Axion-plasmon polaritons in strongly magnetized plasmas

H. Terças, J. D. Rodrigues, and J. T. Mendonça

1 Instituto de Plasmas e Fusão Nuclear, Lisboa, Portugal
2 Instituto Superior Técnico, Lisboa, Portugal

Axions are hypothetical particles related to the violation of the charge-parity symmetry, being the most prone candidates for dark matter. Multiple attempts to prove their existence are currently performed in different physical systems. Here, we anticipate the possibility of the axions coupling to the electrostatic (Langmuir) modes of a strongly magnetized plasma, by showing that a new quasi-particle can be defined, the axion-plasmon polariton. The excitation of axions can be inferred from the pronounced modification of the dispersion relation of the Langmuir waves, a feature that we estimate to be accessible in state-of-the-art plasma-based experiments. We further show that, under extreme density and magnetic field conditions (e.g. at the interior of dense neutron stars), the axion-plasmon polariton becomes dynamically unstable, similarly to the case of the Jeans instability occurring in self-gravitating fluids. This latter result anticipates a plausible mechanism to the creation of axion-like particles in the universe.

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Introduction. The violation of the charge-parity (CP) symmetry is perhaps one of the most fundamental problems in modern physics [1, 2]. Although CP-violation is, by construction, inherent to the Standard Model (e.g. it appears as a mixing angle in the Cabibbo-Kobayashi-Maskawa (CKM) matrix describing quark masses [3, 4], and in the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix mixing lepton flavours [5]), there is no evidence of its manifestation in quantum chromodynamics (QCD). At the origin of the so-called strong CP problem is the anomalous electric dipole moment of the neutron [6], which according to a CP-broken QCD calculation would be of the order of $10^{-8}$ e.m., while experiments point towards a $10^9$ times larger value. An elegant - and probably the most consensual - way to solve the strong CP problem is the Peccei-Quinn (PQ) mechanism, in which the phase violating the CP-symmetry in the QCD Lagrangian is promoted to a complex field [7, 8]. The corresponding Goldstone pseudo-boson is known as the axion, and emerges after breaking the U(1) symmetry in the gluon coupling term [9].

Axion-like particles (ALPs) are hypothetical particles with a extremely small mass (possibly in the meV range) and couple very weakly with quarks, leptons and photons. ALPs have received renewed breath after being indicated as appealing candidates to dark matter in the universe [10, 11]. Several experiments are in operation for at least a decade with the goal of observing ALP signatures, using both laboratory and astrophysical observations [12–14]. Unfortunately, most of the observations are too dubious to confirm the existence of ALPs. For example, the PVLAS experiment - originally design to probe the birefringent properties of the electromagnetic vaccum [15] - advanced preliminary results indicat-
In a plasma, Eqs. (4) and (5) must be closed with the equations for the sources, \( \rho = -e(n_e - n_i) \) and \( \mathbf{u} = -\mathbf{J}/(en_e) \), namely
\[
\frac{\partial n_e}{\partial t} + \nabla \cdot (n_e \mathbf{u}) = 0,
\]
\[
\left( \frac{\partial}{\partial t} + \mathbf{u} \cdot \nabla \right) P = \frac{e}{m_e n_e} (\mathbf{E} + \mathbf{u} \times \mathbf{B}),
\]
where \( P = k_B T_e n_e^2 \) is the electron pressure and \( \gamma \simeq 3 \) is the adiabatic exponent [32]. In the above equations, we have neglected the ions inertia, as we are interested in the excitation of axions via electron plasma waves only.

