

Cross-correlation scheme for quantum optical coherence tomography based on Michelson interferometer

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Quantum optical coherence tomography (QOCT) offers a simple way to cancel dispersion broadening in a sample while also providing twice the resolution compared to classical OCT. However, to achieve these advantages, a bright and broadband source of entangled photon pairs is required. A simple implementation uses collinear spontaneous parametric down-conversion in a Michelson interferometer (MI), yet this autocorrelation scheme suffers from parasitic terms and sensitivity to phase noise. Here, we introduce a cross-correlation MI-based QOCT that overcomes these drawbacks, significantly advancing QOCT toward practical applications.

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I. INTRODUCTION

The resolution of far-field imaging systems is limited by the optical radiation wavelength and the numerical aperture (NA). While the transverse resolution is inversely proportional to NA, the axial resolution is inversely proportional to the square of NA [1]. This limitation is particularly critical in applications where NA is significantly limited, such as retinal imaging through an eye pupil [2].

The technique that addresses this limitation is optical coherence tomography (OCT), introduced in Ref. [3] and widely described in Ref. [2]. In the time-domain OCT scheme, presented in Fig. 1(a), the sample is placed in one arm of a Michelson interferometer (MI), while the other arm's delay τ is varied. Each reflective layer of the sample generates an interferogram $M_{\text{OCT}}(\tau)$, allowing one to reconstruct the structure of the sample—coordinates and reflectivities of the layers. The interferogram obtained from a single layer can be considered as a point-spread function (PSF) of the imaging system. Its width determines the axial resolution, independent of NA but limited by the coherence length of the optical source. Thus, a low-coherence broadband source is required for fine axial resolution. However, any dispersion in the sample or setup components broadens the interferogram, severely

limiting the scanning depth of OCT unless specific sample dispersion properties are known and compensated.

Several dispersion-cancellation strategies for OCT have been demonstrated. Some—like the hardware phase-compensation method [4]—require precise *a priori* knowledge of the sample's dispersion and only cancel dispersion at a single imaging depth. Others rely on elaborate nonlinear optics implementations, for example, phase conjugation or chirped-pulse schemes [5–8], which demand high-power lasers and complex setups. Other dispersion-cancellation approaches rely on intensive numerical spectral postprocessing [9–14], which makes them intrinsically sensitive to interferometric phase and timing instabilities.

Quantum optical coherence tomography (QOCT) [15,16], based on Hong-Ou-Mandel (HOM) interference [17], presented in Fig. 1(b), overcomes the dispersion broadening issue and also provides double resolution. In QOCT, each reflective layer produces a dip $M_{\text{HOM}}(\tau)$, where the width, inversely proportional to the bandwidth of the spontaneous parametric down-conversion (SPDC) [18] field, defines the system's axial resolution. Notably, QOCT is insensitive to group-velocity dispersion if the SPDC pump field is sufficiently narrowband [19], and it provides twice the axial resolution of OCT with the same source spectrum. Additionally, as mentioned in Ref. [17], HOM interference does not require any phase stabilization or subwavelength scanning steps, and the interferogram pattern $M_{\text{HOM}}(\tau)$ is much clearer and easier to process compared to $M_{\text{OCT}}(\tau)$.

Despite its advantages, the practical application of QOCT remains limited due to the long measurement times required by the low brightness of broadband SPDC sources. Bright sources are more effortlessly produced via

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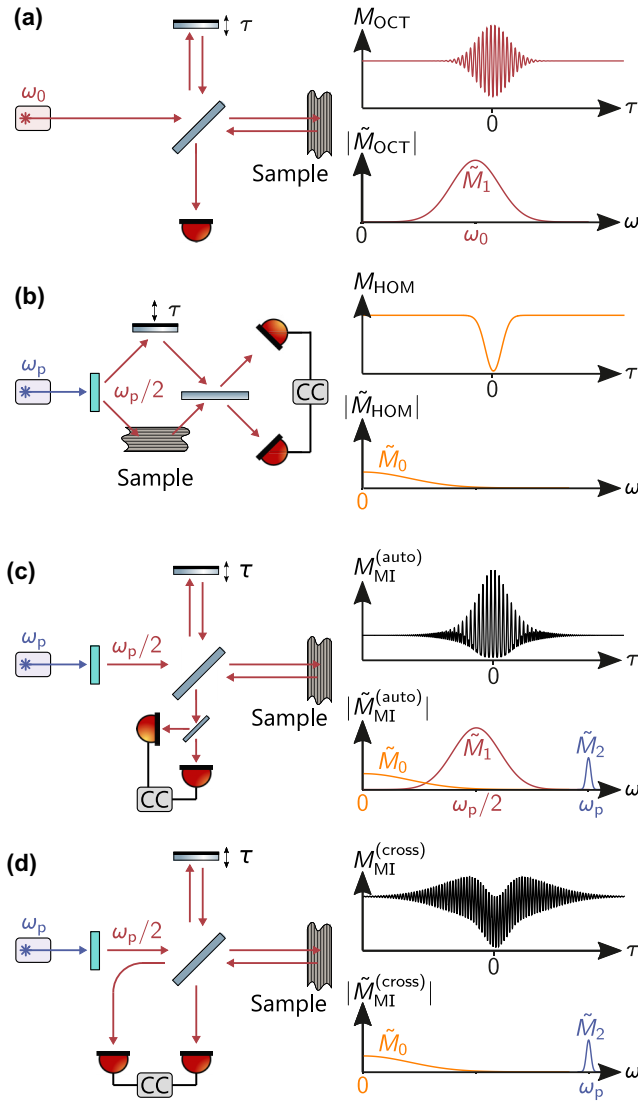


FIG. 1. Optical coherence tomography (OCT) and quantum OCT (QOCT) techniques: (a) time-domain OCT, (b) traditional Hong-Ou-Mandel (HOM)-based QOCT, (c) Michelson interferometer (MI)-based QOCT with autocorrelation measurements, and (d) the studied scheme of MI-based QOCT with cross-correlation measurements. CC, coincidence circuit. The left column presents sketches of the experimental schemes, while the right column shows the corresponding point-spread function (PSF) interferograms $M(\tau)$ (from a single reflective layer) and their Fourier transforms $|\tilde{M}(\omega)|$.

collinear phase matching, which allows for effective use of longer crystal lengths. As mentioned in Ref. [20], for broadband sources, degenerate phase-matching types I or 0 are preferable, though this leads to both photons being in the same optical mode, preventing their deterministic separation for HOM interference.

Recent studies have suggested that QOCT can also be implemented using a brighter, broader [21], and more

adjustable collinear type-0 or type-I SPDC source coupled to an MI with an autocorrelation measurement scheme [22–26]. In this configuration [Fig. 1(c)], the interferometer’s output signal is split and the coincidence count rate interferogram $M_{\text{MI}}^{(\text{auto})}(\tau)$ is measured.

The coincidence interferogram $M_{\text{MI}}^{(\text{auto})}(\tau)$ has a complex structure. In addition to the desired HOM term M_0 , it contains several parasitic contributions. Its Fourier transform $\tilde{M}_{\text{MI}}^{(\text{auto})}(\omega)$ reveals that these terms occupy distinct central frequencies, allowing M_0 to be extracted by postprocessing via Fourier filtering.

However, achieving submicrometer axial resolution [26] requires an ultrabroadband SPDC source, which causes overlap between the target and parasitic components in the Fourier domain. Phase instabilities further broaden these features, making postprocessing more challenging, reducing the usable spectral bandwidth, and ultimately limiting the axial resolution.

