Axion production in unstable magnetized plasmas: an active source of dark-matter

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Axions, the hypothetical particles restoring the charge-parity symmetry in the strong sector of the Standard Model, and one of the most prone candidates for dark matter, are well-known to interact with plasmas. In a recent publication [Phys. Rev. Lett. 120, 181803 (2018)], we have shown that if the plasma dynamically responds to the presence of axions, then a new quasi-particle (the axion plasmon-polariton) can be formed, being at the basis of a new generation of plasma-based detection techniques. In this work, we exploit the axion-plasmon hybridization to actively produce axions in streaming magnetized plasmas. We show that, if we make the plasma unstable via the injection of an energetic electron beam (beam-plasma instability), an appreciable production rate of few axions per minute can be achieved. The produced axions can then be detected by Primakoff decay into photons.

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Introduction. The charge-parity (CP) violation is one of the most fundamental, yet unsolved problems of modern physics [1, 2]. Although CP violation is, by construction, inherent to the Standard Model [3–5], 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]. To solve the problem, the celebrated Peccei-Quinn (PQ) mechanism promotes the CP angle in the QCD Lagrangian to a complex field [7, 8]. The associated particle, the axion, then emerges after breaking the U(1) symmetry in the gluon coupling term [9].

Axions and axion-like particles (ALPs) are predicted with a extremely small mass (possibly in the meV range) and couple very weakly to ordinary matter. For that reason, ALPs became appealing candidates (arguably, the most well theoretically motivated) to fix the dark matter puzzle [10, 11]. Many facilities have been built with the goal of observing axion or ALP signatures, both based on laboratory and astrophysical observations [12–14]. However, given the smallness of the axion-photon coupling, testing the axion is notoriously difficult, rendering most of the experimental observations inconclusive. For example, the PVLAS experiment - originally design to probe the birefringent properties of the electromagnetic vacuum [15] - advanced preliminary results indicating the existence of axions back in 2008, and since then those findings are object of controversial debates (see e.g. [8, 16–18] and references therein). Telescope experiments, such as CAST [19], probing a vast range of axion masses, and ADMX [20–22], IAXO [23] and MADMAX [24], investigating more precise regions of the QCD axion parameter space, are designed to probe axions produced by astrophysical objects. By construction, these experiments rely on a passive approach, in the sense that no axion production is envisaged. It is therefore desirable to look for alternatives, where axions could be actively produced in the lab. This ambitious task might be accessible in the next generation of experiments based on high-power lasers, which are designed to provide conditions to probe QED physics in parameter regimes that are inaccessible to particle colliders [25]. The ELI experiment, for example, will offer the possibility to study the Heisenberg-Euler vacuum (virtual electron-positon pairs) [25, 26] and the quantum recoil due to radiation emission [27]. The wakefield acceleration paradigm gained much breath as it reveals to be an efficient way to accelerate particles [28–30]. Recent results of the AWAKE experiment [31] show that electrons can be boosted up to 2 GeV with the help of a plasma wakefield. Interestingly, recent theoretical studies have pointed out that such wakefields could
ultimately be used to produce ALPs in the lab [32–34]. Along these lines, a scheme based on petaWatt lasers was proposed by one of us a few years ago [35]. An inherent limitation in this scheme, however, stems in the fact that the such high powers are only achieved within the Rayleigh length of the beam (∼ few micrometers), corresponding to an extremely small axion-photon conversion length (and, consequently, to a vanishing axion-photon conversion probability).

