

Broadband interferometry-based searches for photon-axion conversion in vacuum

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A novel experiment is introduced to detect photon-axion conversion independent of the dark-matter hypothesis in a broad mass-range called WISP Interferometer (WINTER). The setup consists of a free-space Mach-Zehnder-type interferometer incorporating an external magnetic field and vacuum in one of the arms, where photon-axion mixing occurs via the Primakoff effect and is detected through changes in amplitude. The expected axion-induced signal is then modulated by polarization changes. The experiment is designed to integrate a Fabry-Pérot cavity with a finesse of 10^5 that will be operated in a vacuum environment, significantly enhancing the sensitivity. It is projected to reach the Dine-Fischler-Srednicki-Zhitnitsky theoretical line with photon-axion coupling sensitivities down to $g_{a\gamma\gamma} \simeq 3.7 \times 10^{-14} \text{ GeV}^{-1}$ for axion masses up to $380 \text{ } \mu\text{eV}$.

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I. INTRODUCTION

Axions are hypothetical particles that arise in various extensions of the Standard Model of particle physics. Originally proposed to resolve the strong CP problem in QCD [1,2], axions have become a prominent candidate for dark matter due to their weak interactions with normal matter and their ability to exist in large quantities in the Universe without being directly detected [3,4]. A key aspect of detecting axions is their coupling to photons in the presence of an external magnetic field. This interaction allows for the conversion of photons into axions and vice versa when they traverse regions with strong magnetic fields.

Several experimental approaches have been developed to exploit this axion-photon coupling. Haloscopes aim to detect dark matter axions by converting them into photons using a low-loss resonant cavity embedded in a strong magnetic field [4]. Although very sensitive, these experiments are highly model dependent, relying on unverified assumptions about the local density and distribution of dark matter. Based on several N-body simulations, the local dark matter distribution is expected to diverge from the currently widely assumed standard halo model [5,6]. This prediction

severely limits the effective sensitivity of haloscope measurements. Experiments using the light-shining-through-the-wall (LSW) technique are model independent and intrinsically broadband, thus covering a large portion of the available axion parameter space [7]. However, they require conversion of photons into axions and vice versa, which in turn limits their sensitivity.

Interferometric experiments [8] have emerged as an innovative approach in the search for axions due to their exceptional sensitivity to very small changes in amplitude, phase, and polarization. The recently proposed Weak Interacting Slim Particle searches on a Fiber Interferometer (WISPF) experiment [9] benefits from resonant conversion in a fiber-based interferometer. WISPF allows for adjusting the refractive index by varying the gas pressure in the hollow core of a photonic crystal fiber so that the resonant condition can be tuned to reach the Dine-Fischler-Srednicki-Zhitnitsky (DFSZ-)predicted photon-axion coupling.

Motivated by this approach, a novel broadband experiment called Weak Interacting Slim Particle Interferometer (WINTER) is proposed in this paper. WINTER uses a free-space Mach-Zehnder interferometer (MZI) to detect the photon flux reduction due to the conversion of photons to axions according to the Primakoff effect [10]. The key element of WINTER is the placement of the sensing arm of the interferometer in vacuum and under a strong magnetic field, which, together with the implementation of a Fabry-Pérot cavity (FPC), boosts the expected sensitivity.

This approach makes WINTER a unique broadband, model-independent, and highly sensitive experiment

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capable of detecting axionlike particles (ALPs)¹ and axions beyond the theoretical axion band and covering a large unexplored portion of the axion parameter space. This region partially aligns with the predictions from lattice QCD, which estimate that axions making up a portion of dark matter in a postinflationary scenario would have masses in the range 50 μeV –1500 μeV [11].

II. PHOTON-AXION MIXING IN VACUUM

In the following, the axion is treated as a classical pseudoscalar field, a , that mixes with the electric field in the presence of an external magnetic field. The mass of the axion, m_a , is assumed to be much smaller than the energy of the photons as expressed by their angular frequency, ω_γ , so that the condition $|m_a|^2 \ll \omega_\gamma^2$ is always satisfied. In the plane-wave approximation, a Helmholtz-like equation for the state $\psi(z) = \begin{pmatrix} E(z) \\ a(z) \end{pmatrix}$ can be derived [12],

$$[\partial_z^2 + M^2]\psi(z) = 0, \quad (1)$$

where z is the direction of propagation. M^2 is the mixing matrix,

$$M^2 = \frac{1}{2\omega_\gamma} \begin{pmatrix} k_\gamma^2 & -G \\ -G & k_a^2 \end{pmatrix}, \quad (2)$$

where $k_\gamma^2 = \omega_\gamma^2 - m_\gamma^2$ and $k_a^2 = \omega_a^2 - m_a^2$ are the photon and axion momenta, respectively, and $G = g_{a\gamma\gamma} B_{\text{ext}}/2$ is the mixing energy which is proportional to the external magnetic field and the coupling between the axion and photon $g_{a\gamma\gamma}$.

The resulting probability of photon-to-axion conversion, $P_{\gamma \rightarrow a}$, is calculated in a base rotated by an angle θ such that

$$P_{\gamma \rightarrow a} = \sin^2(2\theta) \sin^2(k_{\text{osc}} z), \quad (3)$$

where $k_{\text{osc}} = \sqrt{G^2 + \left(\frac{m_a^2}{4\omega_\gamma}\right)^2}$ is the oscillation wave number under the approximation $\omega_\gamma \approx \omega_a \gg m_a$ and $m_\gamma = 0$ (vacuum conditions). The mixing angle, θ , calculated by diagonalizing the mixing matrix, is given by

$$\tan(2\theta) = \frac{4\omega_\gamma G}{k_\gamma^2 - k_a^2}, \quad (4)$$

and the conversion probability can be rewritten as

¹The experiment is also sensitive to photon-scalar coupling.

$$P_{\gamma \rightarrow a} = \frac{G^2}{G^2 + \frac{m_a^4}{16\omega_\gamma^2}} \sin^2 \left(\sqrt{G^2 + \frac{m_a^4}{16\omega_\gamma^2}} z \right). \quad (5)$$

For $k_{\text{osc}} z \ll 1$, the sine function can be expanded, with the resulting (mass-independent) conversion probability

$$P_{\gamma \rightarrow a} \simeq (Gz)^2 = \frac{g_{a\gamma\gamma}^2 B_{\text{ext}}^2}{4} z^2. \quad (6)$$

III. INTERFEROMETRIC DETECTION

The principle behind detecting photon-axion conversion in a MZI has already been discussed in detail for the WISPMI experiment [9]. The experimental scheme of WINTER is shown in Fig. 1. WINTER uses an electro-optical modulator (EOM) for modulating the laser beam in amplitude (EOM-AM) while switching between the two orthogonal states of linear polarization passing through a Pockels cell (EOM-PC). The amplitude modulation is performed using a sinusoidal-amplitude modulated electric field,

$$\mathbf{E}(\mathbf{t}) = \frac{E_0}{\sqrt{2}} \begin{pmatrix} 1 + \beta_m \sin(\omega_m t + \chi) \\ 1 + \beta_m \sin(\omega_m t) \end{pmatrix} e^{i\omega_\gamma t}, \quad (7)$$

with an amplitude E_0 , a modulation depth β_m , an amplitude-modulation frequency ω_m , and a phase offset $\chi \approx \pi/2$ between the two modulation signals. The external magnetic field, B_{ext} , is aligned to be perpendicular to the direction of the wave propagation and parallel to one of the linear polarization modes of the laser beam. The sensing and reference arm lengths of the interferometer (L_{sen} and L_{ref}) are chosen to differ by $\Delta L = L_{\text{sen}} - L_{\text{ref}}$. The total length is given by $L = (L_{\text{sen}} + L_{\text{ref}})/2$.

