

Steps Towards a Tunable Vacuum Ultraviolet Light Source for TR-ARPES

by

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The following individuals certify that they have read, and recommend to the Faculty of Graduate and Postdoctoral Studies for acceptance, the thesis entitled:

Steps Towards a Tunable Vacuum Ultraviolet Light Source for TR-ARPES

submitted by **Martin Cross** in partial fulfillment of the requirements for the degree of **Master of Applied Science** in **Engineering Physics**.

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Abstract

We detail the steps towards a tunable vacuum ultraviolet light source for time- and angle-resolved photo emission spectroscopy (TR-ARPES); specifically, the construction and characterization of a ytterbium-fiber based pump source, and a nonlinear broadening and compression stage using gas-filled hollow-core fiber (HCF). The ultimate goal is generating sub-30 fs pulses with photon energies between 5 eV and 11 eV via resonant dispersive wave (RDW) generation in a succeeding HCF. The interim goal discussed herein is producing 30 fs, infrared ($\sim 1 \mu\text{m}$), pulses with sufficient energy (μW level) for RDW generation. Using second harmonic generation (SHG) frequency resolved optical gating (FROG) we currently measure 174 fs, 4.0 μJ pulses. We will provide explanations for the shortcoming and suggest several possible solutions to address it.

Lay Summary

We detail the steps towards a tunable ultraviolet light source for measuring the properties of new materials. Ultimately, the light should be delivered in very short pulses—on the order of 30 “femtoseconds” in duration, which is to a second, what 4 minutes is to the age of the universe. These pulses act like a camera flash in a dark room, briefly capturing the behavior of the contents. The interim goal discussed herein is producing 30 fs, infrared, pulses with sufficient energy to produce ultraviolet pulses in a succeeding stage. We currently generate and measure 174 fs pulses. We will provide explanations for the shortcoming and suggest several possible solutions to address it.

Preface

The text, figures and tables in this thesis are original, unpublished, works of the author, M. Cross, except Figure D.2 and Figure D.5, which are reproduced with permission from Elsevier, and Figure D.1, which was generated by Edmund Kelleher. The work detailed within is divided as follows:

The Yb-fiber pump source was built by Edmund Kelleher, Evgeny Ostroumov, and me. The data used in this thesis were collected by the three of us.

The pi-cam (Section 2.6.2) was developed by Egor Peshkov and Edmund Kelleher. The data used in this thesis were collected by me.

The “Ref. PSD” (Section 2.6.2) was developed by Jiayi Tang and Arthur Mills. The data used in this thesis were collected by me.

The Spectral broadening stage was built by Edmund Kelleher, Federico Belli, Evgeny Ostroumov, and me. The data used in this thesis were collected by me. Gas cells were provided by John Travers. Initial simulations (Figure D.1) were performed by Edmund Kelleher and John Travers. B-integral calculations for the windows were done by Federico Belli.

All simulations and analyses were carried out by me where not specifically noted. Edmund Kelleher provided a Python–method for solving the generalized nonlinear Schrödinger equation, which was used in my simulations.

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Chapter 1

Introduction

Since the discovery of the photoelectric effect, sources of high-energy radiation — ultraviolet and beyond — have become increasingly important for studying the electronic properties of materials. Fundamentally, the more tightly an electron is bound to a material, the greater the energy required to remove it. In this regard, the study of high work function materials, or core electrons, necessitate higher energy photons. Further requirements arise if in addition we want to explore the momenta and transient characteristics of these electrons. Techniques such as angle resolved photo emission spectroscopy (ARPES) use the preservation of in-plane momentum to spatially disperse electrons photo-emitted from a sample, though crucially, the rate of dispersion (in-plane versus out-of-plane displacement of the electron) is also controlled by the incident photon energy. Being able to tune the photon energy allows for broader, or narrower, electron emission cones, and therefore higher resolution, or greater range, in momentum-space respectively. For this purpose, monochromating discrete spectral lines of a gas-discharge lamp can suffice [29]. That being said, if we seek to reveal the non-equilibrium dynamics of these photo-excited electrons, a pulse of light on the order of tens of femtoseconds in duration is required. The only methods for creating such pulses in the ultraviolet (outside of institutional-scale free electron lasers, which suffer from timing jitter [16]) involve a two-step process, where low-photon-energy light is generated in a laser and subsequently frequency up-converted.

Sum-frequency generation (SFG) in a nonlinear optical (NLO) crystal is perhaps the most straight forward approach. Conversion efficiencies of a few percent allow for modest pulse energies at megahertz repetition rates, limiting space-charge effects while preserving a rapid acquisition rate. The technique is however hindered by reabsorption in the crystal at energies above 6 eV for beta barium borate (BBO), and 7 eV for the more exotic potassium beryllium fluoroborate (KBBF). Unfortunately, these photon energies do not reveal the full Brillouin zone in a variety of materials.

High harmonic generation (HHG) is an alternative technique that limits the effects of reabsorption via a short interaction length in the generating

media, and routinely produces photon energies of 40 eV [12] [27]. Typically, mJ level pulses at kHz repetitions rates [11] are focused through a gas jet only tens of microns in diameter [39]. In contrast to SFG, high pulse energies are required, as the conversion efficiencies are only on the order of 10^{-6} [12]. In recent years, however, advances in ytterbium-based gain media, and novel approaches in the “tight focusing regime,” have allowed for MHz repetition rates [39]. An elegant example is the use of an external femtosecond enhancement cavity (fsEC) to build up the necessary kW level average powers [27]. A more brute-force approach involves parallel amplification, and subsequent coherent combination of several pulse trains [14]. Of course, these advances come at the cost of flexibility and simplicity, with strict restrictions on the repetition rate for fsEC, and precise phase delays necessary for coherent combination of multiple pulses. In a sense, these systems begin to leave the domain of convenient “table-top” sources.

Recently, new avenues for generating deep-UV radiation between 5 eV and 11 eV have opened through the advance of hollow-core photonic crystal fiber (PCF) technology [47]. These micro-structured hollow fibers mimic capillaries with small core diameters, though do so with a two decade reduction in loss [47]. Simply varying the species and pressure of a filling gas allows one to transition from controlled spectral broadening, to more striking soliton dynamics where self-compression can lead to emission of a widely tunable, resonant dispersive wave (RDW) in the ultraviolet. This technique far exceeds the photon energies achieved through SFG in a NLO crystal, and is comparatively simple and flexible with respect to HHG – the repetition rate can be varied arbitrary, and conversion efficiencies on the order of 10^{-2} [6] allow conventionally attainable, μ J-scale, pulses to be used.

Already, a source of this type has demonstrated potential in ARPES, and shows promise for extensibility [6]. In fact, much has yet to be explored in this sub-field: numerical simulations show the possibility for extending photon energies to the extreme-UV region [13], and it appears unknown, though possible, that RDWs track the polarization of the input pulse. Controlling the handedness of the probe could make such sources an asset for spin-resolved photo emission spectroscopy.

With that said, herein we will discuss initial progress toward the construction of a new VUV RDW source, targeting photon energies from 5 eV to 11 eV, repetition rates from 300 kHz to 20 MHz, photon fluxes on the order of 10^{12} photons/s, and hopefully, controllable polarization. Additionally, we aim to use a monochromator to balance time and energy resolution between 30 fs and 10 meV respectively. Table 1.1 compares this source to several alternatives in use.

This thesis details construction of the source up to and including the second compressor in Figure 1.1, and is organized as follows: Chapter 2 covers the first stage of the source, starting at the oscillator and ending at the first compressor. This stage produces pulses of infrared light with sufficient peak power for spectral broadening and RDW emission in the following gas-filled hollow-core fiber (HCF) stages. Such pulses would ordinarily suffer from excessive nonlinear broadening and phase, thus, we have used a grating-based pulse stretcher and compressor to implement a chirped pulse amplification (CPA) scheme. The three amplifiers (Yb-SMF, Yb-PCF, and Yb-Rod) are all constructed of solid-core, Yb-doped fiber, with progressively increasing modal areas to reduce the optical intensities within. In addition, we have used an acousto-optic modulator as a “pulse-picker” to vary the repetition rate by integer division of the fundamental 38 MHz frequency. We will outline the theory of operation and construction of this first stage, and provide typical pulse characteristics (ex. optical power, autocorrelation, optical spectrum) at several points for diagnostic purposes. The final output will be further characterized by second harmonic generation frequency resolved optical gating (SHG-FROG). At a repetition rate of 4.9 MHz, we currently generate up to 4.5 μ J pulses, centered at 1035 nm, with a FWHM duration of 264 fs.

Chapter 3 discusses the second stage of the source, where pump light from the first stage is spectrally broadened in a gas-filled HCF, and subsequently compressed to a shorter duration with a grating pair. The reduction in pulse duration is necessary for efficient RDW emission in the third stage. We will go over the general properties of HCFs, and the spectral broadening/compression scheme. Subsequently, we will use the HCF’s dispersion profile and nonlinear refractive index, as a function of fiber geometry and gas parameters, to generate a model of the pulse’s spectral and temporal evolution. Simulations will be compared to SHG-FROG measurements of the pulse after the HCF, and after the second grating compressor. At a repetition rate of 4.9 MHz, we currently measure pulses with an energy of 4.0 μ J, and a duration of 174 fs after the second compressor. Concluding this chapter is a conceptual discussion of the third stage, where soliton self-compression, together with proper phase matching conditions, leads to RDW emission that can be further monochromated.

Chapter 4 provides schematics, source code, and operating procedures for the peripheral equipment needed to run and safe-guard the source.

Table 1.1: Comparison of typical TR-ARPES sources ^a.

	NLO crystal SFG		HHG		DWG	
Reference	[44]	[11]	[12] ^b	[27] ^c	[6]	Projected
Photon Energy (eV)	6.0	6.28	15-40	8-40	5.5-9	5-11
Pulse Width (fs)	160	67.5	100	150	-	-
Repetition Rate	80 MHz	250 kHz	10 kHz	60 MHz	1 kHz	0.3-20MHz
Photon Flux (photons/s)	-	10^{13}	3.6×10^7	10^{11}	10^9	10^{12}
Energy Resolution (meV)	22	70	90	22	-	10
Temporal Resolution (fs)	163	65	125	190	-	30

^a Reference Zhou et al. [59] for a more thorough comparison of available sources.

^b Multiple available harmonics; specifications are given for the 23rd harmonic at 35.6 eV.

^c Multiple available harmonics; specifications are given for the 21st harmonic at 25 eV.

- These values were unspecified.

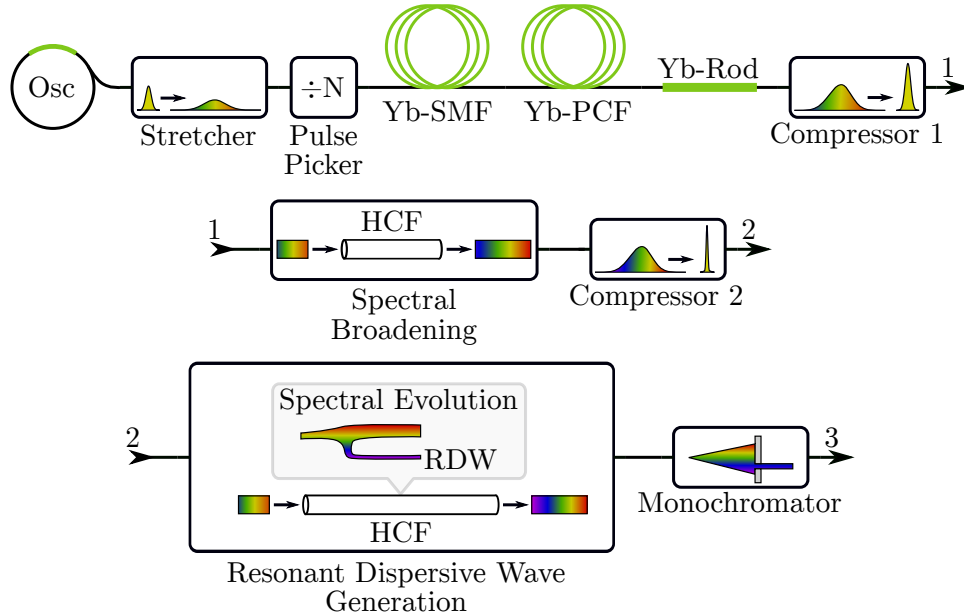


Figure 1.1: Block diagram of the system showing the three stages: infrared ($1\mu\text{m}$) pulse generation and amplification; spectral broadening and pulse compression; resonant dispersive wave generation and monochromatization.

Chapter 2

Ytterbium Fiber Based Pump Source

2.1 Oscillator

2.1.1 Design Principles

A Yb-doped fiber oscillator, with a ring topology, is the start of our system. The oscillator operates in the transition between dispersion-managed, and dissipative soliton regimes [51], where a compromise between pulse energy and temporal compressibility is found. The cavity length, and hence repetition rate, is chosen to allow subsequent frequency division by pulse-picking.

As designed, this oscillator should operate in the dispersion-managed regime, where mode-locking is effected using nonlinear polarization evolution (NPE) in the fiber section. However, we observe operation in the dissipative soliton regime, which requires dissipative (lossy) elements that have not been specifically added. In particular, there is no deliberate spectral filtering, which is crucial to most oscillators operating in the dissipative regime [8, 54].

We will start with a “walk-through” of the oscillator, describing each component’s intended role, then introduce the two pertinent operating regimes, and the “map-strength” parameter that determines which is predominant.

2.1.2 Implementation

It is easiest to describe the operation of the oscillator by following the pulse through one pass of the cavity. In the following, refer to Figure 2.1 for component numbering.

Conceptually, starting just after isolator (1) is simplest, as the pulse is assured to be linearly polarized. Half waveplate (HWP) (2) and quarter waveplate (QWP) (3) act as a polarization controller, and can be used to access any pure polarization state, given the input is linearly polarized [9]. For NPE to occur, we must impart some ellipticity to the beam polarization

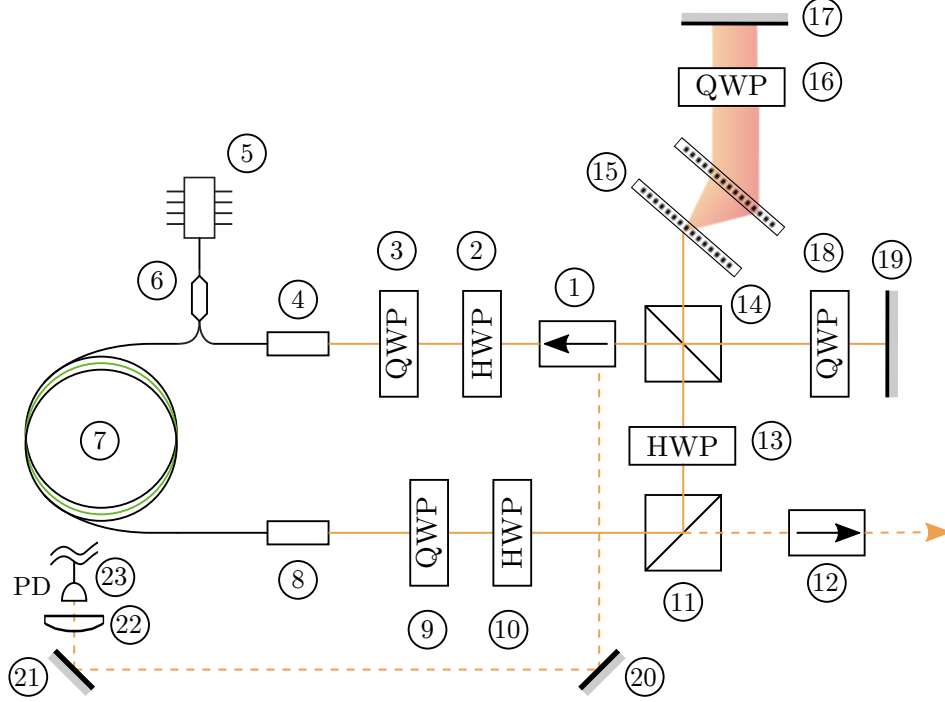


Figure 2.1: Optical layout of the oscillator. HWP, half waveplate; QWP, quarter wave plate; PD, photo diode. We use a ring topology for the cavity. The path traversed within the cavity is indicated by solid yellow lines, and the direction is set by isolator (1). Isolator (1) is rotated to pass P-polarized light, and the rejected portion is sampled with a photo diode for use in the interlock and AOM circuits. The output passes through isolator (12).

before launching it into the fiber section [24].

Collimator (4) couples the beam into the fiber, and soon after it is combined in wavelength division multiplexer (WDM) (6) with pump-light at 976 nm from diode (5). Within the fiber section (7), dispersion, self phase modulation (SPM), and gain, work in tandem to evolve the temporal and spectral profiles of the pulse. In addition, the same nonlinear refractive index responsible for SPM rotates the polarization ellipse in proportion to the intensity; since the pulse intensity varies with time, the polarization is now inextricably time dependent.

The polarization's temporal dependence allows us to select only the most intense portion of the pulse for another round trip of the cavity, and reject the rest. Exactly how polarization varies with time was not determined,

but in theory, QWP (9) and HWP (10) are set to vertically polarize the most intense portion of the pulse, then polarizing beam splitter (PBS) (11) rejects the rest through isolator (12). The output of the laser is in fact the rejected portion, though given there are 38 million cycles of the cavity per second, an equilibrium is quickly established and the output closely mirrors the circulating pulse.

The recirculated portion passes through HWP (13), which allows the beam to pass directly through PBS (14), and into grating pair (15). Given the fiber section is entirely normally dispersive, we use grating pair (15) to introduce a compensating amount of anomalous dispersion, otherwise the pulse would dissipate further with each cycle. As before, QWP (16) simply acts as a directional switch, so that the beam returned by mirror (17) reflects off PBS (14). QWP (18) and mirror (19) act in exactly the same fashion, and couple the beam back into the fiber section.

2.1.3 Mode-Locked Operating Regime

“Mode-locking” is the act of fixing the phase between various modes of the oscillator cavity, forming a pulse train. In the dispersion-managed regime, locking is effected by conservative processes—dispersion and self phase modulation (SPM) [15]. The first is a linear addition of phase and the second is nonlinear, i.e. intensity dependent. In the dissipative regime, modes are additionally grown and shed on each pass of the cavity. Eliminating some of the modes “resets” the pulse; otherwise, it would evolve without bound [54]. Pulses that form in both regimes are broadly called “solitons,” a word derived from the solitary, unchanging nature, of the fundamental form of these pulses.

The transition from dispersion-managed to dissipative regime is largely determined by the map strength

$$S = \frac{|\text{GDD}_+ - \text{GDD}_-|}{\min(T_{\text{FWHM}}^2)}, \quad (2.1)$$

where $\text{GDD}_{+/-}$ are lumped positive and negative group delay dispersions, and T_{FWHM} is the full-width-half-max duration of the pulse in the cavity. For net-normal dispersion cavities ($\text{GDD}_+ > \text{GDD}_-$), the oscillator will operate in the dissipative regimes if the map-strength is below a critical value of 3.9 [4].

Conceptually, the map-strength represents the degree to which the pulse duration and bandwidth change as the pulse traverses the cavity. Typically

net-anomalous dispersion is required to balance the normal-like phase imparted by SPM. However, if anomalous dispersion is present in regions where the pulse spectrum is widest, and conversely, normal dispersion is present where the spectrum is narrowest, there can be zero net-phase accumulation in a net-normal dispersion cavity [7]. This is because second-order phase accumulation is proportional to the square of the frequency offset from the center frequency, $(\omega - \omega_0)^2$. Reference Figure 2.2 for an example dispersion map (a piece-wise function of position giving the local dispersion), pulse duration and bandwidth, showing round-trip self consistency in a net-normal cavity.

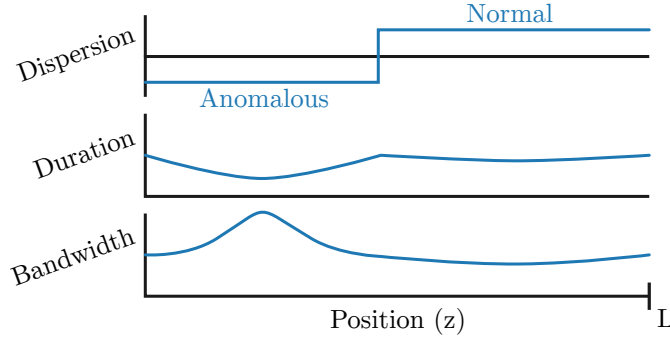


Figure 2.2: Cartoon dispersion map, showing self-consistency of the pulse duration and bandwidth over one round-trip of the oscillator cavity (L is the cavity length). At $z = 0$ the pulse is positively chirped (has a normal-like phase) and enters a region with anomalous dispersion where it begins to de-chirp. As the positive chirp is removed, the pulse duration decreases and nonlinear effects (self phase modulation) lead to broadening of the spectral bandwidth. A minimum pulse duration is reached once the pulse is fully de-chirped. Immediately after this point, the pulse rapidly becomes negatively chirped, due to the large bandwidth. Note that a negative chirp in the presence of self phase modulation leads to a decrease in spectral bandwidth. The same process occurs in the normal dispersion section, but the initial chirp is negative and larger in magnitude. A more detailed description is given by Woodward [54].

2.1.4 Characterization Procedure

The operating regime was classified based on the length and dispersion of elements in the oscillator cavity. A direct measurement of the cavity dispersion, using the method given by Knox [20], supports the length-based estimate of a net-normal cavity. We ran the oscillator in a quasi continuous wave (CW) regime and recorded the frequency shift of the opto-electrically

detected pulse train’s 10th harmonic, as the center wavelength was swept from 1034.94 nm to 1040.56 nm. From these measurements, the group delay and dispersion were extracted. A qualitative comparison between the measured pulse characteristics and those in the literature was also made.

Pulses from the oscillator were characterized temporally with an autocorrelator (Femtochrome FR-103MN) and spectrally with an optical spectrum analyzer (OSA) (HP 70952B); the results are shown in Figure 2.5. Output power was measured with a photodiode based power meter (Coherent OP-2IR). Additionally, a fast photo diode (home-built, ~ 1 GHz bandwidth) was used as an opto-electric converter, allowing us to assess the pulse train stability on an oscilloscope (Agilent DSO5034A) and RF spectrum analyzer (Tektronix RSA3303A).

2.1.5 Characterization Results

Estimate and Qualitative Assessment of the Operating Regime

The main contributions to the dispersion map come from the fiber lengths and intra-cavity grating pair. The fiber section of the cavity is composed of 4.47 m of HI-1060, and 0.30 m of YB-1200-4/125 Yb-doped gain fiber, with a dispersion of $21.1 \text{ fs}^2/\text{mm}$ and $22.3 \text{ fs}^2/\text{mm}$ respectively¹; therefore, $\text{GDD}_+ = 0.101 \text{ ps}^2$. The intra-cavity gratings are separated by 16 mm and have a line density of 10001/mm, providing $\text{GDD}_- = -0.098 \text{ ps}^2$. Given the free space section of the cavity is 0.97 m long, we find the average dispersion, $(\text{GDD}_+ + \text{GDD}_-)/l_{\text{total}}$, is $0.52 \text{ fs}^2/\text{mm}$.

Group delay as a function of CW center frequency is plotted in Figure 2.3. From the slope we extract a net GDD of $(0.016 \pm 0.004) \text{ ps}^2$. Given the total cavity length is 5.71 m, the cavity-averaged dispersion is $(2.8 \pm 0.7) \text{ fs}^2/\text{mm}$. The value is greater than that found based on fiber lengths, perhaps because it includes the dispersion of other optical components (waveplates, isolator, etc.). Regardless, the group delay clearly increases with frequency, reaffirming our cavity has net-normal dispersion.

Ordinarily, we would expect a cavity with a net-normal dispersion of this magnitude to operate in a dispersion-managed soliton regime. However, qualitatively, the pulse spectrum (Figures 2.4, 2.5) has distinct steep edges, and a letter “M” shape, which are characteristic features of a dissipative soliton [38, 53, 54]. Evidently, the map strength is below the critical threshold of 3.9, required for guidance of the pulse evolution by conservative

¹ values provided by a colleague, based on geometric properties available from Nufern and assuming pure silica for the cladding. The expansion wavelength is 1050 nm.

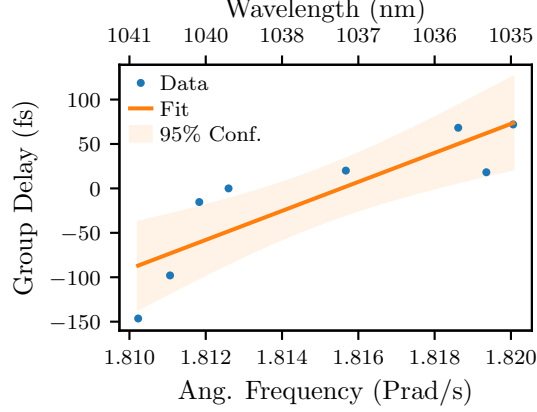


Figure 2.3: Group delay in the oscillator cavity. The grating spacing is 17 mm. The data are plotted with respect to angular frequency, such that a linear fit extracts the cavity GDD $((0.016 \pm 0.004) \text{ ps}^2)$. Based on a total cavity length of 5.71 m, $\beta_2 = (2.8 \pm 0.7) \text{ fs}^2/\text{mm}$. A 95% confidence interval for the fit is shown, assuming the data are normally distributed.

effects. An estimate of the map strength can be had by assuming a Gaussian temporal profile for the emitted pulse. The minimum intra-cavity pulse duration occurs after the grating pair, and is found to be

$$\tau_{\min} = \sqrt{\tau_0^2 - (2\pi\Delta\nu \cdot \text{GDD}_-)^2} = 2.89 \text{ ps}, \quad (2.2)$$

where $\tau_0 = 5.30 \text{ ps}$ is the initial pulse duration (oscillator output), $\Delta\nu = 7.22 \text{ THz}$ is the pulse bandwidth given by the FROG retrieval in Figure 2.4, and GDD_- is the phase imparted by the grating pair (see Appendix A for a derivation). Combined with the value for GDD_+ based on fiber lengths, it follows that the map strength, $S = 0.024$, falls far short of that needed for non-dissipative propagation.

The map strength is small primarily because the pulse is never fully compressed within the cavity ($\tau_{\min}$ is fairly large). Note that outside the cavity we can compress the pulse to 101 fs (Figure 2.5), so we are not limited by bandwidth or higher-order spectral phase. It is possible the lack of intra-cavity compression results from placing the 2m SMF delay line before the gain fiber. In this location, the SMF fiber stretches the pulse before amplification, weakening NPE, and hence, the strength of the effective saturable absorber [8, 23].

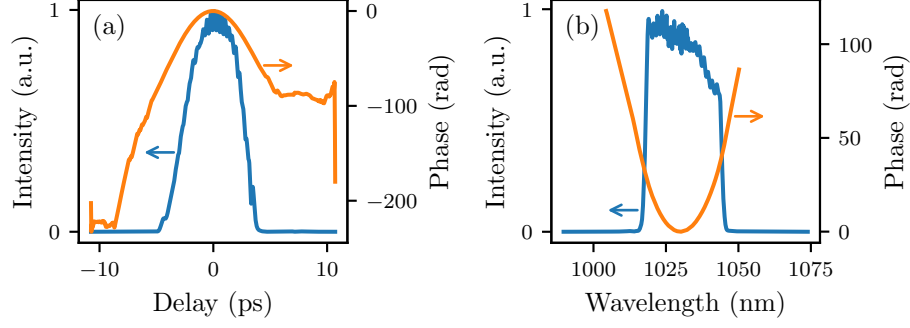


Figure 2.4: FROG retrieval of the uncompressed pulse emitted by the oscillator. (a) Temporal intensity and phase. Intensity FWHM is 5.30 ps. (b) Spectral intensity and phase. Intensity FWHM is 25.9 nm (7.22 THz). The pulse is strongly up-chirped (TBP = 38.82; sign confirmed with an external compressor), which is characteristic of a net-normal dispersion regime.

Typical Pulse Characteristics

Typical operating characteristics for the oscillator are listed in Table 2.1. The oscillator tends to drift with varying environmental conditions, hence, these values are largely a guide for trouble shooting purposes. Figure 2.5 shows a typical pulse spectrum and autocorrelation for the oscillator. Additionally, we have used the measured spectrum—with zero spectral phase—to generate a transform-limited autocorrelation for comparison. Assuming a sech pulse shape, we measure a pulse duration within 27% of the transform limit, which is reasonable for dissipative solitons [52].

Table 2.1: Typical Operating Characteristics

Parameter	Symbol	Test Conditions	Typ.	Unit
Average Power	P_{out}	-	60	mW
Repetition Rate	f_{rep}	-	38.0	MHz
Linewidth	-	1st Harm.	< 10	Hz
Linewidth	-	10th Harm.	< 30	Hz
Sideband Suppression	-	10th Harm.	> 60	dB
Temporal FWHM	τ_{FWHM}	Compressed	100	fs
Spectral Bandwidth	BW	-10 dBm	27	nm

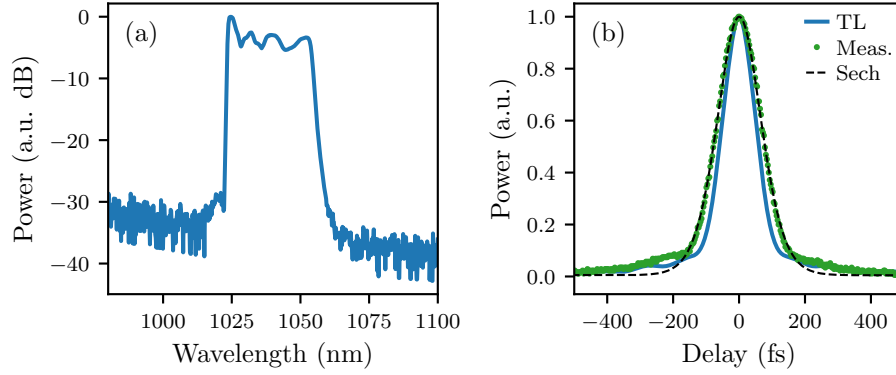


Figure 2.5: Typical spectral and temporal characteristics of the oscillator. (a) Spectrum (blue solid), normalized to peak intensity. (b) Comparison of the measured (green dots, 101 fs FWHM) and transform-limited (blue solid, 74 fs FWHM) auto correlations after the PCF. The duration of the measured AC is found from a sech fit (black dashed).

2.2 Stretcher

2.2.1 Design Principles

Overview

The pulse stretcher imparts a controllable amount of normal dispersion onto the pulse in order to lengthen it in time. This reduces the peak power of the pulse, and thus, the amount of nonlinear phase accumulated in the following amplifier chain. The stretcher is constructed in the “Martinez” configuration [26], which has the benefit that it is an exact counterpart to the “Treacy” grating compressor [40]. In theory, the phase added by the stretcher is entirely removed by the compressor after amplification; in practice, there are residual aberrations induced by the lens in the stretcher that make the process not perfectly reversible, though they can be minimized by using an achromatic lens. The amount of group delay dispersion (GDD), and hence pulse duration, is controlled by the grating pitch, the focal length of the lens, and the relative position between lens and grating. The last parameter can be varied in-situ if the lens is placed on a stage.

Geometry

Before looking at the stretcher as implemented, let us consider an “unfolded” configuration, that clearly demonstrates the principles. In Figure 2.6 the unstretched pulse approaches from the left hand side, and its spectral components are angularly dispersed by the first grating. We may treat each spectral component as a ray with a different angle; here the red and blue rays represent the longest and shortest wavelengths respectively. The unity-magnification telescope is used to image these rays on the second grating, and has two important features: it inverts the angle that incident rays make with the horizontal, and it images the rays on the *reverse* side of the second grating as compared to a Treacy compressor (inset, Figure 2.6). The latter property is made clear if we extend the ray trajectories past the second grating to the focus, where a virtual grating has been placed. If the pulse instead approached from the right, the rays dispersed by the virtual grating would retrace the dotted trajectories and be incident on the opposite side of the second grating (note, to complete the compressor one would need to mirror the roof prism (RP) across the grating).

One can now visually appreciate how the stretcher adds normal dispersion—the red ray takes a shorter path than the blue ray—and how it is the reverse of the Treacy compressor. It is also apparent that the geometry is symmet-

ric between lenses. We can exploit this property to conserve table space by placing a “folding” mirror at the midplane between lenses, effectively halving the length. This folded geometry is the one we have used, and is shown in Figure 2.7.

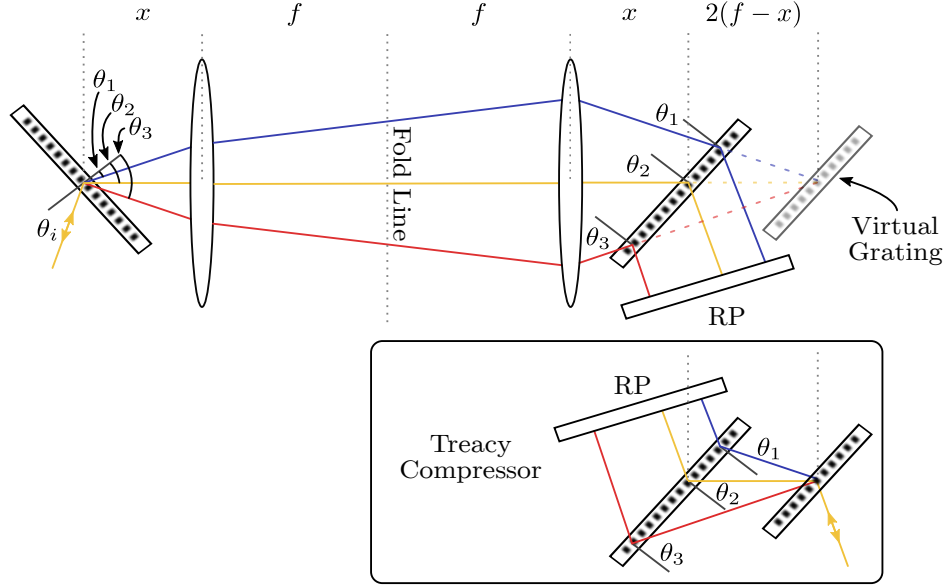


Figure 2.6: Diagram of an unfolded Martinez stretcher (Inset: diagram of a Treacy compressor). RP, roof prism. The unity-magnification telescope preserves ray angles and images the rays on the reverse side of the second grating as compared to a Treacy compressor. The GDD will be equal and opposite to that of the Treacy compressor formed by the second and virtual gratings (redrawn below, with the beam propagating in the opposite direction). Notice the difference in path lengths between red and blue beams is reversed.

2.2.2 Implementation

In the following section refer to Figure 2.7 for component numbering. Our first task is conditioning the beam from the oscillator such that grating (3) operates efficiently. HWP (1) sets the polarization normal to the table, and mirror (2) steers the beam through a gap in roof prism (6) onto grating (3) at an incident angle of 61.1° ; under these nominal conditions, the quoted diffraction efficiency is $> 94\%$ at 1030 nm. The diffraction angles ($\theta_{1-3} \subseteq$

θ_D) are found with the common transmissive-grating equation

$$\sin \theta_D = \sin \theta_i + \frac{m\lambda}{d} \quad (2.3)$$

where θ_i is the incident angle, m is the diffraction order, λ is the wavelength, and d is the grating pitch. Grating (3) has a line density of 17001/mm ($d = 588$ nm), and is optimized for the $m = -1$ order. As is typically the case, θ_D shares the sign convention of refraction, hence, a negative value indicates that incident and diffracted rays are on the same side of the grating normal.

