

Earth as a Transducer for the Detection of Ultralight Bosonic Dark Matter

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Abstract

Ultralight bosonic dark matter (UBDM) that couples to electromagnetism can generate an oscillating magnetic-field signal at the Earth's surface. This is referred to as the “Earth transducer” effect, as the Earth converts UBDM into a detectable magnetic field. Similar DM-induced fields in laboratory experiments typically scale with the size L of the experiment. Because the Earth transducer signal instead scales with the large radius of the Earth, R , it is one of the most powerful direct probes of UBDM with masses $m_{\text{DM}} \lesssim 1/R \sim 3 \times 10^{-14}$ eV. It has many other favorable properties, such as high spatial and temporal coherence and robustness to atmospheric modeling. In this review, we derive the Earth transducer effect and its properties for multiple UBDM models, and we discuss current and future prospects to detect it.

Keywords: ultralight dark matter; axion; dark photon; magnetometry; geophysics

1. Introduction

An abundance of astrophysical and cosmological observations indicate the presence of a poorly understood component of the universe called dark matter (DM). Its existence is inferred from its gravitational impact on a variety of phenomena, including the motion of stars/galaxies [1,2], the merging of clusters [3], and the expansion of the universe [4]. While these observations indicate its approximate abundance and velocity [5–7], very little else is known about DM. Some of its properties of interest include the fundamental mass and spin of its constituents and its possible non-gravitational interactions.

Of the myriad models that have been proposed for DM, one class that has garnered significant interest recently is ultralight bosonic dark matter (UBDM) [8–10]. This class consists of models where DM is composed of constituent particles with masses $m_{\text{DM}} \lesssim 10$ eV.¹ In this regime, DM behaves as a classical field [11–13]. Popular UBDM candidates include dark-photon dark matter (DPDM) [14–17], axion-like particle dark matter (henceforth, simply axion DM) [18–22], and millicharged dark matter (mDM) [14,23–26].

All of these candidates may couple to electromagnetism (EM) [14,27], and a variety of laboratory experiments have been conducted or proposed to search for such interactions [28–51]. The most widely explored strategy is the use of resonant cavities [27,46,49,50]. A key feature of such experiments is that the size L of the cavity should match the Compton wavelength $\lambda_{\text{DM}} \equiv 2\pi/m_{\text{DM}}$ of the DM, which becomes increasingly difficult below $m_{\text{DM}} \sim \mu\text{eV}$. At lower masses, resonant LC circuits can be employed [33,43]. The signal in these systems scales with the size L of the experimental shield when $\lambda_{\text{DM}} \gg L$.



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In general, the presence of a conducting boundary around a laboratory experiment imposes a limitation on detecting these models of UBDM in the extremely-low-mass regime. A natural question is whether any advantage can be gained by removing the experimental shield entirely. In Ref. [52], it was shown that in this case, the Earth sets the conducting boundary of the system. This is because both the ground and ionosphere act as good conductors at extremely low frequencies. The dominant signal in this system is then an oscillating magnetic field at the Earth's surface, whose magnitude scales with the radius R of the Earth. This effect is known as the “Earth transducer” effect, as the Earth can convert UBDM into a detectable magnetic-field signal.

The Earth transducer effect provides a new method for direct detection of UBDM coupled to EM, namely, via correlated measurements from a global array of unshielded magnetometers. The UBDM signal would appear as a sharp peak in the spectrum of each magnetometer, whose amplitude exhibits a prescribed dependence on the magnetometer's geographic location. These unique properties make the signal easily distinguishable from noise. There are a number of completed and ongoing initiatives to utilize existing datasets and undertake dedicated searches to look for this effect. Refs. [53–55] performed searches in publicly available geomagnetic-field datasets maintained by the SuperMAG collaboration [56–58], and Refs. [59–61] searched data collected from the Eskdalemuir observatory [62]. In addition, the SNIPE Hunt [63,64] and GPEX [65] collaborations have taken correlated measurements with unshielded magnetometers and analyzed their datasets to search for the Earth transducer effect. In many regions of parameter space, these searches are the leading direct constraints on certain UBDM models (see Figures 3 and 4).

In this review, we summarize the current progress and outlook for probing UBDM parameter space with the Earth transducer effect. In Section 2, we begin by describing the UBDM models of interest and deriving the Earth transducer effect and its properties for these models. We also comment on methods to extend the signal modeling to higher frequencies for which $\lambda_{\text{DM}} \lesssim R$. In Section 3, we review the searches for this effect in existing datasets and dedicated searches. The results of these searches are summarized in Figures 3 and 4. In Section 4, we discuss the current state of these detection efforts and conclude the main text of the review. In Appendix A, we review the relevant properties of vector spherical harmonics (VSH).

2. Theory

In this section, we review the theoretical framework for the Earth transducer effect. This framework was first developed for DPDM in Ref. [52] and then extended to axion DM and mDM in Refs. [54,66], respectively. We begin by reviewing properties of EM-coupled UBDM models and deriving how they are manifested in experimental contexts. All of these models can be described in a common framework, namely, as an “effective current” J_{eff}^μ . In this language, standard techniques for analyzing EM problems can be applied to understand the effects of UBDM. In particular, we will see that the effects of UBDM generically scale with the size L of the experimental shield, when $\lambda_{\text{DM}} \gg L$.

We then proceed to derive the Earth transducer effect. We briefly discuss the conductive properties of the Earth's interior and atmosphere in order to understand the EM boundary conditions of the system. We then compute the Earth transducer effect for each UBDM candidate in a simple spherical model of the Earth's atmosphere. Importantly, we show that the effect scales with the radius R of the Earth, rather than the much smaller height h of the lower atmosphere. We also discuss the coherence properties of the signal, which make it distinguishable from noise. Although we compute the effect for a simplified model, we argue that for low masses, this result can be made robust to the environmental details of the Earth's atmosphere.

Finally, we review methods to extend the Earth transducer effect to higher masses, where this robustness breaks down. Ref. [67] proposed to perform a differential measurement of $\nabla \times \mathbf{B}$ in order to robustly probe the local effective current, independent of global boundary conditions. Additionally, Refs. [59,61,68] computed the Earth transducer effect in a more refined empirical model of the Earth's atmospheric conductivity profile.

2.1. UBDM Models

A number of intriguing models have been proposed for the fundamental nature of DM. In this review, we will focus on so-called UBDM models [8–10], where the constituent particles are sufficiently light that the DM behaves like a classical field. Astrophysical observations indicate that the local DM in the vicinity of the Solar System has energy density $\rho_{\text{DM}} \sim m_{\text{DM}} n_{\text{DM}} \sim 0.3 \text{ GeV}/\text{cm}^3$ and velocity $v_{\text{DM}} \sim 10^{-3}$ [5–7]. Therefore, if $m_{\text{DM}} \lesssim 10 \text{ eV}$, there exist $N_{\text{DM}} \gtrsim 1$ DM particles within each de Broglie volume $\lambda_{\text{dB}}^3 \equiv (2\pi/m_{\text{DM}}v_{\text{DM}})^3$. This implies two important properties for this DM mass range. The first is that DM candidates in this mass range must be bosonic (as fermions are restricted to occupation numbers $N_{\text{DM}} < 1$). The second is that quantum effects become unimportant and the DM behaves classically.

More specifically, the DM can be described by a classical field, such as a scalar (spin zero) field $\chi(x, t)$, or a vector (spin one) field $X_\mu(x, t)$.² To understand its kinematics, let us first focus on the case of a real scalar field χ . In the absence of interactions, this field satisfies the Klein–Gordon equation

$$\square\chi + m_{\text{DM}}^2\chi = 0. \quad (1)$$

As this is a type of wave equation, UBDM is sometimes also referred to as “wave DM”. Generically, Equation (1) is solved by a sum of plane waves

$$\chi(x, t) = \text{Re}[\sum_n \chi_{0,n} e^{i(k_n \cdot x - \omega_n t)}], \quad (2)$$

where $\omega_n^2 = |\mathbf{k}_n|^2 + m_{\text{DM}}^2$, while $\chi_{0,n}$ are various complex amplitudes. As $|\mathbf{k}_n|/m_{\text{DM}} \sim v_{\text{DM}} \ll 1$, Equation (2) simplifies to the form

$$\chi(x, t) \approx \chi_0 \cos(m_{\text{DM}}t + \alpha_0), \quad (3)$$

for a single real amplitude χ_0 and phase α_0 . In other words, we expect the DM to be spatially homogeneous and oscillating in time with frequency $f_{\text{DM}} \equiv m_{\text{DM}}/2\pi$. Note that the approximation in Equation (3) only applies for distances within a de Broglie wavelength, $|x| \lesssim \lambda_{\text{dB}}$, and times within a coherence time, $t \lesssim T_{\text{coh}} \equiv 1/(f_{\text{DM}}v_{\text{DM}}^2)$. Over larger length or time scales, χ_0 and α_0 will generically vary. The typical size of χ_0 is fixed by the local DM energy density³

$$\langle \chi_0^2 \rangle = \frac{2\rho_{\text{DM}}}{m_{\text{DM}}^2}, \quad (4)$$

where $\langle \cdot \rangle$ represents an average over several de Broglie wavelengths or coherence times.

