

Axion production and detection using a dual NMR-type experiment

Jeff A. Dror,¹ Qiushi Wei¹, and Fengwei Yang^{1,2}

¹*Institute for Fundamental Theory, Physics Department, University of Florida, Gainesville, Florida 32611, USA*

²*Department of Physics and Astronomy, University of Notre Dame, South Bend, Indiana 46556, USA*



(Received 10 November 2025; accepted 17 July 2026; published 4 August 2026)

Axions that couple to nuclear spins via the axial current interaction can be both *produced and detected* using nuclear magnetic resonance (NMR) techniques. In this scheme, nuclei driven by a real oscillating magnetic field in one device act as an axion source, which can drive NMR in a nearby spin-polarized sample interrogated with a sensitive magnetometer. We study the prospects for detecting axions through this method and identify two key characteristics that result in compelling detection sensitivity. First, the gradient of the generated axion field can be substantial, set by the inverse distance from the source. Near the source, it reduces to the inverse of the source’s geometric size. Second, because the generated axion field is produced at a known frequency, the detection medium can be tuned precisely to this frequency, enabling long interrogation times. We show that the experimental sensitivity of a pair of centimeter-scale NMR devices operating over a 15-day integration time can already surpass existing astrophysical bounds on the axion-nucleon coupling. A similar sensitivity can be achieved with 10 centimeter-scale NMR devices with only 1 hour of integration time. These dual NMR configurations are capable of probing a wide range of axion masses, up to values comparable to the inverse distance between the source and the sensor.

DOI: [10.1103/y8rj-rpfg](https://doi.org/10.1103/y8rj-rpfg)

Introduction. Axions are derivatively coupled light scalar fields that appear in many extensions of the standard model. They provide a compelling solution to the strong *CP* problem, serve as viable dark matter candidates, and arise naturally in string theory [1–7]. Their derivative couplings ensure that the axion mass is protected from large quantum corrections by a shift symmetry, broken only by non-perturbative effects. Consequently, axions remain light, and their interactions with the standard model are typically extremely weak.

It is conceivable that axions persist but do not make up a substantial cosmological abundance. If so, their detection must rely on the production of axions in the laboratory. In the case of axion-photon couplings, this has motivated experiments employing light-shining-through-walls experiments [8–16], which utilize radio or optical frequency cavities with background electromagnetic fields. If an axion carries both derivative (*CP*-conserving) and nonderivative (*CP*-violating) interactions with nuclei, then it has been proposed to generate axions with mechanical oscillators or rotators and detect its presence using nuclear magnetic

resonance (NMR) techniques [17–20]. For purely derivative couplings, existing searches for exotic spin-dependent interactions using atomic magnetometers [21,22] focus on the static limit where the axion mass is much larger than the characteristic oscillation frequency ($m_a \gg \omega$).

In this Letter, we present a novel production and detection scheme for axions that relies solely on the axial current interaction between axions and nucleons using a pair of NMR-type experiments, providing the first comprehensive dynamical framework for laboratory-scale axion production. Independent of the ambient axion abundance, axion production is realized by an “inverse” version of an NMR-based axion detector, where the precessing spins of nuclei driven by an external alternating magnetic field generate an oscillating axion field, which can in turn be detected by an additional nearby NMR device. This experimental setup extends current spin-precession axion experiments, such as CASPEr [23–29], to probe axions not sourced cosmologically.

Unlike axion dark matter, which oscillates at its mass and has a finite coherence time, the NMR-generated axion field is produced coherently at the nuclear spin precession frequency, independent of the axion mass. This allows a dual-NMR setup to probe a wide range of axion masses with a single resonance frequency and enables much longer interrogation times than typical dark matter searches. Moreover, the axion field gradient—which drives the NMR detector—is determined by the source geometry

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.

and distance, and can be significantly larger than the corresponding gradient in the dark matter case. Together, these features provide promising sensitivity to the axion-nucleon coupling g_{aN} , even though the generated field amplitude remains much smaller than that of ambient axion dark matter.

In this Letter, we will derive both the production and response in NMR devices, present the projected sensitivities for specific experimental setups, and examine their dependence on experimental parameters.

Axion field production. We start with the production of the axion field from a spin-precessing source. The axial current interaction between an axion a and a nucleon N is [30]

$$\mathcal{L} \supset g_{aN} \partial_\mu a \bar{N} \gamma^\mu \gamma^5 N. \quad (1)$$

This interaction gives the field equation for axions,

$$(\partial_t^2 - \nabla^2 + m_a^2)a(t, \mathbf{x}) = -g_{aN} \partial_\mu j^{5\mu}(t, \mathbf{x}), \quad (2)$$

where $j^{5\mu} = \bar{N} \gamma^\mu \gamma^5 N$ is the nucleon axial current. For a nonrelativistic source medium with spin-1/2 nuclei, j^{5i} is twice the nuclear spin density, while the temporal component j^{50} is negligible. The source term is thus a “spin charge” distribution induced by the spin polarization of the medium, $\rho_s(t, \mathbf{x}) \equiv -\partial_i j^{5i}(t, \mathbf{x})$, exciting the axion field with a coupling strength g_{aN} [31].