Transmissive volume-phase holographic (VPH) gratings, like grating (3), are most efficient when the angle of incidence and diffraction are equal and opposite. The efficiency peak occurs when rays, reflected from periodic planes with alternating refractive index inside the grating, meet the Bragg condition [2]. The height and width of the efficiency peak are controlled by both the thickness of the refractive layers and the contrast between them; conveniently, the location (wavelength) of the peak can then be adjusted by tuning the incident angle. We use this feature, and adjust the angle of grating (3) to minimize the power in the un-diffracted (zeroth-order) beam transmitted through grating (3).

Lens (4) makes up our unity-magnification telescope when double-passed. It is placed 10 cm after grating (3), and has a focal length of 200 mm. We use an achromatic doublet to reduce the effects of chromatic and spherical aberrations, though their presence is still noticeable when coupling into a collimator (Figure 2.8). It is important to note that within the stretcher spherical aberrations map to chromatic aberrations since the beam is spectrally dispersed across the width of the lens.

The beam is displaced vertically as it passes through the stretcher, so that the input and output can be separated (Figure 2.7 (b)). Taking advantage of folding mirror (5), we offset the center of lens (4) with respect to the incoming beam, causing the reflected beam to emerge on the opposite side of the centerline. Roof prism (6) further displaces the beam from the center of the lens, which allows the final beam to be diverted by pick-off mirror (7).

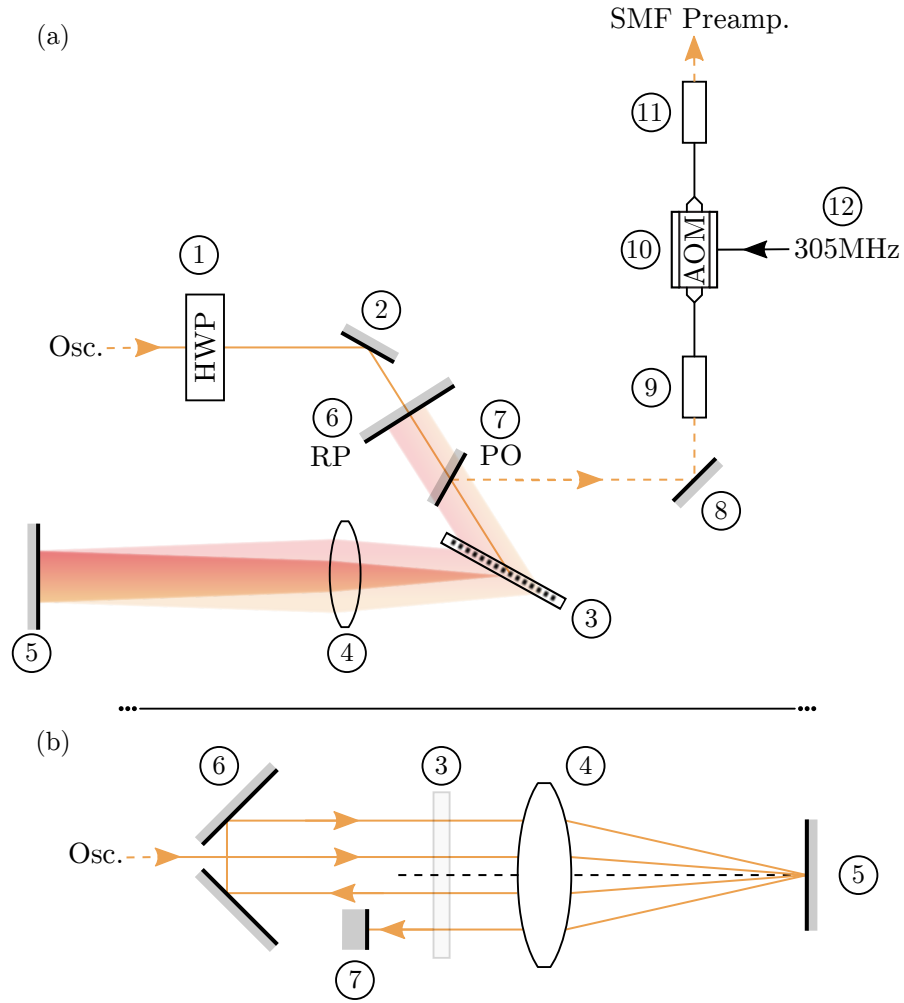


Figure 2.7: Optical layout of the stretcher. HWP, half wave plate; RP, gapped hollow roof prism; PO pick off mirror; AOM acousto-optic modulator. (a) Top down view of stretcher. The pulse picker (AOM) is also shown. (b) Side view showing the vertical displacement of the beam as it traverses the stretcher.

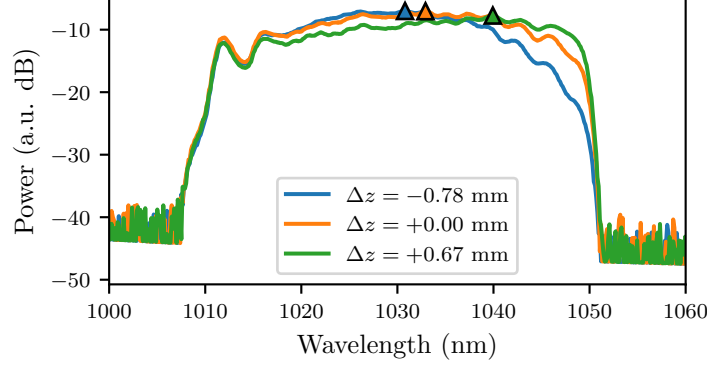


Figure 2.8: Post-stretcher spectra showing spectral filtering when coupled to a collimator. Δz is the displacement of folding mirror (5) (Figure 2.7) from the optimally collimated position, which leads to a divergent or convergent beam depending on the direction. Peak spectral power is indicated with a triangle. A shift in the optimally coupled (peak) wavelength indicates that longitudinal chromatic aberrations are present.

2.2.3 Dispersion

As shown in Section 2.2.1, a Martinez stretcher is essentially the inverse of a Treacy compressor, with a slant separation between gratings of $2(f - x)$. Following the work of Weiner [50], the GDD imparted by the stretcher is

$$\text{GDD}_{\text{Str}} = \psi_2 = L \frac{\partial^2 \beta(\omega)}{\partial \omega^2} = 2(f - x) \frac{m^2 \lambda^3}{\pi c^2 d^2 \cos^2 \theta_D}. \quad (2.4)$$

and the next higher phase component due to third order dispersion (TOD) is

$$\psi_3 = L \frac{\partial^3 \beta(\omega)}{\partial \omega^3} = -2(f - x) \frac{3m^2 \lambda^4}{4\pi^2 c^3 d^2 \cos^2 \theta_D} \left[\frac{1 - \sin \theta_i \sin \theta_D}{\cos^2 \theta_D} \right]. \quad (2.5)$$

Note that ψ_2 and ψ_3 are for the stretcher as implemented, which is often considered to be a “double-pass” of the original Martinez stretcher. The roof prism (RP, Figure 2.6) returns the beam through the stretcher, doubling the accumulated phase and also ensuring the spectral components are spatially recombined.

2.2.4 Pulse Duration

We are unable to measure the stretched pulse duration with locally available equipment; it exceeds the delay range of our autocorrelator and SHG-FROG, and the bandwidth of our fastest photodetector (~ 1 GHz). A reasonably accurate estimate can be had by numerically applying the phase from the stretcher to the measured pulse spectrum. Since dispersion is a linear process, only the initial phase of the pulse will be neglected, and this is relatively small.

Starting with the oscillator spectrum in Figure 2.5, we apply both ψ_2 and ψ_3 , then take the Fourier transform to find the stretched pulse envelope (Figure 2.9-e). For reference, we also use the zero-phase spectrum to retrieve the transform limited (TL) pulse (Figure 2.9-c). Note that the numerically stretched pulse no longer has a Gaussian-like profile. In fact, the temporal profile is now shaped like the spectrum, with time zero corresponding to $\lambda = 1030$ nm, since we defined this as the “central” wavelength when calculating the dispersion. The stretched pulse duration is 439 ps, far greater than that of the transform limited pulse (74 fs).

2.2.5 Simulation of Stretched/Unstretched Pulse Propagation

We were interested to see the benefits of pulse-stretching, so for both the unstretched and stretched pulses, we simulated the propagation dynamics in 20 cm of PM-980 fiber, which makes up the collimator pigtails on the AOM. A fourth-order Runge-Kutta in the interaction picture method was used to solve the generalized nonlinear Schrödinger equation (GNLSE) [17] (this method was implemented in Python by Edmund Kelleher, Technical University of Denmark). The temporal and spectral evolution of the unstretched pulse are shown in Figure 2.9-a/b respectively. A line-plot of the final spectrum for the unstretched and stretched pulses is shown in Figure 2.9-d/f respectively. We have used the dispersive characteristics of Nufern PM-980-XP ($\beta_2 = 21.9 \text{ fs}^2/\text{mm}$, $\beta_3 = 34.1 \text{ fs}^3/\text{mm}$, $\beta_4 = -29.2 \text{ fs}^4/\text{mm}$)², and a typical value of $n_2 = 2.96 \times 10^{-20} \text{ m}^2/\text{W}$ for the nonlinear refractive index of single mode fiber [19].

Even at a low average power of 20 mW, the unstretched pulse has sufficient peak-power for SPM, leading to spectral broadening in the first ~ 7 cm (Figure 2.9-b). The SPM process is quenched by simultaneous temporal

²values provided by a colleague, based on geometric properties available from Nufern and assuming pure silica for the cladding. The expansion wavelength is 1050 nm.

broadening of the pulse (Figure 2.9-a) due to the dispersive nature of the fiber, and so, the spectrum is largely unchanged in the remaining length.

The natural broadening, and quenching of SPM, alludes to why the dynamics of the stretched pulse are quite different. By initially introducing a large amount of dispersion, we immediately lower the peak-power, avoiding spectral broadening in the first section of fiber. The result is that the spectrum of the stretched pulse is essentially invariant over the length of the fiber (Figure 2.9-f), aside from a phase due to material and waveguide dispersion.

2.2.6 Consequences of Not Stretching Before Amplification

As we have seen, nonlinear broadening occurs even at the low powers emitted from the oscillator. Amplification of the pulse worsens the problem, and the effects can become quite dramatic. In our case, the decision to use a pulse stretcher came from necessity, as we were no longer able to compress the pulse to a reasonable form after amplification. In Figure 2.10 we compare running the system with and without the stretcher, at an average power of 25 W and a repetition rate of 20 MHz. With stretching (Figure 2.6-a), the compressed pulse duration after amplification is 186 fs. Assuming a sech pulse shape, the peak-power would be 5.9 MW—over three orders of magnitude greater than out of the oscillator. Inevitably, if the pulse is unstretched there is strong SPM throughout the amplification process, leading to spectral broadening (Figure 2.10-d). Deep modulations develop in the spectrum where two spectral components at the same frequency are π rad out of phase. The irregularity of the modulations suggests the same of the spectral phase. Compressing the unstretched pulse after amplification reveals the irregular phase in a different manner; this time via “wings” on either side of the peak in the autocorrelation (Figure 2.10-c). The wings indicate the spectral phase is substantially non-parabolic (not pure GDD), since it cannot be eliminated by the compressor.

To some extent, it is interesting that SPM leads to an incompressible spectral phase, when it is the main phenomenon used in spectral broadening and compression schemes [41]. Invariably, the generation of higher-order phase components through pure SPM results from an intensity envelope that deviates from a parabolic form. This can be shown by solving the pulse propagation equation in the absence of dispersion and loss,

$$\frac{\partial E}{\partial z} = i\gamma|E|^2E, \quad (2.6)$$

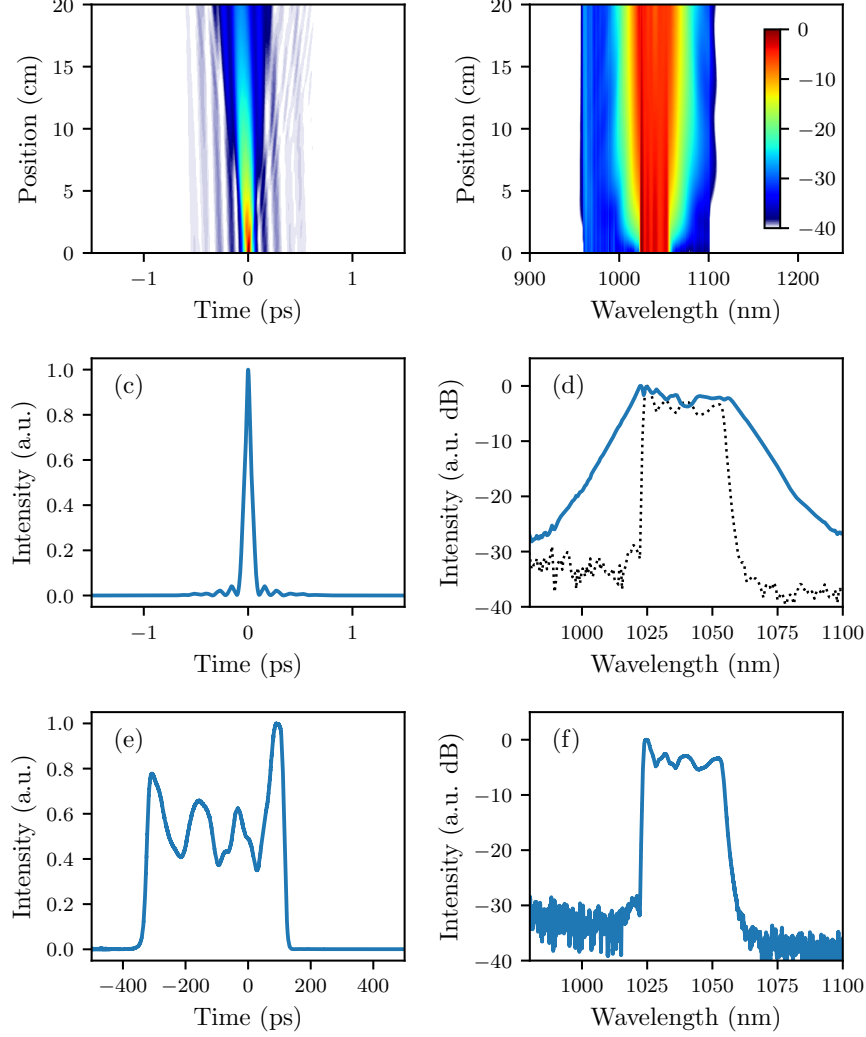


Figure 2.9: Comparison of stretched and unstretched pulse propagation dynamics in 20 cm of PM-980 fiber. (a)/(b) Simulated temporal/spectral evolution of the *unstretched* pulse with propagation distance. Spectral colorbar is logscale, referenced to the peak spectral intensity. SPM occurs in the first ~ 7 cm where there is sufficient peak-power, before the pulse disperses in time. (c) Initial temporal envelope of the unstretched pulse, assuming the transform limit of the initial (measured) spectral magnitude (74 fs FWHM; 5.8 kW pk-pwr). (d) Final spectrum of the unstretched pulse is significantly broadened (dotted initial spectrum for reference). (e) Initial stretched pulse envelope takes on the shape of the spectrum, as spectral components are separated in time (439 ps FWHM; 2.0 W pk-pwr). (f) Final spectrum of the stretched pulse shows little broadening, as the peak-power is three orders of magnitude lower.

where $E(z, t)$ is the complex amplitude (electric field envelope), and γ is a nonlinear parameter [1]. The solution

$$E(z, t) = E(z = 0, t) \exp\left(i\gamma|E(z = 0, t)|^2 z\right) \quad (2.7)$$

reveals that the pulse envelope is static with propagation, but accumulates a temporal phase in proportion to its intensity. For a parabolic intensity envelope, $|E(z = 0, t)|^2 \propto -t^2$. In this case, the phase is parabolic and the chirp is purely linear. Often, pulse intensity envelopes are approximately parabolic about their peaks (Figure 2.10-a), and in this region the nearly parabolic phase can be removed with a grating compressor. This is perhaps why the unstretched pulse (Figure 2.10-c) has a nicely preserved peak, while its flanks are buried in out-of-phase spectral components.

Of course, this discussion has neglected the interplay between SPM and dispersion, which occur simultaneously in a fiber. When dispersion is present, the pulse envelope no longer remains invariant. Take, for example, the stretched pulse in Figure 2.10-e, which acquires an intensity envelope that mimics the spectral one. This transition from a smooth (somewhat parabolic) to structured envelope shows how dispersion can indirectly lead to the generation of higher order phase through SPM.

In general then, it is best to avoid these issues by lowering the peak power, as is done by pre-stretching the pulse.

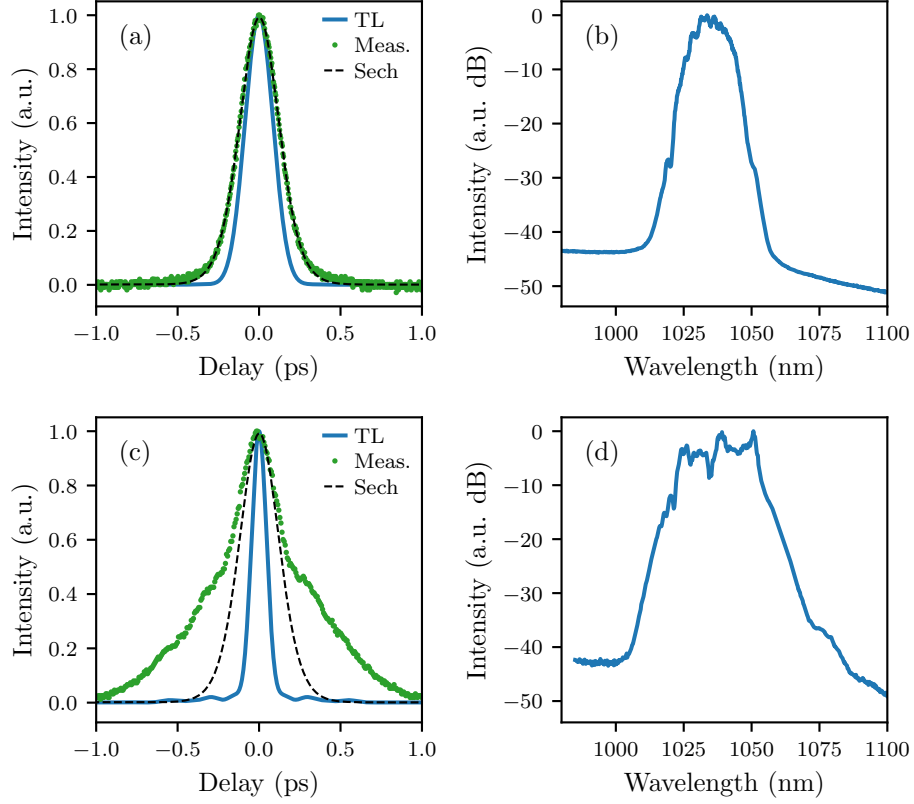


Figure 2.10: Pulse autocorrelations and spectra after amplification, with and without stretching beforehand. (a) Autocorrelation of the stretched pulse (green points). The dashed line is a fit to the data assuming a sech pulse shape (186 fs FWHM). An autocorrelation of the transform limited pulse is shown for reference (blue; 139 fs FWHM). (b) spectrum of the stretched pulse, showing a relatively smooth profile. (c) Autocorrelation of the unstretched pulse (green points) has large “wings” due to newly generated spectral components with higher-order phase. A fit of the *stretched* pulse autocorrelation is superimposed for reference (black dashed). An autocorrelation of the transform limited pulse is also shown (blue; 75 fs FWHM). (d) Spectrum of the unstretched pulse is broader and strongly modulated.

2.3 Acousto-Optic Modulator (AOM)

2.3.1 Design Principles

To adjust the repetition rate of the laser, we use an acousto-optic modulator (AOM) to “pick” every N^{th} pulse. This allows integer division of the fundamental frequency down to arbitrarily low frequencies, limited in practice by excessive amplified spontaneous emission (ASE) and spontaneous giant-pulse formation in the Yb amplifiers. A pulse is picked by momentarily applying a radio frequency (RF) burst to the device, allowing the pulse to pass through; otherwise, the pulse is blocked.

Abstractly, an AOM acts like an optical switch. Light is diverted inside the device by an effective diffraction grating, formed by the periodic compression and rarefaction of a medium when driven by a piezo-electric actuator. The change in refractive index associated with the acoustic pressure wave leads to Bragg diffraction of incident light, just like that in the VPH gratings discussed earlier (Section 2.2.2). The AOM is designed in such a way that only the diffracted beam passes the device; consequently, momentarily driving the piezo actuator allows one to pick a single pulse from a train of pulses.

2.3.2 Implementation and Characteristics

The AOM (10) is shown in Figure 2.7. Two fiber collimators (9 & 11) are used to couple light into and out of the device. We drive the AOM with RF bursts having a carrier frequency of 380 MHz, derived from the 10th harmonic of the photo-detected pulse train emitted by the oscillator (Figure 2.11). The gate signal that modulates the carrier (i.e. defines the duration of the burst) is derived from the first harmonic of the same photo-detected pulse train. By sourcing both signals from the oscillator, we ensure a constant phase-lock between the RF drive signal and the optical pulse traversing the AOM. A fixed phase is desirable, as it eliminates any potential jitter induced by carrier-envelope phase drift in the RF pulse, if the carrier was externally sourced.

Additionally, automatic gain control (AGC) is used to buffer the photo-detected signal, which inherits amplitude fluctuations of the oscillator. The input signal level after AGC is set such that the average power driving the AOM is nominally 2 W in the continuously on state ($N=1$). We maintain this output power to $\pm 15\%$ for input signal levels from -14 dBm to 14 dBm (Figure 2.12-a). Configured in this manner, the AGC protects the AOM

from being overdriven in the event of major oscillator power fluctuations, and maintains a consistent drive power, and hence diffraction efficiency.

Initially, we used the AGC to target a fixed 2 W average power for varying N (up to a 5 W instantaneous envelope power). This approach is advantageous as it maximizes the diffraction efficiency. Unfortunately, the AGC is driven into an indeterminate state when the main RF amplifier is off, or $N > 2$, since the instantaneous power is no longer sufficient to maintain the requisite average power. A comparison of the two AGC implementations is shown in Figure 2.11.

The rise and fall times of the RF burst are 7.8 ns and 9.4 ns respectively (Figure 2.12-b), measured with a 350 MHz bandwidth oscilloscope. Accounting for the rise time of the oscilloscope ($0.35/350 \times 10^6 \text{ Hz} = 1 \text{ ns}$), and that of the AOM (6 ns), we find the rise and fall times of the overall optical gate are 9.8 ns and 11.1 ns respectively. As the optical pulses emitted by the oscillator have sub-nanosecond duration, we have nearly full strength gating for repetition rates below 47 MHz; the chosen repetition rate of 38.0 MHz provides a margin of error.

The width (18 ns), and phase (81.5 ns) with respect to the optical pulse train, of the bursts is controlled by the pulse picker (AlphaNov “Pulse Picker”; see Figure 4.2 for location). The device is clocked by the photo-detected pulse train from the oscillator and generates a rectangular pulse (the gate signal) of controllable width and delay with respect to the clock, at $1/N$ the clock frequency.

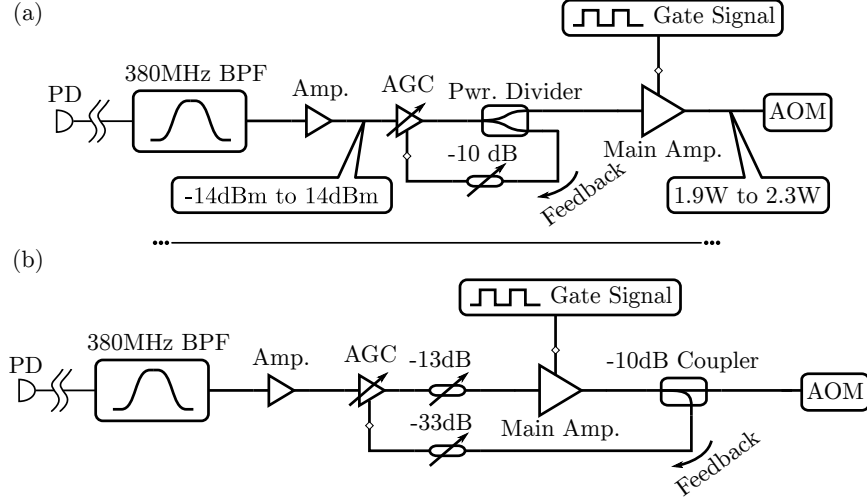


Figure 2.11: Schematic of the AGC and AOM drive electronics. (a) Current implementation. AGC output is used for feedback. The 10 dB attenuator sets the AGC output amplitude, such that the main amplifier is limited to a 2 W peak output power. For an input range between -14 dBm and 14 dBm the output varies between 1.9 W and 2.3 W. (b) Previous implementation. Output from the main amplifier is used for feedback, with the 33 dB attenuator selected to limit the average power to 2 W. The 13 dB attenuator prevents overloading of the main amplifier input.

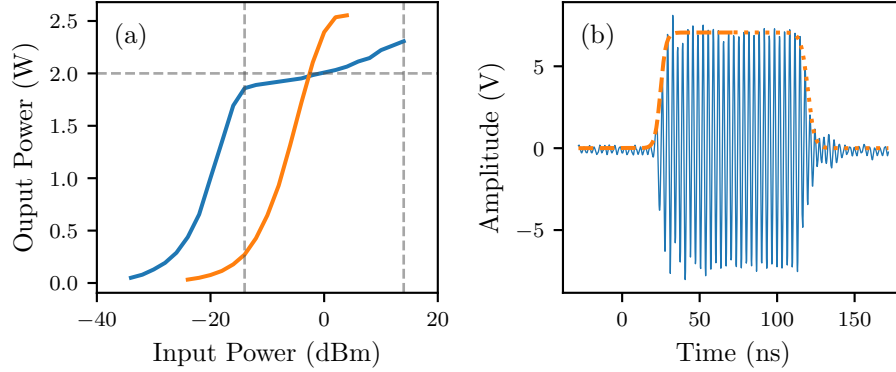


Figure 2.12: Characteristics of the AOM drive electronics. (a) Main RF amplifier gain. The output power with AGC (blue) is leveled to $\pm 15\%$ of 2 W, for input powers from -14 dBm to 14 dBm (dashed vertical lines). Output power without AGC is shown for comparison (orange); input saturates at an output power just above 2.5 W. (b) RF burst from the main amplifier. Rise and fall times are 7.8 ns and 9.4 ns respectively, given 10% and 90% threshold levels.

2.4 Single Mode Fiber Preamplifier

2.4.1 Design Principles

A single mode fiber (SMF) preamplifier is the first stage of amplification after the oscillator. The Yb1200-4/125 fiber can be spliced directly to HI-1060 fiber, making it easy to integrate with commonly available telecom products. The core absorption is very high at 976 nm, nominally 1200 dB/m, which allows short lengths to be used, reducing nonlinearities and material dispersion.

Signal power after the SMF preamp is also of particular importance, as too low a value can allow spontaneous giant pulse formation in the following amplifiers. A minimum average seed power has not been established, we have avoided average powers below 50 mW to be cautious. Given this restriction, it is desirable to run the SMF preamplifier in a way that minimizes output power variation as the pulse-picker division factor N is changed.

2.4.2 Implementation

The schematic for the SMF preamplifier is shown in Figure 2.13. The now pulse-picked signal from the AOM is coupled into the SMF preamp via collimator (3), and steering mirrors (1) and (2). WDM (4) combines the signal with pump light from a laser diode (13). Isolator (12) protects the diode from back-reflections. The signal is then amplified in a 35 cm section of Yb-doped gain fiber. Collimator (6) re-launches the beam into free-space. QWP (7) and HWP (8) are used to maximize the power transmitted by isolator (9), which is rotated such that the passed beam is polarized horizontally with respect to the table. Vertical polarization ensures the majority of the signal is either S or P polarized when reflecting off the following dielectric mirrors, preserving the polarization extinction ratio (PER). Mirrors (11) and (10) guide the beam to the PCF preamplifier. The redundant pick-off mirror (10) is only in place for historical reasons, and should be removed in the future.

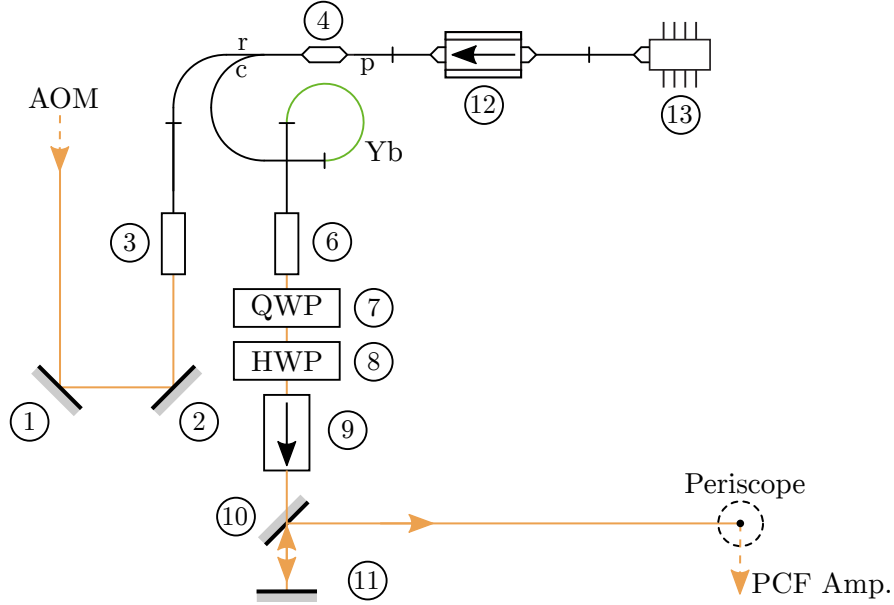


Figure 2.13: Schematic of the SMF preamp. HWP, half wave plate; QWP, quarter wave plate; r,c,p are the reflect (HI1060, 80 cm), common (HI1060, 80 cm), and pass ports of wavelength division multiplexer (WDM) (4). The collimators have 50 cm HI1060 pigtails, and we use a 35 cm length of Yb1200-4/125 gain fiber.

2.4.3 Characterization Procedure

The signal power and spectrum after the SMF preamplifier were measured over a range of pulse picker division factors and pump currents. Initially we tested pump currents from 100 mA to 400 mA and division factors from $N = 1$ to $N = 60$ (Figures 2.14-a/b). As the design evolved, it became clear that operating at higher pump currents is desirable. We extended the power measurements for the $N = 1$ and $N = 60$ case up to 850 mA (Figure 2.14-c). At 850 mA we also recorded the spectrum for division factors $N \in \{1 \text{ to } 8 \text{ by } 1, 10, 15, 19, 38, \infty\}$, where $N = \infty$ corresponds to blocking the input (Figure 2.14-d).

In addition to the single preamplifier, we characterized the signal power and spectrum for two of the same preamplifiers in series. Only the spectral data is provided here, as it contains the most pertinent information (Figure 2.15).

2.4.4 Characterization Results

Single Preamplifier

It can be seen in Figure 2.14-b that higher pump powers lead to proportionally smaller changes in the average output power as N is increased. The decreased variation conveniently coincides with a greater overall power (Figure 2.14-a). It should be noted that the leveling effect realized by increasing the pump power partially results from an increased proportion of ASE (Figure 2.14-c/d), which adds a constant background.

However, ASE is not the only factor at play. In Figure 2.14-c we see that the current threshold (the point after which output power varies linearly with pump current) increases with N . This delayed turn-on is responsible for the initially large $P(N = 1)/P(N = 60)$ ratio. After about 400 mA, the output powers vary linearly with pump current, and the power ratios asymptotically approach the ratio of slopes. Considering that the individual picked pulses are identical, it appears there is a nonlinear effect arising from their repetition rate. In other words, the output power is neither invariant, nor linearly dependent on the rate of pulses. The exact mechanisms for this behavior are complicated and likely follow a transient application of gain saturation; see, for example, Verdeyen [49, Chapter 8]. Suffice to say, it is desirable to operate at pump currents above 400 mA, where even the unseeded (ASE) output is in the saturated (linear) regime.

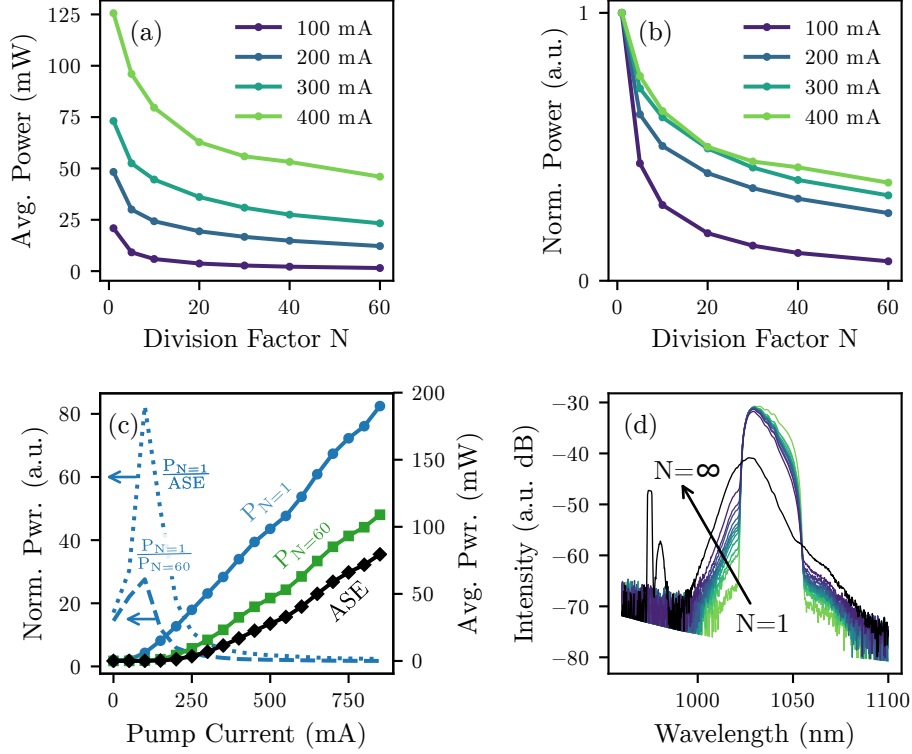


Figure 2.14: Signal characteristics after the SMF preamp for varying pulse-picking division factors N , and pump currents. (a) Average output power increases with pump current, and decreases with division factor. (b) Normalization shows that the average output power varies less with N for larger pump currents. (c) Average output power for $N = 1$ (blue circles) and $N = 60$ (green squares), as a function of pump power. The ASE (black diamonds) also increases with pump current. The proportion of signal to ASE (blue dotted), and the ratio of the powers $P(N = 1)/P(N = 60)$ (blue dashed) decreases with increasing pump power above 100 mA. The input power is 1.8 mW for $N = 1$, and decreases in proportion to N . (d) Spectra for $N \in \{1 \text{ to } 8 \text{ by } 1, 10, 15, 19, 38, \infty\}$. The pulse spectrum is blue-shifted with increasing N , and a larger proportion of the signal is ASE (revealed by blocking the input to the SMF preamp entirely; i.e. $N = \infty$). Here the pump current is 850 mA (spectra courtesy of Evgeny Ostroumov).

