

Coherently enhanced axion-photon conversion via seeded photons for short-pulse axion detection

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We propose a seeded axion-photon conversion scheme to enhance the sensitivity of light-shining-through-a-wall (LSW) experiments for axion detection, where the axions are generated from short-pulse lasers and the usual resonant cavity is not applicable. By injecting a weak, coherent seed electromagnetic (EM) field into the axion-photon conversion region, the axion-induced EM field can constructively interfere with the seed field, amplifying the number of regenerated photons to a level exceeding that of the unseeded scenario. We evaluate the expected signal enhancement, statistical limits from Poisson counting with seed fluctuations and background, and the potential improvement in coupling sensitivity. Compared to a standard LSW setup, the seeded scheme can achieve orders-of-magnitude higher photon yield per axion, potentially surpassing resonance-enhanced experiments in certain parameter regimes. This approach presents a promising pathway to extend the reach of laboratory axion searches, particularly in scenarios where the resonant cavities are impractical.

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I. INTRODUCTION

Axions and axionlike particles (ALPs) are well-motivated hypothetical bosons, originally introduced to solve the strong CP problem [1–5] or extend the standard model [6–11]. Through their two-photon coupling $g_{\gamma\gamma}$, they can be probed in laboratory experiments using high-power lasers and strong magnetic fields. In a typical light-shining-through-a-wall (LSW) experiment [12–17], a high-intensity laser propagates through a strong magnetic field, converting some photons into axions. These axions pass through an opaque wall, then reconvert into photons in a second magnet on the far side of the wall. The regenerated photons would be detected as a signal of axions. Such experiments have been carried out by collaborations like ALPS [12,18], GammeV [14], LIPSS [15,19], OSQAR [16], and BMV [17] and have placed the most stringent laboratory limits on axionlike particle couplings,

excluding photon-axion coupling strengths $g_{\gamma\gamma}$ down to the 10^{-7} – 10^{-8} GeV^{-1} level. However, the sensitivity of single-pass LSW experiments is extremely limited by the tiny photon-axion conversion probability, so more advanced techniques are required to probe weaker couplings.

One successful enhancement is the use of high-finesse Fabry-Perot cavities to resonantly boost the fields in the production and regeneration regions. The ALPS-II experiment, for example, employs a power-buildup cavity on the laser side (raising the effective laser power to ~ 150 kW) and a matched regeneration cavity on the detection side to amplify the reconverted photon number by nearly 4 orders of magnitude [13,20,21]. Together with an interferometric heterodyne readout, ALPS-II is designed to improve sensitivity to $g_{\gamma\gamma}$ down to about 2×10^{-11} GeV^{-1} for axions with mass smaller than 0.1 meV, a dramatic leap over earlier LSW experiments [22].

The rapid development of high-power laser technology and laser plasma studies provided a much stronger background magnetic field, from hundreds to 10^6 T [23–26]. Such strong fields make it possible to generate axions with intense laser pulses, with typical durations ranging from nano- to femtoseconds. However, the axion generated from these intense laser pulses would also be short pulses, making it challenging to enhance the axion regeneration process with the resonant cavities.

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Specifically suited to the wakefield-based axion generation scheme proposed in [27], here we focus on enhancing the photon regeneration process via seeding. In this paper, we study a new approach to amplify the regenerated photon signal in short-pulse axion experiments by introducing a coherent seed electromagnetic field in the regeneration region. The idea is to inject a coherently controlled laser pulse (the seed) into the second magnet simultaneously with the arrival of the axion pulse, such that any axion-induced electromagnetic field will interfere with the seed electromagnetic field. If the seed is tuned to the same frequency and phase as the would-be regenerated photons, the interference will be constructive and effectively boost the detectable photon signal. This is analogous to optical parametric amplification in nonlinear optics [28], where a weak seed signal can stimulate and coherently enhance the conversion from the pump to the signal in a nonlinear crystal. In our case, the seed photons play a similar role to the injected signal light, the axion field acts as the pump, and the background field acts as an existing idler light in the parametric amplification.

