existence of axion-polaritons, \( \phi_\perp \). (b) The DA converts to THz photons via the polariton resonance when \( p^2 = \omega_a^2 = \omega_T^2 \). The AF-TI dielectric bound-states refer to Fig. 3. (c) The axion-polariton dispersion relation [35]. The APs are electric and magnetic fields. For the AQ, we rescale the external field \( \nu \) and tune the applied magnetic field \( B \). The parameter space with significant absorption, however, is excluded by astrophysical constraints. For these reasons we neglect the direct nuclear and electron DA couplings.

**Dynamical axion quasiparticles in antiferromagnets:**

The criteria for generating AQSs in condensed matter as hinted by Wilczek are [55]: (i) effective action in the form of Eq. (2) (ii) realization of the Dirac equation for electrons and (iii) tuneable Dirac masses.

Criterion (i) can be met in general in magnetoelectric materials with nonzero diagonal components of the magnetoelectric polarisability tensor \( \alpha_{ij} \) and \( \alpha_{xy} = 0 \). For the AQ, we rescale \( \theta = 1 / f_a \) (Fig. 1) where \( M, P \) are magnetization and electric polarisation. Since \( \theta \) is odd under spatial inversion \( \mathcal{P} \) and time reversal \( \mathcal{T} \), the physical observables \( \sim e^{i S/k} \) are defined modulo \( 2\pi \), the \( \theta \) term can be nonzero in (a) magnetoelectric materials with a magnetic point groups with broken \( \mathcal{P} \), and broken \( \mathcal{T} \) where \( \theta \) is nonquantized, (b) \( \theta = \pi \) can be taken as a defining property of TIs [56, 51].

Criterion (ii) can be realised in Dirac quasiparticle materials such as TIs where the simultaneous presence of \( \mathcal{P} \) and \( \mathcal{T} \) symmetries protects the Kramers double degeneracy of the bulk Dirac bands, while at the surfaces realise \( \mathcal{T} \) protected 2D Dirac quasiparticle helical states [56]. To satisfy (iii) and generate dynamical axion fields, gradients of \( \theta \) need also be generated dynamically, one possibility being magnetic fluctuations [35, 50, 47].

To simultaneously satisfy all three criteria for AQSs we identify Dirac quasiparticle antiferromagnets as suitable candidates [57, 59]. We consider a Dirac antiferromagnetic insulator with \( \mathcal{P} \) and \( \mathcal{T} \) symmetry broken and thus magnetoelectric point group, but importantly the combination \( \mathcal{PT} \) preserved, with a generic electronic Dirac Hamiltonian \( \mathcal{H}(k) = \sum_{i=1,5} A_i(k) \gamma_i \), where \( \gamma_i \) are Dirac matrices and \( A_i(k) \) parameterise the band structure. The antiferromagnetic coupling in proper basis choice to \( \gamma_5 \) in the Dirac Hamiltonian [55, 59].

As a particular lattice realization we consider the antiferromagnetic Fu-Kane-Mele Hubbard model on the bipartite (orbital degree of freedom \( \tau \) ) diamond lattice with two spins per lattice site \( \sigma \) [59, 61]. The antiferromagnetism breaks \( \mathcal{T} \), and \( \mathcal{P} \), but preserves \( \mathcal{PT} \) as marked by the red ball in Fig. 2(b) and thus preserves the form of the Dirac Hamiltonian. The Hamiltonian with a Hubbard term treated on a mean-field level reads:

\[
\mathcal{H} = \lambda \left( \mathbf{A}(k) - \frac{U \mathbf{m}}{\alpha} \right) \cdot \sigma \tau_z + t \text{Re} f(k) \tau_x + t \text{Im} f(k) \tau_y,
\]

where

\[
S_{\text{CS}} = \sum_{i=\text{D},Q} \frac{\alpha}{\pi} C_i \int d^4 x \theta_i \mathbf{E} \cdot \mathbf{B},
\]
the nearest neighbour hopping on the diamond lattice (cf. Fig. 2(a)) \( f(k) = \sum_{j=1}^{4} (t + \delta t_j) e^{i k d_j} \) (\( d_j \) being the four nearest neighbour vectors), \( A_x(k, U m_x) = 4 \sin k_x^2 + (\cos k_x^2 - \cos k_y^2) \) plus cyclic permutations, \( U \) is Hubbard correlation strength, and \( \delta t_j \) represent the renormalization of the hopping due to the deformation of the AB bond. The AQ has a mean value given by \( \theta_Q = \frac{\pi}{2} [1 + \text{sign}(\delta t_1)] - \arctan \left( \frac{U m_x}{Q} \right) \) [59]. It was shown that the fluctuations in the Néel order parameter \( \mathbf{L} \sim \mathbf{m}_A - \mathbf{m}_B \) (with axis along \( z \)) can be within approximation \( \frac{U m_i}{Q} << 1 \) related to dynamical fluctuations of \( \theta_Q \) [35, 69]:

\[
\delta \theta_Q \sim \frac{2}{3} \sum_{i=x,y,z} U m_i .
\]  

The band structure of our model is shown in Fig. 2(a) for a realistic range of effective exchange coupling \( U m \sim 0-0.40 \) and illustrates the tuning of the Dirac bands with a Dirac point shifted slightly off the \( X \) \((\sim U m/2\lambda)\) point due to the effect of antiferromagnetism. The AQ spin wave (SW) [62, 64] dispersion on the diamond lattice is

\[
\hbar \omega_{Q,1} \approx g \mu_B H_0 \pm \sqrt{(8S J f(0) + g \mu_B H_A)^2 - (8S J f(q))^2},
\]  

where \( g \approx 1 \) is the Landé factor and \( H_E = 8S J \), and \( q \) is the spin wave wave-vector. The AQSW tunes, in a first-order approximation, only the \( z \)-component of the \( L \)-order parameter [35, 59], which therefore tunes the Dirac mass as schematically illustrated by the shaded region in Fig. 2(a).

**Estimation of A-TIs parameters:** No antiferromagnetic bulk dynamical axionic insulator has been yet identified in the lab. Remarkably, however, our model has exactly the same magnetic point group, 6\( \bar{4} \)\( m' \), as the mean-field medium of Fe-doped Bi\(_2\)Se\(_3\). This can be seen by deforming the face centred cubic primitive unit cell (Fig. 2(a)) along the [111] direction to produce the rhombohedral unit cell of tetrahedrally Bi\(_2\)Se\(_3\) (Fig. 2(b)). It can be shown that the antiferromagnetism couples to the same \( \gamma_5 \) matrix as in our model [35], and applying the Neumann principle gives axion-field favourable nonzero diagonal symmetric elements to \( \alpha_{ij} \), and leads to the analogous expression for the AQ spin-wave field, Eq. (3).

The quadratic action for small fluctuations \( \theta_Q \) is given by:

\[
S_{AQ} = \frac{f_Q^2}{2} \int d^4 x \left[ \dot{\theta}_Q^2 - (v_s, \partial_t \theta_Q)^2 - m_A^2 \theta_Q^2 \right],
\]  

where \( f_Q \) in this context quantifies the stiffness of the spin wave. Scanning \( \omega_\pm(B_0) \) (see Fig. 4(c)) requires specifying \( m_A(B_0) \) and \( f_Q(B_0) \). For (Bi\(_{1-x}\)Fe\(_x\))\(_2\)Se\(_3\) using Eq. (1) with doping factor at 3.5% [65], exchange of 1 meV [66] and anisotropy of 16 meV [67], the spin wave mass is \( m_A = [0.12(B_0/2 \text{T}) + 0.6] \) meV. From Ref. [35] we find \( f_Q = 190 \) eV at \( B_0 = 2 \) T, and take \( f_Q^2 \approx 1/m_A \) from the \( \delta \mathbf{L} \) kinetic term in Ref. [40].

**Equations of Motion and driven axion-polariton:** Including the usual Maxwell term, linearizing for small fluctuations in \( E \) and \( \theta_Q \) in the presence of an applied magnetic field, \( B_0 \), and external DA source, we find the system of equations derived from the action take the form (see also Refs. [40, 55, 68]):

\[
\epsilon \dot{\mathbf{E}} - \nabla \mathbf{E} + \frac{\alpha}{\pi} [\mathbf{E} \cdot \nabla \theta_Q - \nabla (\nabla \theta_Q \cdot \mathbf{E})] = \mathbf{A} \cos \omega_\pm t ,
\]

\[
\dot{\theta}_Q - v_s^2 \dot{\theta}_Q + m_A^2 \theta_Q - \frac{\alpha}{4 \pi^2 f_Q^2} B_0 \cdot \mathbf{E} = 0 ,
\]  

where \( \epsilon = \epsilon_{\text{Ti}} \) is the Ti dielectric constant, and \( v_s \) is the spin wave speed.

