

Axion pair creation by accelerated nuclei

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We study the interaction between axions and nuclei by combining the Peccei-Quinn mechanism with results from QCD which imply that the QCD condensates are reduced within nuclear matter. Thus, the effective axion mass is also reduced, implying that nuclear matter acts as an effective medium for the axion field. Consequently, we find many analogies between nuclei and dielectric bodies, including a finite axion-nucleon scattering cross section with a characteristic momentum dependence given by the nuclear structure or a Casimir type attraction between two static nuclei with a specific distance dependence. Moreover, accelerated nuclei such as in ultraperipheral heavy-ion collisions can create entangled pairs of axions in analogy to the dynamical Casimir effect which can also be interpreted as signatures of the Unruh effect.

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I. INTRODUCTION

In the history of physics, open questions and problems have often motivated the prediction of new particles, which have later been observed experimentally. Presently, axions are prominent candidates for such new particles. In the seminal work [1,2] by Peccei and Quinn, they were introduced in order to solve the strong CP problem, i.e., the question of why QCD is, at least to extremely good accuracy, invariant under charge C , parity P , and thus also time T reversal (see below and Refs. [3–9]). In addition, axions are also promising candidates for dark matter—provided that they live long enough [10–14].

So, naturally, many schemes for axion searches have been proposed and realized, but so far without finding conclusive evidence for axions; see, e.g., Ref. [15]. For example, the well-studied anisotropies of the cosmic microwave background radiation [16] and astronomical data regarding the evolution of stars [17–23] or white dwarfs [24,25] could be affected by axions and thus allow us to place restrictions on the possible axion parameter space; see also Refs. [26–30]. Another class of axions searches are Earth-based experiments. Here, prominent examples include the “light shining through wall” experiments [31–41], collider probes [42–48], and “light-by-light scattering” schemes [49–55].

The involved energy scales determine which axion mass ranges these schemes are most sensitive to. In addition, the majority of these schemes do also depend crucially on the

model-dependent coupling strengths between axions and other particles, such as photons or electrons—which are also not known *a priori*—or on additional assumptions, e.g., regarding the stellar evolution or the axion lifetime. To avoid this limitation, we restrict ourselves to a minimal, model-independent set of assumptions in the following. Regarding the axions, we just employ the Peccei-Quinn mechanism, which was the motivation for considering axions in the first place. Combining this idea with results from QCD, we derive the interaction between axions and nuclei. Thus, our results do only depend on the (QCD contribution to the) axion mass and not on other unknown coupling strengths.

II. STRONG CP PROBLEM AND NUCLEAR MATTER

Let us briefly review the basics. The field-strength tensor of the strong interaction in terms of the Yang-Mills fields A_μ^a and the $SU(3)$ structure constants f^a_{bc} is given by $G_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + f^a_{bc} A_\mu^b A_\nu^c$, while the dual tensor reads $\tilde{G}^{\mu\nu a} = \epsilon^{\mu\nu\rho\sigma} G_{\rho\sigma}^a$. In analogy to the two Lorentz invariants $E^2 - B^2$ and $\mathbf{E} \cdot \mathbf{B}$ in electrodynamics, we may form the two contractions $G_{\mu\nu}^a G_a^{\mu\nu}$ and $G_{\mu\nu}^a \tilde{G}_a^{\mu\nu}$ leading to the Lagrangian ($\hbar = c = 1$)

$$\mathcal{L}_{\text{gluons}} = -\frac{1}{4g_{\text{QCD}}^2} G_{\mu\nu}^a G_a^{\mu\nu} + \frac{\theta}{32\pi^2} G_{\mu\nu}^a \tilde{G}_a^{\mu\nu}, \quad (1)$$

where g_{QCD} is the QCD coupling strength [56]. Again, in analogy to electrodynamics, the first term $G_{\mu\nu}^a G_a^{\mu\nu}$ is even under charge C , parity P , and time T reversal, while the second term $G_{\mu\nu}^a \tilde{G}_a^{\mu\nu}$ is odd under CP and T .

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For $\theta = 0$, the theory is CP invariant, and thus expectation values of CP odd quantities such as $G_{\mu\nu}^a \tilde{G}_a^{\mu\nu}$ vanish. Allowing for a small but finite θ , however, this is no longer true. As already expected from stationary perturbation theory (treating θ as the small expansion parameter), these quantities may now acquire nonvanishing expectation values which scale linearly in θ for small θ . Specifically, one finds (to leading order)

$$\langle \hat{G}_{\mu\nu}^a \hat{G}_a^{\mu\nu} \rangle = 32\pi^2 \langle \hat{\Psi} \hat{\Psi} \rangle \frac{m_u m_d}{m_u + m_d} \theta, \quad (2)$$

where m_u and m_d are the masses of the up and down quarks while $\langle \hat{\Psi} \hat{\Psi} \rangle \approx -10^{-2} \text{ GeV}^3$ is the quark condensate; see, e.g., Ref. [6].

In a similar way, the θ -term would also generate other CP -violating expectation values and observables, which are not observed experimentally. Present experimental bounds (e.g., from the electric dipole moment of the neutron [57]) limit this parameter to $|\theta| < 10^{-9}$. The question of why this parameter is so small (or even zero) is usually referred to as the strong CP problem.

A. Peccei-Quinn mechanism

To solve this strong CP problem, the idea is to postulate a massless pseudoscalar axion field ϕ coupled to the $G_{\mu\nu}^a \tilde{G}_a^{\mu\nu}$ term

$$\mathcal{L}_{\text{axion}} = \frac{1}{2} (\partial_\mu \phi) (\partial^\mu \phi) + \frac{g_\phi}{32\pi^2} G_{\mu\nu}^a \tilde{G}_a^{\mu\nu} \phi. \quad (3)$$

Since ϕ is massless, the θ -term in Eq. (1) can be absorbed by a simple redefinition of $\phi \rightarrow \phi - \theta/g_\phi$. After that, the equation of motion $\square\phi = g_\phi G_{\mu\nu}^a \tilde{G}_a^{\mu\nu} / (32\pi^2)$ can be combined with Eq. (2) to show that the expectation value $\phi = \langle \hat{\phi} \rangle$ behaves as a massive scalar field where the effective axion mass m_ϕ is determined by (see also Ref. [58])

$$m_\phi^2 = -g_\phi^2 \langle \hat{\Psi} \hat{\Psi} \rangle \frac{m_u m_d}{m_u + m_d}. \quad (4)$$

Note that the quark condensate $\langle \hat{\Psi} \hat{\Psi} \rangle$ is negative.

Even though this line of arguments is extremely simplified and leaves out many important details, it allows us to grasp the main idea. In analogy to a superconductor where the coupling between the electromagnetic field A_μ and the currents j^μ inside the superconductor generate an effective mass for the photons, the QCD fluctuations coupled to the axion field generate an effective mass m_ϕ term for the latter. Thus, in the ground state, the electromagnetic field vanishes in the superconductor (Meißner-Ochsenfeld effect). Similarly, the CP -violating term vanishes in vacuum

$\phi = \langle \hat{\phi} \rangle = 0$, which is how the Peccei-Quinn mechanism solves the strong CP problem.

B. Nuclear matter

After having sketched the Peccei-Quinn mechanism in vacuum, let us consider the situation in nuclear matter, say, inside a large nucleus [another example could be the interior of a neutron star (see, e.g., Refs. [58–62]), but we do not focus on this case here]. In principle, the QCD fluctuations inside nuclei such as the quark condensate in Eqs. (2) and (4) could be computed exactly via a full fledged nonperturbative QCD calculation of the nuclear structure (e.g., in lattice QCD), but unfortunately this is beyond our present capabilities. Nevertheless, it is by now well established that the quark condensate is reduced inside nuclei. As a first qualitative illustration, one might consider the most simple bag model where the nuclear interior is modeled as a free field, and thus these condensates are even reduced to zero. More elaborate studies of dense nuclear matter yield a partial reduction [58,63], where the most commonly accepted value lies between 30% and 50% (at nuclear saturation density). Note that this reduction does not only apply to nuclear matter inside neutron stars, for example, but also to the interior of nuclei—provided that they are large enough. Studying pion resonances in ^{121}Sn nuclei, for instance, a recent experiment [64] obtained a reduction of 40%, which falls within the uncertainty of theoretical calculations ranging from 35% to 45% [65–67].

