

Geomagnetic Constraints on Millicharged Dark Matter

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Millicharged particles are well-motivated dark matter candidates arising in many extensions of the standard model. We show that, despite their tiny coupling e_m to photons, millicharged dark matter (mDM) in the Earth's geomagnetic field can generate a quasistatic, monochromatic magnetic signal with angular frequency twice the mDM mass. Using null results from the SuperMAG and SNIPE Hunt collaborations, we constrain the effective charge of bosonic mDM in the mass range $10^{-18} - 10^{-14}$ eV. The resulting upper bounds exceed stellar cooling constraints by over 13 orders of magnitude, demonstrating the power of this method.

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Introduction—Identifying the nature of dark matter has been one of the frontiers in modern physics. Millicharged particles, i.e., quantum fields with electric charges much smaller than the one e of the electron, are natural candidates [1–5]. Millicharged dark matter (mDM) can arise from models where the visible and dark sectors are kinetically mixed [3,6–10] or from models where these particles are directly charged under the $U(1)_Y$ in the UV theory [11–13]. They are also widely predicted in string theory compactifications [11–16] and grand unification theories [17–19].

The masses of mDM can span a wide range. Depending on their masses, mDM could have been produced in the early universe either via the freeze-in mechanism [20–25] or the misalignment mechanism [26–30]. In particular, the bosonic mDM can be of a low mass (possibly even sub-eV) [31,32], which might offer a solution to the small-scale structure problems and be relevant to astrophysical

observations [33–35]. For example, mDM with masses on the order of a few tens of MeV could account for the anomaly in the 21-cm signal if a small fraction of the dark matter is composed of them [36–40].

There have been various experiments and proposals capable of searching for mDM. They include cavities filled with strong electric fields to probe Schwinger pair production [41–43], observation of the invisible decays of positronium [44], lamb shift of the hydrogen atom [45], laser polarization experiments [46–49], Cavendish experiment for testing Coulomb's law [50], timing of radio waves from pulsars and fast radio bursts [51,52], millicharged condensates [53], and direct deflection type experiments [54–56]. Meanwhile, astrophysical observations including the SN1987A [57–59] and stellar evolution [60–65] have placed very strong constraints, and even more stringent ones can be obtained from galactic large scale magnetic fields [66] and cosmological measurements [32,67–69], although the latter are highly model dependent (for comprehensive reviews, see also Refs. [70,71]).

In this Letter, we propose a new method to search for ultralight mDM by detecting their magnetic signal on Earth's surface. For sub-eV masses, the high occupation number allows the mDM field to be treated as a classical wave. When the Compton wavelength is much larger than the Earth's radius, the mDM field coherently interacts with the geomagnetic field, producing a quasistatic, monochromatic magnetic signal with angular frequency equal to twice the mDM mass, and width given by $\langle v_\phi^2 \rangle \sim 10^{-6}$ of the

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main signal frequency, where v_ϕ is the mDM velocity. It is important to emphasize that the magnetic field induced by mDM shows a characteristic inverse-square dependence on the mass in the relevant parameter regime,

$$B_\phi \sim 4e_m^2 B_0 R_e^2 \rho / m_\phi^2, \quad (1)$$

where e_m is the millicharge, m_ϕ the mDM mass, B_0 the geomagnetic field, R_e the Earth's radius, and ρ the local dark matter energy density. This strong $1/m_\phi^2$ scaling significantly enhances the sensitivity to ultralight mDM. By contrast, other dark-matter candidates induce magnetic signals with different mass dependence: dark photons, $B_{A'} \sim \epsilon m_{A'} R_e \sqrt{2\rho}$ [72,73]; axions, $B_a \sim g_{a\gamma} B_0 R_e \sqrt{2\rho}$ [74].

This renders mDM very distinct and promising in searches for ultralight dark matter.

