

Hunting axion dark matter with antiferromagnets: A case study with nickel oxide

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We show that nickel oxide, which is already a very promising target to look for sub-MeV dark matter scattering, can be employed to hunt axion dark matter, with masses in the meV range and couplings to electrons allowing them to potentially be QCD axions. We describe the interactions between axions and the collective excitations of nickel oxide in terms of an effective field theory, built solely out of symmetry arguments. The processes of conversion into one or two excitations provide, respectively, a narrow band and a broadband channel for the axion search, and the possibility of varying an external magnetic field up to a phase transition point allows one to cover a large portion of a yet unexplored parameter space, reaching axion masses down to few fractions of an meV. Our results underline nickel oxide as an ideal candidate for a multipurpose target for light dark matter searches.

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I. INTRODUCTION

Direct dark matter searches for both heavy weakly interacting massive particles [e.g., [1–5]] and axions [e.g., [6–16]], have so far reported null results. Concerning the axion, in particular, only a handful of experiments have (or will have) enough sensitivity to probe the QCD axion region, i.e., the right masses and couplings to potentially offer a solution to the strong CP problem. In order to keep expanding the scope of our searches, and probe yet unexplored regions in parameter space, new ideas and detection strategies are needed.

A good deal of work has been recently devoted toward taking advantage of the rich spectrum of collective excitations featured by different condensed matter systems, trying to employ them to look for sub-MeV dark matter scattering and axion absorption, probing new values of their possible mass [e.g., [17–41]]. For a review, see [42,43].

In this work we show that nickel oxide (NiO), a well studied antiferromagnet [e.g., [44–46]], offers an optimal target to probe the electron coupling of axion dark matter

with masses in the unexplored range $m_a \sim \mathcal{O}(0.1\text{--}10)$ meV. (Other proposals valid for similar masses are [33,34,40].) It has already been shown that NiO provides an ideal target also to look for sub-MeV dark matter scattering from spin-dependent interactions [39]. Among other reasons, this is because its vanishing magnetization ensures that the emission of multiple collective excitations (the so-called magnons) is allowed. This enables the transfer of a large fraction of the dark matter energy even at masses as low as the keV.

As we will argue, NiO is also very efficient at absorbing axions, with masses in the same range as the typical energies of its magnon modes. In particular, the process of axion conversion into a single excitation provides (at least) two resonant values for a narrow band search, corresponding to the two lowest lying magnon modes. By tuning an external magnetic field one can also make one of the magnons arbitrarily light, giving access to axion masses even below the meV range. The conversion into two magnons, instead, offers the possibility for a *broadband* channel: the system is able to absorb axions in a wide range of masses by simply adjusting the total energy and relative momentum of the two outgoing excitations.

We work within an effective field theory (EFT) for the description of magnons and their interactions, valid at low energies and long wavelengths [47–50], which we extend to incorporate the rich phenomenology of NiO [51]. In this framework, magnons are the (pseudo-)Goldstone bosons associated with the breaking of nonrelativistic spin symmetry, as induced by the antiferromagnetic ground state. The main advantage of this EFT is that it is completely

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dictated by symmetry, thus bypassing the need for a microscopic description of the system. The intricacies of the short distance physics are encoded in a few effective coefficients, which we match from data. This has a twofold advantage. On the one hand, the EFT formulation allows one to clearly disentangle those phenomena that are universal, i.e., common to a broad class of systems, to those that are specific to a given short distance realization of these systems. On the other hand, it helps formulating the current problem in a language that is closer to high energy physics, thus allowing one to borrow concepts and techniques from this field.

Conventions: We use natural units, $\hbar = c = 1$, and follow an index notation with $i, j, k = 1, 2, 3$ and $a, b = 1, 2$.

II. THE EFT

We start by summarizing the EFT for magnons in NiO, which has been developed in full detail in [51]. After that, we introduce the axion, and describe how to incorporate it into the above mentioned EFT. Throughout our treatment, we will assume to work at zero temperature, as motivated by the standard setup for direct dark matter searches.

A. Magnons in NiO

A magnetic material can be thought of as a collection of spins arising from the elementary spins of the unpaired electrons localized around each lattice site, with the exchange and superexchange interaction between these electrons providing an effective coupling between different spins. In an antiferromagnet, these couplings are such that, below a certain critical temperature, the ground state of the system develops a long-range order, with neighboring spins antiparallel to each other and all aligned along one direction. We take this to be the z axis. For NiO, this happens below the temperature $T_N \simeq 523$ K, called the Néel temperature [e.g., [44,52]].

In this context, one can define an order parameter, the so-called Néel vector, consisting of a staggered sum of all the spins in the material, i.e., $\mathcal{N} \equiv \sum_i (-1)^i \mathbf{S}_i$, where $\mathbf{S}_i$ is the spin of the i -th lattice site and $(-1)^i$ is positive for spins on even sites, and negative for those on odd sites. In the ordered ground state, the Néel vector acquires a nonzero expectation value, $\langle \mathcal{N} \rangle \propto \hat{z}$. From the symmetry viewpoint, this spontaneously breaks the internal spin rotations to the subgroup of rotations that leave the order parameters unchanged, i.e., those around the z axis. The symmetry breaking pattern is then $\text{SO}(3) \rightarrow \text{SO}(2)$, and magnons are nothing but the associated Goldstone bosons. As such, they can be described by an EFT.