The optical path length difference ΔL required for destructive interference at the dark fringe corresponds to a large multiple of half-wavelengths. Experimentally speaking, imperfections in alignment, noise, and locking stability introduce a small deviation from the ideal condition. Therefore, the actual path length difference can be written as

$$\Delta L = (2n + 1) \frac{\lambda_\gamma}{2} + \sigma_\gamma, \quad (8)$$

where $n \in \mathbb{Z}$, λ_γ is the laser wavelength, and σ_γ represents the uncertainty in locking on the dark fringe. Applying ΔL for the destructive interference condition, we obtain the residual phase deviation for the carrier,

$$k_\gamma \Delta L = (2n + 1)\pi + \alpha_\gamma, \quad (9)$$

with $k_\gamma = 2\pi/\lambda_\gamma$ expressing the laser wave number and $\alpha_\gamma = k_\gamma \sigma_\gamma$ quantifying the phase error from imperfect locking. The respective destructive interference condition

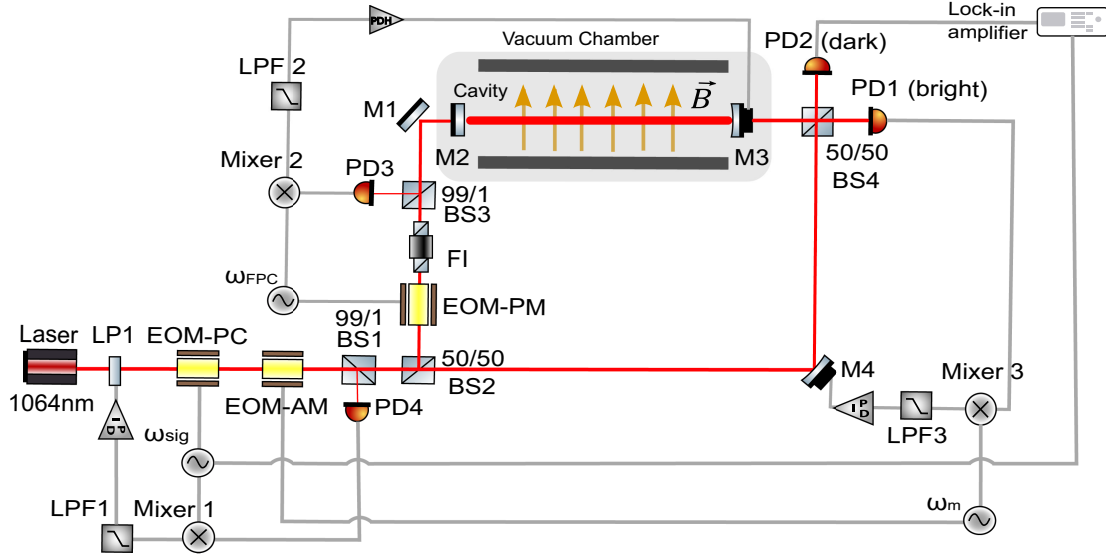


FIG. 1. Schematic view of the experimental setup of WINTER using a free-space MZI for broadband detection. In red, the laser beam in free-space is shown. The sensing arm of the interferometer is placed inside a vacuum chamber with a FPC of 10m length integrated in a dipole magnet with 9T of the same length. The acronyms correspond to: electro-optical modulator (EOM) with amplitude (EOM-AM), phase (EOM-PM), and polarization (EOM-PC) modulation, Faraday isolator (FI), linear polarizer (LP), beam-splitter (BS), photodetector (PD), mirror (M), and low-pass filter (LPF). More details about the setup are given in Sec. IV.

for amplitude-modulated sidebands is then given by

$$k_m \Delta L = (2p + 1)\pi + \alpha_m, \quad (10)$$

where $\alpha_m = \frac{k_m \alpha_\gamma}{k_\gamma}$ is the sideband phase error, $k_m = \omega_m/c$ is the modulation wave number, c is the speed of light, and $p \in \mathbb{Z}$ [with $p \ll n$ decoupling effectively the two conditions in Eqs. (9) and (10)]. Since $k_m \ll k_\gamma$, the sideband phase error, α_m , can be neglected. In the presence of an external magnetic field and nonvanishing photon-axion conversion probability $P_{\gamma \rightarrow a}$, a modulated signal is expected at the dark port originating from the combined interference of the sidebands with each other and with the carrier (for the full derivation, see Appendix A).

To lock on, for example, the upper sideband at $\omega_\gamma + \omega_m$, we first demodulate the resulting power from port 2 (dark port) at the amplitude modulation frequency ω_m . This demodulation scheme is performed using lock-in amplification by mixing the dark port signal with a local oscillator (LO) at ω_m following a homodyne procedure (see Appendix B for details), which stabilizes the interferometer at a fixed working point in amplitude. Terms arising from the self-interference of the carrier and sidebands are omitted, as their contributions are negligible after demodulation and averaging the power over long time. Once the interferometer is locked, the signal itself is demodulated at the polarization modulation frequency, ω_{sig} , to track the slow amplitude variations induced by the polarization changes. In both demodulation steps, the output of the mixer is passed through a low-pass filter with a bandwidth

much lower than the respective modulation frequency, ensuring that only slow or dc variations of the signal are retained. The time-averaged output power of the signal after the mixer is then

$$P_{\text{av}} = \frac{1}{2} P_{\text{tot}} \beta_m P_{\gamma \rightarrow a}, \quad (11)$$

where P_{tot} is the input power of the laser. The effect of the polarization modulation through the EOM-PC is accounted by time averaging the conversion probability (for more details, see Sec. B).

IV. EXPERIMENTAL SETUP

The general setup for the WINTER experiment shown in Fig. 1 is based on a free-space MZI, as introduced above. One of the arms (sensing arm) is embedded inside the bore of a strong dipole magnet and uses a FPC under high-vacuum conditions, while the other one acts as a reference arm. In the following, we assume a magnet similar to the LHC test dipole magnets, like the one used by the CERN Axion Solar Telescope (CAST) experiment [13], with a length of ≈ 10 m and a magnetic field strength of ≈ 9 T.

The FPC increases the total path length of the photons within the magnetic field. Therefore, it increases the conversion probability $P_{\gamma \rightarrow a}$ [Eq. (6)] in a quadratic way and improves the sensitivity because of the effectively larger circulating power. The FPC employed in WINTER comprises of two mirrors (a planar and a concave one) with an aperture radius of 17.5 cm and a high-power continuous-wave laser operating at a wavelength of 1064 nm. With this

optical scheme, diffraction and clipping losses are minimized to not degrade the finesse of the cavity (see Appendix C for details). Given the oscillation wave number, k_{osc} , defined in Sec. II assuming vacuum conditions with the photon-axion mixing energy satisfying $G \ll m_a^2/(4\omega_\gamma)$, the oscillation length of the axion field is estimated to be

$$L_{\text{osc}} = \frac{\pi}{k_{\text{osc}}} \simeq \frac{4\pi\omega_\gamma}{m_a^2} \simeq 10 \text{ m} \left(\frac{m_a}{0.38 \text{ MeV}} \right)^{-2}. \quad (12)$$

For the given length of the magnet, the maximum length of the cavity $L_{\text{FPC}} \approx 10 \text{ m}$, results in a photon-axion conversion taking place up to a mass of $\approx 0.5 \text{ MeV}$.

