

LHC as an Axion-Photon Collider

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Assuming the existence of an axionlike particle (ALP), beams of relativistic particles emit fluxes of quasireal ALPs, analogous to the photon fluxes described by Weizsäcker-Williams-type approximations. Consequently, ALP-ALP and ALP-photon collisions can occur at the LHC. We initiate the study of the LHC as an ALP collider, and show that ALP collisions provide competitive probes of certain ALP couplings. We show that ALP fluxes from heavy ions are suppressed relative to those from protons, unlike their photon counterpart. As a result, the most likely processes are ALP-photon collisions occurring in the proton-ion (pA) ultraperipheral collisions. Using our implementation of ALP fluxes in simulation tools, we show that ALP-photon collisions in pPb efficiently probe ALP couplings to third generation fermions. LHC data with realistic pPb luminosity can constrain the product of ALP couplings to nucleons and top quarks at the level of $O(0.01 \text{ TeV}^{-1})$. Notably, ALP-photon collisions naturally provide the leading probe of ALP flavor-violating couplings to the top quark. We suggest that the 2016 pPb dataset collected at CMS, ATLAS, and LHCb should be explored for evidence of such collisions.

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Introduction—Scalar fields with approximate shift symmetry are ubiquitous in the landscape of theories beyond the standard model (SM). They naturally arise from the compactification of higher-dimensional gauge fields [1–3] and of higher-form fields from string theory [4–6]. They also arise as pseudo-Goldstone bosons resulting from the spontaneous breaking of approximate global symmetries [7], as happens in the Peccei-Quinn mechanism [8]. In this context, the corresponding particle is called the axion, and it is common to refer to any of these scalar excitations as an *axion-like* particle (ALP). Axions and ALPs play an important role in cosmology, such as inflation and the cosmological constant

problem, and are excellent candidates for dark matter [9]. In a sense, ALPs provide a versatile portal to new physics, linking open problems in particle physics, cosmology, and fundamental theory.

Axionlike particles are traditionally searched for in laboratory experiments such as light-shining-through-the-wall setups [10,11], in emission from the Sun (helioscopes) [12] and as dark matter (haloscopes), see Refs. [13–15] for reviews. Other probes arise from astrophysics [16], fifth force searches [17–21], meson properties [22,23], production at colliders [24–27]—where limits have been recently reported by the LHC experiments [28–32], and more.

In this Letter, we investigate a phenomenon that has so far remained unexplored: high-energy ALP collisions. Where could ALP collisions occur in the real world? Consider photons at the Large Hadron Collider (LHC). Although the LHC collides protons (p) and/or ions (A), it can also be viewed as a photon collider. This is because, at the semiclassical level, the relativistic beams of hadrons constitute intense electromagnetic sources [33–35]. This strong flux of photons appears at the quantum level via the factorization of amplitudes with forward kinematics [36]. This remarkable fact has led to a forward physics program at the LHC, enabling elegant probes of QED

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and of physics beyond the SM in ultraperipheral collisions (UPCs) [37,38]. Notably, the phenomenon of light-by-light scattering [39] was observed for the first time at ATLAS in PbPb-induced photon collisions [29].

Following the same logic, if an ALP exists and couples to hadrons, relativistic beams emit a coherent ALP field. This emission of a flux of ALPs, sometimes referred to as “axion bremsstrahlung,” was computed for protons in [40,41], and has been used in a number of ways in searches at colliders [42–44] and beam dumps [45,46]. However, the phenomena of ALP-ALP and ALP-photon collisions, that follow in analogy to photon-photon collisions, have never been exploited. In this Letter, we initiate the investigation of ALP collisions at the LHC. We study where to best find them, and what ALP-SM couplings are best probed by such processes.