In this Letter, we show that axions can be actively produced in an unstable magnetized plasma. If an energetic electron beam is injected in the plasma, unstable oscillations of the plasma electron waves, or plasmons, are produced. This effect is dubbed in the literature as the beam-plasma instability [36, 37]. The growing plasmon perturbation then provide the energy to the growth of axions. A schematic representation of the process is depicted in Fig. 1. In the absence of axions, the plasmons are insensitive to the magnetic field; however, if axions exist, then they admix the plasmons in a coherent fashion, leading to the formation of a hybrid quasi-particle, the axion-plasmon polariton [38]. As such, if the plasmons become dynamically unstable, their small axion component will also grow, leading to an efficient axion production rate of ∼ 1 photon per unit volume can be attained. The produced photons can then be resonantly converted into a transverse photons, escaping the plasma at a rate of ∼ 1 photon per day. These photons can then be probed with the help of single-photon microwave detectors. Some implications in the radio signals emitted by pulsars are also discussed.

Beam-plasma instability in magnetized plasmas. The minimal electromagnetic theory accommodating the axion-photon coupling can be constructed as follows (ℏ = c = 1) [39, 40]

\[
\mathcal{L} = -\frac{1}{4} F_{\mu \nu} F^{\mu \nu} - A_\mu J_\mu^e + \mathcal{L}_\varphi + \mathcal{L}_{\text{int}},
\]

where \( F_{\mu \nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \) is the electromagnetic (EM) tensor, \( J_\mu^e \) is the electron four-current, and \( \mathcal{L}_\varphi = \partial_\mu \varphi^\ast \partial^\mu \varphi - m_\varphi^2 |\varphi|^2 \) is the axion Lagrangian (with \( \varphi \) denoting the axion field). For the QCD axion, \( m_\varphi = \sqrt{2 f_\pi m_{u} / f_\varphi} \), where \( z = m_u / m_{u} \) is the up/down mass ratio, and \( f_\varphi(\pi) \) is the axion (pion) decay constant [7, 9]. Upon integration of the anomalous of the axion-gluon triangle, one obtains \( \mathcal{L}_{\text{int}} = - (g/4) F_{\mu \nu} F^{\mu \nu} + \epsilon^{\mu \nu \alpha \beta} F_{\mu \alpha} F_{\nu \beta} \), where \( F_{\mu \nu} = \epsilon^{\mu \nu \alpha \beta} F_{\alpha \beta} \) denotes the dual EM tensor and \( g \) is the axion-photon coupling. Although motivated for the QCD axion, the model in Eq. (1) is valid for any axion-like particle (ALP). From Euler-Lagrange equations, one obtains the Maxwell’s equations [38], in particular the Poisson equation

\[
\nabla \cdot (\mathbf{E} + g \varphi \mathbf{B}) = \frac{\rho}{\varepsilon_0},
\]

and the Klein-Gordon equation describing the axion field

\[
(\Box + m_\varphi^2) \varphi = g \mathbf{E} \cdot \mathbf{B},
\]

with \( \Box = \partial_t^2 - \nabla^2 \) denoting the d’Alembert operator. In the situation of an electron beam propagating inside the plasma, \( \rho = e (n_e - n_i - n_b) \), where \( n_i, n_e \) and \( n_b \) respectively represent the ion, electron and beam densities. As we are interested in electron plasma waves only, we can assume the ions to be immobile. Thus, the equations governing the dynamics of the plasma and beam electrons are given by

\[
\frac{\partial n_e}{\partial t} + \nabla \cdot (n_e \mathbf{u}_e) = 0,
\]

with \( \alpha = \{e, b\} \), and

\[
\left( \frac{\partial}{\partial t} + \mathbf{u}_\alpha \cdot \nabla \right) \mathbf{u}_\alpha = - \frac{e}{\gamma_\alpha m_e} (\mathbf{E} + \mathbf{u} \times \mathbf{B}),
\]

where \( \gamma_\alpha = (1 - u_\alpha^2)^{-1/2} \) is the Lorentz factor. In the following, we will consider the plasma electrons to be initially at rest (\( \gamma_e \approx 1 \)), while the beam electrons propagate with velocity \( \mathbf{u}_b \). We are interested in describing the electrostatic perturbations along a static, homogeneous magnetic field \( \mathbf{B} = B_0 \mathbf{e}_z \). As such, owing to the quasi-neutrality condition of the plasma, we perturb the densities as \( n_e = n_0 + ˜n_e \) and \( n_b = f n_0 + ˜n_b \) (here, \( f \) is the fraction of the electrons in the beam), and the axion field as \( \varphi = ˜\varphi \) (neglecting the presence of a vev, \( \varphi_0 = 0 \)) to obtain

