

Nuclear magnetic resonance dark-matter searches are sensitive to dark photons and the axion-photon coupling

Carl Beadle¹, Sebastian A. R. Ellis¹, Jacob M. Leedom², and Nicholas L. Rodd^{3,4}

¹*Département de Physique Théorique, Université de Genève,
24 quai Ernest Ansermet, 1211 Genève 4, Switzerland*

²*CEICO, Institute of Physics of the Czech Academy of Sciences,
Na Slovance 2, 182 00 Prague 8, Czech Republic*

³*Theory Group, Lawrence Berkeley National Laboratory, Berkeley, California 94720, USA*

⁴*Berkeley Center for Theoretical Physics, University of California, Berkeley, California 94720, USA*



(Received 15 June 2025; accepted 21 January 2026; published 11 February 2026)

We demonstrate that nuclear magnetic resonance based searches for dark matter (DM) have intrinsic and powerful sensitivity to dark photons and the axion-photon coupling. The reason is conceptually straightforward. An instrument such as CASPER-Gradient begins with a large sample of nuclear spins polarized in a background magnetic field. In the presence of axion DM coupled to nucleons, the spin ensemble feels an effective magnetic field $\mathbf{B} \propto \nabla a$ that tilts the spins, generating a potentially observable precession. If the magnetic field is real rather than effective, the system responds identically. A real field can be generated by a kinetically mixed dark photon within the shielded region the sample is placed or an axion coupled to photons through its interaction with the background magnetic field. We show that all three signals are detectable and distinguishable. If CASPER-Gradient were to reach the QCD axion prediction of the axion-nucleon coupling, it would simultaneously be sensitive to kinetic mixings of $\epsilon \simeq 3 \times 10^{-16}$ and axion-photon couplings of $g_{a\gamma\gamma} \simeq 2 \times 10^{-16} \text{ GeV}^{-1}$ for $m \simeq 1 \mu\text{eV}$. Our analysis demonstrates that CASPER-Gradient has the tantalizing possibility of *simultaneously* measuring two separate couplings of axion dark matter, thereby providing a unique window into the axion's UV completion.

DOI: [10.1103/f9fj-6kkj](https://doi.org/10.1103/f9fj-6kkj)

Particle dark matter (DM) remains the most compelling hypothesis to explain the Universe's missing matter. However, evidence for its existence does not pinpoint a particular mass or strength of coupling to visible matter. Viable particle DM parameter space therefore spans a daunting range of masses. Historically, WIMPs and the electroweak scale have taken preeminence. More recently the community has cast a broader net, exemplified by the efforts to search for ultralight DM ($m \ll 10 \text{ eV}$), where DM transitions to a classical wave description [1].

A question less often considered is what is the full range of physics that our nets may catch. Planned experiments aim to cover enormous swathes of unexplored ultralight DM parameter space. Unexpected signals could emerge. For example, these devices can detect relativistic signals such as a cosmic axion background [2,3] or gravitational waves [4–10]. With this paper we instead demonstrate that

even DM signals have been overlooked: devices designed to search for one type of wave DM can be strongly sensitive to another. More to the point, ongoing efforts to detect axion DM coupled to nucleons with nuclear magnetic resonance (NMR), for example by CASPER-Gradient [11,12], would be highly sensitive to the axion-photon coupling and dark photon DM, as is the case in other torque based searches [13]. Given our ignorance as to the nature of DM, it stands that these situations should be highlighted whenever possible.

To justify our claim we recall the basics of axion NMR. Consider a sample of nuclear spins aligned in a background magnetic field $\mathbf{B}_0$. In conventional NMR, one fires an electromagnetic wave through the sample that generates an interaction according to $\mathcal{H}_B = -\gamma \mathbf{B} \cdot \mathbf{S}$, with γ and $\mathbf{S}$ the nucleon gyromagnetic ratio and spin operator. Consequently, the spins are tipped into the plane transverse to $\mathbf{B}_0$ and begin to precess at the Larmor frequency, $\omega_0 = \gamma B_0$, yielding an oscillatory signal in the transverse plane. CASPER-Gradient exploits the same physical principle. If the spins are immersed in an axion DM wave, then the axion-nucleon coupling, $\mathcal{L}_{aN} = g_N(\partial_\mu a)\bar{N}\gamma^\mu\gamma_5 N$, generates an interaction $\mathcal{H}_{aN} = -2g_N \nabla a \cdot \mathbf{S}$. By analogy to $\mathcal{H}_B$ the axion induces an *effective* magnetic field of the form

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

$\mathbf{B}_{aN} = g_N(2/\gamma)\nabla a$. The axion field has an identical effect: DM tips the spins, which then precess leading to a detectable signal. There is a resonant enhancement when $\omega_0 \simeq m$ and therefore one performs a resonant search by varying B_0 .

We can now explain our core insight: if DM is weakly coupled to electromagnetism and generates a magnetic field, CASPER-Gradient can detect it. Consider first dark photon DM, A' , that mixes with the visible photon via a kinetic mixing parameter ϵ [14]. To mitigate the impact of stray magnetic fields, the spin sample must be placed within a shielded region of geometric size L . The dark photon will generate a magnetic field $\mathbf{B}_{A'} \sim \epsilon m^2 L A'$ [15] within the conductive shielding. Setting $B_{aN} \sim B_{A'}$, the qualitative dark-photon sensitivity in terms of the axion-nucleon coupling is $\epsilon \sim g_N v / \gamma m L$, with $v \sim 10^{-3}$ the local DM speed. We can rephrase the sensitivity as follows. Writing $g_N = c_N m / m_\pi f_\pi$, which depends on the pion mass and decay constant, we have

$$\epsilon \sim 10^{-16} c_N (\gamma_{\text{He}}/\gamma) (10 \text{ cm}/L), \quad (1)$$

assuming a helium-3 sample. Although an ambitious target, the ultimate goal of all axion experiments is QCD sensitivity, where $c_N \sim 1$. Note the above scaling is distinct from ideas to repurpose axion-photon searches for the dark photon, as in, e.g., Ref. [16]. In those cases the dark photon sensitivity does not gain from the large magnetic

field volume that is maximized for the axion sensitivity. Here, however, the scaling of both the axion and dark photon signal is controlled by the density of spins, or effectively the background magnetization.