**Axion-plasmon polaritons.** The effect of the axion field in the Langmuir waves can then be determined as follows: consider a strong, homogeneous magnetic field along the \( z \)-axis, \( \mathbf{B} = B_0 \mathbf{e}_z \). The electrostatic oscillations along the direction of the magnetic field will then provide the EM energy to excite the axion field, as can be seen from Eq. (5). The created axions will then feedback the plasma via the modified Maxwell equations. Mathematically, this effect can be formulated by assuming fluctuations around the plasma quasi-neutrality condition, \( n_e \sim n_0 + \tilde{n} \). Taking the decomposition into Fourier modes, \( (\tilde{n}, \tilde{\varphi}) \sim e^{i k z - i \omega t} \), and keeping linear terms only, we obtain
\[
\left( \omega^2 - \omega^2_p - S_e k^2 \right) \tilde{n} - i \frac{g e c}{m_e} B_0 k \tilde{\varphi} = 0,
\]
\[
\left( -\omega^2 + \frac{\tilde{m}_e^2 c^4 + e^2 k^2}{\hbar^2} \right) \tilde{\varphi} + i \frac{g e c^2}{c \hbar k} B_0 \tilde{n} = 0,
\]
where \( S_e = \sqrt{3 k_B T_e/m_e} \) is the plasma thermal speed, \( \omega_p = \sqrt{e^2 n_0/(\epsilon_0 m_e)} \) the plasma frequency and \( \tilde{m}_e = m_e + h g B_0 / \sqrt{c} \) the effective axion mass in the plasma. Nontrivial solutions to Eq. (7) implies the secular equation
\[
(\omega^2 - \omega^2_p) (\omega^2 - \omega^2_e - g^2 e^3 \omega_e^2/(m_e \epsilon_0)) = 0,
\]
where \( \omega^2_p = \omega_p^2 + S_e^2 k^2 \) and \( \omega^2_e = \tilde{m}_e^2 c^4 / \hbar^2 + e^2 k^2 \) are the Langmuir and axion bare dispersions, respectively, and \( \omega_e = e B_0 / m_e \) represents the cyclotron frequency. Solving Eq. (8) yields the lower (L) and upper (U) polariton modes
\[
\omega^2_{U,L} = \frac{1}{2} \left( \omega^2_p + \omega^2_p \pm \sqrt{\left( \omega^2_p - \omega^2_e \right)^2 + 4 \Omega^4} \right),
\]
where \( \Omega = (g^2 e^3 \omega_e^2/\epsilon_0 m_e)^{1/4} \) represents the Rabi frequency. The dispersion (9) describes the hybridization between the axions and the plasmons. If \( \Omega \) is larger than the decay rate \( \Gamma \) (to be specified below), a new quasiparticle is formed: the axion-plasmon polariton. As such, if axions are excited by the plasma in the presence of the external magnetic field, then the Langmuir
rotating-wave approximation (RWA) and inverse Fourier \( \omega \) onant with the plasma frequency, Since we are interested in the modes that are nearly res-
polariton operators \( \hat{a}_k \) and \( \hat{b}_k \), obeying the usual commutation relations \([\hat{c}_k, \hat{c}^\dagger_q] = \delta_{k,q} (\hat{c}_k = \{ \hat{a}_k, \hat{b}_k \})\), as

\[
\hat{n}(x) = \sum_k A e^{ikx} \left( \hat{a}_k + \hat{a}^\dagger_k \right), \quad \varphi(x) = \sum_k B e^{ikx} \hat{b}_k, \tag{12}
\]

where \( A = k/\sqrt{\omega_p n_0} \) and \( B = c/\sqrt{k \epsilon_0 \omega_p} \) are normalization constants. In terms of these operators, Eq. \( (11) \) can be recast as the Heisenberg equations \( i \dot{\hat{c}}_k = [\hat{c}_k, \hat{H}] \) associated to the Hamiltonian

\[
\hat{H} = \sum_k \omega_p \hat{a}_k \hat{a}^\dagger_k + \sum_k \omega_p \hat{b}_k \hat{b}^\dagger_k + \Omega \sum_k \hat{a}^\dagger_k \hat{b}_k + \text{h.c.} \tag{13}
\]

Full diagnozalization can be obtained by introducing the polariton operators \( \hat{A}_k = u_k \hat{a}_k - v_k \hat{b}_k \) and \( \hat{B}_k = v_k \hat{b}_k + u_k \hat{a}_k \), yielding

\[
\hat{H} = \sum_k \tilde{\omega}_L \hat{A}^\dagger_k \hat{A}_k + \sum_k \tilde{\omega}_U \hat{A}_k \hat{A}^\dagger_k, \tag{14}
\]