By using null results for axion and dark photon searches from the SuperMAG [72–75] and SNIPE Hunt [76] collaborations, we estimate a world-leading bound on ultralight mDM, which surpasses the existing constraints from stellar cooling by more than 13 orders of magnitude. Consequently, it is expected that future dedicated measurements of the magnetic activity on the Earth's surface will open up new avenues for the exploration of mDM of different mass ranges and smaller values of e_m .

mDM electrodynamics—The interaction between the ultralight bosonic mDM field ϕ and the photon field A_μ is described by the following Lagrangian,

$$\mathcal{L} = D_\mu \phi (D^\mu \phi)^* - m_\phi^2 |\phi|^2 - \frac{1}{4} F_{\mu\nu} F^{\mu\nu}, \quad (2)$$

where $D_\mu = \partial_\mu + ie_m A_\mu$ is the covariant derivative and $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ is the electromagnetic strength tensor [77]. For a given external electromagnetic field $\mathcal{A}_\mu = (A_0, -\vec{A})$, we make the decomposition $A_\mu \rightarrow A_\mu + \mathcal{A}_\mu$, where now A_μ corresponds explicitly to the induced electromagnetic signal. The equations of motion are then written as

$$\partial_\mu F^{\mu\nu} + 2e_m^2 |\phi|^2 A^\nu = J_m^\nu - 2e_m^2 \mathcal{A}^\nu |\phi|^2 \equiv J_{\text{eff}}^\nu, \quad (3)$$

$$(\square + m_\phi^2)\phi = -ie_m \partial_\mu \mathcal{A}^\mu \phi - 2ie_m \partial_\mu \phi \mathcal{A}^\mu + e_m^2 \mathcal{A}_\mu \mathcal{A}^\mu \phi, \quad (4)$$

where $J_m^\nu = ie_m (\phi^* \partial^\nu \phi - \phi \partial^\nu \phi^*)$ is the millicharged current; the nonlinear terms are neglected in Eq. (4). In Eq. (3), the effective mass term $2e_m^2 |\phi|^2 A^\nu$ can, in principle, induce a photon parametric resonance [32]. However, we find that this effect is subdominant in the parameter space of our interest, and we can safely neglect it.

Gauge invariance persists in our model at all orders. Indeed, it is evident from the equations of motion, as the above mass term in Eq. (3) and the nonlinear terms in Eq. (4) are neglected. Even when the mass term is taken

into account, the gauge invariance is realized from the significant backreaction on the dark matter field once the previously neglected nonlinear terms in Eq. (4) are taken into account. For convenience, we adopt the Coulomb gauge, i.e., $\vec{\nabla} \cdot \vec{A} = 0$ and $A_0 = 0$.

Once we know the mDM field ϕ at the Earth's location, the electromagnetic signal is found by solving Eq. (3).

Due to the high occupation number of sub-eV mDM, it is typically treated as a free classical wave, $\phi = \phi_0 = (\sqrt{2\rho}/m_\phi) \cos(\vec{k}_\phi \cdot \vec{x} - m_\phi t)$, where $\vec{k}_\phi = m_\phi \vec{v}_\phi$ denotes the mDM momentum. However, as the mDM field approaches the Earth, the Earth's geomagnetic field can alter the field configuration, especially for ultralight masses. To capture this effect, we solve Eq. (4) with the boundary condition $\phi \rightarrow \phi_0$ as $|\vec{x}| \rightarrow \infty$. For the mass range relevant to our Letter, the de Broglie wavelength of mDM greatly exceeds the Earth's radius. This allows us to simplify the problem by modeling the Earth's effect as a Dirac delta potential localized at $\vec{x} = 0$, leading to the solution [80]

$$\phi(t) = \frac{1}{2} \frac{\sqrt{2\rho}}{m_\phi} (1 + i\epsilon) \left(\frac{e^{-i\omega_\phi t}}{1 + i\kappa\epsilon} + \frac{e^{i\omega_\phi t}}{1 - i\kappa\epsilon} \right), \quad (5)$$

where $\omega_\phi = \sqrt{k_\phi^2 + m_\phi^2} \simeq m_\phi$ is the mDM energy, and we define two dimensionless parameters, $\kappa \approx e_m B_0 R_e / k_\phi$ and $\epsilon \approx 8e_m B_0 R_e^2$. A small κ results in a large gyroradius, so that the trajectories of the particles are barely deflected by Earth's magnetic field. Meanwhile, a small ϵ highlights the wavelike behavior of the millicharged particles, leading to significant tunneling through the geomagnetic potential barrier that would otherwise repel them. When the interaction with the Earth's magnetic field is weak, i.e., $\epsilon, \kappa \ll 1$, Eq. (5) reduces to the free-wave form $\phi = \phi_0 \sim \sqrt{2\rho}/m_\phi$. For the parameter region where we set upper bounds on e_m , these conditions are always satisfied. In this regime, the effective current $J_{\text{eff}} \sim e_m^2 \mathcal{A} |\phi|^2$ scales as $1/m_\phi^2$, enhancing the signal at small m_ϕ .