Axion couplings—The interactions of the ALP quantum field $a(x^\mu)$ relevant to this Letter are described at the electroweak scale by the following effective Lagrangian [23]:

$$\mathcal{L} = \frac{(\partial_\mu a)^2}{2} - \frac{m_a^2 a^2}{2} + c_{GG} \frac{\alpha_s a}{4\pi f} G_{\mu\nu}^a \tilde{G}^{a\mu\nu} + \frac{\partial^\mu a}{f} (\bar{u}_L \mathbf{k}_U \gamma_\mu u_L + \bar{u}_R \mathbf{k}_u \gamma_\mu u_R + \bar{d}_L \mathbf{k}_D \gamma_\mu d_L + \dots), \quad (1)$$

where $G_{\mu\nu}^a$ is the gluon field, u_L is a 3-vector encoding the three generations of left-handed quark fields, etc. The $\mathbf{k}_i$ are Hermitian 3×3 matrices in generation space. The fermion fields are in their mass basis. In particular, the $af_i \tilde{f}_i$ flavor-conserving couplings take the form $c_{u_i u_i} = [k_u - k_U]_{ii}$.

Below the QCD scale, the ALP coupling to nucleons $N = (p, n)$ is determined to be

$$\mathcal{L}_{\text{nucleons}} = -ig_N a \bar{N} \gamma_5 N, \quad g_{p,n} = \frac{m_N}{2f} (g_0(c_{uu} + c_{dd} + 2c_{GG}) \pm g_A \Delta c_{ud}), \quad \Delta c_{ud} = c_{uu} - c_{dd} + 2c_{GG} \frac{m_d - m_u}{m_d + m_u}, \quad (2)$$

with $g_0 \approx 0.44$, $g_A \approx 1.25$, using chiral perturbation theory techniques, and assuming $m_a \ll m_\pi$. When m_a is closer to the pion mass, ALP-pion mixing can substantially increase g_{ap} [23].

Axion fluxes from protons and ions—To describe ALP collisions, we need to know how ALPs are emitted from protons and ions. We focus on the elastic component of the fluxes—that will be mostly selected by our requirements on final states. These ALP fluxes are described by collinear approximations in the spirit of the Weizsäcker-Williams-type approximations for photons.

From the QFT viewpoint, we derive the ALP flux from the factorization of the ALP emission cross section in the forward region. Using the ALP-nucleon EFT of Eq. (2), the ALP and photon fluxes from the proton are given by

$$f_{a|p}(x) = \frac{g_p^2}{16\pi^2} x \log\left(\frac{E_p^2}{m_p^2 + \frac{1-x}{x^2} m_a^2}\right), \quad (3)$$

$$f_{\gamma|p}(x) = \frac{\alpha}{2\pi} \frac{1 + (1-x)^2}{x} \log\left(\frac{E_p^2}{m_p^2}\right), \quad (4)$$

where E_p and m_p are the energy and proton mass, respectively. The result Eq. (3) agrees with the flux obtained by integrating the splitting function for axion bremsstrahlung [47]. The difference in x dependence between $f_{a|p}$ and $f_{\gamma|p}$ is tied to the fact that ALP emission requires a helicity flip of the proton while photon emission does not.

For fluxes of particles from heavy ions, a more accurate description is obtained using the semiclassical approach of virtual quanta [33]. In analogy with the electromagnetic case, we treat the ALP field classically. Its equation of motion sourced by a relativistic pointlike charge is

$$\square a(t, \vec{x}) = [Zg_p + (A - Z)g_n] \delta^{(3)}[\vec{x} - \vec{x}'(t)]. \quad (5)$$

In the following, we take $g_p \approx g_n$ for simplicity. We evaluate the field at a fixed transverse distance $\vec{b}$ from the ion trajectory. The number distribution of equivalent ALPs per unit area is then obtained from the energy density of the field in frequency space:

$$\frac{d^2 f_{a/A}}{db^2}(x, \vec{b}) = \frac{A^2 g_p^2}{16\pi^3 x} (xM_A)^2 \left(K_0^2(\xi) + \frac{1}{\gamma^2} K_1^2(\xi) \right), \quad (6)$$

where $\xi = xM_A |\vec{b}|$, M_A , and γ are the ion mass and Lorentz factor, and K_n are the modified Bessel functions of the second kind.