Dual Preamplifiers

Early on in the project we tried leveling the output power by using two SMF preamplifiers in series. The idea being that the first amplifier could drive the latter into saturation. The concept appeared to work to some degree; in Figure 2.15 we see that increasing the pump current of the first preamplifier leads to a smaller variation in the combined peak spectral power as the division factor is increased. A detrimental side effect is a modulation of the spectrum for larger values of N . This is likely the result of SPM in the second preamplifier, as the input peak intensity is now higher (bear in mind, there was no pulse stretcher at this point). After trying various permutations of drive currents, we found it was better to increase the pump current of the second preamplifier, and leave the first at 100 mA. As a whole, the dual preamplifier scheme suffers from secondary amplification of ASE emitted by the first preamplifier. In comparison to the spectrum of the single stage preamplifier (Figure 2.14-d), the spectral shoulders are around 10 dB taller relative to the peak, for all division factors. The problem is significantly worsened with increasing pump current. To rectify this issue, one could filter out the ASE between preamplifiers, and potentially see a net reduction in ASE compared to single stage preamplification [31]. Ultimately, the second preamplifier was removed, as a single stage was more convenient and produced less ASE without requiring filtering; though, it may be worth revisiting the dual SMF preamplifier setup with the pulse stretcher now in-place.

2.5 Yb Photonic Crystal Fiber (PCF) Preamplifier

2.5.1 Design Principles

The second stage of amplification consists of a large mode area (LMA) photonic crystal fiber, with a solid Yb-doped core. This mid-step in gain is again required to prevent spontaneous giant pulse formation in the following PC-rod amplifier. The manufacturers of the PC-rod amplifier suggest an input power of 5 W, which is on the order of 50 times larger than what leaves the SMF preamplifier. The mid-step in amplification also lessens gain narrowing, as the gain is distributed between amplifiers [58].

As we saw in Figure 2.15-b, two amplifiers in series can result in nonlinearities. To mitigate this issue, the PCF used has a mode field area (MFA) of $700\text{ }\mu\text{m}^2$, versus the $15\text{ }\mu\text{m}^2$ of our SMF preamplifier. Note that intensity,

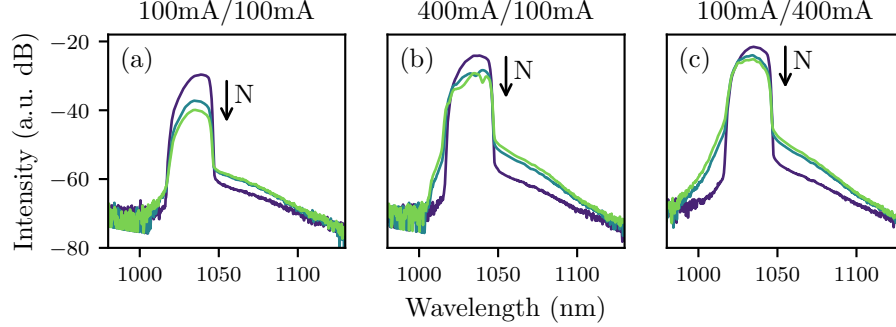


Figure 2.15: Spectral evolution with N , for the combined output of the dual SMF preamplifiers. Arrow indicates increasing $N \in \{1, 20, 40\}$. The pump currents I_1/I_2 are for the first and second preamplifier respectively. (a) Large variation in the peak spectral power at low pump currents. (b) Increasing I_1 aids in leveling the output power, but spectral modulations due to nonlinearities are evident. (c) Increasing I_2 has an improved leveling effect, with reduced nonlinearity.

$I(z)$, is proportional to the reciprocal of MFA, so it is substantially reduced in this PCF.

The benefit of PCF over regular step-index fiber is that for the same MFA the PCF guides in a single mode. Additionally, this particular PCF has integrated stress rods that induce birefringence, making the fiber polarization maintaining (PM). Along with eliminating NPE and other coherent polarization changes, the stress rods introduce a refractive index asymmetry that should reduce ASE, since only the slow-axis is guided for wavelengths below 1200 nm [35] [32].

The PCF preamplifier is backwards pumped, resulting in progressively increasing gain towards the output. This configuration minimizes the accumulated nonlinear phase, or “B-integral”

$$B = \frac{2\pi}{\lambda_0} \int n_2 I(z) dz, \quad (2.8)$$

where n_2 is the nonlinear index, since $I(z)$ is initially small and increases most near the output of the preamplifier. Contrast this with forward pumping, where the gain is greatest at the input, leading to a large pulse intensity over the full length of the preamplifier.

In the following sections we will cover implementation and characterization of the PCF preamplifier.

2.5.2 Implementation

With reference to Figure 2.16 for component numbering, we start by aligning the polarization of the SMF preamp output horizontally to the table with HWP (2). The PCF is designed to coil in a plane parallel to the slow axis, and it is this axis on which we want to launch the input signal (Figure 2.17-a). Minor adjustment of HWP(2) may be necessary to compensate for any rotation of connector (6). Mirrors (3) and (4) orient the beam to be normal with the face of PCF connector (6). Mirror (4) is a short-pass dichroic filter, and serves to both steer the signal and remove any remaining pump light traveling in the opposite direction. Lens (5) ($f = 30$ mm) mode-matches the input signal to the PCF (7), and is fixed in a translating mount with transverse and longitudinal positioners for fine tuning.

Often the alignment procedure is simplified by coupling a small amount of pump light (< 300 mW) into the PCF as a tracer beam, and overlapping it with the signal before mirror (1). This is particularly helpful when the signal power is small, and cannot be distinguished at the output of the PCF. After coarse alignment, one can image the output face of the PCF (Figure 2.17-b) and begin to improve the core/cladding intensity ratio. If the PCF is properly coupled, increasing the pump power will then lead to a rapid increase in this ratio (Figure 2.17-c).

For the pump source, we use an 80 W laser diode (18) centered at 976 nm. Diode (18) is mounted on a thermo electric cooler (TEC) and water cooled chiller plate (19). A controller regulates the temperature of diode (18) by modulating the TEC, both to prevent damage from heating, and to stabilize the wavelength of emission.

Pump light is coupled to the PCF through lens (21) ($f = 25.4$ mm) and lens (9) ($f = 19$ mm), which form a reducing telescope. The 200 μ m diameter core of fiber pig tail (20) is imaged to a 150 μ m spot on the output face of the PCF. Since the PCF has an inner cladding diameter of 200 μ m, it is under-filled by 25%. Under-filling is intentional, as it tends to excite lower order modes that are less susceptible to bend losses, and overlap with the core to a greater extent [3, 46]. Even with under-filling, the pump light will excite some higher-order cladding modes that need to be removed to prevent damage to the outer coating of the PCF. This is done with connector (8), which has an integrated cladding-mode stripper. Connector (8) is mounted in a water cooled copper mount that serves to both dissipate the energy of the cladding-modes, and stabilize the temperature of the output connector, reducing thermally-induced drift in the beam-pointing. Mechanical drawings of the mount can be found in Appendix E.

To collimate the output signal from the PCF, we use the same lens (9) that focuses pump light in the opposite direction. Dichroic mirrors (10) and (11) separate the pump and output, and together with mirror (12), steer the output towards isolator (14). HWP (13) varies the fraction of the output signal that is rejected by isolator (14). The rejected portion is monitored on photo diode (23), and used by the interlock controller to determine the operational status of the PCF preamplifier. HWP(15) rotates the polarization to be vertical to the table, as done after the SMF preamplifier to preserve the PER. The rejected portion is monitored on photo diode (23), and used by the interlock controller to determine the operational status of the PCF preamplifier. HWP(15) rotates the polarization to be vertical to the table, as done after the SMF preamplifier to preserve the PER.

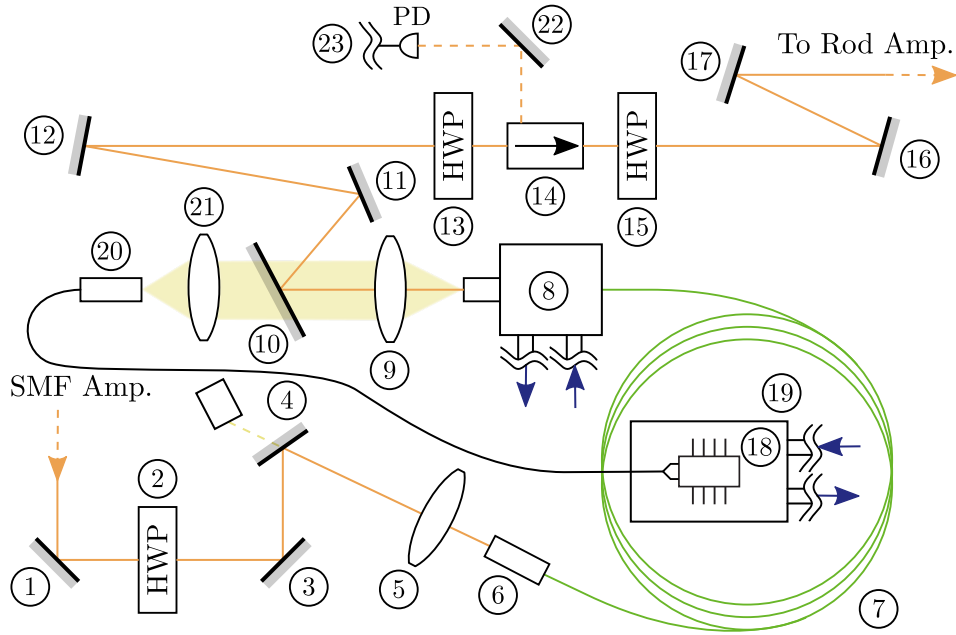


Figure 2.16: Schematic of the PCF preamplifier, and its associated optics and electronics. HWP, half wave plate; PD, photodiode. Blue arrows indicate cooling water.

2.5.3 Mode-Matching

Characterization Procedure

The input beam diameter (from the SMF preamplifier) was measured so that lens (5) could be properly selected for mode-matching to the PCF core. A camera beam profiler and 90-10 knife-edge measurement [45] were used to determine the beam diameter.

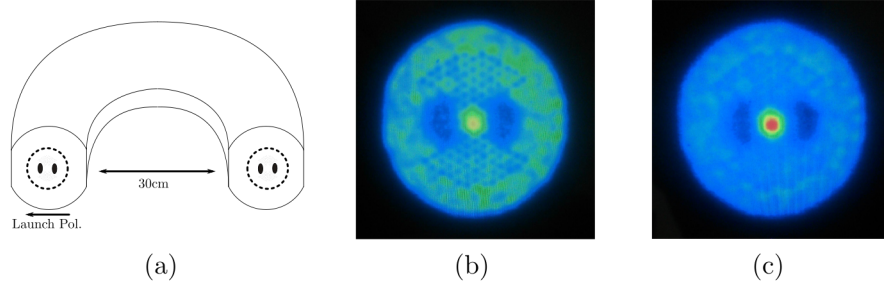


Figure 2.17: Geometry and modal properties of the PCF. (a) “Coil Control” – polished sides on the PCF – ensure that the fiber coils in a single plane, parallel to the slow axis. Stress rods are the two black ovals. See the data sheet [37] for an accurate rendering of the core structure. (b) Imaged output face of the PCF without pumping. Signal is concentrated in the core, however, a significant portion scatters into the cladding. (c) Increasing the pump power strongly promotes the core signal.

Characterization Results

The input beam was found to be slightly elliptical. The intensity $1/e^2$ major/minor axes were 1.54 mm/1.40 mm, as measured with a camera beam-profiler (Coherent LaserCam-HR II), and 1.17 mm/1.09 mm, as found with a 90-10 knife-edge measurement. The mode field diameter (MFD) of the PCF is specified to be 31 μm . Based on these values, the appropriate focal length for mode-matching ranges from 36.4 mm to 25.8 mm. We opted to use a plano-convex lens, with a focal length of 30 mm. Longitudinal translation of this lens will allow some optimization of the coupling, though, a perfect overlap of wavefront curvature and diameter will not be possible. A more sophisticated telescopic focusing arrangement could be used, but we have found the performance of the preamplifier is sufficient for our needs.

2.5.4 Pump Coupling Optics

Characterization Procedure

A Zemax simulation was performed retroactively to determine the expected pump-coupling efficiency. This was spurred by analysis of the all-aspheric pump coupling system for the PC-rod, which revealed the benefits of not using spherical lens (see Appendix B for more details). The same geometric image analysis was performed here.

The object diameter and NA are 200 μm and 0.22 respectively, following

the fiber pigtail specification for the pump diode (nLight pearl P10). The object is collimated with lens (21) (Thorlabs LA1051-B, $f = 25.4$ mm) and focused with lens (9) (Newport KBX043AR.18, $f = 19.0$ mm); reference Figure 2.16 for component numbering. The image aperture and acceptance NA are $200\text{ }\mu\text{m}$ and 0.6 respectively, following the pump cladding specifications for the PCF preamplifier.

The lens positions are optimized for both minimum divergence between lens and maximum coupling efficiency simultaneously. The results are shown in Figure 2.18.

Characterization Results

Like the results for the all spherical coupling system in Appendix B, efficiencies are fairly poor at only 50.0%. The cause can be found at the focal plane (Figure 2.18-a), where spherical aberrations are visible. We therefore suggest changing to aspheric optics, if greater efficiencies are required.

It also should be mentioned that the spacing between lenses was varied from 25 mm to 100 mm with little effect—only a $\pm 1\%$ difference.

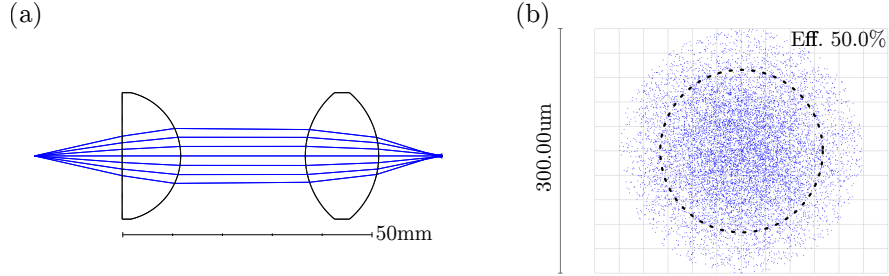


Figure 2.18: Zemax simulation of the PCF preamplifier pump coupling efficiency. Rays are traced left-to-right in (a). The object is a $200\text{ }\mu\text{m}$ circle, with an NA of 0.22, representing the fiber-pigtail core. The image aperture is a $200\text{ }\mu\text{m}$ circle, with an NA of 0.6, representing the PCF pump-cladding. The system is simultaneously optimized for collimation between lenses and maximal coupling efficiency. Note the surface orientations that place the lenses in “near-best-form”. (a) Spherical lens system. Components left-to-right: Thorlabs LA1051-B, ELF of 25.4 mm; Newport KBX043AR.18, EFL of 19.0 mm. (b) Distribution of rays traced to the image plane. The PCF’s $200\text{ }\mu\text{m}$ pump cladding diameter is indicated by a dotted circle, and rays outside a $300\text{ }\mu\text{m}$ diameter are removed.

2.5.5 Output Power

Characterization Procedure

An output power curve for the PCF preamplifier was made by sweeping the pump current from 0.5 A to 3.3 A in 0.1 A steps. For these measurements the repetition rate was 38 MHz, and the input power from the SMF preamplifier was 153 mW. A Coherent LM-10-HTD thermopile head was used.

Characterization Results

The output power curve for the PCF preamplifier is shown in Figure 2.19-a. Optical transparency of the PCF preamplifier—an output power equal to the input power of 153 mW—is reached for a pump power of 5.13 W (1.1 A). For pump powers above 12.9 W, we use a linear fit to find that the slope efficiency is 26%. Extrapolating the fit to intercept the x-axis gives us a threshold pump power of 7.0 W. The maximum output power is 5.1 W, representing a gain of 15 dB. The slope efficiency for the PCF preamplifier is rather poor in comparison to the PC-rod amplifier (Section 2.6), which has a slope efficiency of 75%. The difference likely results from using spherical lenses to coupling the pump-light, as opposed to aspheric lens, like those used for the PC-rod. We have shown in Appendix B that using aspheric lenses, in the case of the PC-rod, can result in a greater than two-fold increase in coupling efficiency. Here the increase could be even greater, as the spherical lens radii are smaller, and the 19 mm focusing lens is bi-convex (the near “best-form” lens for focusing collimated light is plano-convex).

2.5.6 Temporal and Spectral Properties

Characterization Procedure

A portion of the output after the PCF was sampled using the reflection off a Thorlabs PS811 4° round wedge prism. The sampled beam was switched with a flipper mirror between the autocorrelator (Femtochrome FR-103MN), and a fiber-pigtailed collimator connected to the OSA (HP 70952B). A grating compressor was placed before the autocorrelator, and adjusted for minimum spatial chirp using an imaging spectrometer. We used two Wasatch 1250 l/mm gratings separated by 37 cm ($GDD = -5.0 \text{ ps}^2$) to maximally compress the pulse. At the time, our stretcher used the same grating type, with an effective separation of 32 cm ($GDD = -4.3 \text{ ps}^2$).

Characterization Results

The spectrum after the PCF shows significant gain narrowing (Figure 2.19) at a repetition rate of 19 MHz and pump current of 1.5 A (8.7 W). Increasing the pump current to 3.3 A (24.4 W) does not change the spectrum appreciably. Given the significant loss of bandwidth, we expect the maximally compressed pulse duration to have increase from that after the oscillator. The transform limited autocorrelation for the gain narrowed spectrum is compared to what is measured with the Femtochrome autocorrelator after the grating compressor in Figure 2.19-c. The respective pulse durations are 134 fs FWHM, and 194 fs FWHM, where a sech pulse shape has been assumed for the latter measured autocorrelation. As expected, both pulse durations have increased from their corresponding values after the oscillator (Figure 2.5); however, the percent deviation from the transform limit has only increased from 27% to 31%, indicating a similar pulse shape after amplification.

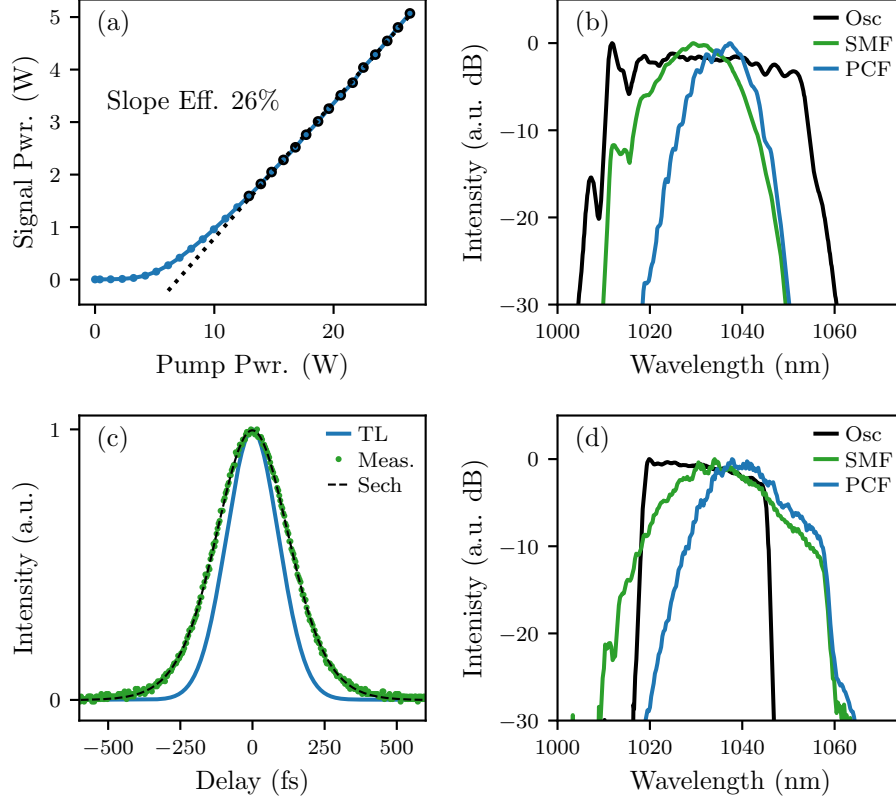


Figure 2.19: Characteristics of the PCF preamplifier. **(a)** Output power of the PCF preamplifier (blue), for an input signal power of 153 mW. The slope efficiency is 26%, given a linear fit (black dots) in the region well past the threshold (black circles). The threshold optical power is 7.0 W, found by extrapolating the linear fit to the x-axis. **(b)** Spectral evolution through the system, normalized to peak intensity. Oscillator spectrum (black) is gain narrowed on the 'blue-side' in the SMF preamp (green), and on the 'red-side' in the PCF preamplifier (blue). **(c)** Comparison of the measured (green dots, 194 fs FWHM) and transform-limited (blue solid, 134 fs FWHM) auto correlations after the PCF. The duration of the measured AC is found from a sech fit (black dashed). **(d)** Spectral evolution through the system before the addition of a stretcher, normalized to peak intensity. Oscillator spectrum (black) is broadened by SPM in the SMF preamp (green), and gain-narrowed on the 'red-side' after the PCF preamplifier (blue).

Reviewing spectral data collected prior to implementing the stretcher (Figure 2.19-d) reveals that the SMF and PCF gain bandwidths overlap to form a bandpass filter. Without the stretcher, SPM in the SMF preamplifier extends the oscillator spectrum to shorter—and more fortuitously—longer wavelengths. After amplification in the PCF preamplifier, the broadened spectrum exhibits the same short-wavelength attenuation seen in Figure 2.19-b, but in contrast, is uniformly amplified at longer wavelengths. Evidently, the PCF gain bandwidth extends to longer wavelengths, preserving the spectral shape at this end. In comparison to the PCF, the SMF preamplifier has a similar, but mirrored effect, attenuating the long-wavelengths. The intersection of these two gain bandwidths gives us the parabolic-looking combination.

2.6 Yb Photonic Crystal Rod Amplifier

2.6.1 Design Principles

The last stage of amplification is carried out in a yet larger mode-area Yb-doped fiber. Here the emphasis is on raising the optical damage threshold, and reducing nonlinearities [30]. The substantial mode field diameter of this fiber (65 μm at a wavelength of 1040 nm) requires a similarly large cladding diameter, making it impractical to coil, so it is packaged as a PC-rod in a water-cooled aluminum jacket. Unlike the PCF preamplifier, the PC-rod does not have stress rods, and is only weakly birefringent. However, as with the PCF preamplifier, the PC-rod is backwards pumped to minimize the accumulated nonlinear phase.

The following are combined in this last stage of amplification: a grating compressor, which removes spectral phase imparted by the stretcher, as well as additional second-order phase due to material dispersion in the amplifier chain; a beam expander and focusing lens that are mode-matched to the HCF in the following spectral broadening stage; and an active stabilization system that is one prong of a two prong solution to maintaining coupling into said HCF. The full system extends over Figures 2.20 to 2.22. We will discuss the physical implementation of all these components simultaneously, however, characterization of the mode-matching optics and active stabilization system will be discussed in separate subsections. The compressor is integral to the temporal characteristics of the amplifier, so they will be discussed together.

2.6.2 Implementation

With reference to Figure 2.20 for component numbering, we lower the beam height, and turn the beam from the PCF preamplifier with periscope (1). The periscope mirrors are set such that the beam leaves at right angles, and is either fully S, or P, polarized. Even so, we have chosen to use silver mirrors to minimize polarization changes due to misalignment. Lens (3) ($f = 250$ mm) and lens (4) ($f = 200$ mm) form a reducing telescope, for mode matching to PC-rod (11). To minimize any potential thermal lensing, lens (3) and (4) are made of fused silica. Mirrors (5) and (6) are used to align the beam such that it is orthogonal to the face of the PC-rod (11). Mirror (7) is a dichroic filter that removes what remains of the counter-propagating pump light passing through water-cooled pinhole (10). We take advantage of the difference in guiding mechanisms to separate the pump and signal with the pinhole: the signal—a guided single-mode—leaves the PC-rod core with an effective NA of only 0.01. In comparison, pump light travels in the multi-mode cladding, which has an NA of 0.5. The large disparity in the respective free-space half-angles (0.5° versus 30°) allows the pinhole to pass the signal while blocking the majority of the unabsorbed pump light. The small signal NA also allows the use of a relatively long focal length lens (9) ($f = 80$ mm) for coupling. This means that even though Lens (9) is aspheric, it likely could be exchanged for a spherical lens without repercussion. A translation stage and mount are used to hold lens (9), allowing three degrees of freedom in positioning the focus.

After emerging from the PC-rod, lens (12) collimates the signal, which then passes through dichroic filter (13). Lens (12) serves the dual purpose of focusing pump light back into the PC-rod. Together, lens (12) ($f = 32$ mm) and lens (16) ($f = 40$ mm) form a reducing telescope that images the $200\text{ }\mu\text{m}$ diameter core of fiber pigtail (17) to a $160\text{ }\mu\text{m}$ spot on the output face of the PC-rod. The pump cladding diameter of the PC-rod is $260\text{ }\mu\text{m}$, and is thus under filled to 60%, falling within the 51% to 64% range suggested by NKT [36]. Though the same de-magnification factor can be achieved with any lens pair having the same ratio of focal lengths, lens (16) is required to have a short focal length in order to collimate the pump light at a diameter that does not clip on dichroic filters (14) and (13). Given fiber pigtail (17) has a NA of 0.22, the collimated beam diameter is 18 mm. Both lens (12) and (16) are mounted with three degrees of freedom. Mirror (14) is an additional dichroic filter used to protect pump diode (18). Unlike for the PCF preamplifier, pump diode (18) is wavelength stabilized by a fiber volume Bragg grating, and is therefore directly mounted to chiller plate (19),

without stringent temperature control. However, it is important to note that the quoted wavelength stability (± 2 nm) only applies to operating currents between 14 A and 19 A, and we typically operate at 7 A (107 W). We did not investigate the wavelength stability of the pump diode at 7 A; however, at this pump current the output power of the rod fluctuated by less than 1% over three hours, indicating the wavelength stability is sufficient.

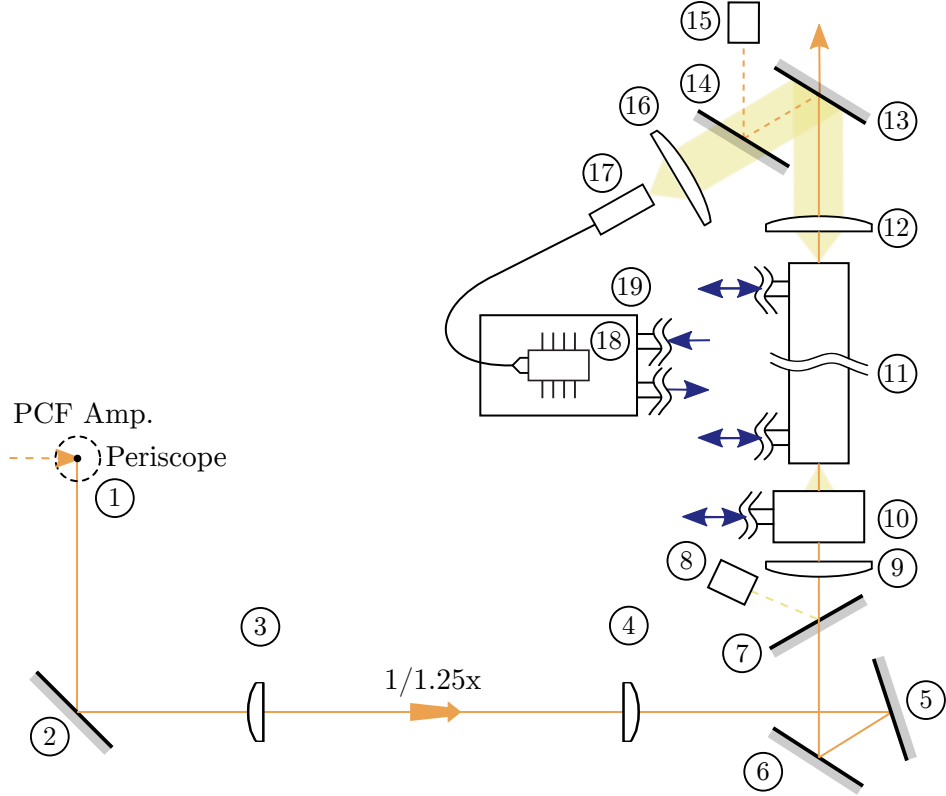


Figure 2.20: Schematic of the PC-rod amplifier. Blue arrows indicate cooling water.

Following collimation and dichroic filtering, the signal after the PC-rod is further conditioned, and temporally compressed. An optical layout is shown in Figure 2.21. With reference to this second schematic, mirrors (1) and (2) steer the signal along a row of holes in the table. They are required as the output face of the PC-rod is angle polished at 2° , refracting the beam at an angle of approximately 0.5° . Lens (3) ($f = -75$ mm) and lens (4) ($f = 300$ mm) form a $4\times$ beam expander to reduce thermal lensing

in the following transmissive components, particularly isolator (6). HWP (5) aligns the signal polarization to pass through the input polarizer of the isolator. By using a HWP, the isolator can be rotated such that the rejected beams are parallel to the table. One of these rejected beams is attenuated with HWP (13) and PBS (14), and sampled by photo diode (15) for use in the interlock. The rejected beam also doubles as a reference for the degree of coupling into the PC-rod. When the PC-rod is mis-coupled, the LP_{11} mode is excited and is predominantly polarized orthogonally to the signal. Thus, reducing the rejected signal power minimizes the LP_{11} component and improves coupling into the fundamental LP_{01} mode. Doing so is critical for achieving highly circular foci when later coupling into the HCF fiber, since the LP_{11} mode tends to make the beam elliptical. Following isolator (6), HWP(7) rotates the polarization vertically to the table, such that the signal is S-polarized with respect to grating pair (10). Mirror (9) adjusts the angle of incidence to be 61.1° on the first grating, which in turn is angled such that the diffracted beam runs along a row of holes in the table. Mirror (11) sends the signal back through the compressor, and is slightly angled so that pick-off mirror (8) can divert the returned signal.

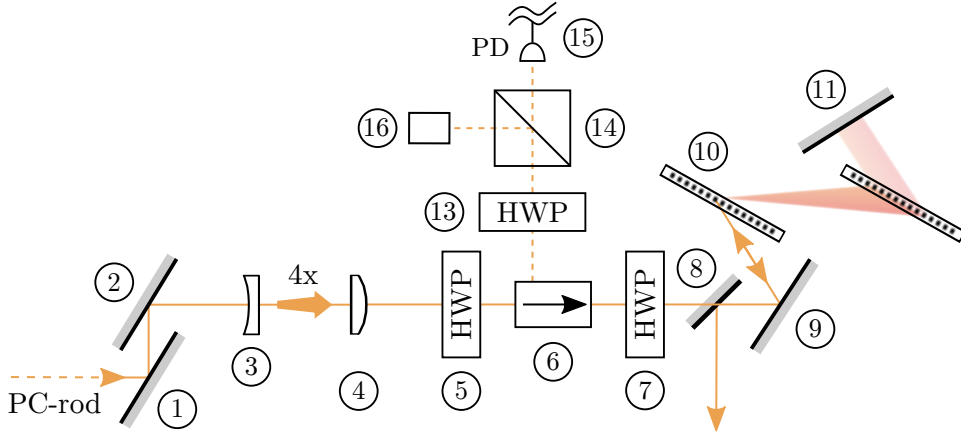


Figure 2.21: Schematic of the beam-conditioning optics and grating compressor after the PC-rod amplifier. From left to right, the beam is leveled, expanded, isolated (and sampled for the interlock), then finally compressed. Half wave plates abbreviated as HWP.

Having been temporally compressed, the last step is stabilizing and focusing the signal into the HCF of the spectral broadening stage. Figure 2.22 continues the schematic after the pick-off mirror, with re-indexed numbering. The signal height is raised with a periscope, which has an active mirror at

the top (PZT1). PZT1 has piezoelectric actuators on both tip/tilt axes, and in combination with PSD1, reduces lateral movement of the beam. PZT2 is back-polished to allow sampling, and mirrors (1) and (2) guide the sampled portion to PSD1. A motion controller attempts to minimize the offset between the beam position and the center of PSD1 by actuating PZT1. It is not possible to change the reference position from the center of the PSD, therefore, mirrors (1) and (2) are adjusted to “zero” the offset. Note that one of the mirrors is redundant, since the PSD is insensitive to beam angle. Mirror (3) simply redirects the beam to the second back-polished mirror (4). Mirror (5) (a 50:50 beam splitter) adjusts the beam position on PSD2. Mirror (6) steers the beam through HWP (7) and thin film polarizer (TFP) (8), which form an attenuator. The attenuated signal is focused into the gas cell, and onto the face of the HCF, with lens (10). The rejected portion of the beam is deposited in beam-dump (9). Also shown are diagnostics for characterizing the stabilization system. Mirrors (11) to (14) form a “lever-arm” that amplifies movement of the beam. Mirror (18) temporarily injects a stable HeNe beam as a reference to check the lever-arm stability. Mirror (15) is back-polished and angled to act as a beam splitter. The divided beam impinges on a third PSD and a Raspberry Pi camera (pi-cam), which are used to monitor fluctuations in the beam position. The pi-cam is largely sensitive to angular beam displacements, since the focusing lens (17) maps beam angle to position on the image sensor. Neutral density filter (16) prevents the image from being saturated.

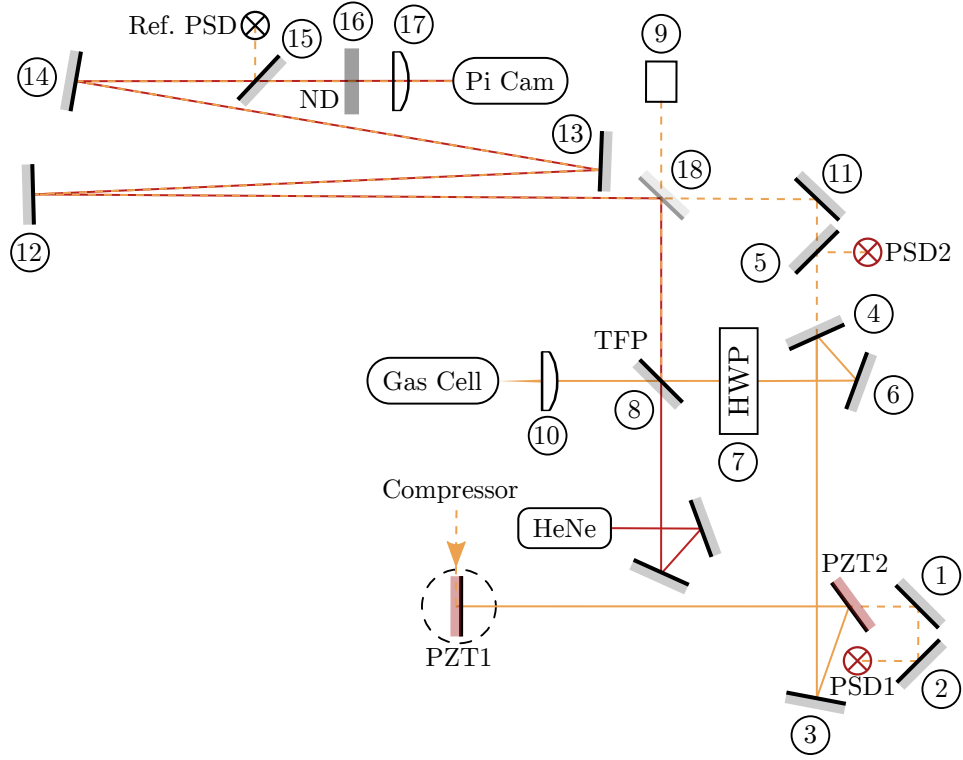


Figure 2.22: Schematic of the active stabilization (AS) system for coupling into the HCF, and the optics used to characterize pointing drift. HWP, half wave plate; PZT piezo-actuated mirror; PSD, position sensitive detector; ND, neutral density filter. The stabilization system consists of two active mirrors (PZT1 and PZT2), two controllers (not shown) and two position sensitive detectors (PSD1 and PSD2). The controllers “walk the beam” with the PZTs until it is centered on both PSDs, which effectively function as a pair of apertures. After AS, the beam is attenuated with HWP (7) and TFP (8), then focused into the HCF with lens (10). Mirrors (11) through (14) amplify beam pointing deviations, which are measured on a third PSD and the Raspberry-Pi based camera. A HeNe laser was used temporarily as a stable reference.