In this review, we will consider UBDM models that couple to EM. All such models can be treated in a common framework. For any interaction $\mathcal{L}_{\text{int}}$ between a new field and the Standard Model (SM) photon A_μ , we can define an “effective current”

$$J_{\text{eff}}^\mu \equiv \partial_\nu \frac{\partial \mathcal{L}_{\text{int}}}{\partial (\partial_\nu A_\mu)} - \frac{\partial \mathcal{L}_{\text{int}}}{\partial A_\mu}. \quad (5)$$

Under a gauge transformation $A_\mu \rightarrow A_\mu + \partial_\mu \alpha$, the interaction term transforms as follows:

$$\mathcal{L}_{\text{int}} \rightarrow \mathcal{L}_{\text{int}} - J_{\text{eff}}^\mu \partial_\mu \alpha. \quad (6)$$

Gauge invariance then implies that the effective current must be conserved, i.e., $\partial_\mu J_{\text{eff}}^\mu = 0$. This effective current behaves similarly to a SM current. In particular, it modifies Maxwell's equations as

$$\nabla \cdot \mathbf{E} = \rho_{\text{eff}}, \quad (7)$$

$$\nabla \cdot \mathbf{B} = 0, \quad (8)$$

$$\nabla \times \mathbf{E} + \partial_t \mathbf{B} = 0, \quad (9)$$

$$\nabla \times \mathbf{B} - \partial_t \mathbf{E} = \mathbf{J}_{\text{eff}}, \quad (10)$$

where $J_{\text{eff}}^\mu = (\rho_{\text{eff}}, \mathbf{J}_{\text{eff}})$. Under the limitation that the DM is very weakly coupled to the SM, backreaction can be neglected and the DM can be treated as a background field. Furthermore, if background EM fields are present, then all instances of A_μ appearing in J_{eff}^μ can take their background values. Assuming these EM fields are static, the effective current will inherit the oscillatory time dependence of the DM. Then to leading order, J_{eff}^μ will behave as a background AC current density, and standard EM techniques can be applied to understand its observable effects.

In this review, we will consider three UBDM models. The first is a massive vector field A'_μ , called a (kinetically mixed) dark photon. This candidate couples to EM via the interaction⁴

$$\mathcal{L}_{\text{int},A'} = \varepsilon m_{A'}^2 A_\mu A'^\mu, \quad (11)$$

where $m_{A'}$ is the mass of the DPDM field,⁵ while $\varepsilon < 1$ parametrizes its coupling to EM. The equation of motion for A'_μ is the Proca equation,

$$\partial_\mu F'^{\mu\nu} + m^2 A'^\nu = 0, \quad (12)$$

where $F'_{\mu\nu} = \partial_\mu A'_\nu - \partial_\nu A'_\mu$ is the DPDM field strength. Taking the derivative of Equation (12) shows that $\partial_\mu A'^\mu = 0$. This implies that $A'_0 \sim v_{\text{DM}} \cdot \mathbf{A}'$, and so the temporal components of A'_μ are suppressed. The spatial components take an analogous form to Equation (3),

$$A'_i(\mathbf{x}, t) = A'_{i,0} \cos(m_{A'} t + \alpha_{i,0}), \quad (13)$$

where $A'_{i,0}$ and $\alpha_{i,0}$ are amplitudes and phases for each component $i = 1, 2, 3$. Importantly, the phases $\alpha_{i,0}$ need not be the same. This allows $\mathbf{A}'(t)$ to exhibit an elliptical polarization (in which case $|\mathbf{A}'(t)| \neq 0$ throughout its entire oscillation). From Equations (5) and (11), we see that the effective current for DPDM is given by

$$J_{\text{eff},A'}^\mu = -\varepsilon m_{A'}^2 A'^\mu. \quad (14)$$

As with A'_μ , the temporal component of Equation (14) is suppressed, i.e., there is no effective charge ρ_{eff} , and the spatial components are uniform in space. Equation (14) also makes it clear that the observable effects of DPDM should scale with $m_{A'}$. This makes detection of DPDM particularly difficult in the low-mass regime.

Our second candidate is a real scalar field a , known as an axion-like particle (or simply axion). As a real scalar, it satisfies Equation (1) and takes the form in Equation (3). This candidate couples to EM via the interaction

$$\mathcal{L}_{\text{int},a} = \frac{1}{4} g_{a\gamma} a F_{\mu\nu} \tilde{F}^{\mu\nu}, \quad (15)$$

where $g_{a\gamma}$ is the axion–photon coupling, while $\tilde{F}_{\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}F_{\rho\sigma}$. From Equation (5), the effective current corresponding to this interaction is

$$J_{\text{eff},a}^{\mu} = -\frac{1}{2}g_{a\gamma}\epsilon^{\mu\nu\rho\sigma}\partial_{\nu}aF_{\rho\sigma}, \quad (16)$$

where $\epsilon^{\mu\nu\rho\sigma}$ is the Levi-Civita symbol, or equivalently,

$$\rho_{\text{eff},a} = g_{a\gamma}\nabla a \cdot \mathbf{B}, \quad (17)$$

$$\mathbf{J}_{\text{eff},a} = -g_{a\gamma}[(\partial_t a)\mathbf{B} + \nabla a \times \mathbf{E}]. \quad (18)$$

Unlike the case of DPDM, the axion effective current depends on the EM field. In order to enhance the observable effects of axion DM, it is therefore beneficial to introduce background EM fields. Because $\nabla a \sim v_{\text{DM}} \cdot \partial_t a$, we see that the largest effect will arise from the first term in Equation (18), and so we should introduce a background magnetic field $\mathbf{B}_0$. In the case of the Earth transducer effect, this role will be played by the Earth's DC magnetic field. In contrast to $J_{\text{eff},A'}$, which is uniform in space, $J_{\text{eff},a}$ will follow the dipolar shape of the geomagnetic field.

Our final model of interest is a complex scalar field ϕ , called a millicharged particle. This field is charged under EM and so couples through its kinetic term

$$\mathcal{L}_{\text{int},\phi} = D_{\mu}\phi(D^{\mu}\phi)^*, \quad (19)$$

where $D_{\mu} = \partial_{\mu} + ie_m A_{\mu}$, while e_m is the electric charge of the mDM. The corresponding effective current is given by

$$J_{\text{eff},\phi}^{\mu} = ie_m(\phi^*\partial^{\mu}\phi - \phi\partial^{\mu}\phi^*) - 2e_m^2|\phi|^2 A^{\mu}. \quad (20)$$

As the effective current depends explicitly on A_{μ} , it is important to take care with the choice of gauge. Physically, we expect $\rho_{\text{eff}} = 0$, so that the mDM does not generate large-scale electric fields throughout the galaxy. This implies the condition

$$\text{Im}[\phi\partial_t\phi^*] = e_m|\phi|^2 A_0. \quad (21)$$

Fixing temporal Coulomb gauge, $A_0 = 0$ and $\nabla \cdot \mathbf{A} = 0$,⁶ we then see that Equations (1) and (21) imply the form⁷

$$\phi(x, t) = \phi_0 \cos(m_{\phi}t + \alpha_0), \quad (22)$$

within the non-relativistic (NR) limit $v_{\text{DM}} \ll 1$. In principle, Equation (21) can be satisfied even for a complex amplitude ϕ_0 , but the phase of ϕ_0 can always be removed by a gauge transformation, and so we will take it to be real. With the form in Equation (22), the spatial components of the effective current become

$$\mathbf{J}_{\text{eff},\phi} = -e_m^2\phi_0^2[1 + \cos(2m_{\phi}t + 2\alpha_0)]\mathbf{A} \quad (23)$$

within the NR limit. This effective current contains both DC and AC contributions. In this review, we will focus on the effect of the AC part. Note that in the mDM case, the AC part of the effective current oscillates at frequency $2m_{\phi}$, as will any response fields. This is because the relevant contribution to $J_{\text{eff},\phi}^{\mu}$ in Equation (20) depends on $|\phi|^2$. Again, to generate an observable effect, we require a background EM field, which will be the geomagnetic field in the case of the Earth transducer effect.