For explicit calculations, from hereon, we shall use the 30% reduction as a conservative value. Retracing the line of arguments above, we see that a reduction of the quark condensate $\langle \hat{\Psi} \hat{\Psi} \rangle$ inside the nuclei would also imply a reduction of the effective axion mass, i.e., $m_\phi^2 \rightarrow m_\phi^2 [1 - \varepsilon]$, where we have denoted the reduction ratio by ε . Employing the aforementioned analogy to superconductors, a reduction of the superfluid density (e.g., by defects or a finite temperature) would lead to a suppression of the effective photon mass and thus a growth of the London penetration depth.

In view of the above arguments, the reduction of the axion mass (or, more precisely, the curvature of the axion potential) is a fairly robust result and largely independent of other model parameters and assumptions. Even though there is a slight uncertainty regarding the precise value of ε , the situation is quite different from other axion searches, where the considered range for the axion-photon interaction strength, for example, can span several orders of magnitude.

Note that the coupling between axions and gluons in Eq. (3) also suggests, via Feynman diagrams containing internal quark loops, for example, the possibility of an effective coupling between axions and photons. However, there are several potential reasons why this effective coupling could be suppressed (which might also explain why we did not see axions yet). As one option, this effective coupling

could be (partially) canceled by another contribution of different origin. As another option, interaction effects such as running coupling or confinement could suppress this effective coupling [68–71]. For example, the axion field could actually be effectively coupled to a suitable spatio-temporal average of $G_{\mu\nu}^a \tilde{G}_a^{\mu\nu}$ such that the coupling strength g_ϕ is most effective at low momentum scales. Together with quark confinement, this could suppress the Feynman diagrams containing quark loops.

In the instanton picture [56], the CP -odd expectation value (2) can be understood as a slight imbalance between instantons (which are self-dual $G_{\mu\nu}^a = \tilde{G}_{\mu\nu}^a$) and anti-instantons (which are anti-self-dual $G_{\mu\nu}^a = -\tilde{G}_{\mu\nu}^a$). To see this imbalance, one has to average over many (anti-)instantons, i.e., over a large volume. Zooming in to shorter length scales [i.e., higher momentum scales, e.g., above the QCD scale $\Lambda_{\text{QCD}} = \mathcal{O}(100 \text{ MeV})$ which sets the typical size of the instantons], one would see single (anti-)instantons instead, which are then either self-dual or anti-self-dual. Thus, at such high momentum scales, the quantum fluctuation corrections to this mean field (2) become significant, indicating that the properties of the axion field are strongly scale dependent; see also the discussion in Sec. III A below.

III. SCATTERING CROSS SECTION

Incorporating the reduction of the effective axion mass inside the nucleus, the axion evolution equation becomes

$$(\square + m_\phi^2)\phi = m_\phi^2 \varepsilon(\mathbf{r})\phi, \quad (5)$$

where $\varepsilon(\mathbf{r})$ vanishes in vacuum and describes the reduction inside the nucleus. To ensure the validity of this effective description, we assume that the energy and momentum scales of the axion field ϕ are much smaller than the characteristic scales of the QCD fluctuations, which are governed by $\Lambda_{\text{QCD}} = \mathcal{O}(100 \text{ MeV})$.

For a fixed axion frequency $\partial_t \phi \rightarrow -i\omega\phi$, the above Klein-Fock-Gordon equation (5) assumes the same form as a stationary Schrödinger equation describing scattering at a spherical well potential $\propto \varepsilon(r)$. Thus, we may apply the same methods and read off the differential scattering cross section $d\sigma/d\Omega$. For axion wavelengths much larger than the nuclear radius, the nucleus acts as a delta potential, and thus we find (to lowest order in ε , see also Refs. [72,73])

$$\frac{d\sigma}{d\Omega} = \frac{m_\phi^4 \varepsilon^2 V_{\text{nuc}}^2}{16\pi^2}, \quad (6)$$

where V_{nuc} denotes the volume of the nucleus and ε the reduction factor inside. For example, for $\varepsilon = 0.3$ [58] and an axion mass of $m_\phi = 10 \text{ MeV}$ (see Refs. [74–76]), we find a cross section of $\sigma = 40 \text{ } \mu\text{b}$ for an iron nucleus. To visualize this quantity, it would correspond to a classical mean free path of 250 m for the axion in solid iron. Thus,

this cross section should be taken into account when considering such heavy axions propagating in stellar environments; see Sec. VI.

If the axion wavelength is much longer than the distance between the iron nuclei, this picture of a classical mean free path would no longer apply, and the whole bulk iron would instead act as an effective medium. Thus, an iron cylinder could act like an optical fiber for very slow axions, or one could imagine a magnifying lens for axions, for example.

If, on the other hand, the axion wavelength becomes comparable to the nuclear radius, we get the well-known resonance effects. For even shorter axion wavelengths, the precise internal structure of the nuclei, i.e., how $\varepsilon(\mathbf{r})$ is switched on and off, starts to play a role—but at that point, one should also be careful and consider the region of validity of the above effective description (5).

A. Comparison to chiral perturbation theory

The ability to take into account the dependence on $\varepsilon(\mathbf{r})$, i.e., the nuclear structure, results from the fact that we are *not* treating the nuclei as point particles. This point-particle approximation is typically associated with low-energy effective theories, such as chiral perturbation theory; see, e.g., Refs. [60,72,73,77].

The goal of such low-energy effective theories is to provide a good approximation in the limit of low axion masses and momenta, while we are mainly interested in the regime of intermediate axion masses and momenta, such as $m_\phi = 10 \text{ MeV}$, which are sufficient to resolve the nuclear structure. As already indicated in our discussion at the end of Sec. II B, one should be very careful when extrapolating results from one energy/momentum scale to another. For example, if one takes the first-order coupling strengths and amplitudes derived in Refs. [78,79] via chiral perturbation theory (based on different assumptions) and directly (or naively) extrapolates them to larger axion masses, one finds that they display divergences. For example, the axion-nucleon coupling in Eq. (13) of Ref. [77] diverges if the axion mass equals the pion mass, while the axion-induced decay amplitude of the kaon in Eq. (16) of Ref. [78] diverges if the axion mass equals a linear combination of the pion and kaon masses.

Casting aside these problems associated with the scale dependence and directly comparing our results with the effective interaction Lagrangian $\lambda g_\phi^2 \phi^2 \bar{\Psi}_{\text{nuc}} \Psi_{\text{nuc}}$ from chiral perturbation theory, we find that the cross sections are roughly comparable, but our predictions are around 1 order of magnitude smaller (for $\varepsilon = 0.3$). In terms of the effective medium picture described in Sec. II B, chiral perturbation theory would correspond to the bag model and does not reproduce the finite reduction factor of $\varepsilon = 0.3$. Given our ignorance of the scale dependence, this comparably slight disagreement (of the naive extrapolation) is probably not too surprising.

Ideally, a full solution of QCD should reproduce chiral perturbation theory on the one hand and the effective medium description employed here on the other hand in suitable limiting cases. Given our lack of knowledge, we have to resort to suitable interpolations between these limiting cases (with all due prudence). Thus, chiral perturbation theory and the effective medium description employed here describe complementary approaches. To avoid double counting, the results of the two approaches should not be simply added but suitably interpolated.