A convenient way of parametrizing them is in terms of local modulations of the order parameter, $\hat{\mathbf{n}} = e^{i\theta^a S_a} \cdot \hat{z}$, where $S_a = \{S_1, S_2\}$ are the broken $\text{SO}(3)$ generators, with $(S_i)_{jk} = -i\epsilon_{ijk}$, and $\theta^a(x)$ are the magnon fields. The most

TABLE I. Parameters for the EFT of magnons in NiO [51]. μ_B is the Bohr magneton, and the mass density, ρ_T is taken from [53]. Λ_{UV} is estimated from [44] as the momentum for which the dispersion relation deviates from linear by 10%. Note that while the EFT is general for any antiferromagnet with the same symmetry breaking pattern, the value of the parameters crucially depends on the microscopic structure of the order parameter and of the crystal.

v_θ	c_1 [MeV/Å]	μ/μ_B	λ_x [meV ²]	λ_z [meV ²]	Λ_{UV} [keV]	ρ_T [g/cm ³]
1.3×10^{-4}	0.58	2.18	9.80	0.17	0.6	6.7

general, lowest order Lagrangian describing magnons at low energies and small momenta is simply given by [e.g., [48]]

$$\mathcal{L}_\theta = \frac{c_1}{2} [(\partial_t \hat{\mathbf{n}})^2 - v_\theta^2 (\nabla_i \hat{\mathbf{n}})^2]. \quad (1)$$

Here v_θ is the magnon propagation speed, and it can be determined from their dispersion relation, obtaining $v_\theta \simeq 1.3 \times 10^{-4}$ [39]. The coefficient c_1 can instead be matched to the parameters of the microscopic spin Hamiltonian using neutron scattering data [39]. We report the values of all the effective coefficients in Table I.

The phenomenology of NiO is, however, richer than this. Specifically, its magnon modes actually present a small but nonzero gap. This is due to tiny explicit breaking effects, turning the magnons into *pseudo*-Goldstones. To the best of our knowledge, the microscopic mechanism generating these gaps has not been completely understood. Data from Raman and Brillouin scattering experiments seem to point to magneto-crystalline anisotropies and dipolar interactions between spins as the most plausible explanation [45]. Yet, the simplest model which correctly reproduces the spectrum of the two lowest lying magnon modes only includes two anisotropic terms with no dipole interactions, as done, for example, in [44,46]. These anisotropies explicitly break the internal $\text{SO}(3)$, and are encoded in the EFT by two additional operators [51],

$$\mathcal{L}_{\text{anis}} = c_1 [\lambda_z \hat{n}_z^2 - \lambda_x \hat{n}_x^2], \quad (2)$$

where $\lambda_z, \lambda_x > 0$. The anisotropic couplings can also be matched to the parameters of the microscopic Hamiltonian [51], and their values are again in Table I. The anisotropies in the x and z directions are often dubbed “hard” and “easy” axis anisotropies, respectively [46].

In order to scan different axion masses, one needs to vary the magnon gap. This can be achieved by placing the NiO sample in a constant magnetic field, which we take to be aligned with the easy axis, $\mathbf{H} = H\hat{z}$. This can be accounted for in the EFT by promoting the ordinary time derivative to a covariant one in the Lagrangian (1), $\partial_t \hat{\mathbf{n}} \rightarrow \partial_t \hat{\mathbf{n}} + \mu \mathbf{H} \times \hat{\mathbf{n}}$,

such that the Lagrangian is now invariant under simultaneous gauge transformations of $\hat{\mathbf{n}}$ and $\mathbf{H}$, with the latter playing the role of a vector potential. Indeed, at the level of the Hamiltonian, this procedure returns the standard Zeeman coupling [50,51,54]. Here μ is the gyromagnetic ratio, again given in Table I.

All together, the EFT for magnons in NiO is given by

$$\begin{aligned}\mathcal{L}_{\text{EFT}} &= \frac{c_1}{2} [(\partial_t \hat{\mathbf{n}} + \mu \mathbf{H} \times \hat{\mathbf{n}})^2 - v_\theta^2 (\nabla_i \hat{\mathbf{n}})^2] + \mathcal{L}_{\text{anis}} \\ &= \frac{c_1}{2} [(\dot{\theta}^a - \mu H \epsilon^{ab} \theta^b)^2 - v_\theta^2 (\nabla \theta^a)^2 \\ &\quad - 2\lambda_z (\theta^a)^2 - 2\lambda_x \delta^{a2} \delta^{b2} \theta^a \theta^b + \mathcal{O}(\theta^4)],\end{aligned}\quad (3)$$

where we expanded in fluctuations around equilibrium.

From the solutions to the linearized equations of motion one finds that there are, as expected, two physical magnon modes, with dispersion relations given by [51]

$$\begin{aligned}\omega_{q,\pm}^2 &= \mu^2 H^2 + \lambda_x + 2\lambda_z + v_\theta^2 q^2 \\ &\quad \pm \sqrt{4\mu^2 H^2 (\lambda_x + 2\lambda_z + v_\theta^2 q^2) + \lambda_x^2}.\end{aligned}\quad (4)$$

The quantization of this theory is complicated by the presence of a one-derivative kinetic term, which does not allow one to diagonalize the quadratic Lagrangian (3) by a local field redefinition. The details of the quantization procedure have been developed in [51].