Quantization. In order to better understand the polariton character of the avoided crossing in Eq. \( (9) \), we proceed to a canonical quantization of the theory. We start by observing that Eq. \( (8) \) allows the following decomposition into fast and slow oscillations

\[
\begin{align*}
(\omega^2 - \omega_p^2) \hat{n} &= (\omega - \omega_p)(\omega + \omega_p) \hat{n} \\
(\omega^2 - \omega^2) \hat{\varphi} &= (\omega - \omega)(\omega + \omega) \hat{\varphi}.
\end{align*}
\tag{10}
\]

Since we are interested in the modes that are nearly resonant with the plasma frequency, \( \omega \simeq \omega_p \), we perform a rotating-wave approximation (RWA) and inverse Fourier transform Eq. \( (8) \) to obtain

\[
\begin{align*}
\left( i \frac{\partial}{\partial t} - \omega_p \right) \hat{n} + i \frac{g e c}{2 m e \omega_p} B_0 k \hat{\varphi} &= 0 \\
\left( i \frac{\partial}{\partial t} - \omega_p \right) \hat{\varphi} - i \frac{g e c}{2 e \omega_p \epsilon} B_0 \hat{n} &= 0.
\end{align*}
\tag{11}
\]

We can represent the plasmon and axion field in terms of bosonic operators \( \hat{a}_k \) and \( \hat{b}_k \), obeying the usual commutation relations \([\hat{c}_k, \hat{c}^\dagger_q] = \delta_{k,q} (\hat{c}_k = \{ \hat{a}_k, \hat{b}_k \})\), as

\[
\hat{n}(x) = \sum_k A e^{ikx} \left( \hat{a}_k + \hat{a}^\dagger_k \right), \quad \varphi(x) = \sum_k B e^{ikx} \hat{b}_k, \tag{12}
\]

where \( A = k/\sqrt{\omega_p n_0} \) and \( B = c/\sqrt{k \epsilon_0 \omega_p} \) are normalization constants. In terms of these operators, Eq. \( (11) \) can be recast as the Heisenberg equations \( i \dot{\hat{c}}_k = [\hat{c}_k, \hat{H}] \) associated to the Hamiltonian

\[
\hat{H} = \sum_k \omega_p \hat{a}_k \hat{a}^\dagger_k + \sum_k \omega_p \hat{b}_k \hat{b}^\dagger_k + \Omega \sum_k \hat{a}^\dagger_k \hat{b}_k + \text{h.c.} \tag{13}
\]

Full diagnozalization can be obtained by introducing the polariton operators \( \hat{A}_k = u_k \hat{a}_k - v_k \hat{b}_k \) and \( \hat{B}_k = v_k \hat{b}_k + u_k \hat{a}_k \), yielding

\[
\hat{H} = \sum_k \tilde{\omega}_L \hat{A}^\dagger_k \hat{A}_k + \sum_k \tilde{\omega}_U \hat{A}_k \hat{A}^\dagger_k, \tag{14}
\]

In order to obtain Eq. \( (14) \), we have neglected the counter-rotating terms \( \hat{a}^\dagger_k \hat{b}_k \) and \( \hat{a}_k \hat{b}^\dagger_k \), a procedure that is well justified provided the smallness of the coupling parameter \( g \). In fact, by retaining these terms, the diagonalization of Eq. \( (13) \) yields the full dispersion relation in \( (9) \). For sufficiently strong magnetic fields (e.g. at the interior of extremely dense neutron stars), the inclusion of the counter-rotating terms may be nevertheless necessary, a question that we will address later on.