To overcome these limitations, we explore and experimentally demonstrate a *cross-correlation* MI-based QOCT scheme [Fig. 1(d)]. This scheme combines the dispersion cancellation, doubled resolution, and robustness to phase instability of the standard HOM approach with the high brightness of collinear SPDC, while eliminating overlap between parasitic Fourier components of the interferogram. By monitoring coincidences between the MI’s two output ports (one of which is recycled to the input via a circulator), we obtain the PSF interferogram $M_{\text{MI}}^{(\text{cross})}(\tau)$. Its Fourier transform $\tilde{M}_{\text{MI}}^{(\text{cross})}(\omega)$ contains only a narrow pump-frequency peak $\tilde{M}_2$ and the target HOM term $\tilde{M}_0$. Because the parasitic peak is well separated from the HOM dip, it can be removed by simple Fourier filtering or by averaging over high-frequency oscillations during a coarse mechanical scan without limiting SPDC bandwidth or requiring subwavelength path control.

The cross-correlation approach was recently demonstrated for high-precision refractive-index measurements in optical fibers [27]; however, in that application, the PSF width is not critical because each scan measures a single path length and the peak can be located with sub-PSF precision. By contrast, here we provide a comprehensive study of a Michelson-based cross-correlation QOCT: we compare its resolution and signal-to-noise ratio (SNR) with both autocorrelation QOCT and classical OCT, analyze how the axial resolution depends on the source spectrum, demonstrate enhanced robustness to phase noise and optical dispersion, and finally apply the method to defectoscopy—a task where PSF narrowing and dispersion cancellation are essential.

The paper is organized as follows. Section II presents the theoretical descriptions of classical OCT and both autocorrelation and cross-correlation MI-based QOCT schemes and highlights the advantages of the latter. Section III describes the experimental setup that implements both

QOCT schemes alongside a classical OCT reference. Section IV compares resolution and SNR across the three techniques and includes a defectoscopy demonstration.

II. THEORY

In this section we follow the approach of Ref. [26] to describe quantum interference in the MI-based QOCT schemes of Figs. 1(a), 1(c), and 1(d) and to compare them theoretically.

Consider interference of a photon pair in the MI. The interferometer scheme is illustrated in Fig. 2. It consists of a 50:50 beam splitter, a reflective sample, and a translated mirror in a reference channel. It connects four input and four output optical modes. Photon creation operators in these modes are denoted as $\hat{a}_{\text{in}}^\dagger$, $\hat{a}_{\text{out}}^\dagger$, $\hat{b}_{\text{in}}^\dagger$, $\hat{b}_{\text{out}}^\dagger$, $\hat{c}_{\text{in}}^\dagger$, $\hat{c}_{\text{out}}^\dagger$, $\hat{d}_{\text{in}}^\dagger$, $\hat{d}_{\text{out}}^\dagger$. The beam splitter connects these operators as follows:

$$\begin{pmatrix} \hat{a}_{\text{in}}^\dagger \\ \hat{b}_{\text{in}}^\dagger \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & i \\ i & 1 \end{pmatrix} \begin{pmatrix} \hat{c}_{\text{in}}^\dagger \\ \hat{d}_{\text{in}}^\dagger \end{pmatrix}, \quad (1)$$

$$\begin{pmatrix} \hat{c}_{\text{out}}^\dagger \\ \hat{d}_{\text{out}}^\dagger \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & -i \\ -i & 1 \end{pmatrix} \begin{pmatrix} \hat{a}_{\text{out}}^\dagger \\ \hat{b}_{\text{out}}^\dagger \end{pmatrix}.$$

Let $H(\omega)$ denote the sample response function and τ the delay in the reference channel. Then input and output modes are connected as follows:

$$\begin{pmatrix} \hat{c}_{\text{in}}^\dagger \\ \hat{d}_{\text{in}}^\dagger \end{pmatrix} = \begin{pmatrix} e^{i\omega\tau} & 0 \\ 0 & H(\omega) \end{pmatrix} \begin{pmatrix} \hat{c}_{\text{out}}^\dagger \\ \hat{d}_{\text{out}}^\dagger \end{pmatrix}. \quad (2)$$

The general description of the single-photon and photon-pair MI interference is presented in Appendix A. Here we make two assumptions.

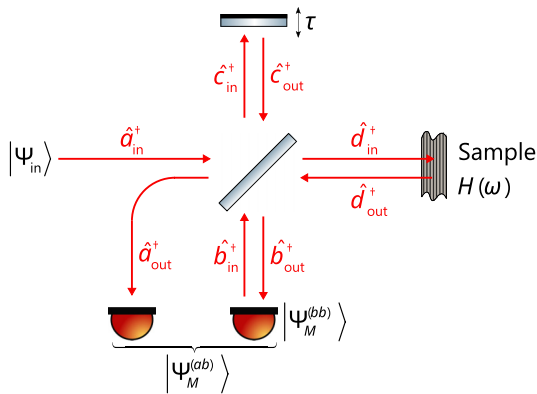


FIG. 2. Michelson interferometer input and output modes, input state $|\Psi_{\text{in}}\rangle$, and measurement projections for autocorrelation $|\Psi_M^{(bb)}\rangle$ and cross-correlation $|\Psi_M^{(ab)}\rangle$.

(1) The sample has a single layer without dispersion, with the response function

$$H(\omega) = r e^{i\omega T}, \quad r^2 \equiv R, \quad (3)$$

where R is the layer reflectivity. Here, $\omega T = kd = n\omega d/c$, with k is a wavevector, d the layer thickness, n the refractive index, and c the speed of light. Thus, $T = nd/c$.

(2) The input state consists of photons with the same polarization, propagating in input mode a and having a Gaussian spectral distribution.

First, we consider a classical or single-photon state with a Gaussian spectral distribution

$$G(\omega|\omega_0, \sigma) \equiv \frac{1}{\sqrt{2\pi}\sigma} \exp\left[-\frac{(\omega - \omega_0)^2}{2\sigma^2}\right], \quad (4)$$

with the central frequency ω_0 and the bandwidth σ (in this section we define the bandwidth as a standard deviation of the Gaussian distribution). Measuring the output intensity (count rate) in mode b as a function of τ yields the following interferogram:

$$M^{(b)}(\tau) = \frac{1+R}{4} - \frac{r}{2} \exp\left[-\frac{(T-\tau)^2}{2\sigma^{-2}}\right] \cos[\omega_0(T-\tau)], \quad (5)$$

presented in Fig. 1(a) (see Appendix A for details). This interferogram represents the PSF of classical OCT.

For QOCT we consider a collinear degenerate SPDC state with type-0 or type-I phase matching,

$$|\Psi_{\text{in}}\rangle = \frac{1}{\sqrt{2}} \iint d\omega_s d\omega_i F(\omega_s, \omega_i) \hat{a}_{\text{in}}^\dagger(\omega_s) \hat{a}_{\text{in}}^\dagger(\omega_i) |\text{vac}\rangle, \quad (6)$$

where $F(\omega_s, \omega_i)$ is a spectral amplitude and

$$|F(\omega_s, \omega_i)|^2 = 2 \underbrace{G(\omega_s - \omega_i|0, \Delta)}_{\text{phase matching}} \times \underbrace{G(\omega_s + \omega_i|\omega_p, \delta)}_{\text{pump}} \quad (7)$$

is a Gaussian approximation of the photon-pair spectral distribution [28,29], visualized in Fig. 3. Here Δ represents the phase-matching bandwidth and $\delta \ll \Delta$ is the pump bandwidth.