The total distribution is obtained by integrating over the possible impact parameters, from the nuclear radius R to infinity. The resulting ALP and photon fluxes from ions are given by

$$f_{a|A}(x) = \frac{A^2 g_p^2}{8\pi^2 \gamma^2 x} y \left[K_0(y) K_1(y) - \frac{y}{2} (1 - \gamma^2) (K_1^2(y) - K_0^2(y)) \right], \quad (7)$$

$$f_{\gamma|A}(x) = \frac{2Z^2 \alpha}{\pi x} y \left[K_0(y) K_1(y) - \frac{y}{2} (K_1^2(y) - K_1^2(y)) \right], \quad (8)$$

where $y = xM_A R$. The fluxes are exponentially suppressed for high frequencies, $y \gg 1$. In the low-frequency regime $y \ll 1$, the ALP flux scales as $f_{a|A} \sim \gamma^{-2} x^{-1} \log(1/y)$, while $f_{\gamma|A} \sim x^{-1} \log(1/y)$. Since $\gamma \gg 1$, the ALP flux from ions is significantly suppressed relative to the photon flux in the ultrarelativistic limit. From the semiclassical view,

this suppression occurs because the scalar field has no transverse degrees of freedom. In contrast, in the electromagnetic field case, the dominant contribution comes from the transverse component of the electric field, which is amplified by the relativistic boost as $E'_\perp \sim \gamma E_\perp$ leading to the relative γ^2 factor between $f_{\gamma|A}$ and $f_{a|A}$.

In the ideal ultraperipheral collision of two nuclei N, N' (with $N = p, A$) emitting respectively quasireal particles a, b , the total cross section is governed by the joint number distribution of the two quasireal particles carrying fractions of the total energy of each nucleus x_a, x_b : $f_{a|N}(x_a)f_{b|N'}(x_b)dx_adx_b$. The total UPC cross section $\sigma_{N,N'}$ for a subprocess $a + b \rightarrow X$ with cross section $\sigma_{a+b \rightarrow X}$ is

$$\sigma_{N,N'} = \int d\hat{s} \frac{dL_{a,b}^{N,N'}}{d\hat{s}} \sigma_{a+b \rightarrow X}(\hat{s}), \quad (9)$$

where the effective differential luminosity for producing a system with invariant mass $\sqrt{\hat{s}}$ is given by

$$\frac{dL_{a,b}^{N,N'}}{d\hat{s}} = \frac{1}{s} \int_{\hat{s}/s}^1 dx_a dx_b f_{a|N}(x_a) f_{b|N'}(x_b) \delta\left(\frac{\hat{s}}{s} - x_a x_b\right). \quad (10)$$

The normalized profiles of $dL^{pp}/d\hat{s}$ and $dL^{pPb}/d\hat{s}$ are shown in Fig. 1. The curves show that aa and $a\gamma$ collisions occur typically at center-of-mass energy higher than the one for $\gamma\gamma$. This is due to the specific structure of the ALP fluxes, stemming from the ALP being scalar.

The validity of the ALP-nucleon EFT imposes that the ALP virtuality Q is bounded from above by the proton mass, i.e., by the inverse proton radius $R_p^{-1} \sim m_p$. Unlike for the photon, the ALP-nucleon form factors that smoothly implement the cutoff on Q^2 are unknown experimentally. Still, since these form factors come from the composite structure of the proton, which is governed by QCD, it is reasonable to assume that the scale of the response of the proton to an ALP probe is similar to a photon probe. We thus use the bound $Q^2 \leq Q_{\max}^2 \approx m_p^2$. This implies an upper bound on x , that is for instance $x \leq 0.62$ when $m_a \ll m_p$. Notice there is some intrinsic uncertainty associated to this cutoff, for instance a variation of $Q_{\max}$ by $\pm 10\%$ results in a variation of the total cross section by $\sim 13\%$.

