Resonant Axion-Plasmon Conversion in Neutron Star Magnetospheres

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Resonant axion-plasmon conversion in the magnetospheres of magnetars may substantially impact the landscape of dark-matter axion detection. This Letter explores how resonant axion-plasmon conversion modify the expected radio signals from axion-photon conversions observed on Earth. Critically, the resonant conversion radius lies within the region expected for axion-photon conversion, introducing a nonradiative power loss that diminishes the anticipated photon flux. Our analysis demonstrates that this effect can reduce radio telescope sensitivities, shifting them into regions excluded by previous experiments. These findings compel a reassessment of experimental constraints derived from radio signatures of axion-photon conversions and highlight the necessity of accounting for plasmon effects in astrophysical axion searches. The presented corrections provide critical insights for refining the detection strategies of future telescope-based dark matter axion experiments.

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Introduction—The charge-parity (CP) violation remains one of the most fundamental, yet unresolved, issues in modern physics [1–5]. At the core of the so-called *strong CP problem* is the anomalous electric dipole moment of the neutron [6]. The Peccei-Quinn (PQ) mechanism provides a compelling resolution by promoting the CP-violating angle in the QCD Lagrangian to a dynamical field [7,8], with its associated pseudo-Goldstone boson—the axion—emerging from the spontaneous breaking of a global U(1) symmetry [9].

Axions and axionlike particles (ALPs) are well-motivated candidates for dark matter, owing to their extremely small mass and weak coupling to standard model particles [10–14]. Various experimental efforts, both terrestrial and astrophysical, have been dedicated to their detection [15–17]. However, the axion-photon coupling is weak, making their detection particularly challenging. Ground-based experiments such as PVLAS have initially generated controversial results [18], but have recently arguably set the best laboratory limits in the search for ALPs [19]; large-scale astrophysical searches, including CAST [20], ADMX [21–23], IAXO [24], and MADMAX [25], have sought axion signals in different mass ranges. These searches rely on passive detection, motivating complementary laboratory-based strategies such as high-power laser interactions [26], plasma wakefield acceleration [27–30], and direct axion production [31–37].

Astrophysical settings provide natural environments for axion-photon conversion, particularly in strong magnetic fields [38–50]. Axions propagating through the magnetospheres of neutron stars are known to resonantly convert into radio photons [51–55], a mechanism akin to the Mikheyev-Smirnov-Wolfenstein effect in neutrino physics. Recent works have highlighted the role of plasma effects in modifying the conversion dynamics [41,43], but these studies focus solely on axion-photon interactions, neglecting the role of plasmons. Standard axion electrodynamics describes photon conversion in vacuum or dilute media via the term $g\varphi F_{\mu\nu}\tilde{F}^{\mu\nu}$. In neutron star magnetospheres, dense plasma introduces collective plasmon modes, requiring modified electrodynamics with direct axion-plasmon coupling. This creates a distinct resonant conversion process at smaller radii than axion-photon conversion (see Fig. 1), reducing the effective photon-production volume and suppressing observable flux.

In this Letter, we show that axion-plasmon conversion crucially dims expected radio signals from neutron stars. Our quantitative framework shows typical conditions reduce radio search sensitivities significantly, potentially shifting viable parameter space into excluded regions. These results necessitate incorporating axion-plasmon interactions in astrophysical searches and refining future observational strategies.

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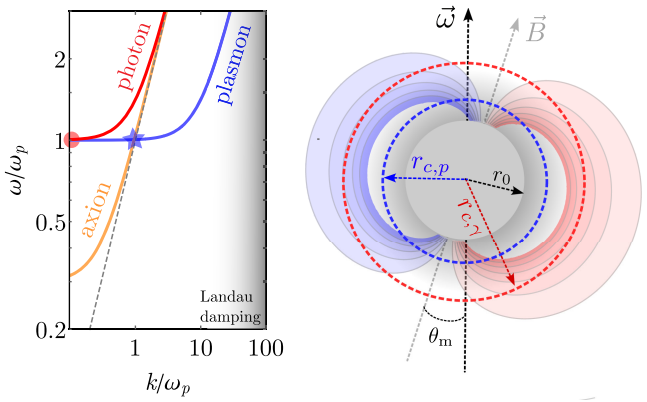


