

Exploring the dark axion portal in the LUXE-NPOD experiment

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The optical dump at the LUXE experiment has the potential to create a large flux of $\mathcal{O}(\text{GeV})$ photons that can be used to look for new physics when directed at a solid material dump. The LUXE-NPOD extension of LUXE, which focuses the hard photons onto a slab of tungsten, offers two interaction points (laser-electron and photon-tungsten), making it well suited to test theories containing two or more new particles. In this paper, we examine the dark axion portal, a scenario involving both dark photons (DPs) and axionlike particles (ALPs) and their mutual interactions, and its implications on the phenomenology at LUXE-NPOD. We focus on the case of strictly massive DPs. To simulate the spectra of particle populations generated at the electron-laser interaction point, we solve a set of extended 1D cascade equations. We recover a photon spectrum consistent with previous analyses and present previously unconsidered DP and ALP spectra. Further, we derive the overall sensitivity of LUXE-NPOD to various parameters of the new particles and show that it is capable of probing potentially uncharted regions in the dark axion parameter space. For ALPs in the 10–500 MeV mass range, and DPs that are either heavier ($\sim\text{GeV}$) or significantly lighter $\mathcal{O}(10^{-18}\text{--}10^{-13})$ eV, we obtain promising constraints on DP kinetic mixing parameters smaller than 10^{-2} and on $\mathcal{O}(10^{-4}\text{--}10^{-3})$ GeV^{-1} ALP-photon couplings. We find that restrictions on $\mathcal{O}(10^{-3})$ kinetic mixing can be extracted for arbitrarily small DP masses. Our discussion aims to be systematic and demonstrates a practical method of analyzing constraints on multidimensional parameter spaces.

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I. INTRODUCTION

The Standard Model (SM) of particle physics, though remarkably successful, falls short of explaining neutrino oscillations, the observed baryon asymmetry in the Universe, and dark matter. These gaps have motivated extensive theoretical and experimental efforts to explore extensions of the SM.

Many proposed extensions introduce a “hidden” or “dark” sector, composed of new particles that are neutral under the SM gauge group. Interactions between this sector and the SM can arise via effective operators known as “portals.” Two well-studied cases are the dark photon (DP), which couples to the SM photon through a vector portal [1–3], and the axionlike particle (ALP), which couples to SM gauge bosons via an axion portal [4–10]. Both scenarios have been constrained by numerous experiments (see, e.g., Refs. [11–60]).

While most studies focus on either the DP or the ALP in isolation, Ref. [61] showed that interactions between DPs, ALPs, and the SM photon can emerge from a unique mechanism that is not merely a combination of vector and axion portals. In some experiments, this so-called dark axion portal enables novel production and decay channels that may dominate the phenomenology in certain regions of parameter space [62–76].

We investigate the implications of the dark axion portal for the proposed LUXE-NPOD experiment [77,78], which uses a stream of hard photons generated at the LUXE experiment [79–81] to probe physics beyond the SM by placing a material dump and forward detector downstream of the electron-laser interaction point. The original NPOD proposal demonstrated that this configuration is sensitive to new regions of the ALP parameter space. Here, we extend the previous analysis to include DPs and the dark axion portal, examining the case in which all three portals (vector, axion, and dark axion) are active. The dual interaction sites—electron laser and photon dump—make LUXE-NPOD particularly suitable for testing models with multiple new particles. We present a systematic phenomenological study of this setup using methods that can be applied to other multiportal frameworks.

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Our results corroborate previous computations of the hard photon spectrum at LUXE and demonstrate the possibility of producing additional DP and ALP spectra in new physics (NP) scenarios. Moreover, they indicate that LUXE-NPOD is sensitive to regions of dark axion portal parameter space that were not covered by existing single-particle analyses; namely, combinations of kinetic mixings in the 10^{-4} – 10^{-2} range, ALP-photons/DPs couplings in the $\mathcal{O}(10^{-4}$ – $10^{-3})$ GeV $^{-1}$ range, an ALP mass of 10–500 MeV, and a much heavier ($\sim$ GeV) or lighter $\mathcal{O}(10^{-18}$ – $10^{-13})$ eV DP.

This paper is organized as follows. In Sec. II, we introduce the vector, axion, and dark axion portals as a model for potential NP probed at LUXE-NPOD. Section III describes the LUXE-NPOD setup and presents the framework used to analyze and compute the expected NP signals. In Sec. IV, we present the projected sensitivity of LUXE-NPOD across various subplanes of the model's parameter space, highlighting regimes where the dark axion portal leads to novel phenomenology. Section V concludes with a summary of our results. Technical details of cascade simulations and rate calculations are provided in the Appendixes.

II. PORTALS TO HIDDEN SECTORS

The vector portal, axion portal, and dark axion portal terms can be written as

$$\mathcal{L}_{\text{vector portal}} = \frac{\varepsilon}{2} F^{\mu\nu} F'_{\mu\nu}, \quad (1)$$

$$\mathcal{L}_{\text{axion portal}} = \frac{g_{a\gamma\gamma}}{4} a F^{\mu\nu} \tilde{F}_{\mu\nu}, \quad (2)$$

$$\mathcal{L}_{\text{dark axion portal}} = \frac{g_{a\gamma\gamma'}}{2} a F^{\mu\nu} \tilde{F}'_{\mu\nu}, \quad (3)$$

where $F'_{\mu\nu}$ is the DP field strength and $\tilde{F}_{\mu\nu} = \frac{1}{2} \epsilon_{\mu\nu\rho\sigma} F^{\rho\sigma}$. In the above, ε is the kinetic mixing parameter, $g_{a\gamma\gamma}$ is the ALP-photon-photon coupling, and $g_{a\gamma\gamma'}$ is the ALP-photon-DP coupling. The dark axion portal operator in Eq. (3) typically arises from a Kim–Shifman–Vainshtein–Zakharov (KSVZ)-type mechanism [61], which introduces an additional ALP-DP-DP term $\frac{g_{a\gamma\gamma'}}{4} a F^{\mu\nu} \tilde{F}'_{\mu\nu}$ that also affects our phenomenological analysis, as we explain in Sec. III.

It is well known that if the DP is massless it completely decouples from the SM and gives rise to so-called milli-charged particles (see, e.g., Ref. [3]). In this work, we consider strictly massive DPs and do not focus on milli-charged particles. Accordingly, Eqs. (1)–(3) are written in the basis where the DP carries a canonical diagonal mass term. Aside from this restriction, we do not assume a dominant coupling or particular mass hierarchy between the ALP and the DP.

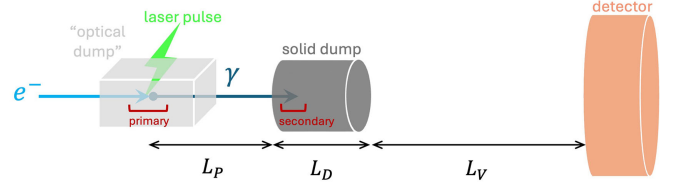


FIG. 1. Schematic of the LUXE configuration and the NPOD extension (inspired by Fig. 2 in Ref. [77]). Electrons collide with a laser pulse (“optical dump”) in the LUXE experiment, generating hard photons which the NPOD extension directs onto a solid dump. A forward detector is stationed after the dump.

III. DARK AXION PORTAL AT LUXE-NPOD

Complete information regarding the configuration of LUXE can be found in Refs. [79,80], and details of the proposed NPOD extension are given in Ref. [77]. We briefly recount the relevant points for our purposes.

LUXE plans to probe QED beyond the Schwinger limit in the strong-field domain where nonperturbative, nonlinear phenomena are relevant [82–87], by aiming an energetic beam of electrons, generated at the European X-ray Free-Electron Laser (XFEL), at a high-intensity laser. The beam-laser collisions result in a collimated flux of $\sim$ GeV photons emerging from the interaction point.

The NPOD proposal suggests directing the stream of photons onto a solid material dump, with a forward detector placed beyond the dump (a depiction of this setup is presented in Fig. 1). In this setup, ALPs can be produced in an elastic Primakoff-like process $\gamma + N \rightarrow N + a$ where photons γ interact with dump nuclei N . After the $\gamma \rightarrow a$ conversion, ALPs that decay into two photons outside the dump and before the detector can be identified by a diphoton signal in the detector. The expected number of signals is computed in Eq. (2) in Ref. [77].

Besides enabling efficient ALP production, the dump is also effective at blocking the SM background from reaching the detector. This fact, combined with appropriate detector technology and signal selection criteria, implies that < 1 diphoton signals from non-ALP sources are expected in one year of data collection. We therefore consider the ALP search to be background free.

A projection of LUXE-NPOD’s sensitivity to the ALP parameters is plotted in Fig. 4 of Ref. [77]. This result demonstrates that LUXE-NPOD has the power to exclude regions of ALP parameter space corresponding to masses of $\mathcal{O}(30\text{--}300)$ MeV and couplings of $\mathcal{O}(1\text{--}300)$ PeV $^{-1}$.

In what follows, we perform a similar analysis of potential NP signals at LUXE-NPOD for a model including DPs, ALPs, and the dark axion portal. We assume the same experimental configuration and technical specifications that are detailed in Ref. [77]. This setup has two separate interaction points—one in the optical dump and another in the solid dump—and we are considering two new particles that interact with the SM as well as with each other.

Consequently, numerous different NP processes may arise. To make sense of the implications on the observed signals, we attempt to work systematically.

A. Allowed processes

We begin by defining some useful terminology. We call the electron-laser collision the “primary” interaction point and the interaction in the solid dump the “secondary” interaction point (see Fig. 1). The production of particles at these points will be referred to as “primary/secondary production.”

Recall that the ALP-only case studied in Ref. [77] considered the conversion of primary photons to secondary ALPs via interactions with the dump nuclei. The tree-level Feynman diagram for this Primakoff-like process can be seen in Fig. 2(a) (left). The inclusion of DPs and the dark axion portal allows for an analogous $\gamma \rightarrow a$ conversion channel mediated by a DP instead of a photon—see Fig. 2(a) (right). Additionally, the optical dump now generates not only photons but also primary DPs, which can undergo similar $\gamma' \rightarrow a$ Primakoff-like conversion involving γ/γ' -mediated tree-level diagrams, or they can pass through the entire dump. There are also various possibilities involving primary ALPs.

We list all the elastic tree-level processes that generate NP particles at LUXE-NPOD to leading order in ε and the ALP couplings. For completeness, we also include DPs/ALPs that are generated from interactions of primary particles with electrons in the dump (instead of just nuclear interactions). Secondary NP production is generated by

$$\text{optical dump} \rightarrow \gamma \rightarrow \text{solid dump} \rightarrow a, \quad (4a)$$

$$\text{optical dump} \rightarrow \gamma \rightarrow \text{solid dump} \rightarrow \gamma', \quad (4b)$$

$$\text{optical dump} \rightarrow \gamma' \rightarrow \text{solid dump} \rightarrow a, \quad (4c)$$

$$\text{optical dump} \rightarrow a \rightarrow \text{solid dump} \rightarrow \gamma', \quad (4d)$$

and for primary NP production, we have

$$\text{optical dump} \rightarrow \gamma' \xrightarrow{\text{solid dump}} \gamma', \quad (5a)$$

$$\text{optical dump} \rightarrow a \xrightarrow{\text{solid dump}} a. \quad (5b)$$

Feynman diagrams corresponding to all of the above processes can be found in Figs. 2 and 3.

After NP production, the propagating DPs/ALPs can decay in the volume between the dump and the detector. A DP may decay to an ALP and a photon (if $m_{\gamma'} > m_a$), registering a monophoton signal. An ALP may decay to two

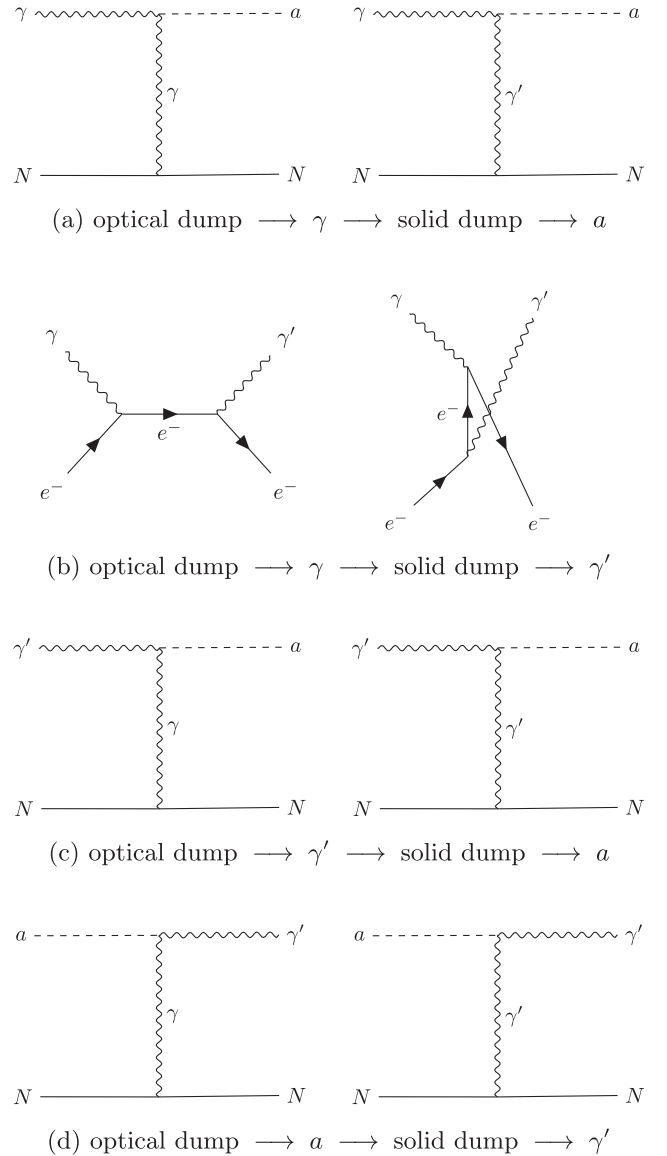


FIG. 2. Tree-level Feynman diagrams for secondary NP production of DPs/ALPs at LUXE-NPOD. The processes are facilitated through interactions with electrons e^- or nuclei N in the solid dump. Each case receives contributions from two mediation channels that interfere with one another. Panels (a)–(d) correspond to the four processes of Eqs. (4a)–(4d), respectively: (a) secondary ALP production from a primary photon; (b) secondary DP production from a primary photon; (c) secondary ALP production from a primary DP; (d) secondary DP production from a primary ALP.

photons or (if $m_{\gamma'} < m_a$) to a DP and a photon, registering either a diphoton or a monophoton signal.