IV. CASIMIR ATTRACTION BETWEEN STATIC NUCLEI

Even in the absence of real axions, their inevitable quantum fluctuations can also induce potentially measurable effects. As a first example, let us study the induced Casimir attraction between two static nuclei. To this end, let us start from the perturbation Hamiltonian of the quantized axion field $\hat{\phi}$,

$$\hat{H}_{\text{int}} = \frac{m_\phi^2}{2} \int d^3r \varepsilon(\mathbf{r}) \hat{\phi}^2. \quad (7)$$

Using ε as a small expansion parameter, we may apply second-order stationary perturbation theory to obtain the ground-state energy shift

$$\Delta E_{(2)} = \sum_{|\Psi\rangle \neq |0\rangle} \frac{|\langle \Psi | \hat{H}_{\text{int}} | 0 \rangle|^2}{E_{|\Psi\rangle} - E_{|0\rangle}}. \quad (8)$$

In view of the structure of the perturbation Hamiltonian (7), we see that two-axion states $|\Psi\rangle = |\mathbf{k}_1, \mathbf{k}_2\rangle$ are the relevant intermediate states.

Let us first consider the nonrelativistic limit where the two axion momenta $\mathbf{k}_1$ and $\mathbf{k}_2$ are well below the axion mass m_ϕ . Then, we may approximate the energy denominator in Eq. (8) by $E_{|\Psi\rangle} - E_{|0\rangle} \approx 2m_\phi$, which yields the energy shift containing the perturbation Hamiltonian (7) in normal ordering $\Delta E_{(2)} \approx \langle 0 | (\hat{H}_{\text{int}})^2 | 0 \rangle / (2m_\phi)$. Using the Wick theorem and focusing on those contributions which depend on the distance between the nuclei, which give rise to the Casimir effect, we find

$$\Delta E_{\text{Casimir}} \approx \frac{m_\phi^3 \varepsilon^2 V_{\text{nuc}}^2}{8} (\langle 0 | \hat{\phi}(\mathbf{r}_1) \hat{\phi}(\mathbf{r}_2) | 0 \rangle)^2, \quad (9)$$

where we have assumed that the nuclear radii are much smaller than all other relevant length scales such that we may approximate the nuclei as point particles at the positions $\mathbf{r}_1$ and $\mathbf{r}_2$.

Here, $\langle 0 | \hat{\phi}(\mathbf{r}_1) \hat{\phi}(\mathbf{r}_2) | 0 \rangle = W(\ell)$ denotes the usual two-point function $W(\ell) = m_\phi K_1(m_\phi \ell) / (4\pi^2 \ell)$ of a massive scalar field depending on the distance $\ell = |\mathbf{r}_1 - \mathbf{r}_2|$, where K_1 is the modified Bessel function of the second kind. At large distances $m_\phi \ell \gg 1$ (which is the case considered above),

it obeys the asymptotic form $W(\ell) \sim \sqrt{m_\phi / (32\pi^3 \ell^3)} e^{-m_\phi \ell}$; see Ref. [80].

In view of the exponential suppression at large separations ℓ , the Casimir force becomes most pronounced at smaller distances. In this regime, however, the approximation $E_{|\Psi\rangle} - E_{|0\rangle} \approx 2m_\phi$ is no longer applicable, and we have to perform the sum over all intermediate two-axion states $|\Psi\rangle = |\mathbf{k}_1, \mathbf{k}_2\rangle$, i.e., the d^3k_1 and d^3k_2 integrations; see Appendix A. For very small distances, the precise functional form of $\varepsilon(\mathbf{r})$ starts to play a role, which depends on the nuclear model. Again, this dependence on $\varepsilon(\mathbf{r})$ is related to the fact that we are not treating the nuclei as point particles.

To provide an example for the order of magnitude, let us take an axion mass of $m_\phi = 50$ MeV and two iron nuclei at a rather small distance of order 10 fm (i.e., just above twice the nuclear radius), for which we obtain a Casimir energy of order 10 eV; see Appendix A. Actually measuring this Casimir energy would be tremendously difficult in view of the strong Coulomb repulsion between the iron nuclei at large distances and the nuclear forces at small distances. This will remain true for other values of the axion mass: for large m_ϕ , the Casimir attraction would be of very short range, while for small m_ϕ , its overall strength becomes very weak.

V. DYNAMICAL CASIMIR EFFECT BY ACCELERATED NUCLEI

After having discussed the static Casimir effect, let us consider a dynamical scenario involving accelerated nuclei. One way of realizing such a scenario would be to have ultraperipheral collisions of nuclei in analogy to the famous Rutherford experiment. In this case, the accelerated nuclei can create entangled axions pairs out of the vacuum, which is often referred to as the dynamical Casimir effect; see, e.g., Refs. [81–84]. The amplitude for this process can be calculated with the same perturbation Hamiltonian (7) but now adapted to the case of moving nuclei. To lowest order in ε , the amplitudes stemming from the two (or more) nuclei would just add up, so it is sufficient to consider one accelerated nucleus. Actually, using the same arguments as in Ref. [85], the axion radiation emitted by such an accelerated nucleus can also be interpreted as a signature of the Unruh effect; see also Refs. [86–89]. Assuming a constant acceleration a for simplicity, a comoving observer would perceive the vacuum as a thermal bath of axions with the Unruh temperature $T_{\text{Unruh}} = a / (2\pi k_B)$. Thus, the nucleus can scatter an axion out of this thermal bath according to the cross section (6). Translated back to the laboratory frame, this scattering process would correspond to the emission of an entangled pair of axions; see also Ref. [90].

To calculate the axion pair-creation amplitude, let us again assume that the nuclear radius is much smaller than all other relevant length scales such as the axion

wavelengths. Then, we may approximate the nucleus as a point particle with the trajectory $\mathbf{r}_{\text{nuc}}(t)$ and derive the axion pair-creation amplitude via time-dependent perturbation theory with respect to $\hat{H}_{\text{int}}$,

$$\mathfrak{A}_{\mathbf{k}_1\mathbf{k}_2} = \frac{\varepsilon m_\phi^2 V_{\text{nuc}}}{2\sqrt{\omega_1\omega_2}} \int dt e^{i(\omega_1+\omega_2)t-i(\mathbf{k}_1+\mathbf{k}_2)\cdot\mathbf{r}_{\text{nuc}}(t)}. \quad (10)$$

Following Ref. [85], let us introduce the new time coordinate $\tau = t - (\mathbf{k}_1 + \mathbf{k}_2) \cdot \mathbf{r}_{\text{nuc}}(t)/[\omega_1 + \omega_2]$. For nonrelativistic velocities $|\dot{\mathbf{r}}_{\text{nuc}}| \ll 1$, we may simplify the amplitude in Eq. (10) to

$$\mathfrak{A}_{\mathbf{k}_1\mathbf{k}_2} \approx \frac{\varepsilon m_\phi^2 V_{\text{nuc}}}{2\sqrt{\omega_1\omega_2}} \frac{\mathbf{k}_1 + \mathbf{k}_2}{\omega_1 + \omega_2} \cdot \tilde{\mathbf{v}}_{\text{nuc}}(\omega_1 + \omega_2), \quad (11)$$

where $\tilde{\mathbf{v}}_{\text{nuc}}$ denotes the Fourier transform of $\dot{\mathbf{r}}_{\text{nuc}}(\tau)$. As a consequence, the axion pair-creation amplitude becomes maximal if both axions are emitted into the same direction $\mathbf{k}_1 \parallel \mathbf{k}_2$ aligned with the acceleration.