A similar argument applies to the axion-photon coupling, $\mathcal{L}_{a\gamma\gamma} = -\frac{1}{4} g_{a\gamma\gamma} a F \tilde{F}$, where axion DM couples to $\mathbf{B}_0$ and induces $B_{a\gamma} \sim g_{a\gamma\gamma} B_0 L (\partial_t a)$, where now L is the scale over which B_0 is generated. The expected sensitivity from $B_{aN} \sim B_{a\gamma}$ is $g_{a\gamma\gamma} \sim g_N v / \gamma B_0 L$, or taking $g_{a\gamma\gamma} \sim c_\gamma \alpha_{\text{EM}} m / 2\pi m_\pi f_\pi$,

$$c_\gamma \sim c_N (10 \text{ T}/B_0) (\gamma_{\text{He}}/\gamma) (10 \text{ cm}/L). \quad (2)$$

This result demonstrates that at the highest frequencies—recall B_0 is tuned to adjust the resonant frequency—an instrument sensitive to the QCD axion value of c_N would also be sensitive to the QCD axion prediction of $c_\gamma \sim 1$. In the event of a detection, this would immediately allow for an inference as to the underlying UV model [17].

In what follows we refine the qualitative results of Eqs. (1) and (2) culminating in the projections shown in Fig. 1. These results show the reach for several sensitivity estimates for CASPER-Gradient that exist in the literature, specifically Refs. [18,19], each shown for two different target nuclei, xenon-129 and helium-3. The results are compared to existing constraints from Refs. [20–41]. The theoretical predictions of Refs. [42–45] are also shown, see Ref. [46].

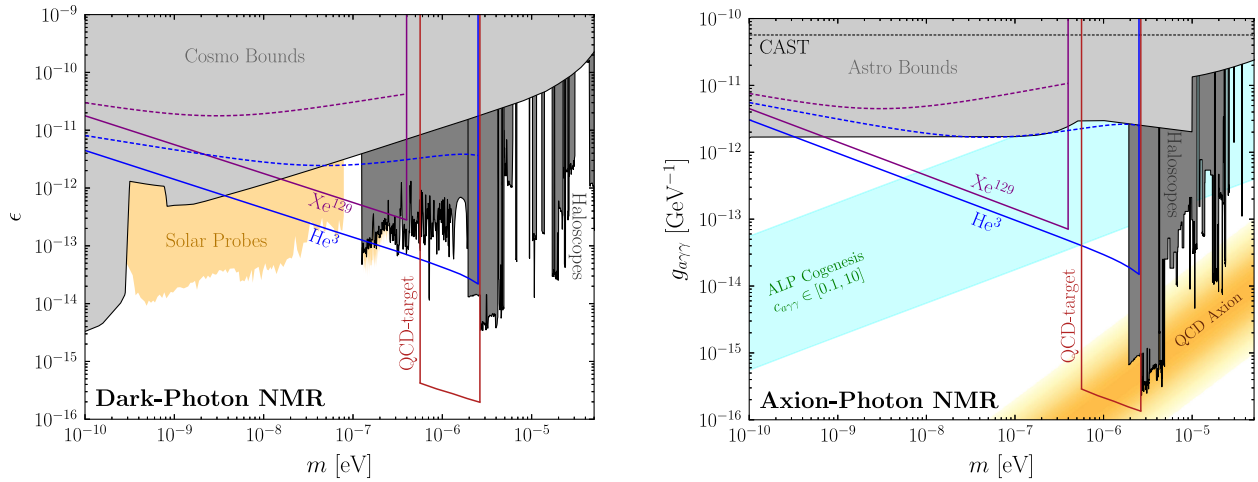


FIG. 1. Projected reach of DM NMR to dark photons (left) and axions (right). Both searches can be performed *simultaneously* along with a search for the axion nucleon coupling. The purple and blue lines correspond to a sample constituted of xenon-129 ($\gamma = 12 \text{ MHz/T}$, $n = 1.3 \times 10^{22} \text{ cm}^{-3}$, $H_{\text{max}} = 10 \text{ T}$) and helium-3 ($\gamma = 32 \text{ MHz/T}$, $n = 2.8 \times 10^{22} \text{ cm}^{-3}$, $H_{\text{max}} = 20 \text{ T}$). Solid curves assume $T_2 = 100 \text{ s}$, whereas dashed curves assume a line width limited by a part-per-million. An integration time of 100 s is assumed in both cases, which corresponds to a scan rate in mass of 7.3 years per decade in mass [18]. A search capable of reaching the QCD axion prediction for g_N would have the sensitivity shown in red for a five year integration [18]. The reach assumes a single TM mode with the sample placed at the optimal radius of $r \simeq 0.77R$. Also shown are constraints on dark-photon DM from Solar measurements in yellow [20,21], constraints of cosmological or astrophysical origin in light gray [22–30], and a combination of haloscope bounds in dark gray [31–40]. The results of the CAST experiment are denoted by a dotted line [41]. The QCD-band is shown in gold [42–44]. Lastly, the cyan region labeled ALP Cogenesis highlights a motivated region of parameter space, where the DM ALPs explain the Universe’s baryon asymmetry [45]. See the text for further details.

Principles of DM NMR. We begin with a self-contained discussion of how DM effects can be detected with NMR. Our analysis considerably simplifies earlier discussions, e.g., Ref. [18], and presents the first simple estimate of the DM sensitivity. Consider first a spin with magnetic moment $\boldsymbol{\mu} = \gamma \mathbf{S}$ in a magnetic field. The spin experiences a torque and therefore precesses according to $\dot{\boldsymbol{\mu}} = \boldsymbol{\mu} \times \gamma \mathbf{B}$. Taking $\mathbf{B} = B_0 \hat{\mathbf{z}}$, the transverse coordinates $\mu_{\pm} = (\mu_x \pm i\mu_y)/\sqrt{2}$ evolve under the precession as $\mu_{\pm}(t) \propto e^{\mp i\omega_0 t}$. NMR extends this picture to a macroscopic sample of spins: $\boldsymbol{\mu}$ is promoted to $\mathbf{M}$, the magnetization within the material, and $\mathbf{B}$ to $\mathbf{H}$. We assume $\mathbf{H} = H_0 \hat{\mathbf{z}}$ up to small corrections, however, the dynamics in the sample is also controlled by spin-spin and spin-lattice interactions. These interactions exponentially damp coherent oscillations in the transverse and longitudinal directions, respectively, and are quantified by corresponding timescales T_2 and T_1 . The system evolves according to the Bloch equations,

$$\dot{\mathbf{M}} = \mathbf{M} \times \gamma \mathbf{H} - \frac{M_x \hat{\mathbf{x}} + M_y \hat{\mathbf{y}}}{T_2} - \frac{(M_z - M_0) \hat{\mathbf{z}}}{T_1}, \quad (3)$$

with M_0 the steady state magnetization generated by H_0 .