In the following, we describe the necessary control loops and associated modulations for the experiment:

- (A) matching the power between the two polarization states (Sec. IV A), using the polarization modulation at frequency ω_{sig} (EOM-PC);
- (B) locking the FPC at its working point (Sec. IV B), using phase modulation at frequency ω_{FPC} (EOM-PM); and
- (C) locking the interferometer near a dark fringe (Sec. IV C), using amplitude modulation at frequency ω_m (EOM-AM).

A. Power matching between the polarization states

A key aspect of the experiment is the modulation of the expected signal from photon-axion mixing at the frequency ω_{sig} [see Eq. (7)], such that it can be reliably distinguished from other instrumental effects. For photons with a lifetime of τ_{ph} inside the FPC, the condition $\omega_{\text{sig}} < 2\pi/\tau_{\text{ph}}$ has to be satisfied (see Sec. IV B for further details). The modulation is achieved by changing the polarization of the light before it enters the interferometer. In the proposed setup (see Fig. 1), this is implemented in a two-step process. First, the laser's linearly polarized light passes through a Proportional–Integral–Derivative (PID)-controlled motorized linear polarizer (LP1), which is set to an angle close to 45 deg relative to the initial polarization direction, such that the resulting field is shared equally between the two linear polarization states. In a second step, the Pockels cell (EOM-PC) is driven at the frequency ω_{sig} to create a modulated signal as given by Eq. (7). To maintain an equal share of power in the two polarization modes, a beam splitter (99/1 BS1) is used to read out the error signal with a photodiode (PD4), which feeds into a PID loop that changes the direction of LP1.

A free-space beam splitter (BS1) is then used to split the modulated beam equally and redirect it to the two interferometer arms.

B. Cavity locking

The FPC used in the WINTER experiment consists of a plane mirror (M2) and a concave piezo-driven mirror (M3) with relative reflectivity R separated by the distance

$L_{\text{FPC}} = 10 \text{ m}$ as shown in Fig. 1. The resulting finesse $\mathcal{F}$ of the FPC is given by

$$\mathcal{F} = \frac{\pi\sqrt{R}}{1-R} = \frac{\text{FSR}}{\Delta f}, \quad (13)$$

which is the ratio between the free-spectral range ($\text{FSR} = c/2L_{\text{FPC}} = 15 \text{ MHz}$) and the line width Δf of the cavity modes. For the planned experiment, a transmission loss value of $1 - R = 30 \text{ ppm}$ is assumed for each mirror, which, accounting for diffraction and clipping losses, it translates into a finesse of $\mathcal{F} \approx 10^5$ and a line width of $\Delta f = 150 \text{ Hz}$. The effects of diffraction in the conversion region and thermal load in the mirrors are discussed in the Appendixes Sec. C.

The locking of the FPC involves adjusting its length using a piezo-driven mirror (M3) to match $L_{\text{FPC}} = N\frac{\lambda}{2}$ where N corresponds to the cavity mode number assuming vacuum conditions. In this case, the FPC becomes transparent for the incoming laser, no light is reflected, and the circulating power reaches its maximum value.

The position of the piezoelectric mirror M3 is adjusted by minimizing the sideband signal reflected from the mirror M2 and measured with the photodiode PD3. The sideband is generated by modulating the phase of the carrier with the EOM-PM, following the well-established Pound-Drever-Hall locking scheme [14]. A FI is placed after the BS2 to prevent any back reflections from the FPC. The modulation frequency driving the EOM-PM for locking the cavity is then chosen to be half of the FSR at $f_{\text{FPC}} = \omega_{\text{FPC}}/(2\pi) = 7.5 \text{ MHz}$ so that it does not interfere with the modulation of the signal and the stabilization of the interferometer.

The modulation frequency, ω_{sig} , of the EOM-PC has to be chosen to be sufficiently small to allow the power in the FPC to build up during the modulation period. This timescale should be larger than the photon lifetime:

$$\tau_{\text{ph}} = \frac{1}{\text{FSR}(1-R)}. \quad (14)$$

Therefore, the resulting frequency $f_{\text{sig}} = \omega_{\text{sig}}/(2\pi)$ should fulfill the requirement

$$f_{\text{sig}} < \text{FSR}(1-R) \approx 480 \text{ Hz}. \quad (15)$$

The final choice of f_{sig} will be optimized to strike a balance between the required power build-up in the cavity and the low-frequency noise. The frequency could be swept to verify that the response is flat in the relevant frequency range. We settle here for a value of $f_{\text{sig}} \approx 20 \text{ Hz}$ where the average power in the polarization mode parallel to the magnetic field is $\approx 39\%$ and a modulation index between the two polarization modes of $\beta_{\text{sig}} \approx 1$.

Note that the input laser beam requires a line width smaller than the FPC line width to guarantee an optimal coupling. For this experiment, the former will be on the

order of 0.1 kHz [15]. This ensures that the coherence length of the laser is sufficiently large to match the total optical path length of the FPC.

C. Interferometer stabilization

The interferometer needs to be phase locked at a working point, which, for WINTER, is chosen to be close to 0.01% of the dark fringe power for PD2. This reduces the corresponding shot noise from the laser and improves the sensitivity (see also Sec. V).

The lock is achieved by driving the piezoelectric mirror M4 to compensate phase changes between the two arms, including residual phase contributions from the optical elements, which are tracked using the error signal generated from PD1, demodulated at the frequency ω_m , the frequency at which the EOM-AM modulates the amplitude of the carrier. The amplitude-modulation frequency is chosen sufficiently high, $f_m = \omega_m/(2\pi) \approx 1$ kHz, to avoid overlap with the polarization modulation frequency $f_{\text{sig}} \approx 20$ Hz and to reduce the impact of $1/f$ noise from the laser.

The axion-induced signal is read out at PD2 once the interferometer is locked near the dark fringe. This is done by demodulating the PD2 signal at the polarization-modulation frequency ω_{sig} , applied via the EOM-PC, which tags the axion-sensitive component of the optical field.

V. SENSITIVITY ANALYSIS

The expected sensitivity of the WINTER setup will be estimated assuming that the MZI is operated close to a dark fringe. The main noise sources are the dark current of the PD2 and the shot noise contribution from the sidebands. The resulting signal-to-noise ratio (SNR) is then given by the ratio of the signal power received at the dark port (port 2) from photon-axion conversion [see Eq. (11)] and the noise power contributed by the PD2 and by the shot noise of the laser in the dark fringe.

For an InGaAs photodiode, the noise-equivalent power from the dark current (NEP_{PD2}) amounts to approximately $0.7 \text{ fW}/\sqrt{\text{Hz}}$ [16]. The estimated contribution of the shot noise from the sidebands is derived from the shot noise from the residual laser signal in the dark fringe integrated over the bandwidth of the sideband. The resulting NEP_{SN} from the shot noise is estimated,

$$\text{NEP}_{\text{SN}} = \sqrt{\frac{2 \cdot P_{\text{dark}} \cdot E_{\text{ph}}}{e}} = 73.4 \text{ pW}/\sqrt{\text{Hz}}, \quad (16)$$

where $P_{\text{dark}} \approx 10^{-4} P_{\text{tot}}$ is the residual power received at the photodetector in the dark fringe. E_{ph} is the photon's energy, and e is the detection efficiency of the photodetector (PD2), which is assumed to be 90%. For a commercial laser with $P_{\text{tot}} = 130 \text{ W}$ and $\lambda_{\gamma} = 1064 \text{ nm}$, the effective sideband noise power for a bandwidth of $\Delta\nu = 10 \text{ Hz}$ is then

$$P_{\text{SN+PD2}} = \sqrt{(\text{NEP}_{\text{PD2}}^2 + \text{NEP}_{\text{SN}}^2) \Delta\nu} = 232.2 \text{ pW}. \quad (17)$$