2.6.3 Output Power

Characterization Procedure

An output power curve for the PC-rod amplifier was made by sweeping the pump current from 0.9 A to 8.0 A. The corresponding optical pump power was determined from a linear fit of the diode output power versus current, which we previously characterized (Figure 2.23-b). A Newport 818T-150 thermopile head was used to measure both the output power of the diode and PC-rod. A power-voltage calibration curve for the Newport thermopile was determined with another factory-calibrated Thorlabs S350C thermopile (Figure 2.23-a). Measurement of the PC-rod amplifier output power was carried out at a repetition rate of 38 MHz, and the input power from the PCF preamplifier was 5.65 W (PCF pump current of 3.3 A).

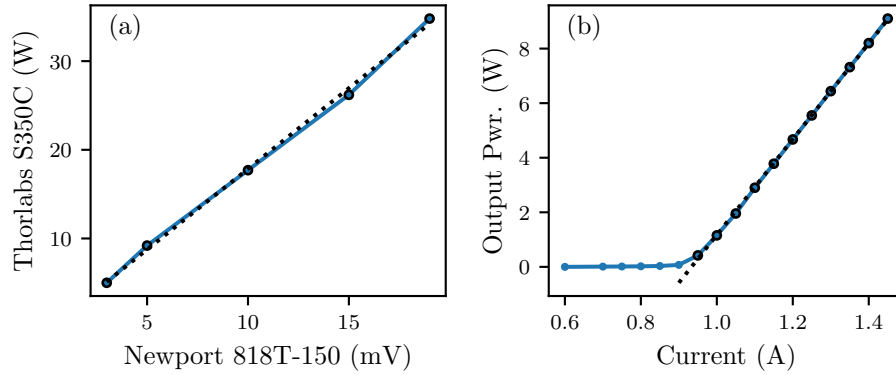


Figure 2.23: Characterization of the nLight element e18.2000976200 diode. (a) Calibration curve for the Newport 818T-150 thermopile. The relationship between optical power and voltage is $P[\text{W}] = 1.82 \times V[\text{mV}] - 0.34$. (b) Output power of the nLight element e18.2000976200 diode used to pump the PC-rod. The relationship between optical power and current in the linear regime is $P[\text{W}] = 17.55 \times I[\text{A}] - 16.38$. The threshold current is 0.93 A.

Characterization Results

The output power curve for the PC-rod amplifier is shown in Figure 2.24-a. Optical transparency is reached at a pump power of 9.95 W (1.5 A). For pump powers above 9.95 W we use a linear fit to find that the slope efficiency is 75%. Extrapolating the fit to intercept the x-axis gives us a

threshold pump power of 57.3 W (at 2.3 A). The maximum output power is 50.6 W (pump power/current of 124 W/8 A), representing a gain of 9.52 dB.

2.6.4 Temporal and Spectral Properties

Characterization Procedure

Evolution of the pulse spectrum was mapped up to and after the PC-rod, in order to assess the effects of gain narrowing. Each amplifier was driven at its typical operational current (SMF preamplifier, 850 mA; PCF preamplifier, 3.3 A; PC-rod amplifier, 7 A). The pulse picker was set in the 'High' gating mode ($N = 1$) to reduce the per-pulse energy, and thus minimize nonlinear broadening through SPM.

In a separate test, at the same operating currents, the repetition rate was lowered to determine if larger pulse energies would lead to excessive nonlinear broadening and higher-order spectral phase through SPM. Both an autocorrelation and spectrum were recorded after the pulse compressor, comprising two gratings with identical specifications to that in the stretcher, and adjusted to minimize the pulse duration.

Finally, the pulse picker division factor was increased to $N = 8$ and the pulse was retrieved with SHG-FROG.

Characterization Results

As noted previously, there is significant gain narrowing leading up to the PC-rod amplifier. In Figure 2.24-b attenuation of the short-wavelengths is evident. At the other end, the spectrum from the oscillator steeply decays just below 1050 nm, hiding the limited gain bandwidth of the SMF preamplifier; though fortuitously, the oscillator spectrum is "gain shaped" such that the spectrum of the SMF preamplifier is fairly rectangular, but peaked where the PCF amplifiers have limited gain. The resulting spectrum out of the PC-rod is therefore flatter than it would be for the oscillator spectrum in Figure 2.19-b. This suggests that further improvements could be made by adjusting the oscillator, or by using a spatial light modulator in the stretcher [21]. We also note that gain narrowing within the PC-rod itself is minimal, with the spectrum closely overlapping that from the PCF preamplifier at long wavelengths, and narrowed by only one or two nanometers at short wavelengths. However, the top of the spectrum does become more structured.

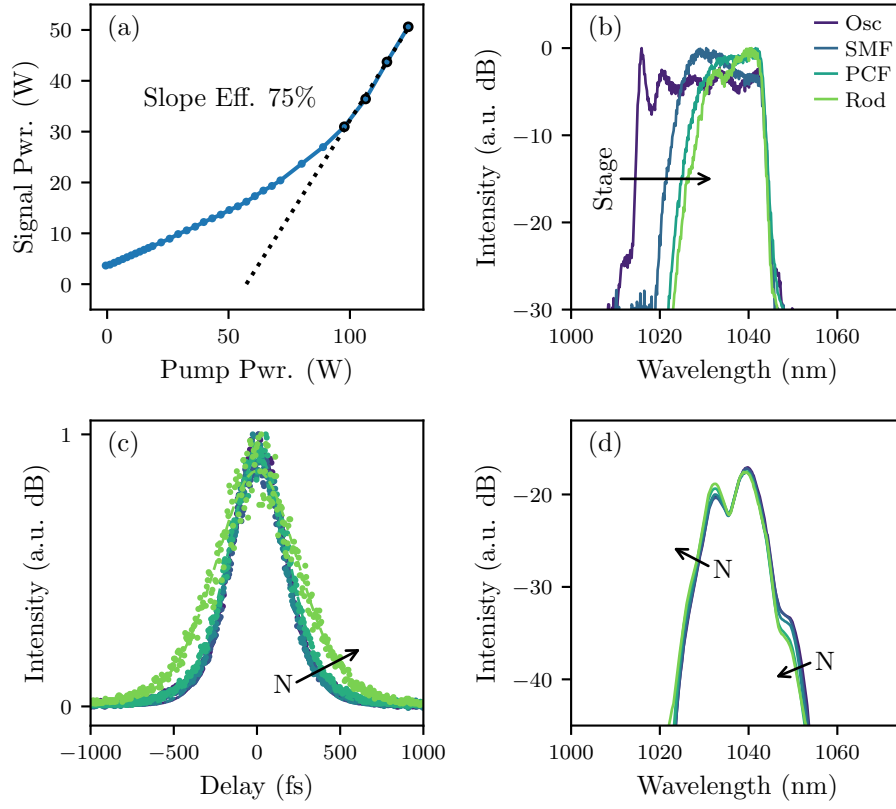


Figure 2.24: PC-rod characteristics. (a) Output power of the PC-rod amplifier (blue solid), with the PCF preamplifier running at 3.3 A (5.65 W). Slope efficiency is 75%, given a linear fit (black dotted) in the region well past the threshold (black circles). The threshold optical power is 57.3 W, found by extrapolating the linear fit to the x-axis. (b) Spectral evolution at 38 MHz showing gain narrowing after each stage of amplification, normalized to peak intensity. Arrow indicates the order of progression [oscillator $\rightarrow$ SMF preamplifier (850 mA) $\rightarrow$ PCF preamplifier (3.3 A) $\rightarrow$ PC-rod amplifier (7 A)]. (c) Autocorrelation of the maximally compressed pulse after the PC-rod, as the division factor N is varied (arrow indicates increasing $N \in \{2, 3, 8, 19, 38\}$). The pulse duration is found from a sech fit, and increases from 252 fs FWHM to 383 fs FWHM with N . (d) Evolution of the spectrum after the PC-rod, as the division factor N is varied (arrow indicates increasing $N \in \{2, 3, 8, 19, 38\}$).

Varying the repetition rate of the pulse train has a tolerable effect on the pulse duration after the PC-rod. The maximally compressed pulse width gradually increases with pulse picker division factor, and there are no major structural changes, such as the development of a pedestal (Figure 2.24-c). One undesirable artifact is an increase in autocorrelation intensity noise with division factor. The source of the noise was not investigated, though its magnitude seems to be correlated to the amplitude of the SHG signal.

The spectrum also shows little change with division factor, aside from a slight blue-shift, and peaking around 1030 nm (Figure 2.24-d). One interesting aspect is the peak at ~ 1042 nm, which is relatively unchanged. The peaks at 1030 nm and 1042 nm generally correspond to the gain peaks of the SMF preamplifier, and PCF pre/main amplifiers, respectively. Thus, peaking at 1030 nm suggests the slight blue-shifting originates in the SMF preamplifier; an issue we previously noted.

To determine the spectral and temporal phase of the pulse, an SHG FROG retrieval was taken after the PC-rod (Figure 2.25). The measured and retrieved SHG-FROG spectrograms are shown in Figure 2.25-a/b respectively. Both the broader shape and fine structure of the retrieval closely resemble the measurement, and the FROG trace error (G) is only 0.5%. This error is the RMS difference between the measured and retrieved spectrograms, divided by the number of data points. Typically, the FROG trace error should be below 1-2%; errors up to 5% can be acceptable for determining the major features of the pulse, if the retrieved and measured FROG traces are closely matched in both general shape and fine structure [34]. In Figure 2.25-c, the retrieved delay marginals (summation over frequency of the retrieved spectrogram) are compared to a conventional SHG intensity autocorrelation. The agreement is excellent, supporting the validity of the retrieval. Our second point of reference is comparing the retrieved spectral intensity to that measured with an optical spectrum analyzer (Figure 2.25-d). Here the overlap is not perfect, though the retrieval mirrors many of the qualitative features of the OSA spectrum, particularly the structure on the short-wavelength side. The most obvious difference is that the FROG spectrum does not cover the longest wavelengths captured by the OSA. Aside from technical errors (misalignment, etc...), or limitations (crystal bandwidth, etc...), the missing spectral components could be incoherent with the amplified pulse. For example, ASE integrated over a pulse period may be substantial, but optical gating would largely eliminate this low peak-intensity background.

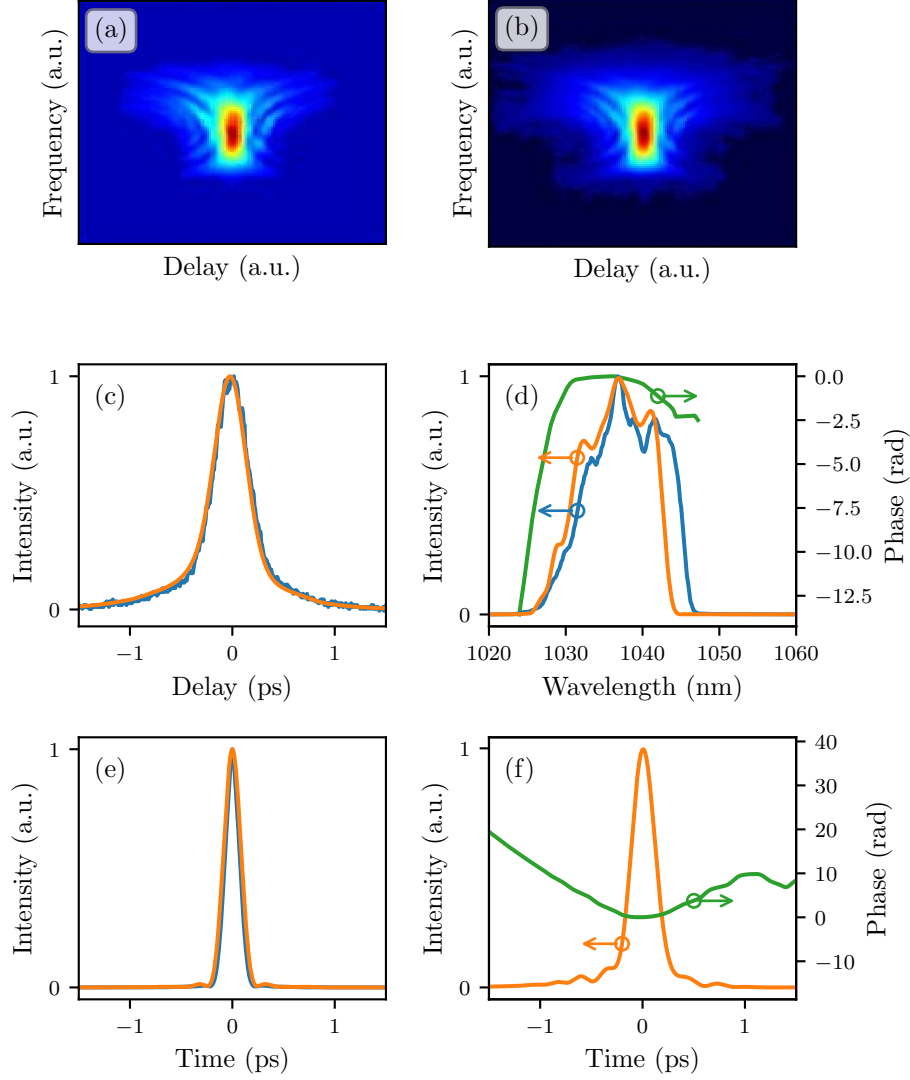


Figure 2.25: SHG-FROG retrieval of the pulse after the PC-rod amplifier. (a)/(b) measured/retrieved SHG-FROG spectrograms. (c) Comparison of the SHG-FROG delay marginals (orange) and an independent SHG based intensity autocorrelation (blue). (d) Retrieved spectrum (orange) and phase (green), compared with the spectral intensity as measured by an OSA (blue). (e) Transform limited pulses for the retrieved (orange) and OSA measured (blue) spectral intensities; the FWHM durations are 183 fs and 159 fs respectively. (f) Retrieved temporal intensity (orange) and phase (green); the FWHM duration is 264 fs.

Nonetheless, assuming the missing spectral components are coherent, their significance on the transform limited pulse duration is investigated in Figure 2.25-e. As expected, the TL duration of the OSA spectrum is shorter, 159 fs versus 183 fs for the FROG spectrum, though generally the two pulses are comparable.

With confidence in the retrieval established, we note that the spectral phase of the pulse is fairly flat, deviating by less than π rad over the FWHM bandwidth. There appears to be a significant second order component, indicating the pulse could possibly be compressed further; though, the resulting addition of third order phase may negate any benefit, as the gratings were adjusted to minimize the autocorrelation duration and pedestal. Fitting higher order polynomials to the phase (Figure 2.26) reveals the presence of third and fourth order components. Beyond fourth order we do not see a significant reduction in the residuals, therefore, these components appear to be negligible.

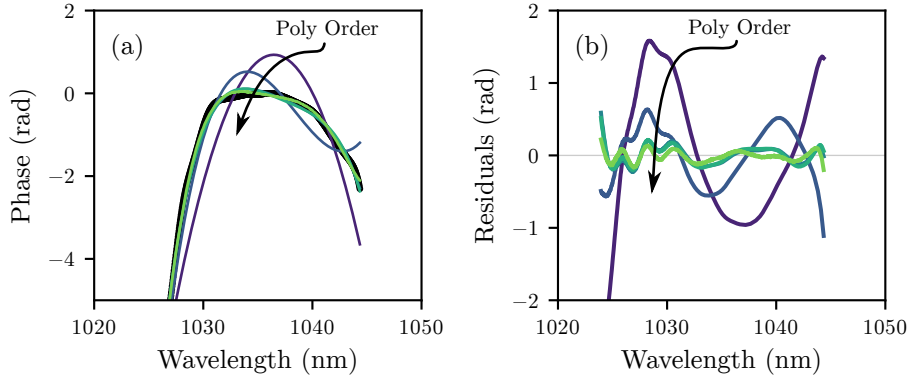


Figure 2.26: Polynomial fit of the retrieved spectral phase after the PC-rod (Figure 2.25-d). (a) Retrieved spectral phase (black points) and polynomial fits (color gradient) of order 2 to 6. Arrow indicates increasing polynomial order. (b) Residuals for the fits. Color gradient shared with panel-(a).

The retrieved temporal intensity and phase are found in Figure 2.25-f. The pulse has an acceptably smooth and symmetric profile, and a strong contrast ratio. The phase is again predominantly parabolic. It is reassuring to see that the retrieved pulse duration of 264 fs FWHM closely matches that determined by deconvolving the independently measured intensity autocorrelation with a sech pulse form (Figure C.1-e; 257 fs FWHM).

It is important to be aware that the system properties as described are

subject to spontaneous changes of the oscillator state. In particular, the oscillator has an affinity for spectrally broadened states, like that in Figure C.1-b. A discussion of this issue can be found in Appendix C.

2.6.5 Mode-Matching and Focusing

Characterization Procedure

A simulation incorporating aberrations was used to determine the correct lens configuration for mode-matching. It was performed in Zemax, using the Physical Optics Propagation (POP) tool. We traced a Gaussian beam with a waist $w_0 = 33 \mu\text{m}$ (based on the manufacturer's specifications) along the propagation axis z , on a 512×512 array of intensities spanning $1.32 \text{ mm} \times 1.32 \text{ mm}$ in the transverse plane. The initial reference plane at $z = 0$ is the location of the waist, and the face of the fictitious PC-rod. A layout of the spacing between optics is found in Figure 2.27. The intensity and phase of the propagated beam are used to calculate a coupling efficiency, based on an overlap integral with a Gaussian mode with $w_0 = 9.8 \mu\text{m}$. The mode waist is assumed, given a HCF physical core diameter of $28 \mu\text{m}$, and a guideline from our collaborators that generally the MFD is 70% of the core diameter. The system is first optimized for minimum divergence after the $4\times$ beam expander (second and third lens in Figure 2.27). Afterwards, only the HCF position is varied to optimize the coupling efficiency. We measure the spot size at the focus, and compare it to the predictions of the simulation. In addition, the effect of grating and PC-rod misalignment on the spot size is shown qualitatively. We also compare the simulated and measured coupling efficiencies.

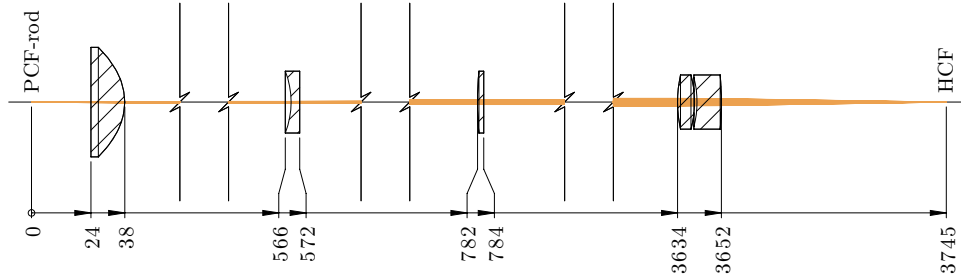


Figure 2.27: Layout of the optics for mode-matching between the PC-rod output, and the HCF fiber in the second stage of nonlinear broadening. Dimensions are in millimeters. Light propagates left to right.

To ensure consistency between the simulations and measurements, we

characterized a HeNe laser, and used it to reproduce a predicted spot size. A Newport KBX076 lens with focal length $f = 200$ mm was placed a known distance from the front of the HeNe laser, and a camera (Coherent LaserCam-HR II) was swept through focus to profile the beam (Figure 2.28 shows the layout). A Gaussian beam profile was fit to the data to determine the location, and E-field $1/e$ beam radius, of the focus. These beam parameters and distances were then back-propagated in Zemax to determine the location, and $1/e$ beam radius, of the laser waist before the lens.

With the HeNe beam parameters known, we optimized a two lens focusing system for our desired beam waist, $w_0 = 9.8 \mu\text{m}$, at the focus (Figure 2.29 shows the layout). Attaining a spot size this small required that we first expand the beam with a divergent lens. A Newport KPC037 ($f = -75$ mm) plano-convex lens, and a Newport KBX061 ($f = 100$ mm) biconvex lens made up our focusing system. A 90-10 knife-edge measurement was used to profile the beam, and the resulting widths were fit with a Gaussian beam profile to determine the location, and $1/e$ beam radius, of the focus. Once the location of the focus was known, a more thorough knife-edge measurement was made at this position, and fit with an error-function (integrated Gaussian) as a secondary measure of the $1/e$ radius. We also used a Coherent LaserCam II to profile the beam at focus, though the $6.7 \mu\text{m}$ pitch between pixels limits the resolution.

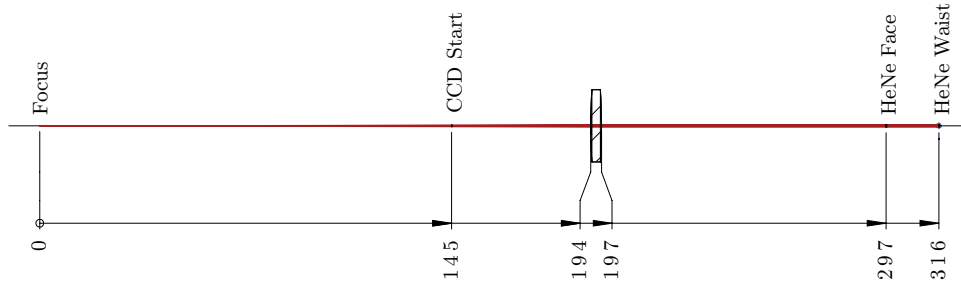


Figure 2.28: Layout of the setup used to characterize the HeNe laser. Positions are in millimeters, and relative to the measured focal position. Light propagates right to left. The start position of the camera CCD before profiling the beam is indicated, along with the physical face of the HeNe laser, and the location of the laser waist before the lens.

Characterization Results

We start by verifying the simulated spot size for the two-lens focusing system after the HeNe laser (Figure 2.29). The paraxial Gaussian beam waist

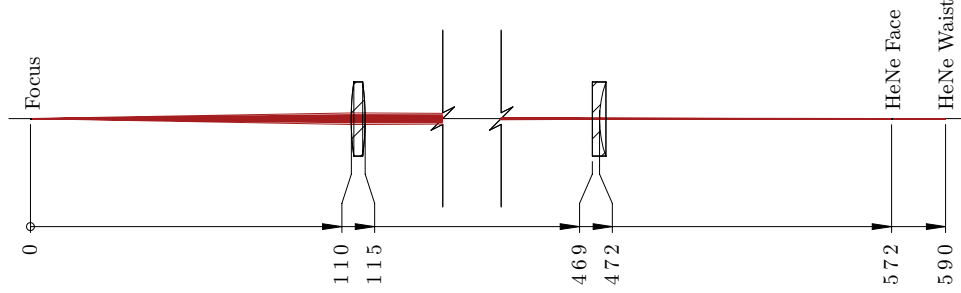


Figure 2.29: Layout of the setup used to demonstrate consistency between simulated and measured spot sizes. Positions are in millimeters, and relative to the simulated focal position. Light propagates right to left.

is optimized to be $w_0 = 9.8 \mu\text{m}$, but the second moment of the simulated spot size is $10.2 \mu\text{m}$, due to spherical aberrations. Note that the second moment width equals the E-field $1/e$ width if the beam profile is Gaussian. The results of three separate measurement techniques are shown in Figure 2.30. Both the longitudinal fit of a Gaussian profile (Figure 2.30-a) and the transverse knife-edge measurement at focus (Figure 2.30-b) yield the same $9.9 \mu\text{m}$ beam waist. These values are in excellent agreement with the model, falling short by only 3%. The third measure, directly imaging the focus on the beam profiler, clearly strains the resolution limits of the camera (Figure 2.30-c). Interpolating, and fitting the image (Figure 2.30-d), however, provides a reasonable upper limit on the waist, at $14 \mu\text{m}$. We expect a larger spot size, as the image is a convolution of the physical spot area and the pixel area. In fact, this last measurement was done largely to ascertain the resolution limit of the camera—knife-edge measurements, while accurate, are quite slow and impractical for real-time alignment. The same setup was also used to verify the resolution and fitting routine of our home-made Raspberry Pi camera beam profiler. The pixel pitch of the “pi-cam” is only $1.1 \mu\text{m}$, and we have found the resolution to be comparable to this.

With Zemax simulations for the HeNe being well reproduced, we next simulated the more complicated optical system for focusing into the HCF (Figure 2.27). We found the minimum divergence after the beam expander is $740 \mu\text{rad}$, and the numerical coupling efficiency is 91.2%. The model neglects reflection and absorption losses, which would occur at the input (coated fused silica) and output (uncoated MgF_2) windows, and off the fiber face; though, reflections off the fiber face will be small due to the hollow core. Ignoring losses, perfect coupling is achieved when the mode area and phase

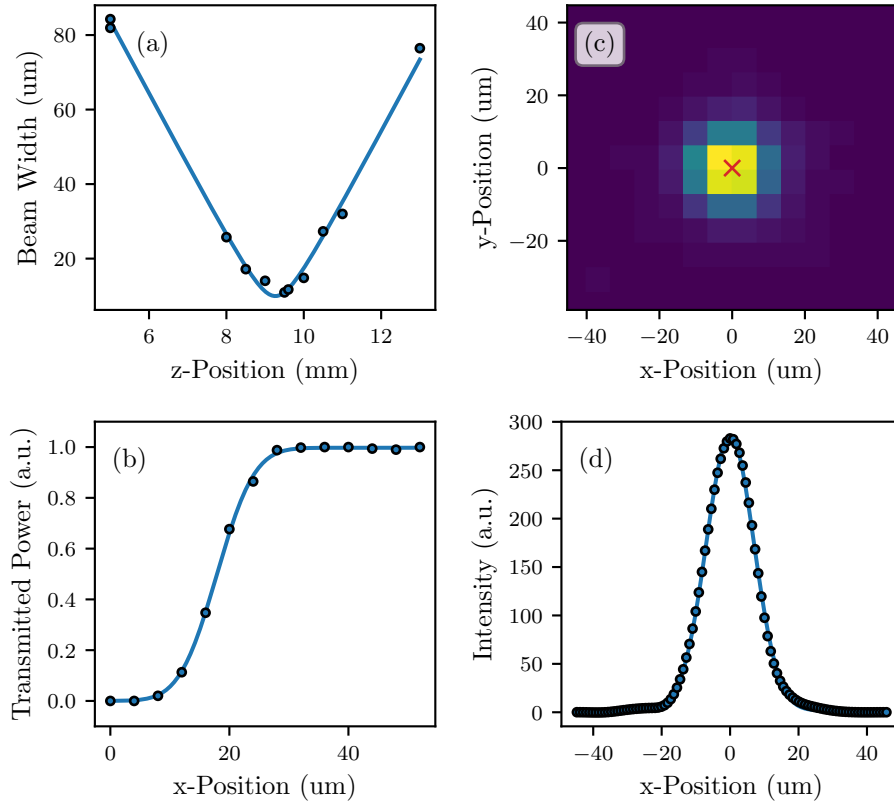


Figure 2.30: Three independent measures of the HeNe spot size, for the setup depicted in Figure 2.29. Zemax simulations predict a waist $w_0 = 10.2 \mu\text{m}$. (a) Measured beam widths (points) at positions along the beam axis, and a fitted Gaussian beam profile (blue), with $w_0 = 9.9 \mu\text{m}$. (b) Knife edge measurement at focus. Transmitted power (points) increases as the knife edge is retracted across the beam. An error-function fit (blue) yields $w_0 = 9.9 \mu\text{m}$. (c) Transverse beam profile imaged with the Coherent LaserCam II. (d) Cubic interpolation of image-(c) along a horizontal line through the red 'x' (points). Gaussian fit giving $w_0 = 14.0 \mu\text{m}$ (blue).

of the beam and HCF are identical. With only one free parameter (position of the HCF) and the presence of spherical aberrations, optimum coupling is found at a compromise between mode area overlap and a flat phase. The simulated 2D beam profile is shown in Figure 2.31-a. A 1D slice along the x-axis of the 2D profile, and the respective phase at each position is provided in Figure 2.31-b. The second moment width is $12.1\text{ }\mu\text{m}$, and importantly the phase deviates by less than $\sim 1/4\text{ rad}$ where the intensity is $> 1\%$ of the peak value.

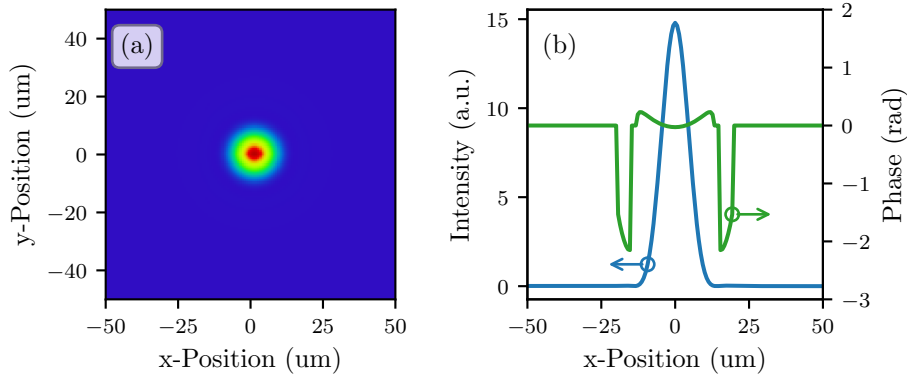


Figure 2.31: Zemax simulation of the focused beam profile on the HCF face. (a) 2D intensity plot of the beam profile. (b) 1D profiles along the x-axis. The beam intensity (blue) has a $1/e$ beam radius $w_0 = 12.1\text{ }\mu\text{m}$. The wavefront phase is also shown (green).

We measured the spot size at the HCF face with the picam (Figure 2.32-a), which allowed us to align in “real-time,” and find a 98.7% circularity spot, with $w_0 = 15.6\text{ }\mu\text{m}$. Note that the color map in Figure 2.32 has one more color than that in Figure 2.31-a, which increases the dynamic range and makes the spot appear larger. Ellipticity in the beam was minimized by adjusting the coupling into the PC-rod, and reducing spatial chirp after the compressor. An example of the ellipticity caused by spatial chirp is shown in Figure 2.32-b; misalignment into the PC-rod has the same appearance. As noted previously, misalignment of the PC-rod tends to excite the LP_{11} mode. A combination of the LP_{01} and LP_{11} modes results in an oblong beam profile. On the other hand, ellipticity due to spatial chirp is a form of astigmatism, since the beam is effectively divergent (or convergent) in the plane of incidence of the gratings. Therefore, one axis of the beam comes to a focus before the other. One can distinguish the two forms of ellipticity

by moving through focus: if the ellipse rotates, the beam is astigmatic, and spatial chirp is likely at play. Another approach is sweeping a card across the beam in the compressor, where the frequency components are spatially separated; if the focused spot does not uniformly extinguish, then one knows the beam is spatially chirped.

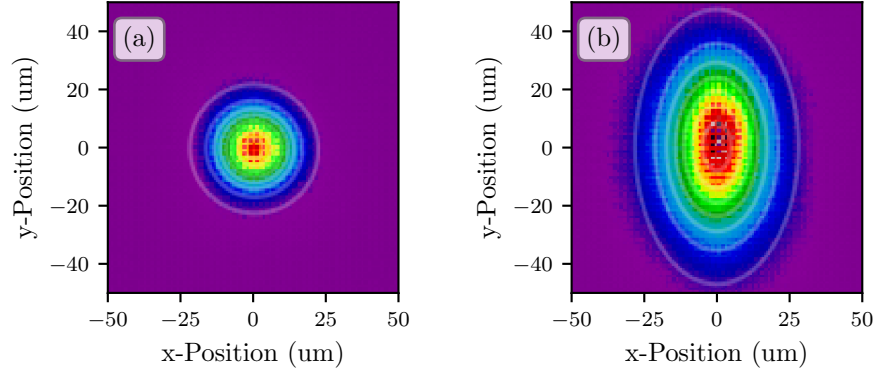


Figure 2.32: “picam” images of the focused beam profile at the HCF face. (a) Well aligned optical system. Spot circularity is 98.7%, and $w_0 = 15.6 \mu\text{m}$. (b) Spatial chirp after the grating compressor. Spot circularity is 59.5%. Along the major/minor axis $w_0 = 35.44 \mu\text{m}/21.07 \mu\text{m}$. A 2D Gaussian fit is used to extract the beam widths.

Disparity between the measured and simulated spot sizes is attributed to higher-order modes and cladding light (which is guided outside the core), as well as misalignment of the collimating lens after the PC-rod. The first issue is accentuated by attenuating the beam with the HWP and thin film polarizer, as is done before coupling to the HCF, or imaging the spot size on the pi-cam. With most of the beam rejected, we are left with the off-polarization components, which tend to be composed of higher-order modes and cladding light. When the attenuation is decreased, we see an increase in the coupling efficiency, leading us to believe that the true beam (with a smaller proportion of cladding light) matches the simulations more closely. The second issue comes about when trying to collimate the beam after the PC-rod. We have found that it is not possible to adjust the collimating lens, such that the beam diameter after the PC-rod matches the simulation for all positions along the propagation axis. This again may be due to cladding light, but regardless, makes the position of the collimating lens uncertain. Compounding the issue is the lens’ short focal length ($f = 32 \text{ mm}$),

which makes visually undetectable (~ 1 mm) changes in its position quite significant.

Nonetheless, we measure coupling efficiencies near the predicted value: routinely percentages in the mid-80s, and on one occasion 90%. Thus, it seems the system is tolerant to the observed deviations.

2.6.6 Beam Stabilization

Characterization Procedure

To directly measure movement of the beam, we use a position sensitive detector (Ref. PSD) and a camera (pi-cam). Their placement on the optical bench is shown schematically in Figure 2.22. Generally, the PSD measures transverse movement of the beam, while the focusing lens before the pi-cam maps angular deviations to position. We also measured coupled power exiting the HCF. This indirectly tracks beam movement, but directly quantifies coupling stability to the HCF.

