Complex conjugate resolved heterodyne swept source optical coherence tomography using coherence revival

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Abstract: We describe a simple and low-cost technique for resolving the complex conjugate ambiguity in Fourier domain optical coherence tomography (OCT) that is applicable to many swept source OCT (SSOCT) systems. First, we review the principles of coherence revival, wherein an interferometer illuminated by an external cavity tunable laser (ECTL) exhibits interference fringes when the two arms of the interferometer are mismatched by an integer multiple of the laser cavity length. Second, we report observations that the spectral interferogram obtained from SSOCT systems employing certain ECTLs are automatically phase modulated when the arm lengths are mismatched this way. This phase modulation results in a frequency-shifted interferogram, effectively creating an extended-depth heterodyne SSOCT system without the use of acousto-optic or electro-optic modulators. We suggest that this phase modulation may be caused by the ECTL cavity optical pathlength varying slightly over the laser sweep, and support this hypothesis with numerical simulations. We also report on the successful implementation of this technique with two commercial swept source lasers operating at 840nm and 1040nm, with sweep rates of 8kHz and 100kHz respectively. The extended imaging depth afforded by this technique was demonstrated by measuring the sensitivity fall-off profiles of each laser with matched and mismatched interferometer arms. The feasibility of this technique for clinical systems is demonstrated by imaging the ocular anterior segments of healthy human volunteers.

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1. Introduction

Optical coherence tomography [1] enables non-invasive, micrometer scale imaging of biological tissues over depth ranges of a few millimeters, and has found widespread use in several biomedical imaging applications, especially ophthalmic [2] and cardiovascular imaging [3]. Swept source optical coherence tomography (SSOCT) provides a dramatic sensitivity advantage over the traditional time-domain counterpart [4], and has recently been demonstrated at A-scans rates exceeding 1 MHz [5]. Unfortunately, this technique suffers from an inherent limited imaging depth due to the finite coherence length of the lasers used and the complex conjugate artifact, which result in a typical imaging range of 1 to 5 mm. While optical attenuation from absorption and scattering typically places an even more restrictive limit on the attainable imaging depth, there are several applications that would benefit from extended imaging depths, most notably ophthalmic imaging of the ocular anterior segment, intrasurgical imaging, and catheter imaging of coronary arteries.

The complex conjugate ambiguity arises from the fact that, in SSOCT, depth profiles of the sample are obtained from the Fourier transform of a spectral interferogram. Because the spectral interferogram is acquired as a real signal, its Fourier transform is always Hermitian symmetric. As a result, positive and negative displacements about the matched pathlength position cannot be distinguished. As the sensitivity of an SSOCT system decays with...
increasing distance from this matched pathlength position, the presence of the complex conjugate effectively halves the usable imaging range.

Resolving this ambiguity to double the SSOCT imaging depth is an area of interest for which a number of techniques have been developed [6–21]. These techniques include phase shifting using a PZT-mounted reference arm [6] or electro-optic phase modulator [7], heterodyne SSOCT [8–10], instantaneous acquisition of phase separated interferograms using 3x3 interferometers [12] or polarization encoding [13], harmonic lock-in detection of phase modulation [14], imparting a phase ramp across a B-scan with either B-M mode scanning [15] or pivot-offset scanning [16–18], and dispersion encoding [21]. The extended imaging range afforded by these techniques comes at a price, either in the form of reduced sensitivity, reduced axial resolution, reduced imaging speed, increased system complexity, increased cost and/or complex post-processing. Furthermore, most of these techniques provide only partial suppression of the complex conjugate artifact, which can result in distracting “ghost” images.

Of particular interest is heterodyne SSOCT [8,9], which resolves the ambiguity by creating a frequency shift that moves the peak sensitivity position away from DC, such that positive and negative displacements from that position can be discerned. A significant advantage of this technique is that it shifts the complex conjugate, rather than attenuating it, and thus does not result in distracting ghost images. Heterodyne SSOCT has previously been implemented by using acousto-optic modulators (AOMs) to apply a differential frequency shifts between the sample and reference arms [8,9]. The drawbacks of the technique mostly stem from the use of AOMs, in that they typically have large insertion losses and restricted optical bandwidths, which results in reduced imaging sensitivity and reduced axial resolution. Furthermore, data processing in traditional implementations of heterodyne SSOCT is significantly more complicated than in traditional SSOCT, requiring either hardware demodulation [9] or complicated post-processing [8].

In this work, we present a novel and extremely simple method of realizing heterodyne SSOCT using coherence revival. This technique exploits the fact that some external cavity tunable lasers (ECTLs) used for SSOCT automatically produce a phase modulated signal when used in an interferometer whose arms are mismatched by an integer multiple of the laser’s cavity length. This technique has a number of advantages over traditional, AOM-based heterodyne SSOCT in that it is simple to implement, causes no reduction of axial resolution, and requires no additional hardware beyond a traditional SSOCT system. Furthermore, the only additional processing step required is the use of a numerical dispersion compensation algorithm, which is an ordinary processing step in many SSOCT systems.