The corresponding marginal single-photon spectral distribution also has a Gaussian form,

$$\int d\omega_i |F(\omega_s, \omega_i)|^2 = G(\omega_s|\omega_p/2, \Delta_+/2), \quad (8)$$

with the bandwidth twice as narrow as $\Delta_+ \equiv \sqrt{\Delta^2 + \delta^2} \approx \Delta$.

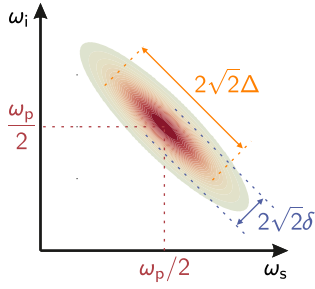


FIG. 3. Two-photon spectral distribution $|f(\omega_s, \omega_i)|$ for the degenerate case.

The traditional MI-based autocorrelation QOCT scheme uses a detector that realizes a projection onto the two-photon state in mode b ,

$$\left| \Psi_M^{(bb)}(\omega_1, \omega_2) \right\rangle = \frac{1}{\sqrt{2}} \hat{b}_{\text{out}}^\dagger(\omega_1) \hat{b}_{\text{out}}^\dagger(\omega_2) |\text{vac}\rangle. \quad (9a)$$

As an alternative, we consider a cross-correlation scheme based on a detector that projects onto the state with two photons in different modes a and b ,

$$\left| \Psi_M^{(ab)}(\omega_1, \omega_2) \right\rangle = \hat{a}_{\text{out}}^\dagger(\omega_1) \hat{b}_{\text{out}}^\dagger(\omega_2) |\text{vac}\rangle. \quad (9b)$$

The probability of registering one photon in mode m and another one in mode n as a function of the delay τ in the reference channel (the two-photon interferogram) is given by

$$M^{(mn)}(\tau) = \iint d\omega_1 d\omega_2 \left| \left\langle \Psi_M^{(mn)}(\omega_1, \omega_2) \mid \Psi_{\text{in}} \right\rangle \right|^2. \quad (10)$$

Substituting Eqs. (1)–(9) into Eq. (10) we obtain the final interferograms for autocorrelation (bb) and cross-correlation (ab) measurement schemes as follows:

$$M_{\text{MI}}^{(\text{auto})} \equiv M^{(bb)} = \frac{M_c + M_0 - M_1 + M_2}{4}, \quad (11a)$$

$$M_{\text{MI}}^{(\text{cross})} \equiv M^{(ab)} = \frac{M_c - M_0 - M_2}{2}. \quad (11b)$$

Here,

$$M_c = \frac{1}{4} (1 + R)^2 \quad (12a)$$

is a constant term. The term

$$M_0(\tau) = \frac{1}{2} R \exp \left[-\frac{(T - \tau)^2}{2\Delta^{-2}} \right] \quad (12b)$$

is related to HOM interference, presented in Fig. 1(b).

The next term

$$M_1(\tau) = r(1 + R) \exp \left[-\frac{(T - \tau)^2}{8\Delta_+^{-2}} \right] \cos \left[\frac{\omega_p (T - \tau)}{2} \right] \quad (12c)$$

is a term related to single-photon Michelson interference. It fits up to a constant factor with the second term of the single-photon interferogram, Eq. (5), for the single-photon Gaussian spectral distribution, Eq. (8). The last term

$$M_2(\tau) = \frac{1}{2} R \exp \left[-\frac{(T - \tau)^2}{2\delta^{-2}} \right] \cos [\omega_p (T - \tau)] \quad (12d)$$

is a similar term, related to the Michelson interference of unseparated photon pairs, equivalent to the interference of narrowband pump photons.

The absolute values of Fourier transforms of the time-dependent terms, presented in Figs. 1(c) and 1(d), are

$$|\tilde{M}_0(\omega)| = \frac{\sqrt{\pi} R}{2} G(\omega | 0, \Delta), \quad (13a)$$

$$|\tilde{M}_1(\omega)| = \frac{\sqrt{\pi} r(1 + R)}{2} G\left(\omega \left| \frac{\omega_p}{2}, \frac{\Delta_+}{2} \right.\right), \quad (13b)$$

$$|\tilde{M}_2(\omega)| = \frac{\sqrt{\pi} R}{4} G(\omega, \omega_p, \delta). \quad (13c)$$

The target M_0 term can be easily extracted if it does not overlap with the M_1 term in the Fourier domain, which requires $\Delta \ll \omega_p/3$. However, to achieve fine axial resolution $dz \sim c/\Delta$, one needs to use a broadband SPDC source. For example, in Ref. [26] the autocorrelation MI QOCT scheme with an SPDC source was used, covering the range 684–992 nm, which corresponds to $\Delta \approx 0.08 \omega_p$. It resulted in non-negligible overlap between $\tilde{M}_0$ and $\tilde{M}_1$ peaks, which degraded the total QOCT resolution and SNR.

However, there is no M_1 term in the cross-correlation interferogram $M_{\text{MI}}^{(\text{cross})}$. It has only one parasitic term M_2 with a spectrally separated peak and a narrow bandwidth δ , which can be easily eliminated either by the Fourier processing of the interferogram or by averaging over the delay $d\tau \sim 2\pi/\omega_p$ during a continuous delay scan.

Additionally, as follows from Eq. (11), the target M_0 term in the interferogram $M_{\text{MI}}^{(\text{cross})}$ has twice the amplitude in comparison with $M_{\text{MI}}^{(\text{auto})}$ and a negative sign, corresponding to the dip shape. The M_1 term in the interferogram $M_{\text{MI}}^{(\text{auto})}$ has a positive sign, corresponding to a peak shape, since it is equivalent to the probability of detecting two photons in one output in the standard HOM interference. Moreover, since two-photon probability is usually measured with a beam splitter and two single-photon detectors, as shown in Fig. 1(c), it leads to an additional

50% loss in coincidence rate; thus, the measured M_1 amplitude in the cross-correlation scheme is 4 times bigger than in the autocorrelation scheme.

III. EXPERIMENT

To support our calculations, we conducted an experiment to benchmark the studied cross-correlation MI QOCT against both the autocorrelation scheme and classical OCT, and to assess its potential applicability to real-world scenarios. The experimental setup is presented in Fig. 4. A degenerate SPDC source (see details in Appendix B) with a 50-nm bandwidth and a central wavelength of 1313 nm was coupled into a Michelson interferometer built with a fiber beam splitter (FBS). The reference and sample arms were taken out into free space. The mirror in the sample arm was translated with a piezo or stepper-motor stage to record interferograms.

To record both cross-correlation $M_{\text{MI}}^{(\text{cross})}(\tau)$ and autocorrelation $M_{\text{MI}}^{(\text{auto})}(\tau)$ interferograms simultaneously, we inserted a fiber circulator at the interferometer input and added a second FBS at the output. Photon detection was performed with three near-infrared single-photon detectors (D1–D3, ID Qube NIR), and coincidences were registered using a Time Tagger Ultra (Swabian Instruments). Autocorrelation interferograms were obtained from D2–D3 coincidences, while cross-correlation interferograms were derived by summing D1–D2 and D1–D3 coincidences. Additionally, single-photon interferograms, related to classical OCT, were obtained from the sum of the D2 and D3 count rates.