Decay rate. The second-quantized formalism is particularly helpful to determine the polariton decay rate, a quantity that is crucial to quantify the strength the plasmon-axion coupling. The mode conversion depicted in Fig. 1 states that, near the crossing point \( k_s \), the polariton character abruptly changes from plasma-like to axion-like, and vice-versa. This “bending” in the dispersion relation is only observable if the decay rate \( \Gamma \) is smaller than the Rabi frequency \( \Omega \), i.e. provided the strong coupling condition holds. Two decay mechanisms are considered here. The first one is the radiative decay of the axion into two photons, which is given at the rate

\[
\Gamma = \frac{\tilde{\omega}_U \Gamma_0}{4} \left( 1 + \frac{\Omega^2}{4 \tilde{\omega}_U} \right),
\]
The second is the decay of axions into plasmons, which can be estimated with the help of Fermi’s Golden Rule,

\[ \Gamma_{\varphi \to \text{pl}} = \frac{2\pi}{\hbar} \sum_{k,q} |M_{k,q}|^2 \delta (\hbar \omega_\varphi - \hbar \omega_{\text{pl}}), \tag{17} \]

where \( M_{k,q} = \hbar \omega \sum_{p} \langle k| (\hat{a}_p^\dagger \hat{b}_p + \hat{a}_p \hat{b}_p^\dagger)|q\rangle \) is the transition amplitude between the states \(|k\rangle = \hat{a}_k^\dagger |0\rangle\) and \(|q\rangle = \hat{b}_q^\dagger |0\rangle\). Considering only transitions near the crossing point \((\omega_\varphi \approx \omega_{\text{pl}})\), and assuming the axions to be at much lower temperature than the plasmons, we obtain

\[ \Gamma_{\varphi \to \text{pl}} \approx \frac{\pi \Omega^2 (1 + n_a(\omega_p))^2}{\sqrt{\omega_p^2 + m_\varphi^2 c^2/\hbar^2}}. \tag{18} \]

leading to the formation of galaxies \([35, 36]\). In the instability region \(k < k_c\), the lower polariton mode is mostly axion-like, i.e. \(|\mathbf{v}_k| \gg |\mathbf{u}_k|\), and therefore this suggests that axions may be formed at the interior of neutron stars, as a consequence of the dynamical instability of the lower polariton mode. While a more rigorous estimate of the instability threshold would imply additional details about i) the renormalization of the electron mass and ii) the electron equation-of-state in extreme conditions, the present cold plasma estimates are reasonable at this stage given the actual uncertainty in the axion parameters \(g\) and \(m_\varphi\).

**Conclusion.** We have shown that the coupling between the axion and the plasmon in a strongly magnetized plasma may constitute a mechanism to observe the presence of axions. In particular, the most important feature is the avoided crossing between the Langmuir and the axionic Klein-Gordon dispersions near the resonance frequency. The consequence of such a mode repulsion is the changing in character of the Langmuir mode, which changes from electrostatic \((\omega \sim \omega_p)\) to an electromagnetic one \((\omega \sim ck)\). The strong coupling condition (i.e. the criterion for the axion-plasmon polariton to be a well-defined quasiparticle) has been calculated for different types of plasmas. Additionally, in extreme situations, such at the interior of neutron stars, the combination of strong magnetic fields and extremely high electronic densities is expected to lead to a dynamical instability in the lower polariton mode, which is mostly axionic in character. This suggests that axion-like particles may be produced as a consequence of an instability mechanism,
similarly to the Jeans instability in self-gravitating fluids leading for the formation of galaxies.