A related but distinct approach was proposed by Beyer *et al.* [29], who used a laser to drive axion-photon scattering and mitigate phase mismatch at higher axion masses. Their enhancement arises from three-wave mixing rather than coherent superposition. In contrast, we work in the low-mass regime and amplify regeneration by letting the axion-induced field interfere coherently with a weak, phase-controlled seed behind the wall.

Figure 1 shows the concept: Because of the smallness of $g_{a\gamma\gamma}$, without a seed, the axion-to-photon reconversion yield is extremely small; when a weak, phase-controlled seed is injected, the axion-induced field interferes constructively with the seed, resulting in a larger regenerated photon number.

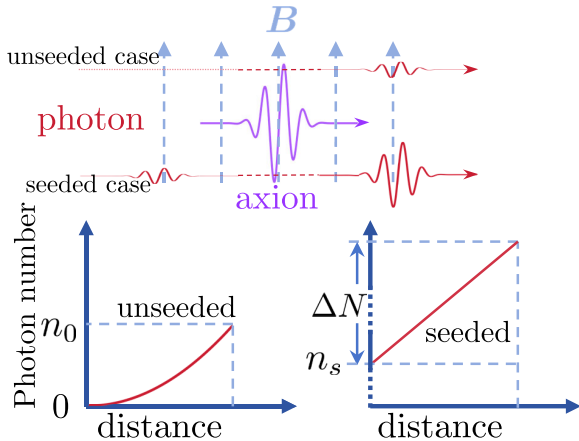


FIG. 1. Conceptual scheme of coherently enhanced axion-photon conversion via seeded photons.

II. AXION-PHOTON CONVERSION WITH A SEED FIELD

In the following, we quantitatively analyze the enhancement of the photon generation by a seed field. We first study the photon regeneration in the absence of a seed, then include a seed field in the process. Consider an axion field $\phi(t, \mathbf{x})$ coupled to the electromagnetic field. In the small mixing and weak-field regime, one can treat the axion-induced electromagnetic field as a perturbation $\mathbf{E}_1$, $\mathbf{B}_1$ on top of any applied background fields (the background being the static magnetic field $\mathbf{B}_0$ in the standard LSW scheme). The modified Maxwell equations are (in natural units $\hbar = c = \epsilon_0 = 1$)

$$\nabla \times \mathbf{B}_1 - \partial_t \mathbf{E}_1 = g_{a\gamma\gamma}(\partial_t \phi) \mathbf{B}_0, \quad (1)$$

along with $\nabla \cdot \mathbf{E}_1 = -g_{a\gamma\gamma} \mathbf{B} \cdot \nabla \phi$, $\nabla \cdot \mathbf{B}_1 = 0$, and $\nabla \times \mathbf{E}_1 + \partial_t \mathbf{B}_1 = 0$. From these, one obtains the wave equation of the vector potential $\mathbf{A}_1$ associated with the induced field,

$$(\partial_t^2 - \nabla^2) \mathbf{A}_1 = g_{a\gamma\gamma}(\partial_t \phi) \mathbf{B}_0, \quad (2)$$

assuming the Coulomb gauge for $\mathbf{A}_1$.

For a monochromatic axion source of angular frequency ω_0 (with corresponding wave number $k_a = \sqrt{\omega_0^2 - m_a^2}$ for axion mass m_a), the axion field can be written as a plane wave $\phi(t, x) = \phi_0 e^{i(k_a x - \omega_0 t)}$ propagating in the $+x$ direction (we take the wall to be at $x = 0$ and the regeneration magnet spanning $0 < x < l_2$). We look for a solution of $\mathbf{A}_1$ with the form $\mathbf{A}_1(t, x) = \hat{\mathbf{e}} \mathcal{A}(x) e^{i(k_1 x - \omega_0 t)}$, where $k_1 = \omega_0$ is the wave number of a photon with the same frequency and $\hat{\mathbf{e}}$ is the polarization unit vector. Here $\mathcal{A}(x)$ is a slowly varying complex amplitude envelope of the wave in the regeneration region. Substituting it into Eq. (2) and assuming $\partial_x \mathcal{A} \ll k_1 \mathcal{A}$ (i.e., the field amplitude $\mathcal{A}$ varies little over a wavelength), we obtain a first-order differential equation,

$$2ik_1 \frac{d\mathcal{A}(x)}{dx} = ig_{a\gamma\gamma} \omega_0 \phi_0 B_0 e^{i(k_1 - k_a)x}. \quad (3)$$