On the other hand, when m_ϕ is very small, such that $\kappa \gg 1$, the $1/m_\phi$ scaling in ϕ is canceled, making J_{eff} finite even as $m_\phi \rightarrow 0$. If the mDM electric charge is large, i.e., $\epsilon, \kappa \gg 1$, Eq. (5) simplifies to $\phi = -i\sqrt{2\rho} v_\phi / (e_m B_0 R_e) \sin(m_\phi t)$. In this limit, the mDM is repelled by the geomagnetic potential, and the effective current saturates at $J_{\text{eff}} \sim \rho v_\phi^2 / (B_0 R_e)$, independent of m_ϕ and e_m , corresponding to a constant magnetic signal of 1.43 pT in our Letter. Further discussion of the mDM wave function behavior is provided in the Supplemental Material [81].

Experimental setup and sensitivity—Let us consider a region of size L that is bounded by a conducting shielding. A dark matter induced oscillating effective current J_{eff} drives an electromagnetic field inside the region that is

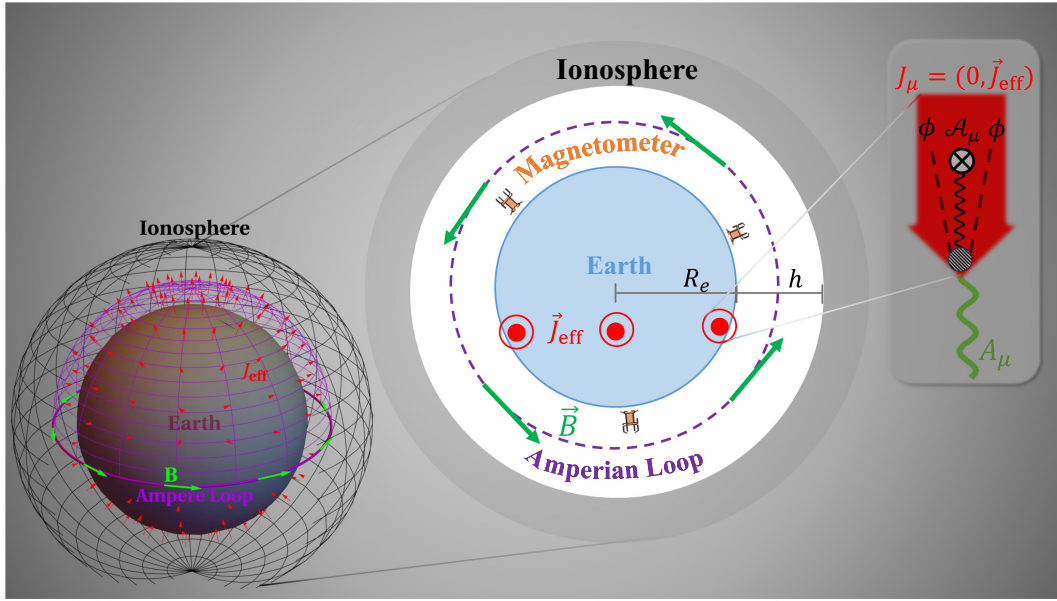


FIG. 1. Middle: Natural Earth's cavity formed between the Earth's surface and the ionosphere. The radial component of the effective current $\vec{J}_{\text{eff}}$ induced by mDM passing through a chosen Amperian loop. The generated magnetic field $\vec{B}$ can be probed by the magnetometers (orange symbols) placed over the surface of the Earth. Lower left: The 3D version of the middle, where the effective current is depicted from Eq. (8). Upper right: The signal magnetic field produced by the millicharged effective current in the geomagnetic field background.

damped in the boundaries. When the frequency of the oscillating current is much smaller than $1/L$, the signal is dominated by a quasistatic magnetic field whose strength is roughly of the form $B \sim J_{\text{eff}}L$. Based on this, novel experimental ideas for axions [82,83] and dark photons [84] were proposed at the laboratory scale for masses between 10^{-10} and 10^{-6} eV, where one can benefit from resonant signal enhancements and strict control of the background noise.