Now that we have unified all of our UBDM models into a common effective current framework, let us discuss the possible effects that could be observed in an experiment. All of our models have $\rho_{\text{eff}} = 0$, and so the most relevant of Maxwell's equations will be the Ampère–Maxwell law, given by Equation (10). In general, the AC effective current will source AC electric fields E_{DM} and/or magnetic fields B_{DM} inside an experimental apparatus. In this review, we will primarily focus on the extremely-low-mass regime, where the Compton wavelength of the DM exceeds the size of the experimental apparatus. More specifically, suppose the experiment is surrounded by some conducting boundary of characteristic length L . In the regime $\lambda_{\text{DM}} \gg L$, there is a simple argument that E_{DM} will generally be suppressed relative to B_{DM} (see Refs. [52,67] for more detailed arguments in the Earth case). Standard EM boundary conditions mandate that (the parallel component of) E_{DM} must vanish at the conducting boundary. As E_{DM} has frequency m_{DM} , it should grow on length scales λ_{DM} . The electric field at the center of the cavity can then be at most $(L^2/\lambda_{\text{DM}}^2)$, whereas we will see that B_{DM} scales as L^1 . Neglecting the electric field contribution, the Ampère–Maxwell law then becomes

$$\nabla \times B_{\text{DM}} = J_{\text{eff}}. \quad (24)$$

Parametrically, this implies $B_{\text{DM}} \sim J_{\text{eff}} \cdot L$. In summary, within the limit $m_{\text{DM}}L \ll 1$, we expect the dominant signal to be an AC magnetic field whose magnitude scales with L . We therefore benefit from utilizing as large an apparatus as possible.

2.2. Earth Transducer Effect

We have seen that the presence of a small conducting shield can suppress the signal from UBDM. A natural line of questioning, then, is as follows: how large can we make our experimental shield, or can we even remove the shield entirely? Inevitably, any experimental setup will have some conducting boundary. If the experiment does not contain a deliberate shield, the room inside which the experiment is conducted may act as a shield [47]. Even if the experiment is conducted outside a laboratory, the Earth itself will set a conducting boundary. This is because the Earth's interior and the upper layers of the Earth's atmosphere act as good conductors at the low frequencies ($7 \times 10^{-4} \text{ Hz} \lesssim f_{\text{DM}} \lesssim 7 \text{ Hz}$) corresponding to our mass range of interest ($3 \times 10^{-18} \text{ eV} \lesssim m_{\text{DM}} \lesssim 3 \times 10^{-14} \text{ eV}$).

The full details of the conductivity profile in/around the Earth are quite complex (see Ref. [52] for more detailed discussion). Here, we review the relevant details for our scenario. A conductive layer of the Earth's interior or atmosphere acts as an efficient shield if its conductivity is higher than the frequency of the DM response, $\sigma \gg m_{\text{DM}}$, and the layer is thicker than its skin depth (damping length), $d \gg \delta \sim \sqrt{2/\sigma m_{\text{DM}}}$. On the other hand, if a layer satisfies $\sigma \ll m_{\text{DM}}$ or $d \ll \delta$, it contributes negligible damping to the DM response and can therefore be treated as vacuum. While the crust of the Earth is not sufficiently thick to constitute an effective shield, the inner layers (mantle and core) of the Earth are certainly thick enough to function as a shield for our entire mass range of interest [70,71]. On the other hand, the lower $\sim 100 \text{ km}$ of the Earth's atmosphere has very low conductivity [72] and can therefore be treated as vacuum. The situation becomes more complicated for the ionosphere—an upper layer of the Earth's atmosphere that is highly ionized by ultraviolet radiation—due to the anisotropy and temporal modulation of its conductivity [73,74]. For masses $m_{\text{DM}} \gg 10^{-16} \text{ eV}$ ($f_{\text{DM}} \gg 0.02 \text{ Hz}$), the ionosphere is certainly thick enough to constitute an effective shield, in which case it sets a relatively spherical outer conducting boundary for the system. For lower masses, the thickness condition is satisfied for certain directional conductivities but not others, and so it becomes unclear whether the ionosphere can effectively damp the DM response. Outside the Earth's atmosphere, the interplanetary medium acts as a collisionless plasma with plasma frequency $\omega_p \sim 10^{-10} \text{ eV} \gg m_{\text{DM}}$ [75].

If the ionosphere does not set the outer boundary of the system, then the interplanetary medium will. It is worth noting that the magnetopause, which delineates the boundary between the Earth's atmosphere and the interplanetary medium, is highly aspherical [76,77], which, in principle, can lead to a complicated geometry for the system.

To begin, let us compute the DM response for a simplified geometry of the Earth–ionosphere cavity, and then we will comment on how this result applies in the more realistic case. We take the Earth to be a perfectly conducting sphere of radius $R = 6400$ km, and we take the ionosphere to be a concentric perfectly conducting shell with inner radius $R + h$, such that the lower atmosphere is a vacuum region of height $h \sim 100$ km $\ll R$. We will solve Equation (24) in this geometry (for the appropriate J_{eff}). An important result will be that $B_{\text{DM}} \propto R$, as opposed to $B_{\text{DM}} \propto h$. Before considering the full computation, let us review a simple geometric argument from Ref. [52], shown in the left panel of Figure 1, which demonstrates this fact in the case of DPDM. Suppose that the DPDM is polarized along the Earth's rotational axis, so that $J_{\text{eff},A'}$ also points in this direction. Consider a Gaussian surface situated just above the Earth's surface, which covers the northern hemisphere. The boundary of this surface is an Ampèrian loop that wraps around the equator. Ampère's law, in integral form, states that

$$J_{\text{eff}} \cdot R^2 \sim \int J_{\text{eff}} \cdot dS = \oint B_{\text{DM}} \cdot dl \sim B_{\text{DM}} \cdot R. \quad (25)$$

It is then clear that $B_{\text{DM}} \sim J_{\text{eff}} \cdot R$. Note that this estimate does not depend on h . In particular, the argument holds even under the limitation $h \rightarrow 0$. This is important not only to show that the Earth transducer effect is enhanced by the large radius of the Earth but also because the result is not sensitive to the exact value of h (which depends on detailed atmospheric modeling). In fact, the argument still applies even if the ionosphere is deformed away from a sphere (as long as E_{DM} can still be ignored in the Ampère–Maxwell law). This demonstrates how the predictions of the Earth transducer effect can be made robust to environmental details of the Earth's atmosphere.

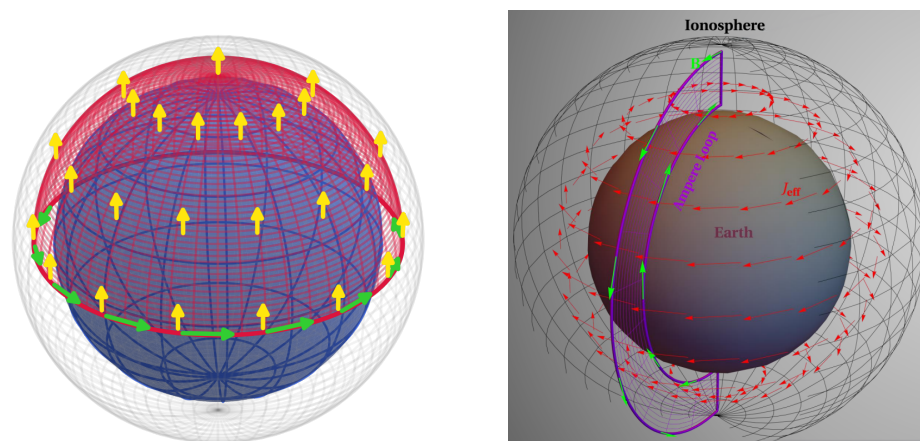


Figure 1. Sketches of idealized spherical Earth–ionosphere cavity and Ampère's law arguments. **Left** (from Ref. [52]): The Earth and ionosphere are represented by the blue and gray spheres, respectively. In red, we show a Gaussian surface covering the northern hemisphere. The yellow arrows represent the DPDM effective current $J_{\text{eff},A'}$, while the green arrows represent the magnetic-field response $B_{A'}$ along the equator. The Ampère's law argument presented in Section 2.2 implies $B_{A'} \sim J_{\text{eff},A'} \cdot R$. **Right** (from Ref. [66]): Similar argument to left panel, but for mDM current $J_{\text{eff},\phi}$ sourced by the final term of Equation (33). Now the Ampèrian loop is shown in purple and $J_{\text{eff},\phi}$ in red. This argument demonstrates that an effective current of this form generates a suppressed response $B_{\phi} \sim J_{\text{eff},\phi} \cdot h$. The dominant response instead comes from the other terms in Equation (33).

Now let us state the full results for each UBDM model. In the DPDM case, solving Equation (24) in the spherical Earth–ionosphere cavity yields the result

$$B_{A'}(\Omega, t) = \sqrt{\frac{4\pi}{3}} \cdot \frac{m_{A'} R}{2 - (m_{A'} R)^2} \cdot \operatorname{Re} \left[\sum_{m=-1}^1 \varepsilon m_{A'} \tilde{A}'_m \cdot \Phi_{1m}(\Omega) \cdot e^{-i(m_{A'} - 2\pi f_d m)t} \right]. \quad (26)$$

Here, $\Phi_{\ell m}$ are VSH (see Appendix A), and $\Omega = (\theta, \varphi)$ denotes geographic coordinates on the Earth's surface.⁸ In the NR limit, only the $\ell = 1$ VSH are relevant for the DPDM signal. The complex amplitudes $\tilde{A}'_m$ are related to the real Cartesian components in Equation (13) by

$$A'_{x,0} e^{-i\alpha_{x,0}} = -\frac{1}{\sqrt{2}} (\tilde{A}'_+ - \tilde{A}'_-), \quad (27)$$

$$A'_{y,0} e^{-i\alpha_{y,0}} = -\frac{i}{\sqrt{2}} (\tilde{A}'_+ + \tilde{A}'_-), \quad (28)$$

$$A'_{z,0} e^{-i\alpha_{z,0}} = \tilde{A}'_0. \quad (29)$$

The presence of $f_d = 1/(\text{sidereal day})$ in the time dependence of Equation (26) arises due to the rotation of Earth. The dark-photon amplitudes $\tilde{A}'_m$ are defined in the inertial celestial frame, while the signal in Equation (26) is measured in the rotating Earth-fixed frame. Translating between these frames induces a polarization-dependent shift in the frequency of the signal. (In frequency space, this will introduce sidebands at $f = f_{A'} \pm f_d$ to the signal).