The difference $k_1 - k_a = \kappa$ represents the momentum mismatch between the axion and photon waves. For $m_a \ll \omega_0$ (nearly massless axion), $\kappa \approx \frac{m_a^2}{2\omega_0}$; if $m_a = 0$, then $\kappa = 0$, meaning a perfect phase matching. Integrating Eq. (3) from 0 to x yields the solution of the photon field amplitude at some position x in the regeneration region,

$$\mathcal{A}(x) = \mathcal{A}(0) + \frac{g_{a\gamma\gamma} \omega_0 \phi_0 B_0}{2ik_1} \left[\frac{e^{i\kappa x} - 1}{\kappa} \right]. \quad (4)$$

The first term is the homogeneous solution, representing the evolution of any initial field $\mathcal{A}(0)$ present at the entrance of the region ($x = 0^+$, immediately behind the wall).

The second term is the particular solution sourced by the axion. In a standard LSW experiment with no seed photons behind the wall, one has the boundary condition $\mathcal{A}(0) = 0$. In that case, Eq. (4) reduces to the usual result: the regenerated photon field amplitude $\mathcal{A}_{\text{axion}}(l_2)$ is proportional to $g_{a\gamma\gamma} B_0 l_2$ (for small κl_2), and the probability for an axion to convert into a photon is

$$P_{a \rightarrow \gamma} = \frac{|\mathcal{A}_{\text{axion}}(l_2)|^2}{|\phi_0|^2} \approx \frac{1}{4} g_{a\gamma\gamma}^2 B_0^2 l_2^2 \mathcal{S}(m_a)^2, \quad (5)$$

where $\mathcal{S}(m_a) = \frac{\sin(\kappa l_2/2)}{\kappa l_2/2}$ is a suppressing factor for nonzero m_a arising from phase matching (for $m_a = 0, \kappa = 0, \mathcal{S} = 1$). This reproduces the standard LSW result [30–32]. If N_a axions enter the regeneration region, the expected number of photons produced is simply $N_\gamma = N_a P_{a \rightarrow \gamma}$.

Now consider the scenario with a seed field. We prepare an initial electromagnetic field $\mathcal{A}(0) \neq 0$ behind the wall, consisting of n_s photons at frequency ω_0 (we assume the seed field copropagates collinearly with the axion wave and has a polarization parallel to $\mathbf{B}_0$ to maximize axion-photon conversion). This seed field could, for instance, be derived from a split-off portion of the same laser used to generate the axions with an appropriate time delay, ensuring it enters the regeneration magnet coincidentally with the axion pulse. The presence of $\mathcal{A}(0)$ modifies the solution for $\mathcal{A}(x)$ to Eq. (4).

The total intensity of the electromagnetic wave at $x = l_2$ is proportional to $|\mathcal{A}(l_2)|^2$. To obtain the number of photons, we integrate the intensity over the interaction volume. The result can be written as

$$N'_\gamma = \omega_0 \int_V |\mathbf{A}_1|^2 dV \propto |\mathcal{A}(l_2)|^2, \quad (6)$$

up to normalization constants that cancel out when comparing with the unseeded case. Using Eq. (4) and expanding the magnitude squared, we find

$$N'_\gamma = |\mathcal{A}(0)|^2 + 2 \text{Re}[\mathcal{A}(0)^* \mathcal{A}_{\text{axion}}(l_2)] + |\mathcal{A}_{\text{axion}}(l_2)|^2, \quad (7)$$