In this Letter, we propose to search for bosonic mDM by applying the same principle mentioned above to the Earth's natural conducting cavity. As illustrated in Fig. 1, the Earth can play the role of a natural cavity shielded from below by the Earth's surface and from above by the ionosphere (or interplanetary medium). We set the Earth's radius as $R_e = 6371.2$ km; therefore, the proposal applies for masses below $1/R_e \sim 10^{-14}$ eV. For such low masses, we benefit in sensitivity due to the Earth's size and the $1/m_\phi^2$ scaling of the signal. Using high-precision magnetometers situated over the surface of the Earth, it is possible to probe ultralight mDM in a wide, unexplored parameter space.

The profile of the vector potential $\vec{A}$ above the Earth's surface is determined by the profile of the geomagnetic field $\vec{B}$ extended over the whole Earth region. We use the updated IGRF model [85] to describe the geomagnetic field from the atmosphere into the mantle, while for the inner and outer cores, we model it by resorting to the electric current that produces it. The most popular theory about the

origin of the Earth's magnetic field is the geodynamo, in which the geomagnetic field is generated by electric currents $\vec{J}$ in the Earth's outer core [86,87]. Currents in the solid inner core and the mantle are vanishing. To get the profile of $\vec{J}$, we use polynomial functions that fit the rms value of the geomagnetic field B over the outer core and respect appropriate boundary conditions. Then, the magnetic field $\vec{B}$ and vector potential $\vec{A}$ can be calculated by solving $\nabla \times \vec{B} = \vec{J}$ and $\nabla \times \vec{A} = \vec{B}$, respectively, together with the conditions $\vec{\nabla} \cdot \vec{B} = 0$ and $\vec{\nabla} \cdot \vec{A} = 0$. See all the details of this calculation in Supplemental Material [81].

As a result, the external vector potential above the Earth's surface, i.e., $r \geq R_e$, is given by

$$\vec{A} = -B_0 R_e \left(1.01\beta \left(\frac{R_e}{r} \right)^3 (2\vec{Y}_{10} - \vec{\Psi}_{10}) - \sqrt{\frac{4\pi}{3}} \left(\frac{R_e}{r} \right)^2 \vec{\Phi}_{10} \right), \quad (6)$$

where $B_0 = 2.94 \times 10^{-2}$ mT is the total geomagnetic field at the equator of the Earth's surface, and $\vec{Y}_{10}$, $\vec{\Psi}_{10}$, $\vec{\Phi}_{10}$ are vector spherical harmonics. The parameter β accounts for the uncertainty in the rms value of the geomagnetic field B in the outer core. The value of β ranges between 1 and 4.6.

Knowing the external vector potential $\vec{A}$, we can analytically solve Eq. (3). For our mass range of interest,

where $m_\phi \ll 1/R_e$ is satisfied, Eq. (3) can be approximated to the magneto quasistatic equation

$$\vec{\nabla} \times \vec{B} = \vec{J}_{\text{eff}}, \quad (7)$$

where $\vec{B} = \vec{\nabla} \times \vec{A}$ is the magnetic field signal and $\vec{J}_{\text{eff}}$ the dark matter effective current given by

$$\begin{aligned} \vec{J}_{\text{eff}} = & \frac{2e_m^2 \rho B_0 R_e}{m_\phi^2} \left(1.01\beta \left(\frac{R_e}{r} \right)^3 (2\vec{Y}_{10} - \vec{\Psi}_{10}) \right. \\ & \left. - \sqrt{\frac{4\pi}{3}} \left(\frac{R_e}{r} \right)^2 \vec{\Phi}_{10} \right) e^{-2im_\phi t}. \end{aligned} \quad (8)$$

In our calculation, we use the fact that both the Earth's surface and the ionosphere layers are good conductors. Then, Eq. (7) can be solved with the boundary condition $\vec{B}_r = 0$ at $r = R_e$ and $r = R_e + h$, where $\vec{B}_r$ is the radial component of the magnetic field and h is the ionosphere height. We note that, at very low frequencies and depending on solar activity and other factors, the ionosphere may not be a good conductor that can damp the electromagnetic active mode completely. In such a case, the upper boundary would be the interplanetary medium. For more details of the ionosphere and interplanetary medium conducting properties, see Ref. [72].

