

Search for QCD Coupled Axion Dark Matter with Data from the MICROSCOPE Space Experiment

Jordan Gué^{1,*}, Peter Wolf², and Aurélien Hees²

¹*Institut de Física d'Altes Energies (IFAE), The Barcelona Institute of Science and Technology, Campus UAB, 08193 Bellaterra (Barcelona), Spain*

²*LTE, Observatoire de Paris, Université PSL, Sorbonne Université, Université de Lille, LNE, CNRS, 61 Avenue de l'Observatoire, 75014 Paris, France*



(Received 2 April 2025; revised 17 June 2025; accepted 3 September 2025; published 21 October 2025)

Axion dark matter coupled via QCD induces a nonzero differential acceleration between test masses of different composition. Tests of the equivalence principle, like the recent MICROSCOPE space mission, are sensitive to such a signal. We use the final released data of the MICROSCOPE experiment to search for this effect. We find no positive signal consistent with the dark matter model, and set upper limits on the axion-gluon coupling that improve existing laboratory bounds by up to 2 orders of magnitude for axion masses in the 10^{-17} to 10^{-13} eV range.

DOI: [10.1103/fwzv-gl1jg](https://doi.org/10.1103/fwzv-gl1jg)

Introduction—Almost 100 years after its first observational evidence [1], the fundamental nature of dark matter (DM) is still to this day one of the biggest mysteries in fundamental physics [2]. Out of all the possible DM candidates, ultralight fields (whose mass is below the eV) are a compelling alternative to historically dominant models, such as weakly interacting massive particles (WIMPs). In these models, DM is described by a non-relativistic classical wave oscillating at its Compton frequency [3,4]. One of the most prominent ultralight dark matter (ULDM) candidates is the QCD axion [5–7], originally postulated 50 years ago to solve the so-called strong CP problem [8]. More general pseudoscalars, more commonly referred to as axionlike particles (ALPs), which do not solve the strong CP problem, can still account for the galactic DM and, therefore, are actively searched for [9–11].

In the simplest scenario, the mixing between the (QCD) axion a and neutral pions (with strength $1/f_a$) naturally solves the strong CP problem, leading to the so-called linear relationship between the axion mass and $1/f_a$ [12,13]. However, some models allow for such mixing without the need of this relation, i.e., at fixed coupling value, lower mass axions are feasible [14,15].

As a consequence of the QCD-axion coupling, the axion potential becomes dependent on the light quarks and pion

masses after the QCD phase transition [12,16]. In particular, to preserve the shift symmetry of the axion field, this potential is periodic and is quadratic in a at the leading order in $a/f_a \ll 1$. Because of their mixing with axions, pions share the same potential, and as a consequence, their mass depends quadratically on the axion field, see, e.g., Refs. [12,17]. Among other observables, this can lead to very strong forces mediated by the axion between very dense astrophysical objects, such as neutron stars [18], a change of the composition of white dwarfs [19] and neutron stars [20], or a model-independent axion production channel in supernovae [21].

The mass of the nucleons depends on the pion mass, and this dependence has been computed at the one-loop level, see, e.g., Refs. [17,22–24]. Similarly, the nucleon g factor at the chiral order also depends on the pion mass [17,24] such that atomic transition frequencies are dependent on the axion field [17]. Finally, the atomic nuclear binding energy at chiral order is also dependent on the pion mass; see Ref. [25] for a detailed derivation. As a consequence, the atomic rest mass becomes quadratically dependent on the axion field [26,27]. This relationship is nonuniversal in the sense that it depends on the atomic number and on the atomic mass, which leads to a breaking of the universality of free fall (UFF) and a violation of the equivalence principle [26].