In the case of axion DM, the effective current requires a background magnetic field B_0 . For the Earth transducer effect, the Earth's DC magnetic field plays the role of B_0 . We model the geomagnetic field using the IGRF-13 model [78], which provides coefficients $g_{\ell m}$ and $h_{\ell m}$ for a multipole expansion

$$B_0 = \sum_{\ell, m} C_{\ell m} \left(\frac{R}{r} \right)^{\ell+2} [(\ell+1)Y_{\ell m} - \Psi_{\ell m}], \quad (30)$$

where

$$C_{\ell m} = (-1)^m \sqrt{\frac{4\pi(2 - \delta_0^m)}{2\ell + 1}} \frac{g_{\ell m} - ih_{\ell m}}{2}. \quad (31)$$

The expansion in Equation (30) is dominated by the C_{10} term, but subsequent terms lead to a correction of $\sim 10\%$. If we plug Equation (30) into Equation (18), then the solution to Equation (24) in the spherical Earth–ionosphere cavity is

$$B_a(\Omega, t) = \sum_{\ell, m} \frac{(\ell+1)m_a R}{\ell(\ell+1) - (m_a R)^2} \cdot \operatorname{Im} \left[g_{a\gamma} \tilde{a}_0 C_{\ell m} \cdot \Phi_{\ell m}(\Omega) \cdot e^{-im_a t} \right], \quad (32)$$

where $\tilde{a}_0 = a_0 e^{-i\alpha_0}$ is the complex amplitude of the axion field. Note that while the DPDM signal includes only dipole modes, the axion signal in principle receives contributions from higher modes, due to the geomagnetic field's deviation from an exact dipole. Moreover, the axion signal does not exhibit the sidebands seen in the DPDM case. This is because the direction of $J_{\text{eff},a}$ is determined by the geomagnetic field, which co-rotates with the Earth, whereas the direction of $J_{\text{eff},A'}$ is determined by the DPDM polarization, which does not.

Finally, we address the case of mDM. As with axion DM, the presence of the geomagnetic field is crucial for the Earth transducer effect. In the mDM case, however, $J_{\text{eff},\phi}$ is

proportional to the vector potential A_0 of the Earth's field [as in Equation (23)]. In temporal Coulomb gauge, the vector potential corresponding to Equation (30) takes the form

$$A_0 = \sum_{\ell,m} \frac{C_{\ell m} R^{\ell+2}}{\ell r^{\ell+1}} \left(\frac{\kappa_{\ell m} R}{r} [(\ell+1) \mathbf{Y}_{\ell m} - \mathbf{\Psi}_{\ell m}] - \mathbf{\Phi}_{\ell m} \right). \quad (33)$$

The final term in Equation (33) is fixed by observations of the geomagnetic field. The coefficients κ_{ℓ} , however, are not, as they lead to no observable magnetic field. Instead, these coefficients are determined by the internal currents in the Earth that source the geomagnetic field, and so they are subject to modeling uncertainties of the Earth's interior. Ref. [66] estimated the range of the dominant dipolar coefficient to be $0.5 \lesssim \kappa_{10} \lesssim 2.3$. It is clear from the form of Equation (33) that the terms proportional to $\kappa_{\ell m}$ will give a similar magnetic-field response as Equation (32). In particular, it will depend only on R , not on h . On the other hand, the magnetic-field response generated by the final term in Equation (33) will be suppressed by h . This can be seen from an Ampère's law argument similar to the one for the DPDM case, shown in the right panel of Figure 1, using a different Ampèrian loop. The dominant contribution to B_{ϕ} , therefore, comes from the model-dependent portion of Equation (33),

$$B_{\phi}(\Omega, t) = \sum_{\ell,m} \frac{(\ell+1)R^2}{\ell^2(\ell+1) - \ell(m_{\phi}R)^2} \cdot \text{Re} \left[e_m^2 \tilde{\phi}_0^2 \kappa_{\ell m} C_{\ell m} \cdot \mathbf{\Phi}_{\ell m}(\Omega) \cdot e^{-2im_{\phi}t} \right], \quad (34)$$

where $\tilde{\phi}_0 = \phi_0 e^{-i\alpha_0}$ is the complex mDM amplitude. As expected, B_{ϕ} oscillates at frequency $2m_{\phi}$ (and does not exhibit sidebands from daily modulation).

Let us comment on some of the important properties of the Earth transducer signal. As we have emphasized, it is proportional to the large radius of the Earth, R , as opposed to the small height of the atmosphere, h . It is also spatially coherent, that is, Equations (26), (32) and (34) all have a prescribed dependence on Ω , so that at a fixed time t , the signals at different locations on the Earth are correlated. This is in contrast to noise, which would be uncorrelated between different locations. Combining data from many locations therefore allows us to better discriminate our signal from noise. In this way, we can leverage datasets with many simultaneous measurements across the Earth. This signal is also temporally coherent, that is, it has an exact oscillatory dependence on t . This allows us to also leverage datasets that span a long period of time. As mentioned below Equation (3), this coherence remains valid for $t \lesssim T_{\text{coh}}$, after which the amplitude and phase of the signal will change. Note that for the axion and mDM signals, the overall amplitude and phase will change from one coherence time to the next, but the spatial dependence will not.⁹ By contrast, the spatial dependence of the DPDM signal may change, as the polarization of the DPDM can vary from coherence time to coherence time.

Finally, we discuss the robustness of the results that we have derived. Equations (26), (32) and (34) were derived for an idealized spherical model of the Earth and ionosphere, even though, as discussed at the beginning of this subsection, the Earth's real atmospheric profile can be much more complicated. Nevertheless, these results can be made robust to environmental details. Let us write the full magnetic-field response as

$$B_{\text{DM}} = B_{\text{sph}} + B_{\text{env}}, \quad (35)$$

where B_{sph} is the result derived for a spherical boundary, while B_{env} is the additional contribution required to satisfy the true boundary conditions of the complicated near-Earth environment. As argued earlier in Section 2.1, as long as $\lambda_{\text{DM}} \gg R$, the electric field in the Ampère–Maxwell law can be neglected and the resulting signal B_{DM} satisfies Equation (24).

This implies that $\nabla \times \mathbf{B}_{\text{env}} = 0$. In terms of VSH, it is clear from Equations (A12)–(A14) that $\mathbf{B}_{\text{env}}$ must be composed only of $\mathbf{Y}_{\ell m}$ and $\mathbf{\Psi}_{\ell m}$ contributions. In particular, the $\Phi_{\ell m}$ contribution will come solely from Equation (26), (32), or (34), depending on the UBDM model. Because VSH form an orthonormal basis for vectorial functions on the sphere, the $\Phi_{\ell m}$ contribution can be separated from the $\mathbf{Y}_{\ell m}$ and $\mathbf{\Psi}_{\ell m}$ contributions through an appropriate projection. In this way, we can isolate a part of the signal that is robust to atmospheric modeling. The important properties of the Earth transducer signal, including how they differ between UBDM models, are summarized in Table 1.

Table 1. Summary of important properties of Earth transducer signal and how they vary for each UBDM model.