For the ultraperipheral collision scenario, $\dot{\mathbf{r}}_{\text{nuc}}(\tau)$ is constant asymptotically and changes during a characteristic acceleration timescale $\Delta\tau$. Thus, for small frequencies $\omega \ll 1/\Delta\tau$, the Fourier transform $\tilde{\mathbf{v}}_{\text{nuc}}$ is basically that of a Heaviside step function $\tilde{\mathbf{v}}_{\text{nuc}}(\omega) \sim 1/\omega$. For high frequencies $\omega \gg 1/\Delta\tau$, on the other hand, the Fourier transform $\tilde{\mathbf{v}}_{\text{nuc}}$ falls off very rapidly since $\dot{\mathbf{r}}_{\text{nuc}}(\tau)$ is a smooth function.

The total axion pair-creation probability $\mathfrak{P}$ is obtained by integrating $|\mathfrak{A}_{\mathbf{k}_1\mathbf{k}_2}|^2$ over the two momenta $\mathbf{k}_1$ and $\mathbf{k}_2$. Its scaling depends on the axion mass m_ϕ in comparison to the acceleration timescale $\Delta\tau$. For long timescales $\Delta\tau \gg 1/m_\phi$, the prefactor m_ϕ in Eq. (11) might be large, but the Fourier transform $\tilde{\mathbf{v}}_{\text{nuc}}(\omega_1 + \omega_2)$ would be strongly suppressed. Thus, a larger pair-creation probability $\mathfrak{P}$ is obtained for short timescales $\Delta\tau \ll 1/m_\phi$. In this case, the created axions are predominantly ultrarelativistic, and thus we may approximate $\omega_1 \approx |\mathbf{k}_1|$ and $\omega_2 \approx |\mathbf{k}_2|$. Integrating $|\mathfrak{A}_{\mathbf{k}_1\mathbf{k}_2}|^2$ up to the cutoff scale $\omega_{\text{cut}} \approx |\mathbf{k}_{\text{cut}}| \sim 1/\Delta\tau$, we find the scaling

$$\mathfrak{P} = \mathcal{O}\left(\frac{\varepsilon^2 m_\phi^4 V_{\text{nuc}}^2}{(\Delta\tau)^2}\right). \quad (12)$$

If we assume the hierarchy $R_{\text{nuc}} \propto V_{\text{nuc}}^{1/3} \ll \Delta\tau \ll 1/m_\phi$ in terms of the effective nuclear radius R_{nuc} as defined by $V_{\text{nuc}} = 4\pi R_{\text{nuc}}^3/3$, this is a small number $\mathfrak{P} \ll 1$. However, considering many collisions, the total probability might become non-negligible. Rewriting the probability $\mathfrak{P}$ in the form $\mathfrak{P} = \mathcal{O}(\varepsilon^2 [m_\phi R_{\text{nuc}}]^4 [R_{\text{nuc}}/\Delta\tau]^2)$, we may use the nonrelativistic assumption $R_{\text{nuc}}/\Delta\tau < 1$ to infer the bound $\mathfrak{P} < \mathcal{O}(\varepsilon^2 [m_\phi R_{\text{nuc}}]^4)$. Pushing our analysis to the limit where $V_{\text{nuc}}^{1/3} \sim \Delta\tau \sim 1/m_\phi$, we would expect $\mathfrak{P} = \mathcal{O}(\varepsilon^2)$, i.e., a probability on the percent level.

The created axions could be detected in analogy to other axion or dark matter searches. For example, one could envisage a large volume $\mathcal{O}(\text{m}^3)$ of liquid xenon as in the XENONnt [91] or LZ [92] experiments. In comparison to solid iron, the classical mean free path would be even a bit smaller, 220 m for $m_\phi = 10$ MeV, leading to a scattering probability on the percent level for a single axion traversing the xenon tank. As in the XENONnt or LZ experiments, the interaction could be detected via the induced recoil of xenon, which has a kinetic energy of several keV and is in the correct regime to be measured. Collision rates of $\mathcal{O}(10^9/\text{s})$ between nuclei as in LHC or Relativistic Heavy Ion Collider (RHIC), together with a long run time of the experiment (as XENONnt or LZ), allow us to probe a comparably large range of parameter space; see Appendix C.

Depending on the axion-photon coupling $g_{\phi\gamma}$, another option would be to look for the decay of the created axions into photon pairs. For small $g_{\phi\gamma}$, this approach could be advantageous in comparison to other axion searches considering both the creation and the decay of the axions via pairs of photons. In our case, the creation probability scales as Eq. (12), while the decay probability scales as $g_{\phi\gamma}^2$, whereas the total (creation and decay) probability is suppressed as $g_{\phi\gamma}^4$ for the other photon-based schemes.

VI. CONCLUSIONS

We combine the Peccei-Quinn mechanism with results from QCD. Since the quark condensate changes inside nuclear matter, we deduce an analogous shift of the effective axion mass. In this way, we obtain an effective axion interaction with nuclei or other objects that generate inhomogeneities in the condensate. This description is quite robust and largely independent of other parameters (such as the axion-photon coupling strengths) or model assumptions. Note that we only deduce a change of the axion mass, i.e., the curvature of the axion potential—the expectation value of the axion field is not modified (in contrast to previous works [58–62] which consider a change of the expectation value inside neutron stars, for example).

For a first application, we derive scattering cross sections σ between axions and nuclei, which only depend on the axion mass m_ϕ and the nuclear radius. For example, for an iron nucleus and $m_\phi = 10$ MeV (see Refs. [74–76]), we find $\sigma = 40 \mu\text{b}$. Note that we consider comparably heavy axions here, whereas many contemporary models focus on much lighter axions. However, as explained above, there are several options or models for which the existence of heavier axions is not excluded. As long as we do not know the properties of the axions yet (or whether they exist at all), we have to consider all possibilities.

These results should also be taken into account when analyzing axion bounds from astronomical data. For example, considering the energy loss of a white dwarf or asymptotic giant branch star core [25] with a radius of

TABLE I. Example values for possible experimental and observational signatures of the coupling between nuclei and axions depending on the axion mass. Note that the last row gives the probability per collision, which should be multiplied with the collision rate, such as $\mathcal{O}(10^9/\text{s})$ at LHC or RHIC.

Axion mass m_ϕ	100 keV	1 MeV	10 MeV
Scattering cross section with iron nucleus	0.4 pb	4 nb	40 μb
Mean free path in solid iron	2.5×10^7 km	2.5×10^3 km	250 m
Mean free path inside white dwarfs	$\mathcal{O}(10^3$ km)	$\mathcal{O}(100$ m)	$\mathcal{O}(1$ cm)
Casimir energy between two protons at 2 fm distance	$\mathcal{O}(10^{-16}$ eV)	$\mathcal{O}(10^{-10}$ eV)	$\mathcal{O}(10^{-4}$ eV)
Creation probability per heavy-ion collision	$<\mathcal{O}(10^{-10})$	$<\mathcal{O}(10^{-6})$	$<\mathcal{O}(10^{-2})$
Mean free path in liquid xenon	2.2×10^7 km	2.2×10^3 km	220 m
Detection probability in 1 m liquid xenon	$<\mathcal{O}(10^{-20})$	$<\mathcal{O}(10^{-12})$	$<\mathcal{O}(10^{-4})$

10^7 m, composed of oxygen with a density of 10^5 of that of ordinary matter, axions with a mass above 300 keV would have a mean free path below 10^4 m. Even though the scattering derived here would be elastic (to lowest order), this should open up many channels for the axions to decay or to lose energy before they can leave the white dwarf. Even heavier axions with a mass above 10 MeV would have a mean free path below 1 cm. Assuming the usual diffusive scaling based on a random walk, these axions would be trapped for a timescale of order years instead of leaving the white dwarf within a fraction of a second.

Similar arguments apply to other dense stellar media, e.g., in supernovae [93–99]. Note that special care is required for neutron stars, since there the underlying assumption of well-isolated and separated nuclei is not such a good approximation anymore.