Supposing the source medium is uniform and the nuclear spins are precessing in alignment, then $j^{5i}(t, \mathbf{x}) \simeq n_N \sigma^i(t)$ inside the source, with n_N the nuclei density and $\sigma(t)$ a unit vector pointing toward the direction of nuclear spins at time t . In this case, the “spin charge” only distributes at the boundary of the source medium as surface charge, $\Sigma_s(t, \mathbf{x}) \equiv n_N \hat{\mathbf{n}} \cdot \sigma(t)$, where $\hat{\mathbf{n}}$ is the unit normal vector to the surface of the source at position $\mathbf{x}$.

If the nuclear spins precess coherently at frequency ω_0 , the oscillating “spin charge” distribution will generate a monochromatic oscillating axion field at the same frequency. We separate temporal and spatial dependence using a complex representation, writing $\sigma(t) = e^{-i\omega_0 t} \sigma$ and $a(t, \mathbf{x}) = e^{-i\omega_0 t} a(\mathbf{x})$, where the physical quantities are obtained by taking the real parts.

The axion field produced by such a source, solved from Eq. (2), is

$$a(\mathbf{x}) = g_{aN} n_N \oint d^2 x' \frac{\hat{\mathbf{n}}' \cdot \sigma}{4\pi |\mathbf{x} - \mathbf{x}'|} e^{ik_0 |\mathbf{x} - \mathbf{x}'|}, \quad (3)$$

where $k_0 \equiv \sqrt{\omega_0^2 - m_a^2}$ (if $\omega_0 \geq m_a$) or ik_0 with $\kappa_0 \equiv \sqrt{m_a^2 - \omega_0^2}$ (if $\omega_0 < m_a$), $\hat{\mathbf{n}}'$ is the unit normal vector to the source surface at $\mathbf{x}'$, and the integration is performed over the source surface. In writing down Eq. (3), we assumed $\omega_0 > 0$, as it no longer represents the retarded solution if $\omega_0 < 0$ [31]. The axion solution is divided into

two regimes, depending on the relative magnitudes of the frequency ω_0 and axion mass m_a . For small axion masses ($m_a \leq \omega_0$), the axion field is a propagating wave outward from the “spin charge” with a wave number k_0 and an amplitude decreasing with distance. For large axion masses ($m_a > \omega_0$), the exponential decay factor $e^{-\kappa_0 |\mathbf{x} - \mathbf{x}'|}$ indicates that the field is evanescent, decaying exponentially over a distance $1/\kappa_0$.

If the size of the source as well as the source-detector distance are both small enough such that $\omega_0 |\mathbf{x} - \mathbf{x}'| \ll 1$, for all masses $m_a \lesssim 1/R$ the axion field satisfies the near-field limit ($k_0 |\mathbf{x} - \mathbf{x}'| \ll 1$ or $\kappa_0 |\mathbf{x} - \mathbf{x}'| \ll 1$) and reduces to

$$a(\mathbf{x}) = g_{aN} n_N \oint d^2 x' \frac{\hat{\mathbf{n}}' \cdot \sigma}{4\pi |\mathbf{x} - \mathbf{x}'|}. \quad (4)$$

In this regime, $a(\mathbf{x})$ is independent of its mass m_a and the precession frequency ω_0 (except for the harmonic time dependence factor we factor out). The produced axion field can be considered as sourced instantaneously by the “spin charge” distribution. For a GHz precession frequency, this can be achieved with centimeter-scale NMR experiments.

The near-zone field, after including the harmonic time dependence back into Eq. (4), can be intuitively understood by drawing an analogy with electrostatics. For any instant t , the axion field and its gradient are analogous, respectively, to the electric potential Φ and the negative electric field $-\mathbf{E}$ induced by a static electric charge distribution identical to the “spin charge” distribution, $\Sigma_s(t, \mathbf{x}) = n_N \hat{\mathbf{n}} \cdot \sigma(t)$, at that moment. This correspondence is particularly useful for understanding the configuration of $\nabla a(t, \mathbf{x})$ in the near zone, which drives the NMR in the detection medium—at any moment, positive “spin charge” draws a component that points toward it, while negative “spin charge” leads to a component that points away from it, as sketched in Fig. 1 (left).

As the axion source, we consider a cylindrical system, as shown in Fig. 1 (left), where the nuclear spins are polarized by a static magnetic field $\mathbf{B}_0 = B_0 \hat{\mathbf{z}}$ in the $+z$ direction and the NMR is driven by an alternating magnetic field $\mathbf{B}_1(t)$ in the xy plane. The frequency of spin precession and the produced axion field ω_0 is given by the Larmor frequency of the nuclei, $\omega_0 = |\gamma B_0|$, where γ is the gyromagnetic ratio of the nuclei. We further assume the source sample is hyperpolarized and the nuclear spins maintain a 90° tilt and precess completely in the xy plane, such that $\sigma = \hat{\mathbf{x}} \mp i\hat{\mathbf{y}}$ ($-$ for $\gamma > 0$ and $+$ for $\gamma < 0$) [41]. The resulting near-field $\nabla a(t, \mathbf{x})$ rotates in the opposite direction as the bulk nuclear spins at the frequency ω_0 , inferred from the rotating surface “spin charge” on the lateral surface of the source cylinder.