Considering that the beam traverses a path length of $L_{\text{FPC}} = 10 \text{ m}$, in a FPC with $\mathcal{F} = 10^5$ finesse, embedded into a 9 T LHC-type dipole magnetic field of equal length, and assuming an operation time of $t = 1$ year, the resulting sensitivity for $g_{\text{a}\gamma\gamma}$ in the broadband range is given by

$$\begin{aligned} g_{\text{a}\gamma\gamma} \approx & 3.7 \times 10^{-14} \text{ GeV}^{-1} \left(\frac{\text{SNR}}{3} \right)^{1/2} \left(\frac{B_{\text{ext}}}{9 \text{ T}} \right)^{-1} \left(\frac{L_{\text{FPC}}}{10 \text{ m}} \right)^{-1} \\ & \times \left(\frac{C_{\text{dut}}}{0.024} \right)^{-1/2} \left(\frac{P_{\text{tot}}}{130 \text{ W}} \right)^{-1/2} \left(\frac{\beta_m}{1} \right)^{-1/2} \left(\frac{\beta_{\text{sig}}}{1} \right)^{-1} \\ & \times \left(\frac{t}{1 \text{ year}} \right)^{-1/4} \left(\frac{\text{NEP}_{\text{SN+PD2}}}{73.4 \text{ pW}/\sqrt{\text{Hz}}} \right)^{1/4} \left(\frac{\Delta\nu}{10 \text{ Hz}} \right)^{1/8} \\ & \times \left(\frac{E_{\text{ph}}}{1.16 \text{ eV}} \right)^{1/4} \left(\frac{\mathcal{F}}{1 \times 10^5} \right)^{-1/2}, \end{aligned} \quad (18)$$

where C_{dut} is a constant value related to the duty cycle accounting for the 39% average power of the polarization mode parallel to the magnetic field inside the FPC, as well as the factor $1/16$ from Eq. (B7) resulting from the time-averaging of the sideband modulation ω_m , the time-averaging of the signal modulation ω_{sig} , and the sine expansion of conversion probability.

The expected sensitivity for WINTER is shown in Fig. 2 where the resulting limit on the axion-photon coupling is on the level of $g_{\text{a}\gamma\gamma} \gtrsim 3.7 \times 10^{-14} \text{ GeV}^{-1}$ for axion masses up to $380 \mu\text{eV}$. This level of sensitivity allows WINTER to probe the DFSZ model region, representing a significant advancement in exploring the QCD axion parameter space at these masses.

Using instead the string of dipole magnets employed in the ALPS-II experiment, providing a magnetic field strength of $B_{\text{ext}} = 5.3 \text{ T}$ over a total magnetic length of $L = 200 \text{ m}$ [17] and assuming an optical resonator with a reduced finesse of $\mathcal{F} = 1.2 \times 10^4$, the sensitivity can be rescaled according to the explicit $B_{\text{ext}}^{-1} L^{-1} \mathcal{F}^{-1/2}$ dependence in Eq. (18). The finesse is chosen to maximize intra-cavity power buildup, while accounting for the increased cavity length, which leads to a longer photon storage time and correspondingly tighter constraints on the cavity line width and optical losses. Despite the reduced magnetic field strength compared to the LHC setup, the substantially increased interaction length leads to a significant improvement in the achievable sensitivity. Assuming identical experimental parameters, the projected sensitivity on the axion-photon coupling reaches $g_{\text{a}\gamma\gamma} \simeq 9.2 \times 10^{-15} \text{ GeV}^{-1}$, with conversion maintained up to an axion mass cutoff of $m_a \simeq 85 \mu\text{eV}$, determined by the total magnetic length.

Presently, a prototype table-top version is already under construction, with a 2 W 1550 nm laser and a static dipole magnetic field of 1 m in length and about 1.2 T

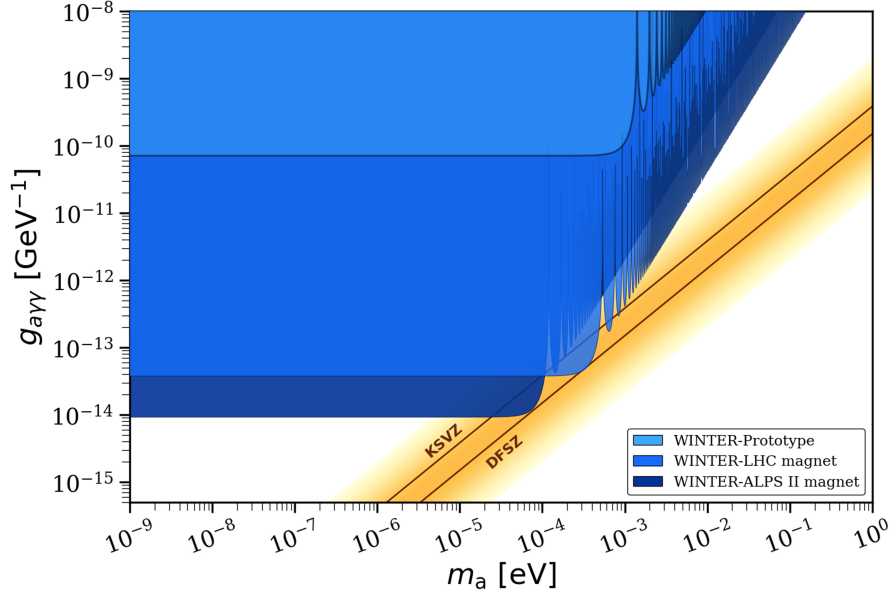


FIG. 2. Projected sensitivity for WINTER experiment using an LHC and ALPS-II magnet configurations, as well as a prototype currently under construction at the University of Hamburg. The specific parameters used for the calculation of all three setups are shown in Table I. Note that the differences in the cutoff axion masses are due to the different wavelengths and magnet length along which the FPC is integrated.

in strength. The custom-made magnet panel is constructed from Neodymium permanent magnets separated by 4 mm defined by the required beam width, which is limited by diffraction losses (see Appendix C). The Neodymium magnets are aligned to provide a close to uniform dipole magnetic field distribution along the entire length. This prototype setup will use a FPC of 1 m in length and a finesse of $\mathcal{F} = 3 \times 10^4$. The corresponding FSR and line width are 150 MHz and 5 kHz, respectively. For a total measuring time of 30 d and a detector with a NEP of $2 \text{ fW}/\sqrt{\text{Hz}}$ [18], the expected exclusion limit of

$g_{\gamma\gamma} \gtrsim 7.1 \times 10^{-11} \text{ GeV}^{-1}$ for ALP masses up to 996 μeV would already improve on existing limits for the axion photon coupling.

VI. CONCLUSION

WINTER is a proposal for a novel experimental setup for broadband searches of axions and ALPs via photon-axion conversion in vacuum. It utilizes a free-space MZI with one arm placed in vacuum and subject to a strong transverse magnetic field, where a FPC enhances the effective

TABLE I. Comparison of the WINTER-prototype and the full WINTER setup with the LHC and ALPS-II magnet configurations.

Parameter	WINTER		
	Prototype	LHC	ALPS-II
B_{ext}	1.2 T	9 T	5.3 T
L_{FPC}	1 m	10 m	200 m
λ_γ	1550 nm	1064 nm	1064 nm
$\mathcal{F}$	3×10^4	10^5	1.2×10^4
t	30 days	1 year	1 year
P_{tot}	2 W	130 W	130 W
$P_{\text{dark}}/P_{\text{tot}}$	1%	0.01%	0.01%
NEP_{PD2}	$2 \text{ fW}/\sqrt{\text{Hz}}$	$0.7 \text{ fW}/\sqrt{\text{Hz}}$	$0.7 \text{ fW}/\sqrt{\text{Hz}}$
$\text{NEP}_{\text{SN+PD2}}$	$75.6 \text{ pW}/\sqrt{\text{Hz}}$	$73.4 \text{ pW}/\sqrt{\text{Hz}}$	$73.4 \text{ pW}/\sqrt{\text{Hz}}$
$g_{\gamma\gamma}^{\text{a}}$	$7.1 \times 10^{-11} \text{ GeV}^{-1}$	$3.7 \times 10^{-14} \text{ GeV}^{-1}$	$9.2 \times 10^{-15} \text{ GeV}^{-1}$
m_a^{b}	996 μeV	380 μeV	85 μeV

^aMinimum projected mass-independent coupling.