FIG. 1. Dispersion relation and mode conversion. Left panel: plasmon (solid blue line), photon (solid red), and axion (solid orange) lines, alongside the light cone $\omega = k$ (gray dashed line). The resonant, nonadiabatic axion-plasmon (adiabatic axion-photon) conversion point is marked by a star (circle). The Landau damped region $k \gtrsim \omega_p/S_e$ is marked for reference. Right panel: schematics of a pulsar, with a misalignment angle θ_m between the magnetic and rotation axis, displaying the resonant axion-plasmon conversion radius $r_{c,p}$ and the adiabatic axion-photon conversion radius $r_{c,\gamma}$. The effective conversion volume $4\pi/3(r_{c,\gamma}^3 - r_{c,p}^3)$ must be considered for the correct estimation of the photon flux density at Earth.

Axion-plasmon dispersion relation—The minimal electromagnetic theory accommodating the axion-photon coupling, while casting eventual QED effects in the plasma due to the strong magnetic fields of the magnetar, can be constructed as follows [53,56–58]:

$$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} - A_\mu J^\mu + \mathcal{L}_\varphi + \mathcal{L}_{\rm int} + \beta\left[\mathcal{F}^2 + \frac{7}{4}\mathcal{G}^2\right] + \sigma[(\partial_\mu F^{\mu\nu})(\partial_\alpha F^\alpha_{\ \nu}) - F_{\mu\nu}\Box F^{\mu\nu}], \quad (1)$$

where $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ is the electromagnetic (EM) tensor, $\Box = \partial_t^2 - \nabla^2$ is the d'Alembert operator, $\mathcal{F} = F_{\mu\nu}F^{\mu\nu}$, and $\mathcal{G} = F_{\mu\nu}\tilde{F}^{\mu\nu}$, with $\tilde{F}^{\mu\nu} = \epsilon^{\mu\nu\alpha\beta}F_{\alpha\beta}$ denoting the dual EM tensor, and J^μ is the four current. The last two terms in Eq. (1) describe the Euler-Heisenberg (one-loop) QED corrections that are relevant for the strong magnetic fields of the magnetars. Here, $\beta = \alpha^2/(90m_e^4)$, with $\alpha = e^2/(4\pi) \simeq 1/137$ being the fine structure constant, and $\sigma = (2/15)\alpha/m_e^2$. Moreover, $\mathcal{L}_\varphi = \partial_\mu\varphi^*\partial^\mu\varphi - m_\varphi^2|\varphi|^2$ is the axion Lagrangian (with φ denoting the axion field). For the QCD axion, $m_\varphi = \sqrt{z}f_\pi m_\pi/f_\varphi$, where $z = m_u/m_d$ is the up and down mass ratio, and $f_{\varphi(\pi)}$ is the axion (pion) decay constant [7,9]. Upon integration of the anomalous axion-gluon triangle, one obtains $\mathcal{L}_{\rm int} = -(g/4)\varphi F_{\mu\nu}\tilde{F}^{\mu\nu}$, where g is the axion-photon coupling. Although motivated for the QCD axion, the model in Eq. (1) is valid for any ALP. From Euler-Lagrange equations, one obtains the inhomogeneous Maxwell's equations [58–60] as

$$(1 + 2\sigma\Box)\partial_\mu F^{\mu\nu} = J^\nu + 2\beta\partial_\mu\left[\mathcal{F}F^{\mu\nu} + \frac{7}{4}\mathcal{G}\tilde{F}^{\mu\nu}\right] + g\partial_\mu\left(\varphi\tilde{F}^{\mu\nu}\right). \quad (2)$$

In the context of the axion-plasmon polariton, the only relevant equation is the Poisson equation ($\nu = 0$), which reads

$$\nabla\cdot(\boldsymbol{\mathcal{E}} + g\varphi\mathbf{B}) = \rho, \quad (3)$$

where the effective electric reads

$$\boldsymbol{\mathcal{E}} = [1 + 2\sigma\Box + 4\beta(E^2 - B^2) + 7\beta(\mathbf{E}\cdot\mathbf{B})\mathbf{B}]\mathbf{E}. \quad (4)$$