At cosmological scales, the axion oscillates in its self-potential at its Compton frequency [9], and these oscillations may comprise DM. In this scenario, the oscillation's amplitude is stochastically related to the DM energy density, and the pion masses, atomic frequencies, and atomic rest masses all become time dependent [12,17,26]. These oscillations have been searched for in several atomic oscillators [28,29], but recently it was suggested that they can also be

*Contact author: jgue@ifae.es

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.

searched for in classical UFF tests or using atom interferometers [26].

In addition, quadratic axion-pion coupling induces more complex signatures similarly to the ones arising for quadratically coupled pure scalar fields [30]. In particular, in the vicinity of massive bodies, the backaction of matter on the axion field can strongly impact its oscillation amplitude, a property that was originally derived under the spherical symmetry assumption in [30] where the small field gradient was neglected. Recently, those solutions have been extended to account for the field propagation [31,32].

The matter-axion coupling being quadratic implies that in the presence of a massive body the breaking of the UFF induced by an ALP has two distinct signatures [30]: (i) oscillations at twice the axion Compton frequency and (ii) a static differential acceleration directed toward the center of the massive body, which corresponds to the one searched for in standard UFF tests. Currently, the best UFF test is provided by the MICROSCOPE experiment [33,34], a space-based experiment which monitored the differential acceleration between two test masses made of platinum and titanium and which provided a constraint on the Eötvös parameter at the level of 10^{-15} [33].

In this Letter, we reanalyze the publicly available MICROSCOPE data [33,34] to search for axion DM through its coupling to gluons. To do so, we recombine the individual measurement session results given in [34] in a way that is optimized for our DM search. We find no evidence of axion DM coupled to gluons, and this allows us to set a new upper bound on the coupling strength $1/f_a$. Our constraint covers 7 orders of magnitude in mass from 10^{-20} to $\sim 10^{-13}$ eV and improves current laboratory bounds by up to 2 orders of magnitude for axion masses between 10^{-17} and $\sim 10^{-13}$ eV, reaching $1/f_a \sim 10^{-16}$ GeV $^{-1}$ at 10^{-20} eV.

Breaking of the universality of free fall by an axion field—As mentioned above, the axion-gluon coupling leads to a quadratic dependence of the pion mass on the axion field and subsequently to a violation of the universality of free fall. We will present the theoretical modeling of the axion induced UFF breaking. A detailed derivation and discussion are presented in Sec. A of Supplemental Material [35].

We start from the Lagrangian [43],

$$\mathcal{L} = -\frac{1}{2\kappa} g^{\mu\nu} \partial_\mu a \partial_\nu a - \frac{m_a^2 c^2}{2\hbar^2 \kappa} a^2 + E_P \frac{g_3^2}{32\pi^2 f_a} G_{\mu\nu}^a \tilde{G}^{a,\mu\nu}, \quad (1)$$

where $\kappa = 8\pi G/c^4$ is the Einstein gravitational constant, $g_{\mu\nu}$ is the spacetime metric, a is the dimensionless ALP of mass m_a , $E_P = \sqrt{\hbar c/\kappa}$ is the reduced Planck energy, f_a is the ALP-gluon coupling (which corresponds to the Peccei-Quinn spontaneous symmetry breaking scale [8]), $G_{\mu\nu}$ is the QCD strength tensor, and g_3 is the strong coupling.

This Lagrangian leads to a mixing between the ALP and the neutral pion, as both are pseudo-Goldstone bosons, such that the pion mass m_π becomes quadratically dependent on $\theta \equiv E_P a/f_a$ [12,17],

$$\frac{\delta m_\pi^2}{m_\pi^2} = -\frac{m_u m_d}{2(m_u + m_d)^2} \theta^2, \quad (2)$$

where $m_{u/d}$ are the up/down quark masses. In the chiral perturbation theory, the mass of the nucleons [17,22–24] and the nuclear binding energy [25] are both dependent on the pion mass. When combined together, this implies that the rest mass of an atom A depends quadratically on the ALP field, with an atom-dependent coefficient, the axionic charge Q_M^A , i.e.,