Initially we used the pi-cam (a Raspberry-Pi based camera) and a power meter (Thorlabs PM16-121) to record the relative change in coupling/beam-pointing stability when the active stabilization system was engaged. The center position of the beam on the pi-cam is found with a 2D Gaussian fit. A HeNe laser was used as a stable reference to ensure deviations in beam position were due to pointing instabilities, not movement of the optics used to measure it. The output beam was sampled with a B-coated wedge; therefore, power measurements are self-referenced (not absolute coupled power). Prior to recording data, the amplifiers ran at standard operating drive currents (PCF, 3.3 A; PC-rod, 7 A) for over one hour. Coupling to the HCF was optimized at low power, then increased to 18 W over the span of 1 min. We ran the system for 2.5 hours with AS engaged, then turned off the AS and left the system running open-loop for an additional 2.8 hours. During the closed-loop portion of the test, minor transverse adjustments were made to the three-axis stage holding the HCF. Focal distance adjustments were not required. The results are shown in Figure 2.34. Note, the HCF was held in a non-conventional mount, and protruded only 1 mm from the front face (Figure 2.33-a).

Additional temporal and spectral characterization of the beam-pointing noise-floor, with and without AS, was performed after adding the reference PSD (Figure 2.22). A short time-window centered about the point where the AS was engaged is shown in Figure 2.35-a. Analog signals from the PSD and a photo diode (Thorlabs DET10C) after the HCF (indirect sampling with a

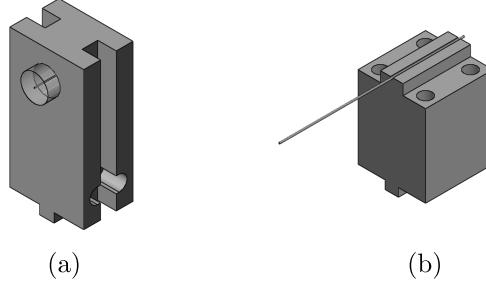


Figure 2.33: Comparison of fiber chuck designs. (a) A new variation of the HCF mount, or “fiber-chuck”. The HCF is held in the bore of a thick-walled capillary and protrudes only ~ 1 mm from the front face. This design is more rigid than the typical cantilevered, thin-capillary shown in panel (b). Ideally, most of the uncoupled, light passes unabsorbed through the thick-walled capillary, without hitting the rest of the mount. In practice, the cantilevered approach is less susceptible to heating, and related thermal expansion.

B-coated wedge) were recorded with a dynamic signal analyzer (SR780) in the frequency-domain (Figure 2.35-b/c).

Finally, we reverted to the conventional, cantilevered, thin-capillary mount for the HCF (Figure 2.33-b). Reasoning for the change can be found in Figure 2.33 and section 3.1.2. Pointing instability was recorded with the pi-cam. Concurrently, we directly measured the coupled power with a thermopile (Thorlabs S425C-L). Prior to recording data, the amplifiers ran at standard operating drive currents (PCF, 3.3 A; PC-rod, 7 A) for over one hour. Coupling to the HCF was optimized at low power, then increased in steps to 15.3 W over the span of 1 hour. Minor adjustment (all transverse) of the three-axis stage holding the HCF was performed with each step increase in power, and 20 min after reaching 15.3 W. Subsequently, the source was unperturbed for 2 hours. The results are shown in Figure 2.36.

Characterization Results

Initial testing and validation of the active stabilization (AS) system is shown in Figure 2.34. As suspected, there are both high-frequency (tens of hertz) and low-frequency (sub-hertz) beam-pointing instabilities that are not present on the reference HeNe laser beam (further characterization on a short time scale and in the frequency domain is shown in Figure 2.35). Similar instabilities occur for the coupled power as well; however, reoccurring and rapid degradation of the coupling, or “thermal runaway,” is more troubling, and

is in fact what prompted implementation of the active stabilization system. Evidently, while active stabilization is effective at reducing pointing instability, it does not eliminate thermal runaway. A clear example occurs at around 1.4 hours, where the beam is relatively stable and thermal runaway initiates nonetheless. Contrast this with the time period between 3 hours and 5 hours where the beam wanders substantially, yet there is no rapid loss of coupling. All this is not to say that pointing instabilities cannot lead to thermal runaway—at 2.5 hours the AS is disengaged and the beam position jumps, triggering a thermal runaway event—but, there is evidence of a secondary source of coupling loss.

We believe that along with pointing instabilities, thermal expansion of the mount holding the HCF (fiber-chuck) causes a loss of coupling—essentially the HCF becomes a moving target. In the first hour (Figure 2.34), the fiber chuck and beam pointing are still equilibrating and apparently balance each other. At 1 hour we reach a new maximally coupled configuration, as indicated by reduced high-frequency noise content in the coupled power (we have found that exciting higher order modes increases the amplitude of high-frequency noise on the coupled power; therefore, when only the lowest order mode is excited [optimally coupled], high-frequency noise is reduced). After 1 hour the beam-pointing stabilizes, but we believe the fiber chuck continues to expand, decoupling the signal from the HCF. The process continues unbound, since decoupling leads to more heating and thus faster decoupling—in other words, thermal runaway.

To minimize heating of the fiber-chuck, our collaborators cantilever the HCF ~ 25 mm in front of it (see Figure 2.33 and Section 3.1.2 for more details). We implemented the same design, then repeated the beam stability measurements using the pi-cam and power meter. The results in Figure 2.36 show that combining active and passive stabilization can yield stable coupling for over 2 hours. For a sense of scale, compare the $\pm 20\%$ power levels with and without (Figure 2.34) proper passive stabilization. It should be noted that gradually approaching the final coupled power contributed to improved stability. Whether the system reaches the same degree of stability in less than 1 hour remains to be seen, but these initial results are promising.

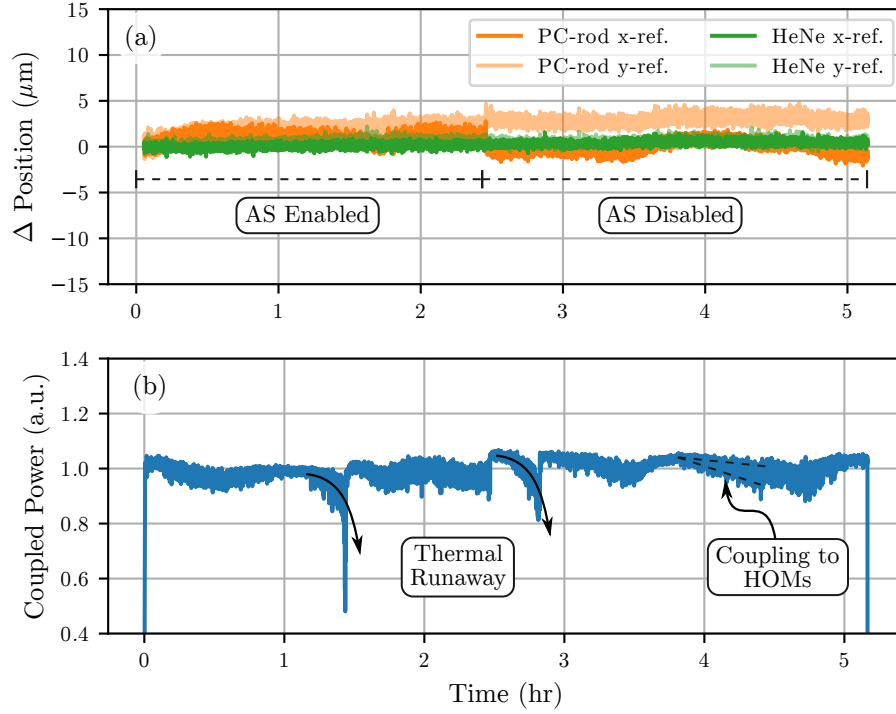


Figure 2.34: Relative change in coupling/beam-pointing stability with the active stabilization (AS) system engaged, as measured with the pi-cam and power meter. After 2.5 hours AS was disengaged and we left the system running unstabilized for an additional 2.8 hours. Minor transverse adjustments were made to the three-axis stage holding the HCF at 1.4 hours and 2.5 hours. Focal distance adjustments were not required (a) Pi-cam measured beam position. A HeNe laser was used as a stable reference. There is a significant increase in pointing instability with AS disengaged. (b) Normalized power coupled to the HCF. Rapidly degrading coupling at 1.4 hours is attributed to thermal expansion of the fiber-chuck, since the beam pointing is fairly stable. Decoupling leads to further heating and thermal runaway. Thermal runaway at 2.5 hours is initiated by pointing drift when AS is disengaged. Noise amplitude increases as the coupling worsens and higher order modes (HOMs) are excited.

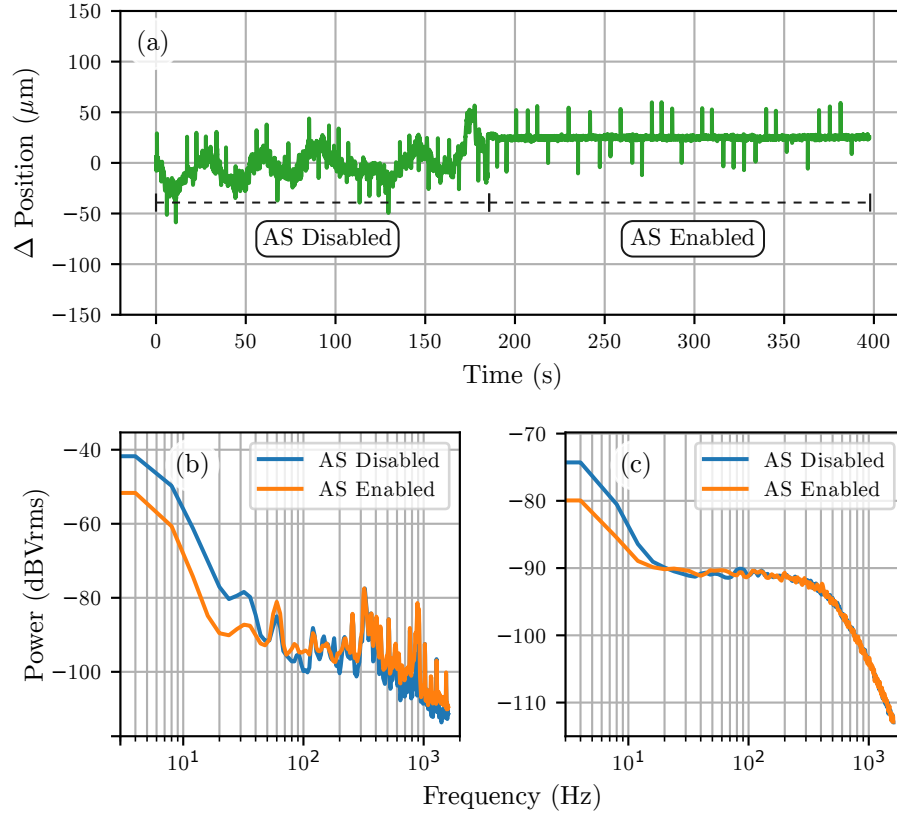


Figure 2.35: Relative change in coupling/beam-pointing stability with the active stabilization (AS) system engaged, as measured with the PSD and power meter. (a) PSD measured beam position. Engaging AS substantially decreased the amplitude of low-frequency pointing noise. High-frequency spikes were not attenuated, though, they may not even be pointing related (e.g. electrical noise on the measured signal). (b) PSD measured spectral power density of pointing noise. With AS engaged, the noise below 40 Hz is attenuated by ~ 10 dB. (c) Photo-diode measured spectral power density of noise in the coupled-power. With AS engaged, the noise below 20 Hz is attenuated. Coupling to the HCF appears to smooth-out high-frequency pointing noise.

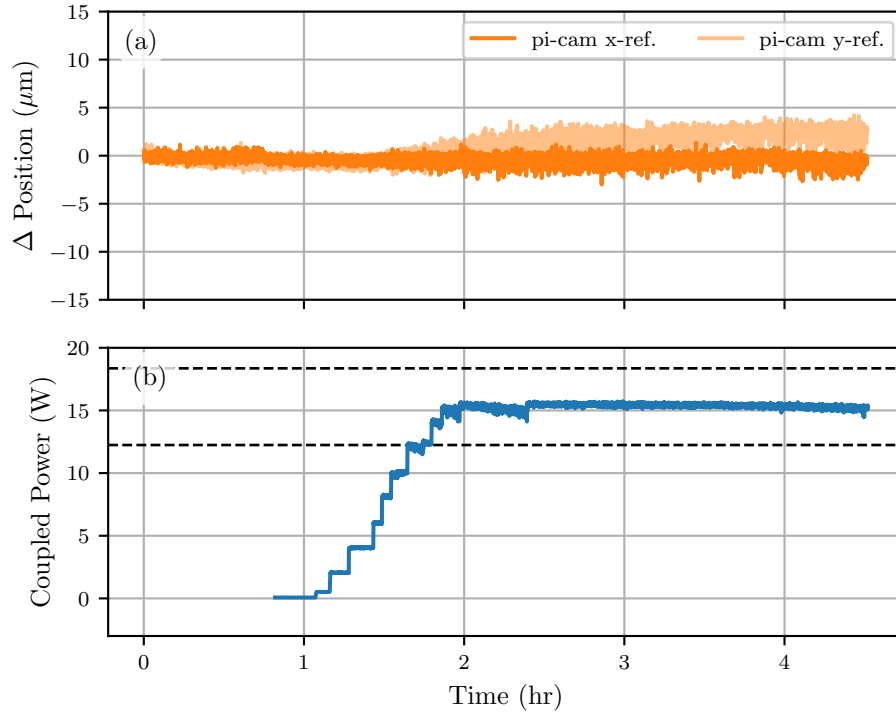


Figure 2.36: Coupling/beam-pointing stability with active stabilization (AS) and passive stabilization (conventional cantilevered thin-capillary mount for the HCF), as measured with the pi-cam and power meter. The coupled power was gradually ramped over ~ 2.5 hour to 15.3 W. Minor transverse adjustments were made to the three-axis stage holding the HCF with each step-increase in the coupled power, and at 2.4 hours. Focal distance adjustments were not required. Afterwards, the HCF was stably coupled for 2 hours (a) Pi-cam measured beam position. We believe pointing noise increases with coupled power because there is more back-scattered light. Note, we adjust the coupled power down-stream of the sampling points for AS feedback (position sensitive detectors), so it should not directly effect the beam stability. (b) Power coupled to the HCF. Dashed black lines are $\pm 20\%$ of 15.3 W. The rate of thermal runaway is likely reduced because of the gradual approach to the final coupled power and the improved passive stability of the fiber-chuck.

Chapter 3

Spectral Evolution in Gas-Filled Hollow-Core Fiber

3.1 Spectral Broadening and Compression

3.1.1 Design Principles

Spectral bandwidth ultimately determines the lower limit on pulse duration. To obtain shorter pulse durations the spectrum must first be broadened, necessitating a nonlinear interaction between a medium and the electric field to generate new spectral components. In our case, this interaction occurs in a gas-filled hollow core fiber, where high intensities can be sustained for extended lengths. Subsequently, a grating or chirped-mirror based compressor is used to remove the phase added by the broadening process and reduce the temporal pulse width.

The particular HCF used in this stage has a “single-ring,” or “revolver,” geometry. A diagram indicating the key geometric parameters of revolver fiber and a scanning electron microscope (SEM) image of the actual fiber used are shown in Figure 3.1. The thin walls, and negative curvature of the micro-structure about the fiber core have been shown to lower propagation losses at smaller core radii [55, 56]. Here, a negative curvature indicates the inner surface of the cladding is convex and curves away from the core; in contrast, the inner surface of a single-wall capillary is concave. We chose to use this fiber largely due to its availability from collaborators (John Travers, Heriot Watt University), but it does represent leading edge hollow-core fiber technology. Yu and Knight [55] provide an excellent historical overview of progress towards the single-ring geometry.

The predominant mechanism for spectral broadening in this stage is SPM, the rate of which depends on pulse intensity and the nonlinear parameter γ . For the revolver geometry, the fundamental mode is almost entirely contained in the central gas-filled core [55] and has an effective

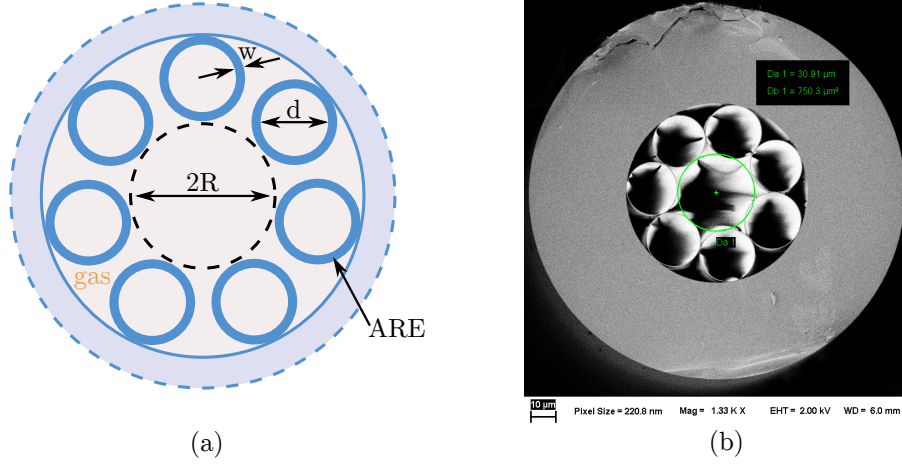


Figure 3.1: Cross section of the hollow core fiber used in the spectral broadening stage. (a) Diagram indicating the key geometric parameters of a single-ring hollow core fiber. The seven interior capillaries, also known as “anti resonant elements” (AREs), reduce propagation losses in comparison to single-walled capillary. A gap may be left between AREs to enhance the effect. See [57] for a more detailed discussion. (b) Scanning electron microscope (SEM) image of the fiber. The measured parameters are: $2R = (30 \pm 2) \mu\text{m}$; $w = (300 \pm 10) \text{nm}$; $d = (20 \pm 2) \mu\text{m}$. Note that there is some non-uniformity in the AREs. SEM image recorded by Mike Chang, University of British Columbia.

area of $A_{\text{eff}} \simeq 1.5R^2$ [47], where R is the core radius as defined in Figure 3.1. Since the mode is confined to the core, it is reasonable to ignore the nonlinear index n_2 of the fiber material and consider only that of the gas [47, 55]. We elected to use the value $n_2 = 1.4 \times 10^{-23} \text{m}^2/\text{W}$ for argon at $p_0 = 1 \text{bar}$ and $T_0 = 300 \text{K}$, measured with a 120 fs pulse at 800 nm [28], following previous work by Emaury et al. [10] on a comparable laser system. As noted in their article and elsewhere [5, 10, 28], the nonlinear index varies depending on which phenomena are considered and what source is used to measure it. This is particularly true when considering molecular gases with rotational/vibrational degrees of freedom, but even distilling the problem down to a single elementary phenomenon—the third-order hyperpolarizability, or equivalently, the instantaneous electronic response—of monatomic gases (like argon) leaves discrepancies. For example, we calculated $n_2 = 8.2 \times 10^{-24} \text{m}^2/\text{W}$ and $n_2 = 1.5 \times 10^{-23} \text{m}^2/\text{W}$ (at $p_0 = 1 \text{bar}$ and $T_0 = 300 \text{K}$), given the experimental values found by Lehmeier et al. [22] and Shelton [42] respectively. Therefore, in choosing n_2 some skepticism is required.

The initial value for the nonlinear index, $n_2(p_0, T_0)$, can be scaled in proportion to the gas density [22], or equivalently

$$n_2(p, T) = \frac{pT_0}{p_0T} n_2(p_0, T_0). \quad (3.1)$$

The nonlinear parameter can then be defined as

$$\gamma \equiv \frac{2\pi}{\lambda} \frac{n_2(p, T)}{A_{\text{eff}}} \quad (3.2)$$

and for a 30 μm core diameter, filled to 20 bar with argon, at a central wavelength of 1035 nm, we find $\gamma = 5.0 \times 10^{-6} (\text{mW})^{-1}$.

Conventionally, spectral broadening is best carried out in a normally dispersive regime [47]. In this way, soliton dynamics are avoided, and the temporal phase imparted by SPM is predominantly parabolic (a more detailed discuss of the process can be found in Section 2.2.6). That being said, spectral broadening in the anomalous dispersion regime can be carried out successfully if the process is terminated before soliton effects, such as modulation instability or soliton fission occur [10]. In fact, one may leverage the first stage of higher-order soliton evolution, where SPM and anomalous dispersion work in tandem to spectrally broaden and temporally self-compress the pulse. The resulting temporal phase will only be parabolic near the center of the pulse [1], but the spectrum is broader than that for an equal length of normally dispersive fiber, and others have compressed these pulses to within 5% of the transform limit [10].

For a given fiber geometry, varying the species and pressure of the filling gas sets the dispersion regime. Transitioning between the anomalous and normal regime is possible because the fiber is intrinsically anomalously dispersive—this property can be seen in the dispersion profiles for the evacuated fiber in Figure 3.2-a. The well studied Marcatili–Schmeltzer (MS) model [25] and a first order perturbation of the MS model that accounts for resonances in the thin-walled cladding [57] were used to calculate the HCF dispersion profiles. It is clear that resonances have a significant effect on the dispersion, which will be discussed in the following sections. For now, the MS model sufficiently captures the net dispersion and its trend with varying gas pressure. It is used to generate a set of dispersion curves for argon from 0 bar to 25 bar, depicted in Figure 3.2-b, where we see the ZDW shifts to longer wavelengths with increasing pressure. This trend holds for other noble gases, with atomically heavier gases hastening the evolution. We chose to use argon for two reasons: it is relatively inexpensive and the

ionization potential is higher than that of krypton or xenon. The latter property is particularly important in the anomalous regime, where soliton self compression can lead to extreme peak powers, and hence, ionization of the gas. Given we are technically limited to a pressure less than 24 bar, our pulse at ~ 1035 nm will propagate in the anomalous regime, where avoiding the perturbative effects of ionization prolongs coherent spectral broadening through SPM.

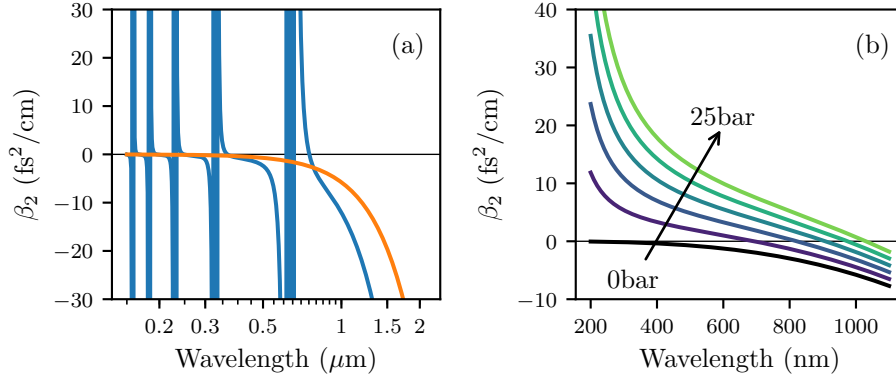


Figure 3.2: Dispersion profiles of the HCF used in the spectral broadening stage, which has a core diameter $2R = 30$ μ m and an ARE wall thickness $w = 300$ nm. (a) Comparison of the Marcatili-Schmeltzer (MS) model (orange) and a perturbative extension of the MS model that accounts for cladding resonances (blue). Here, the fiber is evacuated and dispersion results from the structure of the waveguide. Note that resonances both create and shift zero dispersion wavelengths (ZDWs). (b) MS model of the dispersion as the HCF is pressurized with argon from 0 bar to 25 bar (5 bar steps). The ZDW moves to longer wavelengths with greater pressure.

In the following sections we will cover implementation of the spectral broadening stage and the optical elements used to characterize the output. We then will compare the measured results to simulations and conclude with an assessment of the possible next steps, including a short conceptual discussion on resonant dispersive generation in the final stage.

3.1.2 Implementation

Stabilization of the beam and coupling to the HCF were covered in Section 2.6. Here we will detail the gas cells, and how the HCF is affixed. Optics for characterizing spectral broadening in the HCF will be discussed in the procedural section that follows.

The HCF is mounted between two gas cells in order to remove atmospheric contaminants (through evacuation and purging), and regulate the pressure of the filling gas (a top-down view of one gas-cell is found in Figure 3.3-a, and a schematic of the assembly is shown in Figure 3.4). For ease of reassembly, the two gas cells are connected with a 1/8 inch stainless-steel tube, and the HCF rests in its bore. The small diameter tube is flexible, allowing the gas cells to move independently, yet rigid enough to prevent the HCF from sagging. Alternatively, one could forgo the tube, seal the HCF/gas-cell interface with glue, and tension the HCF so it remains straight while suspended. The latter technique may be used in the future, if having a pressure gradient across the HCF is desired.

The input and output ends of the two-cell assembly are sealed with windows to allow the beam to pass (Figure 3.3-b). The minimum window thickness to avoid fracture and optical distortions from bowing, for a 15 mm clear aperture at 30 bar, is 1 mm for MgF_2 and 1.5 mm for fused silica (FS) (calculations courtesy of F. Belli, Heriot Watt University). We use a 3 mm thick, coated FS window at the input, and an uncoated 1 mm thick MgF_2 window at the output. MgF_2 is used at the output, as it is less dispersive and has a lower nonlinear refractive index; the accumulated nonlinear phase, or “B-integral”

$$B = \frac{2\pi}{\lambda_0} \int n_2 I(z) dz, \quad (3.3)$$

where n_2 is the nonlinear index and $I(z)$ is the intensity, is negligible at only 22 mrad. Here, we assume a 5 μJ , 30 fs FWHM duration pulse, centered at 1030 nm, with Gaussian temporal and spatial form, and a $1/e$ E-field radius of 0.5 mm in the window. Although marginally worse, a 3 mm coated FS window at the output would likely be sufficient, since the B-integral for the same output pulse is only 310 mrad, and significantly, its use would lower the reflection losses from $\sim 2.5\%$ per surface for uncoated MgF_2 , to $< 0.5\%$ – this change may be made in the future.

Inside the gas cell, the stripped end of the HCF (cladding, 150 μm OD; acrylate coating, 350 μm OD) is held by a sliding fit in the bore of a FS capillary (320 μm OD, 200 μm ID), which in turn is bonded (with Bondic UV cured adhesive) to an aluminum pedestal, or “fiber chuck”. To reduce the flux of uncoupled light impinging on the front face of the fiber chuck, the capillary is cantilevered ~ 25 mm in front of it. Offsetting the focal point in this way allows the uncoupled light to diverge and partially avoid the fiber chuck. As discussed in Section-2.6.6, thermal expansion of the fiber chuck is believed to be the primary cause of sub-hertz coupling degradation; therefore, minimizing the absorption of scattered light should reduce the

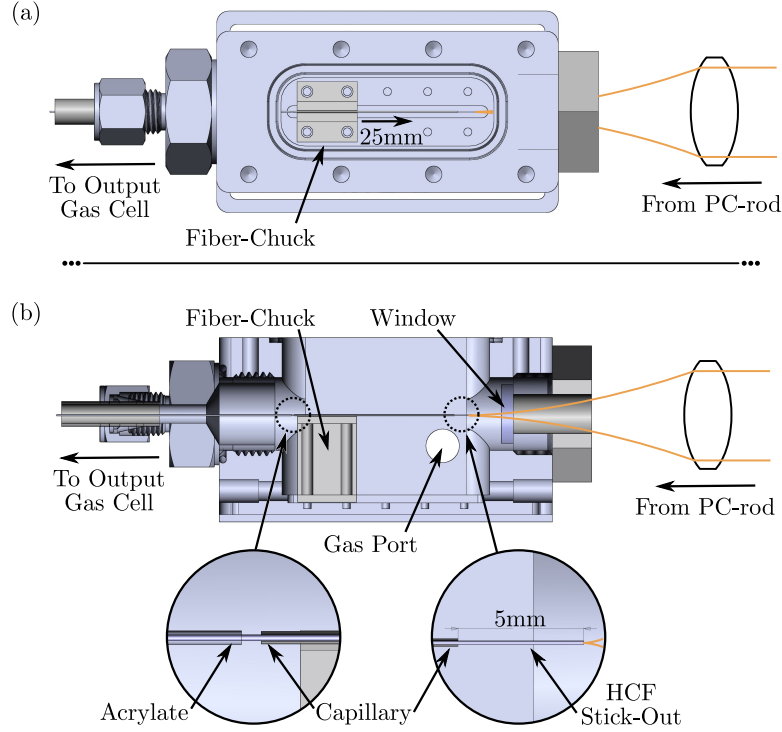


Figure 3.3: Detailed view of the input gas cell, and the fiber chuck used to hold the HCF. (a) Top view of the gas cell, with the cover removed. The beam from the PC-rod is focused on the HCF face. A capillary (320 μm OD, 200 μm ID) is cantilevered ~ 25 mm from the fiber chuck it is glued to. The stripped HCF (cladding, 150 μm OD; acrylate coating, 350 μm OD) slides within the capillary, and extends an additional ~ 5 mm (see detail in (b)). (b) Sectioned side-view of the gas cell. The beam is focused through the 10 mm fitting aperture, and window at right. At left the HCF runs through a stainless-steel tube connecting the input and output gas cells. Also shown is the 5 mm HCF stick-out from the capillary and the “butt joint” where capillary and acrylate coating meet.

severity of this problem. Similarly, to lessen the coupling of scattered light into the capillary wall, the HCF protrudes 5 mm from its face (see detail in Figure 3.3-b). Light coupled to the capillary wall is guided by total internal reflection (TIR), but can escape evanescently where glue partially index-matches the fiber chuck, resulting in additional heating. However, a greater concern arises at the rear end of the capillary, where it is loosely butt-coupled to the acrylate coating on the HCF (see detail in Figure 3.3-b). When the TIR-guided stray-light exits the capillary, it strikes the acrylate and can not only burn the coating, but lead to melting and complete failure

of the HCF. Increasing the gap in the butt-joint is one way to alleviate the problem; however, critically, no glue should be used to secure the HCF to the capillary within the input gas cell, as this exacerbates absorption and heating. To prevent the HCF from moving when the gas cells are evacuated or pressurized, a small drop of glue (Bondic; UV cured) is placed at the interface where the HCF enters the capillary within the output cell. Thankfully, we have found that the 5 mm offset and chosen glue location eliminates the heating issue for the tested average powers, which are below 30 W.

Apart from the gas cell infrastructure surrounding the HCF, a collimating lens after the output gas cell is the only component specific to this stage—we have covered the input coupling optics in the preceding PC-rod stage (Section 2.6) and the remaining optics after the output cell are temporary diagnostics. Currently we use a plano-convex lens, with a focal length $f = 100$ mm.

3.1.3 Characterization Procedure

All measurements were made with an AOM division factor $N = 8$, or equivalently, a repetition rate of 4.9 MHz. The PCF preamplifier and PC-rod were allowed to equilibrate for at least one hour, at a pump current of 3.3 A and 7 A respectively, before measurements were recorded. The active stabilization system was engaged, and the compressed output of the PC-rod was maximally attenuated with HWP (1) and TFP (2)—refer to Figure 3.4 for component numbering in this section. While the system equilibrated, the active stabilization set points were manually monitored, and repositioned at midrange. Concurrently, a ~ 200 mW portion of the signal was allowed to pass TFP (2) by rotating HWP (1) and used to adjust the coupling into the HCF. A power meter (PM) behind wedge (8) monitored the coupling at higher powers. The power reading after wedge (8) is 1/1.20 times that after the collimating lens, accounting for losses off the front and rear surface. The ratio was determined empirically by measuring the power after and before wedge (8), which were 15.3 W and 18.3 W respectively.

The gas cells and HCF between them were purged by two cycles of evacuation and filling with argon: a roughing pump connected to the output cell held a vacuum of ~ 1 mTorr for 20 min, then the two cells were filled with argon to ~ 500 mbar over 5 min, via the input cell. At the end of the second cycle, the filling pressure was ramped gradually (~ 2 bar/min) to the set point. We emphasize that increasing the pressure too rapidly can damage the delicate internal structure of the HCF. Moreover, large changes

in pressure tend to decouple the signal, thus, incrementally increasing the pressure and adjusting the coupling is beneficial.

Spectral broadening was investigated for pulse energies from 1 μ J to 4.0 μ J, and pressures from 0 bar to 25 bar. The input pulse duration was (260 ± 20) fs for all measurements. The pulse energy and pressure were increased primarily to raise nonlinearity and strengthen SPM; in this respect, dispersion was merely a dependent variable. An optical spectrum and SHG-autocorrelation were recorded at each step; a FROG retrieval was taken at the final set point.

Before increasing the pulse energy, a reference spectrum, autocorrelation, and FROG retrieval of the input to the HCF (compressed output of the PC-rod) at 200 mW (41 nJ/pulse) were recorded by diverting the beam with flip-mounted mirror (3). Note that both the diverted-input and output of the HCF follow the same beam path after uncoated wedge (9). Before merging at this point, the HCF output was attenuated by wedges (8) and (9). To explore nonlinear dynamics in the HCF, the entire range of input pulse energies are required, necessitating additional attenuation. It is important to note that BK-7 wedges at an incident angle of 45° are polarizing elements, with an S to P reflectivity ratio of $\sim 10:1$; therefore, HWP (7) was set to maximize the PER after wedge (19), which corresponds to an S-polarized signal.

A beam reducer, consisting of bi-convex lens (11) ($f = 500$ mm) and plano-concave lens (12) ($f = -150$ mm), prevented beam-clipping in the following diagnostics. Flip-mounted mirrors (13) and (18) switch in a grating compressor, which was only used for the spectrally broadened output of the HCF. Gratings (14) have a line density of 1250 l/mm, and their separation was adjusted to maximally compress the pulse, though we were limited by the minimum spacing of (1.0 ± 0.5) mm. Retro-reflecting mirror (15) is a roof prism, which vertically displaces the beam so that pick-off mirror (16) can redirect it. Compressor-induced spatial chirp was reduced by focusing the compressed output with a $f = 150$ mm lens, and minimizing astigmatism. Mirror (17) is simply a turning mirror, and wedge (19) provides another decade of attenuation. HWP (21) and PBS (22) further polarize the beam, but more importantly, were used to polarization-resolve the spectrally broadened output. Mirrors (23) through (28) were used to align the three diagnostics, and HWP (30) ensures S-polarization, as required by the SHG-autocorrelator.

