A Few-Cycle Cr$^{4+}$:YAG Laser and Optical Studies of Photonic Crystals

by

Daniel Jacob Ripin

B.S. Physics, Emory University, 1995

Submitted to the Department of Physics in partial fulfillment of the requirements for the degree of Doctor of Philosophy in Physics

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Abstract

A prismless Cr\textsuperscript{4+}:YAG laser was used to generate 20 fs pulses at 1450 nm with a bandwidth of 190 nm FWHM. Intracavity group velocity dispersion was compensated with double-chirped mirrors. Pulse spectrum was observable from 1140 to >1700 nm. Broadband saturable Bragg reflectors were designed and used to ensure self-starting of 35 fs pulses in the ultrafast Cr\textsuperscript{4+}:YAG laser or to generate picosecond pulses tunable from 1400 to 1525 nm. The mirrors were a 7-pair GaAs/Al\textsubscript{x}O\textsubscript{y} quarter-wave dielectric stack, and the absorber consisted of an InGaAs quantum well centered in a half-wave InP layer. Transmission was measured through a photonic bandgap crystal microcavity resonant near 1550 nm. Cavity quality factors as high as 360 were observed for cavities with a modal volume of only 2(λ/2n)\textsuperscript{3}. Photonic crystals were used to enhance the total emission of a light emitting diode at 980 nm by 8-fold. At particular wavelengths, collected photoluminescence enhancements larger than 100 were observed.

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TABLE OF CONTENTS

ACKNOWLEDGMENTS .......................................................... 5

LIST OF FIGURES ........................................................... 13

LIST OF TABLES ............................................................. 29

INTRODUCTION ............................................................... 31

1.1 Time and Length Scales of Light ..................................... 31
1.2 Ultrashort Pulse Generation ........................................... 32
1.3 Wavelength-Scale Light Manipulation ............................... 33
1.4 Optical Communication ................................................ 33
1.5 Thesis Outline .......................................................... 34

ULTRAFAST LASER DESIGN ............................................... 37

2.1 Introduction ............................................................ 37
2.2 Background ............................................................ 37
2.3 Cr^{4+}:YAG ............................................................. 38
2.4 Four-Level System ..................................................... 41
2.5 Modelocking ........................................................... 44
2.6 Gain Filtering .......................................................... 47
2.7 Group Delay Dispersion .............................................. 48
BROADBAND SATURABLE ABSORBERS FOR
ULTRAFAST Cr⁴⁺:YAG LASERS

4.1 Introduction ........................................... 155
4.2 Background ........................................... 155
4.3 Two Level System .................................... 158
4.4 Saturable Bragg Reflector Structure ............... 160
4.5 Mirror Design ........................................ 161
4.6 Saturable Absorber Design ......................... 165
4.7 Fabrication ........................................... 171
4.8 Self-Starting Cr⁴⁺:YAG Laser .................... 173
4.9 Future Work ........................................ 178
4.10 Conclusions ...................................... 178

PHOTONIC BANDGAP MICROCAVITY ............................ 181

5.1 Introduction ........................................... 181
5.2 Background ........................................... 181
5.3 Photonic Crystals ..................................... 182
5.4 Light Confinement .................................... 186
5.5 Microcavity Design ................................... 188
5.6 Device Fabrication ................................... 196
5.7 Experimental Setup ................................... 199
5.8 Waveguide Coupling and Loss ..................... 202
5.9 Optical Microcavities ............................... 206
5.10 Future Work ....................................... 211
5.11 Conclusions ...................................... 212

PHOTONIC CRYSTAL LIGHT EMITTING DIODE ......................... 213

6.1 Introduction .......................................... 213
6.2 Background .......................................... 213
6.3 980 nm Light Emitting Diode ..................... 215
6.4 LED Fabrication ................................... 217
6.5 Experimental Setup ......................................................... 219
6.6 Unpatterned LED .......................................................... 222
6.7 LED with Random Holes .................................................. 224
6.8 Triangular Photonic Crystal .............................................. 230
6.9 Output Coupling with Photonic Crystal ............................ 234
6.10 Power Dependence ....................................................... 241
6.11 Wavelength Tuning ........................................................ 242
6.12 Angular Dependence ..................................................... 245
6.13 Pump Input Coupling with Photonic Crystal ..................... 248
6.14 Future Work ............................................................... 252
6.15 Conclusions ............................................................... 254

CONCLUSION ........................................................................... 255

7.1 Conclusions ................................................................. 255
7.2 Shorter Cr$^{4+}$:YAG Pulses ............................................. 256
7.3 Toward Single-Cycle Pulses ............................................ 257
7.4 Pump-Probe Spectroscopy .............................................. 258
7.5 Frequency Chains .......................................................... 258
7.6 Tunable, Switchable, and Active Photonic Crystals .......... 259
7.7 Near-field Scanning Optical Microscopy .......................... 260

REFERENCES ......................................................................... 261
LIST OF FIGURES

Figure 1.1 ........................................................................................................... 32
Minimum pulsewidth and fractional bandwidth of pulses generated directly from a Ti:Sapphire laser over time.

Figure 2.1 ........................................................................................................... 41
Measured photoluminescence (PL) of a 2 cm Cr$^{4+}$:YAG laser crystal.

Figure 2.2 ........................................................................................................... 42
Diagram of the four-level system model of a Cr$^{4+}$:YAG laser. A pump laser drives the transition from state 1 to state 4 at a rate $W_p$, but also results in excited state absorption (ESA) from state 3 and state 4 at rates $W_{ESA3}$ and $W_{ESA4}$. Electrons in states 2, 3, and 4 relax to lower energy levels with relaxation rates $\gamma_{ij}$. Stimulated transitions occur at the lasing transition at a rate $W_L$.

Figure 2.3 ........................................................................................................... 46
Intensity waveforms calculated for 5 and 25 cavity modes with random phases.

Figure 2.4 ........................................................................................................... 46
Intensity waveform from a laser consisting of 5 and 25 modes with locked phases. As the number of modes increases, the pulse width decreases and the instantaneous intensity increases.

Figure 2.5 ........................................................................................................... 53
Index of refraction and group velocity dispersion (GVD) for YAG material.

Figure 2.6 ........................................................................................................... 55
Schematic of self phase modulation (SPM) in a Kerr media. (a) Electric field profile of a Gaussian pulse envelope with a carrier frequency. (b) Electric field profile of the Gaussian pulse after propagation through a Kerr media. As the peak of the pulse travels to the right, it sees a larger index of refraction than the pulse wings and therefore has a lower phase velocity. An overall chirp develops
within the pulse. The chirp has the same sign as normal dispersion with higher frequencies delayed relative to lower frequencies.

Figure 2.7 ................................................................................................................... 58
Schematic of optical elements within an actively modelocked laser.

Figure 2.8 ................................................................................................................... 58
Time domain picture of pulses generated by active modelocking. Pulses (A(t)) are generated where loss (l) is smaller than gain (g).

Figure 2.9 ................................................................................................................... 59
Schematic of active modelocking in frequency space. Lines correspond with cavity modes at v_n. Dashed arrows represent coupling from each mode to its neighbors by a sideband generating optical modulator. This process will couple the mode’s phases together.

Figure 2.10 ................................................................................................................ 62
Time domain picture of pulses generated by slow saturable absorber passive modelocking. Pulses (A(t)) are generated where loss (l) is smaller than gain (g). The loss modulates the leading edge of the pulse while gain saturation modulates the trailing edge.

Figure 2.11 ................................................................................................................ 64
Time domain picture of pulses generated by fast saturable absorber passive modelocking. Pulses (A(t)) are generated where loss (l) is smaller than gain (g). The loss is modulated by the intracavity pulses themselves.

Figure 2.12 ................................................................................................................ 66
Schematic of optical elements within a fast saturable absorber passively modelocked laser.

Figure 2.13 ................................................................................................................ 69
Pulse shortening rate (PSR) of active, slow saturable absorber passive, and fast saturable absorber passive modelocking, and pulse broadening rate (PBR) due to the finite gain bandwidth.

Figure 2.14 ................................................................................................................ 70
Schematic of Kerr-lens modelocking. The Kerr-effect creates an intensity dependent optical mode profile within a laser cavity. An aperture is placed within the cavity to generate higher loss for low intensity (cw) modes than for high intensity (pulsed) modes.

Figure 2.15 ................................................................................................................ 73
Schematic of the Kelly side-band phase matching process for GDD and TOD.

Figure 2.16 ................................................................................................................ 73
Schematic of a dispersion managed laser cavity. The cavity is considered to be
dispersion managed if each element within the laser cavity broadens the pulse significantly.

Figure 2.17........................................................................................................ 74
Approximate pulse stretching ratio for pulses in a Cr^{3+}:YAG laser consisting of
a 2 cm long gain crystal.

Figure 2.18........................................................................................................ 75
Schematic of a two mirror laser cavity. The mirrors have a radius of curvature
of R_1 and R_2, and a separation d.

Figure 2.19........................................................................................................ 78
Stability diagram of the two mirror cavity resonator. Stable cavity modes exist
in the gray region. Dotted lines indicate constant radii of curvature for the cav-
ity mirrors.

Figure 2.20........................................................................................................ 79
Schematic of (a) a Z-fold cavity and (b) its representation in terms of mirrors
and lenses that can be easily modeled with ABCD matrices. The cavity consists
of a flat end mirror, a distance L_1, a radius of curvature f_1 = R_1/2 mirror, a dis-
tance d_1, a laser crystal of thickness t, a distance d_2, a radius of curvature f_2 =
R_2/2 mirror, a distance L_2, and another flat end mirror.

Figure 2.21........................................................................................................ 81
Calculated beamwaist (ω_0) for a symmetric Z-fold laser cavity as a function of
the curved mirror separation d = d_1 + d_2 + t.

Figure 2.22........................................................................................................ 82
Calculated beamwaist (ω_0) for an asymmetric Z-fold laser cavity as a function
of the curved mirror separation d = d_1 + d_2 + t.

Figure 2.23........................................................................................................ 87
Schematic of a pulsewidth measuring autocorrelator. An input beam of laser
pulses is split by a beam splitter. One beam (ω_1) travels a length L, while a sec-
ond beam (ω_2) travels a distance L+ΔL, which results in a relative time delay
Δτ = ΔL/c. The beams are recombined with a second beamsplitter, and focused
through a nonlinear crystal. The nonlinear crystal creates sum frequency gen-
eration from path 1 and path 2, and simple harmonic generation from two path
1 or two path 2 photons.

Figure 2.24........................................................................................................ 90
(a) Schematic of interferometric, (b) intensity, and (c) background-free inten-
sity autocorrelation detection methods.

Figure 2.25........................................................................................................ 91
Momentum conservation for simple harmonic generation (SHG) and sum fre-
quency generation (SFG) in a nonlinear autocorrelator. The SHG (2k_1 and
2k₂) travel a path spatially separated from the SFG (k₁ + k₂).

Figure 2.26 ................................................................................................................. 92
Calculated interferometric, intensity, and background-free intensity autocorrelation functions for a 20 fs Gaussian pulse at 1500 nm.

Figure 2.27 ................................................................................................................. 93
Schematic of the time-ambiguity introduced by the noncollinear autocorrelator geometry. The input beams, separated by a distance D, are focused by a lens of focal length f into a nonlinear element. The autocorrelator cannot resolve time delays shorter than Δτ.

Figure 3.1 ................................................................................................................. 97
Schematic of the Z-fold cavity used to test Cr⁴⁺:YAG crystals and mirrors. Mirrors M1 - M3 are all high-reflector mirrors and OC is an output coupling mirror.

Figure 3.2 ................................................................................................................. 99
Reflectivity of (a) highly reflecting mirrors and (b) output coupling mirrors used in Cr⁴⁺:YAG laser cavities.

Figure 3.3 ................................................................................................................. 101
Knife-edge measurement of the Nd:YVO₄ beam waist. Data (black squares) agree well with a fit function (gray line) assuming a 1.2 mm beam waist.

Figure 3.4 ................................................................................................................. 102
Measurement of the transmission through a 1 cm (unfilled circles) and 2 cm (filled shapes) long Cr⁴⁺:YAG crystal for different pump power densities. The operating power density is indicated by a vertical dotted line.

Figure 3.5 ................................................................................................................. 104
Measurements of the output power of a Cr⁴⁺:YAG laser as a function of input pump power absorbed by the gain media for combinations of a 1 cm or 2 cm crystal and several output couplers.

Figure 3.6 ................................................................................................................. 105
Output power from a 1 cm crystal laser as a function of time after the laser is turned on.

Figure 3.7 ................................................................................................................. 106
Three methods of tuning a Z-fold cavity laser with broadband gain using (a) a birefringent tuning filter, (b) a single prism and a rotating end mirror, and (c) two prism and a translating slit.

Figure 3.8 ................................................................................................................. 108
Tuning curve of a Cr⁴⁺:YAG laser with a 2 cm crystal and a birefringent tuning plate.

Figure 3.9 ................................................................................................................. 109
Group delay dispersion (GDD) of 4 cm of Cr$^{4+}$:YAG (from reference [144]) and undoped YAG.

Figure 3.10 .................................................................................................................. 110
A standard geometry used to insert dispersion compensating prisms into a laser cavity.

Figure 3.11 .................................................................................................................. 111
The group delay dispersion (GDD) of several potential prism materials.

Figure 3.12 .................................................................................................................. 112
Transmission through Lucent's All-Wave fiber (with low water content) and standard single mode fiber (with higher water content) (from H. Kogelnik).

Figure 3.13 .................................................................................................................. 113
The group delay dispersion (GDD) of bulk prism materials with low water content.

Figure 3.14 .................................................................................................................. 114
Group delay dispersion (GDD) of Cr$^{4+}$:YAG and fused silica glass of different thicknesses.

Figure 3.15 .................................................................................................................. 115
Schematic of the Z-fold cavity with a 2 cm Cr$^{4+}$:YAG crystal and two fused silica prisms. Mirrors M1 - M3 are all high-reflector mirrors and OC is an output coupling mirror.

Figure 3.16 .................................................................................................................. 116
Typical modelocked spectrum from Cr$^{4+}$:YAG laser using fused silica prisms for dispersion compensation.

Figure 3.17 .................................................................................................................. 117
Autocorrelation measurement (black line) and sech fit (gray dashed line) of 54 fs pulses from a fused silica prism dispersion compensated Cr$^{4+}$:YAG laser.

Figure 3.18 .................................................................................................................. 119
Schematic of Bragg reflecting quarter wave stacks, chirped mirrors, and double-chirped mirrors (from reference [134]).

Figure 3.19 .................................................................................................................. 120
Index of refraction of the layers comprising the double-chirped mirrors (DCMs) designed for use in a Cr$^{4+}$:YAG laser.

Figure 3.20 .................................................................................................................. 121
Measured reflectivity of double-chirped mirror (DCM) designed for use within a Cr$^{4+}$:YAG laser.

Figure 3.21 .................................................................................................................. 122
Designed (gray line) and measured (black line) group delay dispersion (GDD) of the double-chirped mirrors (DCMs) used in a Cr\(^{4+}\):YAG laser.

Figure 3.22 ........................................................................................................ 123
Net cavity group delay dispersion (GDD) of laser cavities consisting of a number of double-chirped mirror (DCM) reflections and either a (a) 1 cm or (b) 2 cm Cr\(^{4+}\):YAG crystal.

Figure 3.23 ........................................................................................................ 124
Group delay dispersion (GDD) of double-chirped mirrors for off-axis reflection. The GDD oscillations grow as a function of angle.

Figure 3.24 ........................................................................................................ 126
Transmission spectra through air exhibiting large water absorption lines between 1300 and 1500 nm. (From reference [148])

Figure 3.25 ........................................................................................................ 126
The group delay dispersion (GDD) from water vapor in a Cr\(^{4+}\):YAG laser. (From reference [149])

Figure 3.26 ........................................................................................................ 127
Power as a function of time after the onset of nitrogen purging of the laser beam path.

Figure 3.27 ........................................................................................................ 128
Purged (gray line) and unpurged (black line) spectrum of saturable absorber modelocked Cr\(^{4+}\):YAG laser.

Figure 3.28 ........................................................................................................ 129
Schematic of a prismless Cr\(^{4+}\):YAG laser used to generate short pulses. The cavity consists of a 2 cm laser crystal, an output coupler (OC), 3 double-chirped mirrors (DCMs) adding 6 DCM bounces per cavity roundtrip (M1-M3, in black), and one unchirped high reflector (M4, in gray).

Figure 3.29 ........................................................................................................ 130
Picture of the prismless Cr\(^{4+}\):YAG laser.

Figure 3.30 ........................................................................................................ 131
Group delay dispersion (GDD) of 2 passes through a 2 cm Cr\(^{4+}\):YAG crystal (light gray), 6 reflections from double-chirped mirrors (DCMs) (gray), and all optical elements in the laser cavity (black). The net cavity GDD curve includes reflections from 6 DCMs, a single unchirped quarter-wave stack high reflector, and an output coupler.

Figure 3.31 ........................................................................................................ 132
Optical power spectrum of a Cr\(^{4+}\):YAG pulse. The black line corresponds to a linear scale (left axis) and the gray line corresponds to a logarithmic scale (right
axis). The full-width at half maximum is 190 nm, with a peak at 1450 nm.

Figure 3.32 .................................................................................................................. 133
Measured autocorrelation function from an interferometric two-photon absorption autocorrelator (line) and fit by a pulse-retrieval algorithm (dots). A pulse width of 19.5 fs is calculated by the pulse-retrieval algorithm, 18.3 fs by assuming sech shaped pulses, and 17.0 fs by assuming gaussian shaped pulses.

Figure 3.33 .................................................................................................................. 134
HeNe laser interference fringes measured by a Si p-i-n photodiode through the autocorrelator interferometer. The peak of each interference fringe is marked by a black circle.

Figure 3.34 .................................................................................................................. 135
Plot mapping the time delay introduced by a speaker in the autocorrelator from scope units to fs. The data was fit by a third-order polynomial.

Figure 3.35 .................................................................................................................. 136
Schematic of the Cr$^{4+}$:YAG autocorrelation experiment. Output light from the laser is chirp compensated by a pair of BaF$_2$ prisms. The pulsewidth is then measured with a Si p-i-n detector two photon absorption based autocorrelator.

Figure 3.36 .................................................................................................................. 137
Series of interferometric autocorrelation traces taken with different quantities of extracavity group delay dispersion (GDD) compensation. As the chirp of the pulses are compensated, the wings of the autocorrelation traces become flat.

Figure 3.37 .................................................................................................................. 138
Autocorrelation measured pulse widths, fit assuming a sech pulse shape, for different values of extracavity group delay dispersion (GDD) introduced by two BaF$_2$ prisms and one double-chirped mirror (DCM) bounce.

Figure 3.38 .................................................................................................................. 139
Group delay dispersion (GDD) of output couplers and unchirped high reflector mirrors used within the ultrafast prismless Cr$^{4+}$:YAG laser.

Figure 3.39 .................................................................................................................. 140
Spectra and corresponding interferometric autocorrelation of ultrashort pulses generated by the prismless Cr$^{4+}$:YAG laser with different combinations of output couplers and unchirped high reflector mirrors. (a) and (b) correspond to a laser using output coupler OC3 and high reflector HR1, (c) and (d) correspond with OC3 and HR2, and (e) and (f) correspond with the use of OC5 and HR1.

Figure 3.40 .................................................................................................................. 141
(a) Optical spectrum and (b) interferometric autocorrelation of pulses from a prismless Cr$^{4+}$:YAG laser containing 7 double-chirped mirror (DCM) bounces per cavity roundtrip.
Figure 3.41

Schematic of a prismless Cr\textsuperscript{4+}:YAG laser used to generate short pulses and eliminate the effects of high angle incidence from double-chirped mirrors (DCMs). The cavity consists of a 2 cm laser crystal, an output coupler (OC), 2 unchirped cavity fold mirrors (M1 and M2, in gray), 3 DCMs adding 6 DCM bounces per cavity roundtrip (M3-M5, in black), and one unchirped high reflector (M6, in gray).

Figure 3.42

Optical spectrum from modelocked prismless Cr\textsuperscript{4+}:YAG laser.

Figure 3.43

Measured interferometric autocorrelation (black line) and 32 fs sech fit (gray dots).

Figure 3.44

Schematic of a Cr\textsuperscript{4+}:YAG laser used in attempts to generate ultrashort optical pulses. The cavity consists of a 1 cm laser crystal, an output coupler (OC), 2 double-chirped mirrors (DCMs) (M1 and M2, in black), and one unchirped high reflector (M3, in gray).

Figure 3.45

Picture of the Cr\textsuperscript{4+}:YAG laser.

Figure 3.46

Tuning curve of the 1 cm Cr\textsuperscript{4+}:YAG crystal based laser.

Figure 3.47

Group delay dispersion (GDD) of two possible cavity designs including a 1 cm Cr\textsuperscript{4+}:YAG crystal, 4 (gray lines) or 5 (black lines) double-chirped mirror (DCM) bounces, and 15 mm (gray line) or 25 mm (black line) of propagation through BaF\textsubscript{2} prisms.

Figure 3.48

Gain profile for a Gaussian pump beam and several levels of pump absorption saturation.

Figure 3.49

Spectra from the Cr\textsuperscript{4+}:YAG and Ti:Sapphire lasers. Together, the spectra span two octaves in wavelength.

Figure 4.1

Absorption fraction for transmission through a typical saturable absorber material.

Figure 4.2

Schematic of a two level system. The two levels have an energy difference of
$E_2 - E_1$, have a stimulated excitation rate of $W_{12}$, a stimulated emission rate of $W_{21}$, a spontaneous excitation rate of $w_{12}$, and a spontaneous emission rate of $w_{21}$.

Figure 4.3 .................................................................................................................. 160
Schematic of two possible saturable Bragg reflector (SBR) designs for Cr$^{4+}$:YAG lasers. Each SBR consists of a high-dielectric contrast quarter wave stack mirror, augmented with a half wave layer containing a saturable absorber.

Figure 4.4 .................................................................................................................. 163
Calculated reflectivity of a high-dielectric contrast mirror consisting of a 7 pair GaAs/Al$_x$O$_y$ quarter-wave dielectric stack.

Figure 4.5 .................................................................................................................. 164
Calculated group delay dispersion (GDD) of a high-dielectric contrast mirror consisting of a 7 pair GaAs/Al$_x$O$_y$ quarter-wave dielectric stack.

Figure 4.6 .................................................................................................................. 165
Measured and calculated reflection for a 4 pair Si/SiO$_2$ Bragg mirror centered at 1475 nm.

Figure 4.7 .................................................................................................................. 166
Electric field energy amplitude and index of refraction of the designed Bragg reflector (SBR) mirror consisting of a GaAs/Al$_x$O$_y$ high-index contrast mirror and an InGaAs/InP quantum well.

Figure 4.8 .................................................................................................................. 168
Photoluminescence (PL) from a broadband SBR.

Figure 4.9 .................................................................................................................. 169
Schematic of the pump-probe apparatus used to study the saturation dynamics of the saturable Bragg reflector.

Figure 4.10 ............................................................................................................... 170
Pump-probe traces of the broadband saturable Bragg reflector at 1540 nm.

Figure 4.11 ............................................................................................................... 172
Scanning electron micrograph images of an (a) unoxidized and (b) oxidized SBR structure. After oxidation, the AlAs layers are converted to polycrystalline Al$_x$O$_y$, which appears to be granular.

Figure 4.12 ............................................................................................................... 173
Scanning electron micrograph of an unsuccessfully oxidized SBR structure. Half of the AlAs layers have been converted to Al$_x$O$_y$ while the remainder have not.

Figure 4.13 ............................................................................................................... 174
Schematic of a Cr$^{4+}$:YAG laser cavity consisting of 3 10 cm radius of curvature
double-chirped mirrors (M1 - M3), an output coupler (OC), and an saturable Bragg reflector (SBR) end mirror.

Figure 4.14 ........................................................................................................ 175
Spectra of the saturable absorber modelocked Cr$^{4+}$:YAG laser tuned from 1400 to 1525 nm with a birefringent tuning plate.

Figure 4.15 ........................................................................................................ 176
Pulse spectrum from a self-started Cr$^{4+}$:YAG laser.

Figure 4.16 ........................................................................................................ 177
Interferometric autocorrelation of a self-started Cr$^{4+}$:YAG laser.

Figure 5.1 ........................................................................................................ 184
(a) A schematic of a typical electronic crystal. Each circle represents the placement of an ion in the crystal lattice. (b) A schematic of a possible three dimensional photonic crystal. Each color represents a material with a different index of refraction.

Figure 5.2 ........................................................................................................ 185
(a) The electronic bandstructure of Si. (b) The photonic band structure of a three-dimensional photonic crystal.

Figure 5.3 ........................................................................................................ 188
Schematic of a (a) monorail and (b) airbridge geometry photonic crystal microcavity. The microcavities consist of a defect region of length $a_d$ between one-dimensional photonic crystal lattice of air holes of spacing $a$ and diameter $d$ within a GaAs waveguide of width $w$ and thickness $T_{GaAs}$. The waveguides rest upon a Al$_x$O$_y$ cladding layer of thickness $T_{Oxide}$ on top of a GaAs substrate.

Figure 5.4 ........................................................................................................ 189
Electric field intensity of the single optical mode in a 200 by 500 nm GaAs waveguide on a Al$_x$O$_y$ pedestal. (Courtesy of Christina Manolatou).

Figure 5.5 ........................................................................................................ 191
Schematic of a GaAs waveguide (a) without and (c) with a photonic crystal, and the corresponding dispersion relations of the waveguide modes (b) and (d). TE-like modes are indicated by black triangles, while TM-like modes are indicated by hollow squares (from reference [186]).

Figure 5.6 ........................................................................................................ 192
Finite-difference time-domain calculation of transmission through a (a) monorail and (b) airbridge geometry waveguide microcavity. In each case, a resonant defect state is designed to be at 1550 nm (from reference [235]).

Figure 5.7 ........................................................................................................ 193
Electric field intensity inside a waveguide microcavity outside and within the
bandgap of the photonic crystal (from reference [235]).

Figure 5.8 ............................................................................................................ 194
Electric field intensity inside a waveguide microcavity on resonance with a de-
fect state within the photonic crystal (from reference [235]).

Figure 5.9 ............................................................................................................ 195
Quality factor (Q) and transmission through a photonic crystal microcavity with
different number of holes (from reference [235]).

Figure 5.10 ............................................................................................................ 198
Scanning electron micrograph (SEM) of a (a) monorail and (b) airbridge
waveguide microcavity.

Figure 5.11 ............................................................................................................ 199
Experimental set-up designed to study the transmission through waveguide de-
vices.

Figure 5.12 ............................................................................................................ 200
Schematic of a NaCl:OH\(^-\) color center laser used as an optical source for trans-
mittance measurements.

Figure 5.13 ............................................................................................................ 201
Calibration of the color center laser output wavelength to birefringent tuning
plate stepper motor position.

Figure 5.14 ............................................................................................................ 201
Color center laser tuning curve.

Figure 5.15 ............................................................................................................ 203
Top views of light propagating through a high-dielectric contrast waveguide.
Each image is acquired from an infrared sensitive vidicon camera. The light
that is imaged is scattered from the waveguide by defects. The top image shows
large scattering at the input of the waveguide from modal mismatch between the
waveguide and fiber coupling lenses.

Figure 5.16 ............................................................................................................ 204
Bird’s eye and output end views of light output from a high-dielectric contrast
waveguide as imaged by an infrared sensitive vidicon camera.

Figure 5.17 ............................................................................................................ 205
Silicon CCD camera images of two photon absorption (TPA) pumped photolu-
minescence (PL) within a high-index contrast waveguide. Images are taken (a)
without and (b) with light coupled into the waveguide. The CCD camera is sen-
sitive to PL from GaAs, but not to the input light at 1.55 \(\mu\)m.

Figure 5.18 ............................................................................................................ 207
Transmission spectra through three monorail devices. Each monorail has a dif-
ferent defect length, resulting in a different transmission peak wavelengths. The maximum transmission of each peak is normalized to unity.

Figure 5.19 .............................................................................................................. 208
Transmission spectra measured through two different PBG air-bridge micro-
cavities. The long and short wavelength resonances have defect center-to-center
lengths of 632 nm and 703 nm respectively. The maxima of the resonance
peaks are normalized to unity.

Figure 5.20 .............................................................................................................. 209
Top view of an illuminated airbridge microcavity resonator. (a) Off resonance
and within the bandgap; the photonic crystal reflects the light. (b) On reso-
nance; light is transmitted through the microcavity resonator. The light is scat-
ttered out of the waveguides primarily from surface roughness, and is viewed
with an infrared vidicon camera.

Figure 5.21 .............................................................................................................. 210
Transmission spectrum of a PBG air bridge sample where the long wavelength
band edge has been shifted into the experimentally accessible wavelength
range. The band edge is shown on a log plot, and exhibits a 27 dB suppression
of transmission within the bandgap compared to outside the bandgap.

Figure 5.22 .............................................................................................................. 211
Transmission spectrum of an airbridge microcavity exhibiting both a resonance
and a band edge within the experimentally accessible wavelength range. Trans-
mission outside the bandgap is normalized to unity, revealing 72% relative
transmission on resonance, with a cavity Q of 230.

Figure 6.1 .............................................................................................................. 215
Schematic of extraction of emitted light from a high index of refraction sub-
strate into air. Most light is trapped within the substrate by total internal reflec-
tion.

Figure 6.2 .............................................................................................................. 216
Schematic of a light emitting diode (LED) designed to emit light at 980 nm. A
single InGaAs quantum well is sandwiched between InGaP cladding layers. A
low index Al_xO_y spacer layer separates the high index quantum well region
from a 6 period GaAs/Al_xO_y distributed Bragg reflector (DBR) mirror. A pho-
tonic crystal, or other geometric structure, is embedded within the LED device.

Figure 6.3 .............................................................................................................. 217
Emission of light emitting diode (LED) near 980 nm and calculated reflectance
of GaAs/Al_xO_y high index contrast distributed Bragg reflector (DBR).

Figure 6.4 .............................................................................................................. 218
Rate of oxidation of AlGaAs into Al_xO_y for different temperatures and Ga con-
centrations (from reference [260]).
Figure 6.5 Scanning electron micrographs (SEMs) of fabricated light emitting diode (LED) structures. (a) A 50 μm diameter circular mesa, with a 12.5 by 12.5 μm square photonic crystal region. (b) Side view of LED layers. (c) Angled top view of photonic crystal lattice. (d) Side view of photonic crystal lattice holes.

Figure 6.6 Schematic of experimental set-up used to image photoluminescence (PL) from an LED sample and measure PL spectrum as a function of position on the sample.

Figure 6.7 Schematic of a tunable Ti:Sapphire laser. The laser is in an X-fold configuration. Mirrors M1 - M3 are all broadband high reflector mirrors and OC is an output coupler. The laser crystal is pumped by an Argon ion laser at 488 nm. A rotatable birefringent tuning plate provides the wavelength dependent loss necessary to tune the laser.