For illustration of the magnitude of the produced axion field, we take the example of such a spin source with a height-to-radius ratio $H/R = 3$ and its center located at the origin. We assume a number density of

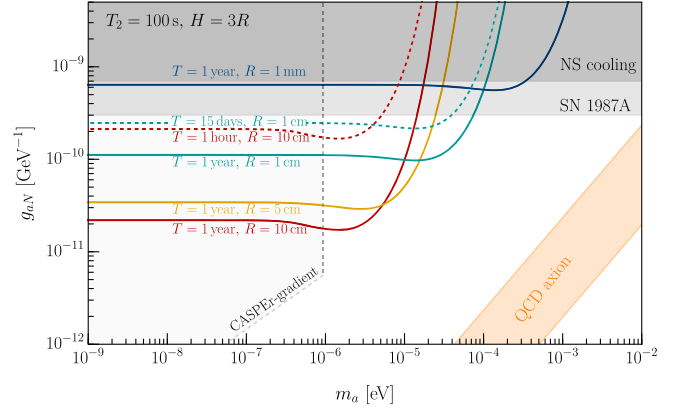
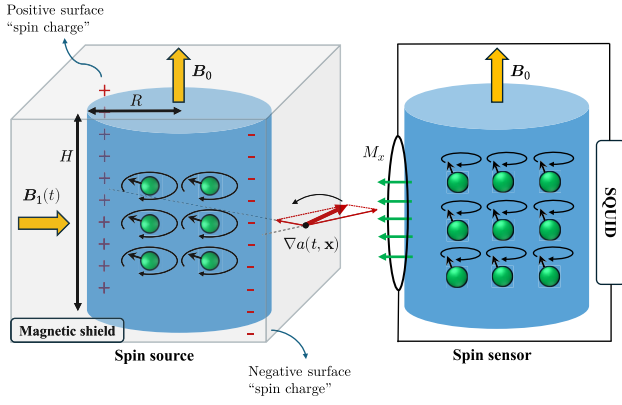


FIG. 1. Summary of the experimental proposal. Left: the sketch of the experiment. Right: the projection of the 95% C.L. upper limit of axion-nucleon coupling g_{aN} as a function of axion mass m_a . The material in both the spin source and spin sensor is ^{129}Xe , with a number density $n_N = 1.3 \times 10^{22} \text{ cm}^{-3}$ and a Larmor frequency $\omega_0 = 0.74 \text{ GHz}$, which corresponds to a static magnetic field of $B_0 = 10 \text{ T}$, and we take the transverse relaxation time to be $T_2 = 100 \text{ s}$. For ^{129}Xe , $g_{aN} \simeq 0.600g_{an} + 0.046g_{ap}$, where g_{an} and g_{ap} are axion-neutron and axion-proton couplings. The source and the detector cylinders' radii are varied by $R = 10 \text{ cm}$, 5 cm , 1 cm , and 1 mm , with their heights fixed at $H = 3R$. The two cylinder axes are separated by $3R$ and the pickup loop is placed at the near end of the sensor medium. The observation time for the solid lines is $T = 1 \text{ yr}$, while the dashed lines use the same setups as the solid lines in the same colors, but run for different observation times. The astrophysical bounds on g_{aN} from neutron star cooling [38] and SN 1987A [39] are shown in gray-shaded regions, the same as those taken in Ref. [40]. The QCD axion parameter space is shown in the orange-shaded region. For comparison, the CASPER-gradient projection [26] is overlaid, which relies on the assumption of axions as 100% dark matter.

$n_N = 1.3 \times 10^{22} \text{ cm}^{-3}$, with both the field and the gradient scaling linearly with the density. Carrying out the numerical integral of Eq. (4), the near-field solution outside the source and its gradient display a power-law behavior along the radial distance on the xy plane,

$$a(x\hat{\mathbf{x}}) \simeq 2.4 \text{ } \mu\text{eV} \left(\frac{x}{R}\right)^{-1.6} \left(\frac{g_{aN}}{10^{-10} \text{ GeV}^{-1}}\right) \left(\frac{R}{10 \text{ cm}}\right), \quad (5)$$

$$\partial_x a(x\hat{\mathbf{x}}) \simeq -7.6 \text{ } \mu\text{eV}^2 \left(\frac{x}{R}\right)^{-2.6} \left(\frac{g_{aN}}{10^{-10} \text{ GeV}^{-1}}\right). \quad (6)$$

Although the axion field is significantly smaller than that of any cosmological population of axion dark matter field, its gradient is relatively enhanced by $|\partial_x a/a| \sim 1/R \sim 1 \text{ } \mu\text{eV}(R/10 \text{ cm})^{-1}$ from its spatial inhomogeneity, which is about 4 orders of magnitude larger than that for axion dark matter with μeV masses [42]. Note that while the experimental setups discussed in this paper are around the near-field limit, we use the full-form solution, Eq. (3), to make our final projections.

Axion field detection. For a nonrelativistic nucleus, the interaction Hamiltonian arising from the Lagrangian in Eq. (1) is

$$H_{\text{int}} = -2g_{aN}\nabla a(t, \mathbf{x}) \cdot \mathbf{S}, \quad (7)$$

where $\mathbf{S}$ is the nuclear spin operator. Comparison with the Zeeman interaction of spins in a magnetic field suggests

interpreting the influence of the axion as inducing an effective magnetic field,

$$\mathbf{B}_a(t, \mathbf{x}) = \frac{2g_{aN}}{\gamma} \nabla a(t, \mathbf{x}). \quad (8)$$

The oscillatory gradient of the axion field derived in the above section will generate this effective magnetic field that can drive another NMR device nearby as a spin sensor.