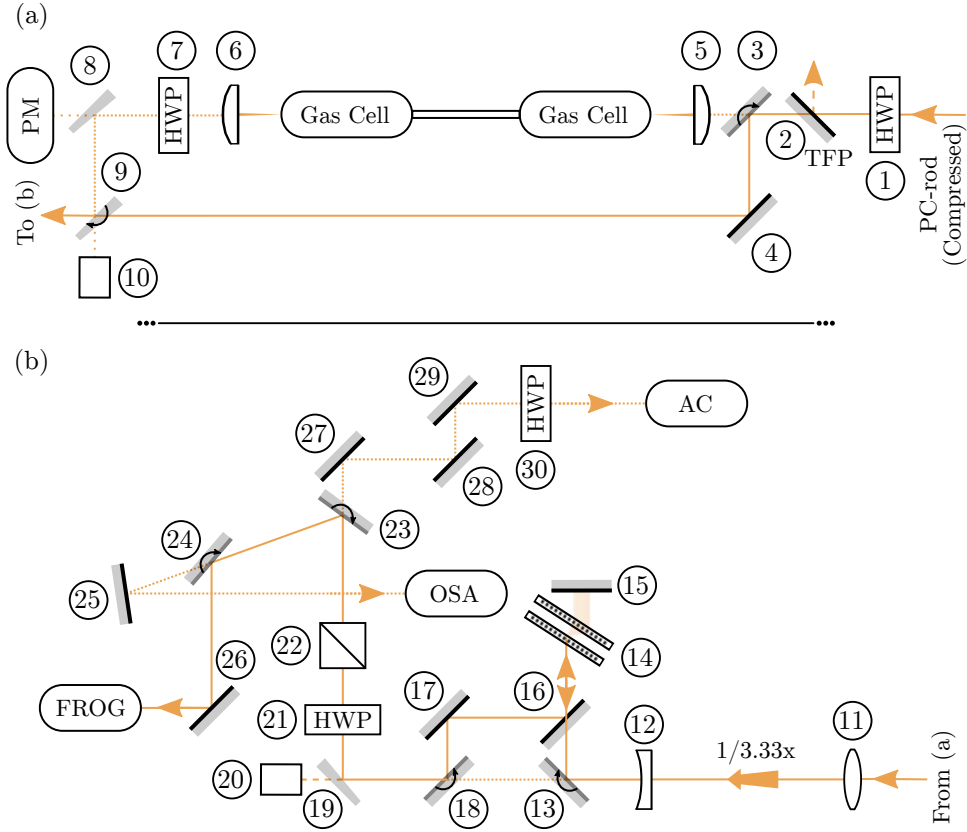


Figure 3.4: Schematic of the diagnostics used to characterize the PC-rod and HCF-based spectral broadening stage. PM, power meter; HWP, half wave plate; TFP, thin film polarizer; AC, autocorrelator; OSA, optical spectrum analyzer. Flip-mounted optics for switching are indicated with a curved arrow. (a) Stabilized input from the PC-rod enters at right, and is attenuated. Flip-mirror (3) allows the HCF to be bypassed. The undivided pulse power is necessary for fully exploring nonlinear dynamics in the HCF, thus, wedges (8) and (9) are used to attenuate the output. Coupling efficiency is monitored with a power meter, after accounting for reflection losses. (b) Continuation of the panel above. A beam reducer prevents clipping in the following diagnostics, and the selectable grating pair compresses the spectrally broadened output of the HCF. The beam is polarization resolved by HWP (21) and polarizing beam splitter (22).

3.1.4 Characterization Results

We carried out three separate broadening trials to better understand the pertinent parameters and effects stifling spectral broadening. Results for the preliminary trials and analyses dispelling higher-order nonlinear effects can be found in Appendix D. Here, we will discuss the results of the final trial, which included SHG-FROG retrievals of the pulse before and after spectral broadening, and after compression with a grating pair. Given the relative similarity between trials, we believe the last trial is generally representative of the three.

For the final broadening trial, we used a 41 cm length of HCF. The coupling efficiency and PER were 83% and 16 dB respectively. The filling pressure was held at 20 bar while the pulse energy was increased. The spectra and autocorrelations are shown in Figure 3.5. The -40 dB spectral bandwidth at $4.0 \mu\text{J}$ and 20 bar is 1.9 times greater than that of the initial pulse coming from the rod. This magnitude of broadening is consistent with the previous trials, which ranged from 2 to 3 times the initial spectral width under similar conditions. Strong spectral modulations that develop on only the short-wavelength side of the spectrum (particularly clear on a linear scale in Figure 3.5-b) indicate an asymmetric pulse intensity envelope, and therefore, asymmetric spectral phase.

The pulse duration increases monotonically with pulse energy, which is somewhat surprising given we are operating in the anomalous dispersion regime and expect soliton self compression. The same behavior was noted in one of the previous trials. In the other, there appeared to be slight self-compression, though poor PER rendered the results inconclusive.

SHG-FROG retrievals of the input, uncompressed output, and compressed output are found in Figure 2.25, Figure 3.6, and Figure 3.7 respectively; they provide crucial insight into the temporal and spectral phase of the pulse. As we supposed earlier based on the asymmetric broadening, the temporal intensity—and therefore, spectral phase—of the uncompressed pulse is also asymmetric. Initially, we explored higher-order nonlinear effects that could lead to this asymmetry (Appendix D), but as will be discussed shortly, our simulations lead us to believe it is actually inherited from the asymmetric phase of the input pulse.

The FROG retrieval of the compressed pulse (Figure 3.7) shows an over two-fold decrease in pulse duration ($399 \text{ fs} \rightarrow 174 \text{ fs}$). However, we were unable to pass through the minimum duration point, since the grating pair could not be brought close enough together (the minimum distance was $(1.0 \pm 0.5) \text{ mm}$). Given this shortfall, it is likely that the pulse is *over com-*

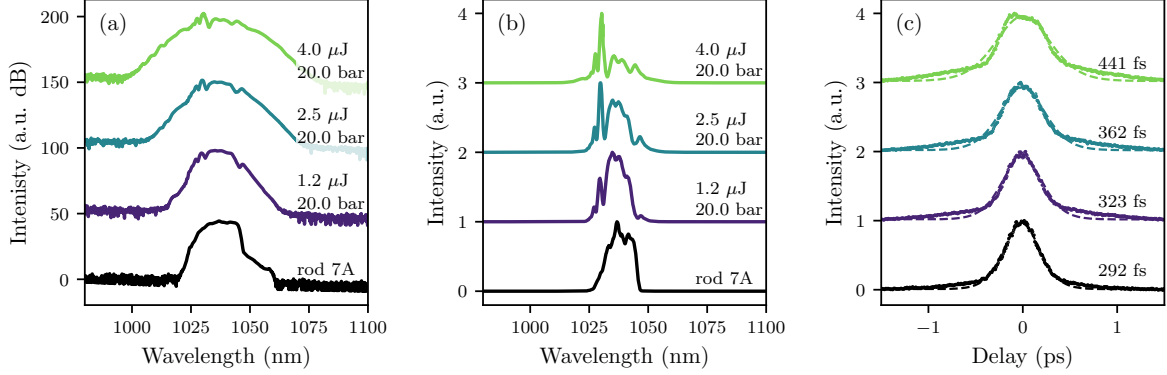


Figure 3.5: Measured spectra and autocorrelations after a 41 cm length of HCF with a satisfactory PER (16 dB). Geometric properties of the fiber are listed in Figure 3.1. (a)/(b) log/linear scale spectra showing broadening as the pulse energy and filling pressure are increased. The spectrum develops an asymmetry, being modulated at short wavelengths and smoothly broadened at long ones. The modulations are more pronounced on a linear scale. (c) Corresponding autocorrelations of the output, with the deconvolved FWHM pulse durations labeled. A sech pulse form is assumed. The pulse duration increases monotonically with pressure and pulse energy, likely tracking a shift towards the normal dispersion regime, in addition to an increase of nonlinear phase. Note there are no satellite peaks with an improved PER. The 441 fs duration trace shows some asymmetry; it is unclear if this is simply from misalignment. External compression was not attempted.

pressed and the minimum duration could be less than 174 fs. To investigate, we numerically compressed the FROG-retrieval of the uncompressed pulse leaving the HCF. The results for three grating separations are shown in Figure 3.8, and they appear to corroborate our suspicion. A minimum duration pulse is found at a grating separation of 0.6 mm, which is less than the minimum grating spacing in the lab. A pulse duration of 178 fs—which is comparable to the FROG retrieved duration of the compressed pulse—is found for a grating separation of 1.2 mm. This separation is nearly equal the minimum grating separation; therefore, we believe the pulse is indeed over compressed. A less dispersive grating pair, or a set of “chirped” mirrors, could likely be used to attain a shorter pulse.

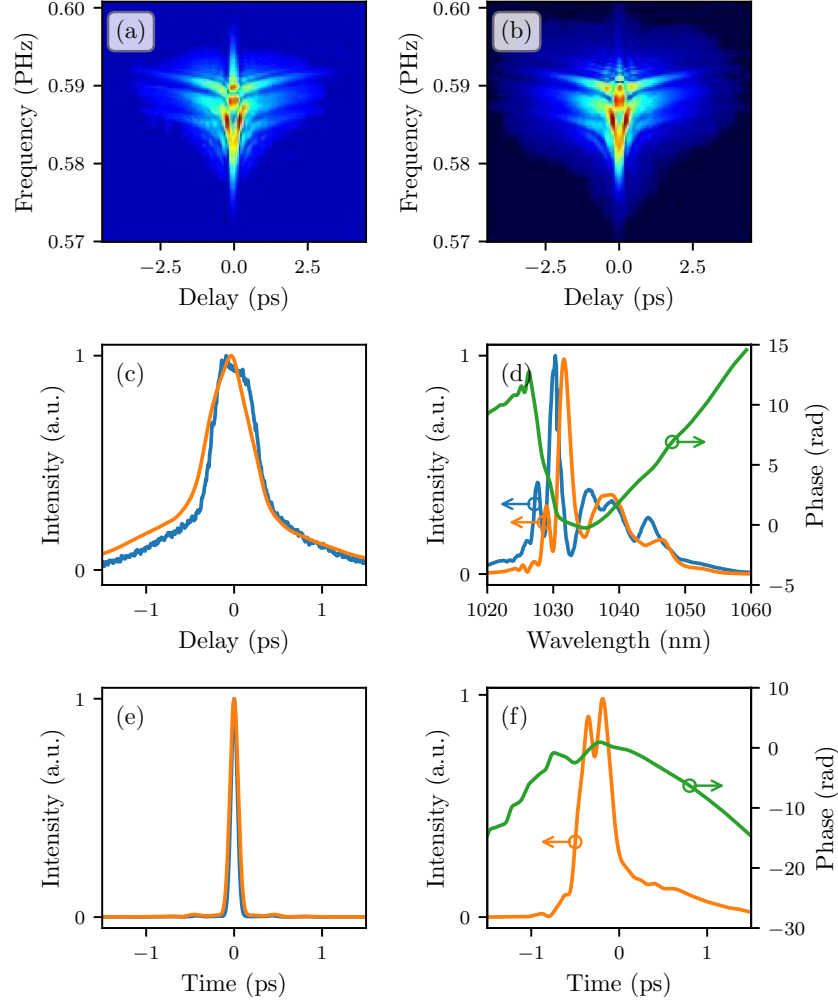


Figure 3.6: SHG-FROG retrieval of the spectrally broadened and uncompressed pulse after the HCF, at a filling pressure of 20 bar and pulse energy of $4.0 \mu\text{J}$. Reference Figure 2.25 for an SHG-FROG retrieval of the input pulse. (a)/(b) measured/retrieved SHG-FROG spectrograms; linear color map. The main features are captured, but FROG trace error (G) is significant at 2.7%. This error is the RMS difference between the measured and retrieved spectrograms, divided by the number of data points. (c) Comparison of the SHG-FROG delay marginals (orange) and an independent SHG based intensity autocorrelation (blue). (d) Retrieved spectrum (orange) and phase (green), compared with the spectral intensity as measured by an OSA (blue). (e) Transform limited pulses for the retrieved (orange) and OSA measured (blue) spectral intensities; the FWHM durations are 101 fs and 80 fs respectively. (f) Retrieved temporal intensity (orange) and phase (green); the FWHM duration is 399 fs.

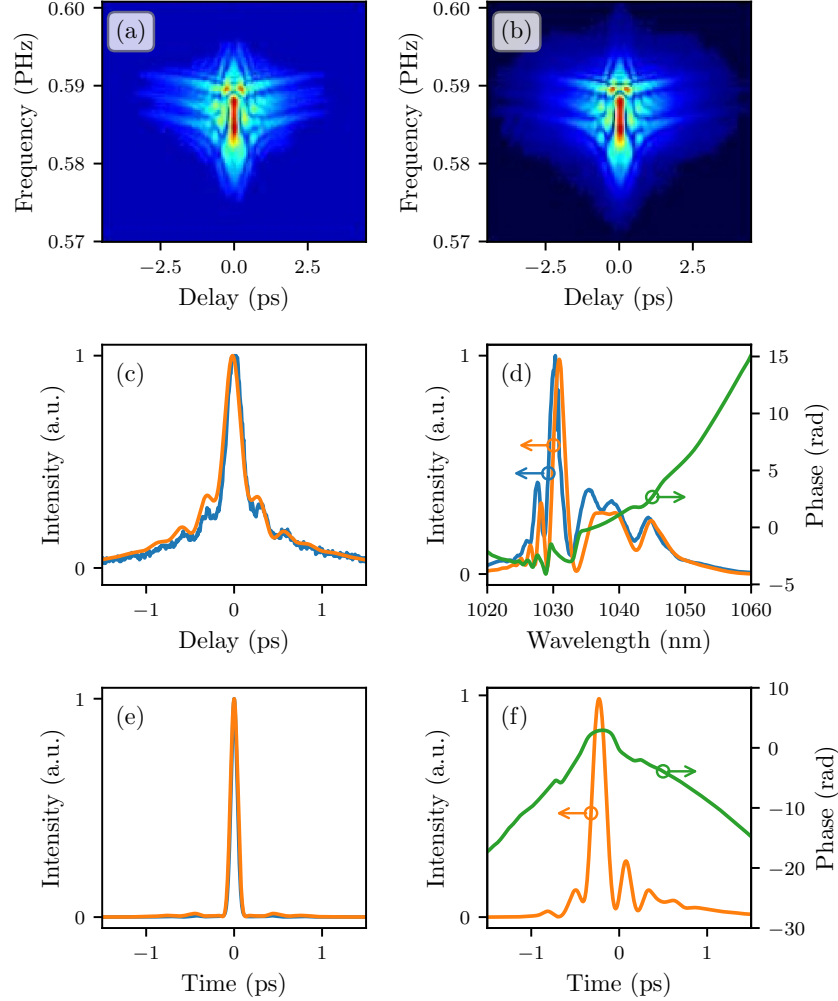


Figure 3.7: SHG-FROG retrieval of the spectrally broadened and compressed pulse after the HCF, at a filling pressure of 20 bar and pulse energy of $4.0 \mu\text{J}$. Reference Figure 2.25 for an SHG-FROG retrieval of the input pulse. (a)/(b) measured/retrieved SHG-FROG spectrograms; linear color map. The retrieval is excellent, with a FROG trace error (G) of only 0.5%. This error is the RMS difference between the measured and retrieved spectrograms, divided by the number of data points. (c) Comparison of the SHG-FROG delay marginals (orange) and an independent SHG based intensity autocorrelation (blue). (d) Retrieved spectrum (orange) and phase (green), compared with the spectral intensity as measured by an OSA (blue). (e) Transform limited pulses for the retrieved (orange) and OSA measured (blue) spectral intensities; the FWHM durations are 92 fs and 80 fs respectively. (f) Retrieved temporal intensity (orange) and phase (green); the FWHM duration is 174 fs.

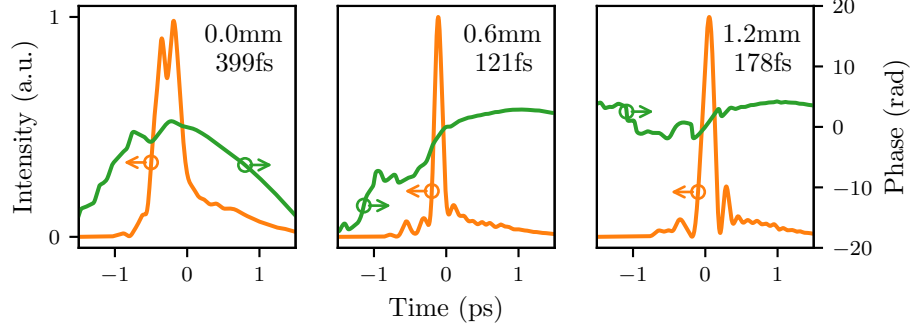


Figure 3.8: Numerical pulse compression of the SHG-FROG retrieved pulse after the HCF. See Figure 3.6 for more information regarding the retrieval. We have applied the theoretical phase imparted by a 12501/mm grating pair, at an incident angle of 40.1 deg, with 1030 nm being the zero-phase reference. The temporal intensity (orange) and phase (green) of the pulse is shown as the grating separation (normal distance) is increased. The minimum pulse duration occurs at a grating separation of 0.6 mm, afterwards the pulse is *over compressed* and the duration increases.

3.1.5 Simulations

Starting with initial conditions derived from the SHG-FROG retrieval after the PC-rod (Figure 2.25), we simulated propagation in 41 cm of HCF. The HCF, gas, and pulse parameters mirror those in the final spectral broadening trial. Namely, we used the perturbative MS dispersion model for HCF with the geometry listed in Figure 3.1, filled with argon to 20 bar, and assumed an input pulse energy of 4.0 μJ . After normalizing for energy, the peak power of the initial pulse is $P_0 = 12.6 \text{ MW}$. The nonlinear parameter, $\gamma = 5.0 \times 10^{-6} (\text{mW})^{-1}$, is found as described in Section 3.1.1.

Importantly, these values for P_0 and γ are treated as initial conditions. Based on our analysis of higher-order nonlinear effects (Appendix D), we determined that our pulse is still in the first-order, SPM-dominated regime of propagation. Given SPM is strongly dependent on γ and P_0 , agreement between simulated and experimental results hinges on the accurate determination of these values. Following the method of Emaury et al. [10], the simulations were run iteratively, with the product γP_0 used as a free parameter. After reducing γP_0 to 17% of its initial value (a multiplication factor of 0.17), the simulated bandwidth was comparable to that measured with an OSA. The results are compiled in Figure 3.9.

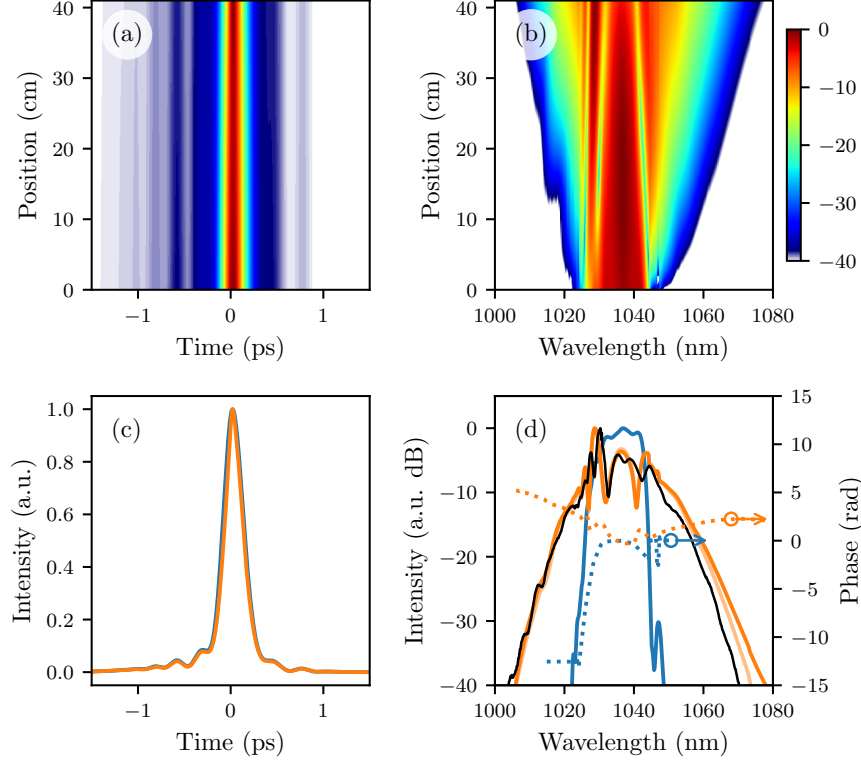


Figure 3.9: Simulated pulse propagation in 41 cm of HCF, starting with initial conditions derived from the SHG-FROG retrieval after the PC-rod (Figure 2.25). The HCF, gas, and pulse parameters mirror those in the third spectral broadening trial. Specifically, we used the perturbative MS dispersion model for HCF with the geometry listed in Figure 3.1, filled with argon to 20 bar, and assumed an input pulse energy of $4.0 \mu\text{J}$. Here, the factor γP_0 is used as a free parameter and reduced to 17% of its initial value ($\gamma P_{0\text{ initial}} = 5.0 \times 10^{-6} \text{ m}^{-1} \text{ W}^{-1} \times 12.6 \text{ MW}$, and we multiply by a factor of 0.17). (a) Temporal evolution of the pulse; color map is linear in intensity. As expected for weak dispersion, the pulse envelope is essentially unchanged. (b) Spectral evolution of the pulse, showing the development of an asymmetric spectrum like that measured after the HCF (Figure 3.5). (c) Temporal slice at 0 cm (blue) and 41 cm (orange), showing the very minor effect of dispersion. It must be pointed out that the simulations do not capture the temporal broadening and structure seen in the uncompressed FROG retrieval (Figure 3.6). (d) Spectral slice at 0 cm (blue) and 41 cm (orange bold), along with the respective spectral phase (blue dotted and orange dotted). The simulated output bears a striking resemblance to the OSA measured spectrum (black). It appears the spectral modulations at short wavelengths correspond to a kink in the initial phase. We have also plotted the simulated output for the case of zero dispersion (orange light), to affirm the negligible role dispersion plays.

Remarkably, the measured asymmetric spectral broadening and modulations are also well reproduced. As suspected, initial spectral phase plays an important role, with the strongest modulations occurring near regions with second or higher order phase components. To rule out the influence of dispersion—particularly resonant behavior in the perturbative MS model—we ran the simulations a third time with dispersion completely neglected. As one can see in Figure 3.9-d, dispersion has only a minor influence at the longest wavelengths, reaffirming our previous conclusions that pulse evolution remains in a predominantly SPM dominated regime.

Is it reasonable to reduce γP_0 to 17% of its initial value? We believe so for two reasons: firstly, as stated in Section 3.1.1, experimental values of γ vary over an order of magnitude; secondly, we naively divide the average output power by the repetition rate to derive the pulse energy, when in fact, a significant fraction of the average power may exist between pulses in the form of ASE. Disentangling the individual contribution of γ and P_0 to their product γP_0 should be the near-term focus, so that simulations can accurately guide development of the source. Given the excellent qualitative agreement between the simulated and experimental spectra, determining only one of γ or P_0 is necessary and the other can be found through fitting. We propose it is easiest to find P_0 , which will require assessing the magnitude of the ASE contribution, perhaps using an AOM gating scheme as done by Pavlov et al. [33].

There is also preliminary evidence for ASE in the SHG-FROG-retrieved spectrum after the rod, reproduced here in Figure 3.10. The retrieved spectrum largely overlaps that measured by the OSA, and generally captures the structure as well, but there is a section of longer-wavelength spectral components not captured by the retrieval—we believe this could be ASE. The rationale follows from the difference between measurement methods: SHG-FROG inherently suppresses constant, low-level, background noise (e.g. ASE) via temporal gating and nonlinear discrimination [18]; in contrast, the OSA continually integrates and averages the spectral power. Given the period between pulses is on the order of 10^5 times the duration of the pulse, a constant low-level signal between pulses can integrate to a substantial component of the total spectral power. The integrated spectrum would not be representative of the pulse spectrum if this were the case, and the SHG-FROG retrieved spectrum would be more accurate.

Of course, the SHG-FROG spectrum is not directly measured, but rather retrieved with an inversion algorithm; therefore, a more thorough investigation is required to determine the statistical significance of the missing long-wavelength spectral components.

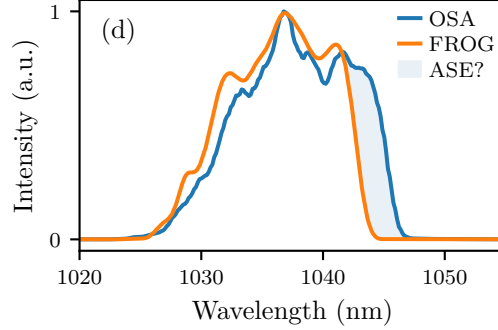


Figure 3.10: Potential signs of amplified spontaneous emission (ASE) after the PC-rod amplifier. The SHG-FROG retrieval lacks the shaded portion of the spectrum captured by the optical spectrum analyzer (OSA). We believe this could be ASE, as SHG-FROG inherently suppresses ASE via temporal gating and nonlinear discrimination [18]. This is in contrast to the OSA, which continually integrates and averages the spectral power.

3.2 Resonant Dispersive Wave Generation

For completeness, we will give a brief conceptual overview of resonant dispersive wave (RDW) generation, which is ultimately the final stage of the vacuum ultraviolet (VUV) source. The phenomenon has already been seen in the simulations (Figure D.1, Figure 3.9) and is a solitonic property allowed by third or higher order dispersion. Key to understanding the process is the term “resonant.” Dispersive waves coexist with solitons at all wavelengths guided by the HCF, however, at certain wavelengths the soliton and dispersive wave are phase-matched, or “resonantly coupled”. The particular wavelengths are determined by plotting the propagation constant for the dispersive wave, $\beta(\omega)$, and the soliton, $\beta_s(\omega)$, and finding their intercepts. To be clear, β is the real component of the wavevector k , and they are equal if loss is ignored. A generic example of the graphical method is provided in Figure 3.11. The soliton propagation constant is intrinsically linear in frequency, since all spectral components must share the same group velocity for the pulse to propagate unchanged. Explicitly,

$$\beta_s(\omega) = \beta_0 \Big|_{\omega_s} + \beta_1 \Big|_{\omega_s} (\omega - \omega_s) + \frac{1}{2} P_s \gamma \quad (3.4)$$

where ω_s is the soliton frequency, and P_s is the peak power of the fundamental soliton. The last term in Equation 3.4 accounts for nonlinear

phase, and is positive in most ordinary materials. In a sense, it is a byproduct of the natural requirement for nonlinear phase to cancel the effect of higher order dispersion and leave only the linear term. In the anomalous dispersion regime this small offset has a significant impact, since the linear (i.e. not nonlinear) dispersion relation up to second-order is concave-down in frequency. From a purely geometric perspective, we see that the linear and solitonic dispersion relations will never intersect, and therefore, there are no points where phase matching occurs. To have an intersection there must be third or higher order components in the linear dispersion relation that introduce an inflection point. In fact, we are guaranteed at least one phase-matched wavelength if the highest order dispersion term is odd.

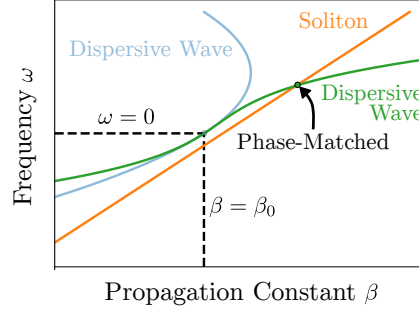


Figure 3.11: Generic illustration of phase matching between solitonic and dispersive radiation. The pulse envelope evolves such that all spectral components of the soliton share the same group velocity (constant slope). For a positive nonlinear index, solitonic radiation travels more slowly, hence the increased propagation constant at $\omega = 0$. The linear dispersion relation up to second order (blue, $\beta_2 < 0$) and third order (green, $\beta_3 > 0$) show the necessity of higher order dispersion for phase-matching in the anomalous regime.

As an example, let us predict the wavelength of the dispersive wave in the fission simulations (Figure D.6). The soliton propagation constant just discussed technically applies to fundamental solitons, but can be extended to higher order solitons by considering only the self compression point; $P_s \rightarrow P_c$, where P_c is the self compressed pulse peak power. Empirically, it has been found that $P_c \simeq 4.6NP_0$, where N is the soliton order, and P_0 is the initial peak power [47]. For our simulation $P_c \simeq 170$ MW. The phase-matched wavelength is best found by looking for zero-crossings of the propagation constant mismatch $\Delta\beta(\omega) = \beta(\omega) - \beta_s(\omega)$. In Figure 3.12 the zero-crossing occurs at 607 nm, which is comparable to where we see the dispersive wave form in the simulations (~ 620 nm). The difference results

from our analytic approximation of the higher order soliton as a fundamental soliton at the compression point.

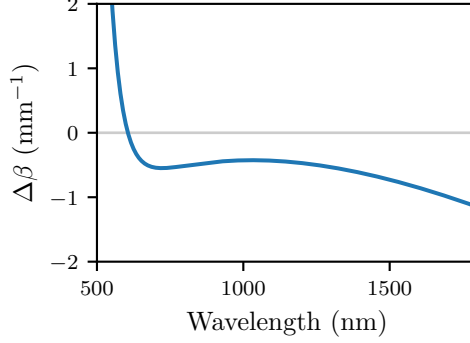


Figure 3.12: Propagation constant mismatch, $\beta(\omega) - \beta_s(\omega)$, between dispersive and solitonic radiation for the fission simulation in Figure D.6 at the self-compression point. The zero-crossing at 607 nm indicates the phase-matched wavelength where RDW generation should occur, and is comparable to the value found in the simulations (620 nm). The slight difference results from our analytic approximation of the pulse as a fundamental soliton at the self-compression point.

It should be pointed out that the location of the RDW coincides with the first cladding resonance of the HCF, captured by the perturbative MS dispersion model (Figure 3.2-a). However, this model does not account for losses, leading to artificially impenetrable barriers in the dispersion profile. As such, our RDW will always be bound to the resonance and it would seem that generation in the VUV range is not possible. In reality these dispersion barriers can be overcome. To get a sense of the shortest wavelength RDW it is better to use the standard MS model (Figure 3.2-b). Further, we ultimately would like to produce 30 fs, $\sim 3 \mu\text{J}$, pulses after the spectral broadening stage, so we will take the liberty of using such a pulse for this conceptual section. With these assumptions, varying the pressure of the filling gas allows for RDWs from $\sim 6 \text{ eV}$ to $\sim 10 \text{ eV}$ for argon and from $\sim 18 \text{ eV}$ to $\sim 23 \text{ eV}$ for helium (Figure 3.13). Phase-matching again happens about a resonance – in this case the first absorption resonance of the gas – except here we can tune the width of the resonance by varying the pressure. Therefore, if the pulse after the spectral broadening stage is sufficiently intense and has spectral components reaching the range of phase-matched wavelengths, VUV RDWs can be generated.

At present, experiments using helium-filled, fibers have reached photon energies of ~ 11 eV, limited by ionization, and higher-order linear and nonlinear effects at sub-femtosecond pulse durations [10, 47, 48].

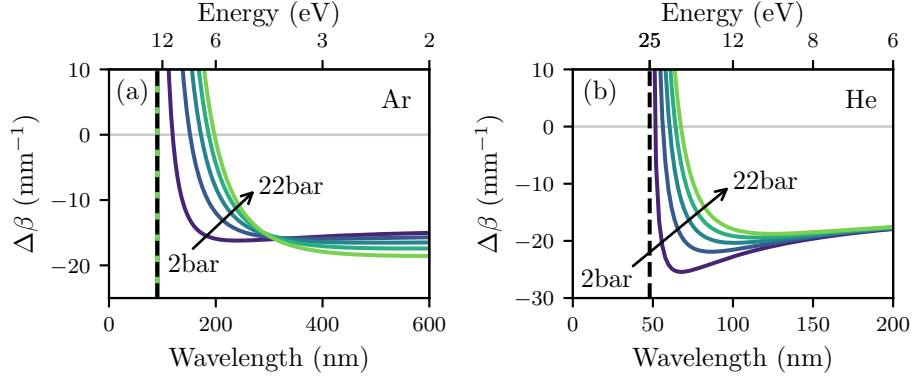


Figure 3.13: Propagation constant mismatch, $\beta(\omega) - \beta_s(\omega)$, between dispersive and solitonic radiation for a 30 fs, 3 μ J, sech pulse propagating in HCF with a 30 μ m core diameter. The filling gas pressure is varied from 2 bar to 22 bar in 4 bar steps. (a) Filling species is argon, which has an absorption resonance at 90 nm, indicated by the dotted black line. The phase-matched wavelength can be pressure-tuned from ~ 6 eV to ~ 10 eV. (b) Filling species is helium, which has an absorption resonance at 49 nm, indicated by the dotted black line. The phase-matched wavelength can be pressure-tuned from ~ 18 eV to ~ 23 eV.

Chapter 4

Peripheral Equipment

4.1 System Interlock

4.1.1 Design Principles

The interlock is in place to prevent self-induced damage to the amplifier chain. Gain media in all stages of amplification will store energy if the rate of absorption exceeds the rate of emission. This rate imbalance is particularly severe when the input signal to the amplifier is blocked, and hence, there is no source to effect stimulated emission other than the amplification of spontaneous emission (ASE). With multi-watt level pump powers in the last two stages of amplification, multi-millijoule energies build within the $\sim$ millisecond upper-state lifetime of the ytterbium gain medium and can be released spontaneously in an avalanche-like “giant pulse”. These giant pulses can cause irreparable damage to the amplifier and other optical components and must be avoided.

The interlock monitors signal levels between amplifier stages, and if they fall out of bounds, disables the requisite pump diode driver before avalanching occurs. Practically, interlocking should occur on a 100 μ s time scale. A signal is within bounds if the amplitude is great enough to be detected, and the number of pulses in a given time period stays within a prescribed range. Counting of the pulses and logic for enabling/disabling the pump diode is embedded in an FPGA (field programmable gate array). A system schematic is shown in Figure 4.1.

4.1.2 Hardware Implementation

Optically, the locations of the three photodetectors can be found in Figures 2.1 (oscillator), 2.16 (PCF preamplifier) and 2.21 (PC-rod amplifier). We used a home-built photodetector with a $\sim$ 1 GHz bandwidth after the oscillator, since it also serves as a clock-source for the pulse-picking AOM, requiring we preserve the 10th harmonic (380 MHz). The other two detectors (Thorlabs PDA10A) have a 150 MHz bandwidth, which is sufficient for preserving the fundamental harmonic, and hence, counting pulses.

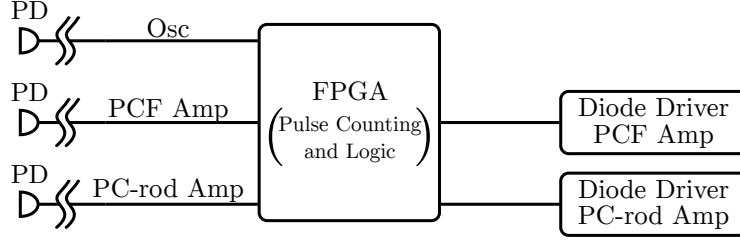


Figure 4.1: Schematic of the system interlock. The oscillator, PCF preamplifier and PC-rod amplifier outputs are sampled and the pulses are detected on photo-diodes (PD). Digital inputs on the FPGA (Field programmable gate array) “square” the analog pulse train, and pulses exceeding the logic-high threshold are counted. When the number of pulses counted falls outside the accepted range, the FPGA will disable the requisite pump diode driver.

Also due to its dual purpose, the photodetected signal after the oscillator takes a circuitous route through the pulse-picker (AlphaNov “Pulse Picker”) before entering the FPGA (Figure 4.2). The input comparator of the pulse-picker “squares” the signal at a threshold level of 1.2 V—that is, the smoothly varying analog signal is converted to a digital rectangular wave. Inputs above 1.2 V are output High (5 V) and those below are output Low (0 V). We feed this digital signal into the FPGA for counting. The photodetected signals after the PCF and PC-rod amplifiers are fed directly into the digital IO (input/output) of the FPGA, which similarly squares the analog signal at the 3.3 V TTL logic-high threshold.