In total, we conducted the following three experiments:

(1) Proof-of-concept. A mirror sample was used to measure the PSF of the system to confirm the absence of the M_1 term in cross-correlation measurements.

(2) Dispersion cancellation test. A 5-mm-thick silicon window was inserted before the mirror in the sample arm to demonstrate QOCT's robustness against dispersion.

(3) Defectoscopy application. Two silicon windows served as a dispersive sample, and we resolved the air gap between them. To compensate for potential nonparallelism of the window surfaces, the SPDC radiation was focused by lens L. This also potentially allows transverse scanning of the sample, as in a three-dimensionally imaging QOCT system.

IV. EXPERIMENTAL RESULTS AND DISCUSSION

A. Proof-of-concept experiment

The results of the proof-of-concept experiment are shown in the light-green block of Fig. 5.

The single-count interferogram, mimicking classical OCT, is presented in Fig. 5(a) as a thin light-purple solid

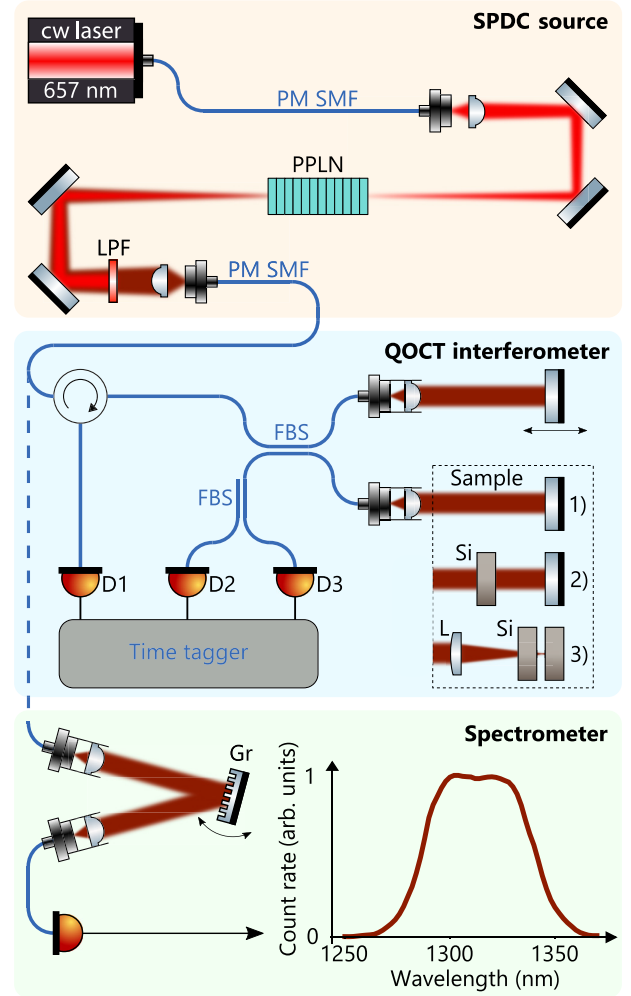


FIG. 4. Experimental setup: A collinear degenerate spontaneous parametric down-conversion (SPDC) source [type-0 periodically poled lithium niobate (PPLN) crystal pumped by a cw 657-nm laser; the pump is blocked by a 1200 nm long-pass filter LPF] is coupled into a polarization-maintaining single-mode fiber (PM SMF) and fed into a quantum optical coherence tomography (QOCT) interferometer built with fiber beam splitters (FBSs) and two free-space arms. A mirror in the reference arm is translated to record interferograms from the sample in the second arm. Samples include (1) a mirror, (2) a mirror with a silicon window, and (3) lens (L) and two silicon windows separated by an air gap. Three single-photon detectors (D1, D2, D3) connected to a time tagger measure both autocorrelation $M_{\text{MI}}^{(\text{auto})}$ and cross-correlation $M_{\text{MI}}^{(\text{cross})}$ interferograms. The SPDC output is also sent to a home-built spectrometer [two collimators, rotating grating (Gr), and a single-photon detector] to record the single-count-rate spectrum.

line. To estimate the corresponding OCT resolution (PSF width), we fitted its envelope with a sinc function (thick dark-purple dashed curve) and obtained a value of $16.6 \pm 0.1 \mu\text{m}$. Here and below, uncertainties correspond to fitting errors.