where $\mathcal{A}_{\text{axion}}(l_2)$ denotes the particular solution [the second term in Eq. (4)] evaluated at $x = l_2$. Identifying $n_s = |\mathcal{A}(0)|^2$ as the number of seed photons and $n_0 = |\mathcal{A}_{\text{axion}}(l_2)|^2$ as the number of photons that would be regenerated in the absence of a seed [i.e., $n_0 = N_a P_{a \rightarrow \gamma}$ from Eq. (5)], the above expression becomes

$$N'_\gamma = n_s + n_0 + 2\sqrt{n_s n_0} \cos \delta, \quad (8)$$

where δ is the relative phase between the seed field and the axion-induced field at $x = l_2$. In the optimal case, the experiment is configured such that the axion-induced photons emerge in phase with the seed ($\delta = 0$, i.e.,

perfectly constructive interference). In this case, the interference term is maximized and we have

$$N'_\gamma = n_s + n_0 + 2\sqrt{n_s n_0}. \quad (9)$$

This result can be intuitively understood: the axion converts into a photon that is coherently added to an existing electromagnetic field. The $2\sqrt{n_s n_0}$ term represents the cross term where one photon arises from the axion converted field and one from the seed field. Equation (9) serves as the core formula governing the enhancement of seeded regeneration.

We are often interested in the net generated photons attributable to axions. Since n_s seed photons were injected, one may subtract the expected contribution from the seed field and define the increase in the number of signal photons as

$$\Delta N \equiv N'_\gamma - n_s = n_0 + 2\sqrt{n_s n_0}. \quad (10)$$

In the absence of any seed ($n_s = 0$), ΔN correctly reduces to n_0 , but for $n_s \gg n_0$, $\Delta N \approx 2\sqrt{n_s n_0} \gg n_0$. The seed field thus amplifies the detectable photon number increase by a factor

$$\mathcal{E} \equiv \frac{\Delta N}{n_0} = 1 + 2\sqrt{\frac{n_s}{n_0}}. \quad (11)$$

For $n_s \gg n_0$, this simplifies to $\mathcal{E} \approx 2\sqrt{n_s/n_0}$. In other words, the fractional increase in photon signal scales as the square root of the seed photon number (assuming the seed is strong compared to the original signal). Nonetheless, since n_0 is extremely small, in practice, even a modest $n_s \sim 1$ can yield a very large $\sqrt{n_s/n_0}$ factor.

Equation (11) indicates that the benefit of seeding is greatest when $n_s \gg n_0$. For instance, consider a 100 J, 800 nm pulse (about 4×10^{20} photons) traversing a total generation field integral $B_1 l_1 = 3000 \text{ T} \cdot \text{m}$. For $g_{a\gamma\gamma} = 10^{-9} \text{ GeV}^{-1}$ this gives a total number of produced axions of $N_a \approx 8.8 \times 10^8$. With a $B_2 l_2 = 500 \text{ T} \cdot \text{m}$ regeneration region and $m_a \approx 0$, Eq. (5) yields $P_{a \rightarrow \gamma} \approx 6.1 \times 10^{-14}$, so the unseeded expectation per shot is $n_0 = N_a P_{a \rightarrow \gamma} \approx 5.4 \times 10^{-5} \ll 1$. Even a single seed photon ($n_s = 1$) would produce an interference term $\Delta N \approx 2\sqrt{n_0 n_s} \approx 1.4 \times 10^{-2}$, corresponding to an enhancement by about 3 orders of magnitude over the unseeded case.

III. STATISTICAL CASE STUDY

For a realistic seed source, the number of seed photons may fluctuate from shot to shot, so an additional measurement is required to determine the seed photon number n_s . To illustrate how such fluctuations influence the attainable sensitivity, we consider three representative cases summarized in Table I. In all cases, N_{shot} denotes the total

TABLE I. Summary of the statistical configurations used to evaluate the median sensitivity on the axion-induced photon number n_0 .