Figure 6.8 Silicon CCD image of photoluminescence (PL) from an unpatterned light emitting diode (LED) mesa.

Figure 6.9 Image of the reflection of pump light at 810 nm from a light emitting diode (LED).

Figure 6.10 Spectrally resolved photoluminescence (PL) from an unpatterned light emitting diode (LED).

Figure 6.11 Schematic of extraction enhancement process from an a disordered microstructure on top of a planar waveguide.

Figure 6.12 Scanning electron micrograph of disordered array of holes in a light emitting diode (LED) mesa.

Figure 6.13 Silicon CCD image of photoluminescence (PL) from a 50 by 50 μm light emitting diode (LED) mesa. A significant enhancement in PL is visible from a 30 by 30 μm region with a disordered array of holes.

Figure 6.14 Photoluminescence (PL) from an unpatterned light emitting diode (LED) mesa (no holes), and from a region patterned with a disordered array of holes (random holes). There is a significant enhancement in the intensity of extracted light.
from the patterned region.

Figure 6.15 ........................................................................................................ 229
Spectral enhancement, the ratio between the photoluminescence (PL) intensity of the enhanced patterned region and a reference unpatterned region, as a function of wavelength.

Figure 6.16 ........................................................................................................ 230
Silicon CCD image of the pump reflection from a 50 by 50 μm light emitting diode (LED) mesa, with a 30 by 30 μm patterned region with a disordered array of holes.

Figure 6.17 ........................................................................................................ 231
(a) Scanning electron micrograph of a photonic crystal embedded within a light emitting diode (LED). The photonic crystal consists of a triangular lattice of holes. (b) Schematic diagram of a triangular lattice in real space (black dots) and the corresponding triangular lattice in reciprocal space (gray dots). The central dot is part of both lattices. The first order Brillioun zone is indicated with a dotted line.

Figure 6.18 ........................................................................................................ 232
Calculated dispersion relations for a triangular lattice photonic crystal. Waveguided modes are indicated by black lines. Radiation modes exist in a region above the light line (the dispersion line for freely propagating light) because radiating modes can carry energy propagating orthogonal to the waveguide plane. Waveguides in the region of radiation modes are called leaky guided modes, and can phase match radiation with a lifetime.

Figure 6.19 ........................................................................................................ 233
Electric field profile of the degenerate optical mode at the Γ-point for a triangular lattice photonic crystal.

Figure 6.20 ........................................................................................................ 234
Schematic of an array of light emitting diode (LED) mesas. The border mesas do not contain any microstructures, while the central mesas contain embedded triangular lattice photonic crystals. In one direction, the lattice constant (a) is varied, while in the second direction the hole diameter (d) is varied.

Figure 6.21 ........................................................................................................ 235
Silicon CCD image of (a) photoluminescence and (b) pump reflection from and light emitting diode (LED) mesa with a triangular photonic crystal lattice designed to couple light from waveguided modes in the LED into radiation.

Figure 6.22 ........................................................................................................ 236
Spectrally resolved Silicon CCD images of photoluminescence (PL) from a light emitting diode (LED) mesa. Each image was taken with a 10 nm full width at half maximum chromatic filter to isolate PL of a particular wavelength. Par-
particularly strong extraction efficiency enhancements from the central photonic crystal region are observed near 925 and 950 nm.

Figure 6.23 ........................................................................................................ 237
Spectrally resolved Silicon CCD images of photoluminescence (PL) from a light emitting diode (LED) mesa. Each image was taken with a 10 nm full width at half maximum chromatic filter to isolate PL of a particular wavelength. In each image, the PL was optically excited by a small pump spot on the left-hand side of the unpatterned border region of each mesa.

Figure 6.24 ........................................................................................................ 238
Concatenation of Silicon CCD images of photoluminescence (PL) from a 100 by 100 μm light emitting diode (LED) mesa. In each of the 4 images, the mesa was excited in the unpatterned border region close to the central 50 by 50 μm photonic crystal region. Light is coupled from the unpatterned region guided modes into the patterned region's leaky guided modes. The propagation directions of the light coupled into the photonic crystal region reflect the 6-fold symmetry of the triangular lattice.

Figure 6.25 ........................................................................................................ 239
Photoluminescence (PL) from an unpatterned light emitting diode (LED) mesa (no holes), and from a region patterned with a photonic crystal (with triangular lattice). There is a significant enhancement in the intensity of extracted light from the patterned region.

Figure 6.26 ........................................................................................................ 240
Spectral enhancement, the ratio between the photoluminescence (PL) intensity of the enhanced patterned region and a reference unpatterned region, as a function of wavelength. A spectral intensity over 100 is observed at 925 nm.

Figure 6.27 ........................................................................................................ 241
Power dependence of the photoluminescence (PL) from mesa (4,4).

Figure 6.28 ........................................................................................................ 243
Photoluminescence (PL) from several unpatterned light emitting diode (LED) mesas patterned with photonic crystals of with different hole diameters.

Figure 6.29 ........................................................................................................ 244
Spectral enhancement calculated from the ratio of photoluminescence (PL) from several unpatterned light emitting diode (LED) mesas patterned with photonic crystals of with different hole diameters and an unpatterned reference LED mesa.

Figure 6.30 ........................................................................................................ 245
Integrated enhancement of photoluminescence (PL) from several light emitting diode (LED) mesas. Mesa (4,1) is unpatterned, mesas (4,2) through (4,9) have a photonic crystal of increasing hole diameter, and mesa (1,4) contains a disor-
edered array of holes.

Figure 6.31 .................................................................................................................. 246
Angular dependence of photoluminescence (PL) from an unpatterened reference light emitting diode (LED) mesa. PL spectra are measured with the LED mesa rotated 0, 10, 15, and 25 degrees from normal.

Figure 6.32 .................................................................................................................. 247
Angular dependence of photoluminescence (PL) from a light emitting diode (LED) mesa patterned with a photonic crystal. PL spectra are measured with the LED mesa rotated 0, 10, 15, and 25 degrees from normal.

Figure 6.33 .................................................................................................................. 248
Angular dependence of spectral enhancement of photoluminescence (PL) from a light emitting diode (LED) mesa patterned with a photonic crystal. PL enhancements are measured with the LED mesa rotated 0, 10, 15, and 25 degrees from normal.

Figure 6.34 .................................................................................................................. 249
Scanning electron micrograph (SEM) of a photonic crystal fabricated to increase the pumping efficiency of a light emitting diode (LED).

Figure 6.35 .................................................................................................................. 250
Band structure of photonic crystal designed to increase pumping efficiency at 810 nm.

Figure 6.36 .................................................................................................................. 251
Silicon CCD images of a light emitting diode (LED) with a photonic crystal designed to increase optical pumping efficiency at 810 nm. Because 810 nm light is coupled into the waveguide modes, there is a dark spot in the pump reflection in the presence of the photonic crystal.

Figure 6.37 .................................................................................................................. 252
Composite of 8 Silicon CCD images of a small pump spot size measurement of photoluminescence (PL) from a photonic crystal light emitting diode (LED). Each image was taken with a small pump spot within the photonic crystal region. Bright patterns of PL are visible with the 6-fold symmetry corresponding with the triangular lattice photonic crystal.

Figure 6.38 .................................................................................................................. 253
Soccer ball geometry for combining pump input and light emitting diode (LED) output photonic crystal regions.
LIST OF TABLES

Table 2.1 ......................................................................................................................... 40
Material properties of several gain media used to generate ultrashort laser pulses.

Table 2.2 ......................................................................................................................... 89
Values of G2(\tau) and B(\tau) for Gaussian and sech shaped pulse envelopes [108].

Table 3.1 ......................................................................................................................... 104
Tabulation of data from pump power, output power plots.

Table 4.1 ......................................................................................................................... 167
Properties of four low-index contrast SBR designs tried with a Cr^{4+}:YAG laser.

Table 4.2 ......................................................................................................................... 171
Characteristics of all saturable Bragg reflector growths in the effort to realize a
broadband structure.

Table 5.1 ......................................................................................................................... 206
Dimensions of fabricated waveguide microcavity devices.
CHAPTER 1: INTRODUCTION

1.1 Time and Length Scales of Light

Physical phenomena in the macroscopic world are typically characterized in terms of meters, kilograms, and seconds, the accepted SI (Système International) units for length, mass, and time. These units, while not of any fundamental significance, are particularly convenient for descriptions of events in everyday life. Light, however, does not operate on our time-scale.

Electromagnetic radiation operates over a wide range of length and time-scales, fundamentally related by:

\[ c = \lambda \nu \]  

(1.1)

where \( c \) is the speed (\( 3 \times 10^8 \) m/s), \( \lambda \) is the wavelength, and \( \nu \) is the frequency of light. Because \( c \) is a universal constant for massless particles such as photons, the length and time-scales of freely propagating light are proportional. This thesis discusses research into the generation of few-cycle optical pulses and photonic crystal manipulation of light on the time and length-scale (femtoseconds and microns) of near-infrared light.
1.2 Ultrashort Pulse Generation

Lasers can be designed to generate pulses of light with short time durations, large peak intensities, and broad spectral bandwidths. It becomes more difficult to create shorter pulses as the pulsewidth approaches the natural time-scale of light. In fact, radiative optical pulses have a minimum pulsewidth of a single optical cycle (one wavelength long in space). To propagate, electromagnetic radiation must obey Maxwell’s wave equation. Solutions of the wave equation are restricted to oscillatory sine-waves and their Fourier sums. Sub-wavelength pulses can be generated in the temporal near-field, but will contain a non-zero dc electric field component that cannot radiate. The pulse will transform into a single-cycle or longer pulse in the far-field. The time-scale associated with a single wavelength cycle pulse, $\tau \sim \lambda/c$, thus represents a lower limit to optical temporal pulsewidths. The corresponding spread in the light’s frequency or energy is given by the uncertainty principle $\Delta\nu \Delta\tau > 1/4\pi$. The difficulty of controlling light over such a large bandwidth can limit the ability to generate short pulses.

Nevertheless, ultrafast lasers are now able to generate pulses in the few optical-cycle regime by broadband control of intracavity group velocity dispersion. A plot of the record pulsewidth and corresponding bandwidth of Ti:Sapphire lasers over time is shown in Figure 1.1. As pulsewidths approach the single-cycle limit, it is necessary to manipulate optical

![Figure 1.1](image) Minimum pulsewidth and fractional bandwidth of pulses generated directly from a Ti:Sapphire laser over time.
bandwidths covering over an octave (λ to 2λ) of spectrum. The principle technological advancements that have allowed the generation of shorter optical pulses have been the development of broadband dispersion compensating and reflective optical elements. The increase in bandwidth of optical elements has generally relied on an improvement in the techniques used to fabricate high index of refraction contrast material systems.

1.3 **Wavelength-Scale Light Manipulation**

In addition to use within broadband optical elements, high index contrast materials can be employed to manipulate light over smaller length-scales. Structures consisting of a spatially-dependent index of refraction are commonly used to modify the propagation of light. Glass fibers, for example, guide light in an optical mode supported by a high index core. Periodic dielectric stack mirrors exhibit high reflection over a broad bandwidth with extremely low loss. The length-scale over which light is confined in a waveguide or resonator and the interaction length of light reflecting from a periodic structure are determined by the index contrast of the dielectric structure.

In the infinite index of refraction contrast limit, resonators can theoretically confine light to volumes as small as (λ/2n)^3, where n is the cavity index of refraction. A realistic resonator surrounded by periodic dielectric mirrors or low-index cladding materials will have a mode which has a finite penetration into the mirror or low-index region. These modes will therefore have a volume greater than the theoretical minimum. A goal of light confinement is to contain radiation within a resonant cavity of volume ~ (λ/2n)^3 for as long as possible without suffering appreciable loss to radiation. This becomes difficult as the modal volume approaches the half wavelength cubed because light confined to such small volumes increasingly phase-matches to radiative free-space modes. On the other hand, if radiation is the desired decay mechanism, periodic dielectric structures can be designed to phase-match these processes efficiently.

1.4 **Optical Communication**

Optical communication systems require short pulses of light to interact with matter in a controlled manner over useful length-scales without appreciable loss. Glass fiber-based communication systems commonly transmit data at wavelengths between 1530 to 1580 nm because of the extremely low fiber loss and the availability of erbium-doped fiber amplifiers.
(EDFAs) covering the spectral range. Recently, by removing water absorption lines, fibers with low loss from 1300 to 1600 nm have become commercially available, opening a substantially larger bandwidth for high bit-rate data transmission.

The maximum transmission rate of a given data channel will be determined by the repetition rate of the data pulses (1's and 0's for binary coding schemes). The pulsewidth must be short enough to prevent pulse cross-talk \( t_{\text{pulse}} < t_{\text{rep-rate}} \) and long enough to prevent filtering due to the channel's allocated bandwidth. Data from different sources can then be multiplexed in together in time using time division multiplexing (TDM). Alternatively, the same amount of bandwidth can be split among several lower repetition rate data channels transmitted at different wavelengths by a scheme known as wavelength division multiplexing (WDM). Because nonlinear interactions develop from interactions between different data streams or from the self-interactions of a single data stream, the amount of data transmitted over a fiber optic system cannot be increased indefinitely by using larger optical powers. To transmit information faster, communication systems are migrating toward larger WDM channel counts, faster transmission rates over TDM channels, or some combination of the two.

Implementation of inexpensive, compact communication systems requires that complex light manipulation be performed over small volumes on fast time-scales. The wavelength of light currently used for fiber optic communication is 1.5 μm. The corresponding time-scale for a single-cycle optical pulse is 5 fs. Semiconductor devices can be used to generate, detect, filter, or switch light in a small volume. Because of their widespread use in electronics, semiconductor devices can be inexpensive and densely integrated with a combination of optical and electrical functions. Silicon and gallium arsenide, common semiconductor materials with low intrinsic optical loss for energies below their bandgap, have indices of refraction of ~3.5 at 1.5 μm. The characteristic length-scale of light in these materials is therefore less than 500 nm.

### 1.5 Thesis Outline

This thesis presents the results of research toward the development of a few-cycle optical pulse source with spectrum from 1300 to 1600 nm and the study of photonic crystal semiconductor devices for wavelength-scale light manipulation.

Ultrafast laser design theory is discussed in detail in Chapter 2. Cr\(^{4+}\):YAG is shown to be suitable as a broadband laser material with gain from 1300 to 1500 nm and capable of generating femtosecond pulses. Modelocking is reviewed as a technique to convert the broad-
band gain of a laser material into ultrashort optical pulses. Design considerations for the creation of a suitable laser cavity are discussed, and autocorrelation is introduced as a method to measure the pulsewidth of short optical pulses.

Pulses as short as 20 fs were generated by a Cr$^{4+}$:YAG modelocked solid-state laser, with spectrum extending from 1140 to > 1700 nm. Chapter 3 discusses the laser used to generate these ultrashort pulses, as well as the careful group velocity dispersion performed by double-chirped mirrors (DCMs) that was required to produce the pulses. To the author’s knowledge, these four-cycle pulses are the shortest produced to date by a Cr$^{4+}$:YAG laser.

It is often difficult to initiate ultrashort pulse generation in a passively modelocked lasers such as the 20 fs Cr$^{4+}$:YAG. Chapter 4 reports on the development of a broadband saturable absorber mirror that is capable of self-starting ultrafast optical pulses. The saturable absorber mirrors were based upon a novel 7 pair GaAs/Al$_x$O$_y$ high-index-contrast broadband mirror, and supported 35 fs pulses from the Cr$^{4+}$:YAG laser, which are the shortest self-started pulses for this laser system.

Chapter 5, shows how 1.5 μm light is confined within a high-index-contrast GaAs/Al$_x$O$_y$ waveguide microcavity resonator using one-dimensional photonic bandgap (PBG) crystals. Resonant transmission is studied through microcavities with modal volumes as small as $2(\lambda/2n)^3$ and measured quality factors as high as 360. These are among the smallest optical cavities ever studied and could be used in the future to modify the emission rate of active materials.

Two-dimensional photonic crystals can alter the coupling between planar-waveguide and radiative free-space modes. Chapter 6 discusses the use of a triangular lattice photonic crystal to increase the output coupling efficiency of optically pumped light emitting diodes (LEDs). Eight-fold overall enhancements and over 100-fold single wavelength enhancements of collected photoluminescence (PL) were observed from a photonic crystal LED. The inverse effect is shown to enhance input coupling of the LED’s optical pump.

Finally, Chapter 7 reviews conclusions reached from the preceding chapters and proposes future directions in the development and application of short pulse sources near 1.5 μm and the study of ultrasmall photonic crystal optical devices.
CHAPTER 2: ULTRAFAST LASER DESIGN

2.1 Introduction

A Cr\textsuperscript{4+}:YAG laser has been developed for the generation of both broadly tunable cw light and ultrafast laser pulses near 1.5 μm. Cr\textsuperscript{4+}:YAG is chosen as the gain media because of its broadband gain and material properties that are favorable for modelocking. Modelocking and Kerr lens modelocking in particular are discussed as techniques to generate short pulses. A Z-fold cavity was designed to optimize the modelocking strength. Autocorrelation is reviewed as a method for measuring the pulsewidth of ultrafast pulses.

2.2 Background

Modelocked lasers are used to generate optical pulses with short time durations, broad bandwidths, and high average powers. Short-pulse sources are ideal for time-resolved studies of ultrafast phenomena and devices [1], as optical clocks with precise timing at the cavity repetition rate [2], and as sources for ultra-high speed optical communications [3]. Coherent broadband sources can be exploited for spectroscopy [4], to generate synchronized multi-wavelength optical sources [5], or to provide a comb of regularly spaced spectral lines for use as a frequency standard [6 - 10]. Finally, the high instantaneous power of short pulses allows one to perform nonlinear microscopy [11], micromachine devices [12, 13], study intense light-matter interactions [14 - 16], and even accelerate electrons to relativistic velocities [17].
Because each ultrafast laser application has specific wavelength, pulse duration, and pulse energy requirements, the number of ultrafast sources under development is as diverse as the number of applications.

This chapter begins by motivating the use of Cr$^{4+}$:YAG as a gain medium for generating ultrashort laser pulses. Modelocking is introduced as a general technique for generating such short pulses in a laser system, and Kerr lens modelocking (KLM) is discussed as a particular realization of modelocking that can produce short pulses from a Cr$^{4+}$:YAG laser. The pulsewidth generated by a modelocked laser will be determined by the balance between pulse broadening and pulse shortening mechanisms. Short optical pulses will only be generated if intracavity group delay dispersion, the tendency of short optical pulses to spread in time due to chromatic dispersion, is carefully compensated. Propagation of optical pulses through dispersive media is briefly reviewed. The design of a suitable optical resonator to maximize the strength of KLM is outlined. Several calculations of optical resonators are presented and were performed with the assistance of Hanfei Shen. Once short pulses are generated, the pulsewidth must be measured. Autocorrelation as a pulsewidth measurement technique is reviewed. The autocorrelator for the experiments discussed were designed and built with the help of Juliet Gopinath.

2.3 Cr$^{4+}$:YAG

Transition-metal-doped materials are popular as ultrafast laser media because they have large gain bandwidths, a Stokes-shifted gain spectra, and a high nonlinear coefficient which enables efficient Kerr lens modelocking [18, 19]. Ti:Sapphire, the most common example of such a laser, has become the workhorse of ultrafast optics. Ti:Sapphire has an exceptionally large gain bandwidth centered around 800 nm and a multitude of desirable material properties. Pulses as short as 5 fs, with a correspondingly broad spectrum from 580 to 1160 nm, have been generated directly from Ti:Sapphire lasers [20]. It is likely that the Ti:Sapphire laser spectrum is cut off at wavelengths above 1200 nm because of the limited cavity mirror bandwidth.

Several methods are used to generate short pulses in the wavelength range from 1300 to 1600 nm, which are of particular importance for optical communication applications. Parametric down-conversion of photons from a Ti:Sapphire pump can be used to generate short pulses in the near infrared. The operation of optical parametric oscillators (OPOs) and optical parametric amplifiers (OPAs) are based upon this principle. Commercial OPOs are sold
which can generate 100 fs pulses from a synchronous Ti:Sapphire pump [21]. The shortest pulses reported near 1.5 μm to date are 14.5 fs pulses generated by parametric down-conversion of 18 fs Ti:Sapphire pulses [22]. This approach, which requires an amplified Ti:Sapphire pump, was only able to produce a 1 kHz repetition rate. With the proper gain media, ultrashort optical pulses can be generated at a much higher repetition rate directly from a laser. Both color-center [23] and Er-doped glass lasers [24] have been used to generate short pulses near 1500 nm. Color-center lasers, however, must be chilled to liquid nitrogen temperatures to avoid loss of the gain-providing color-center defects and must be kept in a vacuum to avoid corrosion from water in air. Er-doped glass is a gain media used to generate short pulses from fiber-lasers [25] and bulk solid-state lasers [24] but has a limited gain bandwidth. Pulses as short as 54 fs have been produced by externally compressing fiber laser pulses with a dispersion-decreasing fiber through adiabatic soliton compression [26].

Cr$^{4+}$:YAG is a promising ultrafast laser crystal within the wavelength range from 1300 to 1600 nm because of its exceptionally large gain bandwidth and relatively good material properties. There are significant advantages of Cr$^{4+}$:YAG over either color center or Er-doped glass lasers. Unlike color-center lasers, Cr$^{4+}$:YAG can be operated at room temperature in air, and can be easily cooled with flowing water. The gain bandwidth is substantially greater than that for Er:Glass. Finally, Cr$^{4+}$:YAG has a convenient absorption peak between 950 to 1100 nm, which can be pumped by readily available all-solid-state Nd:YAG or Nd:YVO$_4$ lasers. Table 2.1 summarizes some of the characteristics of Cr$^{4+}$:YAG, Ti:Sapphire, and Cr:Forsterite [27].

A combination of spectroscopic measurements and theoretical modeling have elucidated the electronic energy levels of carriers in Cr$^{4+}$:YAG crystals [28 - 31]. The active light emitting elements are believed to be tetrahedrally coordinated Cr$^{4+}$ ions. It is difficult to dope a high ratio of active tetrahedrally coordinated ions over Cr$^{3+}$ and Cr$^{4+}$ ions bound with a six-fold symmetry in the YAG host. Tetrahedrally coordinated Cr$^{4+}$ within a YAG host crystal has the symmetry of the D$_{2d}$ group. The lasing transition is broadened by coupling between the electronic and vibrational states within the crystal. Lasers with this type of broadening are sometimes referred to as vibronic lasers.

To demonstrate the broadband gain possible from Cr$^{4+}$:YAG, the photoluminescence (PL) spectrum from a 2 cm laser crystal was measured and is shown in Figure 2.1. The PL spectrum is measured with the crystal cooled to 13 C and pumped by 11 W from a Nd:YVO$_4$ laser running cw at 1064 nm. PL is collected by a 10 cm radius-of-curvature broadband mirror placed 5 cm from the crystal and subsequently focused into a multimode fiber. Light in the optical fiber is directed into an optical spectrum analyzer (OSA) which then measures the PL.
<table>
<thead>
<tr>
<th>Material</th>
<th>Cr$^{4+}$:YAG</th>
<th>Ti:Sapphire</th>
<th>Cr:Fosterite</th>
</tr>
</thead>
<tbody>
<tr>
<td>Gain Bandwidth (nm)</td>
<td>1300 - 1600</td>
<td>650 - 1050</td>
<td>1200 - 1400</td>
</tr>
<tr>
<td>Absorption Bandwidth (nm)</td>
<td>900 – 1150</td>
<td>400 – 600</td>
<td>850 - 1200</td>
</tr>
<tr>
<td>Upper-state lifetime (µs at 300K)</td>
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<td>3.2</td>
<td>3.6</td>
</tr>
<tr>
<td>Emission Cross section (cm$^2$)</td>
<td>4.5 x 10$^{-19}$ at 1420 nm</td>
<td>3.5 x 10$^{-19}$</td>
<td>1.4 x 10$^{-19}$</td>
</tr>
<tr>
<td>Saturation Fluence (J/cm$^2$)</td>
<td>0.4</td>
<td>0.7</td>
<td>0.8</td>
</tr>
<tr>
<td>Index of refraction</td>
<td>1.81</td>
<td>1.76</td>
<td>1.65</td>
</tr>
<tr>
<td>Nonlinear Index of refraction n$_2$ (cm$^2$/W)</td>
<td>6.9 x 10$^{-16}$</td>
<td>3 x 10$^{-16}$</td>
<td>~ 10$^{-16}$</td>
</tr>
<tr>
<td>Thermal Conductivity (W/K cm)</td>
<td>0.12</td>
<td>0.34</td>
<td>0.08</td>
</tr>
<tr>
<td>Shortest Pulses Generated (fs)</td>
<td>20</td>
<td>5</td>
<td>14</td>
</tr>
</tbody>
</table>

Table 2.1 Material properties of several gain media used to generate ultrashort laser pulses.

Cr$^{4+}$:YAG exhibits broadband emission from 1220 to 1650 nm, with a maximum at 1365 nm. Despite the fact that the PL peak wavelength is 1365 nm, most results from Cr$^{4+}$:YAG lasers have been at wavelengths greater than 1500 nm.

Despite its many attractive properties, several unfavorable properties make Cr$^{4+}$:YAG difficult to use within a modelocked laser. A large thermo-optic coefficient and low thermal conductivity create significant thermal lensing effects that modify the laser cavity mode. Thermal lensing effects are difficult to model and make modelocking highly sensitive to pump power. Both the laser emission and pump wavelengths are depleted by excited state absorption (ESA) transitions [32, 33]. The pump ESA leads to a reduced power conversion efficiency, while the emission ESA acts as an intracavity loss element. As a result, the net gain of Cr$^{4+}$:YAG is lower than that of many other ultrafast laser gain materials. Finally, the Cr$^{4+}$:YAG gain spectrum is centered in a spectral region plagued by water absorption lines. Water vapor in the air introduces both intracavity loss and uncontrollable spikes in the dispersion.
2.4 Four-Level System

The active electronic energy levels in a Cr\textsuperscript{4+}:YAG laser can be modeled as a four-level system [34]. A schematic diagram of the four level system is shown in Figure 2.2. Light amplification requires a population inversion between the lasing transition states 3 and 2. If a population inversion can be achieved, stimulated emission will be greater than stimulated absorption, and gain will be present.
Figure 2.2  Diagram of the four-level system model of a Cr$^{4+}$:YAG laser. A pump laser drives the transition from state 1 to state 4 at a rate $W_p$, but also results in excited state absorption (ESA) from state 3 and state 4 at rates $W_{ESA3}$ and $W_{ESA4}$. Electrons in states 2, 3, and 4 relax to lower energy levels with relaxation rates $\gamma_{ij}$. Stimulated transitions occur at the lasing transition at a rate $W_L$.

In order to create a population inversion, carriers are first excited to state 4 from the ground state (state 1) by a pump source. The stimulated pump absorption rate is given by [34]:

$$W_p = \frac{\sigma_p I_p}{h\nu_p}$$

(2.1)

where $\sigma_p$ is the pump absorption cross section, $I_p$ is the pump intensity, and $h\nu_p$ is the pump photon energy. In the ideal four level system, there will be no loss of carriers from state 4 through either excited state absorption or stimulated emission induced by the pump laser back to the ground state. The condition for this is $\gamma_{43} \gg W_p, W_{ESA4}$. Cr$^{4+}$:YAG has a significant pump excited state absorption component, which removes pump photons that would have otherwise been available to excite carriers from state 1 to state 4. It is also possible that some car-
riers excited by ESA will relax to the ground state by nonradiative mechanisms. The nonradiative losses of carriers from state 4 will reduce the quantum efficiency of the laser. In order to avoid pump absorption saturation, carriers should relax quickly to state 3. Furthermore, the relaxation to state 3 should be significantly quicker than relaxation to state 2 or state 1. In order for this to be possible, \( \gamma_{43} \gg \gamma_{42}, \gamma_{41} \).

State 3 should have a long enough lifetime such that a population inversion will develop between state 3 and state 2 (\( \gamma_{32} \ll \gamma_{21} \)). State 3 is often called metastable because of its long lifetime. Carriers that transition from state 3 to state 2 should relax quickly from state 2 to the ground state in order to maintain the population inversion between state 3 and state 2. Cr\(^{4+}\):YAG also has a significant excited state absorption at the laser wavelength. This ESA acts as a loss element for the lasing transition.

In materials that meet all of the conditions discussed above, it should be possible to create a population inversion and observe lasing. By solving rate equations and making all of the lifetime approximations discussed above, the population inversion is [34]:

\[
\frac{\Delta N_{32}}{N} = \frac{W_p \tau_{32}}{1 + W_p \tau_{32}} \tag{2.2}
\]

where \( \Delta N_{32} = N_3 - N_2, N_3 \) is the population of state 3, \( N_2 \) is the population of state 2, \( N \) is the total number of electrons, and \( \tau_{32} = 1/\gamma_{32} \) is the lifetime of state 3. The threshold condition is that gain in the laser should exceed loss. In this case, light will be amplified until the laser transition is saturated. The gain of a laser with a population inversion \( \Delta N \) is:

\[
g = \Delta N \sigma (\nu) \tag{2.3}
\]

where \( \sigma (\nu) \) is the frequency dependent emission cross section. From these relationships, it is possible to estimate the pump power necessary to achieve lasing. At threshold, gain equals cavity loss, which is primarily due to output coupling from the end mirror and scattering losses within the cavity. A Cr\(^{4+}\):YAG laser typically uses a 1\% output coupler, and has scattering loss on the order of 0.5\%. The threshold gain is calculated by the relationship [35]:

\[
g_{th} = -\frac{1}{2L} \ln (r_{OC}^s) \tag{2.4}
\]

where \( L \) is the gain length, \( r_{OC} \) is the output coupler reflectivity fraction, and \( s \) is the fractional loss due to scattering. For a 2 cm crystal, \( g_{th} = 3.8 \times 10^{-3} \text{ cm}^{-1} \). The threshold pumping power can then be computed by the relationship [35]:

ULTRAFAST LASER DESIGN

43
\[
\left( \frac{I}{V} \right)_{\text{threshold}} \approx \frac{\hbar \nu_{41} \Delta N_{\text{th}}}{\tau_{32}}
\]  

(2.5)

where \( I \) is the pump intensity and \( V \) is the volume of the pumped region. Assuming an average mode beam waist of 200 \( \mu \text{m} \), the threshold absorbed pump power is 1.1 \( \text{W} \). As will be shown later in Chapter 3, this is quite close to the experimentally measured threshold pump power.

### 2.5 Modelocking

Without proper cavity design and alignment, lasers with broadband gain will not emit radiation in short pulses. Short optical pulses will only be generated from a laser cavity if cavity conditions are set such that short light pulses experience higher gain than continuous wave (cw) light. A general technique to create this condition is known as modelocking. Modelocking refers to a class of methods which generate short optical pulses by phase-locking periodically spaced frequency domain laser cavity modes. Laser cavity modes are periodically spaced because of longitudinal Fabry-Perot etalon effects. Phase-locking of longitudinal frequency modes is provided by with the introduction of a periodically time-dependent loss modulation in the time domain. A time dependent loss within a laser cavity can be created by either an intracavity optical modulator (active modelocking) or by nonlinear interactions between an optical pulse and the gain media or other materials within the cavity (passive modelocking). The final pulsewidth generated will be determined by the interplay of modelocking pulse shortening mechanisms and intracavity pulse broadening mechanisms.