Importantly, as the magnetic field increases, the gaps of the $+$ and $-$ modes in Eq. (4) respectively increase and decrease. This feature can be used to scan over different axion masses. At $\mu_0 H = \mu_0 \sqrt{2\lambda_z}/\mu \simeq 4.62$ T,¹ the lightest modes becomes gapless and, for larger fields, the system undergoes a (first order) “spin-flop” phase transition to a different ground state, characterized by $\langle \hat{\mathbf{n}} \rangle = \hat{\mathbf{y}}$ [e.g., [46,51]]. In this work, we are not interested in this regime. However, we notice that thermal and radiative corrections are expected to change both the magnon dispersion relation, as well as the value of the critical magnetic field, but not too dramatically [55]. They are also expected to induce a small gap even at the phase transition point.

It is also interesting to note that in the absence of magnetic fields, the magnon gaps are $\omega_{0,+/-} \sim \sqrt{\lambda_x/\mu}$, corresponding to $\mathcal{O}(\text{meV})$ energies. This is in contrast to ferrimagnets [33], where the gaps are $\mathcal{O}(\mu\text{eV})$, and hence are more suitable to probe axion masses in that range.

Finally, our EFT breaks down for momenta larger than a strong coupling scale, Λ_{UV} , when the details of microscopic system become relevant. This cutoff can be estimated, for example, as the momentum for which the dispersion relation sensibly deviates from linearity, indicating that higher

derivative terms in the effective Lagrangian become relevant. In the rest of this work, we always ensure that momenta and energies are below, respectively, Λ_{UV} and $v_\theta \Lambda_{\text{UV}}$, the latter being the typical microscopic energy scale.

B. Axion-magnon coupling

We are interested in the axion-electron coupling, which is ultimately responsible for the interaction between axions and magnons. The corresponding Lagrangian is

$$\begin{aligned}\mathcal{L}_a &= \frac{g_{aee}}{2m_e} \partial_\mu a \bar{e} \gamma^\mu \gamma^5 e \xrightarrow{\text{n.r.}} -\frac{g_{aee}}{m_e} \nabla a \cdot \left(e_{\text{nr}}^\dagger \frac{\boldsymbol{\sigma}}{2} e_{\text{nr}} \right) \\ &\xrightarrow{\text{IR}} -\frac{g_{aee}}{m_e} \nabla a \cdot \mathbf{s},\end{aligned}\quad (5)$$

where we first took the nonrelativistic limit as described in [39] (see also [56]), with e_{nr} the nonrelativistic electron field, neglecting terms suppressed by the ratio of the axion mass to the electron mass. In going to the second line, we used the fact that $e_{\text{nr}}^\dagger \boldsymbol{\sigma} e_{\text{nr}}/2$ is the electron spin density. At low energies, it must be expressed in the terms of the relevant degrees of freedom, which in this case are the magnons. Note that we have neglected an additional term, which scales as the electron velocity. For most materials, this is subleading with respect to the coupling included here. For a more detailed discussion we refer the reader to [40,57]. The inclusion of this additional effect is beyond the scope of the present work.

Now, the spin density is nothing but the Noether current associated with the original $\text{SO}(3)$ symmetry which, given the Lagrangian (3), reads as

$$\begin{aligned}s^i &= c_1 [(\partial_t \hat{\mathbf{n}} \times \hat{\mathbf{n}})^i + \mu (\mathbf{H} \cdot \hat{\mathbf{n}}) \hat{n}^i] \\ &\simeq c_1 [\delta^{ia} (\dot{\theta}^a - \mu H \epsilon^{ab} \theta^b) - \delta^{i3} \theta^a (\epsilon^{ab} \dot{\theta}^b + \mu H \theta^a)],\end{aligned}\quad (6)$$

where we omitted constant terms and neglected higher order contributions in the small fluctuations.

With this at hand, and given the expression for the canonical fields reported in [51], one easily finds the amplitudes, $i\mathcal{M}$, for the axion-magnon conversions:

$$p \rightarrow q, \alpha = i \frac{g_{aee} \sqrt{c_1}}{m_e} p_a [i\omega_{q,\alpha} \mathcal{Z}_{q,\alpha}^a + \mu H \epsilon^{ab} \mathcal{Z}_{q,\alpha}^b]^*, \quad (7a)$$

$$\begin{aligned}p \rightarrow q, \alpha &= i \frac{g_{aee}}{m_e} p_z [i(\omega_{q,\alpha} - \omega_{k,\beta}) \epsilon^{ab} \mathcal{Z}_{q,\alpha}^a \mathcal{Z}_{k,\beta}^b \\ &\quad + 2\mu H \mathcal{Z}_{q,\alpha}^a \mathcal{Z}_{k,\beta}^b]^*,\end{aligned}\quad (7b)$$

where $\alpha, \beta = \pm$ run over the physical magnon states, with dispersion relations as in Eq. (4). The coefficients

¹In the presence of a material, the magnetic field H is typically reported in oersted (Oe). In this work, we use $\mu_0 = 10^{-4}$ T/Oe to convert to Tesla, which is more common in the high energy literature.

$\mathcal{Z}_{q,\alpha}^a \equiv \langle 0 | \theta^a(0) | \mathbf{q}, \alpha \rangle$ are the overlap functions between the magnon fields and the physical magnon states. They are given by [51]

$$|\mathcal{Z}_{q,\pm}^1| = \sqrt{\frac{1}{2} \pm \frac{2\mu^2 H^2 - \lambda_x}{2\sqrt{4\mu^2 H^2 (v_\theta^2 q^2 + \lambda_x + 2\lambda_z) + \lambda_x^2}}}, \quad (8a)$$

$$\mathcal{Z}_{q,\pm}^2 = i \left(\frac{v_\theta^2 q^2 - \mu^2 H^2 + 2\lambda_z - \omega_{q,\pm}^2}{2\mu H \omega_{q,\pm}} \right) \mathcal{Z}_{q,\pm}^1, \quad (8b)$$

where the common overall phase is arbitrary.