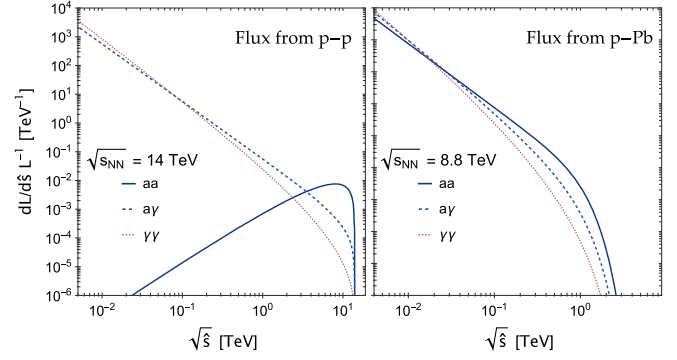
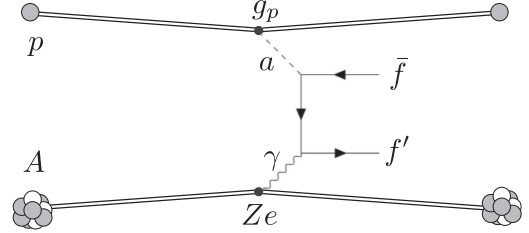


FIG. 1. Normalized differential luminosities of the $aa, a\gamma, \gamma\gamma$ collisions triggered from the pp and pPb UPCs occurring at the LHC. $\sqrt{s_{NN}}$ is the center-of-mass energy per nucleon.

We have interfaced the ALP fluxes with MadGraph5 [48].

Axion collisions at the LHC—Where are ALP collisions most likely to happen at the LHC?

Let us first select the collision mode. A key feature of our fluxes is that ALP emission from ions is suppressed by an extra $(1/\gamma^2)$ factor. Therefore, while the yields for $\gamma\gamma$ from the modes pp, pA, AA have relative scaling $1:Z^2:Z^4$, the yields from $a\gamma$ and aa are, respectively, $1:Z^2 + (A^2/\gamma^2):(Z^2 A^2/\gamma^2)$ and $1:(A^2/\gamma^2):(A^4/\gamma^4)$. Since $\gamma = O(10^3)$ at the LHC, we have $Z^2 \gg (A^2/\gamma^2)$, which is enough to conclude that the pA mode is the most interesting, with a emitted from the proton and γ from the ion. Focusing on the production of fermion-antifermion $\bar{f}f'$ pairs, the full ultraperipheral process is



Upon integration by parts in Eq. (1), the non-derivative ALP couplings scale with the fermion masses, hence the third generation is the most favorable. We thus focus on the top quark sector.

The key cross sections to compare the collision modes are as follows:

UPC	pp	pO	pPb
$a\gamma \rightarrow t\bar{t}$	$3.2 \times 10^{-3} (g^2 c_u^2 / f^2)$ pb	$10.3 (g^2 c_u^2 / f^2)$ pb	$1.4 \times 10^4 (g^2 c_u^2 / f^2)$ pb
$a\gamma \rightarrow t\bar{c}$	$3.4 \times 10^{-3} (g^2 c_{tc} ^2 / f^2)$ pb	$9.9 (g^2 c_{tc} ^2 / f^2)$ pb	$2.3 \times 10^4 (g^2 c_{tc} ^2 / f^2)$ pb
$\gamma\gamma \rightarrow t\bar{t}$	1.8×10^{-4} pb	1.6×10^{-3} pb	3.0×10^{-2} pb

with the ALP coupling scale f in units of TeV, and $c_{tc} = [k_u - k_U]_{23}$. We assume the nominal c.m. energies $\sqrt{s_{NN}} = 14, 9.9, 8.8$ TeV, respectively, and require $p_T(j) > 20$ GeV for the jets.

We can see that the cross sections for all processes increase with the ion mass. However, the $a\gamma$ cross sections increase much more than the $\gamma\gamma$ ones, an effect related to the larger fraction of energy carried by the ALP. In the following, our analyses focus thus on the p Pb collisions, which offer both a large signal and the best signal-to-background ratio.