B. Expected number of signals

To find the number of signals we expect at LUXE-NPOD due to NP, we generalize the expression of Eq. (4) in Ref. [77] such that for any particular NP process

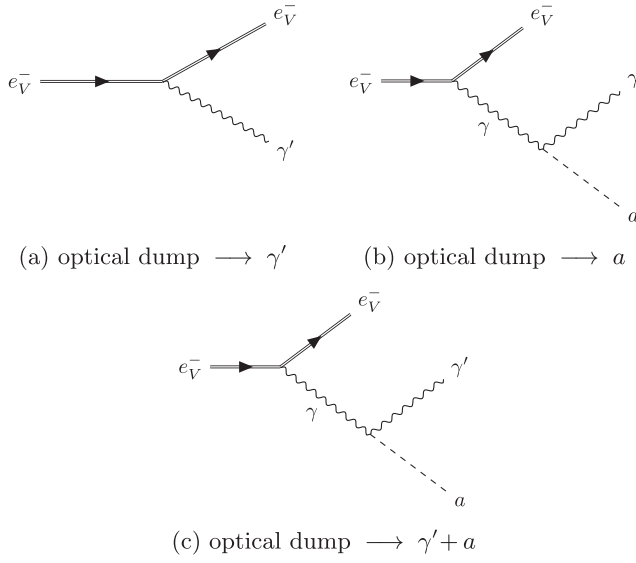


FIG. 3. Tree-level Feynman diagrams for primary NP production of DPs/ALPs at LUXE-NPOD. The double line represents a dressed Volkov-state electron [88]. (a) DP emission, $e_V^- \rightarrow e_V^- \gamma'$; (b) off-shell emission of an ALP together with a photon, $e_V^- \rightarrow e_V^- \gamma a$; (c) off-shell emission of an ALP together with a DP, $e_V^- \rightarrow e_V^- \gamma' a$.

$$\text{optical dump} \longrightarrow Y \longrightarrow \text{solid dump} \longrightarrow X, \quad (6)$$

as in Eqs. (4)–(5),¹ the number of s -photon signals ($s = 1, 2$) seen in the detector is given by

$$N_{sY}^{Y \rightarrow X} = \int dE_Y \frac{d\mathcal{N}_{Y \rightarrow X}}{dE_Y} \mathcal{P}_{X\text{-decay}} \left(\begin{array}{c} \text{after dump} \\ \text{before detector} \end{array} \right) \times \text{BR}(X \rightarrow s\gamma) \Theta(E_{sY} - E_{\text{thresh}}) \mathcal{A}, \quad (7)$$

where $\int dE_Y$ integrates over all energies of primary Y 's generated in the electron-laser collisions, $d\mathcal{N}_{Y \rightarrow X}/dE_Y$ is the differential number of $Y \rightarrow X$ events as a function of the Y energy, $\mathcal{P}_{X\text{-decay}} \left(\begin{array}{c} \text{after dump} \\ \text{before detector} \end{array} \right)$ is the probability that X decays in the detectable volume, $\text{BR}(X \rightarrow s\gamma)$ is the branching fraction of X -decay to an s -photon final state, $\Theta(E_{sY} - E_{\text{thresh}})$ is a step function keeping only photon signals with energies above the detection threshold, and $\mathcal{A}$ is the angular acceptance of the detector. We elaborate on each of these elements below.

The differential number of $Y \rightarrow X$ events is

$$\frac{d\mathcal{N}_{Y \rightarrow X}}{dE_Y} = \mathcal{L}_Y \frac{dN_Y}{dE_Y} \sigma_{Y \rightarrow X}. \quad (8)$$

In the above, we have the effective Y luminosity, given by $\mathcal{L}_Y = N_e N_{\text{BX}} \frac{\rho_W}{A_W m_0} \ell_{\text{eff}}$ —where $N_e = 1.5 \times 10^9$ is the XFEL electron bunch population, $N_{\text{BX}} = 10^7$ is the total number of

bunch crossings in a one-year run of the experiment, $\rho_W = 19.3 \text{ g/cm}^3$ and $A_W = 183.84$ are the density and mass number of tungsten respectively (taken from Ref. [89]), and ℓ_{eff} is the effective luminosity length for Y -dump interactions. For $Y = \gamma$, we take $\ell_{\text{eff}} = \frac{9}{7} X_0$ (since photons lose most of their energy in the first $9/7$ tungsten radiation lengths $X_0 = 0.3456 \text{ cm}$ —see Ref. [89]), but for $Y = \gamma', a$, we have $\ell_{\text{eff}} = L_D = 1 \text{ m}$ (as DPs/ALPs are weakly interacting and not substantially stopped by the dump).

We compute the differential Y flux per incident electron dN_Y/dE_Y using the extended cascade equations of QED [90] (see Appendix A for details). Schematically, the cascade equations are a set of coupled first-order partial differential equations describing the evolution of the spectrum $I_Y(E_Y, t)$ of population Y at time t , i.e.,

$$\frac{\partial I_Y}{\partial t} = \sum_{X, E_X} r_{X \rightarrow Y}(E_Y, E_X, t) I_X(E_X, t) - \sum_{Z, E_Z} r_{Y \rightarrow Z}(E_Z, E_Y, t) I_Y(E_Y, t), \quad (9)$$

where $r_{X \rightarrow Y}(E_Y, E_X, t)$ [similarly $r_{Y \rightarrow Z}(E_Z, E_Y, t)$] describes the rate to produce particle Y (Z) at energy E_Y (E_Z) from processes undergone by particle X (Y) at energy E_X (E_Y). Equation (9) approximates the cascade to occur in a single spatial dimension. This approximation is justified by the near collinearity of the Y particles produced by the electrons and the lack of straggling due to feedback of the nonelectron species. We solve Eqs. (9) numerically to find the spectra of particles $dN_Y/dE_Y = I_Y(E_Y, \infty)$ produced by a beam of 16.5 GeV electrons propagating through the optical dump. For the beams and lasers of interest at LUXE, the spatial extent of the beam and the laser must be taken into account. The electron beam at LUXE is a 3D Gaussian pulse with deviation $5 \mu\text{m}$ in directions transverse to propagation and deviation of $24 \mu\text{m}$ in the direction of propagation. The square of the potential describing the laser profile is spatially distributed according to

$$|A|_{\text{rms}}^2 = \frac{|A|^2 f^2(k \cdot x)}{1 + (z/z_R)^2} \exp\left(-\frac{2r_\perp^2}{w_0^2(1 + (z/z_R)^2)}\right), \quad (10)$$

where $|A|$ is the magnitude of the laser intensity, w_0 is the beam waist (taken to be $10 \mu\text{m}$), $z_R = \pi w_0^2/\lambda$ is the Rayleigh range, $r_\perp$ is the distance perpendicular to the axis of propagation, and z is the coordinate along the axis of propagation. The pulse shape $f(k \cdot x)$ is Gaussian with width 25 fs in LUXE's 45 TW phase-0 run and 120 fs in LUXE's 350 TW phase-1 run. The laser is oriented at 17.2° relative to the electron beam. The optical dump is therefore wide enough for the locally monochromatic approximation to hold [91]. To capture the spatial extent of the beam-laser collision, we discretize the incident beam distribution, simulate the propagation of each point through the spatially varying pulse, and sum the resulting spectra.

¹We accommodate the primary NP production processes in this expression by allowing $Y = X$.

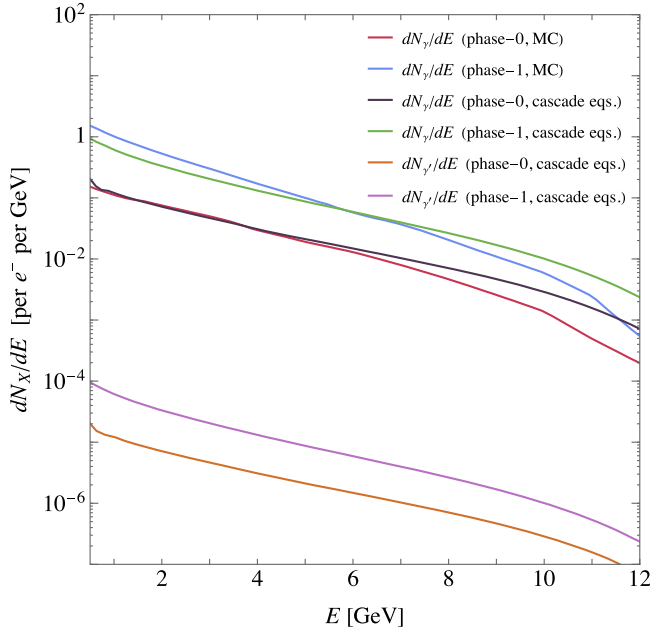


FIG. 4. Particle spectra generated in the electron-laser interaction. Numerically computed spectra of visible and dark photons produced during phase-0 (black and orange, respectively) and phase-1 (green and pink, respectively) of the LUXE experiment after collision with a high-intensity laser, assuming kinetic mixing $\varepsilon = 0.01$ and a DP mass much smaller than that of the electron $m_{\gamma'} \ll m_e$. The photon spectra computed using Monte-Carlo methods are also plotted.

Figure 4 shows the results of our computation for phase 0 and phase 1, along with the photon spectrum computed in Ref. [77] using strong-field QED Monte-Carlo simulations. We see that for small kinetic mixing ε , the straggling of electrons due to DP production is small compared to the straggling due to photons, so the primary DP flux is approximately equal to the photon spectrum suppressed by the kinetic mixing squared $dN_{\gamma'}/dE \approx \varepsilon^2 dN_{\gamma}/dE_{\gamma}$. The primary production of ALPs is dominated by the $e_{\bar{\nu}} \rightarrow e_{\bar{\nu}} \gamma^{(\prime)} a$ processes depicted in Figs. 3(b) and 3(c) (see Appendix A for details); we therefore take a similar crude estimate for the ALP spectrum using the production rate ratio $\Gamma_{e_{\bar{\nu}} \rightarrow e_{\bar{\nu}} a} / \Gamma_{\gamma}$ calculated in Ref. [77] (Fig. 10), i.e., $dN_a/dE_a \approx 4 \times 10^{-16} [(g_{a\gamma\gamma}/\text{TeV}^{-1})^2 + (g_{a\gamma\gamma'}/\text{TeV}^{-1})^2]$. Note that both of these spectra are only present for sufficiently light DPs/ALPs $m_{\gamma',a} \lesssim \mathcal{O}(100)$ keV, since the primary production of NP particles is greatly suppressed above this mass—see Fig. 8 (bottom).

Finally, $\sigma_{Y \rightarrow X}$ is the inclusive cross section, computed via numerical integration over phase space. The differential cross sections $d\sigma_{Y \rightarrow X}/dt$ for each process in Eqs. (4)–(5), where t is the square of 4-momentum transfer, can be found in the Appendix B.

Next, the probability that X decays in the detectable volume depends on the point z where it is produced. Assuming the distance from the optical dump to the solid

dump is $L_P = 13$ m and the distance from the end of the dump to the detector is $L_V = 2.5$ m (recall the tungsten dump has a length $L_D = 1$ m), we may write

$$\mathcal{P}_{X\text{-decay}} = \int_0^{L_P+L_D+L_V} dz p_{X\text{-prod}}(z) \times \left(e^{-\frac{L_P+L_D-z}{L_X}} - e^{-\frac{L_P+L_D+L_V-z}{L_X}} \right), \quad (11)$$

where the X decay length L_X is defined by $L_X = \tau_X \sqrt{E_X^2 - m_X^2}/m_X$ (τ_X being the X lifetime, E_X being its energy, and m_X being its mass) and $p_{X\text{-prod}}(z)$ is the probability density that X is produced at a point $z \in [0, L_P + L_D + L_V]$.

If X is generated in primary production, i.e., at the primary interaction point $z = 0$, we take² $p_{\text{prim}X\text{-prod}}(z) = 2\delta(z)$. If X is generated in secondary production triggered by a photon, we take $p_{\gamma X \rightarrow \text{prod}}(z) = \delta(z - L_P)$. This is again due to the fact that photons lose most of their energy within $\frac{9}{7}X_0 \approx 0.44$ cm, which is orders of magnitude shorter than the full length of the tungsten dump $L_D = 1$ m. Therefore, any secondary NP particles originating from a primary photon are effectively created at a single point at the start of the dump. Lastly, if X is generated in secondary production triggered by a DP or an ALP—which are weakly interacting—we take a simple uniform distribution over the length of the dump $p_{\text{NP} \rightarrow X\text{-prod}}(z) = \Theta(z - L_P)\Theta(L_P + L_D - z)/L_D$.

We note that in computing the particle lifetime τ_X we only take into consideration couplings and interactions that were explicitly presented in Sec. II (namely, SM content and DP/ALP interactions through the vector, axion, and dark axion portals). Contributions of DP/ALP interactions purely within the hidden sector are ignored.

The branching fractions depend on the mass hierarchy in our dark axion model. ALPs can always decay to two photons $a \rightarrow \gamma\gamma$ and, depending on their mass, may also decay to a photon and DP $a \rightarrow \gamma\gamma'$ or to two DPs $a \rightarrow \gamma'\gamma'$. The widths of these decays are given, for example, in Ref. [61], and we can use them to compute $\text{BR}(a \rightarrow s\gamma)$ for $s = 1, 2$. DPs, depending on their mass, may decay to charged leptons, hadrons, or a photon and an ALP. We use the decay widths for $\gamma' \rightarrow \text{SM}$ and $\gamma' \rightarrow a\gamma$ in Refs. [3,61], respectively, to compute $\text{BR}(\gamma' \rightarrow 1\gamma)$. Here, we neglect any decay channels of DPs/ALPs into pure hidden sector states.