\[
\frac{\partial^2}{\partial z^2} ˜\varphi + \frac{e n_0 \partial E}{m_e \gamma_e} = 0,
\]

\[
\left( \frac{\partial}{\partial t} + u_0 \frac{\partial}{\partial z} \right) ˜n_e + \frac{e n_0 \partial E}{\gamma_0 m_e} = 0,
\]

\[
\left( \frac{\partial^2}{\partial t^2} - \nabla^2 + M_\varphi^2 \right) ˜\varphi - g B_0 \frac{\partial E}{\partial z} = 0,
\]

where \( ˜\varphi = \partial_\varphi / \partial z \). After Fourier transforming, this allows us to write Eq. (2) as \( \partial_z (\epsilon(k, \omega) E) = 0 \), where

\[
\epsilon(k, \omega) = 1 - \frac{\omega_p^2}{\omega^2} - \frac{f}{\Omega^4} \left( \frac{\omega_p^2}{\omega^2} - \frac{\omega_\gamma^2}{\omega_\varphi^2} \right) - \frac{f}{\Omega^4} \left( \frac{\omega_\gamma^2}{(\omega - k u_0)^2} - \frac{\omega_\varphi^2}{(\omega - k u_0)^2} \right)
\]

is the dielectric permittivity, \( \omega_p = \sqrt{e^2 n_0/(\epsilon_0 m_e)} \) is the plasma frequency, and \( \omega_\varphi^2 = M_\varphi^2 + k^2 \), with \( M_\varphi = \sqrt{m_\varphi^2 + g^2 B_0^2} \) being the axion effective mass in the plasma. Here,

\[
\Omega = \sqrt{g B_0 \omega_p} \sim 1 \text{ (1.2 Hz)} \sqrt{g \times 10^{13} B_0 \omega_p \text{ GeV}^{-1} \text{ T GHz}}
\]
is the axion-plasmon coupling parameter (Rabi frequency). In the absence of the beam \( f = 0 \), Eq. (7) yields the lower (L) and (U) polariton modes which, within the rotating-wave approximation (RAW) valid for \( \Omega \ll \omega_p \), read [38]

\[
\omega_{L(U)} = \frac{1}{2} \left( \omega_p + \omega_p \mp \sqrt{(\omega_p - \omega_p)^2 + 4\Omega^2} \right)
\]

The Hopfield coefficients \( u_k \) and \( v_k \), that give the corresponding axion or plasmon fraction of a certain hybridized mode \( k \), satisfy the relation \( |u_k|^2 + |v_k|^2 = 1 \). For example, for the U− mode, the axion fraction is given by

\[
u_k = \frac{\omega_{U(p)} - \omega_{L(p)}(\omega_p + \omega_p) \sqrt{(\omega_p - \omega_p)^2 + 4\Omega^2}}{\omega_p + \omega_p}
\]

Conversely, in the absence of axions, Eq. (7) describes the celebrated two-stream instability effect in plasmas. For modes satisfying the condition \( k \leq k_c \), with

\[
k_c = \frac{\omega_p}{u_0} \left( 1 + \nu^{1/3} \right)^{3/2},
\]

with \( \nu = f/\gamma_0 \), being the cut-off wavevector, the plasma and the beam (with dispersion \( \omega = f\omega_p/\gamma_0 + u_0k \)) coalesce and the resulting dispersion relation becomes complex. In the unstable region, the dispersion relation of the plasma reads \( \omega \simeq u_0k(1 - \nu^{1/3}) + i\gamma_p \), where \( \gamma_p \) is the plasma growth rate [36, 41]

\[
\gamma_p = \frac{\nu^{1/3}}{\sqrt{3(1 + \nu^{2/3})}^{3/2}} \frac{u_0^2k^2}{\omega_p} \left[ \frac{\omega_p^2}{u_0^2k^2} \left( 1 + \nu^{2/3} \right)^3 - 1 \right]^{1/2}.
\]