2. Theory

2.1. Coherence revival

Coherence revival refers to the phenomenon where interference fringes are observed in an interferometer illuminated by a light source with a comb-like spectrum not only when the reference and sample arms are matched in delay, but also when the two arms are mismatched at periodic intervals. These intervals can be several orders of magnitude longer than the source coherence length.

This phenomenon occurs if the light source in the interferometer is a laser simultaneously oscillating at multiple longitudinal modes. The period at which each set of interference fringes is observed is then equal to the reciprocal of the mode spacing, which is also equal to the roundtrip delay of the laser cavity. This phenomenon has been used to measure the mode spacing of multi-mode diode lasers, and is discussed extensively in [22].

Briefly, if such a laser oscillates at multiple longitudinal modes simultaneously, even if these modes have random phase relationships with respect to each other (i.e. the laser is not mode locked), the multi-mode field emitted from the laser has a periodic waveform. This periodicity stems from the fact that the mode spacing is constant, or, equivalently, that the laser cavity length is fixed. The field outside the cavity is thus periodic with a period equal to the roundtrip cavity delay [22].
2.2. Coherence revival in SSOCT

This phenomenon is important for SSOCT as many currently available and emerging commercial swept source lasers are external-cavity tunable lasers (ECTLs) that exhibit this behavior. These lasers typically contain a semiconductor gain chip inside a long external cavity (typically tens of millimeters). The long cavity provides very fine mode spacing, and because the semiconductor gain media is inhomogeneously broadened, several of these longitudinal modes can oscillate simultaneously [23]. These lasers then sweep by employing a tunable filter, located inside the cavity, that creates large loss at all but a small subset of these modes. As the filter tunes, the laser mode-hops between these finely spaced longitudinal modes, but because many of the finely spaced modes are excited simultaneously, the tuning appears smooth on a macro-scale.

An important consequence of this phenomenon is that interference fringes can be observed when the sample and reference arm are mismatched by an integer multiple of the laser cavity length. This effect can be understood as arising from the interference of sequential pulses emanating from a pulsed laser, where the first emitted pulse travels through the reference arm, and the second emitted pulse travels through a shorter sample arm. Both pulses arrive at the receiver simultaneously and with a high degree of mutual coherence. Therefore, by mismatching the interferometer arms by one cavity length, the optical path delay of the laser cavity is effectively applied in the sample arm. This concept can be extended to place any integer number of virtual cavities in the sample arm.

2.3. Phase modulation in the virtual cavity

An important consequence of this virtual cavity effect is that the optical path delay of the cavity is effectively applied in the sample arm (under conditions of coherence revival). In an ordinary SSOCT system, effects such as dispersion and phase modulation that occur in the laser cavity are common to light propagating in both the sample and the reference arm, and thus do not affect the SSOCT signal. However, when coherence revival is used to place a virtual cavity in only one arm of the interferometer, this symmetry is broken, and the optical path delay of the cavity is applied in the sample arm only. Thus, any dispersion or phase modulation that is created in the laser cavity will then affect the SSOCT signal.

One of the challenges of previous heterodyne SSOCT systems is that the AOMs used are expensive, lossy, dispersive, and difficult to implement [8,9]. A significant advantage of the virtual cavity effect is that it allows for the placement of a phase modulator directly inside the laser cavity. In fact, we have observed that at least two different models of commercially available swept source lasers create phase modulation automatically when employed in a coherence revival configuration. We suggest that the source of this phase modulation is a frequency shift due to variation of the optical pathlength (OPL) of the laser cavity over the course of the laser sweep. This frequency shift may be due to a change in the physical length of the cavity, likely as part of the tuning mechanism, or a modulation of the refractive index of some element in the cavity, perhaps due to carrier-induced changes of the refractive index of the gain media [24].

To demonstrate mathematically how a variation in the laser cavity OPL results in phase modulation, we derive an expression for the SSOCT signal in a system where the OPL difference between the reference and the sample varies during the scan. The interferometric cross term of the SSOCT signal in a system where the length of one arm changes over the course of the sweep is given by [25,26]

\[ i_s(t) \propto \cos\left(2k(t)[z_r(t) - z_n(t)]\right) \]  

where \( i_s(t) \) is the time dependent photocurrent due to the \( n^{th} \) sample reflector, \( k(t) \) is the wavenumber that is swept in time, and \( z_r \) and \( z_n(t) \) are the axial positions of the reference mirror and \( n^{th} \) reflector. The axial position of the sample reflector is allowed to vary in time during the sweep.
We initially assume that the change in the cavity OPL varies linearly with the instantaneous central wavelength of the laser sweep, \( \lambda_c \). The reflector position can then be cast as a function of \( \lambda_c \):

\[
z_n(\lambda_c) = z_{n0} + M(\lambda_0 - \lambda_c)
\]