Moreover, variation with respect to φ yields the Klein-Gordon equation describing the axion field

$$(\Box + m_\varphi^2)\varphi = g\mathbf{E}\cdot\mathbf{B}. \quad (5)$$

The plasma density is $\rho = e(n_i - n_e)$, where n_i and n_e, respectively, represent the ion and electron densities. As we are interested in electron plasma waves only, we can assume the ions to be immobile. Thus, the equations governing the plasma electrons read

$$\frac{\partial n_e}{\partial t} + \nabla\cdot(n_e\mathbf{u}_e) = 0, \quad (6)$$

$$\left(\frac{\partial}{\partial t} + \mathbf{u}_e\cdot\nabla\right)\mathbf{u}_e = -\frac{e}{\gamma_e m_e}(\mathbf{E} + \mathbf{u}\times\mathbf{B}) - \frac{\nabla P_e}{n_e m_e}, \quad (7)$$

where $\gamma_e = (1 - u_e^2)^{-1/2}$ is the Lorentz factor. In the following, and without any loss of generality, we will consider the plasma electrons to be nonrelativistic ($\gamma_e \simeq 1$). At first, we are interested in describing the electrostatic perturbations along a static, homogeneous magnetic field $\mathbf{B} = B_0\mathbf{e}_z$. We retain first-order terms in the perturbations only, taking $n_e = n_0 + \tilde{n}_e$ and $\varphi = \tilde{\varphi}$ (neglecting the presence of a vev, $\varphi_0 = 0$). Further, we assume the compressions in the electron fluid to be adiabatic, $\nabla\tilde{P}_e \simeq \gamma T_e\nabla\tilde{n}_e$, with γ being the adiabatic index. Putting Eqs. (3)–(7) together, we obtain the following eigenvalue problem:

$$\left(\omega^2 - \frac{\omega_p^2}{\mathcal{D}(\omega,k)} + S_e^2k^2\right)\tilde{n}_e - \frac{gB_0\omega_p}{\mathcal{D}(\omega,k)}\tilde{\psi} = 0, \quad (8)$$

$$\left(\omega^2 - k^2 + m_\varphi^2 + \frac{g^2B_0^2}{\mathcal{D}(\omega,k)}\right)\tilde{\psi} - gB_0\omega_p\tilde{n}_e = 0, \quad (9)$$

where $\tilde{\psi} = ikm_e\tilde{\varphi}/e$ is an auxiliary field, $\omega_p = \sqrt{e^2n_0/m_e}$ is the plasma frequency, and $S_e = \sqrt{\gamma T_e/m_e}$ is the electron

thermal speed. The quantity

$$\mathcal{D}(\omega,k)=1+3\beta B_0^2-2\sigma(\omega^2-k^2) \tag{10}$$

accounts for the QED effects, which relevance strongly depends on the energy scale at which the axion-plasmon mixing takes place. Considering resonant processes only ($\omega\simeq\omega_p$), relativistic axions may be produced by plasmons if $m_\varphi\ll\omega_p$, making resonance possible near the light cone $\omega_p\simeq k$, for which $\mathcal{D}(\omega,k)\simeq 1+3\beta B_0^2$. On the other hand, plasmons resonantly convert into nonrelativistic axions provided the condition $m_\varphi\simeq\omega_p$, and thus $\mathcal{D}(\omega,k)\simeq 1+3\beta B_0^2-2\sigma\omega_p^2$. For the typical conditions found at the polar caps of magnetars, $B_0\sim 10^9$ T and $\omega_p\sim 10^{-3}$ eV, which leads $\mathcal{D}(\omega,k)\simeq 1$ for both cases, allowing us to safely neglect QED effects in the resonant plasmon-axion conversion. Although instructive, this discussion could be avoided by noticing that the dispersive QED term $\sim\sigma$ is only relevant at the energy scales $k\sim m_e$, which lie well above the plasma frequency ($\omega_p\ll m_e$). In Supplemental Material [61], we have performed a complementary analysis to persuade the reader that Euler-Hesenberg corrections to the plasma frequency are indeed negligible.