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[1] J. H. Christenson, J. W. Cronin, V. L. Fitch, and R. Turlay, Phys. Rev. Lett. 13, 138 (1964).
[2] A. Alavi-Harati, I. F. Albuquerque, T. Alexopoulos, M. Arenton, K. Arisaka, S. Averitte, A. R. Barker, L. Bellantoni, A. Bellavance, J. Belz, R. Ben-David, D. R. Bergman, E. Blucher, G. J. Bock, C. Bowin, S. Bright, E. Cheu, S. Childress, R. Coleman, M. D. Corcoran, G. Corti, B. Cox, M. B. Crisler, A. R. Erwin, R. Ford, A. Glazov, A. Golossanov, G. Graham, J. Graham, K. Hagan, E. Halkiadakis, K. Hanagaki, S. Hidaka, Y. B. Hsiung, V. Jejer, J. Jennings, D. A. Jensen, R. Kessler, H. G. E. Kobrak, J. LaDue, A. Lath, A. Ledoveskoy, P. L. McBride, A. P. McManus, P. Mikelsons, E. Monier, T. Nakaya, U. Naumenberg, K. S. Nelson, H. Nguyen, V. O’Dell, M. Pang, R. Pordes, V. Prasad, C. Qiao, B. Quinn, E. J. Ramberg, R. E. Ray, A. Roodman, M. Sadamoto, S. Schnetzer, K. Senyo, P. Shanahan, P. S. Shawhan, W. Slater, N. Solomey, S. V. Somalvar, R. L. Stone, I. Suzuki, E. C. Swallow, R. A. Swanson, S. A. Taeger, R. J. Tesarék, G. B. Thomson, P. A. Toale, A. Tripathi, R. Tschirhart, E. W. Wah, J. Wang, H. B. White, J. Whitmore, B. Winsten, R. Winston, J.-Y. Wu, T. Yamanaka, and E. D. Zimmerman (KTeV Collaboration), Phys. Rev. Lett. 83, 22 (1999).
[3] N. Cabibbo, Phys. Rev. Lett. 10, 531 (1963).
[4] M. Gell-Mann and M. Lévy, Il Nuovo Cimento (1955-1965) 16, 705 (1960).
[5] B. Pontecorvo, Sov. Phys. JETP 7, 172 (1958). [Zh. Eksp. Teor. Fiz.34,247(1957)].
[6] J. M. Pendlebury, S. Afach, N. J. Ayres, C. A. Baker, G. Ban, G. Bison, K. Bodek, M. Burghoff, P. Geltenbort, K. Green, W. C. Griffith, M. van der Grinten, Z. D. Grujic, P. G. Harris, V. Héline, P. Iaydjiev, S. N. Ivanov, M. Kasprzak, Y. Kermaidic, K. Kirch, H.-C. Koch, S. Komposch, A. Kozela, J. Krempel, B. Lauss, T. Lefort, Y. Lemière, D. J. R. May, M. Musgrave, O. Naviliat-Cuncic, F. M. Piegsa, G. Pignol, P. N. Prashanth, G. Quéméner, M. Rawlik, D. Rebreyend, J. D. Richard- son, D. Ries, S. Roccia, D. Rozpedzik, A. Schnabel, P. Schmidt-Wellenburg, N. Severijns, D. Shiers, J. A. Thorne, A. Weis, O. J. Winston, E. Wursten, J. Zejma, and G. Zsigmond, Phys. Rev. D 92, 092003 (2015).
[7] R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. 38, 1400 (1977).
[8] J. E. Kim and G. Carosi, Rev. Mod. Phys. 82, 557 (2010).
[9] S. Weinberg, Phys. Rev. Lett. 40, 223 (1978).
[10] P. Sikivie, Phys. Rev. Lett. 51, 1415 (1983).
[11] S. J. Asztalos, G. Carosi, C. Hagmann, D. Kinion, K. van Bibber, M. Hotz, L. J. Rosenberg, G. Rybka, J. Hoskins, J. Hwang, P. Sikivie, D. B. Tanner, R. Bradley, and J. Clarke, Phys. Rev. Lett. 104, 041301 (2010).