Furthermore, the mean free path of axions with a mass of 10 MeV in solid iron of around 250 m shows that the effects considered here may also become relevant for Earth-based axion search schemes. For very low-energy axions, one might even imagine interesting possibilities such as axion wave guides or magnifying lenses for axions.

Even in the absence of real axions, their inevitable quantum vacuum fluctuations would generate a Casimir type attraction between nuclei, where the range of the force is governed by the axion mass. For the parameters considered here (see Appendix A), the Casimir energies lie well below the Coulomb repulsion, rendering a detection tremendously difficult. However, speculating about larger values of ϵm_ϕ^2 and depending on the precise shape of the nuclei, this situation might change and the Casimir contribution could perhaps become non-negligible for understanding nuclear interactions (e.g., the deuteron binding energy of around 2 MeV). Note that the dependence of the Casimir energy on the nuclear shape can only be studied because we are *not* approximating the nuclei as point particles (as is done in other approaches.)

Finally, axion creation is also possible, depending on the space-time dependence of the quark condensate. We show that in ultraperipheral heavy ion collisions the acceleration of nuclei, i.e., of the dips in the quark condensate, may generate entangled axion pairs akin to the dynamical

Casimir effect or as signatures of Unruh radiation [85]. The axions created in this way could be detected in analogy to other axion or dark-matter searches such as in liquid xenon experiments [91,92].

Typical examples of possible experimental and observational signatures and their strong dependence on the axion mass are provided in Table I. We see that for the values considered there the axion-induced Casimir attraction between nuclei is most probably too small to be measurable. However, the other effects might well have observational or experimental consequences—provided that axions in the corresponding mass range exist.

Moreover, ultrastrong magnetic fields generated by such ultraperipheral collisions [100,101] may change the quark condensate by a few percent too [102–106], a non-perturbative in external photon extension of the recently developed quadratic photon-axion coupling [107,108], which can also lead to the creation of axion pairs. The production rate scales analogously to Eq. (12) and applies to any field that shifts the quark condensate, including $\mathbf{E} \cdot \mathbf{B} = 0$ configurations; see Appendix B.

Generalizing our results to axionlike particles whose masses m_ϕ can also receive non-QCD contributions (see, e.g., Refs. [15,109]), the reduction of m_ϕ would only apply to the part of m_ϕ which is associated to the Peccei-Quinn mechanism; i.e., the value of ϵ would be correspondingly lower. (Determining how ϵ changes in detail will require resolving the behavior of the non-QCD mass contributions, e.g., whether they are generated before or after the Peccei-Quinn symmetry restoration; see Appendix C).

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APPENDIX A: CALCULATION OF CASIMIR ENERGY

For simplicity, we modeled $\epsilon(\mathbf{r})$ by a Gaussian. For two iron nuclei, the results are shown in Fig. 1 as a function of axion mass m_ϕ and distance $\ell = |\mathbf{r}_1 - \mathbf{r}_2|$.

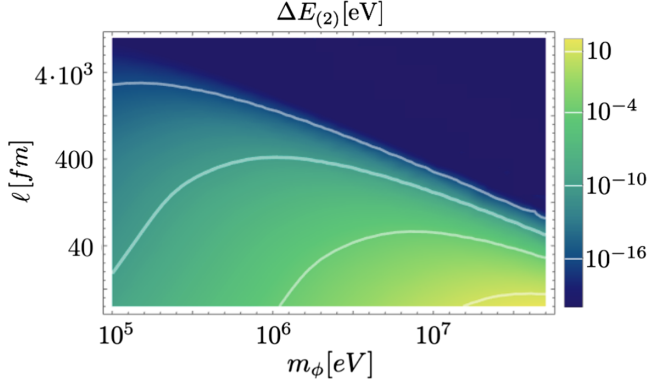


FIG. 1. Casimir energy between two iron nuclei, as a function of axion mass and distance.

Going away from a Gaussian profile $\varepsilon(\mathbf{r})$ and considering sharper (i.e., more steplike) switching functions, one would expect to obtain larger Casimir forces.

Of course, the strong Coulomb repulsion between the iron nuclei of order 100 MeV would render the measurement of such a small attractive force tremendously difficult. For smaller axion masses, which would lead to Casimir forces of longer range (though much smaller amplitude), one could consider iron atoms instead, where the Coulomb repulsion is shielded by the electrons. This would be very similar to “fifth-force” measurements; see, e.g., Refs. [73,77,110–114].

For another option, we could push our analysis to the extreme limit and consider two protons or neutrons, together with slightly heavier axions up to 100 MeV mass. For very small distances between 2 and 3 fm, the Casimir energy would be of the same order, between $\mathcal{O}(10 \text{ eV})$ and $\mathcal{O}(100 \text{ eV})$. However, the Coulomb repulsion between two protons would be much smaller than in the iron case above, such that the Casimir energy might amount to a correction in the percent to per-mill range. For two neutrons, the dominant force is the magnetic dipole-dipole interaction. Here, the Casimir energy may induce a correction also in the percent to per-mill range, provided that the dipoles are aligned. For random dipole orientations, the dipole-dipole interactions would average out such that only the Casimir effect remains.

Note that our scattering potential applied to neutrons is a bit similar to the quadratic axion-nucleon coupling considered previously [72,73,77] which for light axions has astrophysical and fifth force implications.

APPENDIX B: DYNAMICAL CASIMIR EFFECT FOR STRONG MAGNETIC FIELDS

In addition to accelerating nuclei, modifications to the quark condensate driven by strong and time-dependent electromagnetic fields may also provide a mechanism for axion production. Peripheral relativistic heavy ion collisions are capable of generating electromagnetic fields exceeding the Schwinger limit $\mathcal{O}(\text{MeV}^2)$ by up to 5 orders

of magnitude [101]. We focus in particular on the magnetic-dominated collision center. This magnetic field, of order $|\mathbf{B}| \sim 0.1 \text{ GeV}^2$, is expected to induce catalysis [102–104], i.e., growth of the quark condensate by a few percent [105,106]. Applying the same steps in Sec. II B, i.e., replacing the $\sim 30\%$ condensate reduction in nuclei with a percent-order increase in magnetic fields in vacuum, we obtain an axion mass modified by $\varepsilon_B \sim 0.01\text{--}0.05$, which enters in (5) in the same way as for nuclei. Note our resulting cross sections proportional to ε^2 are unaffected by sign of the mass shift (increasing or decreasing condensate).

To show how the axion pair production scales, we consider a volume V_b (subscript b labeling impact parameter) enclosing the region around the collision center where fields are magnetic dominated. The characteristic time for such collisions is $\Delta\tau \sim 10^{-25} \text{ s} \sim 100 \text{ MeV}^{-1}$ [101]. We apply the same time-dependent perturbation theory as in the case for accelerating nuclei in section V, except now in the interaction Hamiltonian we take $\varepsilon_B(\mathbf{r}, t) = \varepsilon_B V_b \delta^3(\mathbf{r} - \mathbf{r}_{\text{coll}}) e^{-t^2/(\Delta\tau)^2}$. Here, $\mathbf{r}_{\text{coll}}$ is the magnetic-dominated center of the collision, and we use a Gaussian to approximate the field’s short duration. The amplitude for axion pair creation, without a preferred direction, reads

$$\mathfrak{A}_{\mathbf{k}_1 \mathbf{k}_2} \approx \frac{\varepsilon_B m_\phi^2 V_b \Delta\tau}{2\sqrt{\omega_1 \omega_2}} e^{-(\omega_1 + \omega_2)^2 (\Delta\tau)^2 / 2} e^{i(\mathbf{k}_1 + \mathbf{k}_2) \cdot \mathbf{r}_{\text{coll}}}, \quad (\text{B1})$$

noting the factor $\Delta\tau$ from the Fourier transform of $\varepsilon_B(\mathbf{r}, t)$. We consider axions with masses $m_\phi \ll 1/\Delta\tau$ and repeat the above steps integrating $|\mathfrak{A}_{\mathbf{k}_1 \mathbf{k}_2}|^2$. As the momentum cutoff we again have $1/\Delta\tau$, giving the scaling

$$\mathfrak{P} = \mathcal{O}(\varepsilon_B^2 m_\phi^4 V_b^2 / (\Delta\tau)^2). \quad (\text{B2})$$