Modelocking requires the existence of a set of equally spaced cavity modes in the frequency domain. The optical modes inside laser cavities are spaced periodically in frequency with a mode spacing of:

\[
\Delta \nu_c = \frac{c}{2n(\nu)L}
\]

(2.6)

where \( c \) is the speed of light, \( L \) is the length of the optical cavity, and \( n(\nu) \) is a frequency dependent effective group velocity index of refraction of the optical cavity. The frequency dependence of the refractive index is from dispersive intracavity elements such as the gain medium, prisms, and mirrors. Because dispersion plays a major part in the pulse dynamics of an ultrafast laser, control of dispersion is critical to generate short pulses. It is clear, for
example, that a large, nonlinear frequency dependence of the index of refraction will spoil the longitudinal mode periodicity.

At any point in time, the oscillating electric field inside a laser cavity will be a superposition of modes:

\[ E(t) = \sum_{m} a_m e^{i(\Delta \omega_m t + \phi_m)} \]  

(2.7)

where \( m \) is the mode number, \( a_m \) is the amplitude of the electric field in mode \( m \), \( \Delta \omega_c = 2\pi \Delta v_c \), and \( \phi_m \) is the phase of mode \( m \). Because the frequency space electric field distribution is periodic with frequency \( \Delta v_c \), the Fourier transform related time domain electric field waveform will repeat every \( \tau_c = 1/\Delta v_c \). The characteristic time \( \tau_c \), known as the cavity roundtrip time, is the length of time for an optical pulse to travel around the laser cavity. The inverse of the cavity roundtrip time, \( \Delta v_c \), is known as the cavity repetition rate.

In free-running operation, the intracavity field energy will be concentrated within a few modes near the laser cavity gain peak. These modes will generally be excited with no relationship between their phases. Such a laser is said to be running in a continuous wave (cw) mode. For lasers with many incoherently excited modes, random noise spikes of a duration corresponding to the increased spectral bandwidth become apparent in the time domain. A plot of intensity envelopes generated by 5 and 25 optical modes with random phases is shown in Figure 2.3. The calculated time-dependent intensity profiles hold the integrated intensity (output energy) constant, while allowing the instantaneous intensity to vary depending on whether the modes constructively or destructively interfere. The time-dependent amplitude noise spikes have a characteristic time scale of \( 1/\Delta v_{\text{tot}} \), where \( \Delta v_{\text{tot}} = N \Delta v_c \), is the total bandwidth of all of the modes, and \( N \) is the number of modes. As long as the relative phases between the modes stay the same, the resulting waveform will repeat with a frequency of \( \Delta v_c \), the cavity repetition rate.
Figure 2.3  Intensity waveforms calculated for 5 and 25 cavity modes with random phases.

Modelocked lasers generate a periodic pulsetrain of short pulses by locking together the phases of each cavity mode. A plot of the laser intensity envelopes from 5 and 25 locked optical modes is shown in Figure 2.4. The waveform repeats with a frequency $\Delta v_c$ (the cavity

Figure 2.4  Intensity waveform from a laser consisting of 5 and 25 modes with locked phases. As the number of modes increases, the pulse width decreases and the instantaneous intensity increases.
repetition rate. The pulsewidth will depend on the number of modes phaselocked together (N), the spacing between those modes (Δνc), and the shape of the modal amplitude envelope E(ω), but will generally be on the order of magnitude of 1/Δνtot. In a 20 fs Cr4+:YAG laser there are approximately 3 x 10^5 modes within the full-width half-maximum bandwidth.

The total width of the electric field spectrum in frequency space is called the bandwidth of the laser pulse. The pulsewidth and bandwidth of a pulse are related by the Heisenberg uncertainty principle, which constrains the product of the bandwidth (ΔE = hΔν) and the pulsewidth (Δτ) to be greater than some number:

\[ ΔEΔτ ≥ \frac{h}{4π} \]  

(2.8)

where h is Plank’s constant. In the Heisenberg uncertainty principle, ΔE and Δτ are the 1/e widths. In laser physics, it is more convenient to consider the time-bandwidth product (tbwp), which relates the full width at half maximum of the pulsewidth and the bandwidth to be:

\[ ΔνΔτ ≥ tbwp \]  

(2.9)

where tbwp is pulse shape dependent. For Gaussian shaped pulses, tbwp = 0.4413, while for sech pulses tbwp = 0.3148 [36].

The ultimate pulsewidth generated by an ultrafast laser is determined by the interplay between several intracavity pulsesshaping effects. Modelocking, for example, is used to create strong pulse shortening effects. In actively modelocked lasers, the pulse shortening is dominated by a time dependent loss from an active modulator. Passively modelocked lasers, on the other hand, achieve modelocking from self amplitude modulation. Pulses in passively modelocked lasers often become short enough that the pulse shaping becomes dominated by soliton or dispersion managed soliton dynamics. Pulses shortening by modelocking will eventually be limited by some pulse broadening effect. Pulses broadening effects include intracavity dispersion, a limited gain bandwidth, and self phase modulation.

2.6 Gain Filtering

Short pulses will have a correspondingly broad spectrum determined by the time-bandwidth product. All other factors held constant, lasers with broadband gain will support shorter pulses than lasers with narrow gain. For most pulse shapes, gain acts as a spectral filter that broadens pulses during amplification. The shape of the gain bandwidth depends on the principle gain broadening mechanism. Purely homogeneously broadened gain media will
have a Lorentzian lineshape, whereas purely inhomogeneously broadened gain medial will have a Gaussian lineshape. In either case, the gain can be approximated near the center of the gain curve by replacing the true lineshape with an inverted parabola. In this parabolic approximation, the wavelength dependent gain is given by:

$$g(\omega) = g\left(1 - \frac{(\omega - \omega_0)^2}{\Omega_g^2}\right)$$

(2.10)

where $g$ is the maximum gain, $\omega$ is the frequency, $\omega_0$ is the center gain frequency, and $\Omega_g$ is the gain bandwidth. After propagating through the gain media, the pulse shape is given by:

$$E'(\omega) = (1 + g(\omega))E(\omega)$$

(2.11)

where $E'(\omega)$ is the pulseshape after the gain media and $E(\omega)$ is the pulse shape before the gain media. By taking the Fourier transform of equation (2.11), $g(\omega)$ can be converted into the time domain operator:

$$E'(t) = 1 + g\left(1 + \frac{1}{\Omega_g^2}\frac{d^2}{dt^2}\right)E(t)$$

(2.12)

For a general pulse shape and gain spectrum, the laser amplifier gain can either broaden or shorten pulses. For example, gain spectrally offset from the pulse spectrum, could amplify the wings of the pulse spectrum to yield an overall larger pulse spectrum. Within a laser where the pulse spectrum is centered around the gain bandwidth, the gain profile will act as a filter and broaden the pulse. The gain filtering becomes particularly strong as the pulse spectrum grows to be on the order of the gain bandwidth. Gain filtering in combination with a Raman redshift could also limit pulsewidths [37]. Despite this gain filtering, it is possible to generate pulse spectrum broader than the gain bandwidth through spectral broadening effects such as self phase modulation.

### 2.7 Group Delay Dispersion

Dispersion is a critical pulse shaping mechanism within many ultrafast laser cavities. A dispersive media is any material that has a frequency dependent dielectric constant, $\varepsilon(\omega)$. The phase velocity of monochromatic light through a dispersive media is therefore $v_{ph} = c/n(\omega)$, where $n(\omega) = \varepsilon(\omega)^{1/2}$. The effects of dispersion on the propagation of an optical pulse are more complicated. A pulse envelope traveling through a dispersive media will usually
propagate with a speed equal to the group velocity \( v_g = \frac{d\omega}{dk} \), and will be shaped according to the group velocity dispersion (GDD).

Light propagating in the z-direction through a dispersive media will have an electric field that propagates according to the wave equation:

\[
\frac{\partial^2 E(z, t)}{\partial z^2} - \frac{n^2(\omega)}{c^2} \frac{\partial^2 E(z, t)}{\partial t^2} = 0
\]  

(2.13)

Assuming solutions to this equation of the form:

\[
E(z, t) = A(z, t) \exp[-i\omega t + i\beta(\omega)z]
\]  

(2.14)

where \( A(z, t) \) is called the pulse envelope. The propagation constant \( \beta(\omega) \) can be expanded around the carrier frequency \( \omega_0 \):

\[
\beta(\omega) = \sum_{n=0}^{\infty} \frac{1}{n!} \frac{d^n}{d\omega^n} \beta(\omega) \bigg|_{\omega_0} (\omega - \omega_0)^n = \sum_{n=0}^{\infty} \frac{1}{n!} \beta_n(\omega)(\omega - \omega_0)^n
\]  

(2.15)

It is possible to reduce Maxwell's wave equation to a nonlinear Schrödinger equation using equations (2.14) and (2.15). In order to derive this equation, \( A(z, t) \) is assumed to be slowly varying compared with the exponential term. \( A(z, \omega) \), the Fourier transform of \( A(z, t) \), can be written, keeping terms to third order in dispersion, as:

\[
A(z, \omega) = A(0, \omega) \exp[i(\beta_0 + \beta_1(\omega - \omega_0) + \beta_2(\omega - \omega_0)^2 + \beta_3(\omega - \omega_0)^3)z]
\]  

(2.16)

To look at how the pulse frequency \( A(z, \omega) \) evolves as it propagates, we take a derivative with respect to \( z \), which in the slowly varying envelope approximation is:

\[
\frac{\partial}{\partial z} A(z, \omega) = i[\beta_0 + \beta_1(\omega - \omega_0) + \beta_2(\omega - \omega_0)^2 + \beta_3(\omega - \omega_0)^3]A(z, \omega)
\]  

(2.17)

Taking the Fourier transform, we arrive at the nonlinear Schrödinger equation for the evolution of the pulse envelope with dispersion given by:

\[
\left[ \frac{\partial}{\partial z} + i\beta_0 + \beta_1 \frac{\partial}{\partial t} + \frac{i}{2} \beta_2 \frac{\partial^2}{\partial t^2} - \frac{1}{6} \beta_3 \frac{\partial^3}{\partial t^3} \right] A(z, t) = 0
\]  

(2.18)
In the absence of other nonlinear effects, a pulse with an envelope shape of \( A(z,t) \) will be shaped by dispersion according to equation (2.18). The term \( \beta_0 \) results in a phase-shift, and is related to the phase velocity by \( v_p = \omega / \beta_0 \).

First-order dispersion results in a group velocity that differs from the speed of light. For \( \beta_2 = \beta_3 = 0 \), solutions to equation (2.18) will be of the form \( A(z,t) = A(z + \delta z, t + \delta t) \) for where \( \delta z / \delta t = 1 / \beta_1 \). The group velocity is therefore defined as \( v_g = 1 / \beta_1 \). Equation (2.18) can be rewritten in the reference frame of the optical pulse, \( z = v_g t \) to eliminate the first order dispersion term. Simply put, first order dispersion creates a time delay for a pulse propagating through a media.

The effects of higher-order dispersion on a pulse shape can be determined by solving the nonlinear Schrodinger equation including the higher-order dispersion terms. In mode-locked lasers, the pulse envelope \( A(z,t) \) is often a Gaussian or sech function. For Gaussian envelope with a pulsewidth of \( \tau_0 \), the pulsewidth as a function of propagation distance through a dispersive media with second order dispersion can be calculated exactly to be [38]:

\[
t(z) = \tau_0 \sqrt{1 + \left( \frac{z}{L_{D2}} \right)^2}
\]

(2.19)

where \( z \) is the propagation distance and \( L_{D2} = \tau_0^2 / |\beta_2| \) is the dispersive length for second order dispersion. For Gaussian pulses, the full width at half maximum pulsewidth is related to the Gaussian 1/e width by:

\[
\tau_{fwhm} = 1.665 \tau_0
\]

(2.20)

During propagation through the dispersive media, the pulse will broaden and simultaneously develop a chirp (a time-dependent frequency). For a \( \beta_2 > 0 \), the long wavelengths of the pulse will be in the leading edge and the short wavelengths in the trailing edge of the pulse. Because the low frequency component of the pulse spectrum travels faster than the high frequency component, positive \( \beta_2 \) is called normal second-order dispersion. The regime where \( \beta_2 < 0 \) is known as anomalous second-order dispersion and adds chirp such that the pulse has a high frequency leading edge and a low frequency trailing edge. The parameter \( \beta_2 \) is called the group velocity dispersion (GVD). For propagation through a dispersive material of thickness \( d \), the parameter \( d\beta_2 \) is known as group delay dispersion (GDD).
2.8 Material Dispersion

Dispersion within a dielectric media is a result of the interaction between light and matter. In a vacuum, there is no wavelength dependence of the index of refraction. While propagating through a media, however, light will interact strongly with electrons at the material resonance wavelengths. A simple classical model of the interaction between light and a material, called the Lorentz model, is able to predict the form of the wavelength-dependent index of refraction for many materials [39]. This model considers the motion of electrons in a material to be governed by a forced damped oscillator equation. The electrons are forced by the oscillating photon electric field and damped by electron-lattice interactions. The electronic equation of motion is:

\[
\frac{d^2}{dt^2}x(t) + \gamma \frac{dx(t)}{dt} + \omega_0^2 x(t) = \frac{q}{m} E_0 e^{-i\omega t} \tag{2.21}
\]

where \(x(t)\) is the electron position, \(\gamma\) is the damping constant, \(\omega_0\) is the natural resonance frequency of the electron in its potential, \(q\) is the electron charge, \(m\) is the electron mass, \(E_0\) is the magnitude of the electric field, and \(\omega\) is the photon drive frequency. From this equation, it is possible to solve for the frequency-dependent polarization of a more realistic material with several resonance frequencies. For a material with an oscillator strength \(f_j\) for each transition of resonance frequency \(\omega_0j\):

\[
P(\omega) = Nq \sum_j f_j x_j(t) = \frac{q^2 N}{m} \left( \sum_j \frac{f_j}{\omega_0^2 - \omega^2 - i\gamma_j \omega} \right) E_0 e^{-i\omega t} \tag{2.22}
\]

The frequency-dependent index of refraction can then be calculated from the real part of the propagation constant. Far away from resonance, damping terms can often be neglected. In transparent materials such as glasses or crystals, for example, the damping coefficient \(\gamma_j\) can be neglected for visible or infrared wavelengths because the resonant frequencies are often in the ultraviolet. In this approximation the frequency dependent index of refraction can be calculated to be:

\[
n(\omega) = \frac{c}{\omega} k_{Re} = \left[ 1 + \frac{q^2 N}{m \varepsilon_0} \left( \sum_j \frac{f_j}{\omega_0^2 - \omega^2} \right) \right]^{-\frac{1}{2}} \tag{2.23}
\]

The index of refraction can be written instead as a function of wavelength as:
\[ n(\lambda) = \sqrt{1 + \sum_j \frac{A_j \lambda^2}{\lambda^2 - B_j}} \]  

(2.24)

This form of the index of refraction is known as the Sellmeier equation; the Sellmeier coefficients \( A_j \) and \( B_j \) are tabulated for many materials [40].

As an example, it is worth considering the dispersion of YAG. In Cr\(^{4+}\):YAG, the index of refraction will be dominated by the contribution from the host crystal YAG. The index of refraction is:

\[ n(\lambda) = \sqrt{1 + \frac{2.293\lambda^2}{\lambda^2 - (0.1095)^2} + \frac{3.705\lambda^2}{\lambda^2 - (17.825)^2}} \]  

(2.25)

where \( \lambda \) is in microns. From the index of refraction, the GVD of YAG can be calculated by the formula:

\[ \text{GVD}_{\text{YAG}} = \frac{\lambda^3}{2\pi c^2} \frac{d^2}{d\lambda^2} n(\lambda) \]  

(2.26)

A plot of the YAG index of refraction and GVD is shown in Figure 2.5. The GVD is normal for wavelengths below 1.6 \( \mu \text{m} \) and anomalous above. It is noteworthy that the zeroth -order chromatic dispersion is normal in both wavelength ranges.
2.9   **Kerr Effect and Self Phase Modulation**

It has been shown that a frequency-dependent index of refraction, $n(\omega)$, introduces group velocity dispersion and contributes to pulse shaping. Materials also have a nonlinear intensity dependent index of refraction $n(I)$ which contributes to pulse shaping at high intensities. In particular, the optical Kerr effect is an intensity dependent index of refraction of the form:

$$n(I) = n_0 + n_2 I$$  \hspace{1cm} (2.27)

where $n_0$ is the low intensity index and $n_2$ is the nonlinear Kerr index. Pulses of light propagating through Kerr media may undergo an effect called self-focusing. An optical pulse will typically have a pulse intensity envelope which is a function of position and time. For example, an optical pulse might have a spatially-dependent Gaussian intensity profile:
\[
I(x) = I_0 \exp \left[ -\frac{x^2}{x_0^2} \right]
\]  

(2.28)

The pulse itself would then create a spatially-dependent index of refraction given by:

\[
n(x) = n_0 + n_2 I_0 \exp \left[ -\frac{x^2}{x_0^2} \right]
\]  

(2.29)

A Gaussian index of refraction profile acts as a lens. A pulse can therefore create an intensity-dependent lens and focus itself. Self-focusing, in combination with a spatially-dependent gain or loss, can be used to couple the time and position profiles of an optical pulse, and forms the basis of Kerr lens modelocking.

In addition to the spatially dependent intensity, optical pulses by definition have a time dependent instantaneous intensity. In the absence of dispersion or other pulse shaping mechanisms, a pulse with a time dependent electric field envelope \(A(z,t)\) will obey the nonlinear equation:

\[
\frac{\partial A}{\partial z} = i \delta |A|^2 A
\]  

(2.30)

where \(\delta\) is called the nonlinear coefficient. The nonlinear coefficient is related to \(n_2\) by:

\[
\delta = \frac{n_2 \omega}{c A_{\text{eff}}}
\]  

(2.31)

where \(A_{\text{eff}}\) is the effective modal area. For a Gaussian beam shape with beamwaist \(r_0\), \(A_{\text{eff}} = \pi r_0^2\). Solutions to this equation describe an effect called self phase modulation (SPM). Because the phase velocity of light at a given frequency is proportional to the index of refraction, an optical pulse will develop a frequency-dependent phase shift which in turn is a function of intensity. The spatially-dependent intensity will translate into a spatially-dependent phase shift. An unchirped Gaussian pulse, for example, will develop a chirp after transmission through a Kerr media of thickness \(z\) [38]:

\[
\delta \omega(t,z) = \frac{2z t}{\tau_0^2 L_{\text{NL}}} \exp \left[ -\frac{t^2}{\tau_0^2} \right]
\]  

(2.32)

where \(\delta \omega(t,z)\) is the frequency shift as a function of time (\(t = 0\) is the center of the pulse) and \(L_{\text{NL}}\) is the nonlinear length given by \(L_{\text{NL}} = 1/\delta |A|^2_{\text{max}}\). This SPM chirp can be explained
graphically, as shown in Figure 2.6. The high intensity pulse peak will see a higher index of

refraction than the lower intensity pulse wings. As a result, the pulse will develop an overall chirp. The high intensity pulse center will lag relative to the low intensity pulse wings, causing the field oscillations in the pulse trailing edge to bunch up to higher frequencies and the field oscillations in the pulse leading edge to stretch to lower frequencies. The chirp has the same sign as normal dispersion where higher frequencies will be delayed relative to lower frequencies. By choosing a media with SPM and anomalous GVD, it is possible for pulses to propagate without pulse distortion. These pulses are known as solitons.

In addition to developing a chirp, SPM will actually create new frequencies in the pulse spectrum. The maximum spectral shift from the carrier frequency \( \omega_0 \) of a Gaussian pulse with initial frequency bandwidth of \( \Delta \omega \) after propagation through a distance \( z \) of a Kerr media is given by [38]:

\[
\delta \omega_{\text{max}}(z) = (0.86 \delta |A^2| \Delta \omega) z
\]

(2.33)

SPM induced spectral broadening is crucial for the generation of ultrashort pulses. In ultrafast Ti:Sapphire lasers, it has been possible to generate pulse spectrum broader than the gain bandwidth through SPM and SAM [20].
2.10 **Pulse Shortening Rate**

Several approaches have been used to modelock lasers and generate short pulse generation. Each type of modelocking can be viewed in the time domain as adding a periodic time dependent loss. To compare the relative strength of each modelocking technique in different pulsewidth regimes, it is useful to calculate a pulse shortening rate (PSR) [41]:

\[
\text{PSR} = \frac{\Delta \tau}{\tau} \tag{2.34}
\]

where $\Delta \tau$ is the change in pulsewidth per pass through the loss modulation and $\tau$ is the pulsewidth. The PSR represents the fractional change in pulsewidth per pass through the cavity modulation element. A large PSR is an indication of strong modelocking.

To calculate the PSR for a given modelocking technique, the change in pulsewidth through the modulator is calculated. The transmission through a periodic modulating element is given by:

\[
T = 1 - M(t) \tag{2.35}
\]

where $M(t)$ is a periodic loss introduced by the cavity modulator. $M(t)$ can in many cases be approximated by a parabola:

\[
M(t) \approx M(\tau)t^2 \tag{2.36}
\]

where the coefficient $M(\tau)$ is given by the particular modulation function and may depend on the pulsewidth. To first order, the pulse shape can be assumed to be an inverted parabola:

\[
E(t) \approx E_0 \left(1 - \frac{t^2}{\tau^2}\right) \tag{2.37}
\]

where $\tau$ is the initial pulsewidth. After transmission through the modulation element, the new pulse shape is given by:

\[
E'(t) \approx E_0 \left(1 - \frac{t^2}{\tau'^2}\right) = T \cdot E(t) \tag{2.38}
\]

where $\tau'$ is the new pulsewidth after transmission through the modulator. Keeping terms of order $t^2$, and ignoring the linear loss 1, it is possible to get the relation:
\frac{1}{\tau'^2} = \frac{1}{\tau^2} + M(\tau) \quad (2.39)

Solving for the new pulsewidth gives:

\tau' = \frac{1}{\tau(1 + M(\tau)\tau^2)^2} \approx \tau \left( 1 - \frac{M(\tau)\tau^2}{2} \right) \quad (2.40)

where the approximation that the modulation loss is small compared to the total loss has been made. The PSR can now be calculated to be:

\text{PSR} = \frac{\tau - \tau'}{\tau} = \frac{M(\tau)\tau^2}{2} \quad (2.41)

Active and passive modelocking are two general approaches used to couple the cavity mode phases and generate short optical pulses. Active modelocking utilizes an optical modulator placed within the optical cavity. This modulator either introduces a time-dependent loss (amplitude modulation) or a time-dependent phase (phase modulation). Passive modelocking uses an optical pulse to modulate its own amplitude or phase. In order to have a pulse modulate itself, saturable absorption or intensity-dependent loss, must be introduced to the cavity. This saturable absorption can be broadly classified as either slow or fast saturable absorption. In slow saturable absorption, intensity-dependent loss modulates the leading edge of the optical pulse, and gain saturation modulates the trailing edge of the optical pulse. In fast saturable absorption, however, the intensity-dependent loss is fast enough to modulate both the leading and trailing edges of the pulse.

2.11 Active Modelocking

Modelocked lasers can be modeled by considering the effect of all pulse shaping mechanisms encountered by an optical pulse in a laser cavity. The modelocking itself is performed by some time-dependent loss element. The pulse will be subsequently shaped by gain, loss, GVD, and self phase modulation. Actively modelocked laser cavities, in particular, contain a gain element, linear loss, and an optical modulator. A schematic of such a laser cavity is shown in Figure 2.7. Group velocity dispersion has been ignored in this model, because most actively modelocked lasers are limited in minimum pulsewidth by bandwidth limitation before GVD effects become significant.
In the time domain picture, an optical modulator creates a time-dependent loss, as shown schematically in Figure 2.8. Laser emission will be generated at times when the net loss from the modulator and linear sources is lower than the net cavity gain. When the loss
from the modulator is higher than the intracavity gain, laser emission will be suppressed. Each cavity roundtrip, the modulator will trim off the temporal edges of the light pulse. For pulsewidths larger than the modulation window, the modulator will strongly shape the pulse. Eventually the pulsewidth will be shorter than the modulation window, and loss modulation from the active modulator will not significantly decrease the pulsewidth. For this reason, active modelocking easily initiates the creation of pulses ($\tau_{\text{pulse}} \gg \tau_{\text{modulator}}$) but is not as effective in shortening pulses which are already shorter than the modulator window ($\tau_{\text{pulse}} \ll \tau_{\text{modulator}}$).

In frequency space, the optical modulator has the effect of generating sidebands on each carrier frequency separately. Intracavity photons can either gain or lose energy at the modulator frequency. When the frequency of the modulator drive corresponds to the cavity repetition rate, the modulator will scatter photons from a cavity mode into its nearest neighboring cavity modes. A schematic of this process is shown in Figure 2.9. This modulation

\[
\begin{align*}
\ldots & \quad V_{n-2} \quad V_{n-1} \quad V_n \quad V_{n+1} \quad V_{n+2} \quad \ldots \\
\end{align*}
\]

Figure 2.9  Schematic of active modelocking in frequency space. Lines correspond with cavity modes at $V_n$. Dashed arrows represent coupling from each mode to its neighbors by a sideband generating optical modulator. This process will couple the mode's phases together.

has the effect of coupling, or 'locking', the phases of adjacent modes. The entire chain of modes will become completely phaselocked as energy is fed from each optical mode to its neighbors throughout the pulse spectrum. It is noteworthy that modelocking feeds energy away from the spectral peak of the gain curve. This outward flow of energy helps generate a pulse spectrum broader than the gain bandwidth in some modelocked lasers.

It is possible to increase the repetition rate of an actively modelocked laser by driving the modulator at integer multiples of the fundamental cavity repetition rate [42]. An optical
cavity mode will couple to modes some multiple of the cavity repetition rate away, rather than the adjacent mode. In this case, called harmonic modelocking, several independent chains of locked cavity modes can exist. In the time domain, it is clear that an active modulation at a frequency some multiple of the cavity repetition rate will result in multiple pulsing. For example, if the modulator is driven at twice the cavity repetition rate, two pulses can be present in the cavity at once.

An analytical solution for the resulting pulse shape and width can be solved by using the slowly varying envelope approximation. This approximation assumes that the pulse propagation is completely determined by the pulse energy envelope $A(t)$ without regard to the phase of the underlying field $E(t) = A(t)e^{i\omega t}$. This approximation holds for pulses as short as two optical cycles. A Master equation can be written and solved which describes the propagation of a pulse through an actively modelocked laser cavity. The equation is solved by requiring the self-consistent condition that there should be no change to the optical pulse after travelling through the cavity. The Master equation for actively modlocked lasers is then [41]:

$$\left[ g \left( 1 + \frac{1}{2} \frac{d^2}{\Omega_g^2 dt^2} \right) - l - M(1 - \cos \Omega_m t) \right] A(t) = 0 \quad (2.42)$$

where $g$ is the small signal gain, $l$ is the loss, $M$ is the modulator strength, and $\Omega_m$ is the modulator frequency. If the gain and modulation is approximated by a parabola centered at the peak of the gain, this equation has an exact solution, given by the Gaussian:

$$A(t) = A_0 \exp \left[ -\frac{t^2}{\tau^2} \right] \quad (2.43)$$

where $A_0$ is the amplitude and $\tau$ is the pulsewidth, given by:

$$\tau = \left( \frac{8g}{M\Omega_g^2 \Omega_m^2} \right)^{\frac{1}{4}} \quad (2.44)$$

The resulting pulsewidth will be smaller for large gain bandwidths, high modulation frequencies (small temporal modulation windows), and large modulation strengths. Larger gain leads to longer pulsewidths because of the increased strength of the gain filter.

It is useful to determine what the rate of pulse compression is for an optical pulse of a given pulsewidth for active modelocking. The PSR is calculated using the modulation function in the Master equation and equation (2.41), and is given by:
\[ \text{PSR}_{\text{active}} = \frac{M(\tau)\tau^2}{2} = \frac{M\Omega_m\tau^2}{4} \]  

(2.45)

As can be seen by the fact that the PSR \( \sim \tau^2 \), active modelocking will rapidly compress the pulsewidth of long pulses but will fizzle in strength when compressing already short pulses. This is because the modulation width of the time-dependent loss is a constant and does not decrease with the pulsewidth. Passively modelocked pulses, in contrast, create their own modulation window and are much more effective in shortening the pulsewidth of already short pulses.

### 2.12 Passive Modelocking

In passively modelocked lasers, the optical pulse in a laser cavity is used to create a time-dependent loss. This time-dependent loss will then shorten the pulse through a process called self amplitude modulation (SAM). SAM is often realized by converting a time-dependent intensity pulse shape into a time-dependent loss through an intensity-dependent loss. Short pulses containing the same integrated energy as longer pulses will have much higher peak intensities, which are then converted into stronger time-dependent losses. Physical elements with an intensity-dependent loss are known as saturable absorbers. Saturable absorbers add intensity-dependent loss by absorbing cw light at some linear level while absorbing higher intensity pulses at a much lower saturated level. Saturable absorbers can either be 'real' (actual elements with saturable absorbance) or consist of a combination of optical elements which together act as an 'effective' saturable absorber.

Once a saturable absorption is saturated, the absorption will relax back to the linear level with some time constant [43]. Some saturable absorbers respond slowly compared to the pulsewidth and therefore only shape the leading edge of the pulse. These slow saturable absorbers require carefully balanced gain saturation dynamics to shape the trailing edge of the pulse. Fast saturable absorbers, however, respond to the pulse intensity quickly enough to modulate both the leading and trailing edges of the pulse. In either case, pulses short enough to generate additional spectrum and nonlinear chirp through self-phase modulation (SPM) can be created. This SPM, when balanced by dispersion, allows the generation of even shorter pulses through soliton-like pulse shaping effects.
A schematic of the modelocking process in a slow saturable absorber laser is shown in Figure 2.10. As in the case of active modelocking, the pulse is created by a time-dependent loss. Continuous wave light will see a constant level of loss \( l \). A short pulse, however, will saturate (decrease) the net loss through absorption saturation effects. Therefore, short pulses are favored over cw generation. For slow saturable absorbers (\( \tau_{\text{pulse}} \ll \tau_{\text{recovery}} \)), the amount of saturable absorption induced by a given pulse will depend on the pulse energy rather than the instantaneous intensity. The time-dependent loss for such a laser can be approximated by [41]:

\[
l(t) = l_1 \exp\left(-\sigma_a \int_{-\infty}^{t} |E(t)|^2 \, dt\right)
\]

where \( l(t) \) is the time dependent loss, \( \sigma_a \) is the absorption cross section, and \( E(t) \) is the instantaneous electric field. In such a laser system, the saturable absorber will recover on a timescale faster than the cavity repetition rate but not fast enough to modulate the trailing edge of the pulse. In the absence of other pulse shaping mechanisms, the ultimate minimum pulse-width will be limited by the saturable absorber recovery time. ‘Real’ saturable absorbers often fall into this category of slow saturable absorbers.
Much shorter pulses can be generated from slow saturable absorber lasers if the intracavity pulse’s trailing edge is modulated by gain saturation. A pulse encountering the gain element will deplete the gain. The pulse’s trailing edge will be suppressed by the saturated gain. The saturable gain is of the same form as the saturable loss:

$$g(t) = g_i \exp \left(-\sigma_g \int_{-\infty}^{t} |E(t)|^2 dt \right)$$  \hspace{1cm} (2.47)$$

where \(g(t)\) is the time dependent gain, and \(\sigma_g\) is the gain saturation cross section. In order for slow saturable absorber modelocking to be effective, \(\sigma_g < \sigma_a\). Otherwise, the gain will saturate before a pulse has a chance to develop within the laser cavity. Gain saturation is effective for dye lasers, but is not typically effective for solid-state gain media with large upper-state lifetimes.