At this point, it is also important to rule out corrections in the plasma dynamics due to pair production, as it may take place in the conditions of magnetars [68–70]. While pair creation can occur in magnetar magnetospheres, its impact on axion-plasmon conversion is negligible because (i) in conversion regions the background plasma densities $n_0\sim 10^{16}$–10^{19} cm^{-3} dominate over pair densities $n_{e^\pm}\sim 10^{10}$–10^{15} cm^{-3}, making $\delta\omega_p/\omega_p<10^{-3}$ [68,71]; (ii) Debye screening localizes pair effects to microscopic scales; and (iii) longitudinal modes ($\mathbf{k}\|\mathbf{B}$) decouple from pair-creation dynamics [72].

In the absence of QED effects, $\mathcal{D}(\omega,k)\sim 1$, Eq. (9) can be recast in a more compact form as $\boldsymbol{\nabla}\cdot[\epsilon(k,\omega)\mathbf{E}]=0$, where

$$\epsilon(k,\omega)=1-\frac{\omega_e^2}{\omega^2}-\frac{\Omega^4}{\omega^2(\omega^2-\omega_\varphi^2)} \tag{11}$$

is the dielectric permittivity, $\omega_e^2=\omega_p^2+S_e^2k^2$ is the plasmon dispersion, $\omega_\varphi^2=M_\varphi^2+k^2$ is the axion dispersion, and $M_\varphi=\sqrt{m_\varphi^2+g^2B_0^2}$ is the axion effective mass in the plasma. Here, $\Omega=\sqrt{gB_0\omega_p}$ is the axion-plasmon coupling parameter (Rabi frequency), which reads

$$\Omega\simeq 2\pi\times 1.2\sqrt{\frac{g}{10^{-12}\ \mathrm{GeV}^{-1}}\frac{B_0}{10^{10}\ \mathrm{T}}\frac{\omega_p}{\mathrm{GHz}}}\mathrm{MHz}. \tag{12}$$

Setting Eq. (11) to zero, we obtain the lower (L) and (U) polariton modes

$$\omega_{\mathrm{L(U)}}^2=\frac{1}{2}\left[\omega_\varphi^2+\omega_e^2\pm\sqrt{(\omega_e^2-\omega_\varphi^2)^2+4\Omega^4}\right]. \tag{13}$$

At this point, we emphasize an important distinction: *plasmons* are purely longitudinal electron oscillations in the magnetosphere, while axion-plasmon *polaritons* are hybrid states formed through resonant coupling between axions and plasmons [59,73]. As shown in Eq. (13), these polaritons exhibit an avoided crossing (energy gap Ω) in their dispersion. This distinctive feature enables resonant conversion at smaller radii than standard axion-photon processes [41], as depicted in Fig. 1.

In the vicinity of the resonant mode k_c set by the condition $\omega_\varphi=\omega_e\simeq\omega_p$ (the point marked with a star in Fig. 1), Eq. (11) can be approximated as

$$(\omega-\omega_\varphi)(\omega-\omega_e)\simeq\frac{\Omega^4}{4\omega_p^2}=\frac{g^2B_0^2}{4}. \tag{14}$$

For the subsequent analysis, it is crucial to note that in neutron star (NS) magnetospheres, both the magnetic field and electron density are inhomogeneous, following Goldreich-Julian (GJ) profiles [41,43,75]. The radial component of the magnetic field reads

$$B(r)=B_0\left(\frac{r_0}{r}\right)^3 f,\quad \text{with}$$
$$f\equiv f(\theta,\theta_\mathrm{m},\psi)=\cos\theta_\mathrm{m}\cos\theta+\sin\theta_\mathrm{m}\sin\theta\cos\psi. \tag{15}$$