It might be illuminating to compare this *model-independent* process to the well-studied *model-dependent* photon-axion conversion $\gamma + \gamma \rightarrow \phi$, which has been applied to such collision scenarios [43]. We note that here the change in condensate ε_B scales quadratically in B [106]; thus, the dynamical Casimir production probability (B2) scales with the same power in field (or photon density) as the photon-axion conversion process [43]. Moreover, aside from the change in quark condensate relative to its in vacuum value and the axion mass under consideration, there is no coupling term suppressing the production as one would usually encounter in axion-photon conversion, e.g., $g_{\phi\gamma\gamma} \sim 10^{-2} \text{ GeV}^{-1}$ for a 10 MeV mass axion. As a major difference, *pairs* of axions are created by the dynamical Casimir effect, while single axions are produced by the $g_{\phi\gamma\gamma}$ coupling term. As another difference, the $g_{\phi\gamma\gamma}$ coupling term is associated with the electromagnetic pseudoscalar $\mathbf{E} \cdot \mathbf{B}$ (see, e.g., Ref. [115]), while the dynamical Casimir effect does also apply to scenarios where $\mathbf{E} \cdot \mathbf{B} = 0$.

Another strong field environment worth mentioning is found on the surface of highly magnetized neutron stars, or magnetars. Here, $\mathbf{B}$ fields reaching up to 10^2 times the Schwinger limit [100,116–118] can grow the quark condensate by $\varepsilon = \mathcal{O}(10^{-8})$, recalling the quadratic field dependence of the condensate shift. While this is a very small shift compared to that in nuclei or in B fields generated in collider experiments, it spans macroscopic distances and thus could be relevant to lighter axions.

APPENDIX C: APPLICATION TO HEAVY AXIONLIKE-PARTICLE PROBES

We return to the axion Lagrangian (3), which we extend to axionlike particles (ALPs). We assume that the ALP still solves the strong CP problem, while relaxing the relation between mass and axion-gluon coupling such that (see Eq. (90.3) of Ref. [15]) $m_{\text{ALP}} > m_\phi = 5.961 \times 10^3 g_\phi \text{ MeV}^2$.

For an example, we consider the possibility that non-QCD mass contributions arise post Peccei-Quinn mechanism, i.e., after the relation between QCD axion mass component and quark condensate in (4) is established,

$$m_{\text{ALP}}^2 = m_\phi^2 + \tilde{m}^2 = -g_\phi^2 \langle \hat{\Psi} \hat{\Psi} \rangle \frac{m_u m_d}{m_u + m_d} + \tilde{m}^2, \quad (\text{C1})$$

where g_ϕ and m_ϕ label the QCD axion coupling and mass and $\tilde{m}$ is the non-QCD mass. The resulting reduction of ALP mass inside nuclear matter can be generalized from Sec. II B,

$$m_{\text{ALP}}^2 \rightarrow m_{\text{ALP}}^2 \left[1 - \varepsilon \frac{m_\phi^2}{m_{\text{ALP}}^2} \right], \quad (\text{C2})$$

where ε is the reduction ratio of the quark condensate. Alternatively, instead of casting the term in brackets in (C2) in terms of ALP mass (with fixed coupling), we can write the condensate reduction effect as a function of ALP coupling (with fixed mass): $m_{\text{ALP}}^2 \rightarrow m_{\text{ALP}}^2 [1 - \varepsilon g_{\text{ALP}}^2 / g_\phi^2]$.

We can now generalize the dynamical (QCD) axion generation method proposal in section V to ALPs. Combining (C2) with (12), the production probability per collision scales according to the QCD axion rate, with an additional factor accounting for the smaller ALP coupling: $\mathfrak{P} \rightarrow \mathfrak{P} g_{\text{ALP}}^4 / g_\phi^4$.

For an example, we take axions dynamically generated by the accelerated nuclei of size $V_{\text{nucl}}^{1/3} = (200 \text{ MeV})^{-1}$, in collisions spanning the timescale of $\Delta t = 100 \text{ MeV}$. This gives a QCD axion production rate on the order of one per 10^8 collisions, which for a $\mathcal{O}(10^9/\text{s})$ collision rate carried out over the course of over a year amounts to $\mathcal{O}(10^8)$ pairs per year. To generalize to the smaller ALP couplings, we multiply the QCD axion result by $g_{\text{ALP}}^4 / g_\phi^4$.

We consider two detection methods introduced in Sec. V: first is to apply this flux to an $\mathcal{O}(\text{m}^3)$ volume liquid xenon target to probe the pseudoscalar coupling to gluons. Using (6), we expect a collision probability of $\mathcal{O}(10^{-6}) \times g_{\text{ALP}}^4 / g_\phi^4$ per MeV mass ALP pair. Going to heavier axion masses, the number of produced ALPs by accelerating nuclei and their collision rate in the xenon detector each increase with factor $m_{\text{ALP}}^4 / \text{MeV}^4$. For $m_{\text{ALP}} > 50 \text{ MeV}$, the mean free path in the xenon liquid shrinks to the point that multiple recoils are possible. In this mass range, the recoil signatures are competitive with beam dump [44] and OPAL [43] constraints. However, a substantial portion of ALPs can decay (depending on the model) into photons prior to entering as well as inside the xenon detector—an analysis of expected recoil events will be presented as a follow up work.

Another method is to probe axion to photon decay ($\phi \rightarrow \gamma\gamma$). We recall the Large Electron-Positron Collider collider based axion probe [119] relies on the process $e^+e^- \rightarrow Z$ producing $\mathcal{O}(10^6)$ Z bosons, which can undergo the decay $Z \rightarrow \gamma + \phi$, followed by the axion's decay into photons. In comparison, the dynamical Casimir effect offers a possible advantage by producing axions directly from accelerating nuclei, without requiring intermediate steps, e.g., Z -boson production. We recall a flux of $\mathcal{O}(10^8) \times g_{\text{ALP}}^4 / g_\phi^4$ axion pairs per year at MeV mass, increasing further at heavier masses. We expect this will allow us to probe new parameter space for laboratory-based experiments in the 40–100 MeV mass range, where OPAL currently sets the lower bound to ALP-photon coupling [43]. In this regime, the lower bound to the ALP-photon coupling is only slightly below the QCD axion (Kim–Shifman–Vainshtein–Zakharov model), and there is thus room for new parameter space to probe, despite suppression of the dynamical effect due to ε scaling with $g_{\text{ALP}}^2 / g_\phi^2$ (in converting QCD axion to ALP production rates).

Determining the exact detection thresholds, i.e., distinguishing the above axion recoils in xenon and $\phi \rightarrow \gamma\gamma$ decay products from background effects, will require solving the colliding nuclei trajectories and computing the spectra of produced axions, which we carry out in upcoming work. It is also important to note that the above considerations, due to the small volume V_{nucl} of the accelerating nuclei, require heavy axions for sizable production rates. However, by accounting also for the dynamical axion creation due to electromagnetic (EM) field induced quark condensate oscillations in the same collisions, lighter axions can be probed. In this case, despite having a smaller quark condensate reduction ε compared to that in nuclei, the volume occupied by the strong EM fields is much larger than V_{nucl} . An analysis of the dynamics of electric and magnetic fields in such collisions will also enter into follow-up work.