^bCutoff axion mass.

circulating power and the total photon path length inside the space pervaded by the magnetic field. The interferometric scheme allows for highly sensitive detection of small amplitude modulations arising from photon-to-axion conversion through the Primakoff effect.

A key advantage of the WINTER setup is its ability to operate near the dark fringe, where shot noise is significantly reduced [19]. This optical noise suppression improves the overall sensitivity, making the system capable of detecting small signals in the presence of dominant background noise. Additionally, the modulation of the polarization state allows the signal to be searched at a selected frequency ω_{sig} , distinguishing it from environmental and instrumental effects.

Unlike haloscopes, WINTER does not rely on assumptions about the local dark matter density or its velocity distribution. This approach is model independent and sensitive to axions regardless of their contribution to the total dark matter abundance, making it a robust probe of photon-axion coupling.

Compared to existing experiments, WINTER offers a unique approach by relying on single-stage photon-to-axion conversion, unlike LSW setups such as Any Light Particle Search II (ALPS II) [20], which require both conversion and reconversion, thereby reducing the signal strength. Additionally, birefringence-based experiments (e.g., Polarisation of Vacuum with Laser (PVLAS) [21]) aim to detect axion-induced changes in polarization, however, without using interferometry techniques. Lastly, the CAST experiment [22] is used to search for solar axions and therefore relies on axion-generation in the Sun's interior.

WINTER, being broadband and model independent, is well positioned to explore previously inaccessible regions of axion parameter space, including the theoretically motivated QCD axion band [11]. In this regime, WINTER achieves sensitivities down to $g_{a\gamma\gamma} \simeq 3.7 \times 10^{-14} \text{ GeV}^{-1}$ when operated with an LHC-type dipole magnet ($B_{\text{ext}} = 9 \text{ T}$, $L = 10 \text{ m}$), while employing the ALPS-II dipole magnet configuration ($B_{\text{ext}} = 5.3 \text{ T}$, $L = 200 \text{ m}$) improves the sensitivity to $g_{a\gamma\gamma} \simeq 9.2 \times 10^{-15} \text{ GeV}^{-1}$ at the expense of a reduced conversion mass range. We emphasize that our scheme searches for relativistic axions converted at the laser frequency and therefore does not directly probe the nonrelativistic halo dark-matter population.

Finally, a prototype table-top version is already under construction at the University of Hamburg. Despite its compact scale, it is expected to demonstrate the feasibility of reaching competitive sensitivity across a broad axion mass range, offering a new and competitive direction in laboratory-based axion searches.

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DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: INTERFEROMETRIC PRINCIPLE

1. Input fields

The input field at port 1 of WINTER's MZI is represented as a superposition of two linear polarization components differentiated by $\chi = \pi/2$, and modulated at a frequency ω_m . Mathematically, this modulated electric field can be expressed as a Jones vector,

$$\mathbf{E}_{\text{in1}}(t) = E_0 \begin{pmatrix} 1 + \beta_m \sin(\omega_m t + \chi) \\ 1 + \beta_m \sin(\omega_m t) \end{pmatrix} e^{i\omega_\gamma t}, \quad (\text{A1})$$

while port 2 has no input field:

$$\mathbf{E}_{\text{in2}}(t) = \begin{pmatrix} 0 \\ 0 \end{pmatrix}. \quad (\text{A2})$$

The input 1 field expression from Eq. (A1) can be reexpressed in terms of the carrier $[\mathbf{E}_{\text{car}}(t)]$ and upper and lower sidebands $[\mathbf{E}_{\text{sid}}(t)]$ in Eq. (A5) as

$$\mathbf{E}_{\text{in1}}(t) = \mathbf{E}_{\text{car}}(t) + \mathbf{E}_{\text{sid}}(t), \quad (\text{A3})$$

where

$$\mathbf{E}_{\text{car}}(t) = E_0 e^{i\omega_\gamma t} \begin{pmatrix} 1 \\ 1 \end{pmatrix}, \quad (\text{A4})$$

$$\mathbf{E}_{\text{sid}}(t) = \frac{E_0 \beta_m}{2i} \begin{pmatrix} e^{i[(\omega_\gamma + \omega_m)t + \chi]} - e^{i[(\omega_\gamma - \omega_m)t + \chi]} \\ e^{i(\omega_\gamma + \omega_m)t} - e^{i(\omega_\gamma - \omega_m)t} \end{pmatrix}, \quad (\text{A5})$$

Following the interferometric scheme from Fig. 1, the input fields then go through a 50/50 BS, which is represented by the unitary matrix

$$\mathbf{U}_{\text{BS}} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & i \\ i & 1 \end{pmatrix}. \quad (\text{A6})$$

Applying that to the input fields $\mathbf{E}_{\text{in}1}(t)$ (port 1) and $\mathbf{E}_{\text{in}2}(t)$ (port 2), we obtain the sensing field $\mathbf{E}_S(t)$ and reference field $\mathbf{E}_R(t)$:

$$\begin{pmatrix} \mathbf{E}_S \\ \mathbf{E}_R \end{pmatrix} = \mathbf{U}_{\text{BS1}} \begin{pmatrix} \mathbf{E}_{\text{in}1}(t) \\ \mathbf{E}_{\text{in}2}(t) \end{pmatrix}. \quad (\text{A7})$$

Note that only the polarization component of the beam, which is parallel to the external magnetic field present in the sensing arm, couples to the axion. Therefore, in the following calculations, we will only focus on this case.

2. Carrier amplitude contribution

Considering the optical fields at the sensing and reference arms after the first 50/50 BS from Eq. (A7), the optical carrier field contributions are estimated, accounting for the photon-axion mixing effect ($P_{\gamma \rightarrow a}$) and the length contributions from each arm, which are given by

$$E_{\text{car},S} = \frac{E_0}{\sqrt{2}} (1 - P_{\gamma \rightarrow a}) e^{i[\omega_\gamma t - k_\gamma (L + \frac{\Delta L}{2})]}, \quad (\text{A8})$$

$$E_{\text{car},R} = \frac{iE_0}{\sqrt{2}} e^{i[\omega_\gamma t - k_\gamma (L - \frac{\Delta L}{2})]}, \quad (\text{A9})$$

where L_S and L_R represent the optical path lengths of the sensing and reference arms, respectively; L is the macroscopical optical path length, $L = \frac{L_R + L_S}{2}$; and $\Delta L = L_R - L_S$ is the path length difference.