(2)

where \( \lambda_0 \) is the central wavelength of the sweep, \( z_{n0} \) is the mean position of the \( n \)th sample reflector, and \( M \) is a parameter that describes the slope of the OPL change with wavelength (e.g., in mm/nm). We combine Eq. (1) and Eq. (2) and recast the photocurrent as a function of the instantaneous central wavenumber, \( k_c \), to yield

\[
i_n(k_c) \propto \cos(2k_c(z_r - z_{n0} - M\lambda_0) + 4\pi M)
\]

(3)

Here, the \( M\lambda_0 \) term represents the axial position shift produced by the phase modulation, and the \( 4\pi M \) term is a constant and unimportant phase shift. Thus, the axial position shift created by the cavity length variation, \( \Delta z \), can be expressed as

\[
\Delta z = M\lambda_0
\]

(4)

It is important to note that this axial position shift is created as a phase delay only. Group delay is given by the derivative of the instantaneous phase shift with respect to frequency, and although the laser cavity OPL changes over the course of the sweep, the cavity length is constant with respect to optical frequency at all times during the sweep. Thus, the cavity OPL variation creates a phase delay without creating an offsetting group delay. It is this separation between phase and group delay that enables the separation between the real image and its complex conjugate [25,26].

2.4. Coherence revival in the Fourier domain

The effects of the swept laser mode structure upon the observed SSOCT signal, including the loss of visibility in coherence revival, are readily understood in terms of simple Fourier relationships (Fig. 1). The length and finesse of the Fabry-Perot resonator cavity determine the spacing and spectral purity of the resonator modes, respectively. The transmission function of the resonator is given by [27]

\[
T_{\text{cavity}}(\omega) = \frac{T_{\text{max}}}{1 + \left(\frac{2F}{\pi}\right)^2 \sin^2\left(\frac{\pi\omega}{\omega_{\text{FSR}}}ight)}
\]

(5)

where \( T_{\text{max}} \) is the peak spectral density, \( F \) is the cavity finesse, and \( \omega_{\text{FSR}} \) is the angular free spectral range given by \( \omega_{\text{FSR}} = \pi c / n g L \). Assuming the laser has at least moderate finesse, \( F \geq 5 \), this expression is well approximated by a series of Lorentzian functions, i.e., a Lorentzian convolved with a comb:

\[
T_{\text{cavity}}(\omega) \approx \left(\frac{T_{\text{max}}}{1 + (\tau\omega)^2}\right) * \left(\sum_{m=\infty}^{\infty} \delta(\omega - m\omega_{\text{FSR}})\right)
\]

(6)

where \( \tau \) is given by \( \tau = 2F / \omega_{\text{FSR}} \) and is inversely proportional to the linewidth of the Lorentzian, and \( \delta \) denotes the Dirac delta function. The ECTL also has a tunable filter placed inside the cavity, with a passband that is much broader than the mode spacing, such that many modes will oscillate simultaneously. If we denote the transmission function of the tunable filter as \( T_{\text{filter}} \), the instantaneous spectrum of this type of laser can be expressed as

\[
S_{\text{inst}}(\omega, \omega_s) = S_{\text{source}}(\omega) T_{\text{filter}}(\omega, \omega_s) \left[\left(\frac{T_{\text{max}}}{1 + (\tau\omega)^2}\right) * \left(\sum_{m=\infty}^{\infty} \delta(\omega - m\omega_{\text{FSR}})\right)\right]
\]

(7)
Fig. 1. Time and Fourier domain representations of coherence revival. The ideal interferogram is convolved with the instantaneous source spectrum to yield the measured spectral interferogram, which is the Fourier transform of the observed A-scan. Equivalently, the ideal sample reflectivity is multiplied by the fall-off function, resulting in the observed A-scan.

where $S_{\text{source}}(\omega)$ is the integrated power spectral density of the laser sweep, and $T_{\text{filter}}$ also depends on $\omega_c$, the instantaneous central frequency of the laser that varies over sweep.

For each spectral channel of an SSOCT A-scan centered at a frequency $\omega_c$, the detected photocurrent is equal to the ideal spectral interferogram multiplied by the instantaneous spectrum $S_{\text{inst}}(\omega, \omega_c)$ and then integrated over $\omega$. This is analogous to convolving the ideal spectral interferogram with the instantaneous spectrum, and thus the sensitivity fall-off profile is related to the Fourier transform of the instantaneous spectrum. The ideal spectral
A coherence revival interferogram is defined as that resulting from the sample structure only, absent any SSOCT instrument function. For simplicity, we assume that $T_{\text{filter}}$ maintains a constant shape across the sweep, and thus the magnitude of its Fourier transform is constant. The fall-off profile is then given by directly taking the normalized magnitude of the Fourier transform of Eq. (7) with respect to $\omega$, and recasting in terms of the pathlength mismatch, $z$:

$$f_{\text{falloff}}(z) = f_{\text{filter}}(z) \ast \left[\exp\left(-\frac{|z|}{\zeta}\right)\left(\sum_{m=-\infty}^{\infty} \delta(z-mn_{\text{eff}}L)\right)\right]$$  \hspace{1cm} (8)

where $z = tc/2$, $L$ is the physical cavity length, $n_{\text{eff}}$ is the effective refractive index, $\zeta$ is the characteristic decay distance given by $\zeta = n_{\text{eff}}LF/\pi$, and $f_{\text{filter}}$ is the Fourier transform of $T_{\text{filter}}$. Because the source bandwidth is much broader than the filter’s spectral bandwidth, after Fourier transformation, the contribution of the source to the fall-off profile is negligible and has thus been dropped.