- [1] R. D. Peccei and H. R. Quinn, *Phys. Rev. Lett.* **38**, 1440 (1977).
- [2] R. D. Peccei and H. R. Quinn, *Phys. Rev. D* **16**, 1791 (1977).
- [3] S. Weinberg, *Phys. Rev. Lett.* **40**, 223 (1978).
- [4] F. Wilczek, *Phys. Rev. Lett.* **40**, 279 (1978).
- [5] J. E. Kim, *Phys. Rev. Lett.* **43**, 103 (1979).
- [6] M. A. Shifman, A. I. Vainshtein, and V. I. Zakharov, *Nucl. Phys.* **B166**, 493 (1980).
- [7] A. R. Zhitnitsky, *Sov. J. Nucl. Phys.* **31**, 260 (1980).
- [8] M. Dine, W. Fischler, and M. Srednicki, *Phys. Lett.* **104B**, 199 (1981).
- [9] R. D. Peccei, *Lect. Notes Phys.* **741**, 3 (2008).
- [10] J. Preskill, M. B. Wise, and F. Wilczek, *Phys. Lett.* **120B**, 127 (1983).
- [11] L. F. Abbott and P. Sikivie, *Phys. Lett.* **120B**, 133 (1983).
- [12] M. Dine and W. Fischler, *Phys. Lett.* **120B**, 137 (1983).
- [13] P. Sikivie, *Phys. Rev. Lett.* **51**, 1415 (1983).
- [14] P. Sikivie, *Lect. Notes Phys.* **741**, 19 (2008).
- [15] A. Ringwald, L. J. Rosenberg, and G. Rybka, *Prog. Theor. Exp. Phys.* **2022**, 083C01 (2022).
- [16] D. Cadamuro and J. Redondo, *J. Cosmol. Astropart. Phys.* **02** (2012) 032.
- [17] R. Bernabei *et al.*, *Phys. Lett. B* **515**, 6 (2001).
- [18] G. G. Raffelt, *Lect. Notes Phys.* **741**, 51 (2008).
- [19] A. Caputo and G. Raffelt, *Proc. Sci. COSMICWISPer* (2024) 041.
- [20] A. Ayala, I. Domínguez, M. Giannotti, A. Mirizzi, and O. Straniero, *Phys. Rev. Lett.* **113**, 191302 (2014).
- [21] I. G. Irastorza and J. Redondo, *Prog. Part. Nucl. Phys.* **102**, 89 (2018).
- [22] A. De Angelis, M. Roncadelli, and O. Mansutti, *Phys. Rev. D* **76**, 121301 (2007).
- [23] J. F. Fortin, H. K. Guo, S. P. Harris, D. Kim, K. Sinha, and C. Sun, *Int. J. Mod. Phys. D* **30**, 2130002 (2021).
- [24] J. Isern, E. García-Berro, S. Torres, and S. Catalan, *Astrophys. J. Lett.* **682** (2008), L109.
- [25] M. J. Dolan, F. J. Hiskens, and R. R. Volkas, *J. Cosmol. Astropart. Phys.* **09** (2021) 010.
- [26] J. Jaeckel and A. Ringwald, *Annu. Rev. Nucl. Part. Sci.* **60**, 405 (2010).
- [27] K. Baker *et al.*, *Ann. Phys. (Amsterdam)* **525**, A93 (2013).
- [28] M. Buschmann, J. W. Foster, and B. R. Safdi, *Phys. Rev. Lett.* **124**, 161103 (2020).
- [29] V. Anastassopoulos *et al.* (CAST Collaboration), *Nat. Phys.* **13**, 584 (2017).
- [30] J. B. Dent, B. Dutta, J. L. Newstead, and A. Thompson, *Phys. Rev. Lett.* **125**, 131805 (2020).
- [31] G. Ruoso *et al.*, *Z. Phys. C* **56**, 505 (1992).
- [32] R. Cameron *et al.*, *Phys. Rev. D* **47**, 3707 (1993).
- [33] M. Fouche *et al.*, *Phys. Rev. D* **78**, 032013 (2008).
- [34] J. Redondo and A. Ringwald, *Contemp. Phys.* **52**, 211 (2011).
- [35] A. S. Chou *et al.* (GammeV (T-969) Collaboration), *Phys. Rev. Lett.* **100**, 080402 (2008).
- [36] A. Afanasev, O. K. Baker, K. B. Beard, G. Biallas, J. Boyce, M. Minarni, R. Ramdon, M. Shinn, and P. Slocum, *Phys. Rev. Lett.* **101**, 120401 (2008).
- [37] K. Ehret *et al.*, *Phys. Lett. B* **689**, 149 (2010).
- [38] R. Ballou *et al.* (OSQAR Collaboration), *Phys. Rev. D* **92**, 092002 (2015).
- [39] T. Inada *et al.*, *Phys. Rev. Lett.* **118**, 071803 (2017).
- [40] K. A. Beyer, G. Marocco, R. Bingham, and G. Gregori, *Phys. Rev. D* **105**, 035031 (2022).
- [41] J. W. D. Halliday *et al.*, *Phys. Rev. Lett.* **134**, 055001 (2025).
- [42] B. Döbrich, J. Jaeckel, F. Kahlhoefer, A. Ringwald, and K. Schmidt-Hoberg, *J. High Energy Phys.* **02** (2016) 018.
- [43] S. Knapen, T. Lin, H. K. Lou, and T. Melia, *Phys. Rev. Lett.* **118**, 171801 (2017).
- [44] M. J. Dolan, T. Ferber, C. Hearty, F. Kahlhoefer, and K. Schmidt-Hoberg, *J. High Energy Phys.* **12** (2017) 094.
- [45] A. Mariotti, D. Redigolo, F. Sala, and K. Tobioka, *Phys. Lett. B* **783**, 13 (2018).
- [46] A. Ariga *et al.* (FASER Collaboration), *Phys. Rev. D* **99**, 095011 (2019).
- [47] D. Banerjee *et al.* (NA64 Collaboration), *Phys. Rev. Lett.* **125**, 081801 (2020).
- [48] S. Gori, G. Perez, and K. Tobioka, *J. High Energy Phys.* **08** (2020) 110.
- [49] L. Maiani, R. Petronzio, and E. Zavattini, *Phys. Lett. B* **175** (1986) 359.
- [50] G. Raffelt and L. Stodolsky, *Phys. Rev. D* **37**, 1237 (1988).
- [51] Y. Semertzidis, R. Cameron, G. Cantatore, A. C. Melissinos, J. Rogers, H. Halama, A. Prodell, F. Nezzrick, C. Rizzo, and E. Zavattini, *Phys. Rev. Lett.* **64**, 2988 (1990).
- [52] A. Ejlli, F. Della Valle, U. Gastaldi, G. Messineo, R. Pengo, G. Ruoso, and G. Zavattini, *Phys. Rep.* **871**, 1 (2020).
- [53] R. Battesti *et al.*, *Eur. Phys. J. D* **46**, 323 (2008).
- [54] X. Fan *et al.*, *Eur. Phys. J. D* **71**, 308 (2017).
- [55] S. Evans and R. Schützhold, *Phys. Rev. D* **109**, L091901 (2024).
- [56] T. Schäfer and E. V. Shuryak, *Rev. Mod. Phys.* **70**, 323 (1998).
- [57] C. Abel, S. Afach, N. J. Ayres, C. A. Baker, G. Ban, G. Bison, K. Bodek, V. Bondar, M. Burghoff and E. Chanel *et al.*, *Phys. Rev. Lett.* **124**, 081803 (2020).
- [58] R. Balkin, J. Serra, K. Springmann, and A. Weiler, *J. High Energy Phys.* **07** (2020) 221.
- [59] A. Hook and J. Huang, *J. High Energy Phys.* **06** (2018) 036.
- [60] R. Balkin, J. Serra, K. Springmann, S. Stelzl, and A. Weiler, *Phys. Rev. D* **109**, 095032 (2024).
- [61] L. Di Luzio, B. Gavela, P. Quilez, and A. Ringwald, *J. High Energy Phys.* **05** (2021) 184.
- [62] F. Anzuini, A. Gómez-Bañón, J. A. Pons, A. Melatos, and P. D. Lasky, *Phys. Rev. D* **109**, 083030 (2024).
- [63] T. D. Cohen, R. J. Furnstahl, and D. K. Griegel, *Phys. Rev. C* **45**, 1881 (1992).
- [64] T. Nishi *et al.* (piAF Collaboration), *Nat. Phys.* **19**, 788 (2023).
- [65] D. Jido, T. Hatsuda, and T. Kunihiro, *Phys. Lett. B* **670**, 109 (2008).
- [66] E. Friedman and A. Gal, *Phys. Lett. B* **792**, 340 (2019).
- [67] S. Hübsch and D. Jido, *Phys. Rev. C* **104**, 015202 (2021).
- [68] P. Jain and S. Mandal, *Int. J. Mod. Phys. D* **15**, 2095 (2006).