We consider a spin sensor analogous to the production source we discussed above—a highly polarized homogeneous medium with gyromagnetic ratio γ , engulfed in a uniform magnetic field $\mathbf{B}_0 = B_0 \hat{\mathbf{z}}$ in the $+z$ direction. We assume the source and the detection medium have the same Larmor frequency ω_0 so that the nuclear spins in the detection medium are resonantly driven.

Following a similar procedure to that in Ref. [40], we solve the Bloch equation of the detection material driven by the axion-induced magnetic field $\mathbf{B}_a$ [31]. The axion-induced transverse magnetization M_x is

$$M_x(t, \mathbf{x}) \simeq g_{aN} M_0 T_2 |\mathcal{A}(\mathbf{x})| (1 - e^{-t/T_2}) \sin(\omega_0 t + \varphi(\mathbf{x})), \quad (9)$$

where $M_0 = n_N \gamma / 2$ is the initial magnetization along the z -axis for a hyperpolarized spin-1/2 sample, T_2 is the transverse relaxation time, $\mathcal{A}(\mathbf{x}) \equiv \pm \partial_x a(\mathbf{x}) + i \partial_y a(\mathbf{x})$ (+ for $\gamma > 0$, − for $\gamma < 0$) is the combination of the derivatives of the axion field that drives M_x , $\varphi(\mathbf{x})$ is

the phase of $\mathcal{A}(\mathbf{x})$, and we have taken the real part of M_x as the physical solution.

In the short-time limit $t \ll T_2$, the amplitude of the transverse magnetization grows linearly in time since $1 - e^{-t/T_2} \simeq t/T_2$, while for $t \gg T_2$, the growth saturates and $M_x(t)$ becomes monochromatic,

$$M_x(t, \mathbf{x}) \simeq g_{aN} M_0 T_2 |\mathcal{A}(\mathbf{x})| \sin(\omega_0 t + \varphi(\mathbf{x})). \quad (10)$$

We analyze this saturated signal in the frequency domain. Following the definition and procedure in Ref. [40], we calculate the power spectrum density (PSD) of M_x in the $t \gg T_2$ limit, for any location [31],

$$P_k^a(\mathbf{x}) \simeq \frac{g_{aN}^2 M_0^2 T_2^2 |\mathcal{A}(\mathbf{x})|^2}{\Delta \omega_k^2 T} \sin^2\left(\frac{\Delta \omega_k T}{2}\right), \quad (11)$$

where T is the observation time, the subscript k corresponds to the frequency index with frequency $\omega_k \equiv 2\pi k/T$, $\Delta \omega_k \equiv \omega_k - \omega_0$ is the frequency difference, and we only retain the term that peaks at $k > 0$. The monochromatic signal lies dominantly in a single bin at $k_* = \omega_0 T/2\pi$ where the PSD peaks.

Experimental sensitivity. For illustration, we consider a double-cylindrical NMR setup consisting of a spin source and a spin sensor with the same geometry, as shown in Fig. 1 (left). To match the Larmor frequencies of nuclear spins in both devices, the source and sensor media are made of the same material and are polarized in the same static magnetic field $\mathbf{B}_0$ aligned with the cylinders' longitudinal axis (z -axis) [43]. A transverse oscillating magnetic field $\mathbf{B}_1(t)$ is applied to drive nuclear spin precession inside the spin source, which is enclosed within a magnetic shield [31] that prevents $\mathbf{B}_1(t)$ from affecting the spin sensor [46] and is appropriately modulated to keep a stably sustained coherent spin precession, and we assume the spins to keep hyperpolarized in the xy plane [47]. The spin sensor is placed adjacent to the spin source (along the x -axis) to maximize the axion-induced signal. The source and sensor could be either within a single solenoid or one could employ two individual solenoids tuned to the same Larmor frequency.

The transverse nuclear magnetization M_x in the sensor medium driven by the axion-induced magnetic field $\mathbf{B}_a$ is measured with a SQUID magnetometer, which probes the magnetic field induced by M_x at its pickup coil. Since the axion field declines with the distance from the spin source, the transverse magnetization M_x in the detection sample and its induced magnetic field are not uniform along the x direction. To optimize signal strength, we consider a SQUID pickup coil at the near end of the sample, as depicted in Fig. 1 (left) [31].

To present the prospective detection sensitivity to g_{aN} , we choose the source and the detection media to

be pure ^{129}Xe , which has a number density of $n_N = 1.3 \times 10^{22} \text{ cm}^{-3}$ and a nuclear gyromagnetic ratio of $\gamma = -7.40 \times 10^7 \text{ s}^{-1} \text{ T}^{-1}$, and we assume the transverse relaxation time of the detection sample to be $T_2 = 100 \text{ s}$ (as chosen in other proposed axion spin-precession experiments [23,40]). We take $B_0 = 10 \text{ T}$, giving a Larmor frequency $\omega_0 = 0.74 \text{ GHz}$ for ^{129}Xe nuclei, and assume that both the source and detection samples are hyperpolarized.

The results are shown in Fig. 1 (right), where we present the projected sensitivities for a pair of cylindrical NMR devices with radius $R = 0.1, 1, 5, 10 \text{ cm}$, height $H = 3R$, and their cylinder axes separated by $3R$. The 95% C.L. upper limits are derived by the Bayesian likelihood analysis [31]. For small axion masses ($m_a < \omega_0$), the experiment is in the near-field regime of the produced axion field [53]. With a 15-day integration time, a pair of $R = 1 \text{ cm}$ NMR devices can achieve sensitivity below axion-nucleon couplings g_{aN} constrained from cooling of neutron stars and observations of supernova 1987A [54]. A similar conclusion can be reached with only a 1 hour integration time and a pair of $R = 10 \text{ cm}$ devices.