We now introduce DM into the system as a small correction to the magnetic field, $\mathbf{H} = H_0 \hat{\mathbf{z}} + \mathbf{H}_{\text{DM}}$. While the aim of CASPER-Gradient is to detect the effective magnetic field $\mathbf{H}_{\text{DM}}^{\text{eff}} = g_N(2/\gamma)\nabla a$, the goal of this paper is to highlight a broader range of possibilities. Taking $H_{\text{DM}}/H_0 \ll 1$, we solve the problem perturbatively. At zeroth order, the steady state solution is $\mathbf{M} = M_0 \hat{\mathbf{z}}$. In general, both transverse components of $\mathbf{M}$ and $\mathbf{H}_{\text{DM}}$ are coupled and must be accounted for carefully. If, however, we define $M_{\pm} = (M_x \pm iM_y)/\sqrt{2}$ and similarly $H_{\text{DM}}^{\pm}$ the first order Bloch equations decouple, leaving

$$\dot{M}_{+} = -T_2^{-1}M_{+} - i(\omega_0 M_{+} - \gamma M_0 H_{\text{DM}}^{+}). \quad (4)$$

Taking $M_{+}(0) = 0$ the solution to this equation is

$$M_{+}(t) = i\gamma M_0 \int_0^t dt' e^{-(t-t')/T_2} e^{-i\omega_0(t-t')} H_{\text{DM}}^{+}(t'). \quad (5)$$

Accordingly, for a perfectly resonant DM field, $H_{\text{DM}}^{+} = H_{\text{DM}} e^{-i\omega_0 t}$, the magnetization grows as

$$M_{+}(t) = i e^{-i\omega_0 t} \gamma M_0 H_{\text{DM}} T_2 (1 - e^{-t/T_2}). \quad (6)$$

For $t \ll T_2$ DM drives a linear growth in the transverse component of the magnetization which continues until the amplitude begins to saturate once $t \sim T_2$. The dynamics are comparable for a more realistic model of a DM field with finite coherence. An exception occurs if the DM coherence time, τ_{DM} , is shorter than T_2 , in which case for $\tau_{\text{DM}} < t < T_2$ the growth of M_{+} slows from linear to $\sqrt{t}$ [18]. Regardless,

for $t \gg T_2$, τ_{DM} the amplitude of the transverse magnetization saturates at $\gamma M_0 H_{\text{DM}} T_2$.

In summary, DM drives the growth of a transverse oscillatory component to the magnetization. To provide a sensitivity estimate we must determine the appropriate backgrounds. First, the spins within the sample generate an intrinsic noise source according to the fluctuation dissipation theorem. Specifically, Eq. (5) defines a susceptibility of $\chi_{++}(t) = i\gamma M_0 e^{-t/T_2} e^{-i\omega_0 t} \Theta(t)$ or $\chi_{++}(\omega) = i\gamma M_0 / [i(\omega_0 - \omega) + T_2^{-1}]$. The imaginary part of this expression dictates the power spectral density (PSD) of the noise,

$$S_{M_{+}}^{\text{SP}}(\omega) \simeq \frac{2\gamma M_0 T_s}{\omega V} \frac{T_2}{1 + (\omega - \omega_0)^2 T_2^2}. \quad (7)$$

This background is none other than the spin projection noise. The noise depends on the bulk magnetization $M_0 = n\gamma\omega_0/4T_s$, with n the spin density in the sample and T_s the spin temperature—from this we confirm the usual result that the form is independent of temperature.

The combined effect of DM and spin-projection noise is a transverse magnetization that leads to a corresponding magnetic field outside the sample. If the magnetization is measured inductively through a SQUID, there is an additional background contribution from the magnetometer itself. Further, thermal noise sources throughout the device contribute additional backgrounds. Nevertheless, for the parameters CASPER intends to employ the dominant effect is spin projection noise.

Accordingly, we can estimate the sensitivity by searching for a DM signal over the spin-projection background. Working at high masses where $T_2 > \tau_{\text{DM}}$, we imagine interrogating the system for a time $T = T_2$ so that the observed power is distributed over a frequency bin of width $2\pi/T_2$. For the bin centered around ω_0 , the average spin projection noise power is $\bar{P}_{\text{SP}} = \gamma^2 n T_2 / 16\pi V$. For the signal, after accounting for the finite coherence, the equivalent expression is $\bar{P}_{\text{DM}} = (\gamma M_0 H_{\text{DM}} T_2)^2 \tau_a / 2\pi$. Assuming the signal and background magnetizations undergo Gaussian fluctuations, the power is distributed exponentially and we therefore expect the sensitivity limit to occur for $\bar{P}_{\text{DM}} = \bar{P}_{\text{SP}}$, or neglecting $\mathcal{O}(1)$ factors

$$H_{\text{DM}} \sim \frac{v}{\gamma} \sqrt{\frac{m}{n T_2 V}}, \quad (8)$$

where we took $M_0 = n(\gamma/2)$, assuming a hyperpolarized sample, and $\tau_a \sim 1/mv^2$. A longer measurement with $T > T_2$ would improve on this sensitivity by $(T/T_2)^{1/4}$; further, when $T_2 > \tau_a$ a gain can be made by accounting for the multiple frequency bins in which the signal appears [18]. For the axion-nucleon coupling, where $H_{\text{DM}} = g_N(2/\gamma)\nabla a \sim g_N(v/\gamma)\sqrt{\rho_{\text{DM}}}$, Eq. (8) gives an estimated sensitivity of $g_N \sim \sqrt{m/\rho_{\text{DM}} n T_2 V}$.