Property	DPDM	Axion DM	mDM
Scaling of $\mathbf{B}_{\text{DM}}$ with length scale	Scales with R	Scales with R	$\Phi_{\ell m}$ parts scale with R $\mathbf{Y}_{\ell m}, \mathbf{\Psi}_{\ell m}$ parts scale with h
Spatial dependence	Only Φ_{1m} (in NR limit) Fixed by DPDM polarization	All $\Phi_{\ell m}$ Fixed by multipoles of $\mathbf{B}_0$	Dominantly $\Phi_{\ell m}$ Fixed by internal currents
Temporal dependence	Oscillates at $f = f_{A'}, f_{A'} \pm f_d$ Phase, amplitude, and orientation change every T_{coh}	Oscillates at $f = f_a$ Phase and amplitude change every T_{coh}	Oscillates at $f = 2f_\phi$ Phase and amplitude change every T_{coh}
Robustness to modeling	Φ_{1m} parts robust	$\Phi_{\ell m}$ parts robust	$\Phi_{\ell m}$ parts robust to atmosphere but depend on interior

2.3. Higher Masses

Let us now discuss techniques to extend the Earth transducer effect to higher DM masses. In Sections 2.1 and 2.2, we focused on the low-mass regime where $\lambda_{\text{DM}} \gg R$. This was an important condition to neglect the electric field contribution to the Ampère–Maxwell law in Equation (10) and ultimately allowed us to argue that the low-mass signal was robust. At masses $m_{\text{DM}} \gtrsim 3 \times 10^{-14}$ eV, this argument no longer holds, and environmental effects can become important. In fact, this is known to be the case. To see how, note that the axion DM signal in Equation (32) exhibits resonances at masses $m_{\text{DM}} = \sqrt{\ell(\ell+1)}/R$. These are the so-called “Schumann resonances” of the Earth–ionosphere cavity [79–81]. The idealized spherical model that we employed in Section 2.2 predicts them at frequencies $f = 10.6$ Hz, 18.3 Hz, 25.9 Hz, ... Measurements of the Schumann resonances, however, find typical values of 7.4–8.0 Hz, 13.7–14.6 Hz, and 19.9–21.1 Hz (depending on time of day and season) for the first three resonances, with widths as small as 1.5–2.0 Hz [81]. The spherical model, therefore, miscalculates the Schumann resonances by more than their widths. As the Earth transducer effect formally diverges at these resonances, the naive spherical model could give predictions for $m_{\text{DM}} \gtrsim 3 \times 10^{-14}$ eV that are off by orders of magnitude! It is thus clear that any computation of the global Earth transducer effect must account for atmospheric effects.

There are two approaches that have been proposed to address the high-mass regime. Ref. [67] proposed to perform a local measurement of the DM effective current by measuring $\nabla \times \mathbf{B}_{\text{DM}}$ rather than $\mathbf{B}_{\text{DM}}$. Meanwhile, Refs. [59,68] included atmospheric effects in the modeling of the global Earth transducer effect. Let us begin with the first approach. In the regime $\lambda_{\text{DM}} \ll R$, the electric-field response is no longer negligible compared to the magnetic-field response, and so the $\partial_t E$ term in Equation (10) cannot simply be dropped. However, because the ground is a good conductor, it follows that near the surface of the Earth, the component parallel to the ground, $E_{\parallel}$, must still vanish! More precisely, it is argued in Ref. [67] that $E_{\parallel}$ is suppressed everywhere inside the Earth–ionosphere cavity as

long as $\lambda_{\text{DM}} \gg h$. Then, for DM masses $3 \times 10^{-14} \text{ eV} \lesssim m_{\text{DM}} \lesssim 2 \times 10^{-12} \text{ eV}$, we can take the parallel component of Equation (10) to find

$$(\nabla \times \mathbf{B})_{\parallel} = J_{\text{eff},\parallel}. \quad (36)$$

Equation (36) gives us a reliable prediction for $\nabla \times \mathbf{B}_{\text{DM}}$ in the directions parallel to the ground. Importantly, this effect is insensitive to atmospheric effects, as it does not rely on solving for a response in the presence of some global boundary conditions; rather, it is a direct probe of the local DM effective current. Note that because SM currents in the lower atmosphere are suppressed, then any environmental magnetic-field noise should satisfy $\nabla \times \mathbf{B}_{\text{noise}} = 0$. Measuring $\nabla \times \mathbf{B}$, therefore, acts as a natural background rejection scheme for environmental noise.

In order to measure $\nabla \times \mathbf{B}_{\text{DM}}$ in a horizontal direction, we require a differential measurement along a vertical gradient. This can be achieved, for instance, near a hill or mountain. Figure 2 shows an illustration of how such a measurement can be performed. Three magnetic-field measurements are taken at positions $\mathbf{r}_0, \mathbf{r}_1, \mathbf{r}_2$. The plane containing these positions must be parallel to the vertical direction. At each position, we measure the component of the magnetic field parallel to the baseline defined by the other two positions. These measurements can then be combined as

$$\Delta \equiv \mathbf{B}(\mathbf{r}_0) \cdot (\mathbf{r}_2 - \mathbf{r}_1) + \mathbf{B}(\mathbf{r}_1) \cdot (\mathbf{r}_0 - \mathbf{r}_2) + \mathbf{B}(\mathbf{r}_2) \cdot (\mathbf{r}_1 - \mathbf{r}_0) \quad (37)$$

$$\approx [(\mathbf{r}_2 - \mathbf{r}_0) \times (\mathbf{r}_1 - \mathbf{r}_0)] \cdot (\nabla \times \mathbf{B}), \quad (38)$$

where the approximation holds within the limit $\mathbf{r}_1, \mathbf{r}_2 \rightarrow \mathbf{r}_0$. This procedure gives a direct measurement of $\nabla \times \mathbf{B}$ in the direction normal to the plane containing the three positions. Note that Δ is proportional to the area of the triangle formed by $\mathbf{r}_0, \mathbf{r}_1, \mathbf{r}_2$, and so it is beneficial to maximize this area. This procedure can be understood as a discretized line integral around the dotted line in Figure 2.

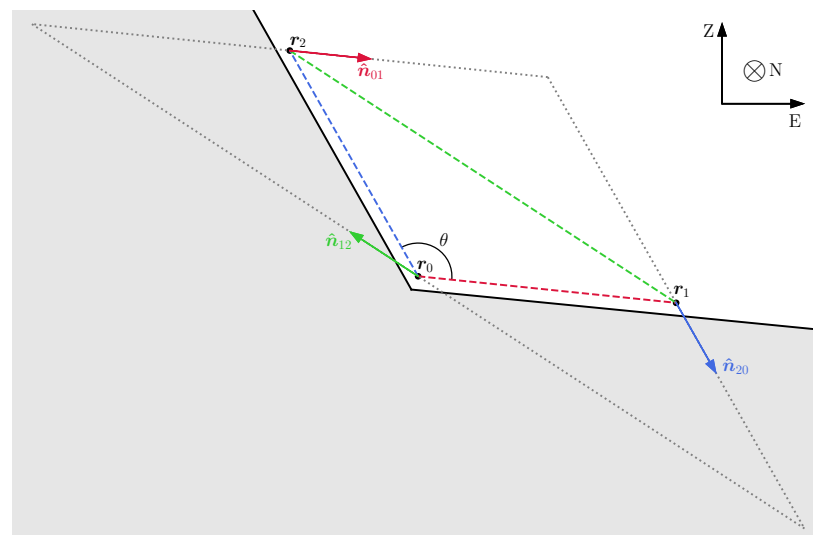


Figure 2. (from Ref. [67]). Illustration of curl measurement scheme described in Section 2.3. Three single-axis magnetic-field measurements are made at $\mathbf{r}_0, \mathbf{r}_1, \mathbf{r}_2$. The direction of each measurement (represented by the unit vectors $\hat{n}_{12}, \hat{n}_{20}, \hat{n}_{01}$) should be parallel to the baseline between the other two locations (shown in matching colored dashed lines). These measurements are combined as in Equation (37), which can be understood as a discretized line integral around the dotted gray line.

Refs. [59,61,68] took an alternative approach to extend the Earth transducer effect to higher masses, which was to incorporate the measured atmospheric conductivity profile into the calculation of the magnetic-field response. This is achieved by modifying Gauss's law and the Ampère–Maxwell law, Equations (7) and (10), as

$$\nabla \cdot (n^2 \mathbf{E}) = \rho_{\text{eff}}, \quad (39)$$

$$\nabla \times \mathbf{B} - \partial_t(n^2 \mathbf{E}) = \mathbf{J}_{\text{eff}}, \quad (40)$$

where $n(x) = \sqrt{1 + i\sigma(x)/m_a}$ is the (spatially varying) index of refraction. Refs. [59,61,68] assumed a spherically symmetric atmospheric conductivity profile $\sigma(r)$. While the upper layers of the Earth's atmosphere exhibit significant anisotropy, for high masses $m_{\text{DM}} \gtrsim 3 \times 10^{-4}$ eV, the magnetic-field response is suppressed above ~ 100 km. Therefore, the calculation is most sensitive to the properties of the lower atmosphere, which is relatively isotropic. Ref. [68] verified this by computing the response for multiple atmospheric models, including both daytime and nighttime models, and found negligible differences in the response. Their resulting B_{DM} exhibits resonant features that agree with the observed central frequencies and widths of the Schumann resonances.