For the detection factors in Eq. (7), we assume a photon energy detection threshold of $E_{\text{thresh}} = 0.5$ GeV and take an angular acceptance of $\mathcal{A} = 1$ for simplicity. Total angular acceptance may seem unjustified; very light X s should be highly boosted and produce very collinear decay remnants that are hard to resolve spatially, and very heavy X s might not be boosted enough so their decay remnants could be ejected at an angle too wide to land inside the disk-shaped

²We use the convention $\int_0^\infty dz \delta(z) = 1/2$.

detector with a radius of $R = 1$ m. However, as we will see, NP events at LUXE-NPOD are dominated by $\mathcal{O}(1)$ GeV ALPs ($X = a$) with masses of $\mathcal{O}(10\text{--}350)$ MeV coupled to DPs that are either heavier than the ALP or in the ultralight regime $m_{\gamma'} \ll m_a$. This implies angular differences of $\mathcal{O}(10\text{--}300)$ mrad for two final photons (or a photon and an ultralight DP). Such angular resolution is achievable with existing detector concepts [92], and the widest angles are roughly compatible with the proposed geometry of the LUXE-NPOD configuration [which has a detection coverage of $\arctan(R/L_V) \approx 400$ mrad]. We therefore claim that the angular aspect of detection has an $\mathcal{O}(1)$ effect in most regions of the parameter space, with potentially stronger effects near the high-mass boundary. In this regime, a dedicated angular acceptance simulation may modify our projected reach.

Putting the pieces together, the total number of s -photon signals ($s = 1, 2$) we expect at LUXE-NPOD is computed via

$$\begin{aligned} N_{s\gamma} &= \sum_{Y,X} N_{s\gamma}^{Y \rightarrow X} \\ &= \sum_{Y,X} \int dE_Y \int dt \int dz \mathcal{L}_Y \frac{dN_Y}{dE_Y} \frac{d\sigma_{Y \rightarrow X}}{dt} P_{X\text{-prod}}(z) \\ &\quad \times \left(e^{-\frac{L_P + L_D - z}{L_X}} - e^{-\frac{L_P + L_D + L_V - z}{L_X}} \right) \\ &\quad \times \text{BR}(X \rightarrow s\gamma) \Theta(E_{s\gamma} - E_{\text{thresh}}) \mathcal{A}. \end{aligned} \quad (12)$$

It is worth stressing that the momentum transfer t depends on the primary energy E_Y , which is why this expression is not

merely a simple product as suggested in Eq. (7)—rather, it is a kind of convolution of all the factors we described.

C. Background estimation

Background signals from SM processes are identical to those presented in Ref. [77]. There are two main sources of background³—photons (from EM interactions near the end of the dump) and neutrons (produced in EM showers in the dump) that may be misidentified by the detector as photons.

The total number of $s\gamma$ background signals can be estimated by

$$\begin{aligned} N_{s\gamma}^{\text{bkg}} &= \sum_b N_{b \rightarrow s\gamma}^{\text{bkg}} \\ &= \sum_b N_{\text{BX}} \mathcal{P}_{b \rightarrow s\gamma}^{\text{bkg}} R_{\text{sel}}^{b \rightarrow s\gamma}, \end{aligned} \quad (13)$$

where the sum is over all background sources that can trigger an $s\gamma$ signal ($b = \gamma, n, 2\gamma, \gamma + n, 2n$), $N_{b \rightarrow s\gamma}^{\text{bkg}}$ is the number of $s\gamma$ background signals resulting from the specific source b , $\mathcal{P}_{b \rightarrow s\gamma}^{\text{bkg}}$ is the probability that process b is detected as an $s\gamma$ signal, and $R_{\text{sel}}^{b \rightarrow s\gamma}$ is a rejection factor achieved by implementing a set of selection criteria on the background processes b that trigger an $s\gamma$ signal. Note that the probability depends on $\mu_{\gamma(n)}$, the rate of background photons (neutrons) that arrive at the detector with energy above the detection threshold, and also on the neutron misidentification rate $f_{n \rightarrow \gamma}$ (which is highly sensitive to the detector technology used).

There are two possibilities to obtain a monophoton background signal:

- (i) detector triggered by a single background photon and 0 neutrons, with probability $\mathcal{P}_{\gamma \rightarrow 1\gamma}^{\text{bkg}} = \Pi(\mu_\gamma, 1) \times \sum_{k_n=0}^{\infty} \Pi(\mu_n, k_n) \mathcal{B}(0, k_n, f_{n \rightarrow \gamma}) = (\mu_\gamma e^{-\mu_\gamma}) (e^{-\mu_n f_{n \rightarrow \gamma}})$, or
- (ii) detector triggered by a single neutron and 0 background photons, with probability $\mathcal{P}_{n \rightarrow 1\gamma}^{\text{bkg}} = \Pi(\mu_\gamma, 0) \times \sum_{k_n=1}^{\infty} \Pi(\mu_n, k_n) \mathcal{B}(1, k_n, f_{n \rightarrow \gamma}) = (e^{-\mu_\gamma}) (\mu_n f_{n \rightarrow \gamma} e^{-\mu_n f_{n \rightarrow \gamma}})$,

where Π and $\mathcal{B}$ are the probability mass functions of the Poisson and Binomial distributions, respectively. The background rates μ_γ, μ_n were estimated in Ref. [77] with a GEANT4 simulation. These rates are smaller and softer during phase-0 of LUXE, but phase-1 levels are conservatively assumed also for phase-0—that is, we use the values $\mu_\gamma = 0.013$ and $\mu_n = 10$ for both phases. The resulting estimates for monophoton background events are

$$N_{\gamma \rightarrow 1\gamma}^{\text{bkg}} = 1.3 \times 10^5 e^{-10 f_{n \rightarrow \gamma}} R_{\text{sel}}^{\gamma \rightarrow 1\gamma} \quad (14)$$

$$N_{n \rightarrow 1\gamma}^{\text{bkg}} = 9.9 \times 10^7 f_{n \rightarrow \gamma} e^{-10 f_{n \rightarrow \gamma}} R_{\text{sel}}^{n \rightarrow 1\gamma}. \quad (15)$$

Estimates for the number of diphoton background events are the same as presented in Eqs. (3)–(5) of Ref. [77].

As noted above, the achievable neutron misidentification rate, $f_{n \rightarrow \gamma}$, is determined primarily by the detector technology. Adopting detector specifications of Ref. [77], we expect $f_{n \rightarrow \gamma} \leq 10^{-3}\text{--}10^{-4}$. In addition, time resolutions of $\mathcal{O}(10\text{--}100)$ ps, position resolutions of $\mathcal{O}(100)$ μm , angular resolutions of $\mathcal{O}(100)$ mrad, and energy resolutions at the level of a few percent—all achievable with existing technologies [92–96]—enable precise reconstruction of diphoton kinematics; this, in turn, allows for strong rejection of diphoton background events $R_{\text{sel}}^{\gamma \rightarrow 2\gamma} \leq 10^{-3}\text{--}10^{-4}$. Overall, we are able to reduce the number of diphoton background events to < 1 (see Eqs. (3)–(5) in Ref. [77]).

³This does not include background of charged particles, as these can be deflected away from the detector with a magnetic field.

For monophoton signals the situation is more challenging. Although time-of-flight information can be used to suppress single-neutron backgrounds, it does not provide discrimination against irreducible single-photon backgrounds. Moreover, in contrast to the diphoton case, there are no additional kinematic handles to exploit. As a result, the selection efficiency $R_{\text{sel}}^{\gamma \rightarrow 1\gamma}$ does not yield a meaningful reduction in event count, and the monophoton background is therefore set by Eq. (14), $N_{1\gamma}^{\text{bkg}} \sim \mathcal{O}(10^5)$.

D. Existing constraints

To properly assess the promise of LUXE-NPOD in probing the dark-axion parameter space, our calculations should be compared to existing bounds on DP and ALP parameters. The rigorous and correct way to do so requires implementation of a proper recasting scheme, such as that developed in Ref. [97]. However, as demonstrated in this manuscript, quantifying the effects of a nonminimal SM extension (that involves both DPs and ALPs) on a given experiment requires a thorough NP treatment and extensive computations, meaning that recasting is much less straightforward. A proper treatment would be akin to replicating the entire scope of the present work dozens of times. While analyzing the dark axion portal in other experiments is an intriguing possibility for future work, for now we simply use the existing bounds “as is,” without recasting. Although this does not provide an accurate indication of LUXE-NPOD’s absolute improvement over existing experiments, it does provide some baseline against which to compare our results.

We consult Ref. [98] for existing constraints on DP masses below ~ 100 keV (with updated versions of these in Ref. [99]). For limits on DP masses above ~ 2 MeV, we refer to Ref. [97]. Explicitly, we plot the LUXE-NPOD sensitivity against limits from tests of the Coulomb interaction, such as Cavendish-like experiments [11], the Plimpton–Lawton experiment [12], atomic spectroscopy [13], atomic force microscopy [12], and the magnetic fields of Earth [14] and Jupiter [15]; constraints on distortions to the cosmic microwave background obtained by COBE and FIRAS [16]; light-shining-through-walls experiments at ALPS [17], Spring-8 [18], UWA [19], ADMX [20], and CROWS [21]; the helioscopes CAST [22] and SHIPS [23]; the neutrino reactor experiment TEXONO [24]; beam dump experiments such as CHARM [25], E137, E141, and E774 [26–29]; anomalous magnetic moment discrepancies [30]; the fixed target experiments LSND [31], NA48/2 [32], WASA [33], APEX [34], and MAMI [35]; and the collider experiments *BABAR* [36], *KLOE* [37], and *FASER* [38]. Note that we ignore limits that rely on DP dark matter properties such as a particular relic abundance, or those obtained from astrophysical sources, e.g., by DP saturation of energy losses.

For ALPs, we similarly draw from Ref. [99], again omitting bounds that attribute dark matter properties to ALPs and also some that are obtained from stellar cooling

constraints. Explicitly, we plot the LUXE-NPOD sensitivity against limits from the CAST helioscope [39,40]; fits of solar data [41]; analyses of globular clusters [42,43]; magnetic birefringence measurements at PVLAS [44]; shining-through-wall experiments by ALPS [17] and EuXFEL [45]; the SAPPHIRES collaboration photon-photon scattering probe [46]; collider experiments including those recast in Ref. [47], NOMAD [48], LEP [49], Belle II [50], BES III [51], OPAL [52], CMS [53], ATLAS [54], and LHC (pp) [52]; beam dump experiments such as those in Refs. [47,55,56], NA64 [57], CHARM [58], and MiniBooNE [59]; and the fixed-target experiment at PrimEx [60]. Several additional restrictions on the ALP-photon-photon coupling have been derived from astrophysical sources for $\sim \text{MeV}$ ALPs [100–109], but their robustness depends on details of the dark sector interactions, and they cease to apply in certain chameleon-like models (see, e.g., the discussion in Ref. [110]).

Lastly, we note that some studies have also derived bounds on the dark axion portal [62–76]. However, most of these extract limits on $g_{a\gamma\gamma'}$ by investigating simplifying benchmark scenarios (for example, Refs. [62,66,71] derive constraints on the $(m_{\gamma'}, g_{a\gamma\gamma'})$ plane under the assumption that the ALP is effectively massless and that the dark axion portal is dominant over the other portals), so we do not strictly adhere to them.

IV. PROJECTED SENSITIVITY

From Sec. III C, we see that the diphoton signals are effectively background free. For the monophoton channel, the large background count $N_{1\gamma}^{\text{bkg}} \sim \mathcal{O}(10^5)$ places us in the Gaussian regime where the asymptotic limit of a profile likelihood ratio test yields a 2σ exclusion criterion of $N_{1\gamma} = 2\sqrt{N_{1\gamma}^{\text{bkg}} + (\delta N_{1\gamma}^{\text{bkg}})^2}$, where δ is the fractional systematic uncertainty on the monophoton background. The statistics-only ($\delta = 0$) 95% C.L. thresholds are

$$N_{1\gamma}|\delta=0 = \mathcal{O}(700), \quad N_{2\gamma} = 3, \quad (16)$$

for monophoton signals and diphoton signals, respectively. Including a benchmark $\delta = 1\%(5\%)$ does not affect the diphoton condition but raises the monophoton criterion to $N_{1\gamma} \approx 2 \times 10^3 (10^4)$. The exclusion regions we present below are statistics-only projections; a full nuisance parameter treatment with dedicated detector-level background simulation is beyond the scope of this work, and estimating the deterioration of LUXE-NPOD’s dark axion reach with varying δ is nontrivial. We nonetheless verified numerically that a $\delta = 5\%$ fractional systematic uncertainty degrades our projected reach by at most a factor of $\mathcal{O}(2)$.

We use the above thresholds to impose limits on the 6 dark axion parameters—two masses m_a , $m_{\gamma'}$, a kinetic mixing strength ϵ , and three ALP couplings $g_{a\gamma\gamma}$, $g_{a\gamma\gamma'}$, $g_{a\gamma'\gamma'}$. In what follows, we methodically evaluate the sensitivity of

LUXE-NPOD to the dark axion portal in all two-dimensional parameter subplanes. Since the ALP-DP-DP coupling is a pure hidden sector parameter, we will not plot $g_{a\gamma'\gamma'}$. This leaves 10 unique pairs of parameters. For each choice, we fix the remaining four parameters and numerically solve Eq. (16) for the chosen pair. We repeat this process while varying the fixed parameters to test a wide array of regimes; for instance, we always test (at least) the three mass hierarchies of heavier DPs $m_{\gamma'} > m_a$, heavier ALPs $m_{\gamma'} < m_a$, and comparable masses $m_{\gamma'} \approx m_a$.

Unless otherwise specified, the plots in this section display merged results from both diphoton and monophoton signals. We implement a consistent color scheme throughout, where the reach of the experiment is indicated by a solid blue line for LUXE's phase 0 and by a solid black line for LUXE's phase 1 (the 95% C.L. regions are emphasized with a transparent coloring). Gray regions indicate parameter space excluded by single-particle analyses (see the discussion in Sec. III D).