III. EVENT RATES

We can now compute the axion absorption event rates per unit target mass. For a target material of density ρ_T (see Table I), this is given by

$$R(\hat{\mathbf{v}}_e) = \frac{\rho_a}{\rho_T m_a} \int d^3v f(|\mathbf{v} + \mathbf{v}_e|) \Gamma(\mathbf{v}). \quad (9)$$

Here $f(v)$ is the dark matter velocity distribution in the Milky Way. To compare with other proposals, we take this as a truncated Maxwellian as in the standard halo model, with dispersion $v_0 = 230$ km/s, escape velocity $v_{\text{esc}} = 600$ km/s, and boosted with respect to the Galactic rest frame by the Earth velocity, $v_e = 240$ km/s [58,59]. We also take the axion dark matter density to be $\rho_a = 0.4$ GeV/cm³ [58,60].

One of the advantages of employing an EFT approach like we do, and thus formulating the problem in terms of local Lagrangians, lies in the fact that the rate Γ can now be computed in the standard way, as explained in any particle physics textbook.

A. One-magnon absorption and directionality

For the conversion into a single magnon, momentum conservation enforces its momentum to be equal to the axion one, $\mathbf{q} = m_a \mathbf{v}$. Since we will be interested in masses $m_a \lesssim 1$ eV, this momentum is always extremely small, and can be set to zero everywhere without any appreciable corrections, as one can check explicitly. In light of this, the absorption rate is

$$\begin{aligned} \Gamma_\alpha &= \frac{\pi m_a |g_{aee}|^2}{2m_e^2} \frac{c_1}{\omega_{0,\alpha}} |v^a V_\alpha^a|^2 \delta(m_a - \omega_{0,\alpha}) \\ &\rightarrow \frac{m_a |g_{aee}|^2}{m_e^2} \frac{c_1}{\omega_{0,\alpha}} |v^a V_\alpha^a|^2 \frac{m_a \omega_{0,\alpha} \gamma_\alpha}{(m_a^2 - \omega_{0,\alpha}^2)^2 + (\omega_{0,\alpha} \gamma_\alpha)^2}, \end{aligned} \quad (10)$$

where $\mathcal{V}_\alpha^a \equiv i\omega_{0,\alpha} \mathcal{Z}_{0,\alpha}^a + \mu H \epsilon^{ab} \mathcal{Z}_{0,\alpha}^b$, and we regulated the energy δ function with a Breit-Wigner distribution, accounting for the magnon finite width, γ_α . This width is roughly $\gamma_\alpha = \mathcal{O}(1-100)$ μeV [61,62], and it depends on various properties of the material, such as Umklapp

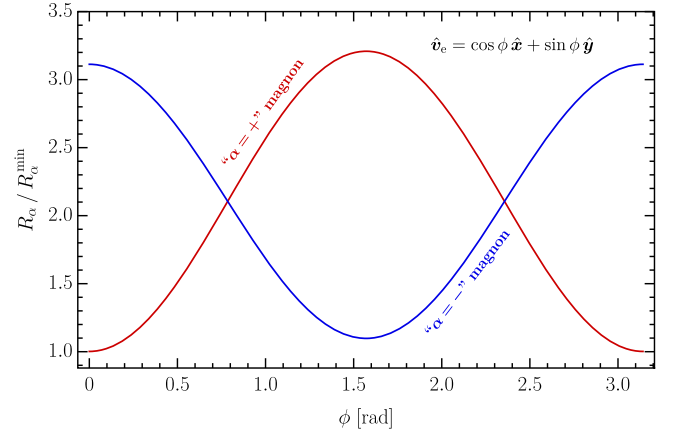


FIG. 1. One-magnon event rate as a function of the direction of the Earth velocity, taken to be on the xy plane, i.e., $\hat{\mathbf{v}}_e = \cos \phi \hat{\mathbf{x}} + \sin \phi \hat{\mathbf{y}}$, at $\mu_0 H = 1$ T. The rate is normalized to its minimum value, achieved for $\hat{\mathbf{v}}_e = \hat{\mathbf{z}}$. Note that the minima of the curve lie very close to $R_+/R_+^{\min} = 1$. This means, for example, that the R_+ rates evaluated at $\mathbf{v}_e = \hat{\mathbf{z}}$ and $\mathbf{v}_e = \hat{\mathbf{x}}$ are very similar. This is because the V_α^a vectors have one component much larger than the other, as a consequence of the large hierarchy between the effects due to the external magnetic field, and those due to the intrinsic anisotropies.

scattering, temperature, and boundary effects [62]. To the best of our knowledge, this information is not available for NiO. We thus follow the conservative approach of [33] and treat $\gamma_\alpha/\omega = \gamma_\alpha/m_a$ as a free parameter.

The event rate has a nontrivial dependence on the direction of the Earth velocity. This is because antiferromagnets feature a preferential direction, as set by the order parameter. This directionality can be exploited as an efficient strategy to discriminate between signal and background [e.g., [63,64]]. As we show, in NiO this feature is quite remarkable, with a modulation following the direction of the Earth velocity of about 300%.