	Signal	Background
Case 1	$N_{\text{shot}} \cdot n_0$	$N_{\text{shot}} \cdot n_b$
Case 2	On $\frac{1}{2} N_{\text{shot}}(n_s + n_0 + 2\sqrt{n_0 n_s})$ Off $\frac{1}{2} N_{\text{shot}} \cdot n_s$	$\frac{1}{2} N_{\text{shot}} \cdot n_b$ $\frac{1}{2} N_{\text{shot}} \cdot n_b$
Case 3	$N_{\text{shot}} \cdot (n_0 + 2\sqrt{n_0 n_s})$	$N_{\text{shot}} \cdot n_b$

number of laser shots and n_b denotes the mean photon number per shot originating from background.

Case 1: This case represents the baseline situation without any seeded photons. The measured counts follow a Poisson distribution with mean $N_{\text{shot}}(n_b + n_0)$. The median sensitivity at 90% confidence level (CL) [33] on n_0 is evaluated using the Feldman-Cousins unified approach [34], which constructs a confidence interval based on likelihood-ratio ordering for Poisson counting with a known background.

Case 2: Seed photons are introduced but their number fluctuates from shot to shot with a Poisson distribution. We have no knowledge of n_s at all and therefore have to determine it statistically. To achieve this, the measurement is divided into an axion-“on” run (axion conversion enabled) and an axion-“off” run (conversion disabled), both using $N_{\text{shot}}/2$ laser shots. The off run measures only the seed and background photons, giving an expectation $\frac{1}{2} N_{\text{shot}}(n_s + n_b)$, while the on run contains the same background plus the axion-induced term, $\frac{1}{2} N_{\text{shot}}(n_s + n_b + n_0 + 2\sqrt{n_0 n_s})$. The difference between these two runs isolates the coherent axion contribution $n_0 + 2\sqrt{n_0 n_s}$, while naturally accounting for the statistical uncertainty in n_s . Because both on/off counts fluctuate independently according to Poisson statistics, the median sensitivity on n_0 is obtained using the likelihood-ratio method of Rolke *et al.* [35], which extends the Feldman-Cousins formalism to include an auxiliary background measurement with finite statistics.

Case 3: Finally, in an idealized situation where the seed photon number is fully known and reproducible, the measurement reduces to a single Poisson process with an expected signal of $N_{\text{shot}}(n_0 + 2\sqrt{n_0 n_s})$ and the same background $N_{\text{shot}} n_b$. The median sensitivity on n_0 is again evaluated following the Feldman-Cousins prescription.

The corresponding results of median sensitivity on n_0 (thus on $g_{a\gamma\gamma}$) with different seed photon numbers and background levels are shown in Fig. 2. To provide analytical insight, we consider a sufficiently large shot number N_{shot} for each configuration. For case 2, under the approximation $n_s \gg n_b$, $N_{\text{shot}} n_s \gg 1$, when both runs yield the same mean background level $\frac{1}{2} N_{\text{shot}} n_s$, the Rolke method yields an upper limit on the signal of approximately $1.64\sqrt{N_{\text{shot}} n_s}$, where $\sqrt{N_{\text{shot}} n_s}$ reflects the

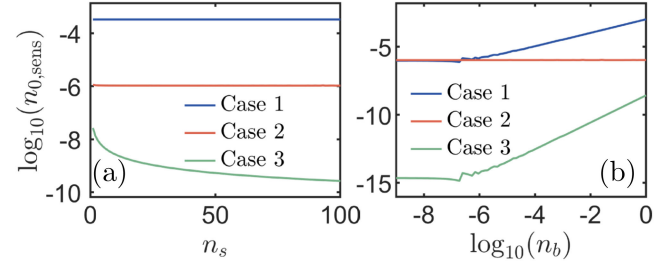


FIG. 2. (a) Median sensitivity (sens) on n_0 at 90% CL versus n_s at fixed $N_{\text{shot}} = 30 \text{ days} \times 1 \text{ Hz}$ and $n_b = 0.1$. (b) Same sensitivities versus n_b at fixed $n_s = 100$ and $N_{\text{shot}} = 30 \text{ days} \times 1 \text{ Hz}$. All sensitivities use the statistical configurations in Table I.