Our results, after exhaustively testing a range of parameter values, indicate that LUXE-NPOD can probe previously unexplored regions of the parameter space when $m_a = \mathcal{O}(10\text{--}500)$ MeV, in two distinct regimes—heavier DPs ($m_{\gamma'} > m_a$) and ultralight DPs ($m_{\gamma'} \ll m_a$).

A. Heavier dark photons $m_{\gamma'} > m_a$

As a consistency check, we consider the DP decoupling limit. Indeed, when we decouple the DP from the ALP, we recover the same sensitivity obtained in Ref. [77]. Moving away from this limit, we introduce a nonzero dark axion portal $g_{a\gamma'\gamma'} > 0$. Despite this additional interaction, LUXE-NPOD's reach in the $(m_a, g_{a\gamma\gamma})$ plane remains essentially unchanged, as shown in Fig. 5 (second row, left). This behavior is expected: the ALP-photon-DP interaction only adds a γ' mediation channel for Primakoff-like ALP production [see Fig. 2(a)], yet this channel is subdominant due to the large DP mass.

Beyond reproducing the same $(m_a, g_{a\gamma\gamma})$ sensitivity as in Ref. [77], we obtain novel reach in other parameter subplanes. One example of this, for the $(m_{\gamma'}, g_{a\gamma\gamma})$ plane, is presented in Fig. 5 (third row, right). After fixing the ALP mass $m_a = 100$ MeV, we see that the experiment is sensitive to $g_{a\gamma\gamma} \sim \mathcal{O}(1\text{--}100)$ PeV^{-1} . This is consistent with our discussion above—heavier DPs do not have much impact on the phenomenology, so the result does not exhibit strong $m_{\gamma'}$ dependence. Indeed, the extremal $g_{a\gamma\gamma}$ values in this curve are exactly equal to the intersections of a constant $m_a = 100$ MeV line with the previous curve we discussed (second row, left). When the DP becomes lighter than the ALP $m_{\gamma'} < m_a$, the ALP lifetime sharply decreases as more ALP decay channels (like $a \rightarrow \gamma\gamma'$ or $a \rightarrow \gamma'\gamma'$) become possible; hence, ALPs with the same m_a and $g_{a\gamma\gamma}$ that previously decayed outside the tungsten dump are now too short lived to escape it, and all sensitivity is lost.

A representative sample of plots for all 10 parameter pairings, in the case of heavier DPs, is shown in Fig. 5. All plots exhibit similar characteristics to the two examples discussed above. Larger m_a , $g_{a\gamma\gamma}$ describe short-lived ALPs that decay in the dump, smaller m_a , $g_{a\gamma\gamma}$ describe long-lived ALPs that escape the detector, and LUXE-NPOD is sensitive to a particular band of intermediate values. Sensitivity to ε and $g_{a\gamma'\gamma'}$ is minimal to none, since the DP is essentially decoupled from the ALP when $m_{\gamma'} > m_a$.

B. Ultralight dark photons $m_{\gamma'} \ll m_a$

The second case where LUXE-NPOD exhibits competitive reach is when DPs are “ultralight.” We refer to this limit symbolically as $m_{\gamma'} \ll m_a$, though it is important to stress that we are referring here specifically to DPs that are light enough to be produced in the electron-laser collisions $m_{\gamma'} \ll \mathcal{O}(150)$ keV [which incidentally means they are also much lighter than $\mathcal{O}(10\text{--}500)$ MeV ALPs].

Three distinctions from the previous case affect the phenomenology. One is that the γ' -mediated Primakoff-like ALP production from Fig. 2(a) is no longer suppressed by a large DP mass; another is that we have the possibility of generating primary DPs, which may contribute to both primary and secondary NP production; the third is that ALPs can now decay into DPs and not just photons (eliciting all the implications this has on the ALP lifetime and branching fractions).

With the previous distinctions in mind, we reexamine the $(m_a, g_{a\gamma\gamma})$ plane in the ultralight DP limit. For a given ALP mass m_a and ALP-photon-photon coupling $g_{a\gamma\gamma}$, the possibility of $a \rightarrow \gamma'$ decays implies a shorter ALP lifetime compared to the ALP-only model (in which ALPs have just a single decay channel, into two photons). Therefore, the detectable volume outside the dump and before the detector corresponds to a downward skewed version of the standard-shaped beam dump projections (like those in Ref. [77] and Fig. 5 of this work). That is, ALPs we might hope to detect are ones with smaller $g_{a\gamma\gamma}$ relative to the previous case. This behavior can be seen in Fig. 6 (second row, left). The existence of $a \rightarrow \gamma'$ decays means we lose a fraction of the diphoton signals we had in the ALP-only model, but this is partially compensated for by our inclusion of monophoton signals in the analysis. We observe this explicitly in Fig. 7, where we separate LUXE-NPOD's sensitivity in the $(m_a, g_{a\gamma\gamma})$ plane into the two distinct decay modes that enable said reach—sensitivity from diphoton signals is depicted with a transparent coloring, and sensitivity from monophoton signals is depicted with diagonal hatch lines. Interestingly, in the $g_{a\gamma\gamma} \lesssim \text{PeV}^{-1}$ region of the plot, it is the ALP production described by Eq. (4c), triggered by primary DPs (rather than photons), that provides the most substantial contribution to the observed signal. Thus, we see that genuinely novel NP production processes facilitate the novel reach in this mass regime.

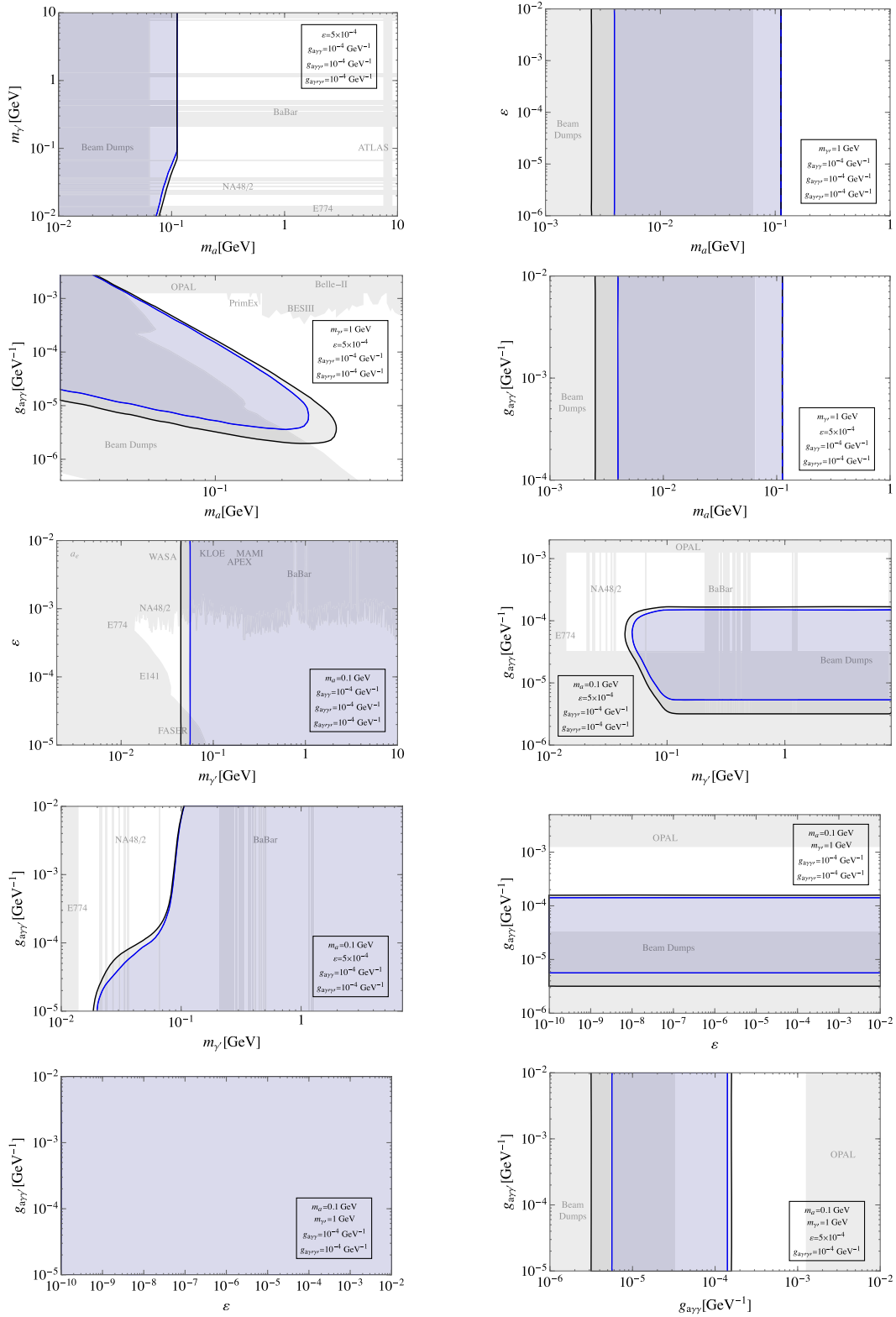


FIG. 5. Projected reach of LUXE-NPOD in the dark axion parameter space for heavy DPs $m_{\gamma'} > m_a$.

Let us also inspect the results obtained in the $(m_{\gamma'}, \epsilon)$ plane, shown in Fig. 6 (third row, left). It seems that LUXE-NPOD has the power to probe the kinetic mixing parameter

ϵ down to arbitrarily low DP masses, without losing any sensitivity. This is in stark contrast to virtually all existing limits on the $(m_{\gamma'}, \epsilon)$ parameters, which generically exhibit

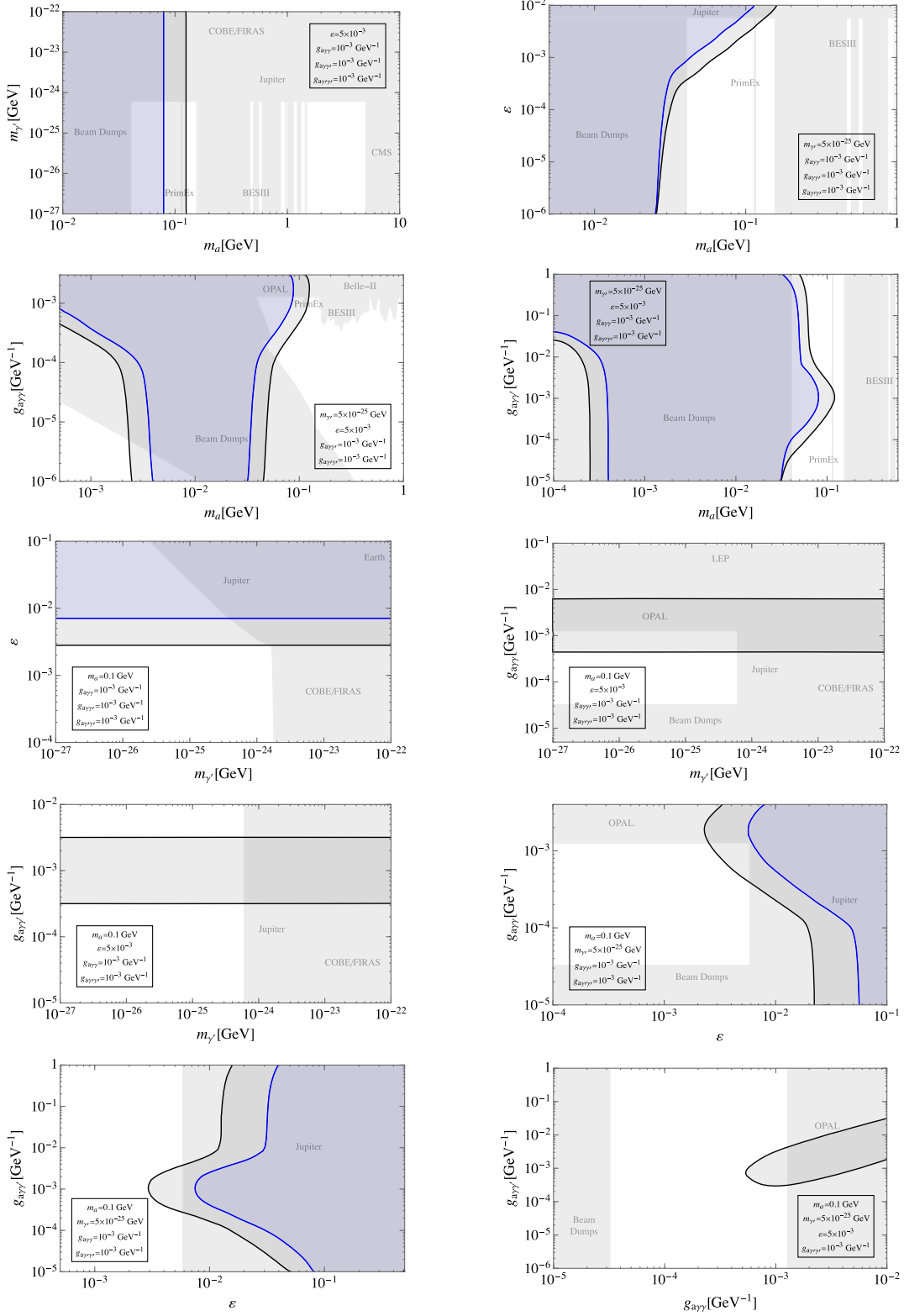


FIG. 6. Projected reach of LUXE-NPOD in the dark axion parameter space when $m_{\gamma'} \ll m_a$.

a maximum sensitivity at some mass scale below which their reach diminishes according to a power law (see, e.g., Fig. 1 in Ref. [98]). The generic behavior of these existing

limits is well understood theoretically; for example, it is known that DPs modify the Coulomb potential $V(r) = \alpha/r + \epsilon^2 \alpha e^{-m_{\gamma'} r}/r$ (see Refs. [1,111]), resulting in an