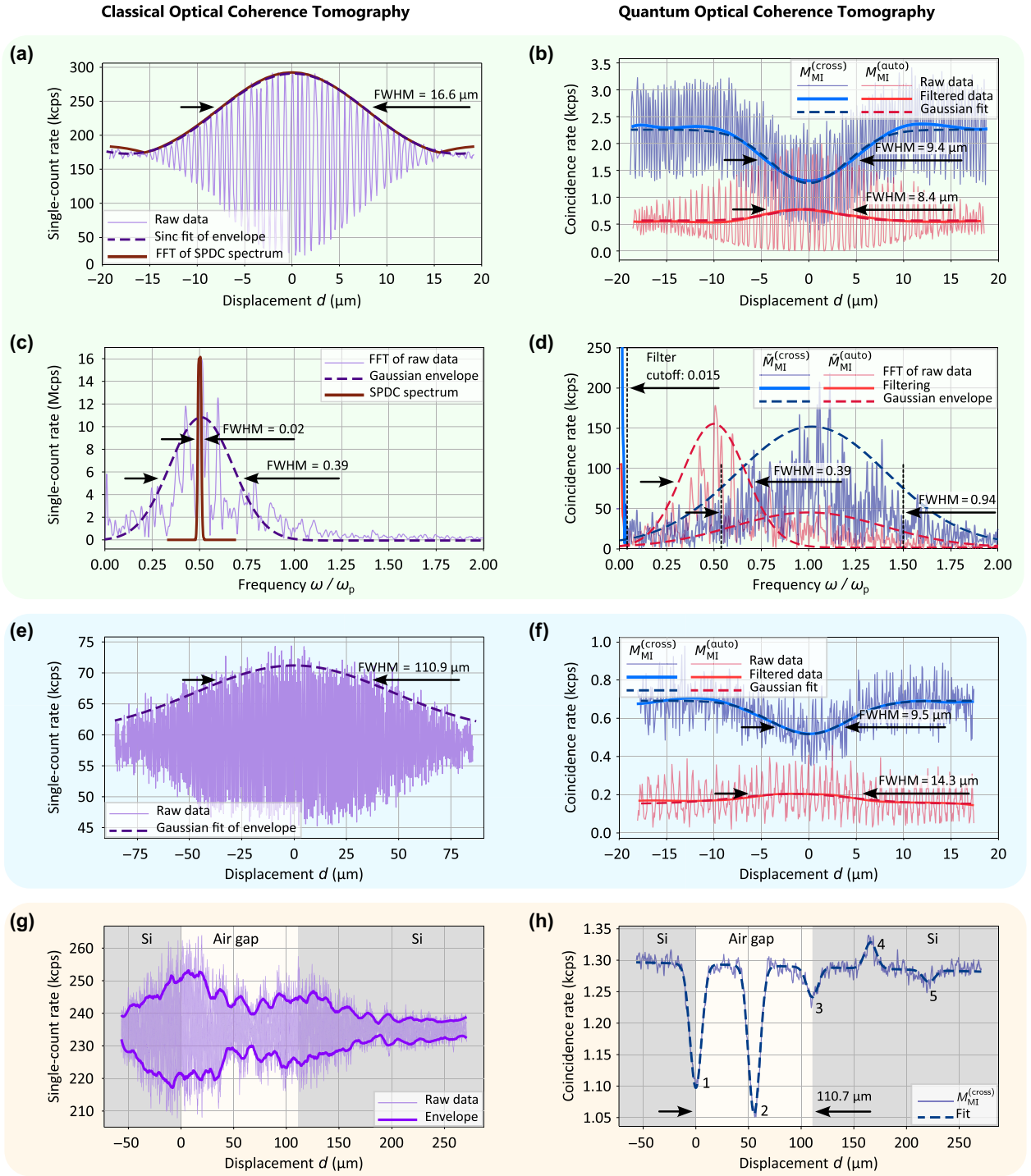


FIG. 5. Experimental results. (a) Single-count point-spread function (PSF) interferogram for optical coherence tomography (OCT): raw data (light-purple), envelope fit (thick dark-purple dashed), theoretical envelope (thick brown) obtained by inverse fast Fourier transform (FFT) of the spontaneous parametric down-conversion (SPDC) source spectrum. (b) Coincidence-count PSF interferograms for quantum OCT (QOCT): raw cross-correlation $M_{\text{MI}}^{(\text{cross})}$ (thin light-blue) and autocorrelation $M_{\text{MI}}^{(\text{auto})}$ (thin light-red) interferograms, filtered Hong-Ou-Mandel dips (thick red and blue), Gaussian fits (thick red and blue dashed). (c) Absolute FFT of the OCT interferogram in (a): raw (light-purple), envelope fit (thick dark-purple dashed), SPDC source spectrum (thick brown). (d) FFTs of the QOCT data: solid blue and red curves, with darker segments indicating the filter region; dashed lines are Gaussian envelopes. (e),(f) Single-count OCT and coincidence-count QOCT PSF interferograms with dispersion introduced by a 5-mm silicon window. (g),(h) Single-count OCT and coincidence-count QOCT PSF interferograms for a sample of two silicon windows separated by an air gap of approximately 110 μm .

The fast Fourier transform (FFT) magnitude of this interferogram is shown in Fig. 5(c). A Gaussian fit to its envelope (thick dark-purple dashed curve) yielded a spectral FWHM of $0.39 \pm 0.02 \omega_p$.

For comparison, the directly measured SPDC spectrum (already shown in the bottom of Fig. 4) is plotted as a brown line in Fig. 5(c), and its inverse FFT is shown as the same line in Fig. 5(a). We see that the interferogram envelope matches perfectly with the inverse FFT of the SPDC spectrum (as expected theoretically). However, the interferogram spectrum is nearly 20 times broader than the SPDC spectrum, indicating strong broadening due to phase noise in the fibers. Thus, although the SPDC spectrum fully determines the OCT resolution, the interferogram bandwidth is dominated by phase fluctuations.

The raw QOCT coincidence interferograms $M_{\text{MI}}^{(\text{cross})}(\tau)$ (light blue) and $M_{\text{MI}}^{(\text{auto})}(\tau)$ (light red) are shown in Fig. 5(b). These were recorded with a 70-nm step size and 100-ms exposure per point.

Their FFT magnitudes are plotted in Fig. 5(d) with the same color coding. The envelopes were fitted with the sums of Gaussian functions (dashed lines) centered at $\omega_p/2$ and ω_p (see details in Appendix D 1). As expected, the autocorrelation spectrum $\tilde{M}_{\text{MI}}^{(\text{auto})}$ exhibits both parasitic peaks $\tilde{M}_1$ and $\tilde{M}_2$. However, in contrast to the theory [Fig. 1(c)], the $\tilde{M}_2$ peak is not narrow but instead twice as broad as $\tilde{M}_1$. This can be explained as follows: the same amplitude of path fluctuation produces different phase modulations at different wavelengths, resulting in a spectral broadening of $\tilde{M}_2$ that is roughly twice that of $\tilde{M}_1$.

These results, shown in Fig. 5(d), demonstrate that even when the SPDC spectrum is relatively narrow and the parasitic peaks are not expected to overlap with the HOM term $\tilde{M}_0$, phase instability can still broaden them significantly, complicating digital spectral filtering.

By contrast, as expected, the $\tilde{M}_{\text{MI}}^{(\text{cross})}(\tau)$ spectrum shows only the $\tilde{M}_2$ parasitic peak (the fitted amplitude of $\tilde{M}_1$ is negligible).

The $\tilde{M}_0$ terms are shown in Fig. 5(d) as thick blue and red curves. To isolate them, we discarded FFT components above a cutoff frequency of $0.015\omega_p$ and then applied an inverse FFT. The resulting processed interferograms, plotted in Fig. 5(b) with the same color coding, agree well with Gaussian fits (dashed curves). The extracted FWHMs are $9.39 \pm 0.25 \mu\text{m}$ and $8.42 \pm 0.35 \mu\text{m}$, respectively—about half of the OCT PSF.

The small discrepancy likely arises because the single-count spectrum is broader than the coincidence-count spectrum due to coupling into the single-mode fiber [19].

Finally, the depth of the cross-correlation dip (995 kcps) exceeds the height of the autocorrelation peak (212 kcps) by more than a factor of 4, demonstrating improved robustness against optical losses.

B. Dispersion cancellation test

The results of the dispersion cancellation experiment are shown in Figs. 5(e) and 5(f). To demonstrate the effect of the 5-mm silicon window on OCT, we first measured the single-count interferogram (OCT PSF) in Fig. 5(e). Its envelope broadened by a factor of 6.68 to a FWHM of $110.9 \pm 6.8 \mu\text{m}$. This value agrees with the calculated value of 6.64 in Eq. (C1b), presented in Appendix C.

The raw QOCT interferograms, recorded with a 70-nm step size and 200-ms exposure per point, are shown in Fig. 5(f). We extracted the HOM terms as in the previous section. The FWHM of the processed QOCT PSFs remains nearly unchanged: $9.52 \pm 0.07 \mu\text{m}$ for cross-correlation and $14.31 \pm 0.39 \mu\text{m}$ for autocorrelation. However, the silicon window introduces additional insertion losses, reducing the coincidence count rates to 0.175 kcps (cross-correlation) and 0.051 kcps (autocorrelation), which made fitting the autocorrelation data challenging.

These results confirm dispersion cancellation in both schemes; however, the cross-correlation scheme is more robust to loss due to its higher SNR and simpler spectral filtering.

To verify that the cross-correlation scheme's SNR is suitable for real-world applications, we have conducted a demonstrative defectoscopy experiment, described in the following.

C. Defectoscopy application

In the previous experiments, we compared the two QOCT measurement schemes by performing subwavelength step-scan measurements on a piezotranslator. This fine scanning allowed us to capture the full interferogram, including high-frequency fringes, and then apply a digital low-pass filter in postprocessing to isolate the low-frequency HOM term. However, this method neither exploits the intrinsic phase-noise insensitivity of HOM interference—since it still requires subwavelength path-length control—nor overcomes the piezo's limited travel (100 μm).

For the defectoscopy application, we switched to hardware filtering. We drove a stepper-motor translator continuously at 16.7 nm/s and binned photon counts into 300-nm intervals, which averages out the high-frequency fringes in real time. Repeating each scan 5 times improved precision, and the stepper motor's 25-mm range allowed us to interrogate the full thickness of the dispersive sample without subwavelength positioning.

Figures 5(g) and 5(h) show the measured interferograms for an air gap between two silicon windows.