$$m_A = m_A^0 \left(1 + Q_M^A \frac{E_P^2}{f_a^2} a^2 \right), \quad (3)$$

where m_A^0 is the unperturbed mass. These charges depend on the atomic mass number A and charge number Z , and they are derived explicitly in [26,44]. The nonuniversality of Q_M^A induces a violation of the UFF [26], which, in the case of quadratic couplings, leads to a differential acceleration between two colocated test masses A and B [30],

$$\Delta \ddot{\vec{x}} = -2(Q_M^A - Q_M^B) \left(\frac{E_P}{f_a} \right)^2 a \left(c^2 \vec{\nabla} a + \vec{v} \dot{a} \right). \quad (4)$$

At cosmological scales, the axion oscillates around the minimum of its potential at its Compton frequency $\omega_a = m_a c^2/\hbar$ [9]. If the axion constitutes DM, its oscillation amplitude a_0 depends on the DM energy density through $a_0 = \sqrt{16\pi G \rho_{\text{DM}}}/\omega_a c$ with $\rho_{\text{DM}} = 0.4$ GeV/cm ^{3} [45]. Locally, close to massive bodies, the backaction of matter on the axion field can strongly impact its oscillation amplitude, which becomes location dependent, i.e.,

$$a(t, \vec{r}) = a_0 A(\vec{r}) \cos[\omega_a t + \Phi(r)], \quad (5a)$$

where $\Phi(r)$ is a phase, which can be position dependent as well (see below and in Supplemental Material [35]).

The solution $A(\vec{r})$ shows a scalarization mechanism [i.e., $A(r)$ can become large close to massive bodies] [30]. Neglecting (i) the backreaction from the metric and (ii) the small contribution from the axion asymptotic gradient $k \sim \omega_a v_{\text{DM}}/c^2 \sim 10^{-3} \omega_a/c$ (with v_{DM} the DM velocity in the Solar System frame) can lead to divergences of the amplitude for certain values of the coupling parameter [30].

Recently, these solutions have been extended to account for k under the form of a multipole extension [31,32]. As pointed out in [31,32], one important contribution of the gradient is to regularize the divergences of the original solution which become resonances. Note that the monopole contribution is the dominant one by a factor

$1/kR_E \gtrsim 10^3 \gg 1$ (see Supplemental Material [35]) such that we consider only this one in our data analysis, i.e., the $A(r)$ function writes

$$A(r) = \frac{|\lambda \cos \lambda \sin(kr - \beta) + \beta \cos(kr - \beta) \sin \lambda|}{kr \sqrt{\lambda^2 \cos^2 \lambda + \beta^2 \sin^2 \lambda}}, \quad (5b)$$

where

$$\beta = kR_E, \lambda = \sqrt{\beta^2 + 12 \left(\frac{E_P}{f_a} \right)^2 |Q_M^E| \frac{GM_E}{R_E c^2}}, \quad (5c)$$

where R_E is the Earth radius modeled as a spherically homogeneous body of mass M_E and of charge Q_M^E ; see Ref. [44].

In standard models of galaxy formation, galactic DM must be virialized [46,47] such that its velocity distribution acquires a characteristic width $\sigma_v \sim 10^{-3}c$ [4,36], and the DM oscillations have a characteristic coherence time $\tau \sim c^2/\sigma_v^2 \times 2\pi\hbar/(m_a c^2)$ [37]. This implies that, for observation times longer than τ , the DM spectral signature is broadened [38,39]. For observation times shorter than τ , the signal is effectively monochromatic, but its amplitude receives a stochastic correction [38,39]. With the virialized DM particles being uncorrelated, all have a different (uniformly distributed) phase. This is mathematically equivalent to introducing a global stochastic amplitude α , following a Rayleigh distribution, and assigning the same phase (uniformly distributed) to all the DM particles, such that the amplitude of the axion wave becomes proportional to α [38–40,48].