The pulse shortening rate for slow saturable absorber modelocking can be determined by finding the modulation window’s dependence on the pulse width. The leading edge of the modulation window is determined by the onset of the optical pulse, while the trailing edge is determined by the end of the optical pulse. The modulation window has a width of \(\tau\), and therefore \(M(\tau) \sim 1/\tau^2\). The PSR is then given by:

$$\text{PSR}_{\text{slow}} = \frac{M(\tau)\tau^2}{2} \propto \text{const}$$  \hspace{1cm} (2.48)$$

Because the PSR for slow saturable absorber modelocking is roughly constant as a function of pulsewidth, this modelocking technique should be effective for generating very short optical pulses.

Fast saturable absorber modelocking uses optical pulses to shape both their own leading and trailing pulse edges. ‘Real’ saturable absorptive processes have recovery times which are often on the time scale of picoseconds and therefore cannot be used to create a fast saturable absorber for femtosecond pulses. It is possible to create an ‘effective’ fast saturable absorber modelocked systems by exploiting fast reactive nonlinear processes in materials. These reactive processes can be converted, through proper cavity design, into fast SAM. The Kerr effect, an intensity-dependent index of refraction, is commonly used as a fast reactive process. The Kerr effect is believed to have a recovery time-scale of \(~1\) fs [44].

Two common realizations of fast saturable absorber modelocking are additive pulse modelocking (APM) and Kerr lens modelocking (KLM) [45, 46]. APM converts a pulse’s
time dependent intensity profile into SAM by creating a nonlinear Mach-Zender interferometer. The intracavity pulse is split into two and recombined after developing nonlinear phase shifts through SPM. The interferometer arm lengths are chosen to constructively combine the pulse center while destructively interfering the pulse wings, resulting in a shorter end pulse. KLM, on the other hand, converts a spatially-dependent intensity profile into SAM with a spatial aperture. KLM is discussed in greater detail in the Section 2.14.

A schematic of the time domain modelocking process in a fast saturable absorber laser is shown in Figure 2.11. Each optical pulse saturates the loss proportionally to the instantaneous intensity of the pulse envelope. The loss modulation window shrinks and the loss modulation amplitude increases as the pulse gets shorter and the instantaneous pulse intensity increases. It is very difficult to get a pulse started in such a cavity, because the loss modulation is weak for small intensity perturbations. This is known as the self-starting problem [47, 48]. Typically, a cavity mirror or prism must be shaken to start modelocking in such a system. Many systems use active modelocking or real saturable absorbers to initiate pulse generation [49 - 51] and then fast saturable absorber modelocking to decrease the pulsewidth further.

Figure 2.11 Time domain picture of pulses generated by fast saturable absorber passive modelocking. Pulses (A(t)) are generated where loss (l) is smaller than gain (g). The loss is modulated by the intracavity pulses themselves.
Alternatively, self-starting lasers can be made from a ring configuration [52] or a well aligned symmetric cavity [53].

It is interesting to consider the frequency domain picture of fast saturable absorber modelocking. Because the modulation window narrows with decreasing pulsewidth, the Fourier decomposition of the modulation function becomes increasingly broad and contains many high harmonics of the cavity repetition rate. These higher harmonics in the modulation function will not only couple each mode to their nearest neighbor but will also couple each mode to an increasing number of distant modes within the frequency comb. Whereas an actively modelocked laser operates by phaselocking modes with their nearest neighbors, fast saturable absorber passively modelocked modes all become increasingly coupled in phase as the pulse width decreases.

The saturable loss in a fast saturable absorber modelocked laser will follow the instantaneous intensity of the light pulse and is given by:

\[ l(t) = l_0 - \gamma |A(t)|^2 \]  \hspace{1cm} (2.49)

where \( l(t) \) is the time dependent loss, \( l_0 \) is the linear loss, \( \gamma \) is the SAM coefficient, and \( A(t) \) is the envelope of the pulse electric field. A Master equation can be written governing fast saturable absorber modelocking. A schematic of all the effects encountered by a pulse in a roundtrip in the laser cavity is shown in Figure 2.12 [46]. Because the pulses generated can be very short, it is necessary to include the effects of GVD, SAM, and SPM. The Master equation assumes the slowly varying envelope approximation and is given by [45]:

\[ \left[ g \left( 1 + \frac{1}{\Omega_g^2 dt^2} \right) - \frac{d^2}{dt^2} + i\Psi + iD \frac{d^2}{dt^2} + (\gamma - i\delta) |A(t)|^2 \right] A(t) = 0 \]  \hspace{1cm} (2.50)

where \( g \) is the gain, \( l \) is the loss, \( \Omega_{g} \) is the gain bandwidth, \( \Psi = (\Delta \omega_0/c) L \) is a phase slip per cavity roundtrip, \( D = (1/2) \beta_2 d \) is the group delay dispersion parameter for a material of length \( d \) and GVD \( \beta_2 \), \( \gamma \) is the SAM coefficient, and \( \delta \) is the SPM coefficient. The solutions to this equation are in the shape of a sech function given by:

\[ A(t) = A_0 \text{sech} \left( \frac{t}{\tau} \right) \exp \left[ i\beta \ln \text{sech} \left( \frac{t}{\tau} \right) \right] \]  \hspace{1cm} (2.51)

where \( \beta \) is a chirp parameter.
Linear Loss & Phase Shift

$l + i\psi$

Kerr SAM & SPM

Gain

$g\left(1+\frac{1}{\Omega^2_g}\frac{d^2}{dt^2}\right)$

GVD

$iD\frac{d^2}{dt^2}$

Figure 2.12 Schematic of optical elements within a fast saturable absorber passively modelocked laser.

The solution (2.51) can be inserted into the Master equation (2.50) and by separating out real and imaginary terms of order sech and sech$^3$, it is possible to derive the four coupled equations. These equations can then be solved for pulsewidth and chirp parameter for different values of gain, gain bandwidth, GDD, SPM, SAM, and pulse energy. The minimum pulsewidth from the fast saturable absorber modelocked laser can be solved to be [45]:

$$\tau_0 = \left(\frac{2g}{W\Omega^2_g}\right)\left[\frac{2 - 3\beta D_n - \beta^2}{\gamma}\right]$$  \hspace{1cm} (2.52)

where $W = 2A^2\tau$ is the pulse energy, and $\beta$ is the chirp parameter given by:

$$\beta = -\frac{3}{2}\left(\frac{\delta + \gamma D_n}{\delta D_n - \gamma}\right) \pm \left\{\left[\frac{3}{2}\left(\frac{\delta + \gamma D_n}{\delta D_n - \gamma}\right)\right]^2 + 2\right\}^{1/2}$$  \hspace{1cm} (2.53)

where $D_n$ is a normalized dispersion given by:

$$D_n = \left(\frac{\Omega^2_g}{g}\right)D$$  \hspace{1cm} (2.54)
In the simple case of zero dispersion and SPM, equation (2.52) reduces to:

\[ \tau_0 = \frac{4g}{\gamma W \Omega_g^2} \]  \hspace{1cm} (2.55)

This zero dispersion pulsewidth gives an order of magnitude estimate of the pulsewidth possible through fast saturable absorber modelocking. The pulsewidth is determined by an interplay between SAM pulse shortening and gain bandwidth filter pulse broadening. As the pulse energy \( W \) is increased, the pulsewidth will decrease because of increased SAM. Either gain or saturation of the nonlinear effect will eventually limit the pulse shortening influence of increased pulse energy. The pulsewidth can be shortened slightly with soliton effects from introduction of proper amounts of anomalous GDD and SPM [54].

The pulse shortening rate for fast saturable absorber modelocking can be determined by considering the effect of pulsewidth on the nonlinear loss modulation. The modulation window is proportional to the pulsewidth, and the modulation intensity is proportional to the instantaneous intensity. The pulse shortening rate is then:

\[ \text{PSR}_{\text{fast}} = \frac{\gamma W}{2\tau} \]  \hspace{1cm} (2.56)

For fast saturable absorber modelocking, the pulse shortening rate varies inversely with pulsewidth.

2.13 Modelocking Comparison

The pulsewidth generated by a laser will be determined by the value where the PSR and the pulse broadening rate due to pulse spreading mechanisms are equal. The limiting pulse broadening mechanism varies from system to system. Many lasers are limited in pulsewidth by the gain bandwidth. The pulse broadening rate (PBR) due to the limited gain bandwidth can be determined in the parabolic approximation. For a pulse with an initial pulsewidth of \( \tau_0 \sim \text{tbwp}/\Omega_0 \), the pulse shape after transmission through a gain element will be given by:

\[ E'(t) = E_0 \left( 1 - \frac{\omega_0^2}{\Omega_0^2} \right) g \left( 1 - \frac{\omega^2}{\Omega_g^2} \right) = g E_0 \left( 1 - \frac{\omega^2}{\Omega_f^2} \right) \]  \hspace{1cm} (2.57)

where terms of order \( \omega^4 \) have been dropped, and the resultant pulse bandwidth is given by:
\[ \Omega_f^2 = \Omega_0^2 \left( \frac{\Omega_g^2}{\Omega_0^2 + \Omega_g^2} \right) \]  

(2.58)

This new bandwidth is always smaller than the initial bandwidth, indicating pulse broadening. The new pulse width can be determined through the time-bandwidth product:

\[ \tau_0 \Omega_0 = \text{tbwp} = \tau' \Omega_f \]  

(2.59)

The pulse broadening rate for a pass through the gain media is then:

\[ \text{PBR}_{\text{gain}} = \frac{\tau' - \tau_0}{\tau_0} = \frac{\sqrt{\text{tbwp}^2 + \Omega_g^2}}{\Omega_g} - 1 \]  

(2.60)

For pulses with a bandwidth smaller than the gain bandwidth \( \Omega_g \), PBR can be approximated as:

\[ \text{PBR}_{\text{gain}} \approx \frac{\text{tbwp}}{2 \Omega_g^2 \tau_0^2} \]  

(2.61)

The steady-state pulse width of a given modelocked laser can be determined by setting PBR = PSR. Figure 2.13 shows a plot of the PSR and PBR for active, slow saturable absorber passive, and fast saturable absorber modelocking. Fast saturable absorber modelocking outperforms slow saturable absorber modelocking and active modelocking.
2.14 Kerr-Lens Modelocking

Kerr-lens modelocking (KLM) is a form of fast saturable absorber modelocking that has been applied to generate short pulses from a number of laser media [55 - 57]. As discussed above, a saturable absorber must generate larger loss for low intensity cw radiation than for higher intensity pulses within a laser cavity. KLM converts a spatially dependent intensity into a spatially-dependent index of refraction through the Kerr effect. Spatial filtering then converts the intensity-dependent modal profile into SAM. KLM was first used to
modelock a solid-state Ti:Sapphire laser [58] and is useful for a wide variety of solid-state gain media which have a relatively large $n_2$.

The optical Kerr effect acts as the basis for the generation of fast SAM in a KLM laser cavity. The Kerr effect dictates that a material’s index of refraction is of the form in equation (2.27). For Cr$^{4+}$:YAG, $n_2$ is a positive number. An optical pulse propagating through a Kerr media with the lowest order spatial mode will have a Gaussian spacial profile. A Kerr media will then develop a Gaussian profile index of refraction called a Kerr lens. The interaction between the optical pulse and the Kerr lens leads to self-focusing.

A schematic of the process to exploit the optical Kerr effect for KLM is shown below in Figure 2.14. Cw radiation in the optical cavity will have a spatial mode profile determined

![Diagram of Kerr-lens modelocking](image)

**Figure 2.14** Schematic of Kerr-lens modelocking. The Kerr-effect creates an intensity dependent optical mode profile within a laser cavity. An aperture is placed within the cavity to generate higher loss for low intensity (cw) modes than for high intensity (pulsed) modes.

by the cavity mirrors and mirror spacings. Higher intensity pulses, on the other hand, will have a slightly different spatial mode profile within the laser cavity due to self focusing from the Kerr lens. This Kerr mode is often characterized by a smaller beam waist within the laser media and can even have a frequency dependence [59]. A spatial aperture is then placed in the cavity at a location where the Kerr beamwaist is smaller than the linear beamwaist. Either ‘hard apertures’, real apertures placed in the laser cavity [60, 61], or ‘soft apertures’, a spatially-dependent gain, can be designed to create more loss for the cw linear mode than for the pulsed mode [62]. Optical pulses, seeing higher gain than the cw radiation, will then be generated.
Proper cavity design is critical for the optimization of KLM strength within a laser cavity. Several approaches have been taken to model KLM in optical cavities. Hard aperture KLM requires proper placement of a slit within the laser cavity, while soft aperture KLM requires precise alignment of the pump profile and cavity mirror positions. In general, soft aperture KLM will be designed to have a smaller beamwaist for pulses than for cw. The soft aperture is then a pump beamwaist which is smaller than the cw mode and more closely matches the modelocked mode. KLM will be strongest where the derivative of the intensity-dependent beamwaist is largest. As will be shown later, this condition is often met near the edges of the stability region of the laser cavity.

2.15 Soliton Modelocking

Soliton effects can lead to pulse shortening in a modelocked laser cavity. A soliton is a stationary wave solution that is stable because of the balance of chirp from anomalous GDD and the chirp from SPM. Solitons are solutions to the nonlinear Schrodinger equation [38]:

\[
-\frac{i}{\hbar} \frac{\partial A}{\partial z} = \delta |A|^2 A - \frac{1}{2} \beta_2 \frac{\partial^2 A}{\partial t^2}
\]

(2.62)

which includes the effects of both SPM and GDD. There are several solutions to this equation, of which the lowest order is:

\[
A(z, t) = A_0 \text{sech}\left(\frac{t}{\tau}\right) \exp \left[-i \frac{\beta_2}{2 \tau^2} z\right]
\]

(2.63)

where the field amplitude \(A_0\) and the pulsewidth \(\tau\) obey the area theorem:

\[
A_0 \tau = \sqrt{\frac{\beta_2}{\delta}}
\]

(2.64)

which can be rewritten for lowest order solitons as:

\[
L_{D2} L_{NL} = 1
\]

(2.65)

The pulsewidth will therefore be shortest for near-zero GVD and high power levels. Close to zero dispersion, higher order dispersion terms will begin to dominate over second-order dispersion, and the soliton solutions will become more complicated. Under the proper conditions, soliton pulse shaping can result in shorter pulses than are predicted by considering
modelocking as the sole source of pulse shortening. For this reason, KLM lasers typically operate in the slightly anomalous dispersion regime. In Ti:Sapphire lasers, for example, it has been shown experimentally that the shortest pulses are generated in this GVD regime [63].

From equation (2.63), the propagation constant of a soliton is given by:

$$k_s = \frac{\beta_2}{2\tau^2} = \frac{\pi}{4Z_0}$$  \hspace{1cm} (2.66)

There is no wavelength dependence to the soliton propagation constant. $Z_0$ is called the soliton period and is the length over which a soliton will accumulate a phase-shift of $\pi/4$. The propagation constant for light propagating freely through a material with second order dispersion, however, is:

$$k_D = \frac{1}{2} \beta_2 (\omega - \omega_0)^2$$  \hspace{1cm} (2.67)

For a propagation distance $z$, the accumulated phase shift is $\phi = kz$. In a laser cavity, solitons experience periodic perturbations after a propagation length $Z_c$, where $Z_c$ is the cavity length. These perturbations will phase-match light in the soliton and in background radiation if:

$$(k_s - k_D)Z_c = m2\pi$$  \hspace{1cm} (2.68)

where $m$ is an integer. Because coupling from the soliton into radiation will be phase-matched at wavelengths where equation (2.68) is fulfilled, equations (2.67) and (2.68) can be solved together to determine the frequency location of cw sidebands, called Kelly sidebands, generated [64]. The spectral offsets of the Kelly sidebands from the spectrum's center frequency in the case of pure second order dispersion are given by:

$$\Delta \omega_2 = \frac{1}{\tau} \sqrt{\frac{4\pi \tau^2 m}{\beta_2 Z_c}} - 1$$  \hspace{1cm} (2.69)

Soliton-continuum phase matching can also occur for higher-order dispersion [65]. Figure 2.15 shows a schematic of the Kelly side-band phase matching process for GDD and TOD. Freely propagating light develops a parabolic spectral phase profile after propagation through a media with GDD. The Kelly side-bands are then symmetric around $\omega_0$. For TOD, however, the side-bands will not necessarily be symmetric. There will often be a single dominant Kelly side-band at $k_s - k_{\text{TOD}} = 0$ on the long wavelength side of the spectrum.
2.16  Dispersion Managed Modelocking

It would appear from equation (2.64) that the pulsewidth can be reduced indefinitely for a given dispersion by increasing the soliton energy. Unfortunately, pulses will break up due to modulation instability when the cavity length is longer than the nonlinear length \( Z_c \gg L_{NL} \). Dispersion managed solitons can maintain a high energy while maintaining a relatively long nonlinear length. A soliton is considered dispersion managed within a laser cavity if the pulsewidth breaths periodically due to intracavity dispersion. The dispersion managed soliton is still considered to be soliton-like if the pulse shape is periodic with the cavity length. Figure 2.16 shows a schematic of the dispersion in a net zero GDD laser cavity. For a pulse prop-

Figure 2.16  Schematic of a dispersion managed laser cavity. The cavity is considered to be dispersion managed if each element within the laser cavity broadens the pulse significantly.

agating through a dispersion managed laser cavity, it is no longer possible to consider the average dispersion in the Master equation [66, 67].
An order of magnitude estimate of the pulsewidth where a laser becomes dispersion managed is:

\[ \text{GDD}_{\text{element}} > \tau^2 \]  \hspace{1cm} (2.70)

For a 20 fs pulse from a Cr\textsuperscript{4+}:YAG laser, an element with a GDD on the order of 400 fs\textsuperscript{2} can produce pulse broadening. Fiber-laser solitons can become dispersion managed for much longer pulsewidths by introducing large swings in fiber dispersion [25, 68]. Dispersion managed cavities can be roughly characterized by the pulse stretching ratio \( \tau_{\text{max}}/\tau_{\text{min}} \), where \( \tau_{\text{max}} \) and \( \tau_{\text{min}} \) are the maximum and minimum pulsewidths within the laser cavity. Figure 2.17

![Graph showing pulse stretching ratio vs. pulsewidth](image)

**Figure 2.17** Approximate pulse stretching ratio for pulses in a Cr\textsuperscript{4+}:YAG laser consisting of a 2 cm long gain crystal.

plots an estimate of the pulse stretching ratio for a 2 cm Cr\textsuperscript{4+}:YAG laser. The ratio is calculated by calculating the pulsewidth of a pulse after propagation through the 2 cm crystal and
dividing it by the initial pulsewidth. Dispersion management becomes significant for sub-20 fs pulses. It has been shown that dispersion management can significantly shape already short pulses. It is believed that the dispersion managed modelocking dynamics begin to become important at a pulse stretching ratio of ~3. For dispersion managed cavities, the pulsewidth will vary within the laser cavity. It therefore becomes important where the output coupler is placed. For dispersion managed cavities, the shortest pulses will be in the anomalous GDD region.

2.17 Two Mirror Cavity

Modelocking, like lasing in general, is quite sensitive to the alignment of laser cavity elements. Proper cavity design is crucial for optimal laser performance. A laser, in its simplest form, consists of a gain media generating light and an optical cavity providing enough feedback into the laser gain media for stimulated emission to dominate over spontaneous emission. Perhaps the simplest laser cavity consists of two mirrors, shown schematically in Figure 2.18. Two mirrors, with radius of curvature $R_1$ and $R_2$ respectively, are separated by a distance $d$. Not all combinations of $R_1$, $R_2$, and $d$ will support optical cavity modes. The existence condition for modes in a given cavity configuration is determined by the following procedure. A general solution for the field profile in an optical cavity is chosen and mathematically propagated through the cavity. After a roundtrip through the cavity, the resultant field is required to be equal to the input field. An optical cavity mode will exist if it is possible to find such self-consistent solutions.

![Figure 2.18 Schematic of a two mirror laser cavity. The mirrors have a radius of curvature of $R_1$ and $R_2$, and a separation $d$.](image)
Solutions for the field profile inside an optical cavity must obey Maxwell's wave equation, given in equation (2.13). This is the general wave propagation equation governing the fields of electromagnetic radiation. Laser radiation, however, is highly directional, allowing for further simplification of the wave equation. Specifically, in the limit where \( k_z \gg k_x, k_y \), the wave equation is simplified into the paraxial equation [69]:

\[
\nabla^2 u - 2i k_z \frac{\partial u}{\partial z} = 0
\]

(2.71)

where:

\[
\nabla^2 u = \frac{\partial^2 u}{\partial x^2} + \frac{\partial^2 u}{\partial y^2}
\]

(2.72)

and:

\[E(x, y, z) = u(x, y, z)e^{ikz}\]

(2.73)

and the second-order spatial derivative with respect to \( z \) has been replaced by a first-order derivative. The general solution to the paraxial equation are the special functions called Hermite-Gaussians, of which the lowest order function is a simple Gaussian of the form:

\[u(x, y, z) = A \exp \left[ -\frac{(x^2 + y^2)}{\omega_0^2} \right]\]

(2.74)

where \( \omega_0 \), the beam waist, is the radius at which the electric field falls to 1/e of its maximum value. The Gaussian beam is completely characterized by a complex q-parameter, defined by:

\[
\frac{1}{q} = \frac{1}{R} - i\left(\frac{\lambda}{\pi \omega_0^2}\right)
\]

(2.75)

where \( R \) is the radius of curvature of the propagating wave phase front, \( \omega \) is the beam waist, and \( \lambda \) is the wavelength. At the beam focus, \( q \) will be completely imaginary with \( \omega = \omega_0 \). The Gaussian q parameter has the convenient property that it transforms by a bi-linear transformation with the same ABCD matrices of ray optics. In ray optics, the position (r) and slope (r') of an optical ray is transformed through an optical system by the equation:

\[
\begin{bmatrix}
    r_f \\
    r'_f
\end{bmatrix} = \begin{bmatrix}
    A & B \\
    C & D
\end{bmatrix} \begin{bmatrix}
    r_i \\
    r'_i
\end{bmatrix}
\]

(2.76)
where the ABCD matrix is a function of the specific optical system. The Gaussian q parameter transforms according to the same ABCD matrix by the equation:

\[ q_f = \frac{Aq_i + B}{Cq_i + D} \]  \hspace{2cm} (2.77)

where \( q_f \) is the transformed q-parameter, and \( q_i \) is the initial q-parameter. The ABCD matrices for several optical elements are listed below:

\[ M_d = \begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} \]  \hspace{2cm} (2.78)

\[ M_R = \begin{bmatrix} 1 & 0 \\ \frac{2}{R} & 1 \end{bmatrix} \]  \hspace{2cm} (2.79)

\[ M_f = \begin{bmatrix} 1 & 0 \\ -\frac{1}{f} & 1 \end{bmatrix} \]  \hspace{2cm} (2.80)

where \( M_d \) is the matrix for translation a distance \( d \), \( M_R \) is the matrix for reflection from a mirror with a radius of curvature \( R \), and \( M_f \) is the matrix for transmission through a lens of focal length \( f \).

To solve for the existence condition of stable cavity modes, the ABCD matrix for a roundtrip pass through the cavity is used. If a stable optical mode exists, \( q_f \) will be equal to \( q_i \) after that roundtrip. For the two mirror cavity,

\[ \text{ABCD}_{\text{roundtrip}} = M_d \cdot M_{R_1} \cdot M_d \cdot M_{R_2} \]  \hspace{2cm} (2.81)

where \( M_d \) is the ABCD matrix corresponding to propagation through vacuum for a distance \( d \), and \( M_{R_1} \) and \( M_{R_2} \) are the ABCD matrices corresponding to reflection off a mirror with a radius of curvature \( R_1 \) and \( R_2 \) respectively. Solving for \( \text{ABCD}_{\text{roundtrip}} \), and setting \( q_i = q_f \), we can solve for the cavity q parameter. Stable cavity modes will exist for:

\[ 0 \leq g_1 g_2 \leq 1 \]  \hspace{2cm} (2.82)

where \( g_i \) is called the g-parameter given by:
\[ g_i = 1 - \frac{d}{R_i} \]  

(2.83)

for \( i = 1,2 \). The stable cavity modes can then be plotted as a function of the \( g \)-parameters as shown in Figure 2.19. The filled gray phase space region contains stable cavity modes. Dotted lines indicate constant radii of curvature for the cavity mirrors.

![Figure 2.19](image)

Figure 2.19 Stability diagram of the two mirror cavity resonator. Stable cavity modes exist in the gray region. Dotted lines indicate constant mirror radii of curvature for the cavity mirrors.

For resonators where \( R_1 = R_2 \), as the cavity length is varied, a single stability region is observed. For all \( R_1 \) not equal to \( R_2 \), two stability regions are separated by an unstable cavity region.

### 2.18 Four Mirror Cavity

It is possible to extend the analysis of stable cavity modes of a two mirror cavity by adding additional mirrors with different radii of curvature. Three and four mirror cavities are
commonly used for ultrafast lasers. These cavities allow the construction of cavity modes with very small beamwaists in the laser crystal. The small beamwaist in turn creates the high nonlinearities that are needed to generate short light pulses.

In this work, a Z-fold cavity, a variant of the four mirror cavity, was constructed for the generation of short pulses from a Cr\textsuperscript{4+}:YAG laser. A schematic of a Z-fold laser cavity and an equivalent representation used for modeling are shown below in Figure 2.20. The cavity consists of a laser crystal of thickness \( t \) surrounded by two cavity fold mirrors. The cavity fold mirrors are separated from the crystal face by a distance \( d_1 \) and \( d_2 \), and have a radius of curvature of \( R_1 \) and \( R_2 \), respectively. After reflection from the curved mirrors, the cavity light is directed to flat end mirrors a distance \( L_1 \) and \( L_2 \) away from the curved mirrors. Astigmatism, generated by the off-axis reflection off the curved cavity fold mirrors and the Brewster-Brew-
ster cut laser crystal, is neglected in this model which is discussed in Section 2.19. The focal length of the effective lenses are half the radius of curvature of the corresponding mirror.

The four mirror cavity is numerically modeled in the same way as the two mirror cavity. In fact, it is possible to map the four mirror cavity into an effective two mirror cavity, and compute the stability regions of the laser cavity using the formulae developed in the two mirror cavity case. This effective cavity is found in the following manner. For each side of the laser cavity, the ABCD matrices for the two mirrors and translations are replaced by a single effective mirror and translation:

\[ M_{\text{eff}} \cdot M_{\text{eff}} \cdot M_{\text{eff}} = M_{d1} \cdot M_{f1} \cdot M_{L1} \cdot M_{fL1} \cdot M_{d1} \]  
(2.84)

By doing this computation, we are able to calculate the values for \( d_{\text{eff}} \) and \( R_{\text{eff}} \), given by:

\[ d_{\text{eff}} = d_1 - f + \frac{f^2}{f - L_1}, \quad R_{\text{eff}} = \frac{f^2}{f - L_1} \]  
(2.85)

where \( f = R/2 \). Using these effective values, the stability regions of the four mirror cavity are identical to those in equation (2.82) for a two mirror cavity consisting of mirrors with radius of curvature \( R_{\text{eff}} \) and \( R_{2\text{eff}} \) and a separation of \( d = d_{\text{eff}} + d_{2\text{eff}} \).

In order to calculate the beam waist for different cavity configurations the ABCD matrix for propagation roundtrip through the laser cavity is computed. The \( q \) parameter is then found self-consistently by setting the initial \( q \) value before the roundtrip propagation equal to the \( q \) value after a roundtrip propagation. For all values of \( q \) that are complex, a beam waist and optical mode will exist. The ABCD matrix for a roundtrip starting and ending at the face of the laser crystal is:

\[ \text{ABCD}_{\text{roundtrip}} = M_t \cdot M_{d1} \cdot M_{f1} \cdot M_{L1}^2 \cdot M_{fL1} \cdot M_{d1} \cdot M_t \cdot M_{d2} \cdot M_{f2} \cdot M_{L2}^2 \cdot M_{f2} \cdot M_{d2} \]  
(2.86)

where \( M_t \) is the matrix for propagation through the laser crystal of thickness \( t \) and index of refraction \( n \); \( M_{d1}, M_{d2}, M_{L1}, \) and \( M_{L2} \) are the ABCD matrices for propagation through a distance \( d_1, d_2, L_1, \) and \( L_2 \) respectively; and \( M_{f1} \) and \( M_{f2} \) are the ABCD matrices for reflection off a mirror of radius of curvature \( R_1 = 2f_1 \) and \( R_2 = 2f_2 \) respectively. As before, \( q_f \) is set equal to \( q_t \) to solve for \( q \). The beam waist is then computed from the imaginary part of \( q \).
Figure 2.21 Calculated beamwaist ($\omega_0$) for a symmetric Z-fold laser cavity as a function of the curved mirror separation $d = d_1 + d_2 + t$.

The beamwaist at the center of the crystal is shown in Figure 2.21 for typical parameters for the Cr\textsuperscript{4+}:YAG laser. In this plot, the beamwaist for a symmetric cavity ($L_1 = L_2$) is computed. For this calculation, $R_1 = R_2 = 10$ cm, $L_1 = L_2 = 60$ cm, $t = 2$ cm, and $d = d_1 + d_2 + t$. The cavity fold mirror separation $L$ is then varied to see how the beamwaist varies within the laser crystal. A single stability region, a region where modes exist, is visible from $d = 10.9$ to $11.8$ cm.