The most unstable mode, occurring at \( k \simeq \omega_p/u_0 \), grows at the rate \( \gamma_p^{max} \simeq 0.69\nu^{1/3}/\omega_p \). These features are depicted in Fig. 2 a).

Given the smallness of \( \Omega \), the instability does not change noticeably in the perspective of the plasma, and therefore the discussion above holds even in the presence of axions. However, and more crucially, the small fraction of axion that participates in the beam-plasma dynamics leads to the production of axions. The normalization condition now reads \( |u_k|^2 + |v_k|^2 + |w_k|^2 = 1 \) where \( w_k \) is the beam fraction in the unstable mode. Each of these coefficients can be determined by solving the eigenvalue problem in Eq. (6) numerically, as illustrated in Fig. 2 b). The axion production mechanism can thus be understood as follows: the beam transfers energy to the plasma, which becomes unstable. Then, the magnetic field mixes the axion and the plasmon modes. Since the upper polariton mode is essentially plasma-like near in the instability region, then the small fraction of axion may grow as well. The corresponding axion growth rate in this situation can then be determined as

\[
\gamma_{p-\varphi} = |u_k|^2\gamma_p \simeq \frac{\Omega^2}{(\omega_p - \omega_\varphi)^2 + \gamma_p^2\gamma_p}.
\]

As such, the axion production is governed by the rate equation \( \dot{N} = 2\gamma_{p-\varphi}N \), where \( N = \int_V |\varphi|^2dV \) is the axion number and \( V \) is the volume of the plasma. The plasmon-axion conversion rate is depicted in Fig. 3 a). As we can see, the more energetic the electron beams, the narrower is the spectral distribution of the produced axions. This effect is particularly important in the consequent photon emission, which can be determined as

\[
\Gamma_{\varphi-\gamma} = \gamma_{p-\varphi}P_{\varphi-\gamma},
\]

where \( P_{\varphi-\gamma} \) is the Primakoff axion-photon conversion probability in the plasma medium [42, 43],

\[
P_{\varphi-\gamma} = \sin^2\Theta \sin^2(\Delta kL).
\]

Here, \( L \) is the plasma column size, \( \tan(\Theta) = gB_0\omega/(m_\varphi^2 - \omega_p^2) \) is the mixing angle and

\[
\Delta k = \sqrt{\omega_p^2 - m_\varphi^2} - \sqrt{\omega_p^2 - \omega_p^2} \simeq \frac{m_\varphi^2 - \omega_p^2}{2\omega_p}.
\]
is the difference in momentum between an axion and a photon with the same frequency $\omega$. Since only photons with frequency larger than the $\omega_p$ escape the plasma, and given that the plasma instability terminates at the cut-off frequency $\omega_c = \sqrt{m_e^2 + k_c^2}$, axion-photon conversion will occur in the range $\omega_p < \omega < \omega_c$. For resonant conversion, $\omega_p \simeq m_c$, the cut-off frequency reads

$$\omega_c = \sqrt{m_e^2 + k_c^2} \simeq \sqrt{2}\omega_p \left[ 1 + \frac{3}{4} \left( \frac{f m_e}{E} \right)^{1/3} \right],$$