The axion field of Eq. (5a) induces two types of signatures on the UFF signal of Eq. (4): (i) a fast oscillating signal (at frequency $2\omega_a$) and (ii) a static position dependent signal. In this Letter, we focus on the static signal, which gives a relative differential acceleration between two test masses of

$$\eta = \frac{|\Delta \vec{x}|}{|\vec{x}|} = \alpha^2 \Delta Q_M \frac{c^2 a_0^2}{2 \frac{GM_E}{r^2}} \left(\frac{E_P}{f_a} \right)^2 \frac{\partial A^2(r)}{\partial r} \equiv \gamma \alpha^2, \quad (6)$$

where η is the Eötvös parameter, $\Delta Q_M = (Q_M^B - Q_M^A)$, and α^2 , which follows an exponential distribution, is due to the DM velocity distribution. The γ function, which characterizes the strength of the signal searched for in the data, is depicted for one axion mass on Fig. 1 (see Supplemental material [35] for more details). The above signal depends on $\vec{k}$; see Eqs. (5b) and (5c), which is a stochastic variable depending on the DM velocity distribution such that one should marginalize over this parameter. However, in practice, the impact of $\vec{k}$ on the signal is negligible (k appears in highly suppressed factors), except on extremely narrow $1/f_a$ coupling bands at resonances (in

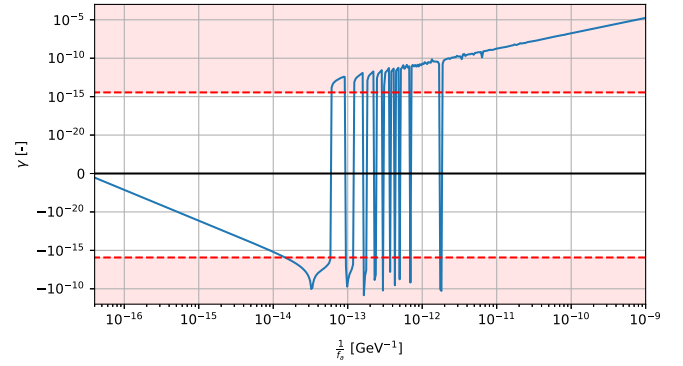


FIG. 1. In blue the evolution of the signal strength γ [cf. Eq. (6)] as a function of the coupling $1/f_a$ for one value of the axion mass ($m_a = 10^{-16}$ eV). The red dashed lines correspond to the lower and upper limits from the MICROSCOPE data analysis, and the red filled surface corresponds to the exclusion area.

practice, never exactly sampled and not visible on the constraint). This has been checked analytically and numerically (see Supplemental Material [35]).

MICROSCOPE data—MICROSCOPE is a space mission designed to probe the UFF in space at the level of 10^{-15} [41] by measuring the differential acceleration between a test mass made of titanium and another one made of platinum. The two free falling bodies are concentric cylindrical test masses controlled in a differential accelerometer. In total, 18 measurement sessions taken between 2016 and 2018 have been used to constrain the Eötvös parameter [34]. The duration of a session lasts between 13 hours and 8 days. Each session has been analyzed independently providing an estimate of the Eötvös parameter η_i , a statistical error $\sigma_{\text{stat},i}$, and a systematic error $\sigma_{\text{syst},i}$ [34], which are provided in Supplemental Material [35]. Note that, in our analysis, we quadratically combine the statistical and systematic uncertainties.

Bayesian parameter inference—The methodology used in this analysis is thoroughly detailed in Sec. B of Supplemental Material [35]. We present here the important steps. We use a Bayesian approach and assume the measurement noise to follow a Gaussian distribution.

For a given axion mass m_a , we compute the axion field coherence time, and we group the MICROSCOPE sessions into subsets where the total duration spanned by a subset of sessions is smaller than the coherence time. Within each subset of sessions, we can, therefore, use the monochromatic modeling of the axion field discussed above. Note that the α stochastic parameter is different and independent for all the subsets. Within each subset of sessions, we combine the individual likelihoods and subsequently marginalize over their α parameter. The combined likelihood of all the measurements is simply provided by the product of the marginalized likelihood from all the subsets.