A second, asymmetric, laser cavity is modeled, and the beamwaist for different cavity fold mirror distances is shown in Figure 2.22. For this computation, $L_1 = 50$ cm and $L_2 = 70$ cm. All other parameters are the same as the previous plot. As can be seen, the single stability region has now split into two, with a small region of no stable optical modes in between.
2.19 Astigmatism Compensation

Astigmatism results from a non-cylindrically symmetric optical element in the light beam path. In a Z-fold cavity, transmission through a Brewster-Brewster cut laser crystal introduces astigmatism. The intracavity optical mode will be characterized by a tangential (parallel to the plane of incidence) and a sagittal (perpendicular to the plane of incidence) beamwaist. In general, the tangential and sagittal beamwaists within a laser cavity can be quite different, resulting in an oval-shaped beam. Because Kerr-lens modelocking is highly sensitive to the alignment and beamwaist profile of the cavity, uncompensated astigmatism can reduce the KLM strength significantly by shrinking the stability region. It is possible to compensate astigmatism from the laser rod by introducing a second counteracting intracavity astigmatic element [70]. In a Z-fold cavity, the laser crystal astigmatism is compensated by rotating the curved cavity fold mirrors surrounding the crystal ($R_1$ and $R_2$ in Figure 2.20). It is
important to note that while it is possible to overlap the sagittal and tangential stability regions, arbitrary positions within the Z-fold cavity will have an astigmatically-shaped optical mode that will have an effect on modelocking [71, 72].

The proper astigmatically compensated cavity design can be found by solving the ABCD matrices for the sagittal and tangential axes and then setting equal the stability regions. First, the sagittal and tangential ABCD matrices for off-axis reflection from a curved mirror are:

\[
M_{\text{Rf}} = \begin{bmatrix}
1 & 0 \\
\frac{2}{R \cos \theta} & 1
\end{bmatrix} 
\]  
(2.87)

\[
M_{\text{Rs}} = \begin{bmatrix}
1 & 0 \\
\frac{2 \cos \theta}{R} & 1
\end{bmatrix} 
\]  
(2.88)

where \( \theta \) is the angle of the off-axis rotation of the curved mirror. The ABCD matrices can be computed for transmission through a Brewster cut rod using equations (2.75) and (2.77). The q-factor for the beamwaist is \( q \sim \omega^2/\lambda \). At the Brewster interface, the effect on the tangential and sagittal beamwaist must be considered. In the sagittal direction, \( \omega \) is the same on both sides of the interface. However, the propagation direction changes by Snell’s law in the paraxial limit from \( \theta \) to \( \theta n \). On the tangential axis the beamwaist is shrunk from \( \omega \) to \( (\cos \theta / \cos \theta_j) \omega \), which at Brewster’s angle is \( n \omega \). Therefore, the ABCD matrices for entering a material at Brewster’s angle are:

\[
M_{\text{int}} = \begin{bmatrix}
1 & 0 \\
0 & \frac{1}{n^3}
\end{bmatrix} 
\]  
(2.89)

\[
M_{\text{ins}} = \begin{bmatrix}
1 & 0 \\
0 & \frac{1}{n}
\end{bmatrix} 
\]  
(2.90)

while the ABCD matrices for exiting a material at Brewster’s angle are related by:

\[
M_{\text{out}} = M_{\text{in}}^{-1} 
\]  
(2.91)
In each case, the ABCD matrix for transmission through the entire Brewster-Brewster rod is:

\[ M_B = M_{\text{out}} M_d M_{\text{in}} \]  \hspace{1cm} (2.92)

where \( M_d \) is the ABCD matrix for propagation through a distance \( d \) in the rod material of index \( n \). The resulting ABCD matrices are then:

\[ M_{BL} = \begin{bmatrix} 1 & \frac{d}{n^3} \\ 0 & 1 \end{bmatrix} \]  \hspace{1cm} (2.93)

\[ M_{BT} = \begin{bmatrix} 1 & \frac{d}{n} \\ 0 & 1 \end{bmatrix} \]  \hspace{1cm} (2.94)

By matching up the sagittal and tangential stability regions, it is possible to determine the optimal astigmatism compensation angle. The length of crystal compensated by an angle \( \theta \) is given by [70]:

\[ d(\theta) = \frac{R n^3 \sin \theta \tan \theta}{n^2 - 1} \]  \hspace{1cm} (2.95)

For example, 1 and 2 cm long Brewster-Brewster cut Cr\(^{4+}\):YAG laser crystals are compensated by 10 cm radius of curvature mirrors rotated by an angle of 11.2 and 15.8 degrees respectively.

2.20 \hspace{0.5cm} KLM Alignment

Kerr lens modelocking is optimal in laser cavities with a beamwaist that is highly dependent on the intracavity intensity. In particular, there should be some spot in the cavity where the beamwaist shrinks strongly as a function of intensity due to Kerr lensing. An aperture can then be placed at that cavity position to add loss for less intense cw light. Because the beamwaist within the laser crystal is often highly sensitive to intensity, gain, or ‘soft’, aperturing can be used to achieve KLM.

KLM requires a critical alignment of cavity parameters. Several groups have used ABCD matrix modelling methods to determine theoretically the proper configuration to maximize KLM strength [73 - 83]. The most common cavity configuration for KLM is the Z-fold
cavity, because it allows tight focusing within the laser crystal. The focusing increases non-linearity within the crystal. Experimental work has shown the complexity of the effects involved in modelocking [84]. Compact cavities [85 - 87], low threshold cavities [88], and variations on the Z-fold cavity [89] have all been shown to support KLM.

The self amplitude modulation (SAM) coefficient for KLM is given by the derivative of the beamwaist with respect to intensity [90]:

$$\gamma = \left( \frac{1}{\omega} \frac{d\omega}{dp} \right)_{p = 0}$$  \hspace{1cm} (2.96)

where $\omega$ is the beamwaist and $p$ is power. The maximum value for the coefficient is given by:

$$|\gamma|_{\text{max}} = \frac{1}{4\sqrt{g_1 g_2 (1 - g_1 g_2)}}$$  \hspace{1cm} (2.97)

where $g_1$ and $g_2$ are the same g-coefficients as for the two mirror cavity. SAM is therefore maximized when $g_1 g_2 = 0$ or 1, which according to equation (2.82) are the edges of the stability region. Graphically, it can be seen in Figure 2.22 that the beamwaist changes dramatically near the edges of the stability region.

A few qualitative design constraints can be used to optimize the Z-fold cavity for KLM. The most sensitive parameters to be varied to achieve modelocking are the curved mirror separations and laser crystal position within the cavity. For an asymmetric cavity (L_1 is not equal to L_2 in Figure 2.20), two distinct stable cavity mode regions will be encountered by varying the curved mirror separation ($d_1 + d_2 + t$). As the curved mirror separation is varied, a strong negative value of $\gamma$ is obtained at the beginning of the second stability region. Asymmetric cavities are optimal when the cavity arm lengths have a ratio of about 5:4, where the longer arm contains any dispersive elements like prisms and the shorter arm contains the output coupler. Any hard apertures added to increase KLM should ideally be near the output coupler and clip the beam in the tangential plane. Because no region exists in the center of the stability region with no cavity modes, it is more difficult to find the KLM region of symmetric cavities ($L_1 = L_2$). On the other hand, KLM is less sensitive to cavity alignment and can permit self-starting. Symmetric cavities have one large stability region as the curved mirror separation is varied, and the optimal position is near the center of the stability curve.

Most of the rules of thumb for optimizing KLM were developed for general Z-fold cavity resonators. It is possible that novel cavity designs could be used to increase the KLM strength. The shortest pulses generated from a Ti:Sapphire laser, for example, were generated
from a double Z-fold cavity designed to independently vary intracavity KLM and SPM [20]. Furthermore, several effects can change the cavity modes dramatically. Thermal lensing, turns a temperature profile in a laser crystal into a radial index profile. Thermal lensing is particularly strong in Cr$^{4+}$:YAG lasers and has even been used to create a stable resonator where in the absence of thermal lensing there would be none [91]. More precise computational modeling could lead to improved cavity design for KLM. Most ABCD matrix techniques approximate short optical pulses as cw light with higher intensities. Full space-time modeling which considers the time-dependent profile of pulses could give better modelocking models [92 - 94].

2.21  Autocorrelation

Direct electronic detection is inadequate to determine the pulsewidth of femtosecond pulses. Typical photodetectors respond on a time-scale of ~ 10 ps to changes in intensity and are therefore much too slow to resolve the intensity profile of femtosecond optical pulses. Several optical techniques have been developed to use short optical pulses to measure their own pulsewidths [95 - 100]. Perhaps the simplest of those techniques, autocorrelation, uses the ultrafast pulse to measure its own pulsewidth through nonlinear interactions [101, 36]. This technique is useful when there is a pulsetrain of identical pulses and when certain simplifying assumptions can be made about the pulse shape. Autocorrelations can be performed in a collinear interferometric, noncollinear noninterferometric, or background-free noncollinear geometry. Either second harmonic generation (SHG) from a phase-matched nonlinear crystal or two photon absorption (TPA) from a semiconductor photodiode can be used as the nonlinear media through which a pulse interacts with itself.

A schematic of a typical interferometric autocorrelator is shown in Figure 2.23. The autocorrelator is essentially a Michelson interferometer with a variable path length in one arm. An input beam consisting of an optical pulsetrain is split into two with a beam splitter. The two resultant optical pulses travel down the respective arms of the interferometer and are retroreflected by two 45 degree mirrors. The optical path length in one arm is varied by a time delay $\Delta \tau$ with respect to the second with either a motorized translation stage or an oscillating speaker. In the collinear geometry, the two pulses are recombined on a second beam splitter. In the noncollinear geometry, one arm is reflected to a path parallel but not overlapped with the second.
Figure 2.23  Schematic of a pulsewidth measuring autocorrelator. An input beam of laser pulses is split by a beam splitter. One beam ($\omega_1$) travels a length $L$, while a second beam ($\omega_2$) travels a distance $L+\Delta L$, which results in a relative time delay $\Delta \tau = \Delta L/c$. The beams are recombined with a second beamsplitter, and focused through a nonlinear crystal. The nonlinear crystal creates sum frequency generation from path 1 and path 2, and simple harmonic generation from two path 1 or two path 2 photons.

The pulses then are focused onto a nonlinear element that performs the second-order correlation between the two pulses. This nonlinear element could be a nonlinear crystal which has high SHG efficiency or a two photon absorption (TPA) photodiode. If a nonlinear crystal is used, the sum frequency generation is detected by a photodiode or photomultiplier tube (PMT). The nonlinear crystal itself must be chosen such that the dispersive length in the crystal is long compared with the crystal thickness. Phase matching should occur over a bandwidth larger than the pulse spectrum in order to fully resolve the pulsewidth.

TPA based detectors have proven to be quite efficient in measuring autocorrelations and are particularly simple to align [102, 103]. The detector's absorption length is generally
shorter than the dispersion length in the photodetector. Also, the bandwidth of a semiconductor photodetector is generally quite large. Pulses as short as 6 fs were measured from a Ti:Sapphire laser with a TPA GaAsP photodiode [104]. Care must be taken to operate in the regime where TPA is the dominant absorption mechanism. The entire laser bandwidth should fit within the frequency range \(2\omega_g >> 2\omega_{bw} >> \omega_g\), where \(\omega_g\) is the bandgap frequency of the photodetector and \(\omega_{bw}\) is any frequency within the bandwidth of the optical pulse spectrum. Si photodiodes and GaAs light emitting diodes are particularly well suited for measuring pulses at 1500 nm [105]. The efficiency of several TPA-based detectors for autocorrelators are compared and analyzed in Reid et. al. [106].

The optical pulsewidth is determined from the two-photon signal detected after the nonlinear crystal or from the TPA photodiode [107]. The sum of the electric fields of two pulses, each with field magnitude \(E(t)\) and delayed by \(\tau\) is:

\[
E_{\omega}(t) = E(t) + E(t + \tau)e^{-i\omega\tau}
\]  
(2.98)

A nonlinear crystal or TPA measures a signal at twice the frequency with a magnitude proportional to the intensity \(E_{2\omega}^2(t)\):

\[
I_{2\omega}(t, \tau) = |E_{2\omega}(t)|^2 \propto E^4(t) + E^4(t + \tau) + (4 + 2\cos 2\omega\tau)|E(t)E(t + \tau)|^2 + 4\cos \omega\tau|E(t)E^3(t + \tau)| + |E^3(t)E(t + \tau)|
\]  
(2.99)

Because realistic photodetectors have a time response slower than the pulsewidth, the measured signal at 2\(\omega\), assuming a symmetric pulse shape, will be given by the time averaged signal:

\[
D(\tau) = \int I_{2\omega}(t, \tau)dt \approx 1 + (2 + \cos 2\omega\tau)G_2(\tau) + 4\cos \omega\tau B(\tau)
\]  
(2.100)

where \(G_2(\tau)\) is the second order autocorrelation function given by:

\[
G_2(\tau) = \frac{\int E^2(t)E^2(t - \tau)dt}{\int E^4(t)dt} = \frac{\int I(t)I(t - \tau)dt}{\int I^2(t)dt}
\]  
(2.101)

and \(B(\tau)\) is given by:

\[
B(\tau) = \frac{\int E^3(t)E(t - \tau)dt}{\int E^4(t)dt}
\]  
(2.102)
The values of $G_2(\tau)$ and $B(\tau)$, dependent on the pulse envelope shape $E(t)$, are tabulated for Gaussian and sech-shaped pulses below in Table 2.2 [108]. Certain properties of the detected autocorrelation $D(\tau)$ can be determined by the forms of $G_2(\tau)$ and $B(\tau)$. Both $G_2(\tau)$ and $B(\tau)$ have a maximum value of 1 at $t = 0$ and a value of 0 at $t = \infty$.

Interferometric, intensity, and background-free intensity geometries are commonly used to measure the autocorrelation function, and are shown schematically in Figure 2.24. The measured signal from each method was fit using equation (2.100) to determine the pulse-width. By using a collinear geometry (Figure 2.24 (a)), and a scan speed slower than the detector response time in order to resolve interference fringes in the cos($\omega t$) terms, it is possible to record an interferometric autocorrelation function. The measured signal is then equal to $D(\tau)$ in equation (2.100). The function has a value of 8 at $t = 0$, has a minimum value close to 0 at $\omega \tau = \pi/2$, and has a background value of 1 for large time delays. The interferometric autocorrelation function is particularly useful for short optical pulses, where it is possible to use the interference fringes to calibrate the scale of the x-axis. The dependence of the interferometric autocorrelation shape on a variety of pulse parameters such as chirp for realistic detectors is discussed in detail in other sources [108, 109].

In intensity autocorrelations (Figure 2.24 (b)), the detector averages over the interference terms $\sim \cos(\omega t)$ or $\cos(2\omega t)$. This can be done by time averaging (choosing a detector that is too slow to vary with the interference terms) or spatially (by using noncollinear beams which spatially average between different phases of the interference terms). In this case, equation (2.100) reduces to:

$$D_{\text{intensity}}(\tau) = 1 + 2G_2(\tau) \quad (2.103)$$

This function has a maximum value of 3 at $\tau = 0$, and a background level of 1.

<table>
<thead>
<tr>
<th>$E(t)$</th>
<th>$t_{\text{fwhm}}/t_p$</th>
<th>$G_2(\tau)$</th>
<th>$B(\tau)$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\exp[-t^2/2t_p^2]$</td>
<td>$2 \ln 2$</td>
<td>$\exp[-t^2/2t_p^2]$</td>
<td>$\exp[-3 \tau^2/8t_p^2]$</td>
</tr>
<tr>
<td>$\text{sech}[t/t_p]$</td>
<td>1.7627</td>
<td>$3[(\tau/t_p)\cosh(\tau/t_p) - \sinh(\tau/t_p)]$</td>
<td>$3[\sinh(2 \tau/t_p) - 2 \tau/t_p]$</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\sinh^{-3}(\tau/t_p)$</td>
<td>$\sinh^{-3}(\tau/t_p)$</td>
</tr>
</tbody>
</table>

Table 2.2 Values of $G_2(\tau)$ and $B(\tau)$ for Gaussian and sech shaped pulse envelopes [108].

ULTRAFAST LASER DESIGN

89
Figure 2.24  (a) Schematic of interferometric, (b) intensity, and (c) background-free intensity autocorrelation detection methods.

The background in the intensity autocorrelation function is from second harmonic generation (SHG) from each individual beam, rather than sum frequency generation (SFG) from the two beams. It is possible to remove this background with a spatial filter in the non-collinear geometry (Figure 2.24 (c)).
Figure 2.25 Momentum conservation for simple harmonic generation (SHG) and sum frequency generation (SFG) in a nonlinear autocorrelator. The SHG (2k₁ and 2k₂) travel a path spatially separated from the SFG (k₁ + k₂).

The propagation direction for the SHG and SFG will be determined by momentum conservation, and shown in Figure 2.25. The SHG from each the individual beam path will travel in the same direction as the beam before doubling, while the SFG from the cross correlation will propagate in between the two SHG signals. The SHG can then be spatially filtered, and the background-free detected intensity autocorrelation function will be:

\[ D_{\text{intensity}}(\tau) = 2G₂(\tau) \]  \hspace{1cm} (2.104)

This function has no background and directly measures the second-order autocorrelation function.

A plot of the calculated intensity, background-free intensity, and interferometric autocorrelations for a 20 fs Gaussian pulse at 1500 nm are shown in Figure 2.26. The interferometric autocorrelation has a peak signal to background to minimum signal ratio of 8:1:0, while the intensity autocorrelation has a signal to background ratio of 3:1. The background-free autocorrelation has no background.
Interferometric autocorrelations are better than intensity autocorrelations for very short pulses because of the limited time resolution of noncollinear autocorrelators. A schematic of the time-ambiguity introduced by a noncollinear autocorrelator is shown in Figure 2.27. Two noncollinear beams are focused to a point at an angle $\alpha = \tan^{-1}(D/f)$, where $D$ is the beam separation and $f$ is the lens focal length. Both pulses have a beamwaist $\omega$ and a length of $d = c\tau$, where $\tau$ is the pulsewidth. Because the beam has a waist $\omega$, there is an uncertainty in the delay given by $\Delta\tau = \Delta L/c = (\omega/c) \sin 2\alpha$. The uncertainty arises because the pulses will overlap over the entire time-span $\Delta\tau$. If $\Delta\tau$ is close to the pulsewidth $\tau$, the autocorrelator will fail to fully resolve the pulsewidth. For a beam separation of 0.5 cm, a 2.5 cm focal length lens, focused to a beam size of 20 $\mu$m, the time resolution will be around 25 fs.
Figure 2.27 Schematic of the time-ambiguity introduced by the noncollinear autocorrelator geometry. The input beams, separated by a distance $D$, are focused by a lens of focal length $f$ into a nonlinear element. The autocorrelator cannot resolve time delays shorter than $\Delta \tau$.

2.22 Conclusions

$\text{Cr}^{4+}:\text{YAG}$ is an ideal material for the generation of ultrafast optical pulses near 1500 nm because of the material's large gain bandwidth. Modelocking can be used to generate pulses rather than cw light from the laser. Kerr lens modelocking in particular is ideal because of the large pulse shortening rate for short pulses and the large Kerr constant of $\text{Cr}^{4+}:\text{YAG}$. A $Z$-fold cavity was designed to create a stable cavity that maximizes KLM. After pulses are generated, autocorrelation traces can be measured and fit to determine the pulsewidth.
CHAPTER 3: Cr$^{4+}$:YAG LASER

3.1 Introduction

Optical pulses shorter than 20 fs with 400 mW average power at a 110 MHz repetition rate have been generated by a Cr$^{4+}$:YAG laser using only double-chirped mirrors for dispersion compensation. The corresponding pulse spectrum has a peak intensity at 1450 nm and extends from 1310 to 1500 nm full-width half maximum. These pulses, the shortest generated to date from a Cr$^{4+}$:YAG laser, are only four optical cycles within the FWHM intensity width. Several additional laser cavity configurations have been constructed to determine the dependence of Cr$^{4+}$:YAG laser performance the intracavity dispersion.

3.2 Background

Cr$^{4+}$:YAG based ultrafast lasers are able to produce short pulses in the wavelength range from 1300 to 1600 nm [110 - 114]. Unlike color-center lasers, Cr$^{4+}$:YAG lasers operate at room temperature and do not require a vacuum to protect the integrity of the gain media. Furthermore, these lasers have a much larger gain bandwidth than either bulk Er-doped glass lasers [24] or Er-doped fiber lasers [25]. Previously, Tong et. al. [115] generated 43 fs pulses from a hard aperture Kerr-lens modelocked (KLM) Cr$^{4+}$:YAG laser with fused silica prisms for dispersion compensation. The minimum pulsewidth in this laser was reportedly limited by third-order dispersion (TOD) [116]. Other groups have used saturable absorber mirrors to
self-start modelocking [117 - 122] or reduced the laser cavity length to create high-repetition rate [91, 123 - 126] Cr⁴⁺:YAG lasers. Because of its large bandwidth, Cr⁴⁺:YAG is also ideal for broadband cw tuning. Tunable single frequency operation was demonstrated in a ring cavity configuration with an intracavity etalon [127]. Diode pumping has also been shown to be suitable for cw operation [128], but has not yet been demonstrated to be suitable for modelocking.

To generate the shortest possible optical pulse, chromatic dispersion within the laser cavity must be controlled over a large wavelength range. In standard modelocked laser cavities, geometric and material dispersion from prism pairs are used to compensate the group delay dispersion (GDD) of the laser gain media, but often limit the minimum achievable pulse width due to higher-order dispersion [129]. As an alternative to prisms, chirped mirrors have been developed to both compensate higher-order dispersion and provide a large mirror bandwidth [130]. More recently, double-chirped mirrors (DCMs) have been introduced to reduce residual oscillations in the dispersion of chirped mirrors caused by spurious reflections within the mirror due to improper impedance matching of the incident optical wave [132 - 135]. DCMs have been used for intracavity dispersion compensation to generate two-cycle optical pulses from Ti:sapphire lasers with pulsewidths as short as 5 fs at 800 nm [20, 136, 137] and three-cycle pulses with 14 fs duration at 1300 nm from Cr:Forsterite [138].

In this chapter, the use of DCMs for broadband higher-order dispersion compensation of a Cr⁴⁺:YAG laser is discussed. Pulses shorter than 20 fs were measured directly from the prismless laser [139]. The pulse spectrum was peaked at 1450 nm and extended from 1310 to 1500 nm full-width half maximum. On a logarithmic scale, spectra from 1140 to > 1700 nm, the limit of our optical spectrum analyzer, was observed.

The results discussed here were made possible by the important contributions of several coworkers. The enabling technology for the short pulse generation is the double-chirped mirror design by Prof. Franz Kärtner. Dr. Uwe Morgner wrote a fantastic code to retrieve the pulse envelope and phase profile of short optical pulses from interferometric autocorrelations. Juliet Gopinath assisted with the autocorrelator and building the laser itself. Drew Kowalevicz assisted with the characterization of the dispersion of the intracavity mirrors. Peter Rakich and Hanfei Shen aided in the experiments, and helped with useful discussions.
Continuous Wave Cr⁴⁺:YAG

A continuous wave (cw) laser was first constructed to evaluate the suitability of several different Cr⁴⁺:YAG crystals and mirrors for use within an ultrafast laser cavity. A schematic of the cw laser cavity is shown below in Figure 3.1.1. The laser cavity is in a Z-fold configuration, which is often used for ultrafast laser cavities because of the small beam waist generated between the two curved fold mirrors (M1 and M2). The small beam waist allows optimal cavity design for nonlinear modelocking techniques such as Kerr lens modelocking (KLM). The Z-fold cavity can also minimize the effects of astigmatism from the Brewster angle cut crystal interfaces. Chapter 2 contains a detailed description of some design issues involved in constructing this laser cavity.

Each Cr⁴⁺:YAG crystal tested was a Brewster-Brewster cut rod supplied by A. V. Shestakov at E.L.S. Company. A Spectra-Physics 11 W 1064 nm diode pumped Nd:YVO₄ laser was chosen to pump the laser crystals because of their high absorption from 950 to 1100 nm. The pump laser was expanded a factor of two by a telescope consisting of a 5 cm focal length lens separated 5 cm from a 10 cm focal length lens, and then focused into the Cr⁴⁺:YAG crystal by second 10 cm focal length lens. The pump beam waist after the telescope was estimated to be 2.4 mm by a knife-edge measurement; the beam waist in free space after the 10 cm lens would then be 15 μm. This particular series of mode-matching lenses was chosen to focus the pump beam slightly tighter than the laser cavity mode and maximize KLM. The laser rod was surrounded by two 10 cm radius-of-curvature cavity fold mirrors (M1 and
In every case the cavity fold mirrors were dichroic mirrors designed to be highly reflective for the lasing wavelength and highly transmissive for the pump wavelength. Both cavity fold mirrors are rotated slightly from normal incidence (11 degrees for a 1 cm crystal, or 16 degrees for a 2 cm crystal) to compensate for astigmatism from the Brewster cut laser crystal. One arm of the laser, 60 cm long, contains a quarter-wave stack highly reflecting mirror (M3), while a second arm of the laser, 50 cm long, contains a broadband output coupler (OC).

Plots of the reflectivity of several output couplers and high-reflectors used within this laser cavity are shown in Figure 3.2 (a) and (b). The reflectivity of each mirror was measured using a Cary spectrophotometer, which takes the ratio of the transmitted light intensity through the mirror and the transmitted light intensity through an empty reference cell. Both high-reflectors (HR1 and HR2) have the broadband reflectivity near 1500 nm, and low reflectivity at the pump wavelength of 1064 nm. Many output couplers (OC1 - OC5) were measured, each with 0.5 - 2% minimum transmission, and slightly different wavelength characteristics. Almost all highly reflective cavity mirrors and output couplers were purchased from CVI Laser Corporation.

Several Cr\textsuperscript{4+}:YAG crystals were tested to identify those of high quality, and a high degree of variation in laser crystal performance was observed. While a few crystals were found to be of suitable quality for use within a laser cavity, most were plagued by spatial inhomogeneities, low gain, or low absorption-saturation intensities. The sparse availability of high quality laser crystals remains a limiting factor for the further development of these lasers. To facilitate cooling of the crystals, each was mounted in a form-fitting copper mount. The crystal was wrapped in Indium foil in order to make good thermal contact between the copper mount and crystal. The copper mount itself was then cooled with cold flowing water, held at 13 C.

Cr\textsuperscript{4+}:YAG laser crystals provide lower gain than the gain media used in many solid-state lasers. Excited state absorption (ESA) of both the lasing wavelength as well as the pump wavelength reduce the gain and quantum efficiency. Therefore 2 cm long crystals, much longer than the \~3 mm crystals used in ultrafast Ti:sapphire lasers, are typically used to provide enough gain for lasing. The long crystal complicates modelocked laser design by introducing large intracavity dispersion and requiring a large astigmatism compensation cavity fold mirror angle. Conventional wisdom for ultrafast laser design dictates that a laser should use the shortest possible crystal as a gain media. While prisms can easily compensate intracavity GDD, they will typically leave higher-order dispersion uncompensated. The amount of uncompensated higher-order dispersion will be proportional to the length of the crystal.
Figure 3.2 Reflectivity of (a) highly reflecting mirrors and (b) output coupling mirrors used in Cr\textsuperscript{4+}:YAG laser cavities.
In principle, all higher-order dispersion can be eliminated by using chirped mirrors for dispersion compensation. Chirped mirrors can compensate both GDD and higher-order dispersion at the same time. The larger astigmatism compensation angle ($\theta_{Ast}$) of the cavity folding mirrors (M1 and M2) creates additional challenges for the generation of short pulses. At off-normal incidence angles, the standard Bragg mirror stopband will be shifted to shorter wavelengths. Furthermore, the double-chirped mirrors (DCMs) antireflection properties will degrade for larger incidence angles. As discussed in the previous chapter, this will lead to increased spectral oscillations in the GDD.

As an initial test of their quality, the pump absorption efficiency of each Cr$^{4+}$:YAG crystal was measured. The wallplug efficiency is the product of the quantum efficiency (the percent of absorbed pump photons converted to emitted lasing photons) and the pump absorption efficiency (the percent of pump photons absorbed). Crystals with either small linear absorptions or low pump-absorption saturation intensities stand little chance of providing enough gain for successful laser operation. The pump absorption through the laser crystals was measured as a function of power density. The power density was varied by sweeping the total pump power for different degrees of pump focusing.

The unfocused Nd:YVO$_4$ pump beamwaist was first measured using a knife-edge technique. A razor blade was translated into the pump beam while simultaneously measuring the transmitted power. It is assumed that the pump laser was emitting a purely TEM$_{00}$ Gaussian mode. The Gaussian intensity profile of the Nd:YVO$_4$ beam is given by:

$$A(x, y) = A_0 \exp\left( -\frac{x^2 + y^2}{\omega_0^2} \right)$$  \hspace{1cm} (3.1)

where $\omega_0$ is the beam waist, $x$ and $y$ are the spatial coordinates orthogonal to the beam propagation direction, and $A_0$ is the amplitude of the intensity profile. The beam is roughly collimated at the output of the laser, and so the beamwaist is assumed to be the minimum value. As a razor-edge is translated through the Gaussian beam, the power transmitted is thus given by the function:

$$P(x) = \int \int A(x, y) dx dy$$  \hspace{1cm} (3.2)

where $P(t)$ is the transmitted power, the razor blade blocks light for $x < x_r$, and the beam is centered at the origin of the coordinate system ($x=0$, $y=0$). This function can be used to fit
measured data to determine the beam waist. Data from a knife-edge measurement of the Nd:YVO₄ beam waist is shown in Figure 3.3. The data fits well to a 1.2 mm beam waist.

![Graph](image)

**Figure 3.3** Knife-edge measurement of the Nd:YVO₄ beam waist. Data (black squares) agree well with a fit function (gray line) assuming a 1.2 mm beam waist.

Using the experimentally determined pump spot size, Gaussian beam optics was used to determine the power density at the beam focus for different lenses. For a lens of a given focal length, the beam was focused to a beam waist:

\[
\omega_0 = \frac{f \lambda}{\pi r}
\]

(3.3)

where \(\omega_0\) is the minimum focused beam waist, \(r\) is the radius of the input beam, and \(f\) is the focal length of the focusing lens. The focusing beam will have a confocal parameter \(b\), the distance over which the focused beam will have a beam waist smaller than \(2^{1/2}\omega_0\), given by:
\[ b = \frac{\pi \omega_0^2}{\lambda} \]  

(3.4)

For a given confocal parameter, the beam waist as a function of position is given by:

\[ \omega(z) = \omega_0 \sqrt{1 + \left(\frac{z}{b}\right)^2} \]  

(3.5)

where \( z \) is the distance from the beam focus. For the computation of the pump absorption, the power density was determined with the minimum value of beam waist. The finite confocal parameter of the focused beam was neglected. This is a rough approximation because the laser crystals are longer than the confocal parameter for most of the lenses used. Thermal lensing effects through the crystal have also been neglected.

A measurement of the transmission fraction of a 1 cm and 2 cm Cr\(^{4+}\):YAG crystal for different pump power densities is shown in Figure 3.4. All power that was not transmitted is

![Figure 3.4](image)

Figure 3.4 Measurement of the transmission through a 1 cm (unfilled circles) and 2 cm (filled shapes) long Cr\(^{4+}\):YAG crystal for different pump power densities. The operating power density is indicated by a vertical dotted line.
assumed to have been absorbed by the crystal. The effects of Fresnel losses from mirrors or air-glass interfaces were taken into account.