The driving term \( \mathbf{A} = 2 \mathbf{B}_0 \gamma_\gamma \sqrt{2/\text{DM}} \) at leading order, and derives from the DA Chern-Simons action, taking the DA as an external DM source, with \( \theta_D \) fixed as described above. Neglecting the AQ dispersion compared to \( E \), we diagonalize Eq. (6) to find \( \phi_\pm \) and \( \omega_\pm^2(k) = (k^2/\epsilon^2 + m_A^2 + b^2) \pm \sqrt{(k^2/\epsilon^2 + m_A^2 + b^2)^2 - 4 k^2 m_A^2/\epsilon^2} \), where \( b^2 = \alpha^2 B_0^2/4 \pi^2 f_Q^2 \) [35], as shown in Fig. 1(c). Similar A-Qs are required for the mixing: in the absence of derivatives, \( \theta_Q \) and \( E \) decouple in Eq. (6). The presence of axion-polaritons can be verified using an inverse “light shining through a wall” [69] experiment: AF-TIs have non-zero transmission to applied lasers, except in the gap frequencies where they are opaque [35]. This can be used to measure the gap size \( \omega_\pm(0) \).

**Detecting Dark Axions from Resonant Conversion:** DA-driven polariton waves in the A-TI are a combination of \( \mathbf{L} \), and \( \mathbf{E} \). When \( T \) symmetry is preserved, TIs...
where $Q$ introduce effective losses, and treat Eqs. (6) for experiments [73], $P$ evidently enhances the forward emission [49]. The concept is detector [72]. A mirror placed behind the AF-TI coherently enhances the forward emission [49]. We have conducting surface states [70]. In the presence of $T$ breaking, the surface states become gapped [35], i.e. insulating. Thus the polariton $E$-field leads to emission of DA-induced photons from the surface of the A-TI, just like a dielectric haloscope, or dish antenna [29, 71]. We propose to detect the emitted photons by using a silicon lens to focus them onto a single photon quantum dot detector [72]. A mirror placed behind the AF-TI coherently enhances the forward emission [49]. The concept is illustrated in Fig. 3.

The power generated by DA-photon conversion can be expressed in analogy to the case of microwave cavity experiments [73], $P_{\text{signal}} = (\omega/Q) \times (\text{energy stored})$. We introduce effective losses, and treat Eqs. (6) for $\theta_Q$ and $E$ classically (the quantum calculation is equivalent [74]) in order to estimate the power output on resonance in steady state:

$$P_{\text{signal}} = \kappa f_+ Q_{\text{sys}} P_0,$$

where $Q_{\text{sys}}$ is the effective loaded quality factor of the system, $\kappa$ is a coupling factor, and $f_+ = b^2/(\omega^2 + b^2)$ is a mode mixing factor. The effective $Q$ is due to the electric field enhancement inside the A-TI due to the resonance in Eq. (6). We assume $Q_{\text{sys}} = 10^2$ (an estimate based on AFMR line widths) such that the polariton in the A-TI is optimally coupled to the free space electromagnetic field at the surface for efficient photon measurement, and material losses due to Gilbert damping and phonon production (additional decay channels in Fig. 1) are of order the photon emission. We absorb into $V_{\text{eff}}$ (see Eq. 1) the relevant form factors, the effect of the A-TI dielectric constant, and any boost factor, $\beta^2$, arising from the geometry [49]. For reference values of $V_{\text{eff}} = 1 \text{ cm}^3$, $g_z = 10^{-10} \text{ GeV}^{-1}$, $\omega_a = 2 \text{ meV}$, $b = 0.5 \text{ meV}$, and assuming $\beta^2 \kappa = 1$ (which could be engineered by appropriate use of coatings and geometry), the total power output is $5 \times 10^{-24} \text{ W}$: about one photon every minute.

Because of the low photon flux, photon detection based on single-photon detector will be more advantageous than power detection [75]. A high confidence level of DA detection requires the dark count rate, $\Gamma_d$, of the single-photon detector to be much smaller than the photon flux. We use $\Gamma_d = 0.001 \text{Hz}$, which has been demonstrated for the quantum dot detector in THz regime at 0.05 K [72]. A wider bandwidth, lower dark count single-photon detector using graphene-based Josephson junction [70] may further improve the sensitivity in the future.