This ends our compact review of DM NMR. A more detailed discussion appears in the Supplemental Material (SM) [47] and we now turn to the different forms of $\mathbf{H}_{\text{DM}}$.

Dark photon NMR. Consider the scenario where DM is a boson of a dark gauged U(1): a dark photon A^μ . (We leave our discussion UV agnostic and refer instead to Refs. [74,75].) This state can couple to the visible sector through the marginal operator $\frac{1}{2}\epsilon F^{\mu\nu}F'_{\mu\nu}$ parametrized by a kinetic mixing ϵ taken to be small throughout. Rotating the two gauge fields to remove the kinetic mixing transforms to the interaction basis, where the dark photon remains decoupled from standard model currents yet provides a source for the visible U(1). In detail, the relevant Lagrangian interaction is $\mathcal{L} \supset \epsilon m^2 A^\mu A_\mu$, from which the dark photon generates an effective current $J_{\text{eff}}^\mu \simeq -\epsilon m^2 A^\mu$.

This dark-photon permeates the shielded regions, such as surrounds CASPER-Gradient. Taking the shielded region to have parametric size L , DM NMR operates in the regime where $mL \ll 1$ and therefore the induced magnetic field dominates [15]. In that limit we can neglect the displacement current implying the relevant equation is $\nabla \times \mathbf{H} \simeq -\epsilon m^2 \mathbf{A}'$, confirming the parametric size of the induced magnetic field as $\epsilon m^2 L \mathbf{A}'$. From Eq. (8), the parametric sensitivity scales as

$$\epsilon \sim \frac{v}{\gamma L} \sqrt{\frac{1}{m \rho_{\text{DM}} n T_2 V}}. \quad (9)$$

When compared to the scaling for g_N this result is consistent with Eq. (1).

To fully quantify the sensitivity we must consider how the modes of the shielded region can be driven by the dark-photon and the geometry of the associated magnetic fields. We treat the shielded region as a cylindrical cavity (deviations away from this will not strongly impact our

projections) and as $mL \ll 1$ the DM driving frequency is always much smaller than the electronic resonant frequencies ($\omega_0 \sim 1/L$), implying we can ignore dissipative effects. Further, the precise composition of the shielding material is not important, as the DM signal is at sufficiently high frequencies that the shield can be treated as a perfect conductor. We can then decompose the cavity field into the conventional transverse electric (TE) and transverse magnetic (TM) modes, each labeled by indices $\ell = (m, n, p)$ (for the azimuthal, radial, and longitudinal indices), from which the amount each mode is excited by the dark photon can be computed directly. The computation is standard and therefore not presented here, although the full details are provided in the SM [47].

One aspect of the calculation that must be mentioned, however, is the spatial profile of the dark-photon induced magnetic field. This dictates the direction of the magnetic field experienced by the spin sample and opens the possibility of distinguishing various DM induced NMR signals; see Fig. 2. To understand the field profile, we first emphasise that it is the lowest accessible cavity modes that are dominantly excited. If $\mathbf{E}_\ell$ represents the field of a given mode, then the induced magnetic field depends on the following integral over the shielded volume,

$$\int dV \mathbf{E}_\ell^* \cdot \mathbf{A}' \simeq \mathbf{A}' \cdot \int dV \mathbf{E}_\ell^*, \quad (10)$$

where the accuracy of the approximation exploits the spatial uniformity of DM over spatial scales $\mathcal{O}(L)$ [1]. The value of the final integral is suppressed for higher modes; indeed, it is an excellent approximation to only retain the lowest nonvanishing modes, TM_{010} driven by A'_z and TE_{111} driven by $A'_{x,y}$. Accordingly, the modes driven are sensitive to the polarization of the dark photon, a fact that could be exploited to not only measure the polarization

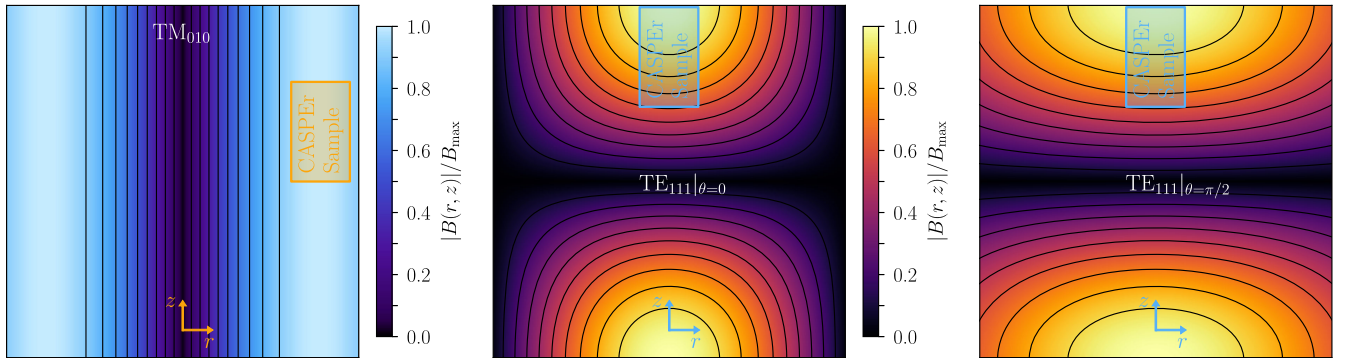


FIG. 2. The profile of the physical magnetic fields DM could excite as compared the size of CASPER's sample of nuclear spins. The contours denote 10% changes in the magnetic field value. The dominant mode the axion-photon coupling can excite is TM_{010} , whereas a dark photon can also excite the TE_{111} mode depending upon its polarization. The figure highlights that as the sample is moved within the magnetic field the response varies, opening a path to distinguishing different DM signals and even mapping out the dark photon polarization.

distribution in the event of a detection but also to distinguish the signal from axion DM.

The spatial profile of the modes determines the ideal placement of the spin sample as shown in Fig. 2. For CASPER-Gradient the sample is sufficiently smaller than the shielded region such that the sample can in principle be moved around and in the case of the TM_{010} mode can inhabit a region of uniform induced field. (For other modes and a larger sample the effect of a nonuniform field across the sample must be accounted for, we consider this scenario in the SM [47].) In detail, we take the shielding to have a radius of $R \simeq 9$ cm and height of 18 cm, whereas the cylindrical sample has a height of 3 cm and radius of 2.5 cm. In Fig. 2 the axis of the sample is taken to be perpendicular to that of the shielding. Generally, the optimal placement is determined by looking for where the modes take their maximum value; for the TM_{010} mode the sample should be placed at a radius of roughly $0.77R$, whereas for the TE_{111} mode the sample should be placed on-axis radially and near the end caps longitudinally.