There were several differences between the 1 cm and 2 cm crystals studied. The 1 cm crystal was more heavily doped than the 2 cm crystal in an effort to maintain high gain for the shorter crystal. Also, the 1 cm rod has a 2 x 5 mm rectangular cross-section, while the 2 cm rod has a 3 mm diameter circular cross section. Both crystals exhibited significant saturation of pump absorption over the pump power density range measured. For this reason, Cr$^{4+}$:YAG is commonly used as a saturable absorber for Nd:YAG and Nd:YVO$_4$ lasers [140]. The 2 cm crystal absorbed 90% of the pump for low power densities, corresponding to a 1.2 cm$^{-1}$ absorption coefficient. The 1 cm crystal absorbed ~ 80% of the pump for low power densities, corresponding to a 1.6 cm$^{-1}$ absorption coefficient. At the pump condition, however, the 1 cm crystal absorption was only ~ 55% ($\alpha = 0.8$ cm$^{-1}$) while the 2 cm crystal absorption was ~ 80% ($\alpha = 0.8$ cm$^{-1}$). In general, the dependence of the Cr$^{4+}$:YAG performance on the Cr doping level is quite complicated, and it is unclear what is the best way to produce high quality short crystals [141].

After the pump absorption efficiency was measured, both crystals were placed in a laser cavity to measure the output power as a function of the input pumping power. Different output couplers were used within the laser cavity for these measurements. A plot of the data measured for the 1 and 2 cm long laser crystal is shown in Figure 3.5. The absorbed pump power was determined by measuring the difference between the input pump power and the transmitted pump power through the crystal. For each crystal and output coupler combination, the threshold pump power and slope efficiency was determined. The measured values are compiled in Table 3.1. The 2 cm crystal outperformed the 1 cm crystal in wallplug efficiency because of its higher pump absorption fraction, and lower pump threshold values.
Figure 3.5 Measurements of the output power of a Cr\textsuperscript{4+}:YAG laser as a function of input pump power absorbed by the gain media for combinations of a 1 cm or 2 cm crystal and several output couplers.

<table>
<thead>
<tr>
<th>Crystal Length (cm)</th>
<th>Output Coupler</th>
<th>Transmission at Lasing Wavelength (%)</th>
<th>Threshold Power (W)</th>
<th>Slope Efficiency (%)</th>
<th>Wavelength (nm)</th>
</tr>
</thead>
<tbody>
<tr>
<td>2</td>
<td>OC1</td>
<td>1.0</td>
<td>1.15</td>
<td>10.6</td>
<td>1532</td>
</tr>
<tr>
<td>2</td>
<td>OC2</td>
<td>0.7</td>
<td>0.92</td>
<td>10.3</td>
<td>1500</td>
</tr>
<tr>
<td>2</td>
<td>OC3</td>
<td>0.8</td>
<td>0.71</td>
<td>12.4</td>
<td>1475</td>
</tr>
<tr>
<td>2</td>
<td>OC4</td>
<td>2.9</td>
<td>1.50</td>
<td>16.9</td>
<td>1447</td>
</tr>
<tr>
<td>1</td>
<td>OC1</td>
<td><strong>No Lasing Possible</strong></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>1</td>
<td>OC2</td>
<td>0.7</td>
<td>3.87</td>
<td>16.8</td>
<td>1500</td>
</tr>
<tr>
<td>1</td>
<td>OC3</td>
<td>0.9</td>
<td>1.82</td>
<td>10.3</td>
<td>1459</td>
</tr>
<tr>
<td>1</td>
<td>OC4</td>
<td><strong>No Lasing Possible</strong></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>1</td>
<td>OC5</td>
<td>0.6</td>
<td>2.25</td>
<td>7.7</td>
<td>1430</td>
</tr>
</tbody>
</table>

Table 3.1 Tabulation of data from pump power, output power plots.
A slight decay of the output power of the 1-cm crystal-based laser over time was observed. This effect was on the time scale of minutes and was not observed in the 2 cm crystal. This effect has been previously observed highly-doped Cr\(^{4+}\):YAG crystals and has been attributed to conversion of Cr\(^{4+}\) ions into inactive Cr\(^{3+}\) ions [142]. A plot of the output power of the laser against time after turning on the pump laser is shown below in Figure 3.6. This reduction in output power is reversible by leaving the crystal unpumped for a similar amount of time as the power reduction time scale. Because of the high level of pump absorption saturation and time decay present in the 1 cm crystal, the 2 cm crystal was chosen for most additional experiments.

3.4 Tunable Cr\(^{4+}\):YAG

The spectrum of ultrashort optical pulses contains a broad bandwidth which is related to the temporal pulsewidth by the time-bandwidth product. An ideal ultrafast laser media should be able to provide gain over as broad as possible a spectral range in order to support
the shortest possible pulses. To confirm that the Cr\(^{4+}:\)YAG crystals and laser cavity provide broad enough gain to support short pulses, a tunable cw laser was built and characterized. In addition to ultrashort pulse generation, broadly tunable lasers are useful for measuring transmission spectra through materials or devices. The tunable Cr\(^{4+}:\)YAG's wavelength range, near 1500 nm, is particularly useful for optical characterization of telecommunication devices.

Without wavelength selective elements, lasers will operate at the spectral peak of their gain curve. This gain curve is determined by a combination of the laser crystal gain and the laser cavity mirror bandwidths. To tune a laser with broadband gain, a wavelength-selective element should be placed within the cavity. Ideal tuning elements insert loss for all unwanted wavelengths, thereby changing the gain curve of the laser and forcing the laser to operate at a new wavelength. The new lasing wavelength can be changed by tuning the loss-minimum to a new wavelength. Three techniques of creating a wavelength selective element to tune a Z-fold cavity with broadband gain are shown schematically in Figure 3.7. In each case, the basic

![Figure 3.7](image)

Three methods of tuning a Z-fold cavity laser with broadband gain using (a) a birefringent tuning filter, (b) a single prism and a rotating end mirror, and (c) two prism and a translating slit.
cavity design is the same Z-fold geometry discussed for the cw laser, with the addition of the wavelength-dependent element for tunability. In the first technique (Figure 3.7 (a)) a birefringent tuning plate is inserted within the laser cavity at Brewster’s angle [143]. The birefringent axis of the plate material will be at some angle to the laser beam polarization. The birefringent plate is designed to rotate the polarization of different wavelengths by different magnitudes. Those wavelengths rotated by $2\pi$ from p-polarization (the polarization with zero reflection at Brewster’s angle) back to p-polarization will propagation without loss, while wavelengths rotated to polarizations other than p will suffer from loss by Brewster’s angle Fresnel reflections off of the laser crystal and birefringent plate. By rotating the plate, it is possible to tune the wavelength which has its polarization rotated by $2\pi$. Birefringent tuning plates are limited by their free spectral range (FSR) and bandwidth. This limitation can be overcome by combining a large FSR and small bandwidth birefringent plate. In the Cr$^{4+}$:YAG laser, insertion of a birefringent tuning plate always resulted in a decrease in output power by a factor of two because the small clear-aperture of the birefringent filter rotation stages used made it impossible to place the filter at Brewster’s angle.

A second technique (Figure 3.7 (b)) for tuning a broadband laser uses a single prism to angularly disperse wavelengths within the laser cavity. A rotating end mirror selects wavelength by back-reflecting a single frequency, while reflecting unwanted frequencies out of the laser cavity. This approach requires adjustment of the cavity high-reflector and precise alignment of the intracavity prism, but adds minimal insertion loss. There is no limiting FSR, and the bandwidth is determined by the separation between the dispersive prism and end mirror. A third approach (Figure 3.7 (c)) employs two intracavity prisms. The first prism angularly disperses different wavelengths, while a second prism collimates the wavelengths to parallel but spatially separate paths. A slit can then be placed in the beam path after the two prisms, and translated orthogonally to select a specific frequency. The two prism approach has no inherent FSR, and the bandwidth is determined by the slit width.

An example of a Cr$^{4+}$:YAG laser tuning curve using a birefringent filter is shown in Figure 3.8. The laser cavity was the same Z-fold cavity as described above, and the 2 cm crystal was used as the gain media. The laser was tunable from 1390 to 1555 nm, with a maximum power of 500 mW. It is likely, in this case, that the tuning range was limited by the bandwidth of the highly-reflecting cavity fold mirrors on the long wavelength side and the output coupler’s bandwidth on the short wavelength side.
3.5 \textit{Cr}^{4+}:YAG Dispersion

Gain over a large bandwidth is essential for the proper performance of an ultrafast laser. The shorter in time the width of an optical pulse is, the broader the corresponding spectrum will be. If, however, each wavelength travels with a different group velocity around the laser cavity, the pulse will broaden in time. The steady-state pulsewidth will be determined by the interplay between pulse shortening mechanisms such as KLM and pulse broadening mechanisms such as dispersion.
The GDDs from 4 cm of Cr$^{4+}$:YAG (a roundtrip pass on a 2 cm crystal) and from undoped YAG crystal are shown below in Figure 3.9. The Cr$^{4+}$:YAG dispersion was measured using a white light interferometer technique by Ishida et. al. [144] to be:

$$GDD = -15296 + 119.83\nu - 0.20541\nu^2$$  \hspace{1cm} (3.6)

where $\nu$ is in THz, while the YAG dispersion was calculated using the YAG Sellmeier coefficients [40]. The Cr$^{4+}$:YAG GDD is very close to that of undoped YAG and is positive (normal GDD) for wavelengths smaller than 1580 nm. To generate short optical pulsewidths, the crystal GDD must be compensated by anomalous dispersion. Anomalous dispersion can be added to the laser cavity by prisms or chirped mirrors.
3.6 Dispersion Compensation with Prisms

Prism pairs are commonly used to compensate dispersion within ultrafast laser cavities. For laser crystals with normal dispersion, anomalous GDD must be introduced into the cavity, while limiting the amount of additional higher-order dispersion. Anomalous GDD can be added to a laser cavity using either normal or anomalous dispersive prism materials. The geometric dispersion from the prism sequence will be anomalous regardless of the sign of the material dispersion. Conversely, only normally dispersive prism materials can be inserted to compensate laser cavities with excess anomalous dispersion.

A standard geometry used to insert dispersion-compensating prisms in a Z-fold cavity is shown schematically below in Figure 3.10 [129]. Two prisms are placed in series; the first prism angularly disperses different wavelengths, while the second prism collimates the now angularly dispersed beam into a laterally dispersed beam. Each prism is designed with a Brewster cut apex angle. Brewster cut prisms are designed such that light entering one prism face at Brewster's angle will exit the second face at Brewster's angle. This design, which minimizes insertion loss from the prism, is the same criteria used to design a Brewster-Brewster cut laser rods, except that the second face is inverted rather than parallel to the first face. The apex angle of a Brewster cut prism should be:

\[
\theta_{\text{apex}} = 2\left(\frac{\pi}{2} - \theta_B\right)
\]  

\hspace{1cm} (3.7)
where $\theta_{\text{apex}}$ is the angle of incidence relative to the normal and $\theta_B$ is Brewster's angle for the prism material. Dispersion will be introduced by the propagation through a dispersive material, as well as from the geometry of the two prisms. For prisms inserted in the geometry shown in Figure 3.10, with a material propagation distance of $L = L_1 + L_2$ and an apex to apex prism separation of $l$, the GDD introduced will be approximately [145]:

$$GDD = \frac{\lambda_i^3}{2\pi c^2} \left[ L \frac{d^2 n}{d\lambda^2} - 4l \left( \frac{dn}{d\lambda} \right)^2 \right] \quad (3.8)$$

where $\lambda_i$ is the carrier wavelength, $c$ is the speed of light, and $n$ is the index of refraction of the prism material.

Choosing the proper prism material is critical for the construction of an optimized ultrafast laser cavity. The dispersion of several prism materials are shown below in Figure 3.11. The GDD in this plot has been calculated for propagation through 1 mm of material.

![Figure 3.11](image)

**Figure 3.11** The group delay dispersion (GDD) of several potential prism materials.

Because of the finite width of the intracavity beam and beveled prism edges, it is difficult to align a laser beam through a prism in order to have a pathlength shorter than 2.5 mm within
the prism material. Therefore two intracavity prisms, each encountered twice per cavity roundtrip, create a lower bound for the material dispersion which corresponds with propagation through at least 10 mm of prism material. Similarly, due to the finite size of prisms, it is difficult align an intracavity beam to pass through more than 15 mm of prism material per prism, or 60 mm per roundtrip. These bounds limit the possible choices of prism materials.

An additional constraint in choosing a proper prism material choice is the presence of water impurities within many glasses commonly used for prisms. This water will cause loss within the laser cavity. Because intracavity powers are often large, water absorption can also lead to local heating within the glass. This effect is exacerbated because most glasses have poor thermal conductivities and cannot conduct away generated heat. There are a series of particularly large water absorption lines in the spectral region between 1300 to 1500 nm. This wavelength range corresponds almost exactly with the Cu^{4+}YAG gain profile. Interestingly, this exact problem has, until recently, limited the usable bandwidth in glass optical fibers for communications. Most fiber optic communications systems transmitted data on wavelengths near 1300 nm (below the water absorption lines in glass fiber) or near 1550 nm (above the water absorption lines). Recently, Lucent Technologies introduced All-Wave fiber, a fiber specially fabricated to overcome this problem by eliminating water impurities. A plot of the transmission through standard single-mode fiber and through Lucent’s All-Wave fiber is shown in Figure 3.12. Some bulk materials with low water content are commercially avail-

![Graph showing optical bandwidth of All-Wave fiber compared to standard fiber.](image)

**Figure 3.12** Transmission through Lucent’s All-Wave fiber (with low water content) and standard single mode fiber (with higher water content) (*from H. Kogelnik*).
able for prisms. The dispersion of a few such materials are shown below in Figure 3.13. Most

glasses have high water contents because the glass melts are heated in a hydrogen furnace. One type of fused silica material called Infrasil, however, is available with low water content. Although glasses are particularly susceptible to water impurities because of their amorphous structure, crystalline materials such as CaF$_2$, BaF$_2$, and YAG naturally exclude water impurities from their regular lattice structures.

Short pulses from Cr$^{4+}$:YAG lasers using fused silica prisms for dispersion compensation have been previously generated by several groups. Tong et. al. was able to produce pulses as short as 43 fs using fused silica prisms [115]. Ishida et. al. used a Brewster-Brewster cut fused silica slab, to reduce the difficulty in alignment, and were able to achieve pulses as short as 60 fs [144]. Fused silica has an anomalous group delay dispersion (GDD) of -32 fs$^2$/mm at 1500 nm and can therefore be used to compensate the normal GDD of a 2 cm Cr$^{4+}$:YAG crystal (11 fs$^2$/mm at 1500 nm). Because fused silica material is anomalous, the Cr$^{4+}$:YAG dispersion can be compensated with either the insertion of material or by the prism separation within the laser cavity. Because the third-order dispersion (TOD) introduced by
adding more prism material is lower than that introduced by increasing the prism separation, it is preferable to minimize the prism separation and rely primarily upon prism insertion. A plot of the GDD for combinations of Cr$^{4+}$:YAG and fused silica material of varying thickness is shown below in Figure 3.14. It is evident that a large TOD (visible as a non-zero linear slope

![Graph showing GDD vs. Wavelength for Cr:YAG and Fused Silica Prism Material](image)

**Figure 3.14** Group delay dispersion (GDD) of Cr$^{4+}$:YAG and fused silica glass of different thicknesses.

of the GDD plot) remains when fused silica is used to compensate the Cr$^{4+}$:YAG GDD. In order to place the zero GDD point near the laser gain, an insertion of somewhere between 10 to 30 mm of fused silica is required.

### 3.7 Prism Compensation Results

As a first step toward producing ultrashort pulses from a Cr$^{4+}$:YAG laser, modelocking was attempted with two infrasil fused silica prisms in a cavity similar to those reported in
previous work [115]. A schematic of the laser cavity is shown below in Figure 3.15. The laser crystal used to provide gain is the 2 cm long, 3 mm diameter Brewster-Brewster cut rod, supplied by A. V. Shestakov at E.L.S. Company. Pump light at 1064 nm from a Spectra-Physics 11 W Nd:YVO₄ laser is focused by a 10 cm focal length lens into the crystal. The laser crystal has a linear absorption of 1.2 cm⁻¹ and is cooled to 13 C during operation. Surrounding the laser crystal are two 10 cm radius-of-curvature mirrors (M1 and M2), rotated 16 degrees from normal incidence for astigmatism compensation. The mirrors, all high reflector type HR1, are highly reflective from 1275 to 1625 nm at normal incidence. Because the cavity fold mirrors are placed at an angle, the reflectivity curve is shifted to shorter wavelengths by an amount given by \( \lambda' = \lambda \cos \theta \). For 16 degrees, a mirror with reflectivity centered at 1500 nm will have its center wavelength shifted by 60 nm. One arm of the laser, 50 cm long, contains an Infrasil prism pair as well as a highly reflecting end mirror (M3). The second arm of the laser, 40 cm long, contains an output coupler (OC1). The output coupler has a minimum transmission of 0.5% at 1585 nm and less than 1.0% transmission between 1530 to 1650 nm. In this particular configuration, the laser cavity was not purged with nitrogen gas to eliminate water absorption, and the output coupler forced all lasing to occur above 1500 nm.

Fused silica prisms were inserted to compensate the GDD of the Cr⁴⁺:YAG gain crystal. The normal GDD of Cr⁴⁺:YAG (11 fs²/mm at 1500 nm) can be compensated by either the insertion of 14 mm of fused silica material, a separation of 65 mm between two fused silica prisms, or by some combination of the material and geometric dispersion. The second prism and high reflector were placed on a sliding translation stage. Modelocking could be initiated

![Figure 3.15 Schematic of the Z-fold cavity with a 2 cm Cr⁴⁺:YAG crystal and two fused silica prisms. Mirrors M1 - M3 are all high-reflective mirrors and OC is an output coupling mirror.](image-url)
by sliding the translation stage and the pulsewidth could be minimized in real time by sliding the same stage slowly.

A typical modelocked pulse spectrum, measured by an optical spectrum analyzer, is shown below in Figure 3.16. The light emission is centered at 1537 nm, with a full-width half maximum of 46 nm. The large shoulder visible on the short wavelength side of the spectrum is caused by the roll-off of the output coupler reflectivity. Long wavelength spikes at 1600 and 1633 nm have been observed in previous work [115] and have been attributed to dispersive solitary waves shed by periodic perturbations within the laser cavity. The sidebands are visible on the long wavelength side of the spectrum because of the large negative GDD for those wavelengths.

The pulsewidth was measured by a noncolinear intensity autocorrelator. The beam was split by a dielectric beamsplitter into two paths. One path consisted of a retroreflector mounted on a speaker for real-time time delay, while a retroreflector in the second arm was mounted on a 1 μm step-size motorized translation stage for calibrated time delay. The beams from each arm are directed parallel to each other with a mirror that lets one arm pass and

Figure 3.16 Typical modelocked spectrum from Cr\(^{4+}\)-YAG laser using fused silica prisms for dispersion compensation.
reflects the second. The parallel beams are then focused through a nonlinear LiIO$_2$ crystal with a 5.8 cm focal length lens. The second harmonic light generated by the two beams was collected through an adjustable iris which blocked uncorrelated simple harmonic generation (SHG) and passed correlated sum frequency generation to measure a background free signal. The correlated sum frequency light was then collected and recorded by a biased photo-multiplier tube (PMT). An autocorrelation trace from the fused silica prism laser cavity is shown in Figure 3.17. The pulse width was measured to be 54 fs, assuming a sech-shaped pulse enve-

![Graph of autocorrelation measurement](image)

**Figure 3.17** Autocorrelation measurement (black line) and sech fit (gray dashed line) of 54 fs pulses from a fused silica prism dispersion compensated Cr$^{4+}$:YAG laser.

lope. Assuming a sech-shaped pulse envelope fits the center of the autocorrelation well, but significant sidelobes are present. Many of these sidelobes were removed from later traces by replacing a particular dielectric stack mirror in the beampath before the autocorrelator. The pulsewidth was measured directly from the laser, and no external compensation of chirp was performed. Assuming a spectral fwhm of 46 nm and a pulse width of 54 fs, a time-bandwidth
product of 0.318 is calculated. This is close to the transform limited value of 0.315 for sech pulses.

It has been reported that the minimum pulse width in the Cr\(^{4+}\):YAG laser compensated by fused silica prisms is limited by large TOD [116]. It is possible to lower the higher-order dispersion introduced from fused silica prisms by using some positive (normal) dispersion prism materials instead. The combination of anomalous dispersion from the prism geometry and the normal material dispersion can be designed to give a net anomalous dispersion with smaller TOD than is possible with fused silica. This approach was attempted with SF58 (PBH71) and SF10 prisms but failed because each material contained large water impurities. Upon insertion inside the laser cavity, significant losses and thermal effects were observed. Even if these materials had low water content and could be used within the laser cavity large amounts of uncompensated TOD would still be present.

### 3.8 Double-Chirped Mirrors

The conventional approach to compensate the intracavity GDD of a gain media in an ultrafast laser is the insertion of either prism pairs or a bulk dispersive material. Chirped mirrors have recently been developed as an alternative method of compensating GDD [130, 131]. Double-chirped mirrors improve the dispersive properties of chirped mirrors by minimizing residual higher-order dispersion introduced by spurious reflections from imperfect anti-reflection [132 - 135]. A schematic of Bragg reflecting quarter-wave stack mirrors, chirped mirrors, and double-chirped mirrors, taken from Natuschek et. al. [134], are shown in Figure 3.18. Standard Bragg dielectric mirrors are optimized for alternating high and low index layers of quarter-wavelength thickness. The bandwidth of these mirrors is determined by the dielectric contrast of the materials and the number of quarter-wave layers used. The dispersion of a typical quarter-wave stack mirror is relatively flat. More detail about how to calculate the reflectivity and GDD of Bragg mirrors is discussed in Chapter 4.
Bragg Mirror: \[\text{SiO}_2 - \text{Substrate} \quad \text{TiO}_2 / \text{SiO}_2 \quad \lambda_b/4 - 	ext{Layers} \quad \text{Air}\]

Chirped Mirror: \[\text{SiO}_2 - \text{Substrate} \quad \text{Bragg-Wavelength} \lambda_B \quad \text{Chirped} \quad \lambda_1 \quad \lambda_2 \quad \text{Negative Dispersion:} \quad \lambda_2 > \lambda_1 \]

Double-Chirped Mirror: \[\text{SiO}_2 - \text{Substrate} \quad \text{Bragg-Wavelength and Coupling Chirped} \quad d_h \neq \lambda_b/4 \quad \text{AR-Coating} \quad \text{Air}\]

Figure 3.18 Schematic of Bragg reflecting quarter wave stacks, chirped mirrors, and double-chirped mirrors (from reference[134]).

Chirped mirrors change the periodicity of the quarter-wave stack as a function of depth into the mirror stack. The mirror will act as a Bragg reflecting mirror for wavelengths at different depths into the mirror. Each wavelength will therefore travel a different optical path length into the structure before undergoing reflection, resulting in dispersion. In principle, the mirror can be chirped over a large range of Bragg periodicities, and can be designed to provide exceptionally broadband reflection. Imperfect anti-reflection from the chirped mirror’s front surface can limit dispersion performance. Light reflecting off this front surface will not travel the same optical path length as the remaining light which travels through the mirror stack. The interference between these two components will create spectral oscillations in the chirped mirror dispersion.

Double-chirped mirrors (DCMs) reduce the detrimental effects of residual reflection off the front surface of the chirped mirror. In addition to chirping the Bragg periodicity of the mirror, DCMs are designed to impedance-match the incoming and transmitted waves at each mirror interface. This essentially combines the broadband reflection and dispersion control of chirped mirrors with a broadband anti-reflection coating. These mirrors are typically designed by a complicated computer optimization technique and are made from relatively high dielectric contrast materials for broadband performance [132]. DCMs are extremely dif-
difficult to properly design and fabricate. The design involves complicated computer optimization algorithms to develop a suitable design. This work was enabled by the design of Cr$^{4+}$:YAG mirrors by Prof. Franz Kaertner.

DCMs, comprised of 48 layers of SiO$_2$ (n = 1.5) and TiO$_2$ (n = 2) were designed to compensate the dispersion of a Cr$^{4+}$:YAG crystal and grown by a state-of-the-art ion-beam sputtering technique [146]. A schematic of the DCM layer structure is shown below in Figure 3.19. In the final optimized design, it is difficult to discern the quarter-wave stack nature of the mirror. The DCM is designed to have a broadband reflection in the region of Cr$^{4+}$:YAG gain and high transmission for the pump laser at 1064 nm. Once fabricated, the mirror reflectivity was measured using a CARY spectrophotometer and is shown in Figure 3.20. The DCM has a high reflection from 1200 to 1700 nm, and has a highly transmissive region from 1000 to 1120 nm for transmission of the pump beam at 1064 nm.

Figure 3.19 Index of refraction of the layers comprising the double-chirped mirrors (DCMs) designed for use in a Cr$^{4+}$:YAG laser.
Figure 3.20  Measured reflectivity of double-chirped mirror (DCM) designed for use within a Cr$^{4+}$:YAG laser.

The GDD was measured by a white light interferometer [147] and is shown below in Figure 3.21 along with the designed GDD. The measured GDD is quite close to the design goal and displays a broad range of flat GDD from 1350 to 1600 nm. In this range, the TOD is roughly constant and has the opposite sign from the Cr$^{4+}$:YAG material. A small oscillation in the DCM GDD is present and is due to imperfect anti-reflection properties of the mirror.
Figure 3.21  Designed (gray line) and measured (black line) group delay dispersion (GDD) of the double-chirped mirrors (DCMs) used in a Cr\textsuperscript{4+}:YAG laser.

The DCMs were originally designed to be used with a 1 cm Cr\textsuperscript{4+}:YAG crystal in combination with SF58 prisms for fine tuning. After the mirrors had already been fabricated, it was determined that the 2 cm laser crystal is more suitable for use because of its high gain, and SF58 prisms contained too much water for use in the laser cavity. Despite the improper design, it was still possible to use these mirrors for the generation of short pulses.
Figure 3.22 Net cavity group delay dispersion (GDD) of laser cavities consisting of a number of double-chirped mirror (DCM) reflections and either (a) 1 cm or (b) 2 cm Cr\textsuperscript{4+}:YAG crystal.
The net cavity GDD of a laser cavity consisting of either a 1 or 2 cm long Cr\(^{4+}\):YAG crystal and a number of reflections off DCMs are shown in Figure 3.22. The DCMs were designed for use with 4 DCM reflections with the 1 cm crystal. The remainder of the GDD would have been compensated by the prisms. Instead, the 2 cm crystal has been used in most laser cavities. The combination of the 2 cm laser crystal and 6 DCM reflections is particularly suitable for a laser cavity with no prisms. For both 1 and 2 cm crystals, an increase in the number of reflections off DCMs results in a corresponding increase in the amplitude of the dispersion oscillations. In all plots, the GDD of DCMs for normal incidence reflection have been used.

Some DCM mirrors within a Z-fold laser cavity must be placed at off-normal angles of incidence. In particular, the cavity fold mirrors surrounding the laser crystal must be placed at a specific angle to compensate for the astigmatism introduced by the Brewster cut laser crystal. A plot of the DCM mirror GDD for several angles of incidence is shown below in Figure 3.23. There is no major difference between the GDD for normal incidence and 5 degree inci-

![Graph showing GDD as a function of wavelength for different angles of incidence.](image)

Figure 3.23 Group delay dispersion (GDD) of double-chirped mirrors for off-axis reflection. The GDD oscillations grow as a function of angle.
dence. For 11 degrees, the astigmatism compensation angle for 1 cm crystals, and 16 degrees, the astigmatism compensation angle for 2 cm crystals, a significant increase in the GDD oscillations is apparent. GDD oscillations increase because the anti-reflection properties of the DCM mirrors are no longer effective for large angles of incidence. In Ti:Sapphire laser, it has been shown that a quarter-wave layer can π-shift the GDD oscillations [135]. DCM pairs can then be used to lower the amplitude of the net cavity GDD oscillations.

3.9 Water Absorption

To date, most efforts at modelocking the Cr\textsuperscript{4+}:YAG laser have focused around generating optical pulses with center wavelengths greater than 1500 nm. This spectral region is interesting because it is currently used for long distance optical fiber communication systems. A significant portion of the Cr\textsuperscript{4+}:YAG gain, however, is at wavelengths below 1500 nm. In order to produce the shortest possible pulses, all available spectrum, both above and below 1500 nm, should be utilized. Furthermore, as mentioned previously, the recent removal of water absorption peaks between 1300 and 1500 nm from optical fibers has opened up additional wavelengths for use in fiber-optic communication systems. An ideal ultrafast source for the study of telecommunication devices or materials should be able to produce pulses with spectrum covering this entire wavelength range.

Water absorption lines add loss and dispersion to air-filled laser cavities. To generate short pulses in the wavelength range of the absorption, the laser cavity must be purged free of water vapor, and optical materials with low water contents must be used in the cavity. Air, for example, contains water vapor and other molecules with significant absorption lines between 1300 to 1500 nm. An example of the transmission spectra of air is shown in Figure 3.24 [148]. The major water absorption lines are due to OH\textsuperscript{-} radicals in water vapor.

Absorption from water vapor in air completely prevents KLM below 1500 nm unless the optical cavity is purged free of water vapor. The water vapor limits KLM primarily through the dispersion and intracavity loss introduced by the water lines. A plot of the GDD introduced by water vapor in air, taken from Conlon et. al. [149], is shown in Figure 3.25. In this experiment, Conlon et. al. measured the dispersion of their laser cavity using a white light interferometer technique [149]. Note in particular that the scale of the GDD due to water absorption is much larger than the GDD scale of the laser crystal, prisms, or DCMs.
Figure 3.24 Transmission spectra through air exhibiting large water absorption lines between 1300 and 1500 nm. (From reference [148])

Figure 3.25 The group delay dispersion (GDD) from water vapor in a Cr^4+:YAG laser. (From reference [149])
To demonstrate the effect of intracavity loss due to water absorption, the cw Cr\textsuperscript{4+}:YAG laser was tuned to 1405.6 nm, a wavelength with large water vapor absorption. The entire cavity was enclosed within a plastic box and filled with dry nitrogen gas slowly over time. The output power of the laser was recorded as the box filled with nitrogen and is shown below in Figure 3.26. With the purging box open to atmosphere, the laser had an output power of 2 mW. After 45 minutes of slow N\textsubscript{2} purging, the output power increased over 30 mW.