The fall-off profile in Eq. (8) is composed of a comb with a period $n_{\text{eff}}L$ that is multiplied by a double-sided exponential function with a characteristic decay distance $\zeta$. The comb is then convolved with $f_{\text{filter}}$, which is the magnitude of the Fourier transform of the tunable filter passband. As with conventional SSOCT, $f_{\text{filter}}$ defines the SSOCT fall-off profile. For coherence revival, this profile applies to each set of fringes, which are separated by the period of the comb. The exponential function, heretofore referred to as the coherence revival fall-off envelope, determines the loss of fringe visibility at increasing multiples of the cavity length. These relationships are all depicted in Fig. 1.

Conventional fall-off profiles are typically specified by the pathlength mismatch that results in a 6dB loss in sensitivity. For comparison, we derived the characteristic distance at which this envelope is reduced by the same amount:

$$\Delta z_{-6dB} \approx 0.44 n_{\text{eff}}LF$$  \hspace{1cm} (9)

This result suggests that, for ideal cavities with large finesse, the coherence revival fall-off envelope would allow the use of many cavity length offsets before fringe visibility is severely degraded.

2.5. Nonlinear cavity length variation

An important consideration that has not yet been addressed relates to the assumption that the cavity OPL changes linearly with wavelength. Equation (3) demonstrates that such a linear relationship would result in a pure phase modulation. In practice, however, the OPL change may not be linear in wavelength, and may instead exhibit a nonlinear relationship. This nonlinear cavity length variation still creates phase modulation, but the modulation frequency is chirped rather than constant. This is analogous to different wavelengths experiencing different OPLs in the sample arm, a phenomenon that is closely related to material dispersion. Therefore, while a nonlinear cavity length variation still creates phase modulation, it also results in distortion in the axial point-spread function (PSF). Fortunately, the well-established numerical techniques used to correct dispersion in SSOCT [28] can also be used to correct this PSF distortion.

3. Methods

3.1. Numerical simulations

To validate the theory developed above, we created a numerical simulation of a coherence revival SSOCT system in MATLAB. The simulation was designed to model an SSOCT system comprising an ECTL with a Gaussian-shaped tuning bandwidth centered at 1050nm, FWHM bandwidth of 50nm, cavity length of 80mm and cavity finesse of 25. These parameters were chosen to mimic the observed behavior of the Axsun laser. The $S_{\text{source}}$ term from Eq. (7) was given by
\[ S_{\text{source}}(\omega) = S_0 \exp \left[ \frac{(\omega - \omega_0)^2}{2\sigma_s^2} \right] \]  \hspace{1cm} (10)

where \( S_0 \) is a normalization term and \( \sigma_s \) is the standard deviation. The cavity transmission function from Eq. (5) was used, except that either the physical cavity length or the cavity index was allowed to vary over the sweep, resulting in the following expressions for \( \omega_{\text{FSR}} \):

\[ \omega_{\text{FSR}}(\omega_c) = \frac{\pi c}{n_{\text{eff}} L(\omega_c)} \quad \text{or} \quad \omega_{\text{FSR}}(\omega_c) = \frac{\pi c}{n_{\text{eff}}(\omega_c) L} \]  \hspace{1cm} (11)

These expressions allow the mode spacing and cavity transmission spectrum to vary as a function of the instantaneous central frequency. Finally, the tunable filter passband was modeled as a Gaussian with a FWHM of 0.05nm, which corresponds to a SSOCT fall-off profile similar to what is observed with commercially available lasers. The tunable filter passband is thus given by

\[ T_{\text{filter}}(\omega - \omega_c) = \exp \left[ \frac{(\omega - \omega_c)^2}{2\sigma_{\text{filter}}^2} \right] \]  \hspace{1cm} (12)

where \( \sigma_{\text{filter}} \) is the standard deviation of the Gaussian and \( \omega_c \) is the instantaneous central frequency of the sweep. Using these formulas, A-scans were simulated using 9216 spectral channels over a bandwidth equal to four times the source FWHM (200nm). For each spectral channel, \( T_{\text{filter}}, \omega_{\text{FSR}} \) and \( T_{\text{cavity}} \) were computed for the corresponding instantaneous central frequency \( \omega_c \), which was stepped linearly over the simulation bandwidth. The instantaneous source term was then computed using Eq. (7).