[12] K. Zioutas, S. Andriamonje, V. Arsov, S. Aune, D. Autiero, F. T. Avignone, K. Barth, A. Belov, B. Beltrán, H. Bräuninger, J. M. Carmona, S. Cebrián, E. Chesi, J. I. Collar, R. Creswick, T. Dafni, M. Davenport, L. Di Lella, C. Eleftheriadis, J. Englhauser, G. Fanourakis, H. Farach, E. Ferrer, H. Fischer, J. Franz, P. Friedrich, T. Geralis, I. Giomataris, S. Gninenko, N. Goloubev, M. D. Hasinoff, F. H. Hein- sius, D. H. H. Hoffmann, I. G. Irastorza, J. Jacoby, D. Kang, K. Königsmann, R. Kotthaus, M. Krčmar, K. Kousouris, M. Kuster, B. Lakić, C. Lasseur, A. Lliolios, A. Ljubičić, G. Lutz, G. Luzon, D. W. Miller, A. Morales, J. Morales, M. Mutterer, A. Nikolaidis, A. Ortiz, T. Pa- paeangelou, A. Placci, G. Raffelt, J. Ruz, H. Riege, M. L. Sarsa, I. Savvidis, W. Serber, P. Serpico, Y. Sernitzidis, L. Stewart, J. D. Vieira, J. Villar, L. Walkers, and K. Zachariadou (CAST Collaboration), Phys. Rev. Lett. 94, 121301 (2005).
[13] M. Fairbairn, T. Rashba, and S. Troitsky, Phys. Rev. Lett. 98, 201801 (2007).
[14] D. S. Akerib, S. Alsum, C. Aquino, H. M. Araújo, X. Bai, A. J. Bailey, J. Balajthy, P. Beltrame, E. P. Bernstein, A. Bernstein, T. P. Biesiadzinski, E. M. Boulton, P. Brä, D. Byram, S. B. Cahn, M. C. Carmona-Beníétèz, C. Chan, A. A. Chiller, C. Chiller, A. Curie, J. E. Cutter, T. J. R. Davison, A. Dobi, J. E. Y. Dokson, E. Druszkiewicz, B. N. Edwards, C. H. Faham, S. R. Fallon, S. Fiorucci, R. J. Gatskille, V. M. Gehman, C. Ghag, K. R. Gibson, M. G. D. Gilchriese, C. H. Hall, M. Hanhardt, S. J. Haselschwardt, S. A. Hertel, D. P. Hogan, M. Horn, D. Q. Huang, C. M. Ignarza, R. G. Jacobsen, W. Ji, K. Kamdin, K. Kazkaz, D. Khaitan, R. Knoche, N. A. Larsen, C. Lee, B. G. Lenardo, K. T. Lesko, A. Lindote, M. I. Lopes, A. Manalaysay, R. L. Mannino, M. F. Marzioni, D. N. McKinsey, D.-M. Mei, J. Mock, M. Moongweluwan, J. A. Morad, A. S. J. Murphy, C. Nehrkorn, H. N. Nelson, F. Neves, K. O’Sullivan, K. C. Oliver-Mallory, K. J. Palladino, E. K. Pease, L. Reichert, C. Rhynes, S. Shaw, T. A. Shutt, C. Silva, M. Solmaz, V. N. Solovov, P. Sorensen, S. Stephenson, T. J. Sumner, M. Szydagis, D. J. Taylor, W. C. Taylor, B. P. Tennyson, P. A. Terman, D. R. Tiedt, W. H. To, M. Tripathi, L. Vrtnikova, S. Uvarov, V. Velan, J. R. Verbus, R. C. Webb, J. T. White, T. J. Whitis, M. S. Withell, F. L. Wolfs, J. Xu, K. Yazdani, S. K. Young, and C. Zhang (LUX Collaboration), Phys. Rev. Lett. 118, 261301 (2017).
[15] E. Zavattini, G. Zavattini, G. Ruoso, E. Polacco, E. Milotti, M. Karuza, U. Gaspalti, G. Di Domenico, F. Della Valle, R. Cimino, S. Carusotto, G. Cantatore, F. Della Valle, R. Cimino, S. Carusotto, G. Cantatore, and M. Bregant (PVLAS Collaboration), Phys. Rev. Lett. 96, 110406 (2006).
[16] R. N. Mohapatra and S. Nasri, Phys. Rev. Lett. 98, 051802 (2007).
[17] J. Barranco, A. C. Monteverde, and D. Delpeine, Journal of Physics: Conference Series 485, 012035 (2014).
[18] M. Fairbairn, T. Rashba, and S. Troitsky, Phys. Rev. D 84, 125019 (2011).
[19] A. Di Piazza, C. Müller, K. Z. Hatsagortsyan, and C. H. Keitel, Rev. Mod. Phys. 84, 1177 (2012).
[20] M. A. Popovici, I. O. Mitu, G. Cta-Danil, F. Negoi, and C. Ivan, Journal of Radiological Protection 37, 176.
[21] D. A. Burton and A. Noble, Contemporary Physics 55, 110 (2014), https://doi.org/10.1080/00107514.2014.886840.