The effect of water vapor on modelocking was investigated by looking at KLM cavities as well as by using semiconductor Bragg reflector (SBR) saturable absorber mirrors to initiate saturable absorber modelocking. Without purging, KLM could not be initiated for wavelengths smaller than 1500 nm. By choosing an SBR with high reflectivity for wavelengths smaller than 1500 nm, it was possible to initiate saturable absorber modelocking. A plot of the purged and unpurged spectrum for a saturable absorber modelocked Cr\textsuperscript{4+}:YAG laser was shown in Figure 3.27.
laser is shown below in Figure 3.27. The saturable absorber used is described in the PhD the-

![Graph showing intensity vs wavelength for No Purging and Purged with N\textsubscript{2}](image)

Figure 3.27 Purged (gray line) and unpurged (black line) spectrum of saturable absorber modelocked Cr\textsuperscript{4+}:YAG laser.

sis of Schibli [150], and consists of a 4 nm InGaAs quantum well layer centered in a half-
wave thick InP layer deposited upon a 30x GaAs/AlAs Bragg mirror stack. The saturable absorber mirror self-started modelocking at 1460 nm with and without N\textsubscript{2} purging. When purging with N\textsubscript{2} gas, however, the spectrum becomes much smoother and more uniform.

### 3.10 Prismless Cavity

Using the design constraints discussed above, an ultrafast laser cavity was built using DCMs to minimize higher-order dispersion and generate ultrashort pulses [139]. Because of the lack of available prism materials with both the proper dispersion characteristics and low water content, a prismless laser cavity using only DCMs for dispersion compensation was constructed. Prismless laser cavities are highly desirable because of their easy alignment and
potential to be shortened to generate higher repetition rate pulse trains. Prismless cavities avoid both the difficult alignment and large insertion loss of intracavity prisms. Cr\(^{4+}\):YAG lasers, in particular, are sensitive to such loss; the laser based on the 2 cm crystal can only withstand about 3% loss per cavity roundtrip. Removing prisms, however, sacrifices the ability to fine-tune the intracavity dispersion. The steady-state pulsewidth is highly sensitive to net cavity dispersion. It is likely that the pulsewidth in a prismless laser will ultimately be limited by this lack of tunable dispersion. Prismless lasers have been used to generate short pulses in several laser systems, including 14 fs pulses at 900 nm from Cr:LiSGaF [151] and 11 fs pulses from Ti:Sapphire [152]. In the case of Ti:Sapphire, prismatic wedges (prisms with very small apex angles) were used for fine tuning of dispersion. The dispersion was varied by changing the wedge insertion. Because of the apex angle, however, it is not possible to align both wedge faces at Brewster’s angle. Wedge pairs of YAG with an angle of 2 degrees already add ~ 2% loss within a laser cavity, which is too high for Cr\(^{4+}\):YAG.

A schematic of the Z-fold laser cavity is shown in Figure 3.28. The 2 cm long 3 mm

![Figure 3.28](image)

Figure 3.28  Schematic of a prismless Cr\(^{4+}\):YAG laser used to generate short pulses. The cavity consists of a 2 cm laser crystal, an output coupler (OC), 3 double-chirped mirrors (DCMs) adding 6 DCM bounces per cavity roundtrip (M1-M3, in black), and one unchirped high reflector (M4, in gray).

diameter Brewster-Brewster cut rod, supplied by A. V. Shestakov at E.L.S. Company, was chosen because of its high gain. Pump light at 1064 nm from a Spectra-Physics 11 W Nd:YVO\(_4\) laser is focused by a 10 cm focal length lens into the crystal. The laser crystal has a linear absorption of 1.2 cm\(^{-1}\) and is cooled to 13 C during operation. Surrounding the laser crystal are two 10 cm radius-of-curvature DCMs (M1 and M2), rotated 16 degrees from normal incidence for astigmatism compensation. One arm of the laser, 70 cm long, contains both
an additional DCM (M3) and an unchirped quarter-wave stack high-reflecter HR1 (M4). The second arm of the laser, 50 cm long, contains the broadband output coupler OC3 which has a minimum transmission of 0.7% at 1515 nm and less than 1.4% transmission between 1420 to 1630 nm. The lasing threshold was at 0.7 W of absorbed pump power and the cw output power was 1 W for 9 W of absorbed pump. The laser beam path was enclosed in plastic tubes and wraps and purged with dry nitrogen gas. A picture of the laser cavity is shown below in Figure 3.29.

Figure 3.29 Picture of the prismless Cr\textsuperscript{4+}:YAG laser.

The laser cavity was designed to have 6 DCM reflections each cavity roundtrip. As can be seen in Figure 3.22, this leads to a flat, broadband, and slightly negative dispersion. A slightly negative (anomalous) GDD is desirable to minimize pulsewidth because it allows to soliton-shaping of the optical pulse. The GDD of Cr\textsuperscript{4+}:YAG [144], of 6 DCM reflections at normal incidence, and of the full laser cavity are shown in Figure 3.30. The 2 cm (4 cm roundtrip) Cr\textsuperscript{4+}:YAG crystal has a GDD of 450 fs\textsuperscript{2} and TOD of 6000 fs\textsuperscript{3} at 1500 nm and zero GDD at 1590 nm. The DCM GDD was measured by a white light interferometer technique [147] and has a roughly linear anomalous dispersion from 1325 to 1575 nm, a GDD value of approximately -125 fs\textsuperscript{2} and a TOD of ~ -1000 fs\textsuperscript{3} at 1500 nm per reflection. A dispersion compensated cavity was created by using 6 DCM reflections in each cavity roundtrip to compensate for the Cr\textsuperscript{4+}:YAG crystal. Effects of the broadband output coupler, unchirped high reflector, and off-normal incidence reflections from DCMs are all included in the net cavity dispersion in Figure 3.30. Oscillations in the GDD of the laser cavity are due to residual spurious reflections caused by imperfect impedance matching of the DCM. Because the DCMs
were designed for normal incidence, these GDD oscillations were enhanced due to the large angle of incidence.

![Graph showing GDD vs Wavelength](image)

Figure 3.30  Group delay dispersion (GDD) of 2 passes through a 2 cm Cr\textsuperscript{4+}:YAG crystal (light gray), 6 reflections from double-chirped mirrors (DCMs) (gray), and all optical elements in the laser cavity (black). The net cavity GDD curve includes reflections from 6 DCMs, a single unchirped quarter-wave stack high reflector, and an output coupler.

The laser was aligned for modelocked by first optimizing cw power. The curved mirror separation and crystal position were then varied until KLM was observed. Modelocking was initiated by tapping one of the end mirrors. The average power of the 110 MHz pulse train was 200 to 400 mW, depending on the alignment, for 9 W of absorbed pump. Water vapor in air introduces intracavity GDD and loss through a series of absorption lines between 1300 and 1500 nm [149]. To remove this water vapor, the optical path was enclosed and purged with dry nitrogen gas. An example of the modelocked pulse spectrum, measured with a calibrated optical spectrum analyzer, is shown on both linear and log scale in Figure 3.31. Significant spectrum is present above the noise floor within the wavelength range of 1140 to
> 1700 nm (the optical spectrum analyzer used has a long wavelength limit of 1700 nm) and has a full width at half maximum of 190 nm, from 1310 to 1500 nm. The spectrum is smooth, with a relatively flat top from 1340 to 1470 nm. The output coupler, which rolls off significantly at wavelengths smaller than 1350 nm, enhances the output spectrum at shorter wavelengths. Long wavelength spectral spikes have been observed previously and attributed to Raman amplified Kelly sidebands [153].

![Optical power spectrum of a Cr\textsuperscript{4+}:YAG pulse. The black line corresponds to a linear scale (left axis) and the gray line corresponds to a logarithmic scale (right axis). The full-width at half maximum is 190 nm, with a peak at 1450 nm.](image_url)

Figure 3.31  Optical power spectrum of a Cr\textsuperscript{4+}:YAG pulse. The black line corresponds to a linear scale (left axis) and the gray line corresponds to a logarithmic scale (right axis). The full-width at half maximum is 190 nm, with a peak at 1450 nm.

The pulsewidth was measured by a fringe-resolved autocorrelator. The beam was split and recombinated by broadband metallic beamsplitters and focused by an off-axis parabolic mirror onto a silicon p-i-n photodiode to perform the autocorrelation function by two-photon absorption. An autocorrelation trace is shown in Figure 3.32. A good 8:1 ratio is visible between the peak of the autocorrelation and the background level. Sidelobes are visible to either side of the central lobe. The periodicity of the sidelobes corresponds to the double peaked flat-top nature of the spectrum shown in Figure 3.31. There are 7 fringes within the
full width at half maximum of the autocorrelation trace, which correspond to a 4-cycle optical pulse.

![Autocorrelation graph](image)

Figure 3.32 Measured autocorrelation function from an interferometric two-photon absorption autocorrelator (line) and fit by a pulse-retrieval algorithm (dots). A pulse width of 19.5 fs is calculated by the pulse-retrieval algorithm, 18.3 fs by assuming sech shaped pulses, and 17.0 fs by assuming gaussian shaped pulses.

The abscissa was calibrated with the interference fringes of a HeNe laser directed through the autocorrelator. The HeNe beam was detected by single photon absorption of the Si p-i-n photodiode. A plot of interference fringes measured through the autocorrelator interferometer is shown in Figure 3.33. The peaks of the interference fringes are found and calibrated to the known wavelength of HeNe lasers. Each fringe separation is found to correspond to a time delay given by \( \Delta \tau = \frac{\lambda_{\text{HeNe}}}{c} \) where \( \lambda_{\text{HeNe}} \) is 632.8 nm. Once the corresponding time delay for each fringe is determined, the time delay can be plotted against the digital scope’s internal units.
Figure 3.33  HeNe laser interference fringes measured by a Si p-i-n photodiode through the autocorrelator interferometer. The peak of each interference fringe is marked by a black circle.

Once the corresponding time delay for each fringe is determined, the time delay can be plotted against the digital scope’s internal units. A typical plot is shown in Figure 3.34. For a translation stage moving linearly with time, the mapping from scope units to fs would be linear. A speaker, however, is likely to oscillate sinusoidally. The data fits well to a third-order polynomial fit, which is then used to calibrate the x-axis.
Figure 3.34  Plot mapping the time delay introduced by a speaker in the autocorrelator from scope units to fs. The data was fit by a third-order polynomial.

Once the x-axis is calibrated, it is possible to fit the autocorrelation of Figure 3.32 and determine the pulse width. Several pulse fits were attempted to arrive at the most accurate measurement of pulse width. The pulse width is estimated to be 18.3 fs for sech-shaped pulses and 17.0 fs for gaussian-shaped pulses. A pulse retrieval algorithm [154] was used to reconstruct the pulse profile from the measured interferometric autocorrelation and spectrum. This pulse retrieval algorithm takes the measured pulse spectrum and iterates through different phase profiles. The spectrum is then Fourier transformed using the phase profile to determine the corresponding temporal pulse shape. The autocorrelation function of the pulse shape is fit to the measured autocorrelation. This entire procedure is iterated until the phase profile gives a pulse shape that converges to the measured autocorrelation. The retrieved pulse intensity envelope has a width of 19.5 fs. A Fourier transform of the optical spectrum, assuming a flat phase profile, indicates a bandwidth-limited pulse width of 17.5 fs.
Pulses travel through a few optical elements between the output of the Cr$^{4+}$:YAG laser and the autocorrelator. In addition to a lens used to focus the beam, pulses travel through a 0.25 thick output coupler substrate. Chirp from the dispersion of these extracavity optical elements was compensated by two BaF$_2$ prisms and one DCM bounce. A schematic of the laser, autocorrelator, and chirp compensation prisms is shown in Figure 3.35. The interferometric autocorrelation trace of a chirped pulse changes qualitatively as the BaF$_2$ prisms are inserted and the chirped pulses are compensated. For pulses with large chirps, the wings of the interferometric autocorrelation traces appear to slope upwards. Examples of interferometric autocorrelation traces for different BaF$_2$ prism insertions are shown in Figure 3.36.

Figure 3.35 Schematic of the Cr$^{4+}$:YAG autocorrelation experiment. Output light from the laser is chirp compensated by a pair of BaF$_2$ prisms. The pulsewidth is then measured with a Si p-i-n detector two photon absorption based autocorrelator.
Figure 3.36  Series of interferometric autocorrelation traces taken with different quantities of extracavity group delay dispersion (GDD) compensation. As the chirp of the pulses are compensated, the wings of the autocorrelation traces become flat.

The pulsewidth corresponding to each trace can be determined by fitting the central fringes of the autocorrelation assuming a sech-shaped pulse envelope. A plot of the fitted pulsewidth as a function of prism insertion is shown in Figure 3.37. Without any prism com-
pensation, a pulse width of 20.5 fs was measured. A minimum pulse width of 18.3 fs was measured for 9 mm of BaF$_2$ material, which corresponds to 127 fs$^2$ of GDD at 1500 nm.

![Graph showing pulse width vs. BaF$_2$ GDD](image)

**Figure 3.37** Autocorrelation measured pulse widths, fit assuming a sech pulse shape, for different values of extracavity group delay dispersion (GDD) introduced by two BaF$_2$ prisms and one double-chirped mirror (DCM) bounce.

### 3.11 Different OCs and HRs

The performance of the ultrafast Cr$^{4+}$:YAG laser is highly sensitive to the properties of intracavity mirrors. In particular, a given optical cavity with DCMs for dispersion compensation can produce very different pulses and spectra if different output couplers and unchirped high-reflectors are used. Not only can the bandwidth of these unchirped mirrors limit the minimum achievable pulsewidth, but their dispersion properties can affect pulse generation. Output couplers can also be used to enhance the wings of the optical spectrum.
Several different output couplers and high reflectors (M4 in Figure 3.28) were used in the prismless laser to study their effects on modelocking. The dispersion profiles of these mirrors, measured by a white light interferometer technique [147], are shown in Figure 3.38. The corresponding mirror reflectivities are shown in Figure 3.2. It was possible to modelock the prismless laser using several combinations high reflectors and output couplers.

Figure 3.38 Group delay dispersion (GDD) of output couplers and unchirped high reflector mirrors used within the ultrafast prismless Cr\textsuperscript{4+}:YAG laser.
Three pairs of spectra and interferometric autocorrelations (IACs), each obtained using a different output coupler and high-reflective combination, are shown in Figure 3.39. In each case, all other cavity parameters of the laser are held constant and are the same as those discussed above for the prismless laser. The first spectrum/IAC pair, Figure 3.39 (a) and (b), are obtained by using output coupler OC3 and high-reflective HR1. This mirror combination
yielded the broadest spectrum and the shortest optical pulses. The short wavelength portion of the spectrum is heavily enhanced by the roll-off of the output coupler reflectivity near 1400 nm. Wings in the autocorrelation correspond with the flat-topped structure of the spectrum. The spectrum is centered at ~1400 nm, and the pulses are measured to be 18.3 fs with a sech fit and 19.5 by a more accurate pulse retrieval fit.

The second spectrum/IAC pair, Figure 3.39 (c) and (d), reveals a much different spectrum. High-reflecter HR2 is used and cuts off the lower wavelength end of the spectrum. The spectrum has a peak at 1500 nm and a shoulder on the long wavelength side. The spectrum in this case has a 140 nm fwhm, and the autocorrelation fits to a 22 fs sech pulse. Sidelobes in the autocorrelation function are visible at around +/- 65 fs time delay and correspond with the long wavelength spectral sideband.

Finally, the third spectrum/IAC pair, Figure 3.39 (e) and (f), uses output coupler OC5 with high-reflecter HR1. This combination allows the generation of a very smooth, singly-humped spectrum centered at 1475 nm with a 96 nm fwhm. The pulse autocorrelation fits extremely well to a 19 fs sech.

As a final test, the high reflector was replaced with an additional DCM. This led to a total of 7 roundtrip reflections off DCMs, which has a larger anomalous GDD and a higher third-order dispersion than the laser cavity consisting of 6 DCM reflections. A plot of the spectrum and autocorrelation from the 7 DCM modelocked laser is shown below in Figure 3.40. The spectrum has an unusual shape and several cw sidebands including a particularly

![Figure 3.40](image)

Figure 3.40  (a) Optical spectrum and (b) interferometric autocorrelation of pulses from a prismless Cr\(^{4+}\):YAG laser containing 7 double-chirped mirror (DCM) bounces per cavity roundtrip.

strong line near 1400 nm. These sidebands are likely a result from higher-order dispersion in
the system. The autocorrelation has a short component of ~ 20 fs fwhm in the center but also has large sidebands off-center.

3.12 Reduction of DCM Angle

To explore the possibility that enhanced spectral oscillations in the DCMs due to non-normal incidence on the cavity fold mirrors limit the pulsewidth, a second prismless cavity was designed and tested. A schematic of this laser cavity is shown below in Figure 3.41. The laser crystal was the same 2 cm long 3 mm diameter Brewster-Brewster cut rod. Pump light at 1064 nm from the 11 W Nd:YVO₄ laser is focused by a 10 cm focal length lens into the crystal. Surrounding the laser crystal are two 10 cm radius-of-curvature mirrors (M1 and M2), rotated 16 degrees from normal incidence for astigmatism compensation. In this particular cavity, the cavity fold mirrors were standard quarter-wave stack dielectric mirrors (HR1). Quarter-wave stack mirrors will not develop oscillations in the dispersion as the angle of incidence is increased but will experience a shift of the mirror properties to shorter wavelengths. For 16 degrees, a Bragg mirror centered at 1500 nm will be shifted by about 60 nm. One arm of the laser, 65 cm long, contains three DCMs (M3-M5) and a single unchirped quarter-wave stack high-reflector end mirror (M6). The three DCMs are aligned so that the angle of incidence is as small as possible (approximately 5 degrees) while still being able to clear the mount on the next pass. The second arm of the laser, 50 cm long, contains a broadband output
coupler with a minimum transmission of 0.7% at 1515 nm and less than 1.4% transmission between 1420 to 1630 nm. Like the previous prismless laser discussed, this cavity has a total of 6 DCM reflections for each cavity roundtrip.

As before, the laser was modelocked by slightly tapping end mirrors. An example of the typical modelocked spectrum is shown in Figure 3.42. A smooth singly peaked spectrum centered at 1473 nm with a fwhm of 68 nm is generated. Strong cw sidebands are present on both the long and short wavelength sides of the spectrum. Three of the sidebands have larger spectral intensities than the modelocked spectral envelope. It is likely that the 60 nm shifted dispersion from the 2 unchirped cavity fold mirrors limits the dispersion compensation of this cavity.

The pulsewidth was measured by a fringe-resolved autocorrelation. The beam was split and recombined by broadband metallic beamsplitters and focused by an off-axis parabolic mirror onto a silicon p-i-n photodiode to perform the autocorrelation function by two-photon absorption. An autocorrelation trace is shown in Figure 3.43. A good 8:1 ratio is visible between the peak of the autocorrelation and the background level. The autocorrelation

![Optical spectrum from modelocked prismless Cr^4+:YAG laser.](Image)
fits well to a 32 fs sech pulse shape. These pulses are significantly longer than the 20 fs pulses produced by the 6 DCM reflection cavity with DCM fold mirrors.

![Figure 3.43](image)

Figure 3.43  Measured interferometric autocorrelation (black line) and 32 fs sech fit (gray dots).

It was not possible to generate shorter optical pulses with the 6 DCM prismless cavity described in Figure 3.41 (unchirped cavity fold mirrors) than with the 6 DCM prismless cavity in Figure 3.28 (DCM cavity fold mirrors). The problem with replacing the cavity fold mirror DCMs with unchirped mirrors is that most unchirped mirrors do not have large enough bandwidth to support short pulses. Therefore, while this replacement removes large oscillations due to the off-axis reflection from DCMs, the 4 reflections off smaller bandwidth unchirped mirrors become limiting. There are a few approaches that may be effective in removing DCM oscillations. DCMs with $\pi$ phase shifts have been proposed and fabricated for Ti:Sapphire lasers [135], which have GDD oscillations which are 180 degrees out of phase. Each reflection off one type of DCM is compensated by a reflection off its complement. A second approach to minimize the GDD oscillations requires the use of shorter laser crystals. Shorter
gain crystals require smaller angles for astigmatism compensation and therefore lead to smaller net cavity dispersion oscillations. Unfortunately, there has been significant difficulty in obtaining high quality short crystals with high gain. A third approach is to design two different sets of DCMs, one designed for normal incidence, and one designed for a specific angle. This approach would be quite expensive to implement.

3.13 Shorter Crystal

In the prismless laser cavity discussed above, the minimum pulse width generated may be limited by the wavelength range of dispersion compensation within the laser cavity. For example, to achieve modelocking at wavelengths above 1500 nm, it is necessary to find a laser configuration that eliminates use of the long wavelength limiting unchirped high reflector. Large oscillations in the GDD, especially caused by the angled cavity-fold DCMs, may also be problematic. Furthermore, while the use of DCMs in a prismless-fold laser cavity allow coarse tuning of the cavity dispersion, it is not possible to easily fine-tune the intracavity GDD.

A 'go-for-broke' try at generating shorter pulses from a Cr⁴⁺:YAG laser was attempted using a 1 cm laser crystal, 5 DCM mirror bounces per cavity roundtrip, and intracavity BaF₂ prisms. The 1 cm crystal requires an 11 degree astigmatism compensation angle, smaller than the 16 degree angle required for 2 cm laser crystals. Cavity fold mirrors placed at this angle have smaller GDD oscillations. The use of an odd number of DCM bounces allows for Z-fold cavities with only DCMs, removing the bandwidth limiting unchirped high-reflector. Dispersion from a combination of 5 DCM bounces and the Cr⁴⁺:YAG laser crystal can be well compensated with low TOD using the normal material dispersion from BaF₂ prisms. The prism insertion can be varied continuously to fine-tune the intracavity dispersion for pulse width optimization. Although many normal dispersion prism materials were found to be inadequate for intracavity use because of high water content, crystalline BaF₂ should have inherently low water content.

A schematic of the laser cavity is shown below in Figure 3.44. The laser crystal was the a 1 cm long, 2 by 5 mm rectangular Brewster-Brewster cut rod. Pump light at 1064 nm from a Nd:YVO₄ laser is focused by a 10 cm focal length lens into the crystal. The crystal had an experimentally measured absorption of 1.6 cm⁻¹. Surrounding the laser crystal are two 10 cm radius-of-curvature DCMs (M1 and M2), rotated 11 degrees from normal incidence for astigmatism compensation. One arm of the laser, 60 cm long, contains both an additional
DCM (M3) and 2 BaF₂ Brewster cut prisms. The second arm of the laser, 50 cm long, contains a broadband output coupler (OC3) with a minimum transmission of 0.7% at 1515 nm and less than 1.4% transmission between 1420 to 1630 nm. The lasing threshold is at 1 W of absorbed pump power. The maximum cw output power, for 5.5 W of absorbed pump, was 400 mW without prisms, 300 mW with 1 prism, and 200 mW with 2 prisms. The laser beam path is enclosed with plastic tubes and wraps, and purged with dry nitrogen gas at all times. A picture of the laser cavity is shown in Figure 3.45.

Figure 3.44 Schematic of a Cr⁴⁺:YAG laser used in attempts to generate ultrashort optical pulses. The cavity consists of a 1 cm laser crystal, an output coupler (OC), 2 double-chirped mirrors (DCMs) (M1 and M2, in black), and one unchirped high reflector (M3, in gray).

Figure 3.45 Picture of the Cr⁴⁺:YAG laser.
One of the major drawbacks of using the 1 cm laser crystal is its relatively low gain. The lower gain results from saturation of the pump absorption. While the linear absorption of the 1 cm crystal is 80%, there is only 50% absorption for the power densities used to pump the KLM cavity due to saturation. To investigate the 1 cm crystal’s potential for broadband gain, a tuning curve was measured using the single prism technique discussed in Figure 3.7 (b). A plot of the 1 cm laser tuning curve is shown below in Figure 3.46. It was possible to tune the laser from 1420 to 1515 nm. Because the laser was not purged with Nitrogen gas during tuning, it is likely that wavelengths below 1420 nm were inhibited from lasing by the strong water absorption lines in that range. Although not as wide as for the 2 cm crystal, the gain bandwidth of the 1 cm crystal, spanning over 100 nm wide, should be suitable for the generation of ultrashort pulses.

Figure 3.46 Tuning curve of the 1 cm Cr³⁺:YAG crystal based laser.
The GDD of two possible dispersion compensated cavity designs using the 1 cm crystal are shown below in Figure 3.47. Highly compensated cavities can be created easily using either 4 or 5 DCM bounces per cavity roundtrip. Plotted are the GDD of a double pass through a 1 cm Cr\(^{4+}\):YAG crystal with 4 and 5 DCM bounces, the material GDD of 15 and 25 mm of BaF\(_2\), a combination of Cr\(^{4+}\):YAG, 4 DCM bounces, and 15 mm of BaF\(_2\), and a combination of Cr\(^{4+}\):YAG, 5 DCM bounces, and 25 mm of BaF\(_2\). The 5 DCM bounce cavity was chosen because it does not require the additional use of a cavity bandwidth limiting unchirped high-reflector. The net cavity GDD is between -250 and 250 fs\(^2\) within the wavelength range from 1300 to 1625 nm. The flat GDD spectral range for the 1 cm crystal is even larger than the corresponding net cavity GDD compensated bandwidth for the prismless 2 cm crystal cavity, which covers the wavelength range from 1325 to 1525 nm.

The laser cavity was constructed, and modelocking was attempted by moving the crystal and cavity fold mirror positions and tapping an end mirror to initiate modelocking. The laser was purged at all times with Nitrogen gas. During these experiments, no signs of mode-
locking were observed. To determine if the laser was limited by the BaF$_2$ material, the prisms were removed. The remaining prismless cavity should behave similarly to the prismless 2 cm laser cavity. Yet again, the cavity parameters were varied, and no modelocking was observed.

There are several factors that could inhibit modelocking. One plausible effect results from the strong pump absorption saturation observed in the 1 cm crystal. Soft aperture KLM requires the spatially-dependent gain profile to provide more gain for pulses than for cw radiation. This generally requires the gain profile to be about a factor of 1.5 to 2 smaller than the cavity beam waist. In the absence of pump absorption saturation, the gain profile should be proportional to the pump profile. With absorption saturation, the gain profile $g(x)$ will vary spatially like:

$$g(x, I_{sat}) \sim \frac{I(x)}{1 + \frac{I(x)}{I_{sat}}}$$  \hspace{1cm} (3.9)

where $I(x)$ is the intensity profile of the pump beam, and $I_{sat}$ is the saturation intensity. The high intensity center of the Gaussian beam will saturate the absorption greater than the low intensity wings of the pump beam. The resulting gain profile will therefore have a larger radius than the pump beam waist. A plot of the gain profile resulting from a Gaussian pump beam for several saturation intensities is shown below in Figure 3.48. It is possible to estimate the saturated width of the gain profile in the case of the 1 cm laser crystal with an order of magnitude calculation. Because the linear absorption is 80%, compared with a 50% saturated absorption, $I_{sat}$ can be estimated by solving:

$$\frac{\int g(x, I_{sat})dx}{\int g(x, \infty)dx} = \frac{5}{8}$$  \hspace{1cm} (3.10)

Once $I_{sat}$ is known, the waist of $g(x, I_{sat})$ can be determined simply. This was done numerically, yielding a gain waist of 1.4 $\omega$, where $\omega$ is the beam waist of the pump beam. This calculation suggests that soft aperture KLM will be ineffective within this cavity. A hard aperture, such as a slit, can be introduced to the cavity to increase the strength of KLM in the absence of, or in conjunction with, soft aperture KLM. The hard aperture must be placed in the cavity in a place where the cw cavity mode has a larger beam waist than the nonlinear cavity mode. The slit can then provide more loss for cw radiation than for pulses. A quick attempt at hard aperture KLM was attempted. Spectrometer slits were placed within the laser cavity to clip the tangential part of the beam. This has been reported as an optimal location for hard aperture KLM slits [60] in similar Z-fold cavities. For varying slit widths, modelocking was
attempted. No signs of pulses were observed. These attempts were nonexhaustive, and further optimization of the cavity for hard aperture KLM could be successful.

![Gain profile for a Gaussian pump beam and several levels of pump absorption saturation.](image)

Figure 3.48  Gain profile for a Gaussian pump beam and several levels of pump absorption saturation.

### 3.14 Future Work

The minimum pulsewidth might be limited by the wavelength range of dispersion compensation within the laser cavity. For example, to achieve modelocking at wavelengths above 1500 nm, it is necessary to find a laser configuration that eliminates use of the long wavelength limiting unchirped high reflector. Large oscillations in the GDD, especially those caused by the angled cavity fold DCMs, may be problematic. The effect of these oscillations can be reduced by using a shorter Cr$^{4+}$$: $YAG crystal which reduces not only the number of DCM bounces required within the cavity but also reduces the astigmatism compensation angle of the cavity folding curved DCMs. Unfortunately, to date, long crystals have been
required to overcome excited state absorption (ESA) and provide enough gain. Alternatively, DCMs can be designed to have smaller GDD oscillations. DCM pairs shifted by a \( \pi \) phase-shift theoretically allow the reduction of dispersion oscillations.

Further decreases in pulsewidth will likely require the use of a tunable dispersive element within the laser cavity. Positive dispersion prisms, such as BaF\(_2\), with low water content could be used in conjunction with DCMs and would likely yield shorter pulses. Here too, it would be useful to have a 1 cm crystal with high gain and low pump saturation. The 1 cm crystal, in conjunction with 5 DCM bounces and BaF\(_2\) prisms is a particularly promising configuration. Wedges with small apex angles could introduce tunable dispersion with easy alignment compared with prisms [152]. Pairs of wedges could be inserted intracavity at Brewster’s angle. Unfortunately, the wedge apex angle needed to achieve GDD tunability leads to small deviations from Brewster’s angle, which add too much loss for the Cr\(^{4+}\):YAG laser.

Improving the intracavity dispersion compensation decreases the pulsewidth of a laser by reducing pulse broadening per cavity roundtrip. Alternatively, the pulsewidth can also be reduced by increasing the cavity pulse shortening mechanisms. The primary pulse shortening mechanism is Kerr lens modelocking. The strength of KLM is highly dependent upon the cavity geometry and pump mode. For all experiments reported above, standard KLM cavities have been employed. These cavity configurations were primarily developed for Ti:Sapphire lasers. There are several factors that suggest that an optimal KLM cavity for Cr\(^{4+}\):YAG should be quite different from that for Ti:Sapphire. The thermal lensing in Cr\(^{4+}\):YAG is much higher than Ti:Sapphire because of its lower thermal conductivity and higher thermo-optic coefficient. Thermal lensing is difficult to model and dramatically alters the cavity mode. An order of magnitude longer crystal is used for Cr\(^{4+}\):YAG than for Ti:Sapphire, and the lasing wavelength is different. Further pulse shortening should be possible by using more sophisticated modeling methods to optimize cavities for KLM. In addition to optimizing for KLM, cavities designed for shorter pulses should consider and exploit the effect of dispersion managed modelocking [66]. KLM models assume that the pulsewidth is constant throughout the cavity roundtrip. For very short pulses, the dispersion of individual cavity elements will cause the pulse to breathe within a cavity roundtrip. The ultimate pulsewidth will then depend on the order and dispersion of the various cavity elements. These dispersion managed effects are reported to dominate the shaping of very short pulses. Taking advantage of this effect, it should be possible to increase the pulse shortening strength.