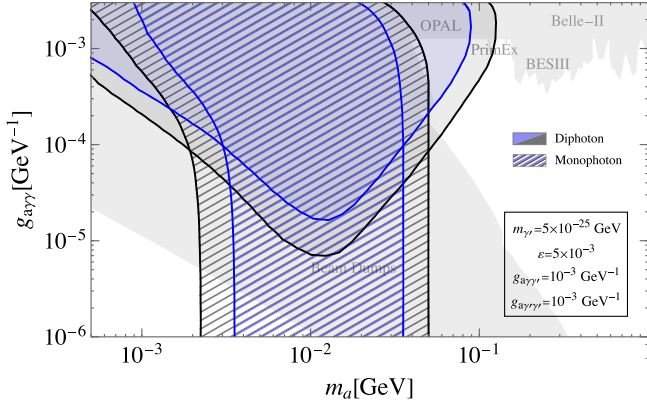


FIG. 7. Projected reach of LUXE-NPOD in the $(m_a, g_{a\gamma\gamma})$ plane when $m_{\gamma'} \ll m_a$, separated into contributions from diphoton signals (solid semitransparent coloring) and monophoton signals (diagonal-hatch-line filling).

appropriate correction $\Delta V \sim \epsilon^2 m_{\gamma'}^2 a r / 2$ when the DP is much lighter than the typical scale $1/r$ of the experiment. The reason we do not observe this power law in LUXE-NPOD is because we are not probing a model with just the DP as a new particle, but a system of DPs *and* ALPs that mutually interact—enabling us to probe DP properties through ALP decays (and vice versa). Support for this claim comes from the fact that the dominant driver of LUXE-NPOD’s reach in the $(m_{\gamma'}, \epsilon)$ plane is the chain $e^- \bar{\nu} \gamma' \Rightarrow \gamma' N \Rightarrow a N \Rightarrow a \rightarrow \gamma\gamma$. In the ultralight DP regime, the contribution of this process to the number of signals scales as

$$\underbrace{\epsilon^2}_{\gamma/\text{flux}} \times \underbrace{g_{a\gamma\gamma'}^2 [1 + \mathcal{O}(m_{\gamma'}^2/|t|)]}_{\gamma t \rightarrow a \text{ conversion}} \times \underbrace{\mathcal{P}_{a\text{-decay}}}_{\text{decay in volume}} \times \underbrace{\mathcal{O}(1)}_{a \rightarrow \gamma\gamma\text{BR}}.$$

Importantly, the expression above does not scale with $m_{\gamma'}$ to leading order ($\mathcal{P}_{a\text{-decay}}$ is only sensitive to the ALP parameters $m_a, g_{a\gamma\gamma}, g_{a\gamma\gamma'}, g_{a\gamma'\gamma'}$ in this regime), unlike the observables of dedicated DP-only searches. This is because the $\gamma' N \rightarrow a N$ conversion remains finite as $m_{\gamma'}$ decreases rather than vanishing as DP-only effects do—a direct consequence of the ALP-photon-DP coupling $g_{a\gamma\gamma'}$.

Finally, just like in the heavy-DP case, there are other pairs of parameters that can be constrained by LUXE-NPOD in the limit of ultralight DPs. Figure 6 shows a representative sample of the projected reach in each of the 10 parameter planes, in this regime. Although there are some planes to which LUXE’s phase 0 is not sufficiently sensitive, we see that phase 1 of the experiment has the potential to probe new territory in all 10 parameter sub-planes. We observe the same trend seen in the heavy-DP case whereby m_a and $g_{a\gamma\gamma}$ affect the ALP lifetime, but this time the results also noticeably depend on ϵ and $g_{a\gamma\gamma'}$, as the DP is no longer decoupled from the ALP. Smaller ϵ means fewer primary DPs are produced, leading to a smaller

number of $\gamma' \rightarrow a$ Primakoff-like conversions in the dump and hence fewer detection events. Regarding $g_{a\gamma\gamma'}$, this parameter determines the magnitude of the $\gamma' \rightarrow a$ conversion cross section, so the number of such conversions (and consequently the number of detected signals) diminishes as $g_{a\gamma\gamma'}$ gets smaller.

V. CONCLUSIONS

In this work, we analyzed the implications of the dark axion model on the observed signals expected at the LUXE-NPOD experiment. The model, which incorporates DPs, ALPs, and their mutual interactions, can be effectively probed by this experiment due to its dual interaction points that allow for novel processes involving both particles.

Our cascade simulations yielded a primary photon spectrum that closely matches previous works and confirmed that primary DP/ALP spectra are well approximated by simple coupling suppression. We showed that LUXE-NPOD is capable of probing novel regions in the six-dimensional parameter space of the dark axion model when the ALP mass falls within the 10–500 MeV range. We are potentially able to place competitive constraints when the DP mass is either $\gtrsim \text{GeV}$ or $\ll \text{eV}$, the kinetic mixing is in the 10^{-4} – 10^{-2} range, and the ALP couplings to photons/DPs are $\mathcal{O}(10^{-4}$ – $10^{-3}) \text{ GeV}^{-1}$. Notably, we found that the existence of a dark axion portal allows for the probing of kinetic mixings for arbitrarily low DP masses, unlike the generic behavior expected from DP-specific searches.

Our analysis demonstrated a systematic approach for constraining a model comprising multiple new particles and several free parameters. The spirit of this method can be replicated and applied to other models that include several NP processes and more than two parametric degrees of freedom.

It is important to note that throughout this work we did not allow for direct decays of DPs/ALPs into solely hidden sector particles. These could have significant effects on particle lifetimes and the expected number of observed signals, though we are currently unaware of a well-established method to quantify such decays based on existing observations.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: CASCADE EQUATIONS IN 1D

To obtain the spectra dN_Y/dE_Y of particles $Y = \gamma, \gamma', a$ emanating from the primary interaction point, we solve the cascade equations for the particle spectra propagating through the laser pulse. To compute transition rates, we work in the framework of nonlinear QED, where the Gaussian laser pulse is treated as a slowly varying background field. The background field A_μ stimulates nonlinear Compton-like scattering (NCS) and nonlinear pair production (NPP). A_μ is taken to be circularly polarized with 4-momentum k_μ , that is,

$$A^\mu(x) = a_1^\mu \cos(kx) + a_2^\mu \sin(kx), \quad (\text{A1})$$

where $a_1^2 = a_2^2 = A^2$ and $a_1 \cdot a_2 = a_1 \cdot k = a_2 \cdot k = 0$.

Electrons propagating through the laser pulse lose energy by emitting photons or NP particles via NCS or

off-shell processes (see diagrams in Fig. 3). The emitted photons, if of high enough energy, may then produce electron-positron pairs via NPP. NP particles may also produce electron-positron pairs (or other combinations of particles), but this effect is small and can be neglected. The interplay of these processes results in a cascade shower of particles. The intensity of the background field controls the rates of these processes and therefore the cascade dynamics. For $\mathcal{O}(1)$ values of the quantum nonlinearity parameter $\eta = e|A|/m_e$ (here, $|A| = \sqrt{|A^2|}$ is the strength of the background field), such as those attainable at LUXE, an electron beam will lose most of its energy while propagating through the dump by producing hard photons and potentially NP particles.

Since the energy of the electron beam E_0 is much larger than the electron mass ($E_0 = 16.5 \text{ GeV} \gg m_e$), the cascade products are emitted through a narrow cone approximately collinear with the incident electrons. The evolution of the particle densities $I_X(E, t)$, where $X = e^-, e^+, \gamma, \gamma', a$, can therefore be effectively described by the 1D cascade-shower equations of QED, extended to include NP particle populations [112]

$$\begin{aligned} \frac{\partial I_\gamma(E_\gamma, t)}{\partial t} = & - \int_{m_e}^{E_\gamma - m_e} d\tilde{E}_e I_\gamma(E_\gamma, t) \frac{dP_{\gamma \rightarrow ee}(\tilde{E}_e, E_\gamma, t)}{d\tilde{E}_e} + \int_{m_e + E_\gamma}^{E_0} d\tilde{E}_e I_e(\tilde{E}_e, t) \frac{dP_{e \rightarrow e\gamma}(E_\gamma, \tilde{E}_e, t)}{d\tilde{E}_e} + \int_{m_e + E_\gamma + m_a}^{E_0} d\tilde{E}_e I_e(\tilde{E}_e, t) \\ & \times \left[\left(\frac{dP_{e \rightarrow e\gamma a}(E_\gamma, \tilde{E}_a, \tilde{E}_e, t)}{d\tilde{E}_a} \right)_{\tilde{E}_a = \tilde{E}_e - m_e - E_\gamma} + \int_{m_a}^{\tilde{E}_e - m_e - E_\gamma} d\tilde{E}_a \frac{d^2 P_{e \rightarrow e\gamma a}(E_\gamma, \tilde{E}_a, \tilde{E}_e, t)}{d\tilde{E}_a d\tilde{E}_e} \right], \end{aligned} \quad (\text{A2a})$$

$$\begin{aligned} \frac{\partial I_e(E_e, t)}{\partial t} = & \int_{E_e + m_e}^{E_0} d\tilde{E}_\gamma I_\gamma(\tilde{E}_\gamma, t) \frac{dP_{\gamma \rightarrow ee}(E_e, \tilde{E}_\gamma, t)}{d\tilde{E}_\gamma} + \int_{E_e}^{E_0} d\tilde{E}_e I_e(\tilde{E}_e, t) \frac{dP_{e \rightarrow e\gamma}(\tilde{E}_e - E_e, \tilde{E}_e, t)}{d\tilde{E}_e} \\ & - \int_0^{E_e - m_e} d\tilde{E}_\gamma I_e(E_e, t) \frac{dP_{e \rightarrow e\gamma}(\tilde{E}_\gamma, E_e, t)}{d\tilde{E}_\gamma} + \int_{E_e + m_a}^{E_0} d\tilde{E}_e I_e(\tilde{E}_e, t) \\ & \times \left[\left(\frac{dP_{e \rightarrow e\gamma a}(\tilde{E}_\gamma, \tilde{E}_e - E_e - \tilde{E}_\gamma, \tilde{E}_e, t)}{d\tilde{E}_\gamma} \right)_{\tilde{E}_\gamma = \tilde{E}_e - E_e - m_a} + \int_0^{\tilde{E}_e - E_e - m_a} d\tilde{E}_\gamma \frac{d^2 P_{e \rightarrow e\gamma a}(\tilde{E}_\gamma, \tilde{E}_e - E_e - \tilde{E}_\gamma, \tilde{E}_e, t)}{d\tilde{E}_e d\tilde{E}_\gamma} \right] \\ & - \int_0^{E_e - m_e - m_a} d\tilde{E}_\gamma I_e(E_e, t) \\ & \times \left[- \left(\frac{dP_{e \rightarrow e\gamma a}(\tilde{E}_\gamma, \tilde{E}_a, E_e, t)}{d\tilde{E}_a} \right)_{\tilde{E}_a = E_e - m_e - \tilde{E}_\gamma} + \int_{m_a}^{E_e - m_e - \tilde{E}_\gamma} d\tilde{E}_a \frac{d^2 P_{e \rightarrow e\gamma a}(\tilde{E}_\gamma, \tilde{E}_a, E_e, t)}{d\tilde{E}_\gamma d\tilde{E}_a} \right] \\ & + \int_{E_e + m_{\gamma'} + m_a}^{E_0} d\tilde{E}_e I_e(\tilde{E}_e, t) \\ & \times \left[\left(\frac{dP_{e \rightarrow e\gamma' a}(\tilde{E}_{\gamma'}, \tilde{E}_e - E_e - \tilde{E}_{\gamma'}, \tilde{E}_e, t)}{d\tilde{E}_{\gamma'}} \right)_{\tilde{E}_{\gamma'} = \tilde{E}_e - E_e - m_a} + \int_{m_{\gamma'}}^{\tilde{E}_e - E_e - m_a} d\tilde{E}_{\gamma'} \frac{d^2 P_{e \rightarrow e\gamma' a}(\tilde{E}_{\gamma'}, \tilde{E}_e - E_e - \tilde{E}_{\gamma'}, \tilde{E}_e, t)}{d\tilde{E}_e d\tilde{E}_{\gamma'}} \right] \end{aligned}$$

$$\begin{aligned}
& - \int_{m_\gamma}^{E_e - m_e - m_a} d\tilde{E}_{\gamma'} I_e(E_e, t) \\
& \times \left[- \left(\frac{dP_{e \rightarrow e\gamma'a}(\tilde{E}_{\gamma'}, \tilde{E}_a, E_e, t)}{d\tilde{E}_a} \right)_{\tilde{E}_a = E_e - m_e - \tilde{E}_{\gamma'}} + \int_{m_a}^{E_e - m_e - \tilde{E}_{\gamma'}} d\tilde{E}_a \frac{d^2 P_{e \rightarrow e\gamma'a}(\tilde{E}_{\gamma'}, \tilde{E}_a, E_e, t)}{d\tilde{E}_{\gamma'} d\tilde{E}_a} \right] \\
& + \int_{E_e + m_{\gamma'}}^{E_0} d\tilde{E}_e I_e(\tilde{E}_e, t) \frac{dP_{e \rightarrow e\gamma'}(\tilde{E}_e - E_e, \tilde{E}_e, t)}{d\tilde{E}_e} - \int_{m_{\gamma'}}^{E_e - m_e} d\tilde{E}_{\gamma'} I_e(E_e, t) \frac{dP_{e \rightarrow e\gamma'}(\tilde{E}_{\gamma'}, E_e, t)}{d\tilde{E}_{\gamma'}}, \quad (\text{A2b})
\end{aligned}$$

$$\begin{aligned}
\frac{\partial I_{\gamma'}(E_{\gamma'}, t)}{\partial t} &= \int_{m_e + E_{\gamma'}}^{E_0} d\tilde{E}_e I_e(\tilde{E}_e, t) \frac{dP_{e \rightarrow e\gamma'}(E_{\gamma'}, \tilde{E}_e, t)}{d\tilde{E}_e} + \int_{m_e + E_{\gamma'} + m_a}^{E_0} d\tilde{E}_e I_e(\tilde{E}_e, t) \\
& \times \left[\left(\frac{dP_{e \rightarrow e\gamma'a}(E_{\gamma'}, \tilde{E}_a, \tilde{E}_e, t)}{d\tilde{E}_a} \right)_{\tilde{E}_a = \tilde{E}_e - m_e - E_{\gamma'}} + \int_{m_a}^{\tilde{E}_e - m_e - E_{\gamma'}} d\tilde{E}_a \frac{d^2 P_{e \rightarrow e\gamma'a}(E_{\gamma'}, \tilde{E}_a, \tilde{E}_e, t)}{d\tilde{E}_a d\tilde{E}_e} \right], \quad (\text{A2c})
\end{aligned}$$

$$\begin{aligned}
\frac{\partial I_a(E_a, t)}{\partial t} &= \int_{m_e + E_a}^{E_0} d\tilde{E}_e I_e(\tilde{E}_e, t) \times \left[\left(\frac{dP_{e \rightarrow e\gamma'a}(\tilde{E}_\gamma, E_a, \tilde{E}_e, t)}{d\tilde{E}_\gamma} \right)_{\tilde{E}_\gamma = \tilde{E}_e - m_e - E_a} + \int_0^{\tilde{E}_e - m_e - E_a} \frac{d^2 P_{e \rightarrow e\gamma'a}(\tilde{E}_\gamma, E_a, \tilde{E}_e, t)}{d\tilde{E}_\gamma d\tilde{E}_e} \right] \\
& + \int_{m_e + m_{\gamma'} + E_a}^{E_0} d\tilde{E}_e I_e(\tilde{E}_e, t) \left[\left(\frac{dP_{e \rightarrow e\gamma'a}(\tilde{E}_{\gamma'}, E_a, \tilde{E}_e, t)}{d\tilde{E}_{\gamma'}} \right)_{\tilde{E}_{\gamma'} = \tilde{E}_e - m_e - E_a} + \int_{m_{\gamma'}}^{\tilde{E}_e - m_e - E_a} \frac{d^2 P_{e \rightarrow e\gamma'a}(\tilde{E}_{\gamma'}, E_a, \tilde{E}_e, t)}{d\tilde{E}_{\gamma'} d\tilde{E}_e} \right]. \quad (\text{A2d})
\end{aligned}$$