Since the produced axion field is coherent and oscillates at the Larmor frequency ω_0 , independently of its mass, one can achieve a broad-band axion search with a single driving frequency and run the experiment for long durations. Thus, we also project the sensitivities for a one-year integration time with the aforementioned setups, which are greatly improved.

The maximal axion mass that can be probed with this configuration is set by the inverse radius of the spin source, $m_a \sim 1/R$, above which the exponential suppression in Eq. (3) becomes significant. Therefore, by decreasing the radius of the cylinder, searches can be readily extended to larger axion masses. With $1/R \simeq 20 \mu\text{eV} (1 \text{ cm}/R)$, one can probe axions with masses above the maximal NMR frequency that can be achieved experimentally ($\sim 1 \mu\text{eV}$), which sets the maximal detectable axion mass for CASPer [55].

The sensitivity scaling of the experiment with various parameters can be understood using the signal-to-noise ratio (SNR). The signal lies predominantly in the frequency bin of $k_* = \omega_0 T/2\pi$, such that $\text{SNR} = P_{k_*}^a/B_{k_*}$, with $B_k = P_k^{\text{SQ}} + P_k^{\text{SP}}$ the background noise PSD, where P_k^{SQ} is the SQUID noise and P_k^{SP} is the spin projection noise [56]. Setting the SNR to a constant threshold and plugging in the expressions for $P_{k_*}^a$ and B_{k_*} (which is dominated by $P_{k_*}^{\text{SP}}$ for the benchmark parameters we consider), we can solve for the scaling behavior of the experimental sensitivity to g_{aN} [31],

$$g_{aN} \propto n_N^{-3/4} T_2^{-1/4} T^{-1/4} V^{-1/4} |\mathcal{G}(\mathbf{x}_0)|^{-1/2}, \quad (12)$$

where V is the volume of the detection sample and $\mathcal{G}(\mathbf{x}) \equiv \mathcal{A}(\mathbf{x})/g_{aN} n_N$ is a dimensionless function that depends on

ω_0 , m_a , and the geometry of the source sample, H/R . In the configuration we consider, $\mathcal{G}(\mathbf{x})$ is evaluated at the near end of the spin sensor, $\mathbf{x}_0 = (2R, 0, 0)$, with the origin at the center of the source cylinder.

The sensitivity scales simply with n_N , T_2 , T , and V , which account for the scaling of curves in Fig. 1 for different choices of device sizes and interrogation times, while its dependence on ω_0 , m_a , and geometry is implicitly contained by $|\mathcal{G}(\mathbf{x}_0)|$. For example, in the near-field limit, with $\mathbf{x}_0 = (2R, 0, 0)$, $|\mathcal{G}(\mathbf{x}_0)|$ attains its maximum at $H/R \simeq 3$, motivating the benchmark source-sensor geometry used for the projections. We provide detailed discussions on the general behavior of $|\mathcal{G}(\mathbf{x}_0)|$ in Supplemental Material [31].

Specific to the 95% C.L. upper limits of g_{aN} derived with a Bayesian likelihood framework in Fig. 1 (right), we take the near-field value, $|\mathcal{G}(\mathbf{x}_0)| \simeq 0.05$, for the optimized height-to-radius ratio $H/R = 3$, which corresponds to the small m_a region for the aforementioned configurations, then

$$g_{aN}^{(\text{UL})} \simeq 2 \times 10^{-11} \text{ GeV}^{-1} \left(\frac{n_N}{1.3 \times 10^{22} \text{ cm}^{-3}} \right)^{-3/4} \times \left(\frac{T_2}{100 \text{ s}} \right)^{-1/4} \left(\frac{T}{1 \text{ yr}} \right)^{-1/4} \left(\frac{V}{0.01 \text{ m}^3} \right)^{-1/4}, \quad (13)$$

where $V \simeq 0.01 \text{ m}^3$ corresponds to the $R = 10 \text{ cm}$ configuration.

Discussion. The sensitivity reach can be further improved in future experiments by leveraging an NMR sample material with a larger nucleus density n_N and a longer transverse relaxation time T_2 , interrogating for an integration time T , or building a spin sensor with a larger effective volume V that is efficiently transversely magnetized by the axion field. Since the axion field produced by the spin source declines as a power law of x/R , the maximal effective volume V that can be achieved should be about the same order of magnitude as the volume of the spin source. In addition, the size of the effective volume $\sim V^{1/3}$ is also limited by the wavelength of the produced axion field $\sim 1/k_0$, beyond which M_x is out of phase and has no positive contribution to the signal [31].

While we adopt $\omega_0 \sim \text{GHz}$ —about the highest NMR frequencies achieved experimentally—to obtain the sensitivities in Fig. 1 (right), a lower ω_0 would neither weaken the sensitivity reach nor reduce the maximal axion mass accessible to the setups. As long as $\omega_0 R \ll 1$, the low-mass sensitivities are governed by the near-field value, $|\mathcal{G}(\mathbf{x}_0)| \simeq 0.05$, and the upper mass limit is set by

$m_a \sim 1/R$, both independent of ω_0 . However, if, with either a larger Larmor frequency ω_0 or a larger size of samples R (which requires a precisely uniform static magnetic field of larger size in larger volume), future spin source setups could achieve $\omega_0 R \simeq 2.5$, where $|\mathcal{G}(\mathbf{x}_0)|$ (for $H/R = 3$) reaches its peak value of $\simeq 0.8$, sensitivities in the small m_a region can be strengthened by a factor of 4 compared with the near-field ones. On the other hand, a larger ω_0 (and hence smaller radius) can extend the sensitivity reach to larger axion masses, since the exponential decay occurs only when m_a exceeds ω_0 and $\kappa_0 = \sqrt{m_a^2 - \omega_0^2}$ exceeds $1/R$.