The range of axion masses accessible to our technique depends on the scaling of material properties with $B_0$, and the attainable range of $B_0$. We take $1 \text{ T} < B_0 < 10 \text{ T}$ with stability $\delta B_0 = 10^{-3} \text{T}$ over the volume, which has been demonstrated [18] [74]. For the parameters of (Bi$_{1-x}$Fe$_x$)$_2$Se$_3$ given above and setting $\omega_a(k=0) = m_a$ we find $0.7 \text{ meV} \leq m_a \leq 3.5 \text{ meV}$, the lower limit being set by the $B_0 = 0$ spin wave mass in the A-TI. Other materials with different anisotropy field strengths can be used to cover a wider range of masses.

Sensitivity to $g_z$ is computed from the signal to noise ratio (SNR), SNR = 3. We take the measurement time on a single frequency $\tau = 10^4 \Gamma_d^{-1} = 10^2 \text{s}$, which allows the full range to be covered in approximately 6 months scan time. The volume of any single, high quality, sample of A-TI is limited to be slightly less than approximately $1 \text{ cm}^3$ to achieve homogenous doping [78]. The sensitivity for $V_{\text{eff}} = 1 \text{ cm}^3$ is shown in Fig. 4 (stage-I).

$V_{\text{eff}}$ can be increased by using multiple A-TI samples. This requires coherent addition similar to methods to boost DA signals from dielectrics [29]. The gain in $V_{\text{eff}}$ can increase linearly with the number of samples, $N$, with wide band response [49], and with $N = 100$ (a feasible total number for solid state synthesis [79]), the sensitivity can be increased as shown in Fig. 4 (stage-II).

A further increase in $V_{\text{eff}}$ can be achieved by surrounding the A-TI samples with a cavity. Permeation of the cavity electric field into the A-TI couples long wavelength cavity modes to the high frequency axion-polariton, so that the axion-polariton has a cavity $E$-field component of TM$_{010}$ type [30]. This method can potentially boost $V_{\text{eff}}$ up to some large fraction of the volume of the cavity, while keeping the A-TI sample volume much smaller. For illustration of this method, we take $V_{\text{eff}} = (0.1 \lambda_{\text{th}})^3$ giving sensitivity as shown in Fig. 4 (stage-III).

In summary, we have shown that A-TIs can host dynamical axionic quasiparticles which can be resonantly driven in the presence of DAs with mass of order 1 meV and emit THz photons which can be detected using a quantum dot photon counter, allowing A-TIs to detect dark matter. We specifically show that antiferromagnetic Fe-doped Bi$_2$Se$_3$ satisfies the three Wilezczek criteria described earlier, and can be used to realize a DA detector.
in the 0.7 to 3.5 meV range. Fig. 4 shows the projected reach of three possible schemes with different effective volumes. Varying the applied $B$ field scans the resonant frequency, giving sensitivity to axion dark matter in a parameter space inaccessible to other methods. Future work on optimizing the material characteristics (such as $T_N$ and anisotropy field strength) can allow for a wider range of DA mass detection.

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![Axion parameter space](image)

**FIG. 4.** Axion parameter space. Vertical lines show the projected sensitivity of our proposal using Fe doped Bi$_2$Se$_3$ at $\sim$5T applied field for $10^5$ s integration time with dark count rate $\Gamma_d = 0.001$ Hz. Staged designs are described in the text. Gray shaded regions assume scanning $1 \text{T} \leq B_0 \leq 10 \text{T}$. The KSVZ and DFSZ axion models are shown as the red band. Existing exclusions from ADMX [17][18], CAST [27], and supernova 1987A [23], are shown as coloured regions.

| Axion Mass, $m_a \text{[eV/c}^2\text{]}$ | Axion Coupling $|g_\gamma| \text{[GeV}^{-1}\text{]}$ |
|-------------------------------------|----------------------------------|
| $10^{-2}$                           | $10^{-12}$                        |
| $10^{-3}$                           | $10^{-13}$                        |
| $10^{-4}$                           | $10^{-14}$                        |
| $10^{-5}$                           | $10^{-15}$                        |
| $10^{-6}$                           | $10^{-16}$                        |

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