Comparing Eq. (9) to Fig. 1 we find that the parametric estimate for fixed T_2 accurately reflects the scaling of the full result, but overstates the sensitivity by roughly a decade. The difference is a combination of $\mathcal{O}(1)$ factors and roughly a factor of ~ 1.8 suppression from the geometry of the cavity. The sensitivity breaks down for $m \gtrsim 1 \mu\text{eV}$ as the cavity modes begin to be driven resonantly. This breakdown is largely irrelevant to our projections because the largest mass value that can be observed is dictated by the externally applied magnetic field and is below the point where the Q -factor of the cavity must be specified. The weaker projections arise by assuming that the scan is limited by a part per million line-width at every frequency due to magnetic field inhomogeneities [76]. The effective T_2 is then frequency dependent, $T_2 \propto 1/m$, so that Eq. (9) predicts a flat sensitivity scaling with mass as approximately observed.

Axion-photon NMR. The considerations for axion DM that couples through $\mathcal{L}_{a\gamma\gamma}$ are almost identical to those for the dark-photon. For axion DM, the axion's coupling to electromagnetism sources an induced current $\mathbf{J}_{\text{eff}} \simeq g_{a\gamma\gamma}(\partial_t a)\mathbf{B}$. The magnetic field in this current is sourced by the background field used to set the Larmor frequency. As that background magnetic field points along the axis of the cylindrical shielded region, from Eq. (10) it is apparent that the axion does not couple to the TE modes, leaving the dominant contribution from the TM_{010} mode, again as shown in Fig. 2. For Fig. 1 we assume the sample is placed optimally. The quantitative sensitivity there can be determined as for the dark photon: $H_{\text{DM}} \simeq g_{a\gamma\gamma} B_0 L(\partial_t a)$. A critical point, however, is that B_0 is not a fixed parameter, but is rather adjusted to scan through various DM masses. If we denote $\bar{B}_0$ as the magnetic field used to scan the largest mass in the search, $\bar{m}$, then $B_0 = \bar{B}_0(m/\bar{m})$. Finally, we can

use Eq. (8) to obtain a parametric sensitivity of

$$g_{a\gamma\gamma} \sim \frac{v\bar{m}}{\gamma\bar{B}_0 L} \sqrt{\frac{1}{m\rho_{\text{DM}} n T_2 V}}. \quad (11)$$

As for the dark photon, this result correctly reproduces the scaling in Fig. 1.

Discussion. Axion NMR is a promising avenue to search for axion DM through its coupling to nucleons. This coupling appears as an *effective* magnetic field that couples to a spin sample, causing it to precess in a measurable manner. Significant experimental efforts are underway by the CASPER-Gradient collaboration [11,12,19,77], and other proposals have been made to conduct similar experiments [78–81]. Axion NMR is also an active area of theoretical research aiming to refine the ultimate limits in sensitivity achievable with this technique [18,82,83].

In this paper we quantified the sensitivity that axion NMR experiments have to specific alternative DM signals that generate magnetic fields in the vicinity of the hyperpolarized spin sample. We demonstrated that CASPER-Gradient can have world-leading sensitivity to dark-photon DM as well as to the axion-photon coupling in the same experimental setup they use to search for the axion-nucleon coupling. Importantly, the three signals can be distinguished by virtue of their symmetry properties. The canonical signal induced by the axion-nucleon coupling is homogeneous throughout the electromagnetically shielded volume in which the spin sample is placed. In contrast, the dark photon and axion-photon coupling generate inhomogeneous signals: Fig. 2 shows that the three signals can be differentiated by modifying the placement of the sample inside the shield. The ability to distinguish between signals provides CASPER-Gradient with the unique prospect of testing two axion couplings simultaneously, thereby probing the UV structure of the axion model. We emphasize that our insight implies that any magnetic field generated by the shielding can be searched for using the spin-sample. Thus, NMR could prove a powerful probe of any new physics that couples to electromagnetism, even beyond the DM models presented here.

The coming years promise revolutionary progress in the search for ultralight DM with masses below a μeV . In particular, as instruments like DMRadio [84,85] or heterodyne SRF haloscopes [86–90] are increased in scale, the goal is to improve the sensitivity by many, many orders of magnitude beyond existing low-mass haloscopes [91–94] to push toward the QCD axion prediction. DM NMR represents a completely distinct experimental strategy to push toward that same goal, and as the present work has highlighted it can probe many forms of DM as it does so. For this reason it will be critical to determine how far the DM NMR approach can ultimately be scaled, in particular

understanding the quantum mechanics-imposed limitations [82] and ways they can be evaded through the use of quantum resources [95].

Acknowledgments. We thank the CASPER collaboration for both feedback and providing technical regarding the CASPER-Gradient instrument. The work of C.B. and S.A.R.E. was supported by Swiss National Science Foundation (SNSF) Ambizione Grant No. PZ00P2 193322, *New frontiers from sub-eV to super-TeV*. J.M.L. was cofunded by the European Union and supported by the Czech Ministry of Education, Youth

and Sports (Project No. FORTE-CZ.02.01.01/00/22 008/0004632), as well as by the Deutsche Forschungsgemeinschaft under Germany’s Excellence Strategy—EXC 2121 “Quantum Universe”—390833306. The research of N.L.R. was supported by the Office of High Energy Physics of the U.S. Department of Energy under Contract No. DE-AC02-05CH11231.