The carrier field components from each arm after recombination from the second 50/50 BS are

$$E_{\text{car},\text{out}1} = \frac{E_0}{2} e^{i\omega_\gamma t} \{ e^{-i[k_\gamma (L + \frac{\Delta L}{2})]} (1 - P_{\gamma \rightarrow a}) - e^{-ik_\gamma (L - \frac{\Delta L}{2})} \}, \quad (\text{A10})$$

$$E_{\text{car},\text{out}2} = i \frac{E_0}{2} e^{i\omega_\gamma t} \{ e^{-i[k_\gamma (L + \frac{\Delta L}{2})]} (1 - P_{\gamma \rightarrow a}) + e^{-ik_\gamma (L - \frac{\Delta L}{2})} \}. \quad (\text{A11})$$

The general condition for destructive interference (dark fringe) for the carrier field in a two-path system based on a MZI is satisfied when the phase difference between the two arms is an odd multiple of π ,

$$k_\gamma \Delta L = (2n + 1)\pi, \quad (\text{A12})$$

where n is an integer representing the fringe order. For infrared lasers (e.g., $\lambda_\gamma = 1064$ nm), the path length differences involved are macroscopic with n being a large number ($n \sim 10^5$). In practice, achieving a perfect destructive interference can become challenging due to adjustment imperfections, locking errors, or noise. The actual path difference can be written as

$$\Delta L = (2n + 1) \frac{\lambda_\gamma}{2} + \sigma_\gamma, \quad (\text{A13})$$

where σ_γ is the experimental uncertainty in locking onto the dark fringe. Substituting into the interference condition from Eq. (A12), we get

$$k_\gamma \Delta L = (2n + 1)\pi + \alpha_\gamma, \quad (\text{A14})$$

where $\alpha_\gamma = k_\gamma \sigma_\gamma$ represents the residual phase error due to imperfect fringe locking.

By applying condition (A14), the resulting output carrier fields become

$$E_{\text{car},\text{out}1} = \frac{E_0}{2} e^{i(\omega_\gamma t - k_\gamma L)} e^{-i[(2n+1)\pi + \alpha_\gamma]} [(1 - P_{\gamma \rightarrow a}) + e^{i\alpha_\gamma}], \quad (\text{A15})$$

$$E_{\text{car},\text{out}2} = i \frac{E_0}{2} e^{i(\omega_\gamma t - k_\gamma L)} e^{-i[(2n+1)\pi + \alpha_\gamma]} [(1 - P_{\gamma \rightarrow a}) - e^{i\alpha_\gamma}], \quad (\text{A16})$$

where output port 1 becomes the bright port and port 2 becomes the dark port.

3. Sideband amplitude contribution

The upper (+) and lower (−) sideband contribution for the sensing and reference arms is given by

$$E_{\text{sid},(\pm),S} = \pm \frac{\beta_m E_0}{2i\sqrt{2}} (1 - P_{\gamma \rightarrow a}) e^{i[(\omega_\gamma \pm \omega_m)t - \phi_S^\pm]}, \quad (\text{A17})$$

$$E_{\text{sid},(\pm),R} = \pm \frac{\beta_m E_0}{2\sqrt{2}} e^{i[(\omega_\gamma \pm \omega_m)t - \phi_R^\pm]}, \quad (\text{A18})$$

where the phase contributions for each arm are

$$\begin{aligned} \phi_S^\pm &= (k_\gamma \pm k_m) \left(L + \frac{\Delta L}{2} \right) \\ \phi_R^\pm &= (k_\gamma \pm k_m) \left(L - \frac{\Delta L}{2} \right). \end{aligned} \quad (\text{A19})$$

The total sideband contribution is described by the sum of the two sidebands:

$$E_{\text{sid}} = E_{\text{sid}(+)} + E_{\text{sid}(-)}. \quad (\text{A20})$$

After recombining the beams through the second 50/50 beam splitter, the field output at port 1 and port 2 is given by

$$\begin{aligned} E_{\text{sid},\text{out}1} &= i \frac{\beta_m E_0}{4} \{ e^{i\psi_+} [-e^{-i\xi_+} (1 - P_{\gamma \rightarrow a}) + e^{i\xi_+}] \\ &\quad - e^{i\psi_-} [-e^{-i\xi_-} (1 - P_{\gamma \rightarrow a}) + e^{i\xi_-}] \} \end{aligned} \quad (\text{A21})$$

$$E_{\text{sid,out2}} = \frac{\beta_m E_0}{4} \left\{ e^{i\psi_+} [e^{-i\xi_+} (1 - P_{\gamma \rightarrow a}) + e^{i\xi_+}] - e^{i\psi_-} [e^{-i\xi_-} (1 - P_{\gamma \rightarrow a}) + e^{i\xi_-}] \right\}, \quad (\text{A22})$$

where

$$\psi_{\pm} = (\omega_{\gamma} \pm \omega_m)t - (k_{\gamma} \pm k_m)L, \quad (\text{A23})$$

$$\xi_{\pm} = (k_{\gamma} \pm k_m) \frac{\Delta L}{2}. \quad (\text{A24})$$

For the sidebands modulation, destructive interference is met when $(k_{\gamma} + k_m)\Delta L/2 = \pi$ for the upper (+) sideband and $(k_{\gamma} - k_m)\Delta L/2 = 0$ for the lower (−) sideband. Analogous to the carrier interference condition from Eq. (A14), we derive

$$k_m \Delta L = (2p + 1)\pi + \alpha_m, \quad (\text{A25})$$

where $p \in \mathbb{Z}(\mathcal{O}(1))$ and $\alpha_m = \frac{k_m \alpha_{\gamma}}{k_{\gamma}}$. Since $k_m \ll k_{\gamma}$, the residual phase error from the sideband locking, α_m , is ~ 6 – 8 orders of magnitude smaller than the one for the carrier α_{γ} , which can be assumed to be dominant, and α_m can therefore be neglected. The resulting phase contributions from Eqs. (A23)–(A24) become

$$\psi_{\pm} = (\omega_{\gamma} \pm \omega_m)t - \frac{L}{\Delta L} \{ [2n + 1 \pm (2p + 1)]\pi + \alpha_{\gamma} \} \quad (\text{A26})$$

$$\xi_{\pm} = \frac{[2n + 1 \pm (2p + 1)]\pi + \alpha_{\gamma}}{2}. \quad (\text{A27})$$

4. Power output estimation

To estimate the power, we expand the complex expressions using Eqs. (A15), (A16), (A21), and (A22) with the phase values from Eqs. (A26), (A27):

$$\begin{aligned} P_{\text{out1}} &= E_{\text{car,out1}}^* \cdot E_{\text{car,out1}} + E_{\text{car,out1}}^* \cdot E_{\text{sid,out1}} + E_{\text{car,out1}} \cdot E_{\text{sid,out1}}^* + E_{\text{sid,out1}}^* \cdot E_{\text{sid,out1}} \\ &= \frac{1}{4} E_0^2 \left\{ 2 + (-2 + P_{\gamma \rightarrow a})P_{\gamma \rightarrow a} - 2(-1 + P_{\gamma \rightarrow a}) \cos \alpha_{\gamma} \right. \\ &\quad + \beta_m^2 [2 + (-2 + P_{\gamma \rightarrow a})P_{\gamma \rightarrow a} + 2(-1 + P_{\gamma \rightarrow a}) \cos \alpha_{\gamma}] \cos^2 \left[\omega_m t - \frac{(2p + 1)\pi}{\Delta L} L \right] \\ &\quad \left. - 2(-1)^p \beta_m (-2 + P_{\gamma \rightarrow a})P_{\gamma \rightarrow a} \cos \left[\omega_m t - \frac{(2p + 1)\pi}{\Delta L} L \right] \right\}, \end{aligned} \quad (\text{A28})$$

$$\begin{aligned} P_{\text{out2}} &= E_{\text{car,out2}}^* \cdot E_{\text{car,out2}} + E_{\text{car,out2}}^* \cdot E_{\text{sid,out2}} + E_{\text{car,out2}} \cdot E_{\text{sid,out2}}^* + E_{\text{sid,out2}}^* \cdot E_{\text{sid,out2}} \\ &= \frac{1}{4} E_0^2 \left\{ 2 + (-2 + P_{\gamma \rightarrow a})P_{\gamma \rightarrow a} + 2(-1 + P_{\gamma \rightarrow a}) \cos \alpha_{\gamma} \right. \\ &\quad + \beta_m^2 [2 + (-2 + P_{\gamma \rightarrow a})P_{\gamma \rightarrow a} - 2(-1 + P_{\gamma \rightarrow a}) \cos \alpha_{\gamma}] \cos^2 \left[\omega_m t - \frac{(2p + 1)\pi}{\Delta L} L \right] \\ &\quad \left. - 2(-1)^p \beta_m (-2 + P_{\gamma \rightarrow a})P_{\gamma \rightarrow a} \cos \left[\omega_m t - \frac{(2p + 1)\pi}{\Delta L} L \right] \right\}. \end{aligned} \quad (\text{A29})$$