valid for relativistic electron beams, $E \gg m_e$. For a typical discharge plasma, $\omega_p \sim 10$ GHz, and an electron beam of relative concentration $f = 0.1$ and energy $E = 35$ MeV, the expected emission spectrum peaks at $\omega \simeq 12.5$ GHz, with a spectral width of $\Delta \omega \simeq 250$ MHz. These features are illustrated in Fig. 3 b). A remarkable emission rate of $\sim 1.5$ photon/day (considering a hypothetical experiment in a $1\text{ m}^3$ plasma column) can compete with the most optimistic predictions of axions emitted from the sun or magnetars. With an observation time of $t_{\text{obs}} \simeq 100$ hr, as it is the case of projected radio telescope experiments [44, 45], we may anticipate a flux that is several orders of magnitude higher than that produced by axion-photon conversion in neutron star magnetospheres. This allows to probe the QCD axion range with a high degree of sensitivity. The technical design of such an experiment, although under current investigation, is out of the scope of this work, as it depends on the detection schemes under choice. We can anticipate, however, that a compact experiment based on microwave single-photon detectors is an appealing and promising option [46].

Our findings can also be interesting to identify axion production via plasma instabilities taking place in magnetar magnetic pole caps. As it is known, alongside with the gamma-ray emission taking place in the region of high-density, boosted plasma $^1$, the beam-plasma instability leads to the formation of plasma bunches that generate radio emission via the curvature effect [47, 48]. During this process, the produced axions may be resonantly converted into photons at the radius $r_c$ related to the Goldreich-Julian magnetosphere density [49, 50]

$$n_c = \frac{2 \pi B_0}{e P} \frac{1}{1 - 4 \pi^2 r_c^2 \sin^2 \theta / P},$$

where $P$ is the pulsar period and $\theta$ is the polar angle with respect to the rotation axis. For $\theta = 90^\circ$, the corresponding plasma frequency is $\omega_p \simeq (1.5 \times 10^2 \text{ GHz}) \sqrt{B/10^{14} \text{ G}}/(1 \text{ sec}/P)$, yielding $\omega_p \sim 1.0$ GHz for the SGR J145-2900 magnetar ($P \simeq 3.76 \text{ s}$, $B_0 \simeq 1.6 \times 10^{14}$ G [51, 52]), a value that is close to a discharge plasma condition. However, the axion-plasma coupling parameter $\Omega \simeq 0.12$ MHz is enhanced by a factor of $10^5$ when compared to the case discussed above, producing axions at a $10^{10}$ times higher rate, as it follows from Eq. (13). Assuming the electron - or positron, depending on the orientation of the acceleration electric field at the polar cap - beam to be much more energetic than the $e^+e^-$ plasma (making the cold plasma model valid), and taking a beam relativistic factor of $\gamma_0 \sim 10^6$, we estimate a cut-off frequency of $\omega_c \simeq 1.4$ GHz. As such, a signal in the range $1.0 \text{ GHz} \leq \omega/2\pi \leq 1.4$ GHz might be expected for the conditions of the experiment proposed in Refs. [44, 45], based on axion dark matter conversion (notice that in our case we do not need a dark matter background). At this stage, however, we can not commit wether or not the axions produced via streaming instability can be detected within the sensitivity of telescopes such as CAST or the Arecibo Telescope for the typical observation periods.

In conclusion, we have shown that a magnetized plasma can be an active source of dark-matter axions. For that, we exploit the two-stream instability - the exponential growth of electron plasma waves - excited by a monoenergetic electron beam to produce axions from the vacuum. The production mechanism is based on the transfer of energy from the electrons to the small axion-plasmon admixture, the later being a consequence of the axion-plasmon polariton coupling occuring in magnetized plasmas [38]. Based on realistic parameters, we estimate a discharge plasma column, injected by 35 MeV-electron
beam to produce roughly 1 axion per minute. The produced axions then may escape the plasma via the resonant axion-photon conversion mechanism, leading to the emission of a narrow microwave spectrum. Our findings will certainly motivate the design of a plasma-based experimental facility aiming at the active production of axions. On the other hand, given the abundance of astrophysical bodies displaying beam-plasma and beam-beam instabilities, we anticipate that a plethora of new exciting phenomena involving the dynamics of axions in plasma may arise in the near future.

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