Poisson fluctuation of the seed counts and the factor 1.64 corresponds to the 90% confidence level. In this regime, the on/off difference targets the effective signal, $\frac{1}{2} N_{\text{shot}}(n_0 + 2\sqrt{n_0 n_s}) \approx N_{\text{shot}} \sqrt{n_0 n_s}$, implying an upper limit (UL) on n_0 of $n_{0,\text{UL}} \approx 1.64^2 / N_{\text{shot}}$ that is essentially independent of n_s . For case 1, when the background is significant ($N_{\text{shot}} n_b \gg 1$), the upper limit on the signal is about $1.64\sqrt{N_{\text{shot}} n_b}$, while the target signal is $N_{\text{shot}} n_0$, hence $n_{0,\text{UL}} \approx 1.64\sqrt{n_b / N_{\text{shot}}}$. In this situation, although case 2 requires on/off measurements, it can still outperform case 1 to some extent.

However, when the background is very low ($N_{\text{shot}} n_b \ll 1$), the Feldman-Cousins method yields an upper limit of about 2.4 events on the signal ($N_{\text{shot}} n_0$) at 90% CL. Thus, case 1 gives $n_{0,\text{UL}} \approx 2.4 / N_{\text{shot}}$. In this regime, the limits from case 1 and case 2 become nearly identical. These analytical upper-limit scalings describe the asymptotic behavior of the expected median sensitivity shown in Fig. 2. In all scenarios, when the seed photon number is perfectly known (case 3), the coherent enhancement term $2\sqrt{n_0 n_s}$ yields the lowest attainable upper limit on n_0 . These comparisons show that the precision in characterizing the seed photon number directly determines how effectively the coherent enhancement can be utilized in setting the axion-photon coupling limit.

It is also important to note that strict shot-to-shot reproducibility of the seed is not essential: what matters is that we know the photon number in each shot with sufficient certainty. Given that the initial photon number of the seed field can be measured with high precision, then shot-to-shot variations can be corrected in the data analysis. This may be done by using correlated photon sources, such as pair production processes (e.g., laser-induced double-photon emission process), where one photon of the pair can be used to monitor the seed photon count, and the other one is injected as the actual seed. In such a scheme, the seed photon fluctuations can be strongly reduced, because the per-shot photon number is tagged by the reference measurement.

Another important consideration is the phase alignment between the seed and axion-induced field. In writing

Eq. (9) we assumed perfect constructive interference ($\cos \delta = 1$). If the relative phase fluctuates from shot to shot, the interference term is reduced after averaging, thereby degrading the seeded enhancement. As a simple model, suppose δ is uniformly distributed in $(-\alpha, \alpha)$. Since the phase statistics are independent of n_0 , the phase averaging can be treated separately, yielding an effective coherence factor

$$f(\alpha) = \frac{1}{2\alpha} \int_{-\alpha}^{\alpha} \cos \delta \, d\delta = \frac{\sin \alpha}{\alpha}. \quad (12)$$

After accumulating N_{shot} shots, the expected net photon-number change becomes

$$\Delta N_{\text{ph}} = N_{\text{shot}}(n_0 + 2\sqrt{n_0 n_s} f(\alpha)) \approx 2N_{\text{shot}} \sqrt{n_0 n_s} f(\alpha). \quad (13)$$