Although the experiment can, in principle, be run continuously for a total observation time due to the coherent nature of the produced axion field, as we show in Supplemental Material [31], the sensitivity of a continuous experiment for an observation time T is equivalent to that of multiple separated experiments with a total integration time T . This makes the experiment more applicable and reliable since one can calibrate the apparatus daily or hourly and conduct multiple experiments simultaneously with multiple sets of devices, without weakening the ultimate detection sensitivity.

In summary, we solved for the full dynamical picture of axion production and detection with nuclear spin precession and found that by combining two NMR experiments, one can achieve considerable sensitivity to axions solely through the axion-nucleon coupling. It is also worth noting that combining NMR devices with resonant cavities or other techniques could allow for the production and detection of axions through both the axion-nucleon and axion-photon couplings, thereby enabling probes of the product of these coupling strengths with possibly better sensitivities.

Acknowledgments. The authors thank Russell Bowers, Peter Graham, Yoonseok Lee, Amalia Madden, and Surjeet Rajendran for useful discussions. The research of J. D. is supported in part by the U.S. Department of Energy Grant No. DESC0025569. J. D. would like to thank the Aspen Center for Physics for hospitality during the completion of part of this work. The research of F. Y. is supported in part by the U.S. Department of Energy under grant DE-SC0022148 at the University of Florida and the NSF Grant No. PHY-2412701 at the University of Notre Dame.