After averaging Eq. (10) over the dark matter velocity distribution, the rate depends on the Earth velocity as

$$R_\alpha(\hat{\mathbf{v}}_e) \propto \frac{\langle \mathbf{v}_\perp^2 \rangle}{2} V_\alpha^a (V_\alpha^a)^* + |v_e^a V_\alpha^a|^2, \quad (11)$$

where $\mathbf{v}_\perp$ is a velocity perpendicular to the order parameter (i.e., on the xy plane), and directionality is due to the second term. If the axion has the right mass to emit an $\alpha = +$ magnon, the rate is maximum when $\hat{\mathbf{v}}_e = \hat{\mathbf{y}}$, while for $\alpha = -$, this happens when $\hat{\mathbf{v}}_e = \hat{\mathbf{x}}$. In Fig. 1 we report the directional dependence of the event rate. As anticipated, the signal modulation is very strong. However, we point out that realistic materials present various domains, each with a different orientation of the order parameter. This effect could decrease the modulation.²

²We are grateful to G. Marocco for pointing this out.

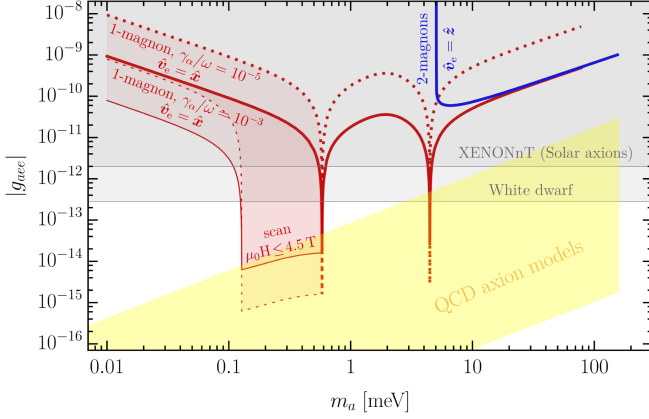


FIG. 2. Projected reach at 95% CL for a kilogram of material and a year of exposure, assuming no background. For single magnon absorption, we treat γ_a/ω as a free parameter. The one- and two-magnon channels have been computed, respectively, for $\hat{\mathbf{v}}_e = \hat{\mathbf{x}}$ and $\hat{\mathbf{v}}_e = \hat{\mathbf{z}}$. The red shaded regions are obtained for a magnetic field varying in the range $\mu_0 H \in [0, 4.5]$ T. The same scan for the higher magnon mode is not visible in logarithmic scale. The projections are interrupted when the energy/momenta involved exceed the EFT cutoff. The yellow band corresponds to the QCD axion models reviewed, for example, in [65,66], and is stopped in correspondence to the cosmological bound determined in [67]. The white dwarf bound is taken from [68], and the XENONnT one from [14] (see also [69]).

In Fig. 2 we instead show the projected reach obtained assuming three events per year for a kilogram of NiO, in the ideal absence of background. The total rate is given by $R = R_+ + R_-$, and we evaluate it in the $\hat{\mathbf{v}}_e = \hat{\mathbf{x}}$ configuration. The result obtained for the $\hat{\mathbf{v}}_e = \hat{\mathbf{y}}$ is analogous, differing by no more than order 1. We show results for both $\gamma_a/\omega = 10^{-3}$ and 10^{-5} , with the second one being rather optimistic and likely achievable only in very pure samples. We also scan different values of the external magnetic field, in the range $\mu_0 H \in [0, 4.5]$ T. Crucially, when this approaches the critical value (see Sec. II A), the one-magnon process probes lower and lower values of the axion mass, as again shown in Fig. 2. This makes the H field scanning in NiO potentially a very powerful avenue to cover a large portion of the parameter space.

B. Two-magnon absorption

As explained also in [39], antiferromagnets offer the possibility of multimagnon events, which are instead strongly suppressed by symmetry in ferrimagnets [27]. In the context of axion absorption, these allow for a broadband channel, since the system can absorb axions with different masses by simply changing the energy and relative momentum of the multiple excitations, without the need for any external tuning. In our EFT language, the matrix element for two-magnon absorption is in Eq. (7b).

In this case the magnons are effectively emitted back to back, $\mathbf{p} \simeq -\mathbf{k}$, since again the axion momentum is

irrelevant. This considerably simplifies the calculation of the event rate, which is now the sum over all possible outgoing magnons $R = R_{++} + R_{+-} + R_{-+} + R_{--}$. Moreover, contrary to the one-magnon case, the maximum rate in this case is obtained for $\hat{\mathbf{v}}_e = \hat{\mathbf{z}}$. The resulting projections are again reported in Fig. 2. As one can see, the absorption into two magnons provides a broadband channel, which can be dominant when the magnons are particularly long lived.

IV. OUTLOOK

NiO can be employed as an optimal target to probe the electron coupling of axion dark matter with masses around the meV and couplings down to the QCD axion ones, a region still experimentally uncharted. The one-magnon channel features a strongly directional signal, which can be key to an efficient background rejection strategy. Moreover, by varying the external magnetic field toward its critical “spin-flop” value, it is possible to probe axion masses that are, in principle, arbitrarily small, thus covering a large portion of the unexplored parameter space. Finally, the excitation of multiple magnons in the final state provides a broadband channel for axion absorption, which is the dominant signature in samples that allow for an especially narrow magnon mode.

Our results, together with the previous proposal to employ NiO as a target for sub-MeV dark matter scattering from spin-dependent interactions [39], make a strong physics case for the employment of this material as a *multipurpose* target for light dark matter searches, covering a brand new region of parameter space.