Each spectral channel of each A-scan was then computed as a single time-domain OCT measurement, as derived by Hee [29], but replacing the original source term with the instantaneous source term, and dropping the high frequency carrier:

\[ I(\omega_c) \propto \text{real} \left\{ \int_{-\infty}^{\infty} S_{\text{inst}}(\omega, \omega_c) \exp \left[ -j(\omega - \omega_c)\Delta \tau_g \right] \frac{d\omega}{2\pi} \right\} \]  \hspace{1cm} (13)

where \( \Delta \tau_g \) is the group delay difference between the sample and reference reflectors. This integral was computed for all 9216 \( \omega_c \) values in the simulation to produce a spectral interferogram, which was then Fourier transformed to yield a depth scan.

Simulations were run for a stationary cavity, a cavity whose pathlength varied linearly with the instantaneous central wavelength, and a cavity whose pathlength was varied linearly with the instantaneous central frequency. In the latter two cases, the total cavity length variation was 100µm over the sweep FWHM (50nm), and 400µm over the entire simulation bandwidth (200nm), yielding a slope parameter \( M \) of approximately 2 µm/nm.

For simulations where the cavity pathlength varied linearly in frequency, numerical dispersion compensation algorithms [28] were also applied to correct the degradation in the axial PSF. For each of these simulations, multiple A-scans were computed, modeling a single sample reflector at various delays. Fall-off profiles were then computed by plotting the peak of each as a function of the delay.

3.2. SSOCT systems at 840nm and 1040nm

To demonstrate coherence revival SSOCT in practice, two SSOCT systems were constructed using different commercially available ECTLs (Fig. 2). The first used a Thorlabs SL850-P16, with a central wavelength of 840nm, tuning bandwidth of 80nm, and repetition rate of 8kHz (forward sweep only). The balanced receiver used was a Thorlabs PDB120A, a Si receiver with 75MHz electronic bandwidth. The second system used an Axsun Technologies swept source laser with a central wavelength of 1040nm, tuning bandwidth of 100nm and repetition...
rate of 100kHz. The balanced receiver used in the 1040nm system was a Wieserlabs WL-BPD1GA, an InGaAs receiver with 1GHz electronic bandwidth. An RF amplifier (HD24388, HD Communications Corp.) was also used with the 1040nm system (not shown). An Alazar Technologies ATS9870 digitizer was used for both systems, operating at 250MS/s and 1GS/s for the 840nm and 1040nm systems, respectively. Both systems had identical topologies, and made use of the spectrally balanced interferometer configuration recently suggested by Klien et al [5]. While the fiber couplers and detectors differed between the two systems, the same digitizer and reference and sample arm optics were used. A very long motorized translator (SGSP46-400X, Sigma Koki) was used in the reference arm.

Sensitivity and fall-off measurements were made with both systems with the sample arms matched, and at various cavity length offsets. The cavity length of each laser was measured by placing an attenuated mirror in the sample arm and translating the reference arm over its entire linear travel (400mm). The distance between the peak fringe visibility positions of each set of interference fringes was determined to be the cavity length. We use the terms +1 or −1 cavity length offset to refer to the situations in which the sample arm was longer or shorter than the reference arm by one cavity length, respectively.

For the 840nm source, fall-off measurements were taken with cavity length offsets of −2, −1, 0, +1 and +2. For the 1040nm source, only the −1, 0 and +1 cavity length offsets were used, because the phase modulation imparted by −2 and +2 offsets exceeded the electronic bandwidth of the digitizer. For each system, fall-off measurements were made using consistent levels of sample and reference power across all cavity lengths offsets, to allow the relative signal levels to be compared.

Finally, to demonstrate the feasibility of this technique for in vivo imaging, the ocular anterior segments of healthy human volunteers were imaged. For these experiments, the powers incident on the patient cornea were 600µW and 1.8mW for the 840nm and 1040nm systems, respectively, which were within the limits of the ANSI Z136.1 standard. To demonstrate the improved imaging depth with coherence revival CCR, both systems were used at both 0 and +1 cavity offsets. The sample arm used consisted of two galvanometers (Cambridge technologies) and a compound objective lens designed to provide sufficient depth of field to demonstrate the extended imaging range of the SSOCT systems [25].

3.3. Wavenumber recalibration and dispersion compensation

As the SSOCT signal was sampled linearly in time, and the lasers swept nonlinearly in wavenumber, the acquired signal required resampling before Fourier transformation. Both lasers contained an internal Mach-Zehnder interferometer clock, whose signal was digitized along with the photoreceiver signal. The zero-crossings of the clock were detected and used to generate a linear-in-wavenumber recalibration vector that was used to resample the SSOCT signal linearly in wavenumber. However, because the clock signals were only intended for
imaging depths of 2.9mm and 3.7mm (for the 840nm and 1040nm systems, respectively), the recalibration vector was first interpolated to increase the achievable imaging depth to 9.4mm and 12.4mm, respectively. The photoreceiver signals were then resampled using this recalibration vector via linear interpolation.