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The diode drivers are enabled/disabled by toggling the respective digital output on the FPGA. For the PCF preamplifier we use an nLight PPS-10-21 diode driver (re-badged Vuemetrix MV-21-01) with an purpose-built interlock input. Output low on the FPGA enables this driver. In similar systems at QMI [27] the fall-time of this driver has been shown to be far less than 1 ms. For the PC-rod amplifier we use an Ostech ls11-la20v30-t12783-

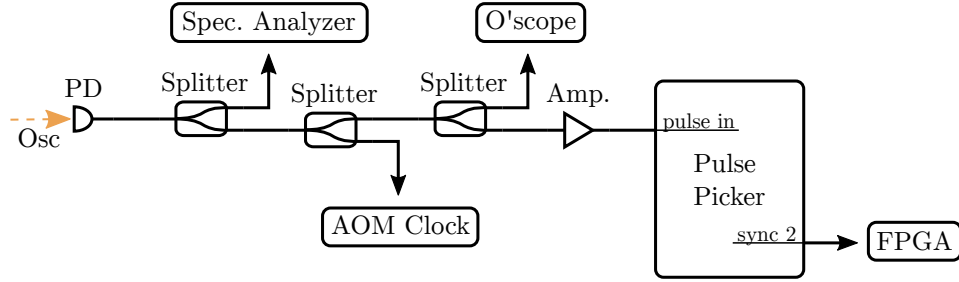


Figure 4.2: Detailed schematic of the signal path between the oscillator photo-diode (PD) and the FPGA. Several diagnostics and the acousto-optic-modulator (AOM) clock/gate signals are derived from the same photodetected pulse train. The final signal input to the FPGA is already a rectangular wave, with 5 V TTL levels.

noise-v0-467-726 diode driver. There is a substantial delay ($\gg 1$ ms) before the output is disable using the purpose-built interlock input on this driver. Therefore, we instead toggle the modulation input to enable (logic high) or disable (logic low) the output. The fall-time of the driver output with respect to the disable signal from the FPGA is shown in Figure 4.3. Including the $20\ \mu\text{s}$ integration period of the FPGA controller/pulse-counter, the system is shut down in $< 200\ \mu\text{s}$ after an interlock event.

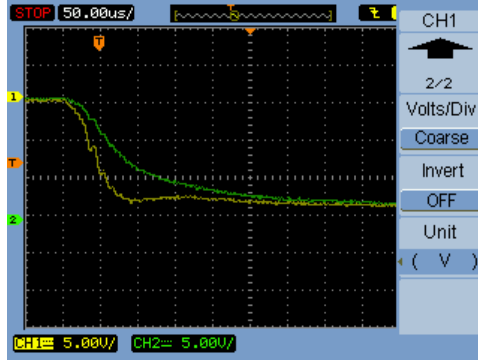


Figure 4.3: Measured fall-time of the Ostec diode driver output (green, $< 100\ \mu\text{s}$) when the modulation input (yellow) is pulled low. The FPGA controller enables/disables the driver via this input. Note there is up to a $20\ \mu\text{s}$ delay between an interlock event and the high-to-low transition of the FPGA output; this allows for counting of the pulse train.

4.1.3 Firmware Implementation

At the heart of the interlock is the FPGA controller. Both pulse counting and logic for enabling/disabling the diode drivers are implemented on this device. Programming the FPGA is very much like writing firmware for a microcontroller, but the design is ultimately implemented in hardware, as a linked array of logical gates. Done in this way, it is practical to design a counter that operates beyond the clock-frequency of the controller. A block-diagram for PC-rod amplifier branch of the firmware is shown in Figure 4.1. The design can largely be broken into two halves: a pulse counter/comparator that determines the repetition rate and checks that it is within the predefined bounds, and a truth table that enables/disables the diode driver based on all other concurrent inputs. Both the PCF preamplifier and oscillator branches share similar topologies, with only the truth table being different. We will discuss the counter/comparator first, then the truth tables for the various branches.

Pulse counting is done in two stages. The first stage, an asynchronous 5-bit Gray counter, acts as a prescaler to reduce the input repetition rate below one-half of the on-board sampling clock frequency (50 MHz). IO circuitry within the FPGA has a bandwidth exceeding 300 MHz, we are only limited by the Nyquist theorem if we take a conventional sampling approach to measuring the repetition rate. Here, the pulse train is the clock for the 5-bit Gray counter—it is not sampled. We use a Gray counter since only one bit changes per clock cycle, which ensures we are at most off-by-one if the Gray counter is in the process of incrementing when it is read by the sampler. The off-by-one error is not cumulative, as the sampler compares the current and previous counts to generate a “delta-count”, or difference between their values. The second counting stage occurs on each rising edge of the sample clock, when a delta-count is generated and added to the 16-bit “accumulator”.

At a slower frequency—the 50 kHz “Status Clock” rate derived through a phase-locked-loop (PLL) from the sampling clock—the counter total is read, cleared and compared to the minimum and maximum allowable value. If it is within the allowed bounds, the “PC-rod Status” bit is asserted. Simultaneously, the rest of the truth table is updated and the PC-rod diode driver is either enabled or disabled accordingly. The status clock period (20 μ s) was chosen to be much shorter than the upper-state lifetime of the gain medium, but also long enough to allow reasonable counting statistics that are robust to a few missed pulses. As implemented, the accumulator total is read, zeroed and halted over the logic-high half-cycle of the status clock, leaving the 10 μ s logic-low half-cycle to count pulses.

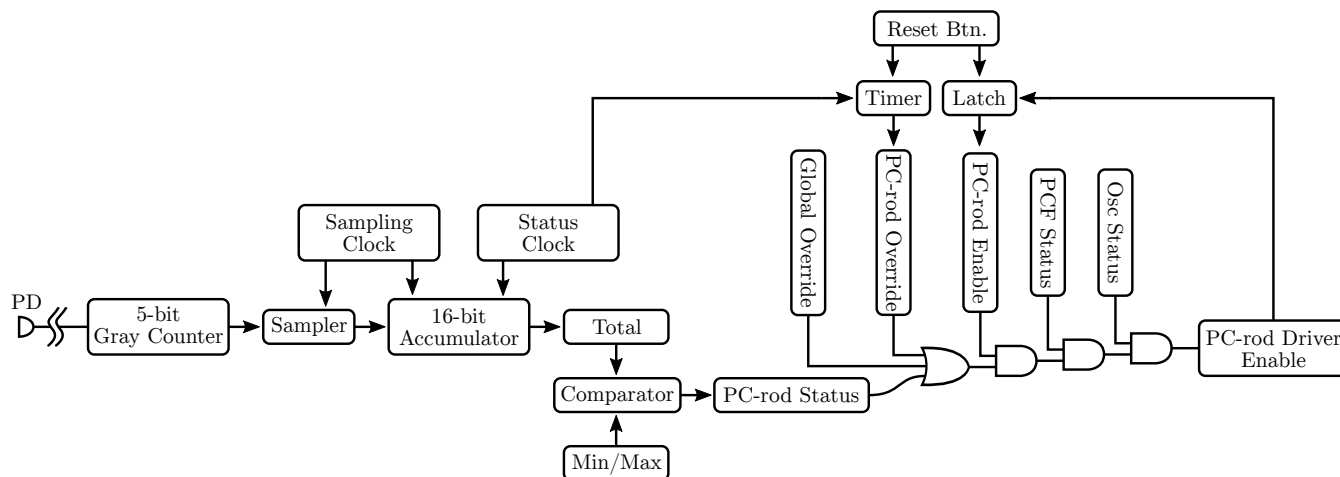


Figure 4.4: Block diagram of the interlock firmware, which is ultimately implemented in hardware on the FPGA. Here we show only the PC-rod amplifier branch of the controller; the oscillator and PCF preamplifier branches have similar topologies. The left half of the structure, up to the “PC-rod Status” bit, is a pulse counter/comparator—if the total count falls within a predetermined range the bit is asserted. Use of an asynchronous 5-bit Gray counter/prescaler allows a 50 MHz sampling clock to count > 200 MHz pulse repetition rates. The right half of the structure composes a truth table, which enables/disables the diode driver based on all other prerequisites. This table is updated at the status clock rate (50 kHz), which is derived through a phase-locked-loop (PLL) from the sampling clock.

The truth table for the PC-rod amplifier is best found by working backwards from the diode driver enable bit. If ‘yes’ is the answer to the following AND’d questions, the diode driver will be enabled:

- Is the oscillator OK?
- AND is the PCF amplifier OK?
- AND has the PC-Rod been enabled?
- AND

Is the PC-rod override still active?

OR is the global override enabled?

OR is the PC-rod OK?

The diode driver enable bit will not be asserted if either the oscillator or PCF preamplifier are off or outside their prescribed operating range. Once the preceding sections of the amplifier chain are operational, the reset button must be pressed to clear the shutdown-latch (PC-rod enable) and reset the startup timer (PC-rod override). The shutdown latch ensures the system remains off in the event of an intermittent failure, and the timer allows one approximately 10s to power on the PC-rod amplifier. After the timer expires, either the PC-rod amplifier must be operational, or the global override switch must be asserted. The logic is identical for the PCF preamplifier, aside from only requiring the oscillator to be operational. For the oscillator we need not worry about self-induced damage, and thus only need to update its status bit—no truth table required.

4.2 Gas Panel

4.2.1 Design Principles

The gas panel is designed for convenient evacuation and filling of the gas cells used in the nonlinear broadening and RDW generation stages, while at the optic bench. Evacuation is necessary to remove atmospheric gases that are easily ionized, or cause other unintended effects, such as Raman scattering. The purging (evacuation and filling to low pressure) procedure we followed can be found in Section 3.1.3. After purging, the gas panel is used to adjust the pressure of the filling gas. Rapid adjustment is particularly important for tuning the phase-matched wavelength in the RDW stage, hence, a self-venting regulator is incorporated into the panel (self venting allows for negative pressure changes). The panel also allows for unregulated filling to the supply pressure, if pressures greater than 34 bar are required (60 bar is the safely permissible limit), and for generating pressure gradients across the fiber.

4.2.2 Implementation

A schematic of the gas panel is shown in Figure 4.5. A pre-regulated gas supply enters through QV1 (quarter-turn valve) and the supply-pressure is monitored on analog gauge G. Pre-regulation at the cylinder is done both as a safety measure, and to minimize pressure droop as the cylinder empties. With RV1 and RV2 (regulating needle valves) closed, the gas cell can be filled either directly from the supply line through RV3, or with additional regulation through PR (pressure regulator) and RV4. In its current configuration, the maximum filling pressure is 24 bar, imposed by the safety relief valve (SRV). Otherwise, the pressure regulator (PR) and gauges operate to 34 bar, and can tolerate pressures up to 60 bar.

Each gas cell (two per fiber stage) is connected to the panel by a coiled 1/16 inch tube to allow some flexibility in its position on the table. The additional valve QV2 can be closed if one wants to create a pressure gradient across the fiber by externally evacuating the respective gas cell. Currently, our gas cells are connected together with a 1/8 inch diameter tube; therefore, we leave QV2 closed and connect only one gas cell to the panel.

For evacuation, the gas supply is shut off with QV1, and the pressure is initially reduced by venting through RV2. The coarse depressurization is done to avoid shocking the vacuum pump. Once the system is near atmospheric pressure, RV2 is closed and RV1 is opened.

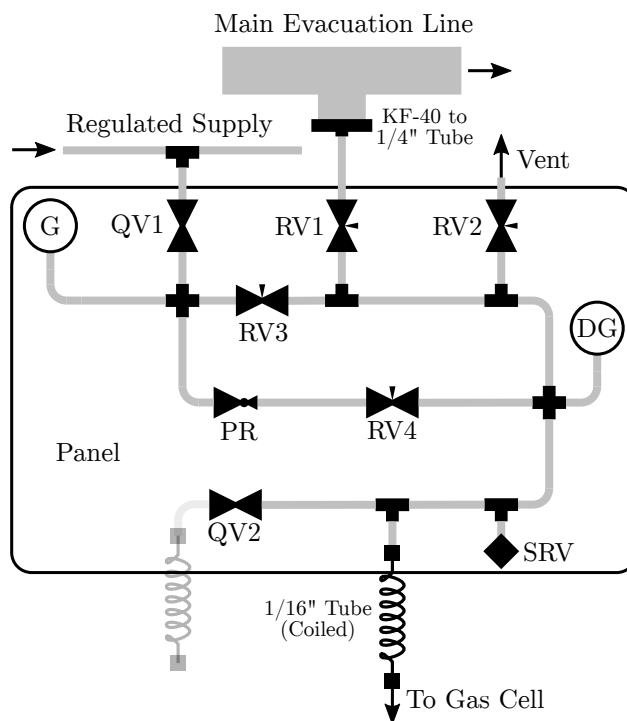


Figure 4.5: Schematic of the gas panel used to fill and evacuate the gas cells. QV, quarter valve; RV, regulating needle valve; PR, pressure regulator; G pressure gauge; DG, digital pressure gauge; SRV, safety relief valve. The gas cells can be filled at the supply pressure (max 60 bar), or to 34 bar with secondary regulation. Currently, however, the safety relief value limits the maximum pressure to 24 bar. The coiled tube after QV2 is not in use, as the cells are connected externally with a 1/8 inch tube. Figure adapted from one provided by John Travers and Federico Belli (Heriot Watt University).

Chapter 5

Summary and Conclusions

A tunable vacuum ultraviolet light source for time- and angle-resolved photo emission spectroscopy (TR-ARPES) is being developed. The design consists of three stages: pump pulse generation with a Yb-fiber based source, nonlinear broadening and subsequent temporal compression in gas-filled hollow-core fiber (HCF), and resonant dispersive wave (RDW) generation in a succeeding HCF. The ultimate goal is generating sub-30 fs pulses with photon energies between 5 eV and 11 eV, with variable repetition rate from 300 kHz to 20 MHz.

The author of this thesis, in collaboration with those listed in the preface, has worked on the ground-up construction and characterization of the first two of three stages, with the interim goal of producing 30 fs pulses with sufficient energy (μ W level) for RDW generation in the third and final stage.

The first stage required E. Kelleher, E. Ostroumov and the author to design and assemble a chirped pulse amplification (CPA) system. At a repetition rate of 4.9 MHz, the retrieved pulse duration using second harmonic generation (SHG) frequency resolved optical gating (FROG) is 264 fs, and the pulse energy is 4.5 μ J, based on dividing the average power (22 W) by the repetition rate.

The second stage was constructed by E. Kelleher, E. Ostroumov, F. Belli and the author, using gas cells provided by J. Travers. After compression with a grating pair, the author has measured 174 fs, 4.0 μ J pulses with SHG-FROG. Mechanical limitations prevented full compression of the pulse. Numerical compression of the same SHG-FROG retrieved pulse reduced the duration to 121 fs, suggesting immediate improvements with minor changes to our compression scheme.

Nonetheless, we have not attained the interim goal of 30 fs. From a comparison of experimental and simulated spectra carried out by the author, it appears the initially assumed value of either γ (the nonlinear parameter), or P_0 (the peak power of the pulse), or both, exceed their true value. To meet the goal, the author suggests carefully assessing the contribution of amplified spontaneous emission (ASE) to the calculated pulse energy, so that P_0 may be accurately determined. Once known, γ can be found through

fitting. Simulations may then be used to determine if modifications to the first stage are required, or if simply using a longer length of HCF in the second stage is sufficient.

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Appendix A

Effect of GVD on Pulse Duration

A.1 Fourier Transform of a Gaussian Pulse

Assuming a Gaussian pulse profile makes for a convenient analytical derivation of the effects of GVD. We start with a transform-limited (TL) Gaussian pulse, and transform it to the frequency domain where it is easy to add spectral phase. The temporal profile of such a pulse is given as

$$E(t) = Ae^{-2\ln(2)\frac{t^2}{\tau^2}}e^{-i\omega_0 t} \quad (\text{A.1})$$

where A is a normalization constant, τ is the intensity FWHM duration, and ω_0 is the central carrier frequency. In the frequency domain we have

$$E(\omega) = \int_{-\infty}^{\infty} Ae^{-2\ln(2)\frac{t^2}{\tau^2}}e^{-i\omega_0 t}e^{i\omega t} dt \quad (\text{A.2})$$

$$= \int_{-\infty}^{\infty} Ae^{-\gamma t^2}e^{i(\omega-\omega_0)t} dt \quad (\text{A.3})$$

$$= Ae^{\gamma\left[\frac{i(\omega-\omega_0)}{2\gamma}\right]^2} \int_{-\infty}^{\infty} e^{-\gamma\left[t-\frac{i(\omega-\omega_0)}{2\gamma}\right]^2} dt \quad (\text{A.4})$$

$$= Ae^{-\frac{(\omega-\omega_0)^2}{4\gamma}} \sqrt{\frac{\pi}{\gamma}} \int_{-\infty}^{\infty} e^{-\pi u^2} du \quad (\text{A.5})$$

$$= \sqrt{\frac{\pi}{\gamma}} Ae^{-\frac{(\omega-\omega_0)^2}{4\gamma}} = \sqrt{\frac{\pi\tau^2}{2\ln(2)}} Ae^{-\frac{\tau^2}{8\ln(2)}(\omega-\omega_0)^2} \quad (\text{A.6})$$

$$E(f) = Be^{-2\ln(2)\frac{(f-f_0)^2}{\Delta\nu^2}}, \quad \Delta\nu = \frac{2\ln(2)}{\pi\tau} \quad (\text{A.7})$$

where the spectral intensity FWHM bandwidth is $\Delta\nu$. The final integral is a normalized error function, which evaluates to 1.

A.2 Adding Spectral Phase

It is a simple matter to add spectral phase in the frequency domain. The amount accumulated is proportional to the propagation distance and the refractive index of the material. A full functional form of the refractive index, and hence the wavevector, is often unknown; we will take a Taylor series approximation about the central carrier frequency. The spectral phase factor and its expansion are then

$$e^{i\phi(\omega)} = e^{ikL} = e^{in(\omega)k_0L} = e^{i\left(\beta_0 + \beta_1(\omega - \omega_0) + \frac{\beta_2}{2}(\omega - \omega_0)^2 + \dots\right)L} \quad (\text{A.8})$$

where k and k_0 are the wavevectors in the medium and vacuum respectively, $n(\omega)$ is the frequency dependent refractive index, L is the propagation distance, and we have truncated the expansion at the second order term, β_2 , which represents the GVD. The lower order component β_0 is a constant phase offset, and β_1 is the reciprocal of group velocity. We add the phase by multiplying, then group the second order terms and rearrange to get the final spectrum of the pulse in a form that can be transformed back to the time domain by inspection. The final intensity FWHM, τ_L , after propagating a distance L through the dispersive material is found as follows:

$$E_L(\omega) = B e^{-\frac{\tau^2}{8\ln(2)}(\omega - \omega_0)^2} \cdot e^{i\left(\beta_0 + \beta_1(\omega - \omega_0) + \frac{\beta_2}{2}(\omega - \omega_0)^2\right)L} \quad (\text{A.9})$$

$$= B e^{-\frac{1}{8\ln(2)}(\tau^2 - i4\ln(2)\beta_2L)(\omega - \omega_0)^2} e^{i\beta_0L} e^{i\beta_1(\omega - \omega_0)L} \quad (\text{A.10})$$

$$\mathcal{F}^{-1} \Downarrow \quad (\text{A.11})$$

$$E_L(t) = A e^{-2\ln(2)\frac{(t - \beta_1L)^2}{(\tau^2 - i4\ln(2)\beta_2L)}} e^{i(\beta_0L - \omega_0t)} \quad (\text{A.12})$$

$$I_L(t) = A e^{-4\ln(2)\frac{(t - \beta_1L)^2}{\tau^2 + \left(\frac{4\ln(2)\beta_2L}{\tau}\right)^2}} \quad (\text{A.13})$$

$$= A e^{-4\ln(2)\frac{(t - \beta_1L)^2}{\tau_L^2}}, \quad \tau_L = \sqrt{\tau^2 + \left(\frac{4\ln(2)\beta_2L}{\tau}\right)^2}. \quad (\text{A.14})$$

If the initial *Gaussian* pulse is not transform-limited, then it must already have acquired some GDD. We may decompose the initial duration τ_0 into TL and GDD components, such that

$$\tau_0^2 = \tau_{\text{TL}}^2 + \tau_{\text{GDD}}^2. \quad (\text{A.15})$$

After isolating the initial GDD component, we simply add it in quadrature to the final pulse duration, giving us

$$\tau_L = \sqrt{\tau_{\text{TL}}^2 + \tau_{\text{GDD}}^2 + \left(\frac{4 \ln(2) \beta_2 L}{\tau_{\text{TL}}}\right)^2} \quad (\text{A.16})$$

$$= \sqrt{\tau_0^2 + \left(\frac{4 \ln(2) \beta_2 L}{\tau_{\text{TL}}}\right)^2}. \quad (\text{A.17})$$

Note that we can redefine $(4 \ln(2)/\tau_{\text{TL}})^2$ as the spectral intensity bandwidth (Eqn. A.6), and therefore the stretched pulse duration can also be written as

$$\tau_L = \sqrt{\tau_0^2 + (\Delta\omega \cdot \text{GDD}_{\text{Str}})^2}. \quad (\text{A.18})$$

Appendix B

Discussion of Spherical Aberrations

B.1 Spherical Aberrations of a Singlet

The spherical aberration of a thin lens is given by Smith [43] as

$$\Delta\left(\frac{1}{v}\right) = \left(\frac{1}{v} - \frac{1}{v_0}\right) = \frac{1}{n(n-1)} \frac{h^2}{8f^3} \left[\frac{n+2}{n-1} s^2 + 2(2n+2)sp \cdots \right. \\ \left. + (3n+2)(n-1)^2 p^2 + \frac{n^3}{n-1} \right] \quad (\text{B.1})$$

$$= \frac{h^2}{f^3} [As^2 + Bsp + Cp^2 + D] = \frac{h^2}{f^3} \cdot S \quad (\text{B.2})$$

where v is the distance between lens and image, and v_0 is the distance between the lens and the point where the marginal ray intercepts the optic axis. The last equality reduces the expression to a ratio of powers between the semi-diameter of the entrance pupil h , and its designed focal length f , multiplied by a scaling factor S , which accounts for the lens' physical construction. The parameters s and p are the “shape” and “position” factors respectively, which are given as follows

$$s = \frac{1 - r_1/r_2}{1 + r_1/r_2}, \quad p = \frac{1 - u/v}{1 + u/v} \quad (\text{B.3})$$

with r_1 the radius of curvature of the object-face of the lens, and u the distance to the object. It is now only a matter of scaling to find the longitudinal and transverse spherical aberrations, expressly

$$\text{LSA} = v^2 \Delta\left(\frac{1}{v}\right), \quad \text{TSA} = hv \Delta\left(\frac{1}{v}\right), \quad (\text{B.4})$$

where it is assumed that the LSA is small, such that $v \simeq v_0$.

As an example, let's determine the LSA and TSA of a 40.0 mm focal length, plano-convex lens (Thorlabs LA1422), with a collimated beam incident on the convex face. This scenario is a simplified and reversed variation of collimating light from a fiber-pigtail, as discussed in section 2.6. Given the geometry of the problem, and the radius of curvature of the convex side (20.6 mm), we find $s = 1$ and $p = -1$. Then knowing the index of refraction for N-BK7 is 1.51 at 980 nm, we find that $S = 0.859$. For the beam diameter discussed in section 2.6 we have $h = 9.0$, providing the last parameter needed to find the spherical aberration $\Delta(1/v) = 1.09 \times 10^{-3} \text{ mm}^{-2}$. It follows that the longitudinal, and transverse spherical aberrations are 1.74 mm and 0.391 mm respectively. These results are supported by a Zemax simulation (shown in Figure B.1), which reports an LSA of 1.64 mm, and a TSA of 0.400 mm.

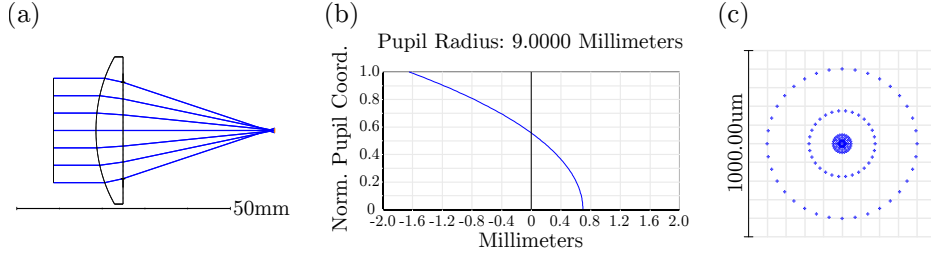


Figure B.1: Zemax simulation for the Thorlabs LA1422 spherical lens. The incoming beam is perfectly collimated and has a diameter of 18.0 mm. The image plane is located at the paraxial focus of the lens. (a) Optical layout of the simulation. Spherical aberrations are clearly present about the focus. (b) Longitudinal aberrations, shown as a distance from the image surface to where a zonal marginal ray crosses the optical axis. (c) Transverse aberrations, shown as a distribution of rays traced to the image plane.

B.2 Effect of Aberrations on Pump-Light Coupling Efficiencies

It is difficult to extend the analytic analysis further, however, in Zemax we can model the exact lens system used for coupling light from the diode into the rod. In our case we use an all-aspheric, two-lens telescope to image the core of the fiber-pigtail onto the face of the rod. Further, the ratio of focal lengths is chosen such that the image falls within the pump cladding diameter. In Figure B.2 we compare the coupling efficiency of an all-spherical

lens system to the all-aspheric one used in the lab. The difference is stark, with more than half the pump light lost using spherical lenses. It is tempting to suggest moving the image plane closer to the last lens in Figure B.2 (a), onto the “plane of confusion,” where the ray bundle appears to have the smallest diameter. Indeed, this would be the optimum location if the object were a point source as shown in the layout, though in reality, our object is of finite size. The true spot size is accurately depicted in Figures B.2 (b) and B.2 (d), where rays from multiple point sources within a $200\text{ }\mu\text{m}$ circle on the object plane have been traced to the image plane.

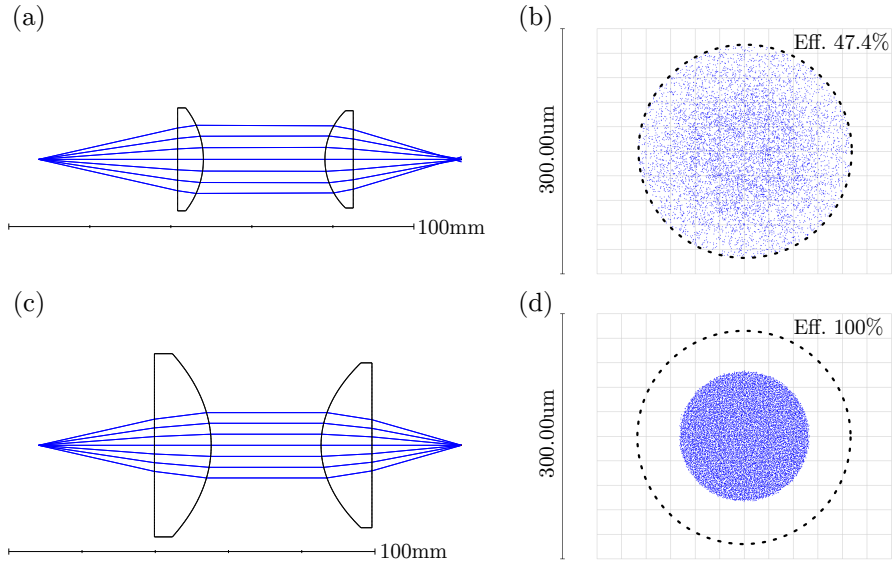


Figure B.2: Zemax simulation comparing an all-spherical and all-aspheric lens system. Rays are traced left-to-right in (a) and (c). The object is a $200\text{ }\mu\text{m}$ circle, with an NA of 0.22, representing the fiber-pigtail core. The system is first optimized for collimation between lenses; afterwards the image plane is moved to the point of maximal coupling efficiency. Note the surface orientations that place the lenses in “near-best-form”. (a) Spherical lens system. Components left-to-right: Thorlabs LA1422, ELF of 40.0 mm; Ross Optics L-PCX175, EFL of 32.0 mm. (b) Distribution of rays traced to the image plane for the all-spherical system. The rod’s $260\text{ }\mu\text{m}$ pump cladding diameter is indicated by a dotted circle, and the vignetted rays are removed. (c) Aspheric lens system. Components left-to-right: Asphericon AHL50-40, ELF of 40.0 mm; Asphericon AHL45-32, EFL of 32.0 mm. (d) Distribution of rays traced to the image plane for the all-aspheric system. Image is a reduced, but near perfect replica of the object, which under-fills the rod.

Appendix C

Effects of Oscillator State Changes

It is important to be aware that the system properties described in Section 2.6.4 are subject to spontaneous changes of the oscillator state. In particular, the oscillator has an affinity for spectrally broadened states, like that in Figure C.1-b. Though counterintuitive, we have found these large-bandwidth states lead to *longer* duration pulses after amplification. Let us consider the two scenarios in Figure C.1, where we will denote the narrow bandwidth initial state as ‘state-A’ and the broader one as ‘state-B.’ As just mentioned, the measured pulse duration after amplification of state-B (Figure C.1-f) is 2.1 times longer than that of state-A (Figure C.1-e). Gain narrowing alone cannot account for this outcome, since the bandwidth after amplification is comparable for either starting state (Figure C.1-a/b). The transform limited pulse durations after amplification corroborate this assertion. In fact, the TL duration after amplification for state-B is 0.8 times that of state-A. Therefore, we are certainly not bandwidth limited by gain narrowing, and higher-order spectral phase must be at play in preventing us from compressing the pulse. We have two chief hypotheses for the source of the higher-order phase: (1) the pulse from the oscillator has higher order phase to begin with, and gain narrowing selects out an incompressible section of the spectrum; (2) stretching the pulse transforms temporal variations in the refractive index to a wavelength dependent phase.

Evidence for the first possibility is found in comparing the measured and transform limited spectra before amplification (Figure C.1-c/d). It is clear that the measured pulse duration for state-A is closer to the transform limit, than that for state-B: the ratio of measured to TL duration for state-A is 1.4, whereas for state-B it is 2.1. If the pulse maintains its form (Gaussian, sech, etc...), then the time-bandwidth product (TBP) is constant, and the pulse duration must scale inversely to the bandwidth ($TBP = \text{time} \times \text{bandwidth}$). Under these circumstances, the ratio of the spectral width before and after amplification determines the expected final pulse duration. Due to the peaked nature of the state-B spectra, a -3 dB definition of spectral width

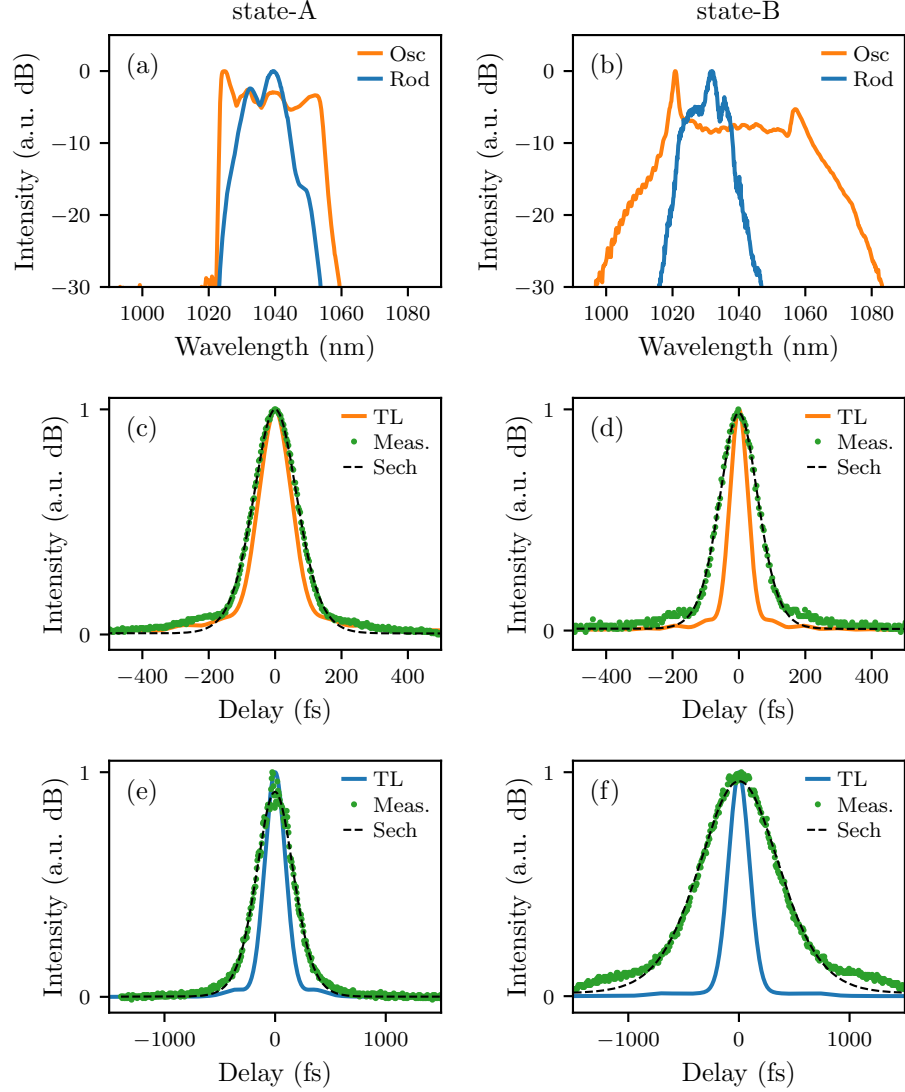


Figure C.1: Effects of gain narrowing on CPA. We compare minimized (column one, 'state-A') and excessive (column two, 'state-B') gain narrowing scenarios. (a)/(b) Spectrum of the oscillator (orange) and PC-rod (blue, 7 A). (c)/(d) Comparison of the measured (green points) and transform limited (orange) autocorrelations for the oscillator. A pulse duration for the measured values is found assuming a sech shape (black dashed). Respective measured/TL pulse durations: (c) 101 fs/74 fs FWHM; (d) 89 fs/43 fs FWHM. (e)/(f) Corresponding autocorrelations for the PC-rod. Measured/TL pulse durations: (e) 257 fs/152 fs FWHM; (f) 552 fs/139 fs FWHM.

would not be representative. Therefore, we use a rectangular spectral analog having the same area and modified first moment of area as the measured spectrum (linear, not log scale). The modification is simply to consider the *absolute value* of the distance from the spectral centroid. The measured and analog spectra are shown in Figure C.2, and their autocorrelations are fairly similar, as desired. With this method, we find the bandwidth of state-B before amplification is 55.9 nm, and afterwards it is 13.2 nm; a reduction of 4.32 times. Given the initial measured pulse duration is 89 fs, we expect that after amplification it will be 376 fs. Clearly, this estimate falls short of the measured value of 552 fs, and we speculate that the difference results from a non-parabolic phase about the sharp spectral peaks of state-B.