Following [39], we use a noninformative Berger-Bernardo reference prior, which in our case is equivalent

to Jeffrey's prior. This noninformative prior is motivated because (i) it ensures invariance of the inference results under a change of variables, (ii) it ensures that the integral of the posterior converges (which is not the case if one chooses naively a uniform prior on γ), (iii) it maximizes the Kullback-Leibler of the posterior with respect to the prior, and (iv) it produces results consistent with the frequentist approach [39]. This prior is estimated numerically (see Supplemental material [35] for more details).

The product of the marginalized likelihood with the prior allows us to estimate the posterior distribution of the γ parameter, which characterizes our signal; see Eq. (6). The ratio of the Bayesian evidence between this model which contains an axion signal and a model which includes only noise is then computed and used to assess if an axion signal is significantly detected within the MICROSCOPE data.

In the case of no positive detection, we use the posterior probability distribution to infer a 95% credible interval on γ . This interval depends on the mass of the axion field m_a and is illustrated by the red dashed curves on the bottom panel from Fig. 1 in the case of $m_a = 10^{-16}$ eV. The nonlinear expression of γ as a function of $1/f_a$ allows us to translate the constraint on γ into a constraint on $1/f_a$.

This methodology is applied for different axion masses.

Results and discussion—We first assess if there is a positive DM signal in the data by computing the Bayes factor $\mathcal{B}$ as the ratio of evidences of the ALP-gluon coupled model and a model assuming no signal. Over the full mass range of interest, $\mathcal{B}$ is < 1 (see Supplemental Material [35]), which indicates no significant detection of DM in the data.

Since no positive signal is detected in the data, we compute an exclusion area on the value of the coupling $1/f_a$. In Fig. 2, we present in red the constraint on $1/f_a$ from our analysis, while existing laboratory constraints from [28,49,50] are shown by the gray-filled area. We find that MICROSCOPE's data constrains new regions of the parameter space, with its most stringent limit being $1/f_a \sim 1.7 \times 10^{-16} \text{ GeV}^{-1}$ at $m_a = 10^{-20}$ eV and reaching $1/f_a \sim 1.3 \times 10^{-10} \text{ GeV}^{-1}$ at $m_a \sim 7.7 \times 10^{-14}$ eV, improving by roughly 2 orders of magnitude the existing laboratory bounds over 4 orders of magnitude of mass. The lower mass part of the constraint of MICROSCOPE depicted in Fig. 2 is a factor 3 less stringent than results obtained previously for different theoretical models (such as [30,51]) because previous studies neglected the stochastic behavior of the field. At larger masses, the stochastic behavior of the signal impacts more significantly the resulting constraints. In particular, the constraint presents several bands. As shown in Fig. 1, the detector is blind to very narrow intervals of values of couplings, around when γ crosses zero. As discussed in Supplemental Material [35], this happens when the axion field reaches an extremum, i.e., when its derivative changes sign [52]. In addition, one can notice a significant difference in smoothness of the