3. Experimental Searches

We have seen that the Earth transducer effect generates a global AC magnetic-field signal at the Earth's surface. The simplest way to detect this effect is with data from unshielded magnetometers. In order to leverage the spatial and temporal coherence discussed in Section 2.2, it is beneficial for this dataset to include measurements from many magnetometers over a long period of time. In this section, we review several searches for the Earth transducer effect in such datasets. A particularly powerful strategy has been to repurpose existing datasets to search for UBDM through this effect. Refs. [53–55] searched a publicly available dataset maintained by the SuperMAG collaboration [56–58], which collates magnetic-field data from a global array of unshielded magnetometers. Likewise, Refs. [59–61] searched magnetic-field data from the Eskdalemuir observatory [62] for the Earth transducer effect. At frequencies $f_{\text{DM}} \gtrsim \text{Hz}$ ($m_{\text{DM}} \gtrsim 4 \times 10^{-15}$ eV), much of the ambient magnetic-field noise is anthropogenic [82], and so noise levels can be reduced significantly by performing measurements in a radio quiet location. For this reason, the SNIPE Hunt [63,64] and GPEX [65] collaborations have performed dedicated high-frequency searches for the Earth transducer effect in remote locations. The properties of these searches are compared in Table 2. All of these searches have placed constraints on UBDM parameter space, as summarized in Figures 3 and 4. In many regions of parameter space, unshielded magnetometers are the most sensitive experimental probes for UBDM.

Table 2. Comparison of various experimental searches for the Earth transducer effect.

Property	SuperMAG	Eskdalemuir	SNIPE Hunt	GPEX
Type	Existing dataset	Existing dataset	Dedicated search	Dedicated search
Total number of sites	~ 500	1	5	1
Total duration	~ 50 yrs	~ 10 yrs	~ 80 h	~ 3 h
Frequency range [Hz]	6×10^{-4} –1	0.02–40	0.5–5	0.1–5
Mass range [eV]	2×10^{-18} – 4×10^{-15}	7×10^{-17} – 2×10^{-13}	2×10^{-15} – 2×10^{-14}	4×10^{-16} – 2×10^{-14}

3.1. Existing Datasets

Low-frequency ground-based magnetometers are frequently employed to study geophysical processes and solar activity. Accordingly, an abundance of data from unshielded magnetometers have already been recorded. One prominent example is the publicly

available dataset maintained by the SuperMAG collaboration [56–58]. This collaboration aggregates and processes data from an array of magnetometer stations dispersed around the globe. They maintain two datasets: a “low-fidelity” (1 min temporal resolution) dataset, which combines data from ~ 500 stations over a span of more than 50 years, and a “high-fidelity” (1 s temporal resolution) dataset, which combines data from ~ 200 stations over more than 20 years. Refs. [53,54] performed searches of the SuperMAG low-fidelity dataset for the Earth transducer signal (for DPDM and axion DM, respectively) at Compton frequencies $f_{\text{DM}} < 1/(1 \text{ min})$, corresponding to masses $m_{\text{DM}} \lesssim 7 \times 10^{-17} \text{ eV}$, and Ref. [55] performed a search of the high-fidelity dataset (for both DPDM and axion DM) for $f_{\text{DM}} < 1 \text{ Hz}$, corresponding to $m_{\text{DM}} \lesssim 4 \times 10^{-15} \text{ eV}$.

Here, we give a cursory overview of the analytical framework utilized in these searches. The ideal search protocol would be to Fourier-transform the magnetic-field time series data from each station, combine them to project onto the VSH combination in Equation (26) or (32), and then search for peaks in the resulting spectrum. The SuperMAG datasets are, however, complicated by the fact that most stations do not actively report data for the entire duration of the dataset. These periods of inactivity make a naive Fourier transform less straightforward. Moreover, the noise levels vary significantly from station to station, such that combining data from different stations with comparable weights may be suboptimal. Instead, Refs. [53–55] employed the following strategy. First, each station’s measurements are weighted by a VSH motivated by the Earth transducer signal, e.g., schematically,

$$X_i^{(n)}(t_j) \sim \Phi_{\ell m}^\alpha(\Omega_i) \cdot B_i^\alpha(t_j), \quad (41)$$

where $B_i^\alpha(t_j)$ represents the α component of the magnetic-field data from station i at time t_j , $\Phi_{\ell m}^\alpha$ represents the α component of $\Phi_{\ell m}$, and Ω_i represents the location of station i . The n quantities $X_i^{(n)}$ weight the data from station i with different VSH components and coefficients. The quantities from each station are then combined into n time series, weighted by the noise levels in each station, e.g., schematically,

$$X^{(n)}(t_j) \sim \frac{1}{W(t_j)} \sum_i w_i X_i^{(n)}(t_j), \quad (42)$$

where w_i are chosen station weights that are inversely proportional to the station’s DC noise level, and the sum runs over all i which report a measurement at time t_j . The total weights $W(t_j) = \sum_i w_i$ can depend on t_j because the number of active stations vary with time. In this way, the framework constructs a small number n of time series with nonzero values at all t . The combination procedure weights station data according to both their noise level and the prediction in Equation (26) or (32).¹⁰ These n time series then become the primary object of study for the search procedure, namely, they are Fourier transformed and compared to noise in a Bayesian framework to determine the significance of a potential signal. This framework requires both a prediction for how the signal would appear in $X^{(n)}$ and an estimate for the noise in $X^{(n)}$. To determine the signal prediction, the same combination procedure that is performed on the data $B_i^\alpha(t_j)$ is applied to the predicted magnetic-field signal in Equation (26) or (32). The noise, on the other hand, is determined by a data-driven analysis directly of the time series $X^{(n)}$.

The analyses in Refs. [53–55] each found a number of signal candidates that exhibited high statistical significance. Many such candidates were the results of transient anomalies in single stations (either from the instrument itself or from post-processing of the data). By contrast, a true UBDM signal would be present at all stations at all times. Refs. [53–55] applied secondary checks, such as geographic uniformity and temporal persistence, to

each of these candidates and ultimately concluded that none constituted robust evidence of UDBM. As a result, they placed constraints on DPDM and axion DM parameter space. The green and blue curves in Figure 3 show the smoothed constraints from the low- and high-fidelity searches, respectively.

Another series of works has analyzed data from the Eskdalemuir observatory in Scotland [62]. Refs. [59,60] searched this dataset for an axion DM signal, and Ref. [61] considered a DPDM signal. Importantly, both of these searches included atmospheric modeling in their signal computation which allowed a robust prediction of the signal up to $f_{\text{DM}} \sim 30$ Hz. In the absence of any robust signal candidates, they also set constraints on DPDM and axion DM parameter space, shown in purple in Figure 3.¹¹ Notice that the sensitivity peaks near the first Schumann resonance $f_{\text{DM}} \sim 7$ Hz, as this is where the Earth transducer signal is maximized.

3.2. Dedicated Searches

At higher frequencies $f_{\text{DM}} \gtrsim$ Hz, ambient magnetic field noise is dominated by anthropogenic sources [82]. This implies that the choice of location is increasingly important for high-frequency measurements of the Earth transducer effect. The SNIPE Hunt collaboration was formed to take dedicated measurements that could fully leverage this observation. SNIPE Hunt has undertaken campaigns in 2022 (SNIPE Run I) [63] and 2025 (SNIPE Run II) [64,69]. The goal of each campaign was to acquire time-synchronized measurements of magnetic fields in regions with minimal human-made magnetic interference, such as arises from power lines and moving vehicles. Table 3 contrasts the relevant characteristics of these two runs. In each case, magnetic field measurements were simultaneously performed at a number of geographically separated sites.¹² The acquired data were assigned timestamps and GPS measurements and written to file for subsequent analysis. The sensors and data acquisition system were powered by rechargeable lithium-ion battery packs, which were periodically replaced during the science runs as they discharged.

Table 3. Characteristics of SNIPE Runs I and II.

Property	SNIPE Run I	SNIPE Run II
Number of sites	3	5
Total (analyzed) duration	~ 60 (~ 25) h	~ 80 (~ 15) h
Sensor type	Giant magnetoresistance	Induction coil
Sensor bandwidth	DC to 80 Hz	< 0.1 Hz to 1.0 kHz
Noise floor	~ 300 pT/ $\sqrt{\text{Hz}}$	~ 1 pT/ $\sqrt{\text{Hz}}$
Detection axes	3 (N–S, E–W, vertical)	2 (N–S, E–W)
Data acquisition system	Laptops with serial connection to Twinleaf VMR	Raspberry Pi + Pi-Plates ADC board

In SNIPE Run I, each site deployed Twinleaf Vector Magnetoresistive (VMR) magnetometers [83]. The noise floor of the measurement was limited by the instrumental noise, which, in this case, was independent of frequency. The analysis of the SNIPE Run I data followed a Bayesian framework similar to the one employed in the SuperMAG searches. Data were modeled as a signal, given by Equation (26) or (32), on top of Gaussian noise. In the Run I analysis, a frequency-independent data-driven estimate was used for the noise.

In SNIPE Run II, two of the five total sites deployed induction-coil magnetometers supplied by Lemi Sensors LLC [84], while the other three sites developed their own sensors and readout electronics in their laboratories. At the heart of each sensor was (10^4) turns of copper wound around a mu-metal core. A flux-feedback approach was used to flatten the response across the desired bandwidth. The noise floor for the sensors used in SNIPE Run

II was set by a combination of instrumental and environmental noise. Consequently, this noise displayed non-negligible frequency dependence and non-stationarity. In the analysis, a frequency- and time-dependent noise model was applied based on a running average of power in different frequency bins.