Thus, for a fixed $g_{a\gamma\gamma}$, the expected signal is reduced by a factor $f(\alpha)$. Conversely, for a fixed observed photon-number change, the corresponding upper limit on $n_0 \propto N_a g_{a\gamma\gamma}^2 \propto g_{a\gamma\gamma}^4$ increases by a factor $1/f(\alpha)^2$, implying that the upper limit on $g_{a\gamma\gamma}$ weakens by a factor $1/\sqrt{f(\alpha)}$. A moderate phase uncertainty only mildly degrades the sensitivity [e.g., $1/\sqrt{f(\pi/2)} \approx 1.25$ for $\alpha = \pi/2$]. However, if the phase is completely randomized over $(-\pi, \pi)$, the interference term averages to zero and the seeded enhancement is entirely lost. Therefore, phase control is mandatory to realize the seeded enhancement. In practice, the axion field phase is set by the phase of the initial laser in the production region and any phase shift accumulated in transit. If the seed is split from the same laser source, one can maintain a fixed phase relationship via interferometric stability over the beam paths.

IV. IMPLICATIONS FOR SENSITIVITY

The ultimate goal of an LSW experiment is to detect or constrain the axion-photon coupling $g_{a\gamma\gamma}$. We therefore convert the photon-counting observables into the median sensitivity on $g_{a\gamma\gamma}$, explicitly incorporating the per-shot background n_b in the sensitivity estimation.

An important advantage of using short-pulse, ultraintense laser drivers for axion generation and detection is that the extremely short temporal window of the interaction effectively suppresses the number of background photons per laser shot. For a typical dark-count rate of about 10 Hz in a single-photon detector, axion-induced photons are expected to appear only within the laser pulse duration of tens of femtoseconds. Consequently, the background photon number per shot can be suppressed down to the level of $n_b \sim 10^{-15}$. Even when taking into account the detector's timing gate (typically of nanosecond duration), the effective background photon number remains as low as $n_b \sim 10^{-8}$ per shot. Therefore, the background $N_{\text{shot}} n_b$ is

negligible even with $N_{\text{shot}} = 30 \text{ days} \times 1 \text{ Hz} = 2.6 \times 10^7$. Under such low-background conditions, cases 1 and 2 yield nearly identical median sensitivities on n_0 (and thus on $g_{a\gamma\gamma}$). In the following discussion, we mainly compare cases 1 and 3. For a null-result experiment, the Feldman-Cousins method gives a 90% CL upper limit of about 2.4 signal photons, which we adopt to evaluate the median sensitivity on $g_{a\gamma\gamma}$.

For a standard unseeded LSW experiment, the median sensitivity on $g_{a\gamma\gamma}$ is given by $N_{\text{shot}} n_0 = 2.4$. Writing $N_{\text{shot}} = RT_{\text{run}}$ (with R the repetition rate of the pulses and T_{run} the total running time), we have $g_{a\gamma\gamma}$ sensitivity (in case of no detection) by

$$g_{a\gamma\gamma}^{\text{sens}} \approx 2.5(RT_{\text{run}} n_L)^{-1/4} (B_1 l_1 B_2 l_2 \mathcal{S}_1 \mathcal{S}_2)^{-1/2}, \quad (14)$$

where n_L is the number of photons per laser shot [related to laser pulse energy $P_L \tau_L / (\hbar \omega_0)$], B_1, l_1 are magnet field strength and length for the axion generation part, B_2, l_2 are those for the photon regeneration, and $\mathcal{S}_{1,2}$ are phase mismatch suppression factors for each stage. Equation (14) essentially comes from $n_0 = \frac{1}{4} g_{a\gamma\gamma}^4 B_1^2 l_1^2 B_2^2 l_2^2 n_L$ and the requirement $RT_{\text{run}} n_0 = 2.4$. Plugging in some representative numbers, where axions are produced in the generation stage via a laser-wakefield interaction [27] and detected through an LSW setup (without resonance in the regeneration cavity): $R = 1 \text{ Hz}$, $T_{\text{run}} = 30 \text{ days}$, $n_L = 4 \times 10^{20}$ (100 J at 800 nm), $B_1 l_1 = 3000 \text{ T} \times 1 \text{ m}$ (as can be realized in a $\sim 1 \text{ m}$ laser-wakefield interaction stage), $B_2 l_2 = 5 \text{ T} \times 100 \text{ m}$ (as can be provided by a superconducting magnet string comparable to ALPS-II), and assuming near-optimal $\mathcal{S}_1 \approx \mathcal{S}_2 \approx 1$, we find

$$g_{a\gamma\gamma}^{\text{sens}} \approx 4.3 \times 10^{-10} \text{ GeV}^{-1}. \quad (15)$$