Data availability. The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

-
- [1] D. Y. Cheong, N. L. Rodd, and L.-T. Wang, Quantum description of wave dark matter, *Phys. Rev. D* **111**, 015028 (2025).
 - [2] J. A. Dror, H. Murayama, and N. L. Rodd, Cosmic axion background, *Phys. Rev. D* **103**, 115004 (2021); **106**, 119902 (E) (2022).
 - [3] T. Nitta *et al.* (ADMX Collaboration), Search for a dark-matter-induced cosmic axion background with ADMX, *Phys. Rev. Lett.* **131**, 101002 (2023).
 - [4] A. Berlin, D. Blas, R. Tito D’Agnolo, S. A. R. Ellis, R. Harnik, Y. Kahn, and J. Schütte-Engel, Detecting high-frequency gravitational waves with microwave cavities, *Phys. Rev. D* **105**, 116011 (2022).
 - [5] A. Berlin, D. Blas, R. Tito D’Agnolo, S. A. R. Ellis, R. Harnik, Y. Kahn, J. Schütte-Engel, and M. Wentzel, Electromagnetic cavities as mechanical bars for gravitational waves, *Phys. Rev. D* **108**, 084058 (2023).
 - [6] V. Domcke, C. Garcia-Cely, and N. L. Rodd, Novel search for high-frequency gravitational waves with low-mass axion haloscopes, *Phys. Rev. Lett.* **129**, 041101 (2022).
 - [7] V. Domcke, C. Garcia-Cely, S. M. Lee, and N. L. Rodd, Symmetries and selection rules: Optimising axion haloscopes for gravitational wave searches, *J. High Energy Phys.* **03** (2024) 128.
 - [8] V. Domcke, S. A. R. Ellis, and N. L. Rodd, Magnets are weber bar gravitational wave detectors, *Phys. Rev. Lett.* **134**, 231401 (2025).
 - [9] K. M. W. Pappas *et al.*, High-frequency gravitational wave search with ABRACADABRA-10cm, [arXiv:2505.02821](https://arxiv.org/abs/2505.02821).
 - [10] R. Tito D’Agnolo and S. A. R. Ellis, Classical (and quantum) heuristics for gravitational wave detection, *J. High Energy Phys.* **04** (2025) 164.
 - [11] P. W. Graham and S. Rajendran, New observables for direct detection of axion dark matter, *Phys. Rev. D* **88**, 035023 (2013).
 - [12] D. Budker, P. W. Graham, M. Ledbetter, S. Rajendran, and A. O. Sushkov, Proposal for a cosmic axion spin precession experiment (CASPER), *Phys. Rev. X* **4**, 021030 (2014).
 - [13] S. Kalia, D. Budker, D. F. J. Kimball, W. Ji, Z. Liu, A. O. Sushkov, C. Timberlake, H. Ulbricht, A. Vinante, and T. Wang, Ultralight dark matter detection with levitated ferromagnets, *Phys. Rev. D* **110**, 115029 (2024).
 - [14] B. Holdom, Two $U(1)$ ’s and epsilon charge shifts, *Phys. Lett.* **166B**, 196 (1986).
 - [15] S. Chaudhuri, P. W. Graham, K. Irwin, J. Mardon, S. Rajendran, and Y. Zhao, Radio for hidden-photon dark matter detection, *Phys. Rev. D* **92**, 075012 (2015).
 - [16] A. Caputo, A. J. Millar, C. A. J. O’Hare, and E. Vitagliano, Dark photon limits: A handbook, *Phys. Rev. D* **104**, 095029 (2021).
 - [17] J. M. Leedom, S. Moldovsky, and H. Murayama, Probing the UV with IR axion dark matter experiments, [arXiv:2504.07188](https://arxiv.org/abs/2504.07188).
 - [18] J. A. Dror, S. Gori, J. M. Leedom, and N. L. Rodd, Sensitivity of spin-precession axion experiments, *Phys. Rev. Lett.* **130**, 181801 (2023).
 - [19] D. F. Jackson Kimball, Overview of the cosmic axion spin precession experiment (CASPER), *Springer Proc. Phys.* **245**, 105 (2020).
 - [20] H. An, S. Ge, J. Liu, and M. Liu, In Situ measurements of dark photon dark matter using parker solar probe: Going beyond the radio window, *Phys. Rev. Lett.* **134**, 171001 (2025).
 - [21] H. An, X. Chen, S. Ge, J. Liu, and Y. Luo, Searching for ultralight dark matter conversion in solar corona using low frequency array data, *Nat. Commun.* **15**, 915 (2024).
 - [22] P. Arias, D. Cadamuro, M. Goodsell, J. Jaeckel, J. Redondo, and A. Ringwald, WISPy cold dark matter, *J. Cosmol. Astropart. Phys.* **06** (2012) 013.
 - [23] A. Caputo, H. Liu, S. Mishra-Sharma, and J. T. Ruderman, Dark photon oscillations in our inhomogeneous universe, *Phys. Rev. Lett.* **125**, 221303 (2020).
 - [24] S. D. McDermott and S. J. Witte, Cosmological evolution of light dark photon dark matter, *Phys. Rev. D* **101**, 063030 (2020).
 - [25] D. Noordhuis, A. Prabhu, S. J. Witte, A. Y. Chen, F. Cruz, and C. Weniger, Novel constraints on axions produced in pulsar polar-cap cascades, *Phys. Rev. Lett.* **131**, 111004 (2023).
 - [26] C. Dessert, D. Dunskey, and B. R. Safdi, Upper limit on the axion-photon coupling from magnetic white dwarf polarization, *Phys. Rev. D* **105**, 103034 (2022).