This also paves the way to a number of further investigations. For instance, our study does not deal with magnetic fields that are arbitrarily close to the critical point, as other effects might become relevant when the magnon becomes sufficiently light. It is also important to understand how the possible directional signal is affected by the presence of multiple domains in a realistic material, or if other material might perform better than NiO (see, e.g., [70]). In this second instance, one will have to work out another EFT, if useful. Moreover, it is known that NiO features a number of additional light excitations [45], and that the magnon gaps can also be varied (quite a lot) by applying an external strain to the sample [e.g., [71]]. Neither of these features are currently included in our EFT, and their implications for light dark matter searches have not been explored. In particular, if not suppressed by possible selection rules, the other light magnon modes could provide additional channels for axion absorption, which would allow one to cover an even larger region in parameter space. Finally, to turn this into a concrete proposal, one must address the question of what are the observable signatures associated with the emission of magnons in the NiO sample. This is also relevant in light of the fact that, once a concrete detection strategy is in

place, the final excluded region could be different from the ideal one presented here. We leave these and other interesting directions for future work.

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DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

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- [1] D. S. Akerib *et al.* (LUX Collaboration), Results from a search for dark matter in the complete LUX exposure, *Phys. Rev. Lett.* **118**, 021303 (2017).
 - [2] R. Agnese *et al.* (SuperCDMS Collaboration), First dark matter constraints from a SuperCDMS single-charge sensitive detector, *Phys. Rev. Lett.* **121**, 051301 (2018); **122**, 069901(E) (2019).
 - [3] E. Aprile *et al.* (XENON Collaboration), Light dark matter search with ionization signals in XENON1T, *Phys. Rev. Lett.* **123**, 251801 (2019).
 - [4] A. H. Abdelhameed *et al.* (CRESST Collaboration), First results from the CRESST-III low-mass dark matter program, *Phys. Rev. D* **100**, 102002 (2019).
 - [5] Qiuhong Wang *et al.* (PandaX-II Collaboration), Results of dark matter search using the full PandaX-II exposure, *Chin. Phys. C* **44**, 125001 (2020).
 - [6] R. Barbieri, C. Braggio, G. Carugno, C. S. Gallo, A. Lombardi, A. Ortolan, R. Pengo, G. Ruoso, and C. C. Speake, Searching for galactic axions through magnetized media: The QUAX proposal, *Phys. Dark Universe* **15**, 135 (2017).
 - [7] N. Crescini *et al.* (QUAX Collaboration), Axion search with a quantum-limited ferromagnetic haloscope, *Phys. Rev. Lett.* **124**, 171801 (2020).
 - [8] Chiara P. Salemi *et al.*, Search for low-mass axion dark matter with ABRACADABRA-10 cm, *Phys. Rev. Lett.* **127**, 081801 (2021).
 - [9] C. Bartram *et al.* (ADMX Collaboration), Search for invisible axion dark matter in the 3.3–4.2 μeV mass range, *Phys. Rev. Lett.* **127**, 261803 (2021).
 - [10] C. Bartram *et al.*, Axion dark matter experiment around 3.3 μeV with Dine-Fischler-Srednicki-Zhitnitsky discovery ability, *Phys. Rev. Lett.* **134**, 111002 (2025).
 - [11] A. Rettaroli *et al.* (QUAX Collaboration), Search for axion dark matter with the QUAX–LNF tunable haloscope, *Phys. Rev. D* **110**, 022008 (2024).
 - [12] K. Altenmüller *et al.* (CAST Collaboration), A new upper limit on the axion-photon coupling with an extended CAST run with a Xe-based Micromegas detector, *Phys. Rev. Lett.* **133**, 221005 (2024).
 - [13] H. Yan, G. A. Sun, S. M. Peng, H. Guo, B. Q. Liu, M. Peng, and H. Zheng, Constraining exotic spin dependent interactions of muons and electrons, *Eur. Phys. J. C* **79**, 971 (2019).
 - [14] E. Aprile *et al.* (XENON Collaboration), Search for new physics in electronic recoil data from XENONnT, *Phys. Rev. Lett.* **129**, 161805 (2022).
 - [15] Francesco Capozzi and Georg Raffelt, Axion and neutrino bounds improved with new calibrations of the tip of the red-giant branch using geometric distance determinations, *Phys. Rev. D* **102**, 083007 (2020).
 - [16] Ciaran O’Hare, Cajohare/axionlimits: Axionlimits, [10.5281/zenodo.3932430](https://zenodo.org/record/3932430) (2020).
 - [17] Wei Guo and Daniel N. McKinsey, Concept for a dark matter detector using liquid Helium-4, *Phys. Rev. D* **87**, 115001 (2013).
 - [18] Yonit Hochberg, Tongyan Lin, and Kathryn M. Zurek, Detecting ultralight bosonic dark matter via absorption in superconductors, *Phys. Rev. D* **94**, 015019 (2016).
 - [19] Katelin Schutz and Kathryn M. Zurek, Detectability of light dark matter with superfluid helium, *Phys. Rev. Lett.* **117**, 121302 (2016).
 - [20] Simon Knapen, Tongyan Lin, and Kathryn M. Zurek, Light dark matter in superfluid helium: Detection with multi-excitation production, *Phys. Rev. D* **95**, 056019 (2017).
 - [21] Yonit Hochberg, Tongyan Lin, and Kathryn M. Zurek, Absorption of light dark matter in semiconductors, *Phys. Rev. D* **95**, 023013 (2017).
 - [22] Simon Knapen, Tongyan Lin, Matt Pyle, and Kathryn M. Zurek, Detection of light dark matter with optical phonons in polar materials, *Phys. Lett. B* **785**, 386 (2018).
 - [23] S. A. Hertel, A. Biekert, J. Lin, V. Velan, and D. N. McKinsey, Direct detection of sub-GeV dark matter using a superfluid ^4He target, *Phys. Rev. D* **100**, 092007 (2019).
 - [24] Sinead Griffin, Simon Knapen, Tongyan Lin, and Kathryn M. Zurek, Directional detection of light dark matter with polar materials, *Phys. Rev. D* **98**, 115034 (2018).
 - [25] David J. E. Marsh, Kin-Chung Fong, Erik W. Lentz, Liboř Smejkal, and Mazhar N. Ali, Proposal to detect dark matter using axionic topological antiferromagnets, *Phys. Rev. Lett.* **123**, 121601 (2019).
 - [26] Francesca Acanfora, Angelo Esposito, and Antonio D. Polosa, Sub-GeV dark matter in superfluid He-4: An effective theory approach, *Eur. Phys. J. C* **79**, 549 (2019).
 - [27] Tanner Trickle, Zhengkang Zhang, and Kathryn M. Zurek, Detecting light dark matter with magnons, *Phys. Rev. Lett.* **124**, 201801 (2020).