The numerical simulations described above demonstrate how a nonlinear-in-wavelength cavity length variation results in axial PSF degradation, analogous to the effects of material dispersion. In fact, we observed such an axial PSF degradation in our experiments. These dispersion-like effects, as well as true dispersion from unmatched optics and fiber lengths in the sample and reference arm, were corrected using a numerical dispersion compensation [28]. Briefly, after resampling to linearize the spectral interferogram in wavenumber, the spectral interferogram was multiplied by a complex phase function, given by:

$$DC(k) = \exp\left(-j\left(a_1(k-k_0)^2 + a_2(k-k_0)^4\right)\right)$$  \hspace{1cm} (14)

where $a_1$ and $a_2$ are fitting parameters and $k_0$ is the central wavenumber of the sweep. Optimal values of $a_1$ and $a_2$ were determined using an optimization algorithm to maximize the peak signal from a mirror.

4. Results

4.1. Numerical simulations results

Results from the numerical simulations are shown in Fig. 3, Fig. 4 and Fig. 5. Figure 3 shows the results from the simulation with a fixed cavity, demonstrating coherence revival without a frequency shift. Figure 4 and Fig. 5 show the results from simulations in which the cavity length varied linearly with the instantaneous central wavelength and instantaneous central

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Fig. 3. (A-C) Fall-off plots from a numerical simulation of a coherence revival SSOCT system with a stationary cavity for 0, +1 and +2 cavity length offsets, respectively. Note that no shift in the peak sensitivity position is observed. (D) Fall-off profile demonstrating coherence revival peaks centered at 80mm and 160mm.
Fig. 4. Fall-off plots from a numerical simulation of a coherence revival SSOCT system with a cavity whose pathlength varies linearly in wavelength with a slope of 2µm/nm, +1 (A) and +2 (B) cavity length offsets. The expected axial position shifts of the peak sensitivity positions of 2.1 and 4.2mm are observed.

Fig. 5. A and B: Fall-off plots from a numerical simulation of a coherence revival SSOCT system with a cavity whose pathlength varies linearly in frequency with an average slope of 2µm/nm, for +1 and +2 cavity length offsets, respectively. The expected axial position shifts of the peak sensitivity positions are observed, but the axial PSFs are severely degraded due to the nonlinear phase modulation. C and D show the same data as A and B after applying the dispersion compensation algorithm from Eq. (14).

frequency, respectively. For the linear-in-wavelength case, the slope of the cavity variation was precisely 2µm/nm, while the linear in frequency case used an average slope of 2µm/nm. From Eq. (4), the expected axial shifts for cavity offsets of 0, +1 and +2 are 0mm, 2.1mm and 4.2mm. To allow for relative comparisons, all simulation figures are plotted on the same vertical scale. Also, as the simulated spectral interferograms are real signals, their Fourier
transform remain Hermitian symmetric. Thus, only positive (frequencies) displacements are shown. Furthermore, the abscissa in the fall-off plots is the axial position that is extracted from the Fourier transform, and thus does not represent the axial position of the reflector for coherence revival peaks. Conversely, the fall-off profile in Fig. 3D was extracted by plotting the peak of the Fourier transformed SSOCT signals against the simulated reflector positions that generated the corresponding signals. Fall-off profiles for the simulations in which the cavity length varied are not shown but were identical to Fig. 3D.

4.2. Experimental fall-off measurements

Fall-off profiles from the 840nm system are shown in Fig. 6. A peak sensitivity of 95dB was measured near the 0 cavity offset with 600µW incident on the sample. This was reduced from the theoretical shot noise limit of 104dB due to coupling losses, unbalanced RIN, and digitization noise due to the low effective number of bits of the digitizer (6.7 ENOB). Sensitivity relative to this peak value is plotted on the ordinates of Fig. 6. Results are shown for the 0, +1 and +2 cavity offsets. Measurements at −1 and −2 cavity offsets were also made, but the results have been omitted here as they were nearly identical to the results from +1 and +2 offsets, respectively. The physical pathlength difference between the peak visibility positions of the 0 and +1 offsets was 66.1mm. The physical pathlength difference between the

![Graph A](image)

![Graph B](image)

![Graph C](image)

Fig. 6. Fall-off measurements from the 840nm system for 0 (A), +1 (B) and +2 (C) cavity length offsets. The physical separations between the peak sensitivity position from 0 to +1 and from +1 to +2 were both 66.1mm.
Fig. 7. Fall-off measurements from the 1040nm system for \(-1\) (A), 0 (B) and \(+1\) (C) cavity length offsets. The physical separations between the peak sensitivity positions of the \(-1\) and 0 and the 0 and \(+1\) offsets were 115.0mm and 114.8mm, respectively.

The peak visibility positions of the 0 and \(+2\) offset was precisely double (within the resolution of the translation stage), at 132.2mm. Exactly the same distances were observed for the negative offsets. These measurements were in good agreement with the manufacturer’s estimate of the cavity length as approximately 50mm of physical pathlength, without accounting for refractive index.