Both SNIPE Runs were able to set constraints on DPDM and axion DM in the frequency range $0.5 \text{ Hz} \lesssim f_{\text{DM}} \lesssim 5 \text{ Hz}$, with the upper bound chosen to avoid atmospheric dependence from the Schumann resonances. These constraints are shown in red (Run I) and orange (Run II) in Figure 3. A similar search was performed by the GPEX collaboration [65] in the desert of XiaoDushan in Gansu Province, China. This collaboration collected data for one hour using an atomic magnetometer manufactured by QUSPIN [85]. The results from their search are shown in yellow in Figure 3. In the future, the SNIPE Hunt collaboration intends to implement the curl technique described in Section 2.3 and illustrated in Figure 2. The dashed pink lines in Figure 3 show the projections for SNIPE Hunt's implementation of the curl measurement, based on Ref. [67].

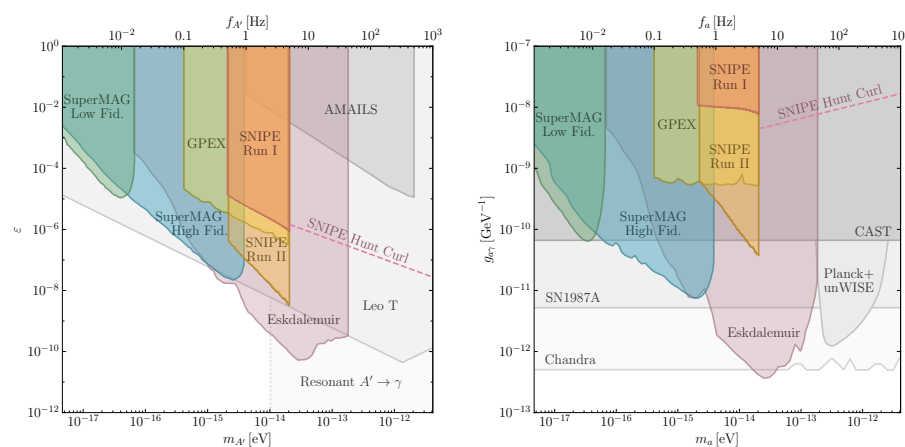


Figure 3. Summary of constraints on DPDM (left) and axion DM (right) from searches for the Earth transducer effect. In cool colors, we show smoothed constraints from searches in existing datasets: the SuperMAG low- and high-fidelity datasets [53–58] and data from the Eskdalemuir observatory [60–62]. In warm colors, we show smoothed constraints from dedicated searches by various collaborations: two data-collecting runs by the SNIPE Hunt collaboration [63,64], and one run by the GPEX collaboration [65]. The dashed pink lines show the sensitivities that SNIPE Hunt could achieve with the curl measurement discussed in Section 2.3 [67]. In gray, we show various existing constraints, with darker shades indicating direct constraints and lighter shades indicating astrophysical/cosmological constraints (see text for details).

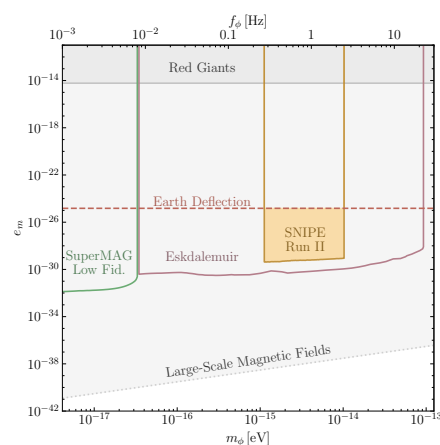


Figure 4. Sensitivity of the Earth transducer effect to mDM. We recast the SuperMAG low-fidelity [54] and Eskdalemuir [60] axion DM constraints from Figure 3 to mDM parameter space via the mapping

in Equation (43), with $\kappa_{\ell m} = \kappa_{10} = 0.5$. These sensitivity curves are not full reanalyses of the corresponding datasets and depend on modeling of the Earth's interior. Nevertheless, they can be interpreted, up to a factor of (1), as constraints on mDM parameter space. In orange, we show a dedicated analysis of the SNIPE Run II dataset for an mDM signal [69]. We shade this region to indicate a more rigorous exclusion. Above the dashed red line, the Earth's magnetic field can deflect mDM significantly (see note 7 above), such that the treatment in this review no longer holds. In gray, we also show existing astrophysical/cosmological constraints (see text for details). Note that the constraint from Large-Scale Magnetic Fields is significantly relaxed if mDM is a subcomponent of DM [86,87].

4. Discussion and Conclusions

In this review, we have described how the Earth can act as a transducer for EM-coupled UBDM, allowing its potential detection by unshielded magnetometers. The effect of UBDM can be parametrized as an effective current J_{eff} , as in Equations (14), (18) and (23) for various models. These currents produce AC magnetic fields via Ampère's law in Equation (24). The resulting magnetic-field signal is enhanced by the large radius of the Earth, R , and it is spatially and temporally coherent. We also showed that, for $m_{\text{DM}} \lesssim 1/R$, the signal prediction can be made robust to environmental details via a VSH projection. We discussed various methods to extend this effect to higher frequencies, such as differential measurements of $(\nabla \times \mathbf{B})_{\parallel}$ or inclusion of atmospheric modeling.

Figure 3 summarizes the results of various searches for the Earth transducer effect. Regions shaded in cool colors represent smoothed constraints from searches in existing datasets (SuperMAG Low Fid., SuperMAG High Fid., and Eskdalemuir), while regions shaded in warm colors represent dedicated searches for the Earth transducer effect (SNIPE Run I, SNIPE Run II, and GPEX). The dashed pink lines show the potential sensitivity of the curl measurement described in Section 2.3. We compare these to existing constraints in various shades of gray, with darker shades representing direct constraints (AMAILS [47] and CAST [35]) and lighter shades representing astrophysical/cosmological constraints (Leo T [88], Resonant $A' \rightarrow \gamma$ [89],¹³ SN1987A [91], Planck + unWISE [92], and Chandra [93]). The Earth transducer effect is the most powerful direct probe of DPDM with masses $m_{A'} \lesssim 10^{-13}$ eV. The constraint from the Eskdalemuir search even surpasses all astrophysical/cosmological constraints on DPDM in some mass ranges. For axion DM, several searches for the Earth transducer effect have surpassed the leading direct constraint from CAST, while the Eskdalemuir search has probed previously unconstrained axion DM parameter space.

In Ref. [66], it was pointed out that all of these searches are also sensitive to mDM. A comparison of Equations (32) and (34) shows that the axion and mDM signals are approximately related by the mapping

$$e_m^2 \sim \frac{\ell m_\phi^2}{\kappa_{\ell m} \sqrt{2\rho_{\text{DM}} R}} \cdot g_{a\gamma} \quad (43)$$

(with $m_\phi = m_a/2$). A precise search for the mDM Earth transducer effect, in principle, differs from a search for the axion DM effect because the linear combinations of VSH in Equations (32) and (34) may not be the same. However, because both expansions are dominated by the $\ell = 1, m = 0$ term, it is sufficient to understand the potential sensitivity to mDM by recasting the axion constraints in Figure 3 using Equation (43) with a fiducial value for κ_{10} . In Figure 4, we show recast constraints for a couple of these searches using $\kappa_{10} = 0.5$. While these are not rigorous reanalyses of the datasets, the constraints should nevertheless be accurate to (1). The SNIPE Hunt collaboration has performed a dedicated analysis of their SNIPE Run II dataset for an mDM signal [69] (considering only the $\ell = 1, m = 0$

term). We show this as a shaded orange region in Figure 4 to demonstrate a more rigorous exclusion. The dashed red line in Figure 4 shows where the Earth’s magnetic field can significantly alter the mDM distribution (see note 7 above). The treatment in this review remains valid below the line. In gray, we also show existing astrophysical constraints from Red Giants [94] and Large-Scale Magnetic Fields [86,87]. Note that the latter constraint is significantly relaxed if the mDM is a subcomponent of DM. In this case, the Earth transducer effect (which is relaxed to a lesser degree) becomes the strongest probe of mDM.

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Conflicts of Interest: The authors declare no conflicts of interest.

Appendix A. Vector Spherical Harmonics

This appendix, which defines the VSH conventions used in this work, is reproduced from Ref. [52] with minor modifications for the convenience of the reader.