Here, r_0 is the NS radius, B_0 is the magnetic field at the pole caps, θ_m is the misalignment angle between $\mathbf{B}$ and the NS angular velocity $\boldsymbol{\omega}_\mathrm{NS}$, $\psi=\phi-\omega_\mathrm{NS}t$, and θ and ϕ are the usual polar and azimuthal angles. The radial electron density is then given by $n(r)=n_0(r_0/r)^3$, with $n_0=2\pi B_0/(e\tau)$ being the electronic density at the NS surface and $\tau=2\pi/\omega_\mathrm{NS}$ being the rotation period. Since these variations are smooth at the scale of $1/k_c$, we can assume that, outside the resonance points, the two independent modes are given by $\omega=\omega_\varphi(r,k)$ and $\omega=\omega_e(r,k)$. At the resonance point $r=r_c$, we have $\omega_\varphi(r_c,k_c)=\omega_e(r_c,k_c)$. To linearize around this point, we can use the prescription

$$r=r_c+\xi,\qquad k=k_c+\delta. \tag{16}$$

Upon substitution in Eq. (14), and using the conversion condition, we obtain [76,77]

$$(a_\varphi\delta+b_\varphi\xi)(a_e\delta+b_e\xi)\simeq\frac{g^2B_c^2}{4}, \tag{17}$$

where $B_c=B(r=r_c)$, and $a_\alpha=\partial_k\omega_\alpha|_{k=k_c}$ and $b_\alpha=\partial_r\omega_\alpha|_{r=r_c}$. Since $\{\xi,\delta\}$ form a canonical pair, $\delta=-i\partial_\xi$, and noticing that the wave amplitude describing an incoming axion ($\xi<0$) and a outgoing plasmon ($\xi>0$) varies as

$\Phi(r) \sim \exp(ik_c r)$, one may show that Eq. (17) can be reduced to the Weber equation (see Supplemental Material for details [61])

$$\frac{\partial^2 \Psi}{\partial \zeta^2} + \left(\nu + \frac{1}{2} - \frac{\zeta^2}{4}\right)\Psi = 0, \tag{18}$$

where $\nu = ig^2 B_c^2/4(a_\varphi b_e)$ is a parameter controlling the axion-plasmon conversion amplitude (see below), $\zeta = \sqrt{(a_\varphi b_e - a_e b_\varphi)/(a_\varphi a_e)}\exp(i3\pi/4)\xi$ is a rescaled radius, and $\Psi(\xi)$ is a rotated wave

$$\Psi(\xi) = \exp(iK\xi^2/2)\Phi(\xi), \qquad K = -\frac{a_\varphi b_e + a_e b_\varphi}{2a_\varphi a_e}. \tag{19}$$

The formal solutions to Eq. (18) are given in terms of the parabolic cylindrical function $D_\nu(\zeta)$ [78]. For $\zeta \lesssim 0$, Eq. (18) describes an incoming axion, for which the solution can be given approximated as [61]

$$\Psi_<(\zeta) \sim \zeta^\nu \exp(-\zeta^2/4), \tag{20}$$

and $\Phi_<(\xi)$ can be obtained from the transformation relation in Eq. (19). Conversely, in the outward region, $\zeta \gtrsim 0$, we may approximate the solution as

$$\Psi_>(\zeta) \sim \zeta^\nu \exp(-\zeta^2/4) - \frac{\sqrt{2\pi}}{(-\nu-1)!}\frac{\exp(i\nu\pi + \zeta^2/4)}{\zeta^{(\nu+1)}}, \tag{21}$$

and a similar expression for $\Phi_>(\xi)$. The first term of Eq. (21) represents the axion mode, while the second represents the plasmon mode. As such, the axion transmission coefficient at the conversion radius r_c is given by the ratio of the axion amplitudes at $r < r_c$ ($\xi < 0$) and $r > r_c$ ($\xi > 0$), $|\exp(i\nu\pi)|$. The square of this quantity gives the axion transmission coefficient, which more explicitly reads $T_{\rm axion} = \exp[-2\pi g^2 B_c^2/(4a_\varphi b_e)]$. From energy conversion arguments, the *resonant* axion-plasmon conversion coefficient, our quantity of interest, is therefore given by

$$P_{\rm conv} = 1 - T_{\rm axion} = 1 - \exp\left(\frac{\pi g^2 B_c^2}{2a_\varphi b_e}\right). \tag{22}$$