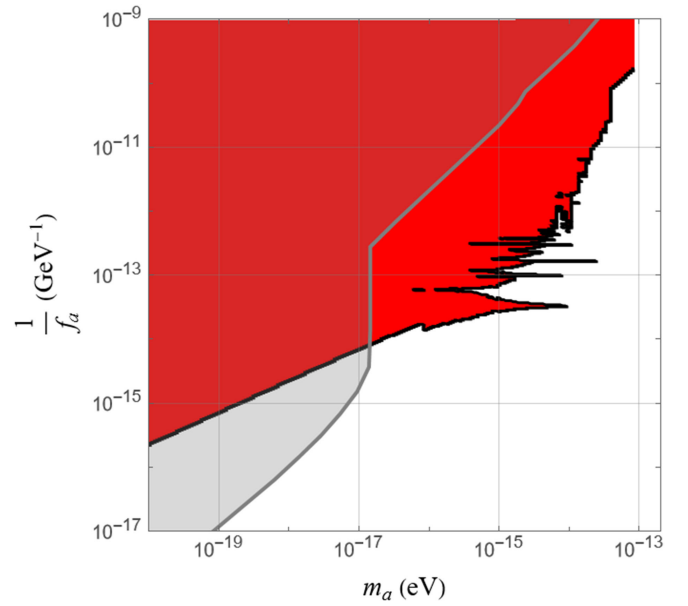


FIG. 2. 95% confidence area in the ALP-gluon coupling $1/f_a$ using final results from MICROSCOPE (in red). In gray are shown the existing constraints from laboratory experiments from [28,49,50] (adapted from [11]).

constraint between low and large masses. As described more thoroughly in Supplemental Material [35], this arises because the various MICROSCOPE sessions need to be grouped in subsets that last one coherence time. Small discontinuities in the constraint plot arises when the session grouping changes. At low masses, the coherence time is larger than the total time spanned by the MICROSCOPE mission, such that a single group including all sessions is possible. At larger masses, the axion coherence time decreases, and above $m_a \sim 10^{-16}$ eV one must necessarily cut the sessions at least in two or more, resulting in the small discontinuities visible in the constraint figure. The small decrease of the constraint visible at $m_a \sim 10^{-16}$ eV can be explained because, on each group of data, the signal to be fitted is proportional to the exponentially distributed parameter α^2 [see Eq. (6)], whose preferred value is 0. For this reason, the probability of one single random draw being close to 0 is larger than when drawing twice or more. Put in other words, the signal is more likely to have a small value when only one group of data is considered and the constraint is larger than when several groups of data are considered. In addition, as the axion mass increases, its coherence time decreases, such that fewer and fewer MICROSCOPE subsets of data can be constructed with the requirements $T_{\text{obs}}^{\text{subset}} < \tau$. As a consequence, the constraint derived from our data analysis decreases.

In conclusion, the ALP-gluon coupling $1/f_a$ leads to a static violation of the UFF, and classical tests like MICROSCOPE can be used to constrain such coupling. In this Letter, we performed a reanalysis of MICROSCOPE

final data, and we have seen no evidence for existence of an axion-gluon coupling for ALP masses between 10^{-20} and 10^{-13} eV, which allows us to derive an exclusion area on its value in this corresponding mass range. Our results improve existing laboratory bounds by up to 2 orders of magnitude and over 4 orders of magnitude of mass, and can be complementary to astrophysical and cosmological bounds [18,53,54]. In the future, these constraints could be improved by new tests of the UFF. In particular, several projects such as MICROSCOPE 2 [55] or STE-QUEST [56] aim at testing the UFF at the 10^{-17} level, improving by more than 2 orders of magnitude the MICROSCOPE results. Considering that $\eta \sim (1/f_a)^{1/4}$, such missions would improve the current constraints by up to a factor 5. In addition, other lab experiments such as atomic comparisons would also be able to probe such regions of the parameter space [17].

Acknowledgments—We thank the CNES and ONERA teams for making available the MICROSCOPE result, and in particular Joël Bergé, Gilles Métris, and Manuel Rodrigues. This work was supported by the Programme National GRAM of CNRS/INSU with INP and IN2P3 cofunded by CNES. IFAE is partially funded by the CERCA program of the Generalitat de Catalunya. J. G. is funded by Grant No. CNS2023-143767. Grant No. CNS2023-143767 is funded by MICIU/AEI/10.13039/501100011033 and by European Union NextGenerationEU/PRTR.

Data availability—The data supporting this study's findings are available within the article and its Supplemental Material.