Analysis framework—The effective Lagrangian is implemented using Universal FeynRules Output (UFO) format and interfaced with MadGraph5. Event samples are generated for $\gamma\gamma \rightarrow t\bar{t}$, $\gamma a \rightarrow t\bar{t}$, and $\gamma a \rightarrow tj$ processes in p -Pb collisions at the same 8.16 TeV colliding energy as the LHC 2016 data taking, with final-state jets required to have $p_T(j) > 20$ GeV. The ALP flux from the proton is modeled in MadGraph5 following Eq. (3) while the photon flux from Pb is described by the electric dipole form factor implemented in gammaUPC [49], as given in Eq. (8). A transverse momentum boost is then applied to the event samples, distributed according to $p_T/(p_T^2 + Q_0^2)^2$, with $Q_0^2 \approx 0.71$ GeV² [50,51]. The modeling is validated against predictions from SuperChic [52].

For event selection, we focus on the exclusive contribution, as nonexclusive events involving proton dissociation and QCD backgrounds can be suppressed by vetoing calorimeter towers above a certain noise [53,54]. No specific selection is applied to the final-state top quarks or jets, in order to assess the raw event yield within the acceptances of the central detectors. Proton tagging is not required in this analysis, although it can be exploited in single proton-tagged events—notably, p Pb collisions at the LHC exhibit negligible pileup, typically with $\langle\mu\rangle \sim 0.1$ [55]. The pseudorapidity distributions for signal and background are presented in Fig. 2, along with the detector acceptances of LHC experiments. We find that 48% (21%) of $a\gamma \rightarrow t\bar{t}$ events and 36% (8%) of $a\gamma \rightarrow t\bar{c}$ events fall within the LHCb (CMS, ATLAS) acceptance.

Tagging or reconstruction efficiencies and systematic uncertainties are not expected to significantly alter the picture. Lepton and jet reconstruction are based on particle-flow algorithms, which offer high efficiency in particle identification. For jets, a typical transverse momentum threshold of 20 GeV is sufficient to ensure high trigger efficiency [56–58]. These algorithms also provide excellent performance in reconstructing high- p_T leptons, including those from highly boosted top quarks with $p_T > 400$ GeV [59]. When combined with jet tagging techniques, the event selection strategy enables efficient reconstruction of $t\bar{t}$ events. In particular, tailored clustering procedures for boosted $t\bar{t}$ final state have shown to be very effective [59]. Regarding systematic uncertainties, jet energy corrections typically remain $\lesssim 10\%$ for jets with $p_T > 20$ GeV.

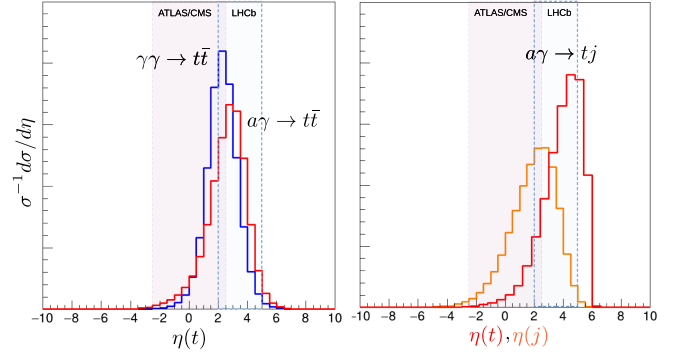


FIG. 2. Pseudorapidity distributions of the final states from the collisions $a\gamma \rightarrow t\bar{t}$, $\gamma\gamma \rightarrow t\bar{t}$ (left) and $a\gamma \rightarrow tj$ (right) in the p Pb mode.

Furthermore, the misidentification rate of light-flavor jets in $t\bar{t}$ events is about 1% [60].

From the above we conclude that p Pb UPCs at the LHC present an opportunity to search for ALP-photon collisions. Notably, a dataset collected in 2016 is already available for analysis.

Projected sensitivity and discussion—Let us evaluate how sensitive the LHC experiments are to the specific ALP-photon collision process investigated here.

At the typical integrated luminosity L_{int} considered, the number of background events is much lower than 1. Thus we apply the simple statistical framework of [61]. Based on the small background, our projected data assume that no event is observed, which sets an upper bound on the signal event rate. A 95% credibility level (CL) corresponds to $\sigma \approx 3L_{\text{int}}^{-1}$ [61].