Data availability. There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

- [1] R. D. Peccei and H. R. Quinn, *CP* conservation in the presence of instantons, *Phys. Rev. Lett.* **38**, 1440 (1977).
- [2] S. Weinberg, A new light boson?, *Phys. Rev. Lett.* **40**, 223 (1978).
- [3] F. Wilczek, Problem of strong *P* and *T* invariance in the presence of instantons, *Phys. Rev. Lett.* **40**, 279 (1978).
- [4] J. Preskill, M. B. Wise, and F. Wilczek, Cosmology of the invisible axion, *Phys. Lett.* **120B**, 127 (1983).
- [5] L. F. Abbott and P. Sikivie, A cosmological bound on the invisible axion, *Phys. Lett.* **120B**, 133 (1983).
- [6] M. Dine and W. Fischler, The not so harmless axion, *Phys. Lett.* **120B**, 137 (1983).
- [7] P. Svrcek and E. Witten, Axions in string theory, *J. High Energy Phys.* **06** (2006) 051.
- [8] K. Van Bibber, N. R. Dagdeviren, S. E. Koonin, A. Kerman, and H. N. Nelson, Proposed experiment to produce and detect light pseudoscalars, *Phys. Rev. Lett.* **59**, 759 (1987).
- [9] F. Hoogeveen, Terrestrial axion production and detection using RF cavities, *Phys. Lett. B* **288**, 195 (1992).
- [10] P. Sikivie, D. B. Tanner, and K. van Bibber, Resonantly enhanced axion-photon regeneration, *Phys. Rev. Lett.* **98**, 172002 (2007).
- [11] K. Ehret *et al.* (ALPS Collaboration), Resonant laser power build-up in ALPS: A ‘Light-shining-through-walls’ experiment, *Nucl. Instrum. Methods Phys. Res., Sect. A* **612**, 83 (2009).
- [12] F. Caspers, J. Jaeckel, and A. Ringwald, Feasibility, engineering aspects and physics reach of microwave cavity experiments searching for hidden photons and axions, *J. Instrum.* **4**, P11013 (2009).
- [13] K. Ehret *et al.*, New ALPS results on hidden-sector lightweights, *Phys. Lett. B* **689**, 149 (2010).
- [14] R. Bähre *et al.*, Any light particle search II—technical design report, *J. Instrum.* **8**, T09001 (2013).
- [15] M. Betz, F. Caspers, M. Gasiot, M. Thumm, and S. W. Rieger, First results of the CERN Resonant Weakly Interacting sub-eV Particle Search (CROWS), *Phys. Rev. D* **88**, 075014 (2013).
- [16] R. Janish, V. Narayan, S. Rajendran, and P. Riggins, Axion production and detection with superconducting RF cavities, *Phys. Rev. D* **100**, 015036 (2019).
- [17] A. Arvanitaki and A. A. Geraci, Resonantly detecting axion-mediated forces with nuclear magnetic resonance, *Phys. Rev. Lett.* **113**, 161801 (2014).
- [18] A. A. Geraci *et al.* (ARIADNE Collaboration), Progress on the ARIADNE axion experiment, *Springer Proc. Phys.* **211**, 151 (2018).
- [19] A. Almasi, J. Lee, H. Winarto, M. Smiciklas, and M. V. Romalis, New limits on anomalous spin-spin interactions, *Phys. Rev. Lett.* **125**, 201802 (2020).
- [20] A. Arvanitaki, J. Engel, A. A. Geraci, A. Madden, A. Hepburn, and K. Van Tilburg, The ferroaxionic force, *arXiv:2411.10516*.
- [21] H. Su, M. Jiang, Y. Wang, Y. Huang, X. Kang, W. Ji, X. Peng, and D. Budker, New constraints on axion-mediated spin interactions using magnetic amplification, *Phys. Rev. Lett.* **133**, 191801 (2024).
- [22] Z. Xu, X. Heng, G. Tian, D. Gong, L. Cong, W. Ji, D. Budker, and K. Wei, Constraints on axion mediated dipole-dipole interactions, *Phys. Rev. Lett.* **134**, 181801 (2025).
- [23] D. Budker, P. W. Graham, M. Ledbetter, S. Rajendran, and A. Sushkov, Proposal for a Cosmic Axion Spin Precession Experiment (CASPER), *Phys. Rev. X* **4**, 021030 (2014).
- [24] P. W. Graham and S. Rajendran, New observables for direct detection of axion dark matter, *Phys. Rev. D* **88**, 035023 (2013).
- [25] C. Abel *et al.*, Search for axionlike dark matter through nuclear spin precession in electric and magnetic fields, *Phys. Rev. X* **7**, 041034 (2017).
- [26] D. F. Jackson Kimball *et al.*, Overview of the Cosmic Axion Spin Precession Experiment (CASPER), *Springer Proc. Phys.* **245**, 105 (2020).
- [27] T. Wu *et al.*, Search for axionlike dark matter with a liquid-state nuclear spin comagnetometer, *Phys. Rev. Lett.* **122**, 191302 (2019).
- [28] A. Garcon *et al.*, Constraints on bosonic dark matter from ultralow-field nuclear magnetic resonance, *Sci. Adv.* **5**, eaax4539 (2019).
- [29] D. Aybas *et al.*, Search for axionlike dark matter using solid-state nuclear magnetic resonance, *Phys. Rev. Lett.* **126**, 141802 (2021).
- [30] Axions can also couple to the electric dipole moment of nucleons through the operator $-ig_d a \bar{N} \sigma_{\mu\nu} \gamma^5 N F^{\mu\nu} / 2$. However, this coupling is stringently constrained by astrophysical observations, and we estimate that the sensitivities of the configurations discussed in this work are insufficient to surpass these constraints.
- [31] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/y8rj-rpfg> for more discussions on the “spin charge” picture; for details on the axion production solution; for details in how to solve the Bloch equation under the axion-induced magnetic field; for the derivation of the PSD of M_x ; for discussions on the effectiveness of the shielding; for details on noise sources, which includes Refs. [32–34]; for details, which includes Ref. [35]; for the relation between the effective axion-nucleon coupling g_{aN} and the axion-proton g_{ap} and axion-neutron g_{an} couplings, which includes Refs. [36,37]; for the geometrical dependence of the signal; for discussions on the magnetic flux from the detection sample; for discussions on the computation of the projected sensitivity using discrete time segments.
- [32] S. M. Anton *et al.*, Magnetic flux noise in dc SQUIDS: Temperature and geometry dependence, *Phys. Rev. Lett.* **110**, 147002 (2013).
- [33] Y. Kahn, B. R. Safdi, and J. Thaler, Broadband and resonant approaches to axion dark matter detection, *Phys. Rev. Lett.* **117**, 141801 (2016).
- [34] M. Braun and J. König, Faraday-rotation fluctuation spectroscopy with static and oscillating magnetic fields, *Phys. Rev. B* **75**, 085310 (2007).
- [35] 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).
- [36] Y. V. Stadnik and V. V. Flambaum, Nuclear spin-dependent interactions: Searches for WIMP, axion and topological defect dark matter, and tests of fundamental symmetries, *Eur. Phys. J. C* **75**, 110 (2015).
- [37] B. S. Hu, J. Padua-Argüelles, S. Leutheusser, T. Miyagi, S. R. Stroberg, and J. D. Holt, *Ab initio* structure factors for