A similar formula for axion-photon conversion based on the Landau-Zenner theory has been obtained in Ref. [79]. For the conditions of a neutron star magnetosphere, the GJ profiles in Eq. (15) yield

$$a_\varphi = \frac{k_c}{\omega_\varphi(r_c, k_c)} \simeq \left(1 - \frac{m_\varphi^2 + g^2 B_0^2 \frac{r_0^6}{r_c^6} f^2}{\omega_p^2}\right)^{1/2}, \tag{23}$$

$$b_e = \frac{\partial \omega_e}{\partial r} \simeq -\frac{3}{2}\omega_p \frac{r_0^{3/2} f^{1/2}}{r_c^{5/2}}, \tag{24}$$

where we have used $k_c \simeq \sqrt{\omega_p^2 - m_\varphi^2 - g^2 B_c^2}$. The maximum probability for the plasmon-axion conversion occurs at the radius $r_{c,p}$ such that $P_{\rm conv} = 1$ and, therefore,

$$r_{c,p} \simeq r_0 \left(\frac{4\omega_p^2 - g^2 B_0^2}{4m_\varphi^2}\right)^{1/6} f^{1/3}. \tag{25}$$

This should be compared with the conversion radius $r_{c,\gamma}$ where an axion decays into a O-mode ($\mathbf{E}\|\mathbf{B}$) photon [41], $r_{c,\gamma} \simeq r_0(\omega_p/m_\varphi)^{1/3} f^{1/3}$. As we show below, this may drastically impact the estimated radiation power due to dark-matter axions converting into photons. The radiated power per solid angle for adiabatic (in opposition to resonant) photon conversion is given by [41]

$$\frac{d\mathcal{P}_\gamma}{d\Omega} \simeq \frac{1}{2}\frac{g^2 B_{c,\gamma}^2}{m_\varphi}\rho_{c,\gamma}^{(\rm DM)} V_{c,\gamma}, \tag{26}$$

with $V_c = (4/3)\pi r_{c,\gamma}^3$ being the resonant conversion volume and $\rho_{c,\gamma}^{(\rm DM)}$ being the dark-matter density calculated at the conversion radius $r_{c,\gamma}$. To this quantity, one must subtract the volume leading to resonant axion-plasmon conversion, taking place at the radius $r_{c,p} < r_{c,\gamma}$, which consists in a nonradiative power loss. Provided the relations imparted by Eq. (25), the corrected radiated power reads

$$\frac{d\mathcal{P}_\gamma^{(\rm corr)}}{d\Omega} \simeq \frac{d\mathcal{P}_\gamma}{d\Omega}\left[1 - \left(1 - \frac{g^2 B_0^2}{4\omega_p^2}\right)^{1/2}\right]. \tag{27}$$

The latter, in turn, corrects the telescope's flux density reaching Earth,

$$S_\gamma = \frac{1}{d^2 \Delta\nu}\frac{d\mathcal{P}_\gamma^{(\rm corr)}}{d\Omega}, \tag{28}$$

with d being the distance from Earth to the magnetar and $\Delta\nu$ as the bandwidth. Considering that the axion dark-matter velocity spread v_0 determines the broadening in the emitted photon frequencies, the bandwidth can be estimated as $\Delta\nu \sim v_0^2 m_\varphi/2\pi$, following the energy conservation arguments worked out in [41]. To project the sensitivity of the radio signals, we have to compare S_γ against the minimal flux to which telescopes are responsive,

$$S_{\gamma,\rm min} = {\rm SNR}_{\rm min}\frac{\rm SEFD}{\sqrt{n_{\rm pol}\Delta\nu\tau_{\rm obs}}}, \tag{29}$$

where $\rm SNR_{min}$ is the minimal signal-to-noise ratio, SEFD is the system-equivalent flux density, $\tau_{\rm obs}$ is the observation