Finally, by solving Eq. (7) together with the condition $\vec{\nabla} \cdot \vec{B} = 0$, the magnetic field signal at the Earth's surface is given by

$$\vec{B}(t) = \frac{e_m^2 \rho B_0 R_e^2}{m_\phi^2} \left(\sqrt{\frac{4\pi}{3}} \frac{h}{R_e} \vec{\Psi}_{10} - 2.02\beta \vec{\Phi}_{10} \right) e^{-i2m_\phi t}. \quad (9)$$

The dominant contribution comes from the $\vec{\Phi}_{10}$ term of the signal. The suppression factor $h/R_e \ll 1$ in the $\vec{\Psi}_{10}$ term can be understood from Ampere's law, $\oint d\vec{l} \cdot \vec{B} = \int d\vec{S} \cdot \vec{J}_{\text{eff}}$, where an integral loop is shown in Fig. S2 in the Supplemental Material [81]. For pure tangential components of the effective current in Eq. (8), i.e., the $\vec{\Phi}_{10}$ term, the integral $\int d\vec{S} \cdot \vec{J}_{\text{eff}}$ gives $\sim \pi R_e h J_{\text{eff}}$, while $\oint d\vec{l} \cdot \vec{B}$ becomes $\sim B R_e$, resulting in $|B| \sim \pi R_e (h/R_e) J_{\text{eff}}$. On the other hand, for radial components of the effective current, such as the $\vec{Y}_{10}$ term in Eq. (8), we can use the Amperian loop shown in Fig. 1 to obtain $\int d\vec{S} \cdot \vec{J}_{\text{eff}} \sim \pi R_e^2 J_{\text{eff}}$, leading to $B \sim \pi R_e J_{\text{eff}}$.

For $R_e = 6371.2$ km, $B_0 = 2.94 \times 10^{-2}$ mT, $\beta = 1$, and $\rho = 0.3$ GeV/cm³, we notice that the amplitude of the magnetic field signal is

$$|B| = 696 \text{ pT} \left(\frac{e_m}{10^{-27}} \right)^2 \left(\frac{10^{-14} \text{ eV}}{m_\phi} \right)^2, \quad (10)$$

meaning that for our mass ranges of interest, which correspond to signal frequencies below 5 Hz, an mDM field can produce a magnetic field signal of hundreds of pT for millicharges down to $e_m = 10^{-27}$. Current unshielded magnetometers can easily reach these sensitivities. To have an idea, in the first campaign of the SNIPE Hunt collaboration, magnetometers with $300 \text{ pT}/\sqrt{\text{Hz}}$ were used [76].

A straightforward concern for a dedicated experimental exploration is the background noise. For the frequencies of interest, the dominant magnetic field interference is of anthropogenic origin [88]. Therefore, to get optimal sensitivity, the experiment can be set up in a location far away from urban environments. Moreover, the mDM signal oscillates at a particular frequency that makes it easy to be distinguished from static sources, such as the strong static of Earth's geomagnetic field. In fact, it was demonstrated in the first SNIPE Hunt expedition that the amplitude spectral density of the measurements were flat and corresponded to the noise floor of the magnetometers, with the exception of a 60 Hz peak associated with the laptop used for data acquisition. While a few narrow peaks at one of the measurement stations are attributed to noise with unknown origin, our signal should be present in all stations at all times. In addition, since the reach of the SNIPE Hunt is now limited by the sensitivity of the magnetometers [76] rather than by the geomagnetic noise, future dedicated experiments using high-precision magnetometers with sensitivity up to order $20 \text{ fT}/\sqrt{\text{Hz}}$ [89–92] will result in stronger sensitivity. Moreover, methods to improve sensitivity and frequency range coverage have already been examined. For instance, it has been pointed out in [93] that the measurement of the curl of the magnetic field can serve as a useful method to reach much higher frequencies.

In Fig. 2, we show constraints for mDM by recasting the null results of searching for axions and dark photons from the SuperMAG [72–75] and SNIPE Hunt [76] collaborations. The SNIPE Hunt collaboration performed coordinated measurements with a network of magnetometers located in the US in three different stations far from human-generated magnetic noise. Each station used a vector magnetoresistive sensor with sensitivity of $300 \text{ pT}/\sqrt{\text{Hz}}$ over a frequency range of 0.1–100 Hz [76]. The SuperMAG data consist of two sets of measurements taken over several tens of years. One set was conceived at one-minute intervals with contributions from $\mathcal{O}(500)$ stations all over the world [72–74]. The other set corresponds to measurements with a one-second time resolution with contributions from $\mathcal{O}(200)$ stations. This “high-fidelity” dataset was analyzed for the cases of dark photon and axion dark matter in Ref. [75]. To obtain the SNIPE Hunt constraints for mDM, we first calculate the magnetic field amplitude $|\vec{B}_a|$ by using the constraints found in Ref. [76] for three different stations, and then, obtain an averaged value. In the same frequency, we equal this average to the counterpart magnetic field signal in Eq. (9) to estimate the bounds on e_m . For the