- [28] Andrea Caputo, Angelo Esposito, and Antonio D. Polosa, Sub-MeV dark matter and the Goldstone modes of superfluid helium, *Phys. Rev. D* **100**, 116007 (2019).
- [29] Andrea Caputo, Angelo Esposito, Emma Geoffray, Antonio D. Polosa, and Sichun Sun, Dark matter, dark photon and superfluid He-4 from effective field theory, *Phys. Lett. B* **802**, 135258 (2020).
- [30] Tanner Trickle, Zhengkang Zhang, Kathryn M. Zurek, Katherine Inzani, and Sinéad M. Griffin, Multi-channel direct detection of light dark matter: Theoretical framework, *J. High Energy Phys.* 03 (2020) 036.
- [31] Sinéad M. Griffin, Katherine Inzani, Tanner Trickle, Zhengkang Zhang, and Kathryn M. Zurek, Multichannel direct detection of light dark matter: Target comparison, *Phys. Rev. D* **101**, 055004 (2020).
- [32] Gordon Baym, D.H. Beck, Jeffrey P. Filippini, C.J. Pethick, and Jessie Shelton, Searching for low mass dark matter via phonon creation in superfluid ^4He , *Phys. Rev. D* **102**, 035014 (2020); **104**, 019901(E) (2021).
- [33] Andrea Mitridate, Tanner Trickle, Zhengkang Zhang, and Kathryn M. Zurek, Detectability of axion dark matter with phonon polaritons and magnons, *Phys. Rev. D* **102**, 095005 (2020).
- [34] So Chigusa, Takeo Moroi, and Kazunori Nakayama, Detecting light boson dark matter through conversion into a magnon, *Phys. Rev. D* **101**, 096013 (2020).
- [35] Andrea Caputo, Angelo Esposito, Fulvio Piccinini, Antonio D. Polosa, and Giuseppe Rossi, Directional detection of light dark matter from three-phonon events in superfluid ^4He , *Phys. Rev. D* **103**, 055017 (2021).
- [36] Konstantin T. Matchev, Jordan Smolinsky, Wei Xue, and Yining You, Superfluid effective field theory for dark matter direct detection, *J. High Energy Phys.* 05 (2022) 034.
- [37] Yining You, Jordan Smolinsky, Wei Xue, Konstantin T. Matchev, Keegan Gunther, Yoonseok Lee, and Tarek Saab, Signatures and detection prospects for sub-GeV dark matter with superfluid helium, *J. High Energy Phys.* 07 (2023) 009.
- [38] Brian Campbell-Deem, Simon Knapen, Tongyan Lin, and Ethan Villarama, Dark matter direct detection from the single phonon to the nuclear recoil regime, *Phys. Rev. D* **106**, 036019 (2022).
- [39] Angelo Esposito and Shashin Pavaskar, Optimal antiferromagnets for light dark matter detection, *Phys. Rev. D* **108**, L011901 (2023).
- [40] Asher Berlin, Alexander J. Millar, Tanner Trickle, and Kevin Zhou, Physical signatures of fermion-coupled axion dark matter, *J. High Energy Phys.* 05 (2024) 314.
- [41] So Chigusa, Dan Kondo, Hitoshi Murayama, Risshin Okabe, and Hiroyuki Sudo, Axion detection via superfluid ^3He ferromagnetic phase and quantum measurement techniques, *J. High Energy Phys.* 09 (2024) 191.
- [42] Yonatan Kahn and Tongyan Lin, Searches for light dark matter using condensed matter systems, *Rep. Prog. Phys.* **85**, 066901 (2022).
- [43] Kathryn M. Zurek, Dark matter candidates of a very low mass, *Annu. Rev. Nucl. Part. Sci.* **74**, 287 (2024).
- [44] Mo T. Hutchings and E.J. Samuelsen, Measurement of spin-wave dispersion in nio by inelastic neutron scattering and its relation to magnetic properties, *Phys. Rev. B* **6**, 3447 (1972).
- [45] J. Milano, L. B. Steren, and M. Grimsditch, Effect of dipolar interaction on the antiferromagnetic resonance spectra of NiO, *Phys. Rev. Lett.* **93**, 077601 (2004).
- [46] Sergio M. Rezende, Antonio Azevedo, and Roberto L. Rodríguez-Suárez, Introduction to antiferromagnetic magnons, *J. Appl. Phys.* **126**, 151101 (2019).
- [47] Premala Chandra, Piers Coleman, and AI Larkin, A quantum fluids approach to frustrated Heisenberg models, *J. Phys. Condens. Matter* **2**, 7933 (1990).
- [48] C. P. Burgess, Goldstone and pseudoGoldstone bosons in nuclear, particle and condensed matter physics, *Phys. Rep.* **330**, 193 (2000).
- [49] Shashin Pavaskar, Riccardo Penco, and Ira Z. Rothstein, An effective field theory of magneto-elasticity, *SciPost Phys.* **12**, 155 (2022).
- [50] Tomáš Brauner, *Effective Field Theory for Spontaneously Broken Symmetry* (Springer, Cham, 2024), 10.1007/978-3-031-48378-3.
- [51] Pier Giuseppe Catinari, Angelo Esposito, and Shashin Pavaskar, companion paper, Effective field theory for anisotropic antiferromagnets: Gapped Goldstone bosons, pseudo-Goldstone excitations, and phase transitions, *Phys. Rev. B* **112**, 064408 (2025).
- [52] J.R. Tomlinson, L. Domash, R.G. Hay, and C.W. Montgomery, The high temperature heat content of nickel oxide, *J. Am. Chem. Soc.* **77**, 909 (1955).
- [53] Karthik Kannan, Devi Radhika, Kishor Kumar Sadasivuni, Kakarla Raghava Reddy, and Anjanapura V. Raghun, Nanostructured metal oxides and its hybrids for photocatalytic and biomedical applications, *Adv. Colloid Interface Sci.* **281**, 102178 (2020).
- [54] H. Leutwyler, Nonrelativistic effective Lagrangians, *Phys. Rev. D* **49**, 3033 (1994).
- [55] F.L.A. Machado, P.R.T. Ribeiro, J. Holanda, R.L. Rodríguez-Suárez, A. Azevedo, and S.M. Rezende, Spin-flop transition in the easy-plane antiferromagnet nickel oxide, *Phys. Rev. B* **95**, 104418 (2017).
- [56] Gordan Krnjaic, Duncan Rocha, and Tanner Trickle, The non-relativistic effective field theory of dark matter-electron interactions, *J. High Energy Phys.* 03 (2025) 165.
- [57] Zihang Wang and Lijing Shao, Axion induced spin effective couplings, *Phys. Rev. D* **103**, 116021 (2021).
- [58] Til Piffl *et al.*, The RAVE survey: The Galactic escape speed and the mass of the Milky Way, *Astron. Astrophys.* **562**, A91 (2014).
- [59] Giacomo Monari, Benoit Famaey, Ismael Carrillo, Tilmann Piffl, Matthias Steinmetz, Rosemary F.G. Wyse, Friedrich Anders, Cristina Chiappini, and Katja Janssen, The escape speed curve of the galaxy obtained from Gaia DR2 implies a heavy Milky Way, *Astron. Astrophys.* **616**, L9 (2018).
- [60] Xiaowei Ou, Anna-Christina Eilers, Lina Necib, and Anna Frebel, The dark matter profile of the Milky Way inferred from its circular velocity curve, *Mon. Not. R. Astron. Soc.* **528**, 693 (2024).
- [61] S. P. Bayrakci, T. Keller, K. Habicht, and B. Keimer, Spin-wave lifetimes throughout the brillouin zone, *Science* **312**, 1926 (2006).
- [62] S. P. Bayrakci, D.A. Tennant, Ph. Leininger, T. Keller, M. C. R. Gibson, S. D. Wilson, R. J. Birgeneau, and B. Keimer, Lifetimes of antiferromagnetic magnons in two and