In the above, $P_{\gamma \rightarrow ee}(E_f, E_i, t)$ is the NPP probability for a photon with energy E_i to produce an electron-positron pair with electron energy E_f , $P_{e \rightarrow e\gamma'}(E_f, E_i, t)$ is the NCS (or NCS-type) probability of an electron with energy E_i to scatter and generate a photon (DP) with energy E_f , and $P_{e \rightarrow e\gamma'(\gamma)a}(E_{f_{\gamma'(\gamma)}}, E_{f_a}, E_i, t)$ is the probability of an electron with energy E_i to undergo an off-shell emission process like in Figs. 3(b) and 3(c), producing a photon (DP) and ALP with energies $E_{f_{\gamma'(\gamma)}}$ and E_{f_a} , respectively. The time dependence of these probabilities is an artifact of the time dependence of the background field felt by the particles.

The cascade equations can be understood as follows: the change in the electron intensity $I_e(E_e, t)$ at time t is determined by the influx of electrons at energy E_e due to more energetic electrons undergoing NCS/NCS-type scattering and off-shell emission processes, the efflux of electrons at energy E_e undergoing similar processes, and the production of electrons at energy E_e due to NPP. Likewise, the change in photon intensity $I_\gamma(E_\gamma, t)$ at a time t is determined by the loss of photons at energy E_γ due to NPP, and the production of photons at energy E_γ due to electrons and positrons at energy $> E_\gamma$ undergoing NCS or off-shell emission. The change in DP and ALP intensities are determined by their production from electrons at higher energies. Note that here we have ignored the feedback of NCS and reverse off-shell processes from DPs, ALPs, and positrons due to their negligible impact on the dynamics.

Implicit in the above equations is the locally monochromatic approximation (LMA), which says that the

background field seen by the particles is slowly varying compared to the timescale of emission [91]. Also neglected above is the small amount of energy provided by the background field to stimulate NCS-type and NPP processes. For example, momentum conservation for the n th emission harmonic in NCS requires $q_0 + n\omega = q'_0 + \omega'$, where $q = p + e^2|A|^2k/(2kp)$ (similarly q') is the initial (final) effective momentum of an electron that is propagating through the background field with momentum p (p'), and $\omega = k_0$ ($\omega' = k'_0$). Since we are interested in electron beams with energy ~ 10 – 100 GeV in a 10^2 – 10^3 nm (1–10 eV) laser, and the emission rates are exponentially suppressed in n , we can approximate the photon emission energy as $\omega' \approx q_0 - q'_0$. The limits of integration in Eq. (A2) are therefore taken to be the initial beam energy, rather than infinity.

The NPP, NCS, and off-shell emission rates are computed via nonperturbative QED calculations similar to those conducted in Appendix B of Ref. [77]. For NPP $\gamma(k') \rightarrow e_V^-(p)e_V^+(p')$, we obtain

$$\begin{aligned}
\frac{dP_{\gamma \rightarrow e^+e^-}}{dq_0} &= \frac{2\alpha m_e^2}{\omega' k_-} \sum_{n > n_{0,\text{pp}}} \left[\left(1 + \frac{r^2}{2} \right) J_n^2 \right. \\
& \left. - \frac{\eta^2}{2} \left(J_n^2 - \frac{J_{n-1}^2}{2} - \frac{J_{n+1}^2}{2} \right) \left(\frac{kp}{kp'} + \frac{kp'}{kp} \right) \right], \quad (\text{A3})
\end{aligned}$$

where α is the fine-structure constant, $n_{0,\text{pp}} = (m_\gamma^2 + 2m_\gamma m_e(1 + \eta^2))/(2kp)$, $r = m_\gamma/m_e$, and $J_n(z)$

are Bessel functions with argument

$$z = e|A| \sqrt{\frac{2n(kk')}{(kq)(kq')} - \frac{m_e^2(1+\eta^2)(kk')^2}{(kq)^2(kq')^2} - \frac{m_\gamma^2}{(kp)(kp')}}. \quad (\text{A4})$$

For NCS $e_V^-(p) \rightarrow e_V^-(p')\gamma(k')$, we similarly have

$$\begin{aligned} \frac{dP_{e \rightarrow e\gamma}}{d\omega'} &= \frac{2\alpha m_e^2}{q_0 q_-} \sum_{n \geq n_{0,\text{CS}}} \left[\left(1 + \frac{r^2}{2}\right) J_n^2 \right. \\ &\quad \left. + \frac{\eta^2}{2} \left(J_n^2 - \frac{J_{n-1}^2}{2} - \frac{J_{n+1}^2}{2} \right) \left(\frac{kp}{kp'} + \frac{kp'}{kp} \right) \right], \end{aligned} \quad (\text{A5})$$

where $n_{0,\text{CS}} = (m_X^2 + 2m_X m_e(1 + \eta^2))/(2kp)$ is the kinematic limit for producing a particle of mass m_X via NCS. The rate for the NCS-type process $e_V^- \rightarrow e_V^-\gamma'$ is straightforwardly obtained by substituting $\alpha \rightarrow e^2\alpha$ and $m_\gamma \rightarrow m_{\gamma'}$ in Eq. (A5) and in the argument of the Bessel functions. For the off-shell emission of ALPs $e_V^- \rightarrow e_V^-\gamma a$, we refer to Eq. (B23) in Ref. [77]. Lastly, we consider the simultaneous off-shell emission of DPs and ALPs $e_V^- \rightarrow e_V^-\gamma' a$; in the limit of light DPs, this rate is obtained from the ALP-only off-shell emission via the substitutions $g_{a\gamma\gamma} \rightarrow g_{a\gamma\gamma'}$ and $m_\gamma \rightarrow m_{\gamma'}$, while for non-negligible $m_{\gamma'}$ the final-state phase space is greatly reduced.

Plugging in the various rates $dP_{i \rightarrow f}/dE$, one can numerically solve Eq. (A2) with the initial conditions

$$\begin{aligned} I_e(E_e, 0) &= \delta(E_e - E_0), \\ I_\gamma(E_\gamma, 0) &= I_{\gamma'}(E_{\gamma'}, 0) = I_a(E_a, 0) = 0. \end{aligned} \quad (\text{A6})$$

The differential Y flux ($Y = \gamma, \gamma', a$) per incident electron is obtained by taking

$$\frac{dN_Y}{dE_Y} = I_Y(E_Y, \infty). \quad (\text{A7})$$

To accurately model the spectra produced at LUXE, we account for the spatial variance of the electron beam by discretizing its distribution into suitably fine bins and computing the spectra resulting from each bin independently. The cascade equations differ across bins beyond simple rescaling due to the differing field profile felt by each point. The beam and optical dump profiles are assumed to collide with intersecting centers of mass at 17.2° . The total spectrum dN_Y/dE_Y is obtained by summing over the spectra produced by each bin, weighted by the beam intensity profile. We neglect dispersion effects and apply both the LMA and collinearity approximation.

Figure 8 shows plots of the NCS, NCS-type, and off-shell emission rates for various parameters ($\eta, m_{\gamma'}, m_a$).

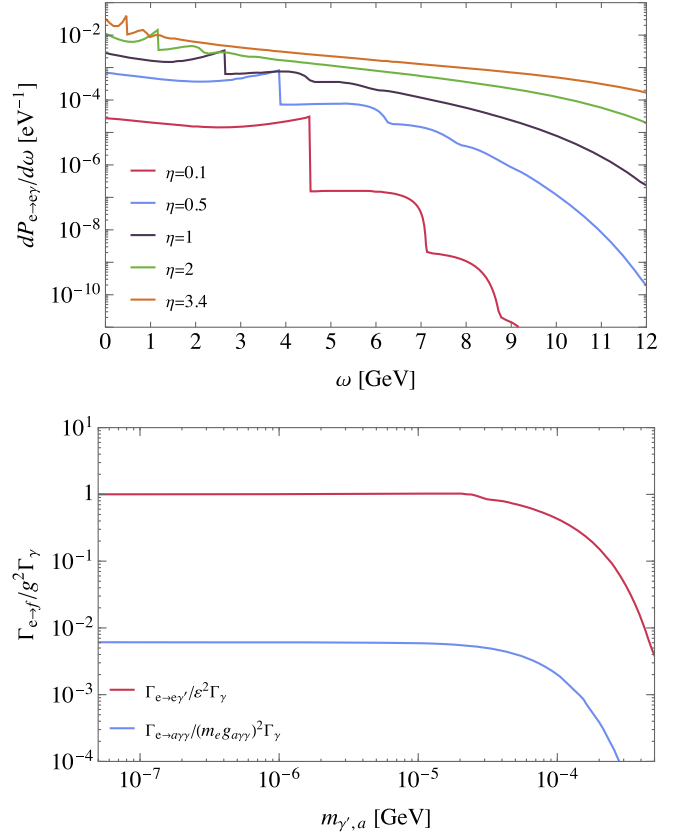


FIG. 8. Rates for processes in strong-field QED. Top: the differential rate per unit energy for nonlinear Compton scattering from an incoming electron of energy 16.5 GeV for various field strengths/quantum nonlinearity parameters. Bottom: new physics production rates for a massive dark photon (red) and an axionlike particle (blue) for varying NP particle mass, normalized by their respective couplings and rate of nonlinear Compton scattering $\Gamma_\gamma^{-1} \equiv \Gamma_{e \rightarrow e\gamma}^{-1} \approx 12$ fs.

APPENDIX B: SECONDARY NP DIFFERENTIAL CROSS SECTIONS

For secondary NP production processes, we perform an exact tree-level computation of the differential cross sections for all the possibilities appearing in Eq. (4). The relevant Feynman diagrams contributing to each process are presented in Fig. 2.

As seen in Fig. 2, all secondary production avenues arise from elastic 2-to-2 scattering processes, where the incoming particles are a primary-produced particle (γ, γ' , or a) and a dump particle (dump electron e^- or dump atom N), and the outgoing particles are a secondary-produced particle (γ' or a) and a dump particle (same as incoming dump particle).

We denote the incoming momenta by k , outgoing momenta by p , and indicate primary (secondary) particles by Y (X) subscripts and dump particles by a D subscript. Thus, for any of the secondary NP processes described above, the laboratory-frame 4-momenta are given by

$$\begin{aligned}
k_Y &= \left(E_Y, 0, 0, \sqrt{E_Y^2 - m_Y^2} \right), & k_D &= (m_D, 0, 0, 0), \\
p_X &= \left(E_X, \sin \theta \sqrt{E_X^2 - m_X^2}, 0, \cos \theta \sqrt{E_X^2 - m_X^2} \right), & p_D &= k_Y + k_D - p_X,
\end{aligned} \tag{B1}$$

where θ is the scattering angle. The transferred momentum is

$$q = k_Y - p_X. \tag{B2}$$

We use the standard Mandelstam variables

$$\begin{aligned}
s &= (k_Y + k_D)^2 = m_Y^2 + m_D^2 + 2m_DE_Y, \\
t &= q^2 = m_Y^2 + m_X^2 - 2E_Y E_X \sqrt{E_Y^2 - m_Y^2} \sqrt{E_X^2 - m_X^2} \cos \theta, \\
u &= (k_D - p_X)^2 = m_D^2 + m_X^2 - 2m_DE_X,
\end{aligned} \tag{B3}$$

together with the identity $s + t + u = m_Y^2 + m_X^2 + 2m_D^2$. Note that we are only considering elastic processes, so the incoming and outgoing dump particles are identical $k_D^2 = p_D^2 = m_D^2$.

With these definitions, the differential cross section $d\sigma_{Y \rightarrow X}/dt$ for any of the 2-to-2 processes we discussed is given by

$$\frac{d\sigma_{Y \rightarrow X}}{dt} = \frac{|\bar{\mathcal{M}}|^2}{16\pi(s^2 + m_Y^4 + m_D^4 - 2m_Y^2 s - 2m_D^2 s - 2m_Y^2 m_D^2)}, \tag{B4}$$

and the integration boundaries (assuming full 2π coverage of the forward direction in the center-of-mass frame) are

$$\begin{aligned}
t_0 &= t_1 + \frac{m_D \sqrt{E_Y^2 - m_Y^2}}{m_Y^2 + m_D^2 + 2m_DE_Y} \sqrt{4m_D^2 E_Y^2 + 4(m_Y^2 - m_X^2)m_DE_Y + (m_Y^2 - m_X^2)^2 - 4m_D^2 m_X^2}, \\
t_1 &= \frac{-2m_D^2 E_Y^2 - (m_Y^2 - m_X^2)m_DE_Y + m_D^2(m_Y^2 + m_X^2)}{m_Y^2 + m_D^2 + 2m_DE_Y}.
\end{aligned} \tag{B5}$$

In the above, $|\bar{\mathcal{M}}|^2$ is the spin/polarization-averaged squared amplitude, which is computed by evaluating the diagrams in Fig. 2.