Combining this statistical method with the experimental analysis framework described above provides the expected sensitivity to the exclusive $a\gamma \rightarrow t\bar{t}$ and $a\gamma \rightarrow tj, \bar{t}j$ processes. This sensitivity is expressed in terms of upper bounds on the ALP couplings. We focus on the $k_u = -k_U$ case, for which there is no CP violation. The probed couplings are c_{tt} and c_{tq} , where $c_{tq} = \sqrt{|c_{tu}|^2 + |c_{tc}|^2}$. The latter combination appears because the processes $a\gamma \rightarrow t\bar{u}$, $a\gamma \rightarrow t\bar{c}$ have, to a very good approximation, the same cross section such that both contribute equally to single-top+jet final state.

We consider the L_{int} corresponding to the 2016 p Pb data collected at ATLAS and CMS (220 nb⁻¹ each) and LHCb (36 nb⁻¹), as well as the realistic projection 1 pb⁻¹ [62]. Putting everything together, we obtain the 95% CL bounds [63]

Det.	ATLAS/CMS		LHCb	
L_{int}	220 nb ⁻¹	1 pb ⁻¹	36 nb ⁻¹	1 pb ⁻¹
(gc_{tt}/f)	(7.2 TeV) ⁻¹	(15.3 TeV) ⁻¹	(4.4 TeV) ⁻¹	(23.1 TeV) ⁻¹
(gc_{tq}/f)	(11.5 TeV) ⁻¹	(24.5 TeV) ⁻¹	(9.9 TeV) ⁻¹	(52.0 TeV) ⁻¹

Let us discuss how these projected bounds compare to the existing bounds on ALP couplings.

A combination of the ALP couplings c_{uu} , c_{dd} , c_{GG} inside the ALP-proton coupling given in Eq. (2) is constrained by kaon decays [22,64]. Specifically, the process $K^+ \rightarrow \pi^+ + a$ leads to anomalous branching ratios that are constrained by the NA62 experiment [65,66]. The exclusion regions are however complex and mass dependent. For example, if the $c_{uu,dd}$ dominate, the kaon bounds vanish completely for $m_a \geq 350$ MeV. Hence we treat the g coupling as an independent parameter.

The c_{tt} coupling is constrained individually in the $m_a \lesssim 5$ GeV region by the B decays [67]. Specifically, the branching ratio $\text{Br}(B \rightarrow K + \text{inv})$ is constrained by BABAR to $(|c_{tt}|/f) \lesssim 1.5 \times 10^{-3} \text{ TeV}^{-1}$ [68]. Notice that $|c_{tt}|$ vanishes if $[k_U]_{33} \approx [k_u]_{33}$.

A strategy to bound the c_{tc} , c_{tu} couplings at the LHC using displaced vertices has been proposed in [69], but applies to ALP masses of order $O(10)$ GeV.

Comparing these bounds to our results, we conclude that the projected sensitivity to c_{tt} from the ALP-photon collision $a\gamma \rightarrow t\bar{t}$ is weaker than the one from B decays. On the other hand, the c_{tc} , c_{tu} couplings are so far unconstrained in the sub-GeV mass range probed in our exclusive processes. We thus find that ALP-photon collisions at the LHC provide a leading probe of the c_{tc} and c_{tu} couplings.

Summary—In this Letter, we point out that if a sub-GeV ALP couples to nucleons, then ALP-ALP and ALP-photon collisions can arise from ultraperipheral collisions at the LHC. This phenomenon provides a new probe of ALPs at TeV-scale energies. Because of the strong suppression of ALP fluxes from ions, the most relevant collision mode is the proton-ion one, particularly $p\text{Pb}$. Focusing on the top quark sector, we demonstrate that ALP-photon collisions from $p\text{Pb}$ UPCs produce a sizable amount of top pairs and top + jet final states, while having vanishing standard model background. The latter channel provides a clean probe of the ALP-top flavor-violating couplings, which remain so far unconstrained for sub-GeV ALPs. These ALP-photon collisions can be readily searched for in the existing 2016 $p\text{Pb}$ datasets collected at ATLAS, CMS, and LHCb.

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Data availability—The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of

preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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