[22] S. Mangles, C. Murphy, Z. Najmudin, A. Thomas, J. Collier, A. Daniger, E. Divall, P. Foster, J. Gallacher, C. Hooker, D. Jaroszynski, A. Langley, W. Mori, P. Norreys, F. Tsung, R. Viskup, B. Walton, and K. Krushelnick, Nature 431, 535 (2004).

[23] C. G. R. Geddes, C. Toth, J. van Tilborg, E. Esarey, C. B. Schroeder, D. Bruhwiler, C. Nieter, J. Cary, and W. P. Leemans, Nature 431, 538 EP (2004).

[24] J. Faure, Y. Glinec, A. Pukhov, S. Kiselev, S. Gordienko, E. Lefebvre, J.-P. Rousseau, F. Burgy, and V. Malka, Nature 431, 541 EP (2004).

[25] J. T. Mendonça, EPL (Europhysics Letters) 79, 21001 (2007).

[26] D. A. Burton and A. Noble, “Plasma-based wakefield accelerators as sources of axion-like particles,” (2017), arXiv:1710.01906.

[27] D. A. Burton and A. Noble, Journal of Physics A: Mathematical and Theoretical 43, 075502 (2010).

[28] D. A. Burton, A. Noble, and T. J. Walton, Journal of Physics A: Mathematical and Theoretical 49, 385501 (2016).

[29] L. Visinelli, Modern Physics Letters A 28, 1350162 (2013).

[30] J. E. Kim, Phys. Rev. Lett. 43, 103 (1979).

[31] M. Dine, W. Fischler, and M. Srednicki, Physics Letters B 104, 199 (1981).

[32] F. F. Chen, Introduction to Plasma Physics (Springer, 2012).

[33] H. Bindslev, Plasma Physics and Controlled Fusion 35, 1615 (1993).

[34] J. K. Vogel, E. Armengaud, F. T. Avignone, M. Betz, P. Brax, P. Brun, G. Cantatore, J. M. Carmona, G. P. Carosi, F. Caspers, S. Caspi, S. A. Cetin, D. Chelouche, F. E. Christensen, A. Dael, T. Dafni, M. Davenport, A. V. Derbin, K. Desch, A. Diago, B. Döbrich, I. Dratchnev, A. Dudarev, C. Eleftheriadis, G. Fanourakis, E. Ferrer-Ribas, J. Galán, J. A. García, J. G. Garza, T. Geralis, B. Gimeno, I. Giomataris, S. Gninenko, H. Gómez, D. González-Díaz, E. Guendelman, C. J. Hailey, T. Hiramatsu, D. H. H. Hoffmann, D. Horns, F. J. Ignaz, I. G. Irastorza, J. Isern, K. Imai, A. C. Jakobsen, J. Jaeckel, K. Jakovic, J. Kaminski, M. Kawasaki, M. Karuza, M. Krcmar, K. Kousouris, C. Krieger, B. Lakic, O. Limousin, A. Lindner, A. Liolios, G. Luzón, S. Matsuki, V. N. Muratova, C. Nones, I. Ortega, T. Papaevangelou, M. J. Pivovaroff, G. Raffelt, J. Redondo, A. Ringwald, S. Russenschuck, J. Ruz, K. Saikawa, I. Savvidis, T. Sekiguchi, Y. K. Semertzidis, I. Shilon, P. Sikivie, H. Silva, H. ten Kate, A. Tomas, S. Troitsky, T. Vafeiadiad, K. van Bibber, P. Vedrine, J. A. Villar, L. Waltkiers, A. Weltman, W. Wester, S. C. Yildiz, and K. Zioutas, Physics Procedia 61, 193 (2015).

[35] Physical Processes in Interstellar Clouds (Nato Science Series C:) (Springer, 2011).

[36] S. Weinberg, Gravitation and Cosmology: Principles and Applications of the General Theory of Relativity (John Wiley & Sons, Inc., 1972).