For processes involving N , we assume the dump atoms to be scalars with electric charge Ze (here, $Z = 74$ is the atomic number of tungsten) and model them using the normalized form factor

$$F_0(t) = \frac{1}{Z} \sqrt{G_2^{\text{el}}(t) + G_2^{\text{inel}}(t)}, \tag{B6}$$

where G_2^{el} , G_2^{inel} are the elastic and inelastic atomic form factors (see Refs. [26,113]) defined by

$$G_2^{\text{el}}(t) = \left(\frac{a^2 t}{1 + a^2 t} \right)^2 \left(\frac{1}{1 + t/d} \right)^2 Z^2, \tag{B7}$$

$$G_2^{\text{inel}}(t) = \left(\frac{a'^2 t}{1 + a'^2 t} \right)^2 \left(\frac{1 + \frac{t}{4m_p^2} (\mu_p^2 - 1)}{(1 + \frac{t}{0.71 \text{ GeV}^2})^4} \right)^2 Z, \tag{B8}$$

with the parameters

$$\begin{aligned}
a &= 111Z^{-1/3}/m_e, \\
d &= 0.164 \text{ GeV}^2 A_W^{-2/3}, \\
a' &= 773Z^{-2/3}/m_e, \\
\mu_p &= 2.79
\end{aligned} \tag{B9}$$

(m_e , m_p are the electron and proton masses, respectively).

The spin/polarization-averaged squared amplitudes are computed in Wolfram's *Mathematica* software with the help of the FeynCalc package. We obtain the following differential cross sections:

$$\begin{aligned}
\frac{d\sigma_{\gamma \rightarrow \gamma'}}{dt} &= \frac{2\pi\alpha^2}{(s+t-m_e^2-m_{\gamma'}^2)^2} \times (s-m_e^2)^{-4} \times \varepsilon^2 \\
&\times [4m_e^6(m_{\gamma'}^2-2s) + 2m_e^8 \\
&+ m_e^4(7m_{\gamma'}^4 - 4(3s+2t)m_{\gamma'}^2 + 12s^2 + 4st + 3t^2) \\
&+ s(s+t-m_{\gamma'}^2)(m_{\gamma'}^4 - 2m_{\gamma'}^2s + 2s^2 + 2st + t^2) \\
&+ m_e^2((12s^2 + 4st + 3t^2)m_{\gamma'}^2 - 8s^3 \\
&- 8s^2t - 2st^2 - t^3 - (6s+5t)m_{\gamma'}^4 + 3m_{\gamma'}^6)],
\end{aligned} \tag{B10a}$$

$$\begin{aligned}
\frac{d\sigma_{\gamma \rightarrow a}}{dt} &= \frac{8\alpha Z^2 F_0^2(t)}{(4m_N^2 - t)^2} \times (s-m_N^2)^{-2} \left(\frac{g_{a\gamma\gamma}^2}{4t^2} + \frac{\varepsilon g_{a\gamma\gamma} g_{a\gamma'\gamma'}}{t(t-m_{\gamma'}^2)} + \frac{\varepsilon^2 g_{a\gamma'\gamma'}^2}{(t-m_{\gamma'}^2)^2} \right) \\
&\times m_N^4(m_a^2 t(m_N^2 + s) - m_a^4 m_N^2 - t((m_N^2 - s)^2 + st)),
\end{aligned} \tag{B10b}$$

$$\begin{aligned}
\frac{d\sigma_{\gamma' \rightarrow a}}{dt} &= \frac{-16\alpha Z^2 F_0^2(t)}{3(4m_N^2 - t)^2} (s^2 + m_{\gamma'}^4 + m_N^4 - 2m_{\gamma'}^2 s - 2m_N^2 s - 2m_N^2 m_{\gamma'}^2)^{-1} \\
&\times \left(\frac{g_{a\gamma\gamma'}^2}{t^2} + \frac{\varepsilon g_{a\gamma\gamma'} g_{a\gamma'\gamma'}}{t(t-m_{\gamma'}^2)} + \frac{\varepsilon^2 g_{a\gamma'\gamma'}^2}{4(t-m_{\gamma'}^2)^2} \right) m_N^4 [m_N^4 t + m_a^4 m_N^2 - m_a^2(m_N^2(2m_{\gamma'}^2 + t) + t(s-m_{\gamma'}^2)) \\
&+ m_N^2(m_{\gamma'}^4 - t(m_{\gamma'}^2 + 2s)) + st(s+t-m_{\gamma'}^2)],
\end{aligned} \tag{B10c}$$

$$\begin{aligned}
\frac{d\sigma_{a \rightarrow \gamma'}}{dt} &= \frac{-16\alpha Z^2 F_0^2(t)}{(4m_N^2 - t)^2} (s^2 + m_a^4 + m_N^4 - 2m_a^2 s - 2m_N^2 s - 2m_N^2 m_a^2)^{-1} \\
&\times \left(\frac{g_{a\gamma\gamma'}^2}{t^2} + \frac{\varepsilon g_{a\gamma\gamma'} g_{a\gamma'\gamma'}}{t(t-m_{\gamma'}^2)} + \frac{\varepsilon^2 g_{a\gamma'\gamma'}^2}{4(t-m_{\gamma'}^2)^2} \right) \times m_N^4 [m_N^4 t + m_a^4 m_N^2 \\
&- m_a^2(m_N^2(2m_{\gamma'}^2 + t) + t(s-m_{\gamma'}^2)) + m_N^2(m_{\gamma'}^4 - t(m_{\gamma'}^2 + 2s)) + st(s+t-m_{\gamma'}^2)].
\end{aligned} \tag{B10d}$$

These are numerically integrated between the limits given in Eq. (B5) to achieve the result in Eq. (12).

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- [1] L. B. Okun, Limits on electrodynamics: Paraphotons?, *Sov. Phys. JETP* **56**, 502 (1982), <https://jetp.ras.ru/cgi-bin/index/e/56/3/p502?a=list>.
 - [2] B. Holdom, Two $U(1)$'s and ε charge shifts, *Phys. Lett.* **166B**, 196 (1986).
 - [3] M. Fabbrihesi, E. Gabrielli, and G. Lanfranchi, *The Physics of the Dark Photon* (Springer, Cham, 2021), 10.1007/978-3-030-62519-1.
 - [4] R. D. Peccei and H. R. Quinn, CP conservation in the presence of pseudoparticles, *Phys. Rev. Lett.* **38**, 1440 (1977).
 - [5] R. D. Peccei and H. R. Quinn, Constraints imposed by CP conservation in the presence of pseudoparticles, *Phys. Rev. D* **16**, 1791 (1977).
 - [6] F. Wilczek, Problem of strong P and T invariance in the presence of instantons, *Phys. Rev. Lett.* **40**, 279 (1978).
 - [7] S. Weinberg, A new light boson?, *Phys. Rev. Lett.* **40**, 223 (1978).
 - [8] J. E. Kim, Weak-interaction singlet and strong CP invariance, *Phys. Rev. Lett.* **43**, 103 (1979).
 - [9] M. A. Shifman, A. I. Vainshtein, and V. I. Zakharov, Can confinement ensure natural CP invariance of strong interactions?, *Nucl. Phys.* **B166**, 493 (1980).
 - [10] A. Hook, TASI lectures on the strong CP problem and axions, *Proc. Sci. TASI2018* (2019) 004.
 - [11] D. F. Bartlett and S. Lögl, Limits on an electromagnetic fifth force, *Phys. Rev. Lett.* **61**, 2285 (1988).

- [12] D. Kroff and P. C. Malta, Constraining hidden photons via atomic force microscope measurements and the Plimpton-Lawton experiment, *Phys. Rev. D* **102**, 095015 (2020).
- [13] J. Jaeckel and S. Roy, Spectroscopy as a test of Coulomb's law: A probe of the hidden sector, *Phys. Rev. D* **82**, 125020 (2010).
- [14] G. Marocco, Dark photon limits from magnetic fields and astrophysical plasmas, [arXiv:2110.02875](#).
- [15] S. Yan, L. Li, and J. Fan, Dark photon limits from magnetic fields and astrophysical plasmas, *J. High Energy Phys.* **06** (2024) 028.
- [16] F. McCarthy, D. Pírvu, J. C. Hill, J. Huang, M. C. Johnson, and K. K. Rogers, Dark photon limits from patchy dark screening of the cosmic microwave background, *Phys. Rev. Lett.* **133**, 141003 (2024).
- [17] K. Ehret *et al.* (ALPS Collaboration), New ALPS results on hidden-sector lightweights, *Phys. Lett. B* **689**, 149 (2010).
- [18] T. Inada, T. Namba, S. Asai, T. Kobayashi, Y. Tanaka, K. Tamasaku, K. Sawada, and T. Ishikawa, Results of a search for paraphotons with intense X-ray beams at SPring-8, *Phys. Lett. B* **722**, 301 (2013).
- [19] S. R. Parker, J. G. Hartnett, R. G. Povey, and M. E. Tobar, Tobar, Cryogenic resonant microwave cavity searches for hidden sector photons, *Phys. Rev. D* **88**, 112004 (2013).
- [20] A. Wagner *et al.*, Search for hidden sector photons with the ADMX detector, *Phys. Rev. Lett.* **105**, 171801 (2010).
- [21] M. Betz, F. Caspers, M. Gasior, M. Thumm, and S. W. Rieger, First results of the CERN resonant weakly interacting Sub-eV particle search (CROWS), *Phys. Rev. D* **88**, 075014 (2013).
- [22] J. Redondo, Helioscope bounds on hidden sector photons, *J. Cosmol. Astropart. Phys.* **07** (2008) 008.
- [23] M. Schwarz, E.-A. Knabbe, A. Lindner, J. Redondo, A. Ringwald, M. Schneide, J. Susol, and G. Wiedemann, Results from the solar hidden photon search (SHIPS), *J. Cosmol. Astropart. Phys.* **08** (2015) 011.
- [24] M. Danilov, S. Demidov, and D. Gorbunov, Constraints on hidden photons produced in nuclear reactors, *Phys. Rev. Lett.* **122**, 041801 (2019).
- [25] S. N. Gninenko, Constraints on Sub-GeV hidden sector gauge bosons from a search for heavy neutrino decays, *Phys. Lett. B* **713**, 244 (2012).
- [26] J. D. Bjorken, R. Essig, P. Schuster, and N. Toro, New fixed-target experiments to search for dark gauge forces, *Phys. Rev. D* **80**, 075018 (2009).
- [27] J. D. Bjorken *et al.*, Search for neutral metastable penetrating particles produced in the SLAC beam dump, *Phys. Rev. D* **38**, 3375 (1988).
- [28] E. M. Riordan *et al.*, Search for short-lived axions in an electron-beam-dump experiment, *Phys. Rev. Lett.* **59**, 755 (1987).
- [29] A. Bross, M. Crisler, S. Pordes, J. Volk, S. Errede, and J. Wrbanek, Search for short-lived particles produced in an electron beam dump, *Phys. Rev. Lett.* **67**, 2942 (1991).
- [30] M. Endo, K. Hamaguchi, and G. Mishima, Constraints on hidden photon models from electron $g - 2$ and hydrogen spectroscopy, *Phys. Rev. D* **86**, 095029 (2012).
- [31] D. Barducci, E. Bertuzzo, G. Grilli di Cortona, and G. M. Salla, Dark photon bounds in the dark EFT, *J. High Energy Phys.* **12** (2021) 081.
- [32] J. R. Batley *et al.* (NA48/2 Collaboration), Search for the dark photon in π^0 decays, *Phys. Lett. B* **746**, 178 (2015).
- [33] P. Adlarson *et al.* (WASA-at-COSY Collaboration), Search for a dark photon in the $\pi^0 \rightarrow e^+e^-\gamma$ decay, *Phys. Lett. B* **726**, 187 (2013).
- [34] S. Abrahamyan *et al.*, Search for a new gauge boson in electron-nucleus fixed-target scattering by the APEX experiment, *Phys. Rev. Lett.* **107**, 191804 (2011).
- [35] H. Merkel *et al.* (A1 Collaboration), Search for light gauge bosons of the dark sector at the Mainz Microtron, *Phys. Rev. Lett.* **106**, 251802 (2011).
- [36] J. P. Lees *et al.*, Search for a Dark Photon in e^+e^- Collisions at BABAR, *Phys. Rev. Lett.* **113**, 201801 (2014).
- [37] D. Babusci *et al.* (KLOE-2 Collaboration), Limit on the production of a light vector gauge boson in ϕ meson decays with the KLOE detector, *Phys. Lett. B* **720**, 111 (2013).
- [38] H. Abreu *et al.* (FASER Collaboration), Search for dark photons with the FASER detector at the LHC, *Phys. Lett. B* **848**, 138378 (2024).
- [39] S. Andriamonje *et al.* (CAST Collaboration), An improved limit on the axion-photon coupling from the CAST experiment, *J. Cosmol. Astropart. Phys.* **04** (2007) 010.
- [40] V. Anastassopoulos *et al.* (CAST Collaboration), New CAST limit on the axion-photon interaction, *Nat. Phys.* **13**, 584 (2017).
- [41] N. Vinyoles, A. Serenelli, F. L. Villante, S. Basu, J. Redondo, and J. Isern, New axion and hidden photon constraints from a solar data global fit, *J. Cosmol. Astropart. Phys.* **10** (2015) 015.
- [42] A. Ayala, I. Domínguez, M. Giannotti, A. Mirizzi, and O. Straniero, Revisiting the bound on axion-photon coupling from globular clusters, *Phys. Rev. Lett.* **113**, 191302 (2014).
- [43] M. J. Dolan, F. J. Hiskens, and R. R. Volkas, Advancing globular cluster constraints on the axion-photon coupling, *J. Cosmol. Astropart. Phys.* **10**, 096 (2022).
- [44] F. Della Valle, A. Ejlli, U. Gastaldi, G. Messineo, E. Milotti, R. Pengo, G. Ruoso, and G. Zavattini, The PVLAS experiment: measuring vacuum magnetic birefringence and dichroism with a birefringent Fabry-Perot cavity, *Eur. Phys. J. C* **76**, 24 (2016).
- [45] J. W. D. Halliday *et al.*, Bounds on heavy axions with an x-ray free electron laser, *Phys. Rev. Lett.* **134**, 055001 (2025).
- [46] K. Homma *et al.* (SAPPHIRES Collaboration), Search for sub-eV axion-like resonance states via stimulated quasi-parallel laser collisions with the parameterization including fully asymmetric collisional geometry, *J. High Energy Phys.* **12** (2021) 108.
- [47] M. J. Dolan, T. Ferber, C. Hearty, F. Kahlhoefer, and K. Schmidt-Hoberg, Revised constraints and Belle II sensitivity for visible and invisible axion-like particles, *J. High Energy Phys.* **12** (2017) 094.