- [27] J. Davies, M. Meyer, and G. Cotter, Constraints on axionlike particles from a combined analysis of three flaring Fermi flat-spectrum radio quasars, *Phys. Rev. D* **107**, 083027 (2023).
- [28] H. Abe *et al.* (MAGIC Collaboration), Constraints on axionlike particles with the perseus galaxy cluster with MAGIC, *Phys. Dark Universe* **44**, 101425 (2024).
- [29] M. Ajello *et al.* (Fermi-LAT Collaboration), Search for spectral irregularities due to photon–axionlike-particle oscillations with the Fermi Large Area Telescope, *Phys. Rev. Lett.* **116**, 161101 (2016).
- [30] J. N. Benabou, C. Dessert, K. C. Patra, T. G. Brink, W. Zheng, A. V. Filippenko, and B. R. Safdi, Search for axions in magnetic white dwarf polarization at Lick and Keck observatories, [arXiv:2504.12377](https://arxiv.org/abs/2504.12377).
- [31] A. Wagner *et al.* (ADMX Collaboration), A search for hidden sector photons with ADMX, *Phys. Rev. Lett.* **105**, 171801 (2010).
- [32] L. H. Nguyen, A. Lobanov, and D. Horns, First results from the WISPDMMX radio frequency cavity searches for hidden photon dark matter, *J. Cosmol. Astropart. Phys.* **10** (2019) 014.
- [33] S. Lee, S. Ahn, J. Choi, B. R. Ko, and Y. K. Semertzidis, Axion dark matter search around $6.7 \mu\text{eV}$, *Phys. Rev. Lett.* **124**, 101802 (2020).
- [34] L. Zhong *et al.* (HAYSTAC Collaboration), Results from phase 1 of the HAYSTAC microwave cavity axion experiment, *Phys. Rev. D* **97**, 092001 (2018).
- [35] K. M. Backes *et al.* (HAYSTAC Collaboration), A quantum-enhanced search for dark matter axions, *Nature (London)* **590**, 238 (2021).
- [36] B. Godfrey *et al.*, Search for dark photon dark matter: Dark E field radio pilot experiment, *Phys. Rev. D* **104**, 012013 (2021).
- [37] S. Ahn *et al.* (CAPP Collaboration), Extensive search for axion dark matter over 1 GHz with CAPP’S main axion experiment, *Phys. Rev. X* **14**, 031023 (2024).
- [38] D. Alesini *et al.*, Search for invisible axion dark matter of mass $m_a = 43 \mu\text{eV}$ with the QUAX- $\alpha\gamma$ experiment, *Phys. Rev. D* **103**, 102004 (2021).
- [39] S. Ahyoune *et al.*, RADES axion search results with a high-temperature superconducting cavity in an 11.7 T magnet, *J. High Energy Phys.* **04** (2025) 113.
- [40] A. P. Quiskamp, B. T. McAllister, P. Altin, E. N. Ivanov, M. Goryachev, and M. E. Tobar, Direct search for dark matter axions excluding ALPogenesis in the 63- to $67\text{-}\mu\text{eV}$ range with the ORGAN experiment, *Sci. Adv.* **8**, abq3765 (2022).
- [41] V. Anastassopoulos *et al.* (CAST Collaboration), New CAST limit on the axion-photon interaction, *Nat. Phys.* **13**, 584 (2017).
- [42] M. Dine, W. Fischler, and M. Srednicki, A simple solution to the strong CP problem with a harmless axion, *Phys. Lett.* **104B**, 199 (1981).
- [43] J. E. Kim, Weak interaction singlet and strong CP invariance, *Phys. Rev. Lett.* **43**, 103 (1979).
- [44] M. A. Shifman, A. I. Vainshtein, and V. I. Zakharov, Can confinement ensure natural CP invariance of strong interactions?, *Nucl. Phys.* **B166**, 493 (1980).
- [45] R. T. Co, L. J. Hall, and K. Harigaya, Predictions for axion couplings from ALPogenesis, *J. High Energy Phys.* **01** (2021) 172.
- [46] C. O’Hare, cajohare/axionlimits: Axionlimits, <https://cajohare.github.io/AxionLimits/> (2020).
- [47] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/f9fj-6kkj> for additional details on the conventions used, a derivation of CASPER-Gradient’s time-scaling, a discussion on cavity modes, a microscopic look at the Bloch equations and a note on the fluctuation-dissipation theorem. It includes additional Refs. [48–73].
- [48] G. Cowan, K. Cranmer, E. Gross, and O. Vitells, Asymptotic formulae for likelihood-based tests of new physics, *Eur. Phys. J. C* **71**, 1554 (2011); **73**, 2501(E) (2013).
- [49] J. W. Foster, N. L. Rodd, and B. R. Safdi, Revealing the dark matter halo with axion direct detection, *Phys. Rev. D* **97**, 123006 (2018).
- [50] D. F. Jackson Kimball, J. Dudley, Y. Li, S. Thulasi, S. Pustelny, D. Budker, and M. Zolotarev, Magnetic shielding and exotic spin-dependent interactions, *Phys. Rev. D* **94**, 082005 (2016).
- [51] D. Gavilan-Martin *et al.*, Searching for dark matter with a 1000 km baseline interferometer, *Nat. Commun.* **16**, 4953 (2025).
- [52] A. Berlin and K. Zhou, Discovering QCD-coupled axion dark matter with polarization haloscopes, *Phys. Rev. D* **108**, 035038 (2023).
- [53] E. U. Condon, Forced oscillations in cavity resonators, *J. Appl. Phys.* **12**, 129 (1941).
- [54] W. B. Smythe, *Static and Dynamic Electricity* (Hemisphere Publishing, New York, NY (USA), 1988).
- [55] R. E. Collin, *Field Theory of Guided Waves*, 2nd ed., IEEE Press Series on Electromagnetic Wave Theory (John Wiley & Sons, Nashville, TN, 1990).
- [56] J. W. Foster, Y. Kahn, R. Nguyen, N. L. Rodd, and B. R. Safdi, Dark matter interferometry, *Phys. Rev. D* **103**, 076018 (2021).
- [57] M. Jiang *et al.*, Search for dark photons with synchronized quantum sensor network, *Nat. Commun.* **15**, 3331 (2024).
- [58] A. Abragam, *The Principles of Nuclear Magnetism*, International Series of Monographs on Physics (Clarendon Press, Oxford, England, 1983).
- [59] N. G. Van Kampen, *Stochastic Processes in Physics and Chemistry* (Elsevier, New York, 2007).
- [60] J. Kowalewski and L. Maler, *Nuclear Spin Relaxation in Liquids*, 2nd ed. (CRC Press, London, England, 2019).
- [61] M. H. Levitt, *Spin Dynamics*, 2nd ed. (Wiley-Blackwell, Hoboken, NJ, 2008).
- [62] A. Redfield, *The Theory of Relaxation Processes* (Elsevier, New York, 1965).
- [63] C. P. Slichter, *Principles of Magnetic Resonance* (Springer, Berlin, Heidelberg, 1990).
- [64] J. D. Jackson, *Classical Electrodynamics* (Wiley, New York, 1998).
- [65] C. Kittel, *Introduction to Solid State Physics*, 8th ed. (Wiley, New York, 2004).
- [66] J. H. Simpson, Chapter 2—instrumental considerations, in *Organic Structure Determination Using 2-D NMR Spectroscopy (Second Edition)* (Academic Press, Boston, 2012), pp. 21–57.
- [67] N. Bloembergen, E. M. Purcell, and R. V. Pound, Relaxation effects in nuclear magnetic resonance absorption, *Phys. Rev.* **73**, 679 (1948).