For perfect dark fringe conditions ($\alpha_{\gamma} = 0$) in both the carrier and the upper (+) modulated sideband, considering $p = 0$, the output powers from each port can be simplified to

$$\begin{aligned} P_{\text{out1}} &= \frac{1}{4} E_0^2 [(-2 + P_{\gamma \rightarrow a})^2 + \beta_m^2 P_{\gamma \rightarrow a}^2 \cos^2(\omega_m t - k_m L) \\ &\quad - 2\beta_m (-2 + P_{\gamma \rightarrow a})P_{\gamma \rightarrow a} \cos(\omega_m t - k_m L)] \end{aligned} \quad (\text{A30})$$

$$\begin{aligned} P_{\text{out2}} &= \frac{1}{4} E_0^2 [P_{\gamma \rightarrow a}^2 + \beta_m^2 (-2 + P_{\gamma \rightarrow a})^2 \cos^2(\omega_m t - k_m L) \\ &\quad - 2\beta_m (-2 + P_{\gamma \rightarrow a})P_{\gamma \rightarrow a} \cos(\omega_m t - k_m L)]. \end{aligned} \quad (\text{A31})$$

It is noted that, as expected, the cross-product between the carrier and the sidebands does not depend on α_{γ} as it cancels out. However, α_{γ} still has an impact on the carrier and sideband self-interference power contributions.

APPENDIX B: DEMODULATION PROCESS

As an example, to lock to the upper sideband at $\omega_{\gamma} + \omega_m$, the resulting power from port 1 (bright port) needs to be demodulated at a frequency ω_m . Note that the output powers resulting from the self-interference of the carrier and sidebands [first and second terms of Eq. (A28)], respectively, are omitted since their contribution is negligible after demodulating and averaging their power over

long periods. In addition, we can approximate by considering only the first-order terms accounting for $P_{\gamma \rightarrow a} \ll 1$.

$$P_{\text{cross,out1}}(t) = E_0^2 \beta_m P_{\gamma \rightarrow a} \cos(\omega_m t - k_m L). \quad (\text{B1})$$

The demodulation approach uses a lock-in amplification, by mixing the signal from the bright port with a LO at the same frequency ω_m , following a homodyne procedure. Regarding that, the output of the mixer is

$$P_{\text{mix}}(t) = P_{\text{cross,out1}}(t) \cos(\omega_m t + \phi_{\text{LO}}), \quad (\text{B2})$$

where ϕ_{LO} is the phase of the LO. The resulting signal power then passes through a low-pass filter to remove high-frequency components ($2\omega_m$). The time-averaged output power of the mixer assuming $\phi_{\text{LO}} = 0$ is

$$P_{\text{av}} = \frac{1}{2} P_{\text{tot}} \beta_m P_{\gamma \rightarrow a}, \quad (\text{B3})$$

where $P_{\text{tot}} = E_0^2$ is the input power of the laser.

After the interferometer is locked, the axion-induced signal is recovered by demodulating the dark port 2 output at twice the polarization-modulation frequency $2\omega_{\text{sig}}$, driven via the EOM-PC. Note in this case that the photon-to-axion conversion probability, assuming the application of an external magnetic field, is proportional to the square of the interaction term, $(\mathbf{E} \cdot \mathbf{B}_{\text{ext}})^2$. If the electric field is modulated along the direction of the magnetic field through the EOM-PC, the dot product can be written as

$$\mathbf{E} \cdot \mathbf{B}_{\text{ext}} = |\mathbf{E}| |\mathbf{B}_{\text{ext}}| \beta_{\text{sig}} \cos(\omega_{\text{sig}} t), \quad (\text{B4})$$

where β_{sig} is the modulation depth. Squaring Eq. (B4) leads to a time-dependent $P_{\gamma \rightarrow a}(t)$:

$$P_{\gamma \rightarrow a}(t) = (\mathbf{E} \cdot \mathbf{B}_{\text{ext}})^2 \propto B_{\text{ext}}^2 \beta_{\text{sig}}^2 \cos^2(\omega_{\text{sig}} t). \quad (\text{B5})$$

Averaging over time gives

$$P_{\gamma \rightarrow a} \propto B_{\text{ext}}^2 \frac{1}{2} \beta_{\text{sig}}^2, \quad (\text{B6})$$

so the time-averaged conversion probability acquires an additional factor $P_{\gamma \rightarrow a} \propto \frac{1}{2} \beta_{\text{sig}}^2$. Therefore, the resulting output power from Eq. (B3), while using Eqs. (B6) and (6), becomes

$$P_{\text{av}} = \frac{1}{16} P_{\text{tot}} \beta_m \beta_{\text{sig}}^2 g_{a\gamma\gamma}^2 B_{\text{ext}}^2 z^2. \quad (\text{B7})$$

APPENDIX C: OPTICAL CAVITY LOSSES

To effectively achieve a high finesse of $\mathcal{F} \approx 10^5$, it is critical to have optimal mirror dimensions and beam aperture to keep the clipping and diffraction losses at a minimum. When accounting for multiple loss mechanisms in the FPC (clipping, diffraction, and absorption), the total loss δ_{total} is the sum,

$$\delta_{\text{total}} = \delta_{\text{clip}} + \delta_{\text{diff}} + \delta_{\text{trans},1} + \delta_{\text{trans},2}, \quad (\text{C1})$$

where the round-trip losses caused by the mirror clipping and diffraction are given by δ_{clip} and δ_{diff} , respectively, and $\delta_{\text{trans},1}, \delta_{\text{trans},2}$ are the transmission losses of the two mirrors. The expression for the finesse shown in Eq. (13) can be simplified when the aforementioned losses are small ($\delta_{\text{total}} \ll 1$) [23,24]:

$$\mathcal{F} = \frac{2\pi}{(\delta_{\text{clip}} + \delta_{\text{diff}} + \delta_{\text{trans},1} + \delta_{\text{trans},2})}. \quad (\text{C2})$$

1. Analytical approximation

Assuming Gaussian beam propagation, the beam expands from its waist at the plane mirror (w_{plane}) to a larger radius at the concave mirror [$w_{\text{conc}}(L_{\text{FPC}})$], calculated from the Rayleigh range as follows:

$$z_R = \frac{\pi w_{\text{plane}}^2}{\lambda_\gamma} \quad (\text{C3})$$

$$w_{\text{conc}}(L_{\text{FPC}}) = w_{\text{plane}} \sqrt{1 + \left(\frac{L_{\text{FPC}}}{z_R}\right)^2}. \quad (\text{C4})$$

Clipping losses arise from the mode mismatch between the beam waist and the effective radius (aperture) of the concave mirror ($a = 17.5$ cm), expressed as

$$\delta_{\text{clip,conc}} = \exp\left(-\frac{2a^2}{w_{\text{conc}}^2}\right). \quad (\text{C5})$$

Since both mirrors truncate the beam with their respective beam waists, the round-trip combination is

$$\delta_{\text{clip}} = 1 - (1 - \delta_{\text{clip,plane}})(1 - \delta_{\text{clip,conc}}), \quad (\text{C6})$$

which for small losses ($\delta_{\text{clip}} \ll 1$) reduces to the additive approximation:

$$\delta_{\text{clip}} \approx \delta_{\text{clip,conc}} + \delta_{\text{clip,plane}}. \quad (\text{C7})$$

2. Numerical estimation

In a stable optical cavity, the fundamental Gaussian mode is well confined by the mirrors, but finite apertures and mirror shapes cause a fraction of the light to diffract outside the ideal mode. This effect is not captured by the clipping loss formulation, since even a beam entirely inside the mirror aperture will gradually lose power due to diffraction at each round trip. The Fox-Li method provides an accurate numerical estimation by finding the optimal eigenmode of the cavity via numerical propagation of the field between mirrors with Fresnel diffraction and application of the mirror aperture and curvature at each step [25].