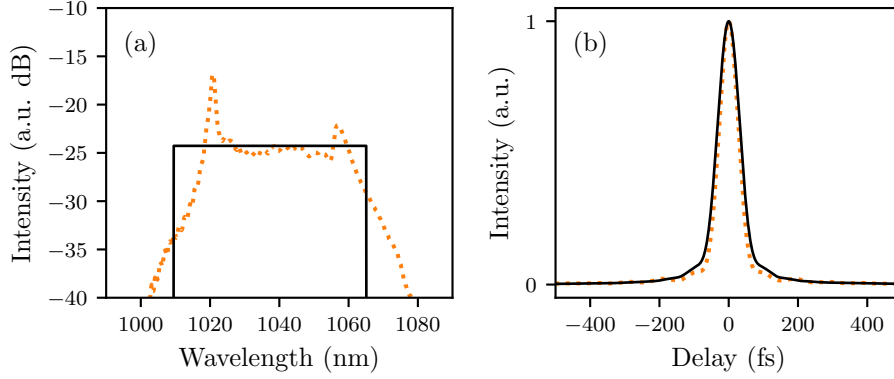


Figure C.2: Characterizing the spectral bandwidth using a rectangular analog spectrum. (a) Measured spectrum (orange dotted) and its rectangular analog (black solid), which preserves area, and the first moment of area. (b) Respective intensity autocorrelations of the transform limited pulses. The measured/rectangular-analog pulse durations are 43 fs/53 fs.

The second possibility comes from a mapping between time and wavelength, which the pulse stretcher imparts. This correlation between time and wavelength links temporal modulations to spectral ones. Initially, gain narrowing was suspected to lead to temporal modulations that impart a non-parabolic spectral phase, and prevent pulse compression. With a stretched pulse, band-pass filtering leads to an attenuation of the leading and trailing edges of the pulse, shortening it in time, and modulating its profile as it traverses the filter (a gain medium in this case). However to first order, gain narrowing simply amplifies select spectral components more than others, and while this does modulate the temporal envelope, it does not add

spectral phase. On the other hand, SPM and gain saturation are effects that lead to a temporally varying refractive index, and do apply spectral phase. For a substantially stretched pulse, we saw in Figure 2.9-e that the temporal envelope of the pulse takes on the shape of the spectrum. With this form, SPM is strongest about the spectral peaks, and thus the phase profile is not uniformly applied to all wavelengths. Gain saturation similarly applies a nonuniform phase, though it is proportional to the integrated intensity—changes in both the gain and refractive index occur as the excited population of the gain medium is depleted. Both effects certainly exist, the questions is in what magnitude, and to what degree they apply differently to state-A and state-B. Given our pulses are stretched to hundreds of picoseconds in duration to avoid SPM, it seems this effect would be limited. In comparison, gain saturation is less sensitive to pulse duration, since the excited state is repopulated at a constant and relatively slow rate between pulses. Importantly, in either case one would have to conjecture that the leading and trailing spectral components effect the phase of, or become, the central ones, as ultimately wavelengths outside the bandwidth of the gain media are attenuated. With further consideration, it then seems less likely that gain saturation at wavelengths outside the bandwidth would influence the final spectral phase of the pulse. We therefore recommend looking at the influence of SPM, and later gain saturation, after thoroughly investigating possibility one.

Much of the gain narrowing mystery could be revealed by tracing the phase evolution of the pulse with SHG-FROG. However, at present we have only retrieved the temporal and spectral phase for state-A after the PC-rod (figure 2.25). This was done in tandem with SHG-FROG retrievals after the HCF fiber, which are discussed in Chapter 3.

Appendix D

Preliminary Results and Higher–Order Nonlinear Effects

D.1 Introduction

We will discuss the simulations, spectral broadening trials, and analyses of higher–order nonlinear effects leading up to the final trial and simulation described in Section 3.1.3. The preliminary simulation provides a context for what was expected before measurements were taken. The stark contrast between it and experimental results is what instigated exploration of higher order effects. Here, we will dispel three of the more common higher–order effects (modulation instability (MI), self–steepening, and soliton fission) as probable causes of the disparity. In doing so, we make the case that broadening in our experiment remains in the “first order” self–phase–modulation (SPM) dominated regime.

The results of two additional broadening trials are provided to give statistical significance to the conclusions drawn from the final broadening trial in Section 3.1.3.

D.2 Simulation

Initial simulations for the expected output of the HCF were performed assuming a generic sech pulse shape, with a total energy of $3.6\,\mu\text{J}$, and a FWHM duration of 250 fs (carried out by Edmund Kelleher, Technical University of Denmark, using code written by John Travers, Heriot Watt University). The spectral evolution of the pulse along a 50 cm length of HCF with a $28\,\mu\text{m}$ core diameter, at 10 bar and 20 bar, assuming $n_2 = 8.38 \times 10^{-24}\,\text{m}^2/\text{W}$ for argon at 1 bar and 300 K, is shown in Figure D.1. Note the modeled core diameter is $2\,\mu\text{m}$ smaller than that measured with the SEM locally and the nonlinear index is $1.67\times$ smaller than that cited in

Section 3.1.1, but these are the initial values provided by our collaborators and are a reasonable starting point. The result for a filling pressure of 10 bar is particularly attractive. For the first ~ 40 cm the spectrum linearly expands while developing symmetric modulations characteristic of pure SPM (cf. Agrawal [1] Figure 4.4). The final rapid spectral expansion results from soliton self-compression, as the pulse propagates in the anomalously dispersive regime. In this case, the self-compressed pulse duration is 20 fs. When the pressure is raised to 20 bar two other interesting phenomena arise: modulation instability (MI), and the generation of a resonant dispersive wave (RDW). The latter will be discussed in Section 3.2.

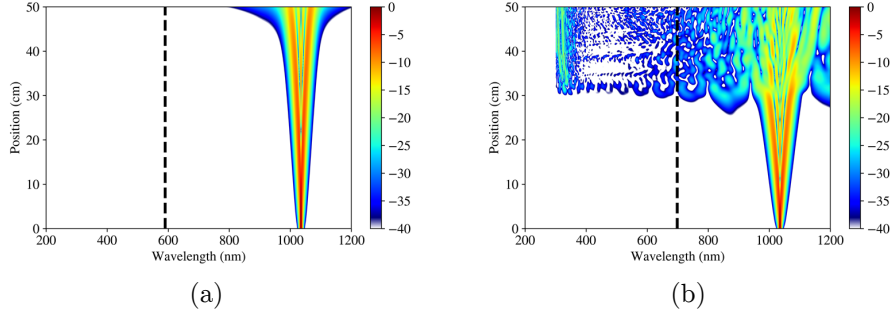


Figure D.1: Simulated spectral evolution of a 3.6 μJ , 250 fs FWHM sech pulse, along a 50 cm length of HCF, with a 28 μm core diameter. The simulation was performed for two different filling pressures of argon. The zero dispersion wavelengths are indicated by dotted lines. (a) 10 bar filling pressure. The initially linear broadening and symmetric modulations are indicative of SPM. Rapid spectral expansion after ~ 40 cm is due to soliton self-compression. By 50 cm the pulse is 20 fs in duration. (b) 20 bar filling pressure. Broadening up to ~ 25 cm appears to be predominantly by SPM. After this point, periodic lobes in the spectrum indicate broadening through modulation instability. A resonant dispersive wave (RDW) is also visible at ~ 350 nm. Simulations carried out by Edmund Kelleher, Technical University of Denmark, using code written by John Travers, Heriot Watt University.

D.3 Modulation Instability (MI)

In a somewhat similar fashion to soliton self-compression, modulation instability is the exponential amplification of noise resulting from an interplay of anomalous dispersion and nonlinearity. When continuous wave (CW) noise and signal co-propagate, their combined envelope is modulated at the beat frequency. Consider that a beat is the most simplistic form of pulse train—

continued addition of higher harmonics, spaced by the magnitude of the beat frequency, would result in a train of narrow, high peak-power, pulses. However, even for only two frequencies there are local regions of increased peak-power, at which SPM can generate new frequencies. Certain beat frequencies lead to a portion of the signal being frequency shifted to spectrally coincide with the noise—this is the first stage of a regenerative feedback loop, since the beating of noise and signal leads to the generation of more noise. Anomalous dispersion is required to “close the loop,” by removing the normal-like phase imparted by SPM—if the newly generated noise components are phase-matched to the signal the process will continue unbounded. The gain spectrum for noise with a angular frequency displacement Ω from the signal is given by [1] as

$$g(\Omega) = |\beta_2 \Omega| (\Omega_c^2 - \Omega^2)^{1/2}, \quad \text{where} \quad \Omega_c^2 = \frac{4\gamma P_0}{|\beta_2|} \quad (\text{D.1})$$

is cutoff frequency above which no gain is experienced, and P_0 is the average signal power. Maximum gain occurs at two locations ($\Omega_{\max} = \pm\Omega_c/\sqrt{2}$) centered about the signal frequency, as shown in Figure D.2. Though partially obfuscated by the long wavelength axis limit, the first two spectral side lobes in the 20 bar simulation (Figure D.1-b) will coincide with the MI gain peaks. Here, SPM provides the “noise” for MI to develop; it is the dynamic range of the color map that makes the side lobes appear noncausal. As we can see, MI leads to a chain of events, where the newly generated spectral components themselves become the signal and grow side lobes. The diminishing distance between lobes can be understood from the $P_0^{1/2}$ dependence of the MI gain peak.

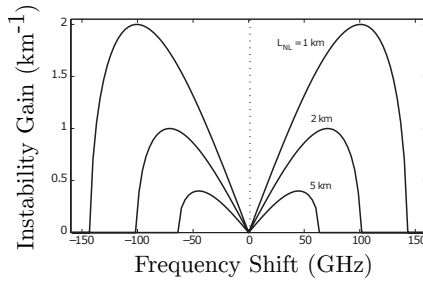


Figure D.2: Gain spectrum of modulation instability. Three different nonlinear lengths ($L_{\text{NL}} = 1/\gamma P_0$) are shown. The gain peaks are symmetric about the signal frequency. Reprinted from [1], pages 132, with permission from Elsevier.

D.4 First Broadening Trial

A first attempt at reproducing the simulations was carried out in a 46 cm length of HCF (here and elsewhere the experimental HCF geometry is the same as listed in Figure 3.1). Evolution of the spectrum and SHG-autocorrelation after the HCF, as the pulse energy and filling pressure are increased, is shown in Figure D.3. Spectrally, we see a smooth broadening without the characteristic modulations of SPM or MI. For the most extreme set point (3.6 μ J, 25 bar) the -40 dB spectral bandwidth is approximately three times that after the PC-rod. The broadening is significant, but falls far short of expectations, being comparable to the 10 bar simulation at a propagation distance of only 30 cm (Figure D.3-b). Autocorrelations in the time domain also leave much to be desired: the pulse does begin to self compress as the pressure, and hence bandwidth, is increased, but the final duration is over an order of magnitude greater than 20 fs.

A clue towards one source of discrepancy is found in the satellite pulses on either side of the autocorrelation peak (Figure D.3-b). We found that the height of these peaks was polarization dependent, and symmetrizing the trace required rotating the HWP before the autocorrelator. To a lesser degree, the flanks of the central peak were also polarization dependent. We verified that the input pulse from the PC-rod was strongly linearly polarized, with a measured polarization extinction ratio (PER) > 21 dB (limited by the resolution of the power meter), as ensured by TFP (2). In contrast to after the HCF, the symmetry of the input pulse autocorrelation was not polarization dependent, pointing to a degradation of the PER after the HCF. Indeed, we found the PER varied from only 9.0 dB to 11 dB over the full range of linear launch polarizations into the HCF (a typical value is 15 dB). A new section of HCF cured the PER degradation problem, but it is important to note that ultimately this HCF section was salvaged by recleaving the input face. The anomalously low coupling efficiency—percentages in the mid 70s rather than mid 80s—was an unheeded sign of a poor cleave. Given the PER and coupling issues, we should not over-analyze deviations from the simulations. The results are provided as a reference for the effects of PER degradation and, as we will see, to support a generally consistent trend in the spectral evolution, regardless of PER.

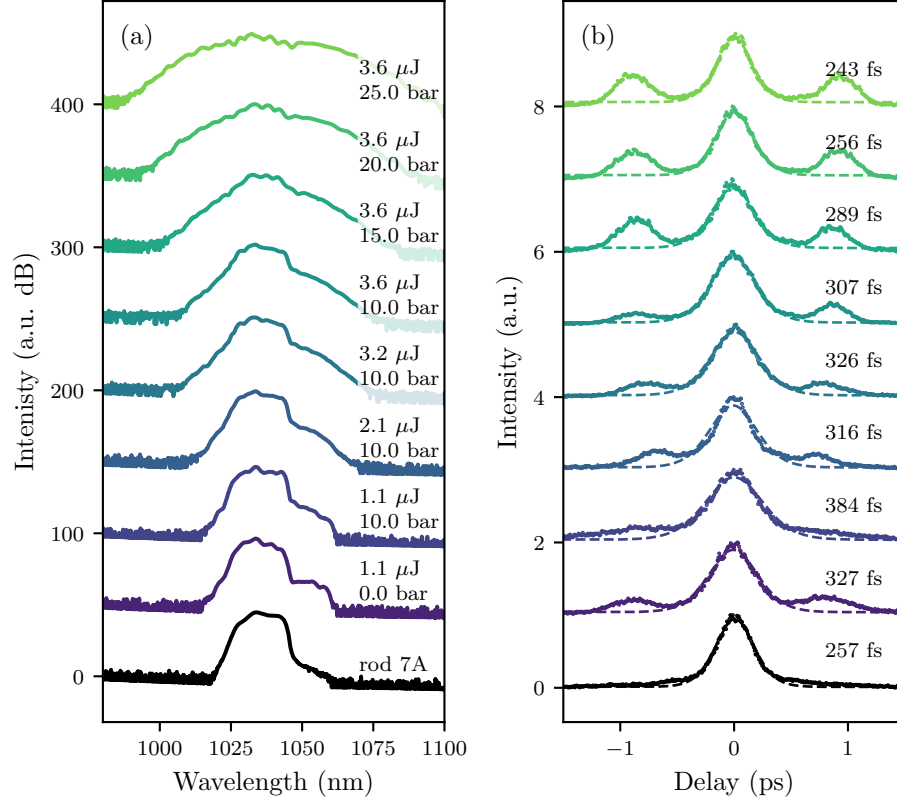


Figure D.3: Measured spectra and autocorrelations after a 46 cm length of HCF with a poor PER (< 11 dB). Geometric properties of the fiber are listed in Figure 3.1. (a) Spectral evolution as the pulse energy and filling pressure are increased. Spectral broadening increases with peak power and nonlinear parameter γ , which grows with gas density. (b) Corresponding autocorrelations of the output, with the deconvolved FWHM pulse durations labeled. A sech pulse form is assumed. The onset of soliton self compression occurs only for the broadest spectra. Two satellite peaks are present, and their relative intensities vary with polarization into the autocorrelator. External compression was not attempted.

D.5 Second Broadening Trial

We repeated the previous broadening trial with a new 27 cm long section of HCF. The coupling efficiency and PER were 81% and 15 dB respectively, which is a marked improvement over the previous fiber. The spectra and autocorrelations are shown in Figure D.4. The satellite peaks in the autocorrelation have disappeared, and we found the traces to remain symmetric as the HWP before the autocorrelator was rotated. The broadening rate remains essentially unchanged from the previous trial. The -40 dB bandwidth grows by only ~ 1.8 times over the 27 cm length. If we assume the rate of broadening is linear with fiber length (reasonable when SPM is the predominant mechanism), then the extrapolated increase in bandwidth over a 46 cm length is $3.1\times$, like the previous trial. A notable difference from the longer HCF is that the pulse duration grows monotonically and does not begin to self compress. Additionally, the short wavelengths are more structured and broaden less than the long wavelengths. The structure is already visible for a modest $1.2\text{ }\mu\text{J}$ pulse energy at 10 bar and seems to grow as the energy and pressure increase. The spectral asymmetry suggests an asymmetric temporal intensity, and by extension, an asymmetric spectral phase. This conclusion is discussed in Section 3.1.3. Here, we will consider two other higher-order effects that can generate these asymmetries (self-steepening and soliton fission) and show that they do not pertain to our experimental results.

D.6 Self-Steepening

Self steepening is similar to SPM in that it results from an intensity dependent refractive index, but in this case the envelope, rather than the phase, is modulated. For a positive nonlinear parameter γ (typical of most materials), the peak of the pulse travels more slowly than the flanks, causing the front to stretch and the rear to compress, “steepening” the pulse (Figure D.5-a). Given SPM generates a blue-shift in proportion to the slope of the trailing pulse edge, self steepening enhances the generation of high frequencies. The effect can be seen in Figure D.5-b, where the high frequencies extend further from the central frequency than the low frequencies. Though the shape of the self-steepened pulse spectrum shares many features with that measured after the HCF (Figure 3.6-d), a critical difference is that they are effectively mirror images, being on frequency and wavelength scales respectively. There is no reason to suppose the nonlinear parameter defies convention and is negative; therefore, another explanation is necessary.

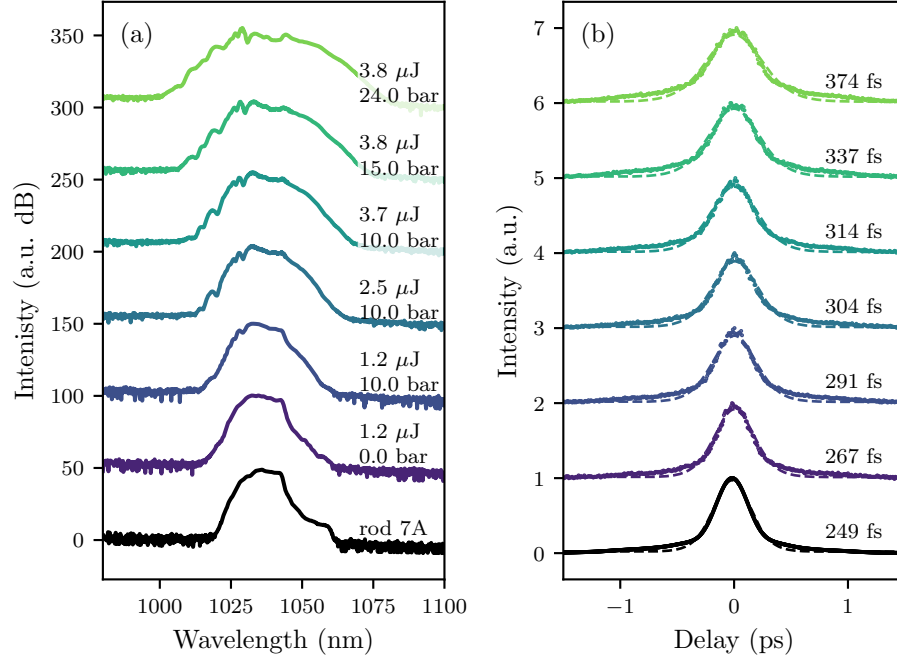


Figure D.4: Measured spectra and autocorrelations after a 27 cm length of HCF with a satisfactory PER (15 dB). Geometric properties of the fiber are listed in Figure 3.1. (a) Spectral broadening as the pulse energy and filling pressure are increased. The spectrum develops an asymmetry, being modulated at short wavelengths and smoothly broadened at long ones. (b) Corresponding autocorrelations of the output, with the deconvolved FWHM pulse durations labeled. A sech pulse form is assumed. The pulse duration increases monotonically with pressure and pulse energy, likely tracking a shift towards the normal dispersion regime, in addition to an increase of nonlinear phase. Note there are no satellite peaks with an improved PER. External compression was not attempted.

D.7 Soliton Fission

“Fission” refers to the splitting of a higher order soliton into several fundamental solitons. This is possible because higher order solitons carry N^2 times the energy of the fundamental solution, where N is the integer order parameter. For sech-shaped pulses

$$N = \sqrt{\frac{\gamma P_0 \tau^2}{|\beta_2|}}, \quad (\text{D.2})$$

where γ is the nonlinear parameter, P_0 is the pulse peak power, τ is the characteristic width, and β_2 is the group velocity dispersion. The number of component fundamental solitons in a higher order solution is, in fact, the order parameter N . The width τ_k and peak power P_k of each sub-soliton is given by

$$\tau_k = \frac{\tau}{2N + 1 - 2k}, \quad P_k = \frac{(2N + 1 - 2k)^2}{N^2} P_0, \quad (\text{D.3})$$

where k varies from 1 to N [1]. When unperturbed, the component solitons have the same group velocity and their superposition (a higher order soliton) evolves cyclically between periods of compression and expansion in the spectral and temporal domains. However, the superposition is in a highly unstable equilibrium and various effects can lead to a fission event, after which the fundamental solitons evolve semi-independently in the temporal and spectral domains. For reference, such an event is depicted in Figure D.5-c, where a soliton of order $N = 2$ breaks into two fundamental solitons, with the narrower and higher intensity one lagging the other. In this case, fission was initiated by self steepening, which shifted the peak to the rear of the pulse. In the spectral domain (Figure D.5-d) we see the now familiar enhancement of high frequencies and also modulations over the spectral region shared by the two separate fundamental solitons.

In our experiment it is unclear what would lead to fission of the pulse. Raman scattering can be ruled out as argon is a monatomic gas. Dispersive waves should not be generated, since even the shortest wavelengths in our trials do not leave the anomalous dispersion regime. Nonetheless, the temporal retrieval of the uncompressed pulse after the HCF (Figure 3.6-f) shows the main peak dividing in two, warranting examination.

We can determine the spectral composition of the temporal peaks from the FROG-retrieved spectral phase (Figure 3.6-d), specifically from its derivative with respect to angular frequency ω , also known as the group delay. Moreover, the slope of the group delay tells us whether the spectral components are solitonic or dispersive: regions with zero slope occur when a

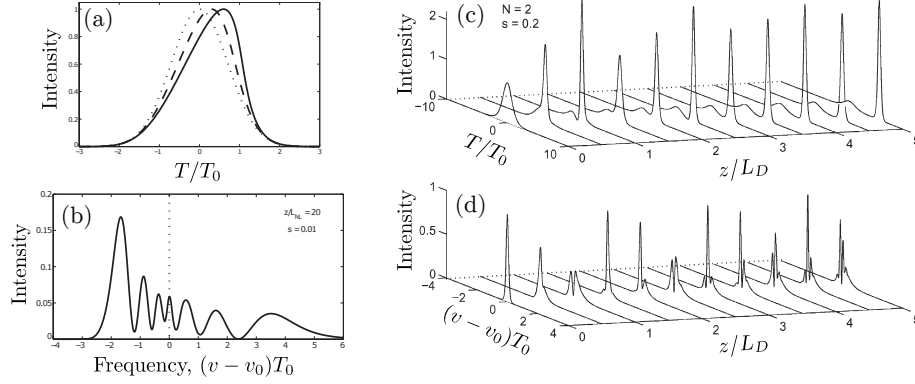


Figure D.5: Potential sources of asymmetric spectra. (a) Example of a self-steepened pulse form. The peak of the pulse shift to the rear as the pulse propagates. (b) Self steepening in the presence of SPM leads to enhanced broadening at higher frequencies. (c) Soliton fission results in two fundamental solitons that evolve semi-independently. Temporally, the two solitons propagate with slightly different velocities. (d) Spectrally, the higher intensity soliton broadens more. Since the solitons share a range of spectral components, a time delay between them modulates the total spectral intensity. In this case, soliton fission was initiated by self steepening. Reprinted from [1], pages 117, 118, 176, with permission from Elsevier.

range of frequencies propagate at the same speed, and therefore indicate the presence of a soliton when the medium is dispersive. This analysis is best demonstrated on an exemplary case of soliton fission before applying it to the more subtle experimental results.

For a clear cut example of soliton fission, we simulated the propagation of a transform limited, 30 fs FWHM duration, 10 MW peak power (0.34 μ J) sech pulse, over 100 cm of the same HCF employed in our broadening stage, filled with argon to 20 bar (Figure D.6). We can identify key features that would indicate soliton fission in our measurements. Let us start with the temporal evolution (Figure D.6-a). The initial pulse forms a soliton of order $N = 3.7$, which self compresses and fissions into three fundamental solitons in ~ 8 cm. At the end of the fiber we can piece together which spectral components belong to a given soliton by looking for zero slope regions in the group delay with an appropriate time offset. For example, there is a broad region from 1100 nm to 1600 nm with a uniform delay of 340 fs (Figure D.6-d) associated with the most prominent soliton (Figure D.6-c). The dispersive wave at ~ 620 nm, as it names suggests, has its spectral components spread in time, which is reflected by the nearly vertical slope of the group delay.

We now apply the same analysis to the uncompressed pulse after the

HCF (Figure 3.6). The spectrum and temporal intensity for this SHG-FROG retrieval are replotted in Figure D.7, along with the group delay. It is immediately evident from the group delay that the structured blue-end of the spectrum is widely dispersed, whereas the red-end has a relatively constant delay. Therefore, it seems that the leading (negative time) peak in Figure D.7-b is solitonic and pulling away from the dispersive elongated tail.

Physically, this is a bit odd. For the given filling pressure, both the MS and perturbative MS models predict an anomalous dispersion across the considered spectral range. Furthermore, since the nonlinear index is positive, greater intensities—like those at the peak of a soliton—will lead to an increase in the refractive index and thus more retardation. These two properties suggest that the shorter wavelengths should lead the longer ones, just as is seen in our simulated example of soliton fission (bar the dispersive wave at the extreme end of the spectrum, which travels more slowly due to third order dispersion) (Figure D.6). The dilemma cannot be resolved by a single SHG-FROG retrieval, due to an inherent time ambiguity in the measurement technique. A second SHG-FROG retrieval after introducing a known *sign* of dispersion is required. In our case, the compressor after the HCF adds some amount of anomalous GDD and we see a reduced pulse duration in the second retrieval (Figure 3.7). This result contradicts our physical intuition, since it shows the longer wavelengths do lead the shorter ones, as asserted by the uncompressed FROG retrieval. It seems then that anomalous dispersion and soliton dynamics play a minor role. The findings corroborate our earlier observation of minimal or nonexistent self compression and significantly under broadened spectra in all three trials.

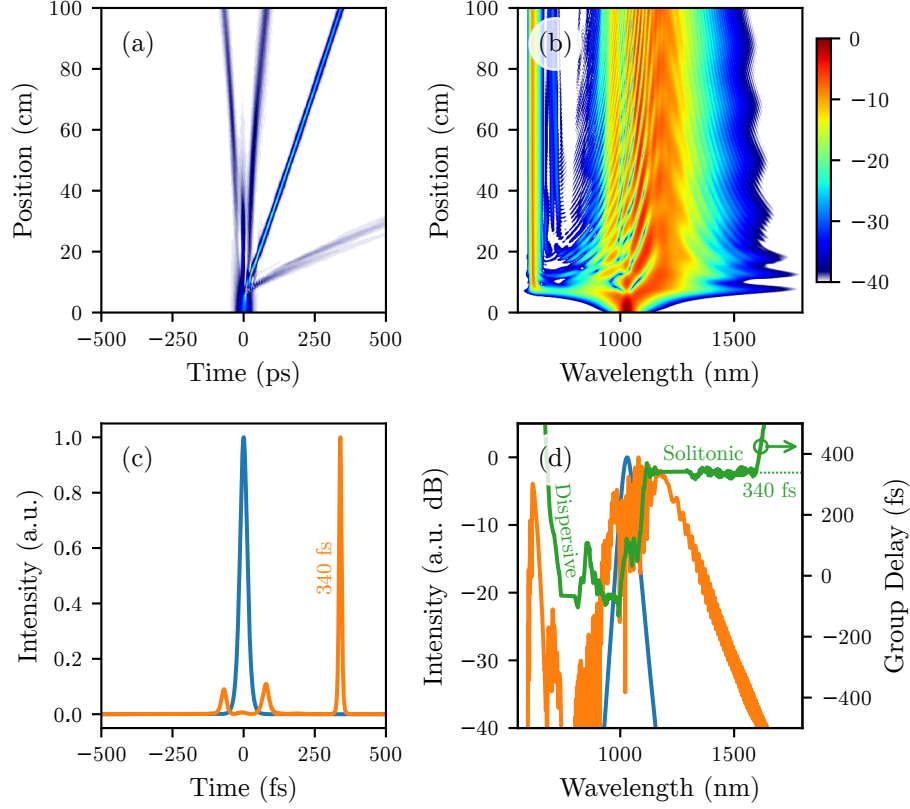


Figure D.6: Simulated fission of a soliton of order $N = 3.7$, providing an example of what can be determined from the group delay. The pulse(s) propagate over 100 cm of HCF that has the geometry defined in Figure 3.1 and is filled with argon at 20 bar. (a) Temporal evolution of the pulse. Color map is linear in power. At ~ 8 cm the pulse self compresses and fissions. Excess energy is shed in a dispersive wave (the most horizontal, fading, streak) in order to form one predominant fundamental soliton and two others that are not yet fully developed. (b) Spectral evolution of the pulse. The fission event at 8 cm is clearly defined by the generation of a dispersive wave centered at ~ 620 nm. (c) Temporal slice at 0 cm (blue) and 100 cm (orange) showing the evolution from a single higher order soliton into three fundamental solitons. (d) Respective spectral slices at 0 cm (blue) and 100 cm (orange). Of particular interest for this demonstration is the group delay (green), which indicates solitonic and dispersive spectral components. Contrast the steep slope of the group delay about the dispersive wave at 620 nm, with the flat region spanning 1100 nm to 1600 nm. Spectral components in the latter region propagated with the same velocity and are all delayed by 340 fs, which is the location of the predominant soliton in (c). Note the group delay has been smoothed to reveal the average behavior.

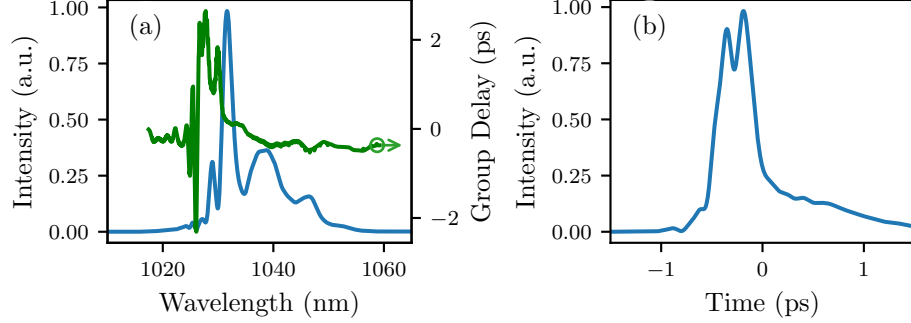


Figure D.7: Group delay of the uncompressed pulse after the HCF, revealing both solitonic and dispersive spectral components and their connection to the temporal intensity. (a) The spectral range above ~ 1040 nm has a fairly constant group delay and thus likely forms a soliton. The delay of these components ranges from -0.5 ps to -0.3 ps placing them in the leading peak of the temporal waveform in (b). Here, “leading” refers to negative times, though physically the direction is ambiguous in an SHG-FROG retrieval. The second temporal peak contains spectral components between 1040 nm and 1035 nm. The small plateau about 1035 nm is potentially solitonic. Below 1035 nm the group delay is steeply sloped, reflecting the widely dispersed spectral components that constitute the elongated tail of the pulse.

D.8 Conclusion

Analysis of higher-order nonlinear effects shows that they have little influence on our experimental pulse. This conclusion forces us to reconsider the influence of initial pulse conditions and model assumptions on low-order nonlinear effects; namely, SPM. In Section 3.1.3 we show that the disparity between initial simulations and experimental results comes from either over estimating the nonlinear parameter γ or the pulse peak power P_0 , and not considering the initial phase of the input pulse.

Appendix E

PCF End-Connector Mount

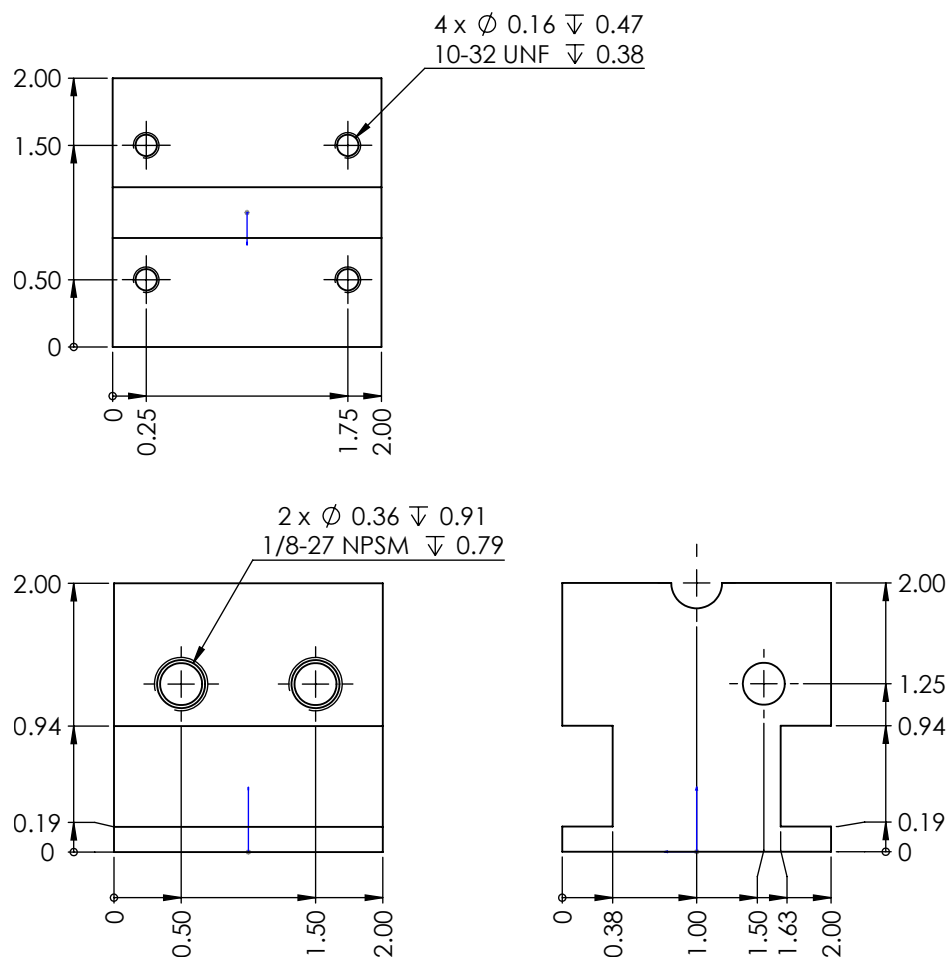


Figure E.1: Base of the water cooled mount used to hold the PCF end-connector. All dimensions are in inches.

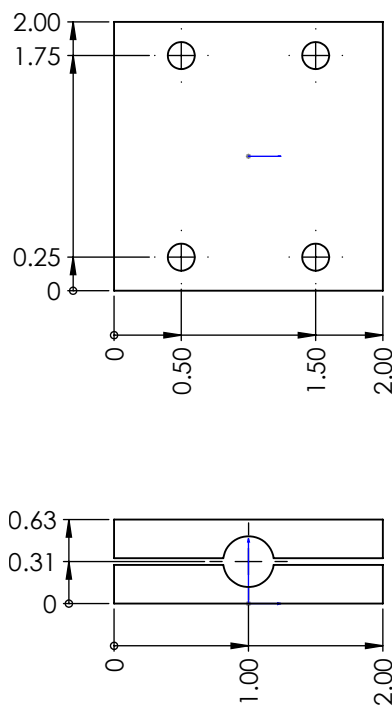


Figure E.2: Top of the water cooled mount used to hold the PCF end-connector. All dimensions are in inches.