Both OCT and QOCT interferograms were measured in parallel. To improve the SNR, the scan was repeated 8 times.

The single-photon OCT interferogram in Fig. 5(g) is severely blurred by dispersion, making layer separation impossible. In contrast, the QOCT coincidence interferogram in Fig. 5(h) exhibits distinct peaks and dips—corresponding to reflections, echoes, and cross-interference terms [30]—allowing us to resolve the window separation. Appendix D3 provides a detailed description of each term enumerated in Fig. 5(h) and the reconstruction algorithm. From this analysis, we extract a gap width of $110.67 \pm 0.11 \mu\text{m}$, in excellent agreement with the known spacer thickness.

V. CONCLUSION

We have studied a cross-correlation QOCT scheme based on a Michelson interferometer that overcomes the main limitations of previous implementations while preserving all QOCT's advantages. Compared with classical OCT, our QOCT scheme delivers a two-fold narrowing of the PSF and remains essentially immune to the sample dispersion that increases the classical OCT PSF by more than a factor of 6.

Compared to the standard autocorrelation QOCT, the cross-correlation measurements achieve a fourfold improvement in SNR. Achieving high axial resolution requires a broadband light source. However, the autocorrelation scheme may not fully benefit from this, as it contains a parasitic peak at half the pump frequency, whose width increases with the source bandwidth. This peak overlaps with the HOM term, making it difficult to extract useful information. A cross-correlation QOCT's interferogram spectrum contains only a single parasitic peak at the pump frequency—well separated from the HOM peak—which greatly simplifies extraction of the target HOM term by postprocessing. This spectral separation reduces sensitivity to phase-instability-induced broadening and allows a compact, fiberized, vibration-insensitive implementation. Moreover, this spectral separation potentially enables combining our cross-correlation scheme with ultrabroadband SPDC sources to approach subwavelength axial resolution [26].

We have also shown that the HOM term can be directly recovered by averaging over the high-frequency oscillations during a continuous coarse mechanical scan, obviating the need for subwavelength path-length control and enabling long-range scanning. Using this approach, we have shown that the SNR of our technique is high enough to localize a $110 \mu\text{m}$ air gap in a highly dispersive silicon medium.

Thus, the presented QOCT scheme enables the use of bright, broadband collinear type-0 or type-I SPDC sources, is immune to optical dispersion and phase instability, and produces clear HOM dips with improved SNR, thus paving the way to real-world QOCT applications.

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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.

APPENDIX A: SINGLE-PHOTON AND PHOTON-PAIR INTERFERENCE IN A MICHELSON INTERFEROMETER: GENERAL DESCRIPTION

Denote the sample response function as $H(\omega)$ and delay in the reference channel as τ . The input and output modes of the interferometer are related according to Eqs. (1) and (2). Combining these equations, we obtain

$$\hat{a}_{\text{in}}^\dagger(\omega) = \alpha(\omega)\hat{a}_{\text{out}}^\dagger(\omega) + \beta(\omega)\hat{b}_{\text{out}}^\dagger(\omega), \quad (\text{A1a})$$

where

$$\alpha(\omega) \equiv \frac{e^{i\omega\tau} + H(\omega)}{2}, \quad \beta(\omega) \equiv i \frac{e^{i\omega\tau} - H(\omega)}{2}. \quad (\text{A1b})$$

1. Single-photon interference

Consider a general case of a mixed single-photon state in input mode a ,

$$\hat{\rho}_{\text{in}} = \iint d\omega d\omega' \rho(\omega, \omega') \hat{a}_{\text{in}}^\dagger(\omega) |\text{vac}\rangle \langle \text{vac}| a_{\text{in}}(\omega'), \quad (\text{A2})$$

where $\rho(\omega, \omega')$ is a cross-spectral density function.

Single-photon measurement in mode b corresponds to a projection to a state

$$|\psi_M^{(b)}(\omega)\rangle = \hat{b}_{\text{out}}(\omega) |\text{vac}\rangle. \quad (\text{A3})$$

Then the total probability of detecting a photon in output mode b as a function of the delay τ gives the interferogram

$$M^{(b)}(\tau) = \int d\omega \langle \psi_M^{(b)}(\omega) | \hat{\rho}_{\text{in}} | \psi_M^{(b)}(\omega) \rangle. \quad (\text{A4})$$

Substituting Eqs. (A1)–(A3) into Eq. (A4) we obtain

$$M^{(b)}(\tau) = \int d\omega \rho(\omega, \omega) |\beta(\omega)|^2. \quad (\text{A5})$$

Note that Eq. (A5) depends on the diagonal part of the crossed-spectral density function only, which describes the

spectral distribution of the state and does not depend on its coherence and purity. So, the same interferogram could be obtained with coherent or thermal light used in classical OCT setups.

For the Gaussian spectral distribution Eq. (4) and the single-layer response function given by Eq. (3), Eq. (A5) transforms to Eq. (5).

2. Photon-pair interference

Consider a photon-pair input state given by Eq. (6) in input mode a . Assume the spectral amplitude $F(\omega_s, \omega_i)$ is symmetrical, so that $F(\omega_s, \omega_i) = F(\omega_i, \omega_s)$, which is true for type-I or type-0 SPDC.

Substituting Eq. (A1) into Eq. (6), we obtain

$$\begin{aligned} |\Psi_{\text{in}}\rangle &= \frac{1}{\sqrt{2}} \\ &\times \left[\iint d\omega_s d\omega_i F_{\text{out}}^{(aa)}(\omega_s, \omega_i) \hat{a}_{\text{out}}^\dagger(\omega_s) \hat{a}_{\text{out}}^\dagger(\omega_i) |\text{vac}\rangle \right. \\ &+ \iint d\omega_s d\omega_i F_{\text{out}}^{(ab)}(\omega_s, \omega_i) \hat{a}_{\text{out}}^\dagger(\omega_s) \hat{b}_{\text{out}}^\dagger(\omega_i) |\text{vac}\rangle \\ &\left. + \iint d\omega_s d\omega_i F_{\text{out}}^{(bb)}(\omega_s, \omega_i) \hat{b}_{\text{out}}^\dagger(\omega_s) \hat{b}_{\text{out}}^\dagger(\omega_i) |\text{vac}\rangle \right], \end{aligned} \quad (\text{A6})$$

where

$$F_{\text{out}}^{(aa)}(\omega_s, \omega_i) \equiv \alpha(\omega_s) \alpha(\omega_i) F(\omega_s, \omega_i), \quad (\text{A7a})$$

$$F_{\text{out}}^{(ab)}(\omega_s, \omega_i) \equiv 2\alpha(\omega_s) \beta(\omega_i) F(\omega_s, \omega_i), \quad (\text{A7b})$$

$$F_{\text{out}}^{(bb)}(\omega_s, \omega_i) \equiv \beta(\omega_s) \beta(\omega_i) F(\omega_s, \omega_i).$$

Here the upper index (mn) is related to the state with one photon in mode m and another in mode n .

The projections for autocorrelation and cross-correlation measurements are given by

$$\begin{aligned} \langle \Psi_M^{(bb)}(\omega_1, \omega_2) | \Psi_{\text{in}} \rangle &= \frac{1}{2} [F_{\text{out}}^{(bb)}(\omega_1, \omega_2) + F_{\text{out}}^{(bb)}(\omega_2, \omega_1)] \\ &= \beta(\omega_1) \beta(\omega_2) F(\omega_1, \omega_2), \end{aligned} \quad (\text{A8a})$$

$$\begin{aligned} \langle \Psi_M^{(ab)}(\omega_1, \omega_2) | \Psi_{\text{out}} \rangle &= \frac{1}{\sqrt{2}} F_{\text{out}}^{(ab)}(\omega_1, \omega_2) \\ &= \frac{1}{\sqrt{2}} \alpha(\omega_1) \beta(\omega_2) F(\omega_1, \omega_2), \end{aligned} \quad (\text{A8b})$$

respectively.