[1] F. Zwicky, *Helv. Phys. Acta* **6**, 110 (1933).
 [2] G. Bertone and T. M. P. Tait, *Nature (London)* **562**, 51 (2018).
 [3] L. Hui, *Annu. Rev. Astron. Astrophys.* **59**, 247 (2021).
 [4] M. Cirelli, A. Strumia, and J. Zupan, *arXiv:2406.01705*.
 [5] J. Preskill, M. B. Wise, and F. Wilczek, *Phys. Lett.* **120B**, 127 (1983).
 [6] L. F. Abbott and P. Sikivie, *Phys. Lett.* **120B**, 133 (1983).
 [7] M. Dine and W. Fischler, *Phys. Lett.* **120B**, 137 (1983).
 [8] R. D. Peccei and H. R. Quinn, *Phys. Rev. Lett.* **38**, 1440 (1977).
 [9] D. J. E. Marsh, *Phys. Rep.* **643**, 1 (2016).
 [10] C. O'Hare, *Proc. Sci. COSMICWISPerS* (2024) 040.
 [11] C. O'Hare, *cajohare/axionlimits: Axionlimits*, <https://cajohare.github.io/AxionLimits/> (2020).
 [12] G. G. di Cortona, E. Hardy, J. P. Vega, and G. Villadoro, *J. High Energy Phys.* **01** (2016) 034.
 [13] S. Weinberg, *Phys. Rev. Lett.* **40**, 223 (1978).
 [14] L. Di Luzio, B. Gavela, P. Quilez, and A. Ringwald, *J. High Energy Phys.* **05** (2021) 184.
 [15] A. Hook, *Phys. Rev. Lett.* **120**, 261802 (2018).

[16] P. Di Vecchia and G. Veneziano, *Nucl. Phys.* **B171**, 253 (1980).
 [17] H. Kim and G. Perez, *Phys. Rev. D* **109**, 015005 (2024).
 [18] A. Hook and J. Huang, *J. High Energy Phys.* **06** (2018) 036.
 [19] R. Balkin, J. Serra, K. Springmann, S. Stelzl, and A. Weiler, *Phys. Rev. D* **109**, 095032 (2024).
 [20] R. Balkin, J. Serra, K. Springmann, S. Stelzl, and A. Weiler, *J. High Energy Phys.* **02** (2025) 141.
 [21] K. Springmann, M. Stadlbauer, S. Stelzl, and A. Weiler, *arXiv:2410.19902*.
 [22] J. Gasser, M. E. Sainio, and A. Švarc, *Nucl. Phys.* **B307**, 779 (1988).
 [23] V. Bernard, N. Kaiser, J. Kambor, and U.-G. Meißner, *Nucl. Phys.* **B388**, 315 (1992).
 [24] S. Scherer and M. Schindler, *A Primer for Chiral Perturbation Theory*, Lecture Notes in Physics, Vol. 830 (Springer, Berlin Heidelberg, 2012).
 [25] T. Damour and J. F. Donoghue, *Phys. Rev. D* **82**, 084033 (2010).
 [26] J. Gué, A. Hees, and P. Wolf, *Phys. Rev. D* **110**, 035005 (2024).
 [27] M. Bauer and S. Chakraborti, *arXiv:2408.06408*.
 [28] E. Madge, G. Perez, and Z. Meir, *Phys. Rev. D* **110**, 015008 (2024).
 [29] X. Zhang, A. Banerjee, M. Leyser, G. Perez, S. Schiller, D. Budker, and D. Antypas, *Phys. Rev. Lett.* **130**, 251002 (2023).
 [30] A. Hees, O. Minazzoli, E. Savalle, Y. V. Stadnik, and P. Wolf, *Phys. Rev. D* **98**, 064051 (2018).
 [31] A. Banerjee, I. M. Bloch, Q. Bonnefoy, S. A. R. Ellis, G. Perez, I. Savoray, K. Springmann, and Y. V. Stadnik, *arXiv:2502.04455*.
 [32] Y. G. del Castillo, B. Hammett, and J. Jaeckel, *arXiv:2502.04456*.
 [33] P. Touboul *et al.* (MICROSCOPE Collaboration), *Phys. Rev. Lett.* **129**, 121102 (2022).
 [34] P. Touboul *et al.*, *Classical Quantum Gravity* **39**, 204009 (2022).
 [35] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/fwzv-gljjg>, which includes Refs. [17,25,26,30–32,34,36–42], for details on the signal modeling and the data analysis.
 [36] N. W. Evans, C. A. J. O'Hare, and C. McCabe, *Phys. Rev. D* **99**, 023012 (2019).
 [37] A. Derevianko, *Phys. Rev. A* **97**, 042506 (2018).
 [38] J. W. Foster, N. L. Rodd, and B. R. Safdi, *Phys. Rev. D* **97**, 123006 (2018).
 [39] G. P. Centers, J. W. Blanchard, J. Conrad, N. L. Figueroa, A. Garcon, A. V. Gramolin, D. F. J. Kimball, M. Lawson, B. Pelssers, J. A. Smiga, A. O. Sushkov, A. Wickenbrock, D. Budker, and A. Derevianko, *Nat. Commun.* **12**, 7321 (2021).
 [40] E. Savalle, A. Hees, F. Frank, E. Cantin, P.-E. Pottie, B. M. Roberts, L. Cros, B. T. Mcallister, and P. Wolf, *Phys. Rev. Lett.* **126**, 051301 (2021).
 [41] P. Touboul *et al.*, *Phys. Rev. Lett.* **119**, 231101 (2017).
 [42] A. Hees and O. Minazzoli, *arXiv:1512.05233*.
 [43] We use a convention where the axion field amplitude is dimensionless; other conventions exist; see, e.g., Refs. [17,26].