These parameter choices are intended as representative values for an illustrative sensitivity estimate. In practice, they can be adapted to specific experimental setups. The sensitivity without the seed photon is indeed weaker than ALPS-II (by almost an order of magnitude). The lack of resonant buildup on both the generation and regeneration cavities contributes to this relatively high floor on $g_{a\gamma\gamma}$.

In the seeded experiment, the criterion for detection changes: we require $RT_{\text{run}} \Delta N = 2.4$ (same upper limit at 90% CL from the Feldman-Cousins method). Using $\Delta N \approx 2\sqrt{n_0 n_s}$ (valid if $n_s \gg n_0$) and noting $n_0 = g_{a\gamma\gamma}^4 C$ where $C = \frac{1}{16} n_L B_1^2 l_1^2 B_2^2 l_2^2$ collects all the experimental constants, we get

$$g_{a\gamma\gamma}^{\text{sens}} \approx 2.2(R^2 T_{\text{run}}^2 n_L n_s)^{-1/4} (B_1 l_1 B_2 l_2 \mathcal{S}_1 \mathcal{S}_2)^{-1/2}. \quad (16)$$

Comparing with Eq. (14), we see that effectively the factor RT_{run} is under a square root instead of a fourth root, and n_s also appears under a fourth root in the denominator. This implies that long integration times and high repetition rates

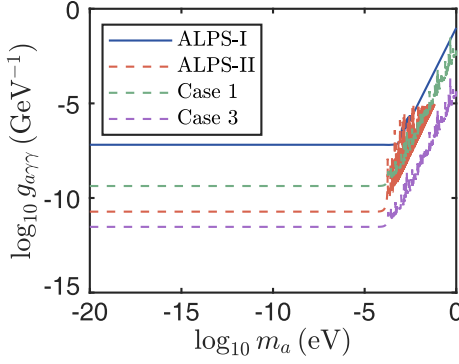


FIG. 3. Projected sensitivity to $g_{a\gamma\gamma}$ for a pulsed LSW experiment with and without seeding, compared to ALPS experiment [13,36].

are more powerful in improving sensitivity for the seeded case than in the unseeded case. Inserting the same numerical values as above and taking $n_s = 100$ photons, Eq. (16) gives

$$g_{a\gamma\gamma}^{\text{sens}} \approx 3.0 \times 10^{-12} \text{ GeV}^{-1}, \quad (17)$$

for the seeded experiment. Figure 3 illustrates this comparison: in the unseeded, single-pass configuration, our projected limit on $g_{a\gamma\gamma}$ remains weaker than the ALPS-II sensitivity; by contrast, introducing a modest seed of $n_s = 100$ photons per shot coherently enhances the regeneration and pushes the reach below the ALPS-II constraint. Of course, this dramatic gain assumes ideal interference and negligible additional noise beyond shot noise. It shows, nonetheless, that seeded regeneration can partially compensate for the lack of resonant enhancement in pulsed experiments.

V. CONCLUSION

We have proposed a seeded photon regeneration scheme for axion searches with high-power laser pulses. We show

that by injecting a coherent seed electromagnetic field into the regeneration region, the coherent interaction between the axion field and the photons can enhance the axion-induced photon signal by orders of magnitude, thereby improving the experimental sensitivity to the axion-photon coupling. The method circumvents the incompatibility of resonant cavities with femtosecond pulses, operates effectively in the time domain, and remains advantageous even under realistic noise constraints, provided that seed intensity fluctuations are monitored via independent means. This approach motivates further studies and has the potential to significantly extend the reach for axions through the pulsed light-shining-through-a-wall experiment.

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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.

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