Let us make a change of variables: $\omega_1 = \omega_0 + \nu_1$, $\omega_2 = \omega_0 + \nu_2$, where $\omega_0 \equiv \omega_p/2$ (half the pump frequency). In

this case, we have

$$\begin{aligned} |\beta(\omega_0 + \nu_1) \beta(\omega_0 + \nu_2)|^2 &= K_c + K_0 - (K_1 + K_{1'})/2 + K_2, \end{aligned} \quad (\text{A9a})$$

$$\begin{aligned} |\alpha(\omega_0 + \nu_1) \beta(\omega_0 + \nu_2)|^2 &= K_c - K_0 + (K_1 - K_{1'})/2 - K_2, \end{aligned} \quad (\text{A9b})$$

where

$$K_c \equiv \frac{1}{4} (|H(\omega_0 + \nu_1)|^2 + 1) (|H(\omega_0 + \nu_2)|^2 + 1), \quad (\text{A10a})$$

$$K_0 \equiv \frac{1}{2} \text{Re}[e^{-i(\nu_1 - \nu_2)\tau} H(\omega_0 + \nu_1) H^*(\omega_0 + \nu_2)], \quad (\text{A10b})$$

$$K_1 \equiv \text{Re}[e^{-i\omega_0\tau} e^{-i\nu_1\tau} H(\omega_0 + \nu_1) (|H(\omega_0 + \nu_2)|^2 + 1)], \quad (\text{A10c})$$

$$K_{1'} \equiv \text{Re}[e^{-i\omega_0\tau} e^{-i\nu_2\tau} H(\omega_0 + \nu_2) (|H(\omega_0 + \nu_1)|^2 + 1)], \quad (\text{A10d})$$

$$K_2 \equiv \frac{1}{2} \text{Re}[e^{-2i\omega_0\tau} e^{-i(\nu_1 + \nu_2)\tau} H(\omega_0 + \nu_1) H(\omega_0 + \nu_2)]. \quad (\text{A10e})$$

Therefore the final interferograms consist of five terms:

$$4M^{(bb)}(\tau) = M_c^{(M)} + M_0(\tau) - \frac{M_1(\tau) + M_{1'}(\tau)}{2} + M_2(\tau), \quad (\text{A11a})$$

$$2M^{(ab)}(\tau) = M_c^{(M)} - M_0(\tau) + \frac{M_1(\tau) - M_{1'}(\tau)}{2} - M_2(\tau), \quad (\text{A11b})$$

where

$$M_j = \iint d\nu_1 d\nu_2 |F(\omega_0 + \nu_1, \omega_0 + \nu_2)|^2 K_j \quad (\text{A12})$$

for $j = c, 0, 1, 1', 2$. It is easy to see that $M_1(\tau) = M_{1'}(\tau)$; therefore, we obtain Eq. (11) for the interferograms for both measurement schemes. Substituting the Gaussian photon-pair spectral distribution Eq. (7) into Eq. (A12), we get the final form of the interferogram terms, Eq. (12).

APPENDIX B: SPDC SOURCE

The SPDC part of the experiment employed a 656.5-nm cw pump laser (linewidth approximately 0.01 nm) to generate photon pairs via type-0 SPDC in a 25-mm periodically poled lithium niobate (PPLN) crystal with a 12.83- μm poling period (HC Photonics) at a stabilized temperature of 50 °C. The pump light was coupled to a polarization-maintaining single-mode fiber (PM SMF) and

focused into the crystal using an aspheric lens while the SPDC light was collected into telecom PM SMF using another aspheric lens. The pump and collection focal parameters [31] $\xi_p = 0.35$ and $\xi_c = 1$ were chosen as a trade-off between heralding efficiency and source brightness [32], and their values were verified using charged-coupled device (CCD) cameras to measure the beam waists at the PPLN location.

APPENDIX C: DISPERSION BROADENING

The effect of dispersion on OCT and QOCT resolution was analyzed in Ref. [26]. From Eqs. (16b,c) and (21b,c) of its Supplemental Material, the FWHM of the HOM term $M_0(\tau)$, which determines the QOCT axial resolution, increases by a factor

$$\alpha_0 = \sqrt{1 + \delta^2 \Delta^2 \kappa^2},$$

while the FWHM of the single-photon Michelson term $M_1(\tau)$, which corresponds to the classical OCT resolution, increases by

$$\alpha_1 = \sqrt{1 + \frac{\Delta^4 \kappa^2}{4}}.$$

Here,

$$\kappa = \frac{l}{c} \frac{\partial n}{\partial \omega} \bigg|_{\omega_p/2},$$

with $n(\omega)$ the refractive index of the dispersive medium and l the total round-trip path length through that medium.

Using the Sellmeier equation for silicon from Ref. [33], one finds

$$\frac{\partial n}{\partial \omega} \bigg|_{\omega_p/2} \approx -1.3 \times 10^{-16} \text{ s}.$$

For our experimental parameters—a pump linewidth of 0.01 nm ($\delta = 2\pi \times 6.9$ GHz), an SPDC bandwidth of 50 nm ($\Delta = 2\pi \times 8.7$ THz), and a silicon plate thickness $l/2 = 5$ mm—this yields

$$\alpha_0 \approx 1 + 5.5 \times 10^{-5} \approx 1, \quad (\text{C1a})$$

$$\alpha_1 \approx \sqrt{1 + \frac{(\delta^2 + \Delta^2)^2 \kappa^2}{4}} \approx 6.64. \quad (\text{C1b})$$

Thus, a 5-mm silicon window negligibly affects the QOCT resolution but broadens the classical OCT point-spread function by more than a factor of 6.

APPENDIX D: DATA PROCESSING DETAILS

1. QOCT spectra fit

Since, in general, the QOCT interferogram spectra contain two parasitic terms, $\tilde{M}_1$ and $\tilde{M}_2$, centered at $\omega_p/2$

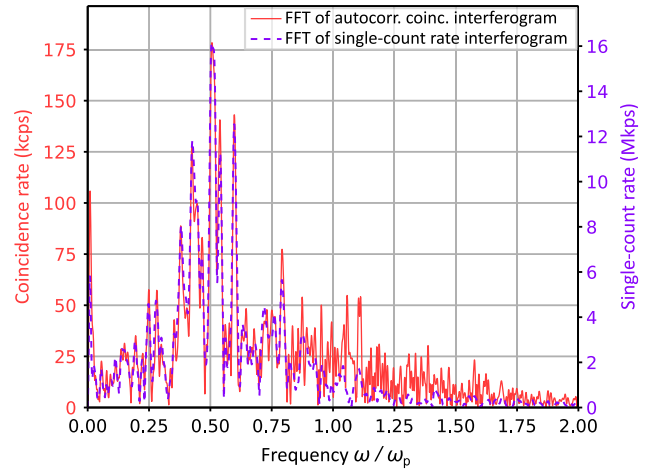


FIG. 6. Overlap of interferogram spectra from Figs. 5(c) and 5(d): fast Fourier transform (FFT) magnitude of the quantum optical coherence tomography (QOCT) coincidence-rate autocorrelation interferogram $\tilde{M}_{\text{MI}}^{(\text{auto})}$ (solid red) and FFT magnitude of the OCT single-count interferogram $\tilde{M}_{\text{OCT}}$ (dashed purple).