- [44] Relevant to this analysis, the axionic charge of the Earth, modeled as a homogeneous SiO_2 ball, is $Q_M^E = -68.442 \times 10^{-3}$ and for the two MICROSCOPE test masses $Q_M^{\text{Pt}} = -69.065 \times 10^{-3}$ and $Q_M^{\text{Ti}} = -68.770 \times 10^{-3}$.
- [45] P. J. McMillan, *Mon. Not. R. Astron. Soc.* **414**, 2446 (2011).
- [46] K. Freese, M. Lisanti, and C. Savage, *Rev. Mod. Phys.* **85**, 1561 (2013).
- [47] A. Pillepich, M. Kuhlen, J. Guedes, and P. Madau, *Astrophys. J.* **784**, 161 (2014).
- [48] H. Nakatsuka, S. Morisaki, T. Fujita, J. Kume, Y. Michimura, K. Nagano, and I. Obata, *Phys. Rev. D* **108**, 092010 (2023).
- [49] I. Schulthess *et al.*, *Phys. Rev. Lett.* **129**, 191801 (2022).
- [50] C. Abel *et al.*, *Phys. Rev. X* **7**, 041034 (2017).
- [51] P. Brax, C. Burrage, J. A. R. Cembranos, and P. Valageas, *Phys. Rev. D* **107**, 095015 (2023).
- [52] Not all such bands are visible on Fig. 2 because its resolution is too low to see the narrower ones (cf. Fig. 1).
- [53] J. Zhang, Z. Lyu, J. Huang, M. C. Johnson, L. Sagunski, M. Sakellariadou, and H. Yang, *Phys. Rev. Lett.* **127**, 161101 (2021).
- [54] K. Blum, R. T. D’Agnolo, M. Lisanti, and B. R. Safdi, *Phys. Lett. B* **737**, 30 (2014).
- [55] J. Bergé, *Rep. Prog. Phys.* **86**, 066901 (2023).
- [56] N. Gaaloul *et al.*, *arXiv:2211.15412*.