- [48] P. Astier *et al.* (NOMAD Collaboration), Search for eV (pseudo)scalar penetrating particles in the SPS neutrino beam, *Phys. Lett. B* **479**, 371 (2000).
- [49] J. Jaeckel and M. Spannowsky, Probing MeV to 90 GeV axion-like particles with LEP and LHC, *Phys. Lett. B* **753**, 482 (2016).
- [50] F. Abudinén *et al.* (Belle II Collaboration), Search for axionlike particles produced in e^+e^- collisions at Belle II, *Phys. Rev. Lett.* **125**, 161806 (2020).
- [51] M. Ablikim *et al.* (BESIII Collaboration), Search for an axion-like particle in radiative J/ψ decays, *Phys. Lett. B* **838**, 137698 (2023).
- [52] S. Knapen, T. Lin, H. K. Lou, and T. Melia, Searching for axionlike particles with ultraperipheral heavy-ion collisions, *Phys. Rev. Lett.* **118**, 171801 (2017).
- [53] A. M. Sirunyan *et al.* (CMS Collaboration), Evidence for light-by-light scattering and searches for axion-like particles in ultraperipheral PbPb collisions at $\sqrt{s_{NN}} = 5.02$ TeV, *Phys. Lett. B* **797**, 134826 (2019).
- [54] G. Aad *et al.* (ATLAS Collaboration), Measurement of light-by-light scattering and search for axion-like particles with 2.2 nb^{-1} of Pb + Pb data with the ATLAS detector, *J. High Energy Phys.* **03** (2021) 243.
- [55] J. Blümlein *et al.*, Limits on the mass of light (pseudo) scalar particles from Bethe-Heitler e^+e^- and $\mu^+\mu^-$ pair production in a proton-iron beam dump experiment, *Int. J. Mod. Phys. A* **7**, 3835 (1992).
- [56] E. M. Riordan *et al.*, Search for short-lived axions in an electron-beam-dump experiment, *Phys. Rev. Lett.* **59**, 755 (1987).
- [57] D. Banerjee *et al.* (NA64 Collaboration), Search for axionlike and scalar particles with the NA64 experiment, *Phys. Rev. Lett.* **125**, 081801 (2020).
- [58] F. Bergsma *et al.* (CHARM Collaboration), Search for axion-like particle production in 400 GeV proton-copper interactions, *Phys. Lett.* **157B**, 458 (1985).
- [59] F. Capozzi, B. Dutta, G. Gurung, W. Jang, I. M. Shoemaker, A. Thompson, and J. Yu, New constraints on ALP couplings to electrons and photons from ArgoNeuT and the Mini-BooNE beam dump, *Phys. Rev. D* **108**, 075019 (2023).
- [60] D. Aloni, C. Fanelli, Y. Soreq, and M. Williams, Photo-production of axionlike particles, *Phys. Rev. Lett.* **123**, 071801 (2019).
- [61] K. Kaneta, H. S. Lee, and S. Yun, Portal connecting dark photons and axions, *Phys. Rev. Lett.* **118**, 101802 (2017).
- [62] P. deNiverville, H.-S. Lee, and M.-S. Seo, Implications of the dark axion portal for the muon $g - 2$, B factories, fixed target neutrino experiments, and beam dumps, *Phys. Rev. D* **98**, 115011 (2018).
- [63] P. deNiverville and H. S. Lee, Implications of the dark axion portal for SHiP and FASER and the advantages of monophoton signals, *Phys. Rev. D* **100**, 055017 (2019).
- [64] D. V. Kirpichnikov, V. E. Lyubovitskij, and A. S. Zhevlakov, Implication of hidden sub-GeV bosons for the $(g - 2)_\mu$, ^8Be - ^4He anomaly, proton charge radius, EDM of fermions, and dark axion portal, *Phys. Rev. D* **102**, 095024 (2020).
- [65] P. deNiverville, H. S. Lee, and Y. M. Lee, New searches at reactor experiments based on the dark axion portal, *Phys. Rev. D* **103**, 075006 (2021).
- [66] A. Hook, G. Marques-Tavares, and C. Ristow, Supernova constraints on an axion-photon-dark photon interaction, *J. High Energy Phys.* **06** (2021) 167.
- [67] S. F. Ge, X. D. Ma, and P. Pasquini, Probing the dark axion portal with muon anomalous magnetic moment, *Eur. Phys. J. C* **81**, 787 (2021).
- [68] J. C. Gutiérrez *et al.*, Cosmology and direct detection of the Dark Axion Portal, [arXiv:2112.11387](https://arxiv.org/abs/2112.11387).
- [69] A. S. Zhevlakov, D. V. Kirpichnikov, and V. E. Lyubovitskij, Implication of the dark axion portal for the EDM of fermions and dark matter probing with NA64e, NA64 μ , LDMX, M^3 , and BABAR, *Phys. Rev. D* **106**, 035018 (2022).
- [70] P. Carena, G. Lucente, and E. Vitagliano, Probing the blue axion with cosmic optical background anisotropies, *Phys. Rev. D* **107**, 083032 (2023).
- [71] K. Jodłowski, Probing some photon portals to new physics at intensity frontier experiments, *Phys. Rev. D* **108**, 115017 (2023).
- [72] H. Hong, U. Min, M. Son, and T. You, A cosmic window on the dark axion portal, *J. High Energy Phys.* **03** (2024) 155.
- [73] A. Hook, G. Marques-Tavares, and C. Ristow, CMB spectral distortions from an axion-dark photon-photon interaction, *J. High Energy Phys.* **05** (2024) 086.
- [74] K. Jodłowski, Dark axion portal at Z boson factories, Dark axion portal at Z boson factories, *J. High Energy Phys.* **08** (2025) 022.
- [75] P. Arias, B. D. Sáez, and J. Jaeckel, Freezing-in the pure dark axion portal, *J. Cosmol. Astropart. Phys.* **06** (2025) 060.
- [76] P. Arias, B. D. Sáez, L. Duarte, J. Jones-Pérez, W. Rodriguez, and D. Z. Herrera, Probing displaced (dark) photons from low reheating freeze-in at the LHC, *J. High Energy Phys.* **01** (2026) 135.
- [77] Z. Bai *et al.*, New physics searches with an optical dump at LUXE, *Phys. Rev. D* **106**, 115034 (2022).
- [78] M. A. Soto *et al.*, Layout optimization for the LUXEN-POD experiment, *Phys. Rev. D* **112**, 112014 (2025).
- [79] H. Abramowicz *et al.*, Conceptual design report for the LUXE experiment, *Eur. Phys. J. Special Topics* **230**, 2445 (2021).
- [80] H. Abramowicz *et al.* (LUXE Collaboration), Technical design report for the LUXE experiment, *Eur. Phys. J. Special Topics* **233**, 1709 (2024).
- [81] H. Abramowicz *et al.*, Input to the ESPPU: The LUXE experiment, [arXiv:2504.00873](https://arxiv.org/abs/2504.00873).
- [82] W. Heisenberg and H. Euler, Folgerungen aus der Diracschen theorie des positrons, *Z. Phys.* **98**, 714 (1936).
- [83] R. Karplus and M. Neuman, The scattering of light by light, *Phys. Rev.* **83**, 776 (1951).
- [84] J. Schwinger, On gauge invariance and vacuum polarization, *Phys. Rev.* **82**, 664 (1951).
- [85] H. R. Reiss, Absorption of light by light, *J. Math. Phys. (N.Y.)* **3**, 59 (1962).
- [86] A. I. Nikishov and V. I. Ritus, Quantum processes in the field of a plane electromagnetic wave and in a constant

- field. I, Sov. Phys. JETP **19**, 529 (1964), <https://jetp.ras.ru/cgi-bin/e/index/e/19/2/p529?a=list>.
- [87] L. S. Brown and T. W. B. Kibble, Interaction of intense laser beams with electrons, *Phys. Rev.* **133**, A705 (1964).
 - [88] D. M. Wolkow, Über eine Klasse von Lösungen der Diracschen Gleichung, *Z. Phys.* **94**, 250 (1935).
 - [89] R. L. Workman *et al.* (Particle Data Group), Review of particle physics, *Prog. Theor. Exp. Phys.* **2022**, 083C01 (2022).
 - [90] Y. S. Tsai and V. Whitis, Thick-target bremsstrahlung and target considerations for secondary-particle production by electrons, *Phys. Rev.* **149**, 1248 (1966).
 - [91] T. Heinzl, B. King, and A. J. MacLeod, Locally monochromatic approximation to QED in intense laser fields, *Phys. Rev. A* **102**, 063110 (2020).
 - [92] W. M. Bonivento, Studies for the electro-magnetic calorimeter SplitCal for the SHiP experiment at CERN with shower direction reconstruction capability, *J. Instrum.* **13**, C02041 (2018).
 - [93] A. Bornheim, Timing performance of the CMS electromagnetic calorimeter and prospects for the future, *Proc. Sci. TIPP2014* (**2015**) 021.
 - [94] A. Lobanov, Precision timing calorimetry with the CMS HGCAL, *J. Instrum.* **15**, C07003 (2020).
 - [95] L. Martinazzoli, Crystal fibers for the LHCb calorimeter upgrade, *IEEE Trans. Nucl. Sci.* **67**, 1003 (2020).
 - [96] Y. Liu, J. Jiang, and Y. Wang, High-granularity crystal calorimetry: conceptual designs and first studies, *J. Instrum.* **15**, C04056 (2020).
 - [97] P. Ilten, Y. Soreq, M. Williams, and W. Xue, Serendipity in dark photon searches, *J. High Energy Phys.* **06** (2018) 004.
 - [98] A. Caputo, A. J. Millar, C. A. J. O'Hare, and E. Vitagliano, Dark photon limits: A handbook, *Phys. Rev. D* **104**, 095029 (2021).
 - [99] C. O'Hare, cajohare/AxionLimits: AxionLimits, <https://cajohare.github.io/AxionLimits/> (2020).
 - [100] O. Ning and B. R. Safdi, Leading axion-photon sensitivity with NuSTAR observations of M82 and M87, *Phys. Rev. Lett.* **134**, 171003 (2025).
 - [101] F. R. Candón, D. F. G. Fiorillo, G. Lucente, E. Vitagliano, and J. K. Vogel, NuSTAR bounds on radiatively decaying particles from M82, *Phys. Rev. Lett.* **134**, 171004 (2025).
 - [102] M. J. Dolan, F. J. Hiskens, and R. R. Volkas, Constraining axion-like particles using the white dwarf initial-final mass relation, *J. Cosmol. Astropart. Phys.* **09** (2021) 010.
 - [103] A. Caputo, H.-T. Janka, G. Raffelt, and E. Vitagliano, Low-energy supernovae severely constrain radiative particle decays, *Phys. Rev. Lett.* **128**, 221103 (2022).
 - [104] G. Lucente, P. Carenza, T. Fischer, M. Giannotti, and A. Mirizzi, Heavy axion-like particles and core-collapse supernovae: Constraints and impact on the explosion mechanism, *J. Cosmol. Astropart. Phys.* **12** (2020) 008.
 - [105] A. Caputo, G. Raffelt, and E. Vitagliano, Muonic boson limits: Supernova redux, *Phys. Rev. D* **105**, 035022 (2022).
 - [106] M. Diamond, D. F. G. Fiorillo, G. Marques-Tavares, and E. Vitagliano, Axion-sourced fireballs from supernovae, *Phys. Rev. D* **107**, 103029 (2023).
 - [107] D. F. G. Fiorillo, T. Pitik, and E. Vitagliano, Energy transfer by feebly interacting particles in Supernovae: The trapping regime, *Phys. Rev. Lett.* **135**, 071005 (2025).
 - [108] P. S. B. Bhupal Dev, J.-F. Fortin, S. P. Harris, K. Sinha, and Y. Zhang, First constraints on the photon coupling of axionlike particles from multimessenger studies of the neutron star merger GW170817, *Phys. Rev. Lett.* **132**, 101003 (2024).
 - [109] M. Diamond, D. F. G. Fiorillo, G. Marques-Tavares, I. Tamborra, and E. Vitagliano, Multimessenger constraints on radiatively decaying axions from GW170817, *Phys. Rev. Lett.* **132**, 101004 (2024).
 - [110] W. DeRocco, P. W. Graham, and S. Rajendran, Exploring the robustness of stellar cooling constraints on light particles, *Phys. Rev. D* **102**, 075015 (2020).
 - [111] V. Popov, On the experimental search for photon mixing, *Turk. J. Phys.* **23**, 943 (1999), <https://journals.tubitak.gov.tr/physics/vol23/iss5/17/>.
 - [112] L. D. Landau and G. Rumer, The cascade theory of electronic showers, *Proc. R. Soc.* **166**, 213 (1938).
 - [113] Y. S. Tsai, Pair production and bremsstrahlung of charged leptons, *Rev. Mod. Phys.* **46**, 815 (1974); **49**, 421(E) (1977).