- [69] E. Masso and J. Redondo, *J. Cosmol. Astropart. Phys.* **09** (2005) 015.
- [70] E. Masso and J. Redondo, *Phys. Rev. Lett.* **97**, 151802 (2006).
- [71] J. Jaeckel, E. Masso, J. Redondo, A. Ringwald, and F. Takahashi, *Phys. Rev. D* **75**, 013004 (2007).
- [72] H. Fukuda and S. Shirai, *Phys. Rev. D* **105**, 095030 (2022).
- [73] K. Van Tilburg, *Phys. Rev. D* **109**, 096036 (2024).
- [74] D. S. M. Alves and N. Weiner, *J. High Energy Phys.* **07** (2018) 092.
- [75] J. Liu, N. McGinnis, C. E. M. Wagner, and X. P. Wang, *J. High Energy Phys.* **05** (2021) 138.
- [76] D. S. M. Alves, *Phys. Rev. D* **103**, 055018 (2021).
- [77] M. Bauer and G. Rostagni, *Phys. Rev. Lett.* **132**, 101802 (2024).
- [78] M. Bauer, M. Neubert, S. Renner, M. Schnubel, and A. Thamm, *Phys. Rev. Lett.* **127**, 081803 (2021).
- [79] E. Goudzovski, D. Redigolo, K. Tobioka, J. Zupan, G. Alonso-Álvarez, D. S. M. Alves, S. Bansal, M. Bauer, J. Brod and V. Chobanova *et al.*, *Rep. Prog. Phys.* **86**, 016201 (2023).
- [80] N. N. Bogolyubov and D. V. Shirkov, *Intersci. Monogr. Phys. Astron.* **3**, 1 (1959).
- [81] G. T. Moore, *J. Math. Phys. (N.Y.)* **11**, 2679 (1970).
- [82] P. C. W. Davies and S. A. Fulling, *Proc. R. Soc. A* **348**, 393 (1976).
- [83] P. C. W. Davies and S. A. Fulling, *Proc. R. Soc. A* **356**, 237 (1977).
- [84] V. V. Dodonov, *MDPI Phys.* **2**, 67 (2020).
- [85] R. Schützhold, G. Schaller, and D. Habs, *Phys. Rev. Lett.* **97**, 121302 (2006); **97**, 139902(E) (2006).
- [86] S. A. Fulling, *Phys. Rev. D* **7**, 2850 (1973).
- [87] P. C. W. Davies, *J. Phys. A* **8**, 609 (1975).
- [88] W. G. Unruh, *Phys. Rev. D* **14**, 870 (1976).
- [89] R. Schützhold, *arXiv:1110.6064*.
- [90] R. Schützhold, G. Schaller, and D. Habs, *Phys. Rev. Lett.* **100**, 091301 (2008).
- [91] E. Aprile *et al.* (XENON Collaboration), *Phys. Rev. Lett.* **131**, 041003 (2023).
- [92] J. Aalbers *et al.* (LZ Collaboration), *Phys. Rev. Lett.* **131**, 041002 (2023).
- [93] M. S. Turner, *Phys. Rev. Lett.* **60**, 1797 (1988).
- [94] G. Lucente, L. Mastrototaro, P. Carenza, L. Di Luzio, M. Giannotti, and A. Mirizzi, *Phys. Rev. D* **105**, 123020 (2022).
- [95] J. H. Chang, R. Essig, and S. D. McDermott, *J. High Energy Phys.* **09** (2018) 051.
- [96] F. Ertas and F. Kahlhoefer, *J. High Energy Phys.* **07** (2020) 050.
- [97] F. Calore, P. Carenza, M. Giannotti, J. Jaeckel, G. Lucente, and A. Mirizzi, *Phys. Rev. D* **104**, 043016 (2021).
- [98] D. I. Dunskey, L. J. Hall, and K. Harigaya, *J. High Energy Phys.* **04** (2024) 130.
- [99] S. Chakraborty, A. Gupta, and M. Vanvlasselaer, *Phys. Rev. D* **110**, 063032 (2024).
- [100] R. Ruffini, G. Vereshchagin, and S. S. Xue, *Phys. Rep.* **487**, 1 (2010).
- [101] C. Grayson, M. Formanek, J. Rafelski, and B. Mueller, *Phys. Rev. D* **106**, 014011 (2022).
- [102] I. A. Shushpanov and A. V. Smilga, *Phys. Lett. B* **402**, 351 (1997).
- [103] S. P. Klevansky and R. H. Lemmer, *Phys. Rev. D* **39**, 3478 (1989).
- [104] H. Suganuma and T. Tatsumi, *Ann. Phys. (Amsterdam)* **208**, 470 (1991).
- [105] M. D’Elia and F. Negro, *Phys. Rev. D* **83**, 114028 (2011).
- [106] P. Watson and H. Reinhardt, *Phys. Rev. D* **89**, 045008 (2014).
- [107] H. Kim, A. Lenoci, G. Perez, and W. Ratzinger, *Phys. Rev. D* **109**, 015030 (2024).
- [108] C. Beadle, S. A. R. Ellis, J. Quevillon, and P. N. Hoa Vuong, *Phys. Rev. D* **110**, 035019 (2024).
- [109] S. Chakraborty, M. Kraus, V. Loladze, T. Okui, and K. Tobioka, *Phys. Rev. D* **104**, 055036 (2021); **108**, 039903 (E) (2023).
- [110] A. Capolupo, S. M. Giampaolo, and A. Quaranta, *Eur. Phys. J. C* **81**, 1116 (2021).
- [111] B. Heacock, T. Fujiie, R. W. Haun, A. Henins, K. Hirota, T. Hosobata, M. G. Huber, M. Kitaguchi, D. A. Pushin and H. Shimizu *et al.*, *Science* **373**, abc2794 (2021).
- [112] S. Barbosa and S. Fichet, *J. High Energy Phys.* **01** (2025) 021.
- [113] H. Kim and G. Perez, *Phys. Rev. D* **109**, 015005 (2024).
- [114] A. Banerjee, G. Perez, M. Safronova, I. Savoray, and A. Shalit, *J. High Energy Phys.* **10** (2023) 042.
- [115] M. D’Elia, M. Mariti, and F. Negro, *Phys. Rev. Lett.* **110**, 082002 (2013).
- [116] C. M. Kim and S. P. Kim, *Eur. Phys. J. C* **83**, 104 (2023).
- [117] S. A. Olausen and V. M. Kaspi, *Astrophys. J. Suppl. Ser.* **212**, 6 (2014).
- [118] T. Enoto, S. Kisaka, and S. Shibata, *Rep. Prog. Phys.* **82**, 106901 (2019).
- [119] J. Jaeckel and M. Spannowsky, *Phys. Lett. B* **753**, 482 (2016).