- [68] I. Solomon, Relaxation processes in a system of two spins, *Phys. Rev.* **99**, 559 (1955).
- [69] D. D. Traficante, Relaxation. Can t_2 be longer than t_1 ?, *Concepts Magn. Reson.* **3**, 171 (1991).
- [70] N. Bloembergen and R. V. Pound, Radiation damping in magnetic resonance experiments, *Phys. Rev.* **95**, 8 (1954).
- [71] H. B. Callen and T. A. Welton, Irreversibility and generalized noise, *Phys. Rev.* **83**, 34 (1951).
- [72] R. Kubo, The fluctuation-dissipation theorem, *Rep. Prog. Phys.* **29**, 255 (1966).
- [73] L. D. Landau and E. M. Lifshitz, Chapter XII—Fluctuations, in *Statistical Physics (Third Edition)*, Vol. 5 of Course of Theoretical Physics (Pergamon, New York, 1980), pp. 333–400.
- [74] D. Cyncynates and Z. J. Weiner, Experimental targets for dark photon dark matter, *Phys. Rev. D* **111**, 103535 (2025).
- [75] W. E. East and J. Huang, Dark photon vortex formation and dynamics, *J. High Energy Phys.* **12** (2022) 089.
- [76] J. Walter, H. Bekker, J. Blanchard, D. Budker, N. L. Figueroa, A. Wickenbrock, Y. Zhang, and P. Zhou, Fast shimless algorithm based on Bayesian optimization for magnetic resonance based dark matter search, *arXiv*: 2309.11614.
- [77] A. Garcon *et al.*, The cosmic axion spin precession experiment (CASPER): A dark-matter search with nuclear magnetic resonance, *Quantum Sci. Technol.* **3**, 014008 (2017).
- [78] C. Gao, W. Halperin, Y. Kahn, M. Nguyen, J. Schütte-Engel, and J. W. Scott, Axion wind detection with the homogeneous precession domain of superfluid helium-3, *Phys. Rev. Lett.* **129**, 211801 (2022).
- [79] J. W. Foster, C. Gao, W. Halperin, Y. Kahn, A. Mande, M. Nguyen, J. Schütte-Engel, and J. W. Scott, Statistics and sensitivity of axion wind detection with the homogeneous precession domain of superfluid helium-3, *Phys. Rev. D* **110**, 115020 (2024).
- [80] C. Brandenstein, S. Stelzl, E. Gutsmedl, W. Schott, A. Weiler, and P. Fierlinger, Towards an electrostatic storage ring for fundamental physics measurements, *EPJ Web Conf.* **282**, 01017 (2023).
- [81] I. M. Bloch, Y. Hochberg, E. Kuflik, and T. Volansky, Axion-like relics: New constraints from old comagnetometer data, *J. High Energy Phys.* **01** (2020) 167.
- [82] D. Aybas, H. Bekker, J. W. Blanchard, D. Budker, G. P. Centers, N. L. Figueroa, A. V. Gramolin, D. F. J. Kimball, A. Wickenbrock, and A. O. Sushkov, Quantum sensitivity limits of nuclear magnetic resonance experiments searching for new fundamental physics, *Quantum Sci. Technol.* **6**, 034007 (2021).
- [83] Y. Zhang, D. A. Tumturk, H. Bekker, D. Budker, D. F. J. Kimball, A. O. Sushkov, and A. Wickenbrock, Frequency-scanning considerations in axionlike dark matter spin-precession experiments, *Ann. Phys. (Berlin)* **536**, 2300223 (2023).
- [84] L. Brouwer *et al.* (DMRadio Collaboration), Proposal for a definitive search for GUT-scale QCD axions, *Phys. Rev. D* **106**, 112003 (2022).
- [85] L. Brouwer *et al.* (DMRadio Collaboration), Projected sensitivity of DMRadio-m3: A search for the QCD axion below $1\ \mu\text{eV}$, *Phys. Rev. D* **106**, 103008 (2022).
- [86] A. Berlin, R. T. D’Agnolo, S. A. R. Ellis, C. Nantista, J. Neilson, P. Schuster, S. Tantawi, N. Toro, and K. Zhou, Axion dark matter detection by superconducting resonant frequency conversion, *J. High Energy Phys.* **07** (2020) 088.
- [87] A. Berlin, R. T. D’Agnolo, S. A. R. Ellis, and K. Zhou, Heterodyne broadband detection of axion dark matter, *Phys. Rev. D* **104**, L111701 (2021).
- [88] A. Berlin *et al.*, Searches for new particles, dark matter, and gravitational waves with SRF cavities, *arXiv*:2203.12714.
- [89] D. Alesini *et al.*, The future search for low-frequency axions and new physics with the FLASH resonant cavity experiment at Frascati National Laboratories, *Phys. Dark Universe* **42**, 101370 (2023).
- [90] S. Ahyoune *et al.*, A proposal for a low-frequency axion search in the $1\text{--}2\ \mu\text{eV}$ range and below with the BabyIAXO magnet, *Ann. Phys. (Berlin)* **535**, 2300326 (2023).
- [91] J. L. Ouellet *et al.*, First results from ABRACADABRA-10cm: A search for sub- μeV axion dark matter, *Phys. Rev. Lett.* **122**, 121802 (2019).
- [92] J. L. Ouellet *et al.*, Design and implementation of the ABRACADABRA-10 cm axion dark matter search, *Phys. Rev. D* **99**, 052012 (2019).
- [93] C. P. Salemi *et al.*, Search for low-mass axion dark matter with ABRACADABRA-10cm, *Phys. Rev. Lett.* **127**, 081801 (2021).
- [94] A. V. Gramolin, D. Aybas, D. Johnson, J. Adam, and A. O. Sushkov, Search for axion-like dark matter with ferromagnets, *Nat. Phys.* **17**, 79 (2021).
- [95] E. Boyers, G. Goldstein, and A. O. Sushkov, Spin squeezing of macroscopic nuclear spin ensembles, *Phys. Rev. D* **111**, 052004 (2025).