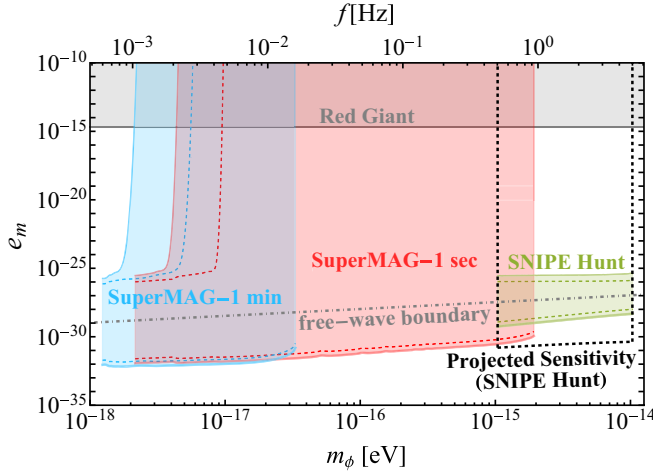


FIG. 2. Sensitivity for millicharged dark matter in the plane of m_ϕ and e_m . The blue, red, and green shaded regions are excluded by the null results from the SuperMAG-1 min, SuperMAG-1 sec, and SNIPE Hunt experiments, respectively. The dashed (solid) lines correspond to $\beta = 1$ (4.6), reflecting the uncertainty on the geomagnetic field rms value over the Earth’s outer core. The lower edge of our constraint is well below the free-wave boundary $\kappa = 1$ (gray dash-dotted line), as given by Eq. (10). We also show the bound from red giant cooling [65] and the projected sensitivity of the SNIPE Hunt experiment using a LEMI-120 sensor with one week of data [94]. This projected reach fully covers the previously unconstrained region, leaving no lower limit in the parameter space.

SuperMAG constraints, since the data are collected with many stations situated over the world, we instead average the $\vec{B}_a$ over the Earth’s surface to get magnetic field expressions independent of the Earth location. Then, we proceed in the same way as done for SNIPE Hunt. As we can see in Fig. 2, the estimated constraints surpass previous bounds from stellar cooling by over 13 orders of magnitude in the case of SNIPE Hunt and over 17 orders of magnitude for SuperMAG.

We note that regions with large e_m , where mDM is repelled significantly by the geomagnetic potential, remain unconstrained by current SuperMAG and SNIPE Hunt data. This is because, at the edge of our sensitivity window for m_ϕ , the minimum detectable signal exceeds the saturation value of 1.43 pT, as discussed below Eq. (5). As a result, mDM with sufficiently large charges becomes undetectable in this regime. These lower bounds are derived using the wave function from Eq. (5) in the signal estimation. Still, within the relevant mass range, enhancing the magnetometer sensitivity by a factor of 3 would suffice to fully cover the large-charge parameter space. This points to a realistic and promising direction for future experimental efforts to close the remaining gap in the mDM search. As shown in Fig. 2, the projected sensitivity of a high-performance LEMI-120 induction coil magnetometer, operating in the 0.5–5 Hz range for one week, can reach this level [94].

Conclusion and outlook—We have presented a novel approach to probe ultralight mDM by measuring their geomagnetic signals. Using classical field theory, we calculated the quasistatic monochromatic magnetic field sourced by the effective current associated with the mDM field, which exhibits a distinctive inverse-square dependence on the mDM mass. By reinterpreting the data from the SuperMAG and SNIPE Hunt experiments, we have derived stringent constraints on the effective charge of bosonic mDM over specific mass ranges. These bounds exceed stellar cooling limits by over 13 and up to 18 orders of magnitude. For two narrow mass windows, unconstrained regions remain at large e_m due to the repulsion of mDM by the geomagnetic potential. Nonetheless, a factor of 3 improvement in magnetometer sensitivity would be sufficient to close these gaps. Finally, a global network of high-precision magnetometers monitoring Earth’s magnetic field could offer a powerful platform for ultralight mDM detection in the future.

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

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