- three dimensions: Experiment, theory, and numerics, *Phys. Rev. Lett.* **111**, 017204 (2013).
- [63] David N. Spergel, The motion of the earth and the detection of wimps, *Phys. Rev. D* **37**, 1353 (1988).
- [64] F. Mayet *et al.*, A review of the discovery reach of directional dark matter detection, *Phys. Rep.* **627**, 1 (2016).
- [65] Igor G. Irastorza and Javier Redondo, New experimental approaches in the search for axion-like particles, *Prog. Part. Nucl. Phys.* **102**, 89 (2018).
- [66] Luca Di Luzio, Maurizio Giannotti, Enrico Nardi, and Luca Visinelli, The landscape of QCD axion models, *Phys. Rep.* **870**, 1 (2020).
- [67] F. Bianchini, G. Grilli di Cortona, and M. Valli, The QCD axion: Some like it hot, *Phys. Rev. D* **110**, 123527 (2024).
- [68] Marcelo M. Miller Bertolami, Brenda E. Melendez, Leandro G. Althaus, and Jordi Isern, Revisiting the axion bounds from the Galactic white dwarf luminosity function, *J. Cosmol. Astropart. Phys.* **10** (2014) 069.
- [69] Andrea Caputo and Georg Raffelt, Astrophysical axion bounds: The 2024 edition, *Proc. Sci. COSMICWISPers* (2024) 041 [arXiv:2401.13728].
- [70] Giacomo Marocco and John Wheeler, Spin-dependent dark matter scattering in quasi-two-dimensional magnets, arXiv:2501.18120.
- [71] Hun-Ho Kim, Kentaro Ueda, Suguru Nakata, Peter Wochner, Andrew Mackenzie, Clifford Hicks, Giniyat Khaliullin, Huimei Liu, Bernhard Keimer, and Matteo Minola, Giant stress response of terahertz magnons in a spin-orbit Mott insulator, *Nat. Commun.* **13**, 6674 (2022).