Fall-off profiles from the 1040nm system are shown in Fig. 7. A peak sensitivity of 98dB was measured near the 0 cavity offset with 1.8mW incident on the sample. The theoretical shot noise limit for this system was also 104dB, and the discrepancy between the measured and theoretical sensitivity was due to the same factors discussed above, as well amplification noise from the RF amplifier. Sensitivity relative to this peak value is plotted on the ordinates of Fig. 7. Results are shown for the \(-1\), 0 and \(+1\) cavity offsets; the \(-2\) and \(+2\) offsets could not be measured as, for those offsets, the phase modulation created by the laser cavity up-converted the spectral interferogram beyond the digitization bandwidth of the digitizer. The physical pathlength difference between the \(-1\) and 0 and the 0 and \(+1\) offsets were 115.0mm and 114.8mm, respectively. Again, this was in good agreement with the pathlength suggested by the manufacturer of approximately 80mm of fiber.
Because the sample and reference powers were kept constant during these measurements, we can determine the loss in peak sensitivity for each cavity length offset from the falloff profiles. Figure 6 shows that the peak sensitivity in the +1 and +2 offsets for the 840nm system were attenuated by 5dB and 10.5dB, respectively, from that of the pathlength matched case. If we define the usable imaging range as the depths over which the signal amplitude is reduced by less than 6dB, the imaging ranges for the 840nm were approximately 2.5mm, 5mm and 4mm for the 0, +1 and +2 offsets.

Similarly, Fig. 7 shows that the loss in sensitivity at the −1 and +1 offsets was only about 1dB, despite the considerably longer cavity length. From this, we infer that the finesse of the 1040nm laser was much higher than the finesse of the 840nm laser. The usable imaging ranges were 9mm, 5.5mm and 9mm for the −1, 0, and +1 offsets, respectively.

It is interesting to note that the optimal dispersion compensation phase functions were nearly identical between the +1 and −1 cavity length offsets for both systems. Furthermore, the phase function parameters $a_1$ and $a_2$ used to optimally correct measurements from −2 and +2 cavity length offsets (for the 840nm system) was precisely double that of the phase function used for the −1 and +1 cavity length offsets.

4.3. Imaging results

Figure 8 shows the results of two images of the same volunteer’s ocular anterior segment for comparison. The image in A was taken with the reference and sample arms matched in pathlength, whereas the image in B was taken with the sample arm one cavity length longer than the sample arm. Both images were acquired on the 840nm system, and each image comprises five averaged frames. The locations of the zero pathlength difference (ZPD) position and the +1 offset position are indicated. Figure 9 shows the results of the same experiment conducted on the 1040nm system. Figure 10 shows two volume projections taken on the same eye with both the 840nm (left) and 1040nm (right) systems.

Fig. 8. Comparison images taken on the 840nm, 8kHz (Thorlabs SL850-P16 laser) system with 0 (A) and +1 (B) cavity length offsets. Both images comprise 1000 (lateral) x 1300 (axial) pixels spanning 13 mm (lateral) x 5.3 mm (axial), the latter scaled to account for refractive index. Each image represents 5 averaged frames obtained over 0.6s. The locations of the ZPD and +1 offset positions are indicated.
Fig. 9. Comparison images taken on the 1040nm, 100kHz (Axsun Technologies laser) system with 0 (A) and +1 (B) cavity length offsets. Both images comprise 2000 (lateral) x 2300 (axial) pixels spanning 14 mm (lateral) x 6.9 mm (axial), the latter scaled to account for refractive index. Each image represents 5 averaged frames obtained over 100msec. The locations of the ZPD and +1 offset positions are indicated.

Fig. 10. Volume projections of the same eye acquired with the 840nm (left) and 1040nm (right) systems. The 840nm volume consisted of 1300 (axial) x 500 x 200 samples, acquired in 12.5 seconds. The 1040nm volume consisted of 2304 (axial) x 500 x 200 samples, acquired in 1 second.

5. Discussion

Coherence revival is an attractive implementation of heterodyne SSOCT, and carries with it a number of advantages over traditional methods employing AOMs. First and foremost, the method is simple and inexpensive; in cases where the laser already exhibits phase modulation, all that is required is an adjustment of the reference arm length and an increase in the digitization speed. Second, while there is an associated loss in sensitivity, the magnitude of this loss depends on the laser design (primarily the cavity finesse). We have shown that, for at
least one commercially available swept source laser, this loss in sensitivity is only about 1 dB. Finally, no additional complicated signal processing or image processing techniques are required. The axial PSF degradation observed in the coherence revival configurations can be managed using numerical dispersion compensation, a common processing step in SS OCT.