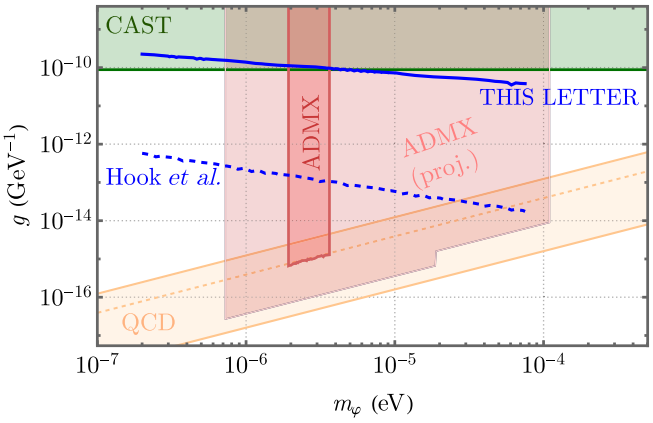


FIG. 2. Impact of axion-plasmon resonance on the sensitivity of photon signals emitted from neutron stars. The dashed line corresponds to project sensitivity for radio signals observations due to axion-photon conversion in the neutron star SGR J1745-2900 (spike configuration, $\theta = 90°$, $\theta_m = 10°$, $\tau = 100$ h), as calculated by Hook *et al.* [41]. Our estimates for the dimmed sensitivity, as a consequence of the radiated power reduction due to resonant axion-plasmon conversion in Eq. (27), is depicted by the solid line.

time, and $n_{\rm pol}$ is the number of photon polarizations (for numerical evaluations, we choose $n_{\rm pol} = 2$). In Fig. 2, we illustrate the correction to the sensitivity to DM axion-photon conversion in the neutron star SGR J1745-2900, located at a distance $d = 8.5$ kpc from Earth and spinning with a period $P \simeq 3.76$ s [41]. Moreover, we assume a DM velocity spreading of $v_0 \sim 10^{-3}$, and a polar magnetic field $B_0 \simeq 1.6 \times 10^{14}$ G [80]. Considering a typical plasma frequency about $\omega_p = 10^{-3}$ eV, an integration time of $\tau = 100$ h and an angle $\theta = 90°$, we project that the sensitivity can be drastically reduced, eventually falling within previously excluded regions. We believe that this effect may significantly alter the landscape of experiments based on radio telescope searches in the years to come.

Naturally, our estimates are prone to uncertainties in the conditions of the atmosphere of magnetars. While our conversion radius estimates use conservative Goldreich-Julian density profiles and dipole magnetic fields, we have quantified the impact of potential perturbations: (i) current sheet density fluctuations ($\delta n_e/n_e \sim 10$–100 [70]) could modify $r_{c,p}$ by $\lesssim 17\%$, but these are confined to equatorial regions and average out over the conversion volume; and (ii) magnetic multipoles ($B_{\rm local}/B_0 \leq 10^3$ [69]) introduce corrections $\delta r_{c,p}/r_{c,p} \lesssim 1\%$ due to the $1/6$ power-law dependence in Eq. (25). Crucially, both effects preserve the hierarchy $r_{c,p} < r_{c,\gamma}$ and remain negligible compared to the global plasma profile that governs axion-plasmon conversion. Finally, unlike the case of plasmons considered here, important corrections to the photon mass due to strong magnetic fields may be in order [81], eventually bringing corrections up to a factor of ~ 2 in the photon mass.

Conclusion—The resonant conversion of axions into plasmons in magnetar magnetospheres introduces a significant correction to the expected photon signals detectable on Earth. Our analysis shows that the radius at which axion-plasmon conversion occurs is smaller than the axion-photon conversion radius, effectively reducing the volume where detectable photons can emerge. This discrepancy, proportional to the difference in the respective conversion volumes, necessitates a reevaluation of experimental expectations based solely on axion-photon conversion mechanisms. The results presented here underscore the importance of accounting for axion-plasmon interactions in strongly magnetized environments, as they could hinder the photon flux reaching Earth and subsequently affect the interpretability of experimental data searching for axion signals. Future observational strategies must incorporate these findings to accurately constrain axion properties and refine detection techniques.

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Data availability—The data that support the findings of this Letter are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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