The vector spherical harmonics are defined in terms of the scalar spherical harmonic $Y_{\ell m}$ by the relations

$$Y_{\ell m} = Y_{\ell m} \hat{r}, \quad \Psi_{\ell m} = r \nabla Y_{\ell m}, \quad \Phi_{\ell m} = \mathbf{r} \times \nabla Y_{\ell m}, \quad (\text{A1})$$

where $\hat{r}$ is the unit vector in the direction of $\mathbf{r}$. Thus, $Y_{\ell m}$ points radially, while $\Psi_{\ell m}$ and $\Phi_{\ell m}$ point tangentially. Some of their relevant properties (and our phase conventions) are

$$Y_{\ell, -m} = (-1)^m Y_{\ell m}^*, \quad (\text{A2})$$

$$\Psi_{\ell, -m} = (-1)^m \Psi_{\ell m}^*, \quad (\text{A3})$$

$$\Phi_{\ell, -m} = (-1)^m \Phi_{\ell m}^*, \quad (\text{A4})$$

$$Y_{\ell m} \cdot \Psi_{\ell m} = Y_{\ell m} \cdot \Phi_{\ell m} = \Psi_{\ell m} \cdot \Phi_{\ell m} = 0, \quad (\text{A5})$$

$$\int d\Omega \mathbf{Y}_{\ell m} \cdot \mathbf{Y}_{\ell' m'}^* = \delta_{\ell\ell'} \delta_{mm'}, \quad (\text{A6})$$

$$\begin{aligned} \int d\Omega \mathbf{\Psi}_{\ell m} \cdot \mathbf{\Psi}_{\ell' m'}^* &= \int d\Omega \mathbf{\Phi}_{\ell m} \cdot \mathbf{\Phi}_{\ell' m'}^* \\ &= \ell(\ell+1) \delta_{\ell\ell'} \delta_{mm'}, \end{aligned} \quad (\text{A7})$$

$$\begin{aligned} \int d\Omega \mathbf{Y}_{\ell m} \cdot \mathbf{\Psi}_{\ell' m'}^* &= \int d\Omega \mathbf{Y}_{\ell m} \cdot \mathbf{\Phi}_{\ell' m'}^* \\ &= \int d\Omega \mathbf{\Psi}_{\ell m} \cdot \mathbf{\Phi}_{\ell' m'}^* = 0. \end{aligned} \quad (\text{A8})$$

For any radially dependent function $f(r)$, the divergences and curls of the VSH are given by

$$\nabla \cdot (f \mathbf{Y}_{\ell m}) = \left(\frac{df}{dr} + \frac{2f}{r} \right) \mathbf{Y}_{\ell m}, \quad (\text{A9})$$

$$\nabla \cdot (f \mathbf{\Psi}_{\ell m}) = -\ell(\ell+1) \frac{f}{r} \mathbf{Y}_{\ell m}, \quad (\text{A10})$$

$$\nabla \cdot (f \mathbf{\Phi}_{\ell m}) = 0, \quad (\text{A11})$$

$$\nabla \times (f \mathbf{Y}_{\ell m}) = -\frac{f}{r} \mathbf{\Phi}_{\ell m}, \quad (\text{A12})$$

$$\nabla \times (f \mathbf{\Psi}_{\ell m}) = \left(\frac{df}{dr} + \frac{f}{r} \right) \mathbf{\Phi}_{\ell m}, \quad (\text{A13})$$

$$\nabla \times (f \mathbf{\Phi}_{\ell m}) = -\frac{\ell(\ell+1)f}{r} \mathbf{Y}_{\ell m} - \left(\frac{df}{dr} + \frac{f}{r} \right) \mathbf{\Psi}_{\ell m}, \quad (\text{A14})$$

with the Laplacians then being

$$\begin{aligned} \nabla^2 (f \mathbf{Y}_{\ell m}) &= \left(\frac{1}{r^2} \frac{d}{dr} \left(r^2 \frac{df}{dr} \right) - \frac{(\ell(\ell+1)+2)f}{r^2} \right) \mathbf{Y}_{\ell m} \\ &\quad + \frac{2f}{r^2} \mathbf{\Psi}_{\ell m}, \end{aligned} \quad (\text{A15})$$

$$\begin{aligned} \nabla^2 (f \mathbf{\Psi}_{\ell m}) &= \left(\frac{1}{r^2} \frac{d}{dr} \left(r^2 \frac{df}{dr} \right) - \frac{\ell(\ell+1)f}{r^2} \right) \mathbf{\Psi}_{\ell m} \\ &\quad + \frac{2\ell(\ell+1)f}{r^2} \mathbf{Y}_{\ell m}, \end{aligned} \quad (\text{A16})$$

$$\nabla^2 (f \mathbf{\Phi}_{\ell m}) = \left(\frac{1}{r^2} \frac{d}{dr} \left(r^2 \frac{df}{dr} \right) - \frac{\ell(\ell+1)f}{r^2} \right) \mathbf{\Phi}_{\ell m}. \quad (\text{A17})$$

The explicit expressions for the spherical harmonics that are relevant to the DPDM signal [see Equation (26)] are

$$\mathbf{\Phi}_{1,-1}(\mathbf{r}) = \sqrt{\frac{3}{8\pi}} e^{-i\varphi} (i\hat{\boldsymbol{\theta}} + \cos\theta \hat{\boldsymbol{\phi}}), \quad (\text{A18})$$

$$\mathbf{\Phi}_{10}(\mathbf{r}) = -\sqrt{\frac{3}{4\pi}} \sin\theta \hat{\boldsymbol{\phi}}, \quad (\text{A19})$$

$$\mathbf{\Phi}_{11}(\mathbf{r}) = \sqrt{\frac{3}{8\pi}} e^{i\varphi} (i\hat{\boldsymbol{\theta}} - \cos\theta \hat{\boldsymbol{\phi}}), \quad (\text{A20})$$

where $\hat{\boldsymbol{\theta}}$ and $\hat{\boldsymbol{\phi}}$ are unit vectors in the directions of increasing θ and φ (see note 8 above for coordinate conventions.)

Notes

¹ We work in natural units $\hbar = c = \varepsilon_0 = 1$.

² Throughout this review, we write x^μ for four vectors and $\mathbf{x}$ for three vectors.

- 3 For a vector field X_μ , the corresponding relation is $\langle |X_{i,0}|^2 \rangle = 2\rho_{\text{DM}}/3m_{\text{DM}}^2$ for each spatial component $i = 1, 2, 3$.
- 4 The DPDM interaction is often expressed instead as a “kinetic mixing”, $\mathcal{L}_{\text{int},A'} = \frac{\varepsilon}{2} F'_{\mu\nu} F^{\mu\nu}$. To leading order, this is equivalent to Equation (11) via a field redefinition $A_\mu \rightarrow A_\mu, A'_\mu \rightarrow A'_\mu + \varepsilon A_\mu$.
- 5 In contexts that are relevant to a particular model, we write the mass as $m_{A'}$, m_a , or m_ϕ . In contexts that apply to all UBDM models, we write m_{DM} .
- 6 In all contexts in this review, there will be no net SM charge density, $\rho_{\text{SM}} = 0$, so it is consistent to impose both of these conditions.
- 7 If e_m is sufficiently large, the mDM can be deflected by the geomagnetic field. This effect can significantly alter the background distribution of the mDM. The distribution in Equation (22) (and the resulting signal in Equation (34)) remain valid as long as $e_m^2 |B_0|^2 R^4 / 4 \ll 1$ [69].
- 8 Note that while φ coincides with the usual definition of longitude, θ is not the usual latitude. The geographic north pole (latitude $+90^\circ$) corresponds to $\theta = 0$, while the geographic south pole (latitude -90°) corresponds to $\theta = \pi$.
- 9 The spatial dependencies of the axion and mDM signals will, however, change as the geomagnetic field slowly drifts over time, which occurs at a rate of $\sim 1\%$ per decade.
- 10 As mentioned in Section 2.2, for low masses $m_{\text{DM}} \ll 10^{-16}$ eV, the complicated geometry of Earth’s atmosphere may introduce $Y_{\ell m}$ and $\Psi_{\ell m}$ contributions into Equation (26) or (32). In principle, these can be removed via a projection onto the $\Phi_{\ell m}$ VSH. However, because the combination procedure in Equation (42) includes the weights w_i , and because the stations are not uniformly distributed across the Earth, the sum in Equation (42) does not constitute an exact VSH projection. In Ref. [53], it was estimated that this effect may alter the signal prediction, as well as the resulting constraint, by a factor of (1) for the low-fidelity search.
- 11 The analyses in Refs. [60,61] discard narrow frequency ranges around integer frequencies and very narrow ranges around multiples of 0.05 Hz. As a result, their constraints do not rule out all frequencies in the regions shown in Figure 3. Rather these constraints should be understood to have many small gaps.
- 12 In contrast to the SuperMAG searches, the SNIPE Hunt runs took measurements at only a few sites, and the GPEX search had only one site. This means SNIPE Hunt and GPEX cannot reliably perform the VSH projection mentioned at the end of Section 2.2. Because these searches focus on masses $m_{\text{DM}} \gg 10^{-16}$ eV, we expect the spherical Earth=ionosphere model to remain valid for their mass range of interest. Therefore, Equations (26), (32) and (34) can be used without the VSH projection.
- 13 This constraint relies on resonant conversion of DPDM into photons during the dark ages. This mechanism was recently called into question in Ref. [90] because nonlinearities may disrupt the resonant conversion. For completeness, we still include this constraint in Figure 3, but represent it with a dotted line.

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