The dispersion compensation algorithm used in our experiments [Eq. (14)] employed two fitting parameters, $a_1$ and $a_2$, which are related to the group velocity dispersion (GVD) and third order dispersion (TOD), respectively. These fitting parameters can be used to quantify the GVD and TOD of the system [30], although it is important to note that the parameters were optimized to correct dispersion-like effects of coherence revival as well as true material dispersion due to unmatched optics and fiber lengths between the sample and reference arm. For the 1040 nm laser, the group delay dispersion and total TOD of the system, when operated in a +1 cavity length offset configuration, were measured to be 9400 fs$^2$ and 55700 fs$^3$, respectively. Assuming this dispersion occurs over the cavity length of 115 mm, and accounting for the index of refraction of the fiber in the cavity (Corning HI 1060, $n \approx 1.47$), this implies a GVD of 60.0 fs$^2$/mm (or $-105$ ps/nm/km) and a TOD of 356 fs$^3$/mm. The manufacturer’s reported value for GVD of HI1060 fiber is $-38$ ps/nm/km, suggesting that there is significantly more dispersion in the system than would be caused by the fiber in the cavity alone. For the 840 nm laser, the group delay dispersion and total TOD of the system, when operated in a +1 cavity length offset configuration, were measured to be 5300 fs$^2$ and 22800 fs$^3$, respectively. Assuming the dispersion occurs over the 66.1 mm cavity length offset and that the cavity is free-space, this implies a GVD of 40.3 fs$^2$/mm (or $-108$ ps/nm/km) and a TOD of 172 fs$^3$/mm.

The required digitization rate depends on the laser sweep speed, desired imaging depth and the frequency of the phase modulation created by the ECTL. As the up-converted frequencies must be Nyquist sampled, the required digitization rate is equal to the digitization rate of a conventionally configured (i.e. non-heterodyne) SS OCT system plus twice the modulation frequency. For the lasers used in our experiments, this resulted in approximately a two-fold increase in required digitization bandwidth as compared to conventional SS OCT.

We observed the generation of two distinct types of image artifacts when using this method. First, with the 840 nm laser, we observed the appearance of faint but sharp “ghost images” in the fall-off plots, only at the deepest end of the imaging depth. These artifacts appeared even when the cavity length offset was zero, and can be clearly seen in Fig. 3A, as faint reflectors between 6 and 10 mm with amplitudes between $-30$ and $-50$ dB. However, these artifacts were not sufficiently bright to appear in biological images. The second type of artifact we observed was seen with both lasers, and was characterized by the appearance of highly dispersed ghost images near the deepest end of the imaging depth. The amplitudes of these artifacts were measured and compared to the amplitude of the desired signal. The relative artifact amplitude depended on the axial position of the true reflector signal, and ranged between $-36.5$ dB and $-42.5$ dB for the 1040 nm laser and between $-25.5$ dB and $-33.5$ dB for the 840 nm laser. Examples of these artifacts can be seen in Fig. 8B and 9B, at the top of the images, as faint ghosts of the pupil (in 8B) and cornea (in 9B).

We attribute these artifacts to two sources. First, nonlinearities in the cavity length variation might give rise to multiple phase modulation frequencies, or even harmonics of the phase modulation frequency that are then aliased into the passband of the system electronics. These higher order modulation frequencies would create additional “ghost” images centered at different depths. Second, the k-clocks used for the 840 nm and 1040 nm sources were designed for imaging depths of 2.9 mm and 3.7 mm, respectively, and were not intended to be interpolated out to 9.4 mm and 12.4 mm. Thus, the artifacts may also be caused by inaccuracies in the wavenumber recalibration.

In practice, these artifacts only appeared at the deepest imaging depths where the sensitivity was poor, and were also so faint that for biological imaging, they were only visible in the presence of very bright reflectors or averaged images. Nevertheless, the wavenumber recalibration issue can be easily addressed in future designs employing the same lasers by...
constructing a Mach-Zehnder interferometer with a longer mismatch, rather than using the lasers’ internal clock. Addressing the nonlinearity of the cavity length variation is a more challenging problem, and may not be necessary as the artifacts were generally unobtrusive.

It is worth noting that the phase modulations used to generate all of the experimental data in this work were likely generated as an unintended by-product of the designs of the two lasers used. One could envision a laser deliberately designed to create a cavity length variation that would exhibit significantly better performance. An ideal laser would create a pure, linear-in-wavelength cavity length variation (and thus create fewer, if any, ghost image artifacts), would have a high finesse (ensuring that the sensitivity loss with increasing cavity length offset is minimized), and might even allow for user control of the axial position shift (by adjusting the slope of the cavity length variation). Such a laser would be valuable for extended depth SSOCT imaging applications, and would even further simplify this technique for resolving the complex conjugate ambiguity in SSOCT.

Although not addressed in this paper, a particularly useful application of coherence revival-based SSOCT is the simultaneous imaging of multiple depths. We have recently demonstrated the use of such a system to simultaneously image the anterior segment and retina of healthy human volunteers [31]. This was accomplished by constructing a dual-path sample arm that matches the retinal imaging path to the reference arm, and offsets the anterior segment imaging path by one cavity length. This technique could, in theory, be extended to encode as many imaging depths as would be allowed by the source’s coherence revival fall-off envelope.

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