The Fox-Li method treats an optical cavity as a round-trip propagation operator $\mathcal{M}$ that maps a transverse field $E(x, y)$ at one mirror surface into the field after one round trip, including free-space Fresnel propagation, phase shifts due to mirror curvature, and amplitude masks representing apertures or mirror edges. By applying $\mathcal{M}$ repeatedly to an arbitrary initial field $E_0(x, y)$, the iteration gradually filters out unstable components, and the field converges to the cavity's dominant eigenmode. Mathematically, this corresponds to

$$E_{n+1} = \mathcal{M}E_n \xrightarrow{n \rightarrow \infty} E_{n+1} = \eta E_n, \quad (\text{C8})$$

where η is a complex eigenvalue whose magnitude determines the round-trip diffraction loss by

$$\delta_{\text{diff}} = 1 - |\eta|^2. \quad (\text{C9})$$

It is noted that the clipping losses average to zero as they are tiny compared to the aperture of the cavity mirrors. For small w_{plane} (typically $< a/3$), the clipping losses are increased due to the high beam divergence, while for $w_{\text{plane}} > a/3$, the clipping losses increase again due to overfilling of the active area of the mirror. The total losses δ_{tot} for the FPC of the WINTER setup are shown in Fig. 3 as a function of the beam waist of the plane mirror (w_{plane}).

The impact of the total losses on the finesse of the cavity is shown in Fig. 4 as a function of the beam waist at the plane

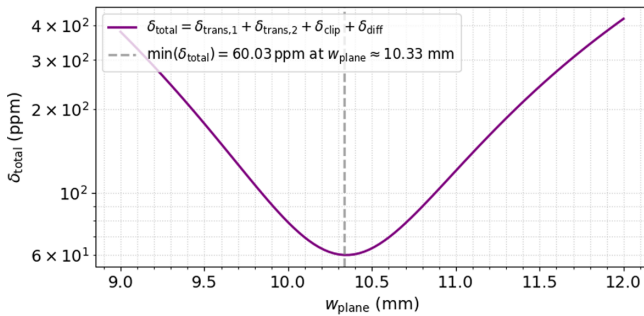


FIG. 3. Representation of the total losses in ppm (in blue) for WINTER experiment as a function of the beam waist at the plane mirror (M2). The selected parameters are shown in Table II.

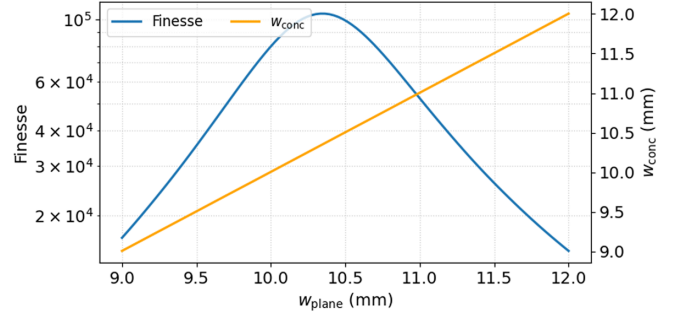


FIG. 4. Representation of the finesse variation (in blue) for WINTER experiment as a function of the beam waist at the plane mirror (M2), accounting for diffraction and clipping losses. At the secondary y axis, the beam waist at the concave mirror (M3) is also shown (in orange). The selected FPC details are shown in Table II.

mirror. For WINTER, the optimal beam waist at the plane mirror is about 10.33 mm with a concave mirror's beam waist of equal value. It is noted that a FPC with similar properties has already been implemented in vacuum conditions and locked from other experiments (see, e.g., ALPS II [20]), demonstrating the feasibility of this approach.

Our calculations primarily focus on the configuration using an LHC dipole magnet. For comparison, we also estimate the parameters for a setup employing the ALPS-II string of dipole magnets. The relevant FPC parameters and corresponding losses for WINTER, with both LHC and ALPS-II dipole magnets, as well as for the WINTER-prototype apparatus, are summarized in Table II.

3. Thermal considerations

WINTER's experimental setup employs a high-finesse optical cavity with mirror coatings and substrates engineered to minimize absorption and scattering, enabling stable long-term operation at high intracavity power. Under the conditions characterized by the parameters listed in Table II, the system achieves intracavity power densities of about 0.66 MW/cm^2 , which lies safely below the upper operational limit of 1 MW/cm^2 demonstrated in the LIGO gravitational-wave detectors under comparable continuous-wave conditions [26]. The mirrors use state-of-the-art dielectric coatings composed of alternating layers of silica (SiO_2 , low refractive index) and titania-doped tantala ($\text{Ti:Ta}_2\text{O}_5$, high refractive index) deposited by ion-beam sputtering, ensuring exceptionally low optical losses. Ultrapure fused silica substrates with optimized radii of curvature support stable Gaussian modes, while the entire system is housed in high vacuum to suppress convective heat transfer and contamination. Residual heat from absorption in the coatings and substrates is conducted away through the mirror mounts, which provide the primary thermal contact in vacuum and transfer heat to the chamber walls, acting as a heat sink. This integrated approach allows WINTER to maintain high intracavity power without thermal distortion or coating degradation.

TABLE II. Comparison of the Fabry-Pérot cavity parameters for the WINTER-prototype and the full WINTER setup with the LHC and ALPS-II magnet configurations.

Parameter	Notation	WINTER		
		Prototype	LHC	ALPS-II
Laser power	P_{tot}	2 W	130 W	130 W
Cavity length	L_{FPC}	1 m	10 m	200 m
Laser wavelength	λ_{γ}	1550 nm	1064 nm	1064 nm
Beam waist at plane mirror	w_{plane}	0.92 mm	10.33 mm	8.28 mm
Beam waist at concave mirror	w_{conc}	1.06 mm	10.34 mm	11.64 mm
Rayleigh range	z_R	1.72 m	315 m	203 m
Radius of curvature	R_{conc}	2 m	5 km	200 m
Aperture radius	a	12.5 mm	17.5 cm	17.5 cm
Transmission losses	$\delta_{\text{trans},1} = \delta_{\text{trans},2}$	100 ppm	30 ppm	200 ppm
Diffraction losses	δ_{diff}	32 ppm	0.03 ppm	122.39 ppm
Clipping losses	δ_{clip}	0 ppm	0 ppm	0 ppm
Total round-trip loss	δ_{total}	232 ppm	60.03 ppm	522.39 ppm
Finesse	$\mathcal{F} = 2\pi/\delta_{\text{total}}$	3×10^4	1×10^5	1.2×10^4

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Correction: Minor errors in Eqs. (A17), (A18), and (B7) and in the text following (B3) have been fixed.