and ω_p , respectively, we approximate $\tilde{M}_{\text{MI}}^{(\text{cross})}$ and $\tilde{M}_{\text{MI}}^{(\text{auto})}$ [Fig. 5(d)] by the sum of Gaussian functions as follows:

$$f_{\text{fit}}(\omega) = c + a_1 \exp \left[-\frac{(\omega - \omega_p/2)^2}{2\sigma_1^2} \right] + a_2 \exp \left[-\frac{(\omega - \omega_p)^2}{2\sigma_2^2} \right]. \quad (\text{D1})$$

Because the single-count spectrum and the part of $\tilde{M}_{\text{MI}}^{(\text{auto})}$ around $\omega_p/2$ are both affected by the same phase fluctuations and therefore coincide (Fig. 6), we fixed the standard deviation of the first Gaussian to $\sigma_1 = 0.16\omega_p$, as extracted from the Gaussian fit of the single-count spectrum [Fig. 5(c), dark-purple dashed curve]. Table I summarizes the extracted values of all other parameters with fitting errors.

The amplitude a_1 extracted for $\tilde{M}_{\text{MI}}^{(\text{cross})}$ is negligibly small ($10^{-13} \pm 10$ kcps), confirming the absence of the $\tilde{M}_1$ peak. The fitted values of σ_2 for the $\tilde{M}_2$ peaks in both $\tilde{M}_{\text{MI}}^{(\text{auto})}$ and $\tilde{M}_{\text{MI}}^{(\text{cross})}$ are $0.40\omega_p$, nearly twice as large as σ_1 . This agrees with the expectation that the same optical path fluctuation amplitude induces twice the phase modulation

TABLE I. Parameters of Gaussian fit for $\tilde{M}_{\text{MI}}^{(\text{cross})}$ and $\tilde{M}_{\text{MI}}^{(\text{auto})}$.

Parameter	$\tilde{M}_{\text{MI}}^{(\text{cross})}$	$\tilde{M}_{\text{MI}}^{(\text{auto})}$
c (kcps)	4.6 ± 1.8	4.6 ± 1.3
a_1 (kcps)	$1.2 \times 10^{-13} \pm 10$	112 ± 7
a_2 (kcps)	148 ± 6	39 ± 4
σ_2 (ω_p)	0.40 ± 0.02	0.40 ± 0.05

depth at twice the frequency, leading to proportionally larger spectral broadening.

2. OCT and QOCT data processing for defectoscopy experiment

During the defectoscopy experiment, we repeated the scan 8 times. Because the source intensity was not perfectly stable within and between runs, we introduced a normalization procedure to eliminate slow drifts in brightness and make the different traces comparable. Each single-count and coincidence trace was normalized by

$$R_{\text{norm}} = \frac{R_1 + \kappa_2 R_2 + \kappa_3 R_3}{\langle R_1 + \kappa_2 R_2 + \kappa_3 R_3 \rangle}, \quad (\text{D2})$$

where R_1 , R_2 , and R_3 are the single-count rates from detectors D1, D2, and D3, respectively; $\langle \dots \rangle$ denotes averaging over the entire trace; and the coefficients κ_2 , κ_3 were chosen such that $\langle R_1 \rangle = 2\kappa_2 \langle R_2 \rangle = 2\kappa_3 \langle R_3 \rangle$.

All single-count interferograms are shown in Fig. 5(g). To align the single-count traces from different runs, we subtracted the mean and divided by the standard deviation for each trace, then rescaled by adding back the global average mean and multiplying by the global average standard deviation. To obtain an approximate envelope, we determined at each point the maximum and minimum values across all measurements (upper and lower envelopes) and then smoothed both curves with a Savitzky-Golay filter (window approximately 17 μm , third-order polynomial).

The cross-correlation interferogram $\tilde{M}_{\text{MI}}^{(\text{cross})}$ [light-blue curve in Fig. 5(h)] was obtained by averaging the coincidence data across all eight measurements.

3. Reconstruction of the sample structure

Following Ref. [30], a QOCT interferogram comprises three types of terms (dips or peaks) with amplitudes V (we define $V > 0$ for dips and $V < 0$ for peaks):

(1) Direct reflections. Reflections from each interface appear at

$$x_1 = 0 \mu\text{m}, \quad x_2 = x_1 + d,$$

with amplitudes

$$V_1 = \frac{A}{2} |\tilde{r}_1|^2, \quad \tilde{r}_1 = r_1, \quad (\text{D3})$$

$$V_2 = \frac{A}{2} |\tilde{r}_2|^2, \quad \tilde{r}_2 = r_1 r_2 t_1^2, \quad (\text{D4})$$

where A is a common amplitude factor, r_1 and r_2 are the complex reflection coefficients of the first and second interfaces, t_1 is the complex transmission coefficient of the first interface, and d is the interface separation.

TABLE II. Experimental amplitudes of the terms in Fig. 5(h) and their theoretical approximations.

Amplitude	Position	Theoretical approximation	Expt. value (cps)	Approx. value (cps)
A_1	x_1	V_1	196 ± 2	196
A_2	$x_1 + d/2$	V_{12}	237 ± 2	237
A_3	$x_1 + d$	$V_2 + V_{13}$	50 ± 2	50
A_4	$x_1 + 3d/2$	$V_{23} + V_{14}$	-38 ± 2	-40
A_5	$x_1 + 2d$	$V_3 + V_{24} + V_{15}$	20 ± 2	17

(2) Echoes. Multiple internal reflections produce echoes at

$$x_k = x_1 + kd, \quad k = 2, 3, 4, \dots$$

with visibility

$$V_k = \frac{A}{2} |\tilde{r}_k|^2, \quad \tilde{r}_k = r_1 r_2 r_{k-1}. \quad (\text{D5})$$

(3) Cross terms. Every pair of dips at x_k and x_l generates a cross term at $x_{kl} = (x_k + x_l)/2$ with amplitude

$$V_{kl} = A \text{Re} [\tilde{r}_k e^{i\omega_0 kd/c} \tilde{r}_l^* e^{-i\omega_0 ld/c}]. \quad (\text{D6})$$

Unlike the previous terms, V_{kl} can be negative (peaks). We omit the pump coherence-length factor since it significantly exceeds the interferogram length.

The interferogram of Fig. 5(h) was first fitted with the following function:

$$f_{\text{fit}}(x) = C - \sum_{j=1}^5 A_j \exp \left[-\frac{(x - x_1 - jd/2)^2}{2\sigma^2} \right] + kx, \quad (\text{D7})$$

with the fitting parameters $A_1, \dots, A_5$ [enumeration corresponds to the term numbers in Fig. 5(h)], σ , x_1 , k , and the desired separation d .

The amplitudes $A_1, \dots, A_5$ were then approximated by the theoretical expressions in Table II, where the free parameters A , r_1 , r_2 , t_1 , and $\phi_0 \equiv \omega_0 d/c$ were adjusted to minimize the discrepancy between experiment and theory.

The approximation shows a good agreement with the experimental data. Finally, fixing the amplitudes to their approximated values, we refitted the interferogram terms to extract the interface separation, $d = 110.67 \pm 0.11 \mu\text{m}$.

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