- spin-dependent dark matter direct detection, *Phys. Rev. Lett.* **128**, 072502 (2022).
- [38] M. Buschmann, C. Dessert, J. W. Foster, A. J. Long, and B. R. Safdi, Upper limit on the QCD axion mass from isolated neutron star cooling, *Phys. Rev. Lett.* **128**, 091102 (2022).
- [39] A. Lella, P. Carenza, G. Co', G. Lucente, M. Giannotti, A. Mirizzi, and T. Rauscher, Getting the most on supernova axions, *Phys. Rev. D* **109**, 023001 (2024).
- [40] J. A. Dror, S. Gori, J. M. Leedom, and N. L. Rodd, Sensitivity of spin-precession axion experiments, *Phys. Rev. Lett.* **130**, 181801 (2023).
- [41] This gives $\sigma(t) = e^{-i\omega_0 t}(\hat{x} \mp i\hat{y})$, which corresponds to the spins rotating in the clockwise (or counterclockwise) direction around $\hat{z}$ when viewed from above if $-$ (or $+$) is taken on, consistent with the dynamics of spins described by the Bloch equation for $\gamma > 0$ (or < 0). Note that in this complex representation, $\text{Re}[\sigma(t)]$ is a unit vector.
- [42] For axion dark matter, the field amplitude is $a = \sqrt{2\rho_{\text{DM}}}/m_a \sim 10^3 \text{ eV}(m_a/1 \mu\text{eV})^{-1}$ and the field gradient is $|\nabla a| \sim m_a v_0 a_0 \sim 10^{-6} \text{ eV}^2$. Additionally, $|\nabla a|/a \sim m_a v_0 \sim 10^{-9} \text{ eV}(m_a/1 \mu\text{eV})$. In these numerical expressions, we assumed an expected local density $\simeq 0.3 \text{ GeV}/\text{cm}^3$ and a characteristic velocity of $v_0 \sim 10^{-3}$.
- [43] The Q -factor of NMR for a polarized sample with $T_2 = 100 \text{ s}$ and Larmor frequency $\omega_0 \sim \text{GHz}$ is $Q \equiv \omega_0 T_2/2 \sim 10^9$, resulting in a resonance width $\Delta\omega/\omega_0 = Q^{-1} \sim 10^{-9}$. The longitudinal field $\mathbf{B}_0$ for the source and the sensor thus needs to be matched with a relative precision of 10^{-9} to satisfy the resonance condition for the sensor medium. Current superconducting magnets with $\mathcal{O}(10) \text{ T}$ magnetic field can achieve 10^{-9} uniformity within centimeter-scale volume with $10^{-8}/\text{hour}$ shift [44], and 10^{-6} uniformity within decimeter-scale volume with $10^{-9}/\text{hour}$ shift [45]. Persistent efforts are being made to develop more GHz-class NMR magnets with field inhomogeneity below 10^{-9} to resolve the narrow spectral lines in structure biology [44]. Operating at a lower ω_0 can relax the requirement on field homogeneity. As we will discuss in the next section, reducing ω_0 does not degrade our projected sensitivity.
- [44] P. Wikus, W. Frantz, R. Kümmerle, and P. Vonlanthen, Commercial gigahertz-class nmr magnets, *Supercond. Sci. Technol.* **35**, 033001 (2022).
- [45] N. Boulant, F. Mauconduit, V. Gras, A. Amadon, C. Le Ster, M. Luong, A. Massire, C. Pallier, L. Sabatier, M. Bottlaender *et al.*, In vivo imaging of the human brain with the isult 11.7-T MRI scanner, *Nat. Methods* **21**, 2013 (2024).
- [46] The $\mathbf{B}_0$ field applied to the spin source can magnetize the shielding material, which induces a residual field, altering the total longitudinal field. This effect is negligible as long as the ratio of the residual field to B_0 remains below Q^{-1} . Otherwise, it can be solved by introducing an identical magnetic shield for the spin sensor or tuning the longitudinal field with a relative precision of $\mathcal{O}(Q^{-1})$.
- [47] Prolonged spin precession has been achieved for $\mathcal{O}(10^4) \text{ s}$ in ^{129}Xe and ^3He gases using spin masers, which utilize optical pumping along the z direction and a rotating transverse magnetic field with a 90° phase delay (compared to the spin precession) generated by a feedback system [48–51]. By appropriately choosing the pumping and feedback parameters, the stabilized transverse magnetization can achieve $\mathcal{O}(1)$ the initial magnetization M_0 [52]. We hope in the future this technique can be further applied to ^{129}Xe or ^3He liquids.
- [48] D. Bear, R. Stoner, R. Walsworth, V. A. Kostelecký, and C. D. Lane, Limit on lorentz and *CPT* violation of the neutron using a two-species noble-gas maser, *Phys. Rev. Lett.* **85**, 5038 (2000).
- [49] M. Rosenberry and T. Chupp, Atomic electric dipole moment measurement using spin exchange pumped masers of ^{129}Xe and ^3He , *Phys. Rev. Lett.* **86**, 22 (2001).
- [50] A. Yoshimi, K. Asahi, K. Sakai, M. Tsuda, K. Yogo, H. Ogawa, T. Suzuki, and M. Nagakura, Nuclear spin maser with an artificial feedback mechanism, *Phys. Lett. A* **304**, 13 (2002).
- [51] A. Yoshimi, T. Inoue, T. Furukawa, T. Nanao, K. Suzuki, M. Chikamori, M. Tsuchiya, H. Hayashi, M. Uchida, N. Hatakeyama *et al.*, Low-frequency ^{129}Xe nuclear spin oscillator with optical spin detection, *Phys. Lett. A* **376**, 1924 (2012).
- [52] M. Feng, L. Wu, and G. Liu, Nonlinear dynamics in an artificial feedback spin maser, *Phys. Rev. A* **111**, 023118 (2025).
- [53] For the $R = 10 \text{ cm}$ setup, the axion field in the spin sensor slightly deviates from the near-field result for small masses.
- [54] Note that the actual magnetic field that fluxes through a pickup loop near the end of the sensor medium is $\mathcal{O}(1)$ smaller than M_x (see Supplemental Material for details). Since $M_x \propto g_{aN}^2$, our sensitivities obtained by treating them as equal are slightly overestimated by an $\mathcal{O}(1)$ factor.
- [55] In the $m_a \gg 1/R$ regime, the magnetization within the detection sample decreases dramatically and the signal is only contributed by a small portion of the sample around the near end. In Fig. 1 (right), we always take V in P_k^{SP} as the volume of the entire detection sample such that the spin projection noise is underestimated in this regime. The actual sensitivities are thus decaying more rapidly when $m_a \gg 1/R$ than in Fig. 1 (right).
- [56] This still holds true with the Bayesian framework we adopt, even though our signal is deterministic rather than stochastic. We provide a detailed discussion of noise sources and the likelihood analysis in Supplemental Material.