By incorporating some of the design considerations discussed above, it should be possible to generate optical pulses near 1500 nm only a couple of optical cycles long. Once the
pulsewidth is shortened to a suitable value, other characteristics of the laser can be improved. For example, either self-starting or high repetition rate Cr\textsuperscript{4+}:YAG lasers could be constructed using DCMs for dispersion compensation. In most KLM lasers, modelocking must be initiated by some sort of external perturbation. Saturable Bragg reflecting mirrors have been used to self-start modelocking in KLM Cr\textsuperscript{4+}:YAG lasers [119, 121] but may limit the pulse width by the mirror bandwidth. Zhang et. al. have generated self-started pulses as short as 44 fs using a broadband saturable absorber mirror and fused silica prisms for dispersion compensation [117, 118]. DCMs might prove useful in shortening the pulses from such a system. Chapter 4 discusses the development of a broadband SBR for use with a DCM compensated cavity. Shortened laser cavities have been used to increase the pulse repetition rate, but are often limited in both cavity length reduction and dispersion compensation by fused silica Littrow prisms [123, 124, 125]. Microchip lasers have recently been reported [91] and either use Gire-Tournier interferometers (GTIs) for dispersion compensation or operate in the normal group delay dispersion (GDD) regime. DCMs could be used for dispersion compensation in these microchip lasers, allowing for the generation of short pulses with a high repetition rate.

A variety of further experiments are feasible with ultrafast Cr\textsuperscript{4+}:YAG pulses. Frequency doubled Cr\textsuperscript{4+}:YAG could be synchronized with a Ti:sapphire laser and phase-locked to permit the extension of optical frequency standard combs to telecommunication wavelengths [6, 7, 8]. Two Ti:Sapphire lasers have already been synchronized by a group at JILA [155] and recently a Ti:Sapphire and a Cr:Forsterite laser were synchronized by interaction in a Kerr medium [156]. Once phase-locked, the Ti:Sapphire and Cr\textsuperscript{4+}:YAG laser pulses could be combined into sub-two cycle optical pulses. A plot of the spectra from the Cr\textsuperscript{4+}:YAG and Ti:Sapphire lasers are shown in Figure 3.49. Together, these two lasers span at least three octaves of spectrum in wavelength, and have a Fourier pulsewidth limit of 3.1 fs. Additional spectrum could be obtained by continuum generation [157]. Some standard telecommunication fibers should have a low enough dispersion to allow continuum generation from nonlinear effects to occur before the pulse spreads. It is probable that enough spectrum could be generated to allow self-referencing of the Cr\textsuperscript{4+}:YAG spectrum to stabilize the optical phase relative to the pulse envelope. Finally, this laser could be used to perform ultrafast spectroscopic measurements on a variety of materials and devices or as a broadband coherent source for imaging techniques.
Figure 3.49 Spectra from the Cr\textsuperscript{4+}:YAG and Ti:Sapphire lasers. Together, the spectra span two octaves in wavelength.

3.15 Conclusion

In conclusion, several ultrafast Cr\textsuperscript{4+}:YAG laser cavities have been built using different combinations of DCMs and prisms for dispersion compensation. Using only DCMs for dispersion compensation, pulses as short as 20 fs have been measured by interferometric autocorrelation. The corresponding pulse spectrum is centered at 1460 nm and spans from 1310 to 1500 nm at full-width half maximum.
CHAPTER 4: BROADBAND SATURABLE ABSORBERS FOR ULTRAFAST CR$^{4+}$:YAG LASERS

4.1 Introduction

Broadband saturable absorber Bragg mirrors (SBRs) have been designed, fabricated, and used to modelock a Cr$^{4+}$:YAG laser. The SBRs are designed to generate picosecond pulses by pure saturable absorber modelocking or to initiate Kerr lens modelocking (KLM). The mirrors are comprised of high-dielectric contrast GaAs/Al$_x$O$_y$ Bragg quarter-wave dielectric stacks and have a bandwidth broad enough to support ultrafast pulses. The SBRs were used to achieve saturable absorber modelocking tunable from 1400 to 1525 nm and to self-start 35 fs Kerr lens modelocked pulses.

4.2 Background

Saturable absorbers are materials with an intensity dependent absorption. The absorption by a typical saturable absorber studied in Erik Thoen’s thesis [158] is shown in Figure 4.1. At low energy densities, light will propagate through the saturable material with some linear absorption fraction. As the energy density is increased, more carriers will be excited to the material’s upper state and the absorption fraction will decrease because of the reduction of the population difference between the ground and excited states. At even higher energy densi-
ties, the absorption fraction will level off and the absorber is said to be completely saturated. Once saturated, only nonsaturable linear loss mechanisms contribute to absorption.

![Absorption vs Energy Density](image)

**Figure 4.1** Absorption fraction for transmission through a typical saturable absorber material.

Saturable absorbers are used to modelock laser cavities and generate short pulses of light. Short pulses, with high peak intensities, will saturate absorption more easily and experience less loss than lower intensity continuous wave (cw) light. In the absence of additional pulse shaping mechanisms, the optical pulsewidth will be determined by the recovery time of the saturable absorber. Semiconductor quantum wells have a recovery times spanning from 1 ps to over 10 ns, but cannot alone produce ultrashort femtosecond pulses. Saturable absorbers, however, can be used in conjunction with Kerr lens modelocking (KLM) to produce extremely short optical pulses.

Saturable absorbers solve the self-starting problem of passive modelocking. As discussed in Chapter 2, fast passive modelocking techniques often have small pulse shortening rates for long pulses. Some external perturbation, such as ‘tapping’ an end mirror or using an
intracavity active modulator is often used to create a transient pulse that is then compressed further by passive modelocking. It is possible, with extremely good cavity design and alignment, to self-start modelocking without external perturbations [47, 48, 52, 53, 159, 160]. This state, however, is difficult to realize in practice. As a simple alternative, a real saturable absorber material that initiates longer picosecond pulses can be placed inside a laser cavity to self-start pulses. Once initiated, a pulse can be easily shortened to the femtosecond regime with passive modelocking techniques such as KLM. The use of intracavity saturable absorbers is generally simple and insensitive to alignment. Pulses close to 5 fs have been generated from a Ti:Sapphire laser using a semiconductor saturable absorber mirror and KLM [137].

A particularly convenient way of introducing saturable absorber materials, such as semiconductor quantum wells or quantum dots, into lasers is to deposit the material on a cavity mirror [49]. When the laser mirror is a Bragg reflector, these saturable absorber mirrors are referred to as saturable Bragg reflectors (SBRs) [50]. Bragg mirrors make excellent laser cavity mirrors because of their low loss. If intracavity dispersion is compensated well enough within a KLM cavity, the bandwidth of the Bragg mirror can become the primary limiting factor inhibiting the production of shorter pulsewidths. In such lasers, it is critical to design a broadband SBR with a mirror bandwidth exceeding the spectral bandwidth of the desired pulses.

Several groups have reported on the design and operation of saturable absorber mirrors for modelocked Cr$^{4+}$:YAG lasers [119 - 122]. These absorbers must be designed with an absorptive element near the Cr$^{4+}$:YAG emission wavelength (∼1500 nm), a low nonsaturable loss (<2%) because of Cr$^{4+}$:YAG's low gain, and a mirror bandwidth broad enough (>200 nm) to support the optical pulsewidth. Each saturable absorber reported used either a InGaAs/InP [119, 121] or InGaAs/InAlAs [120, 122] quantum well that emits near 1550 nm, grown upon GaAs/AlAs Bragg reflecting mirrors. These mirrors typically have a useful bandwidth of ∼100 nm for 30 layer pairs and limit the pulsewidth to ∼60 fs. Zhang et. al. succeeded in generating 44 fs pulses from a laser started by a InGaAs/InAlAs quantum well saturable absorber bonded onto an enhanced gold mirror [117, 118]. While uncoated gold mirrors are extremely broadband, they have too much loss to allow lasing of a Cr$^{4+}$:YAG laser. The gold mirror reflectivity was therefore enhanced by adding three layers of SiO$_2$/TiO$_2$/SiO$_2$.

In Chapter 3, a 20 fs Cr$^{4+}$:YAG laser was demonstrated with a fwhm bandwidth of 190 nm. Modelocking (not self-starting) was initiated by tapping an end-mirror. This chapter describes the development of an InGaAs/InP quantum well SBR with broadband reflection from an oxidized GaAs/Al$_x$O$_y$ Bragg mirror. Self-started pulses were tunable from 1400 to 1525 nm with an intracavity birefringent filter. KLM was optimized and pulses as short as 35
fs were obtained. SBRs were designed and characterized with the help of Juliet Goinath. Dr. Gale Petrich patiently grew several SBR films using gas source molecular beam epitaxy during design optimization. To create the broadband SBRs, GaAs/AlAs layers were oxidized to create GaAs/Al$_x$O$_y$ layers. This oxidation step was performed by Alexei Erchak.

An alternative broadband SBR approach based on Si/SiO$_2$ mirrors and InAs quantum dots [161, 162] is being pursued in parallel. These SBRs potentially have a more robust fabrication process consisting of high temperature chemical vapor deposition of the mirror layers and electron beam deposition of the InAs quantum dots. Si/SiO$_2$ mirrors were provided by Dr. Kazumi Wada, and the InAs quantum dots will be designed and deposited by Rohit Prasankumar. The transmission of the Si/SiO$_2$ mirrors was measured by Aurea Tucay. The performance of the Cr$^{4+}$:YAG laser with SBRs was studied with the adept assistance of Hanfei Shen.

4.3 Two Level System

Saturation of absorption results when a transition is driven with enough intensity to equalize the populations in the ground and excited states. The steady state population of states in an undriven system are distributed according to thermal equilibrium. Absorption will be proportional to the population difference. As more carriers are excited the population difference, and correspondingly absorption, decreases. The qualitative effect of saturable absorption can be modeled by simple two-level rate equations. The derivation of the rate equations closely follows reference [34].

A diagram of a two level system is shown in Figure 4.2. Two electronic levels with energies $E_1$ and $E_2$ are separated by an energy $E_{21} = h\nu_{21} = E_2 - E_1$. Stimulated and spontaneous transitions are allowed between states 1 and 2. The stimulated rates of excitation and emission are given by $W = W_{21} = W_{12} = \sigma I$, where $I$ is the driving intensity and $\sigma$ is the cross section. Spontaneous transitions, caused by thermal radiation, have excitation and emission rates of $\gamma_{12}$ and $\gamma_{21}$ respectively. At thermal equilibrium, $\gamma_{12}$ and $\gamma_{21}$ are related by a thermal distribution:

$$\frac{\gamma_{12}}{\gamma_{21}} = \exp\left(-\frac{E_{21}}{k_B T}\right) \equiv 0$$

(4.1)

where $k_B$ is the Boltzmann constant, $T$ is the temperature, and $E_{21} \gg k_B T$ at room temperature.
Figure 4.2  Schematic of a two level system. The two levels have an energy difference of $E_2 - E_1$, have a stimulated excitation rate of $W_{12}$, a stimulated emission rate of $W_{21}$, a spontaneous excitation rate of $\gamma_{12}$, and a spontaneous emission rate of $\gamma_{21}$.

Because any electrons leaving state 1 enter state 2, and vice versa, the time rate of change of the population in state 1 ($N_1$) and state 2 ($N_2$) is given by:

$$\frac{dN_1}{dt} = \frac{dN_2}{dt} = \sigma I (N_2 - N_1) + \gamma_{21} N_2$$  \hspace{1cm} (4.2)

This two-state rate equation can be transformed into an equation for the more useful variables $N = N_1 + N_2$ (total number of electrons) and $\Delta N = N_1 - N_2$ (population difference). The population difference is defined such that a negative value represents a population inversion. The rate equation can be rewritten as with the new variables to be:

$$\frac{d}{dt} \Delta N = -2\sigma I \Delta N - \frac{(\Delta N - N)}{T_1}$$  \hspace{1cm} (4.3)

where $T_1 = 1/\gamma_{21}$.

In the high intensity limit ($\sigma I >> \gamma_{21}$), stimulated processes will dominate over spontaneous processes, and the steady state solution will be $\Delta N = 0$. The lack of a population difference will result in a total saturation of absorption.

The steady state population difference for intermediate excitation intensities can now be calculated by setting equation (4.3) to be equal to zero. In this case:

$$\Delta N_{ss} = \frac{N}{1 + \frac{I}{I_{sat}}}$$  \hspace{1cm} (4.4)
where $\Delta N_{ss}$ is the steady state population difference and $I_{\text{sat}} = 1/2\sigma T_1$ is called the saturation intensity. Because the absorption coefficient $\alpha$ is proportional to the population difference $\Delta N$, we arrive at a formula for saturated absorption:

$$\alpha = \frac{\alpha_0}{1 + \frac{I}{I_{\text{sat}}}}$$

(4.5)

where $\alpha$ is the absorption and $\alpha_0$ is the small signal absorption. While this absorption saturation has been derived for a simple two-level system, the form of equation (4.5) holds more generally for many multilevel systems.

### 4.4 Saturable Bragg Reflector Structure

The basic design of a saturable Bragg reflector consists of the combination of a broadband reflector and a saturable absorber element with suitable saturation and spectral properties. Schematics of two high-dielectric contrast broadband saturable Bragg reflectors are shown in Figure 4.3. Each SBR is based in a material system with which it is possible to fabricate broadband high-dielectric contrast mirrors. GaAs and Al$_x$O$_y$ have indices of refraction of 3.39 and 1.61 at 1.5 $\mu$m respectively, while Si and SiO$_2$ have indices of 3.48 and 1.45.
With index contrasts this high, it is possible to create mirrors with a reflection over 99.5% for a wavelength range over 1200 to 1800 nm with a 7 period dielectric stack.

Appropriate saturable absorber materials are constrained by the particular choice of mirror substrate. On GaAs/AlₓOᵧ mirrors (initially grown as GaAs/AlAs), it is possible to grow an InGaAs/InP quantum well which can act as a saturable absorber. All Cr⁴⁺:YAG SBRs to date have used either InGaAs/InP or InGaAs/InAlAs quantum well absorbers on GaAs/AlAs mirrors. On Si/SiO₂ mirrors, a sputtered InAs quantum dot layer in a Sapphire matrix is a potential saturable absorber. InAs quantum dots have been shown to self-start an ultrafast Cr:Forsterite laser in a transmissive geometry, but have a saturation intensity that is larger than that of InGaAs/InP by orders of magnitude [163].

4.5 Mirror Design

Highly reflective broadband mirrors are designed with alternating quarter wave thick layers of high and low index materials. The reflectivity and phase shift of a mirror can be calculated by a matrix propagation method [164]. Incident and outgoing electric fields obey:

\[
\begin{bmatrix}
A_2 e^{-i k_{2z} d_2} \\
B_2 e^{i k_{2z} d_2}
\end{bmatrix} = V_{21}(\lambda) \begin{bmatrix}
A_1 e^{-i k_{1z} d_1} \\
B_1 e^{i k_{1z} d_1}
\end{bmatrix}
\]

where A and B are the incoming and outgoing tangential electric fields for TE and tangential magnetic fields for TM, and the propagation matrix \( V_{21}(\lambda) \) is given by:

\[
V_{21}(\lambda) = \frac{1}{2} (1 + p_{21}) \begin{bmatrix}
\exp[-i k_{2z} (d_2 - d_1)] & R_{21} \exp[-i k_{2z} (d_2 - d_1)] \\
R_{21} \exp[i k_{2z} (d_2 - d_1)] & \exp[i k_{2z} (d_2 - d_1)]
\end{bmatrix}
\]

where \( d_1 \) and \( d_2 \) are the thicknesses of layers 1 and 2, \( k_2 \) is given by:

\[
k_{iz} = \frac{2\pi n_i}{\lambda} \cos \theta
\]

where \( \theta \) is the angle of incidence and \( n_i \) is the index of refraction of layer i. \( R_{21} \) is given by:

\[
R_{21} = \frac{1 - p_{21}}{1 + p_{21}}
\]

and \( p_{21} \) is given for TE radiation by:
\[
p_{21}(\text{TE}) = \frac{\mu_2 k_{1z}}{\mu_1 k_{2z}}
\]
(4.10)
and for TM by:
\[
p_{21}(\text{TM}) = \frac{\varepsilon_2 k_{1z}}{\varepsilon_1 k_{2z}}
\]
(4.11)

For normal incidence, it is convenient to use TE, where \( p_{21} = n_1/n_2 \).

Bragg reflectors consist of alternating high and low index of refraction quarter-wave stacks, where \( d_i = \lambda/4n_i \). For multiple layers, the propagation matrix is given by:
\[
V_{\text{tot}} = V_{\text{sh}} \cdot (V_{\text{hl}} \cdot V_{\text{lh}})^N \cdot V_{\text{h0}}
\]
(4.12)
where \( V_{\text{h0}} \) is the propagation matrix for the freespace high index layer interface, \( V_{\text{hl}} \) is the matrix for going from the high to low index, \( V_{\text{lh}} \) is the matrix for going from low index to high index, and \( V_{\text{sh}} \) is the matrix for propagation from the high index layer to the substrate. The complex reflection coefficient from a TE multilayer stack is then given by:
\[
\rho \exp[-i\phi] = \frac{V_{\text{tot}}(2,1)}{V_{\text{tot}}(2,2)}
\]
(4.13)
The reflectivity is then:
\[
R = |\rho|^2
\]
(4.14)
while the phase shift from reflection is \( \phi \). The mirror's group delay dispersion (GDD) can then be calculated from the phase shift by:
\[
\text{GDD} = \frac{\frac{d^2 \phi}{d\omega^2}}
\]
(4.15)
where \( n_H \) is the index of refraction of the high index material, \( n_L \) is the index of refraction of the low index material, and \( N \) is the number of high-low index material layer pairs deposited.

Figure 4.4 shows the calculated reflectivity of a mirror designed for broadband reflectivity centered at 1440 nm with 7 layer pairs of GaAs/Al\(_x\)O\(_y\). Because Cr\(^{4+}\):YAG lasers are sensitive to losses of \( \sim 0.5\% \), a 7 layer pair was chosen. The mirror reflectivity is 99.5% within the wavelength range from 1200 to 1800 nm. It was not possible to measure the reflectivity
directly from the GaAs/AlₓOᵧ mirrors with the Cary spectrophotometer available to us because the oxidation process limits the mirror width to ~ 300 µm.

![Graph showing reflectivity vs. wavelength](image)

Figure 4.4 Calculated reflectivity of a high-dielectric contrast mirror consisting of a 7 pair GaAs/AlₓOᵧ quarter-wave dielectric stack.

The calculated GDD for the 7 layer pair mirror is shown in Figure 4.5. The GDD is close to zero over the highly-reflective wavelength range from 1200 to 1800 nm and, therefore, should not change the net cavity dispersion significantly. Because of the extremely broadband flat GDD, high-dielectric contrast mirrors could be ideal for unchirped high-reflectors or broadband output couplers within ultrafast lasers.
Figure 4.5  Calculated group delay dispersion (GDD) of a high-dielectric contrast mirror consisting of a 7 pair GaAs/AlₓOᵧ quarter-wave dielectric stack.

In the Si/SiOₓ system, a 4 layer pair mirror centered at 1475 nm was chosen because of availability. It is also likely that more Si/SiOₓ layers will cause the layer thickness uniformity to degrade [165]. The reflection of the mirrors, measured using a CARY spectrophotometer, is plotted in Figure 4.6 along with a calculated reflection curve assuming a 4 layer dielectric stack mirror centered at 1475 nm. An advantage of Si/SiOₓ mirrors is that wafer-sized Bragg stacks can be easily fabricated by chemical vapor deposition. No size limiting oxidation step is necessary. The reflectivity is measured to be 99.9% at its maximum and above 99.5% from 1220 to 1840 nm. In addition to a high reflector, the 4 layer mirrors could be an effective broadband output coupler.
Figure 4.6  Measured and calculated reflection for a 4 pair Si/SiO₂ Bragg mirror centered at 1475 nm.

4.6  **Saturable Absorber Design**

Saturable absorbers are primarily characterized by their saturation fluence, saturable loss, and nonsaturable loss. It is important to find an absorber with suitable material properties for the laser. Because Cr⁴⁺:YAG is highly sensitive to loss, it is important to minimize the nonsaturable loss as much as possible. The saturable loss should ideally be less than 0.5% (for a laser with ~ 4% gain) in order to provide a deep modulation depth for initiation of mode-locking. Saturation fluences should be on the order of the energy density within the laser cavity. The Cr⁴⁺:YAG laser typically had an intracavity power of ~ 10 W at a repetition rate of 100 MHz. Each pulse therefore had an intracavity pulse energy on the order of 0.1 μJ. SBRs are placed within the laser cavity after a focusing mirror, usually with a 10 cm radius of curvature. These mirrors focus the light to ~ 50 μm beam waist. The typical energy density incident upon an SBR is thus ~ 1 mJ/cm². A saturable absorber should be chosen to have a
saturation fluence near this value. It is possible to use absorbers with saturation fluences that differ from this value by an order of magnitude by selecting a different intracavity focusing mirror or by proper placement of the saturable absorber in the mirror's standing wave field.

Proper placement of the saturable absorber layer upon the Bragg mirror is critical for optimal operation of the SBR. Light reflected from the mirror will interfere with incident light to create an optical standing wave. The saturable absorber's placement in this field relative to a node or antinode will dramatically change the laser field energy density encountered. Figure 4.7 shows the electric field energy profile of light reflecting from the GaAs/Al$_x$O$_y$

![Electric field energy profile](image)

**Figure 4.7** Electric field energy amplitude and index of refraction of the designed Bragg reflector (SBR) mirror consisting of a GaAs/Al$_x$O$_y$ high-index contrast mirror and an InGaAs/InP quantum well.

broadband SBR. The structure consists of an 7 layer pair of a GaAs/Al$_x$O$_y$ Bragg mirror. The saturable absorber, an InGaAs/InP quantum well, consists of a half-wave thick InP region with a centered 10 nm In$_{0.52}$Ga$_{0.48}$As layer. The InGaAs layer is placed within at the antinode of the standing wave in order to increase the local energy density. The total thickness of the quantum well and cladding layers is chosen to be a half wavelength in order to maintain the dispersion properties of the structure.

For the GaAs based saturable absorbers, the quantum well wavelength, thickness, and placement was chosen so that the absorber would have material properties suitable for the Cr$^{4+}$:YAG laser cavity. Four separate low-index-contrast GaAs/AlAs-mirror-based SBR
designs were compared to determine the proper absorber choice. The characteristics of four previous SBR designs: SBR-I, SBR-II [158, 166], SBR-A2, and SBR-D2 [150], are listed in Table 4.1. Each SBR consisted of an InGaAs/InP quantum well grown upon a low-index-con-

<table>
<thead>
<tr>
<th>SBR</th>
<th>Mirror Bandwidth</th>
<th>Quantum Well Composition</th>
<th>Number of Quantum Wells</th>
<th>Quantum Well Thickness</th>
<th>PL Center Wavelength</th>
<th>Saturable Loss</th>
<th>Saturation Fluence</th>
</tr>
</thead>
<tbody>
<tr>
<td>SBR I</td>
<td>100 nm (tunable)</td>
<td>In$<em>{0.52}$Ga$</em>{0.48}$As</td>
<td>2</td>
<td>4.5 nm</td>
<td>1450 nm</td>
<td>0.4%</td>
<td>~ 100μJ/cm$^2$ @ 1450 nm</td>
</tr>
<tr>
<td>SBR II</td>
<td>100 nm (tunable)</td>
<td>In$<em>{0.52}$Ga$</em>{0.48}$As</td>
<td>2</td>
<td>6.6 nm</td>
<td>1490 nm</td>
<td>0.6%</td>
<td>~ 100μJ/cm$^2$ @ 1470 nm</td>
</tr>
<tr>
<td>A 2</td>
<td>1525 – 1650 nm</td>
<td>In$<em>{0.51}$Ga$</em>{0.49}$As</td>
<td>1</td>
<td>4.0 nm</td>
<td>1610 nm</td>
<td>0.3%</td>
<td>100μJ/cm$^2$ @ 1530 nm</td>
</tr>
<tr>
<td>D 2</td>
<td>1425 – 1550 nm</td>
<td>In$<em>{0.51}$Ga$</em>{0.49}$As</td>
<td>1</td>
<td>4.0 nm</td>
<td>1567 nm</td>
<td>0.3%</td>
<td>100μJ/cm$^2$ @ 1530 nm</td>
</tr>
</tbody>
</table>

Table 4.1 Properties of four low-index contrast SBR designs tried with a Cr$^{4+}$:YAG laser.

trast GaAs/AlAs distributed Bragg reflector with a bandwidth of about 60 nm. SBR-I and SBR-II had two quantum wells placed near the surface of the half-wave InP layer, close to the electric field node. The properties of the SBR were dependent on the portion of the wafer used. For example, the mirror stopband shifted from a center wavelength of 1550 nm (near the wafer center) to 1350 nm (near the wafer edges). Using these SBRs, the Cr$^{4+}$:YAG laser operated in cw mode for most wavelengths and modelocked weakly for wavelengths below 1425 nm. It is likely that the local energy density was not large enough to saturate the quantum wells because the quantum well absorption was low. Pump-probe studies also indicated that the saturation fluence was large near 1450 nm. Furthermore, the PL of the absorbers was at 1450 (SBR-I) and 1490 (SBR-II), and possibly at lower wavelengths toward the edge of the wafers. The low wavelength bandedge probably limited modelocking to wavelengths below 1425 nm.

SBR A2 and D2 are taken from different wafer positions of a growth consisting of an InGaAs quantum well centered inside a half-wave InP layer on top of a GaAs/AlAs mirror. This quantum well was located at the position of an electric field antinode. Furthermore, the quantum well PL was at a longer wavelength to allow modelocking over the entire Cr$^{4+}$:YAG spectral range. The Cr$^{4+}$:YAG laser modelocked for almost all pump power levels with these saturable absorbers, indicating that there was a high enough energy density near the quantum well to saturate absorption with a sufficient modulation depth. For this reason, the saturable absorbers for the oxidized saturable absorber were designed to consist of a single InGaAs/InP quantum well near the electric field antinode, with a bandedge near 1580 nm.
A photoluminescence (PL) spectra of the SBR, studied to experimentally estimate the position of the bandedge, is shown in Figure 4.8. The SBR quantum well is excited with an argon ion laser at 488 nm. PL is collected with lenses and spectrally resolved by a spectrometer. The PL is centered near 1550 nm, and has a full-width half maximum of 110 nm. The center wavelength is slightly lower than the design goal, and may not start modelocking for wavelengths below the bandgap.

To determine the saturation fluence and recovery time of the absorber, pump-probe spectroscopy was performed on the SBR for different energy densities. Pump-probe experiments are able to measure the change in reflectivity of a nonlinear mirror due to saturation or other effects, and the relaxation time scales back to the ground-state. A schematic of the
pump-probe experiment is shown below in Figure 4.9. 150 fs pulses at 1540 nm from an optical parametric oscillator (OPO) are split into two pulses by a beam splitter. One pulse is polarization rotated and the second pulse is delayed with respect to the first by a time $\tau$. The pulses are recombined by a polarization beam splitter (PBS) and focused onto the saturable absorber mirror. The first pulse (the pump) will excite the absorber. The second pulse (the probe), followed by a time $\tau$, will be reflected by the mirror and detected. The intensity of the reflected probe as a function of time delay will be modulated by the saturation state of the SBR. When the absorber is saturated, the probe will be absorbed less than the linear level. Two-photon absorption (TPA) will be observed as a decrease in the probe reflection for time-delays which overlap the pump and probe pulses [167]. More complicated carrier dynamics can also be determined by studying pump-probe [168 - 170].

Several pump-probe traces for varying pump and probe powers are shown in Figure 4.10. The OPO was tuned to 1540 nm for these measurements. The SBR response is characterized by a fast saturation due to spectral hole burning, and a long recombination time of about 40 ps. As the pump fluence is increased, significant two photon absorption (TPA) and free carrier absorption (FCA) reduce the saturable loss. The SBR is not completely saturated for the minimum pump-probe powers used (11 $\mu$J/cm$^2$), but becomes saturated at 22 $\mu$J/cm$^2$. The saturation fluence is therefore estimated to be on the order of $\sim 10$ $\mu$J/cm$^2$. The maxi-
mum saturable loss is 0.3%. At high fluences TPA acts as an inverse saturable absorber, and could limit the minimum pulsewidth for short pulses [167]. TPA due to the quantum well could be lowered in future SBRs by moving the quantum well from of the antinode of the electric field, but TPA from the quantum well cladding layers would remain constant. Further pump-probe studies are necessary to study the operation of the SBR at the wavelength and fluence of the Cr$^{4+}$:YAG laser.

![Graph showing pump-probe traces of the broadband saturable Bragg reflector at 1540 nm.](image)

Figure 4.10  Pump-probe traces of the broadband saturable Bragg reflector at 1540 nm.

Several oxidized broadband SBRs were grown and tested leading up to the most recent SBR growth. The characteristics of these preliminary SBRs are tabulated in Table 4.2.
Table 4.2 Characteristics of all saturable Bragg reflector growths in the effort to realize a broadband

In the Si/SiO₂ mirror system, InAs quantum dots are planned to be deposited in the center of a half-wave Sapphire layer. As with the InGaAs quantum wells, the InAs quantum dots will be placed at the antinode of the electric field in order to maximize the local energy density. This is especially critical for InAs quantum dots, where the saturation fluence is expected to be higher than for quantum wells because of quantum dot size nonuniformity. InAs quantum dots have not been used to modelock a Cr⁴⁺:YAG laser, but have been used to successfully modelock a Cr:Forsterite laser at 1300 nm [163].

4.7 Fabrication

The broadband SBR design was, in large part, dictated by fabrication constraints of GaAs and Si material systems. In each case, high-dielectric contrast materials were used to provide broadband reflection. The mirror layers must be deposited with enough precision to exhibit low loss reflection. High quality saturable absorbers with low saturation fluences and low nonsaturable losses must then be grown onto the broadband mirrors. In the case of GaAs, InGaAs/InP quantum well can grown. In Si, InAs quantum dots are deposited in a nonlattice-matched process.

The GaAs-based SBRs are fabricated using III-V semiconductor-based material growth techniques. First, a GaAs/AlAs multilayer stack is grown by gas source molecular beam epitaxy (GSMBE). The AlAs layers are later converted to AlₓOᵧ through an oxidation process. The mirror stack is grown to be a quarter-wave stack at 1440 nm. The GaAs layers are a quarter-wave thick (106 nm), while the AlAs layer thickness is chosen such that the oxidized AlₓOᵧ layer will be a quarter-wave thick at 1440 nm. Because AlAs layers are believed to shrink ~ 10% upon oxidation, the AlAs layer thickness is chosen to be 240 nm to arrive at 224 nm layers of AlₓOᵧ. On top of the Bragg mirror stack, a 10 nm InGaAs/InP quantum well
is grown in between two 111 nm InP cladding layers. The quantum well and cladding layers have a total optical path length of a half wavelength. The InGaAs thickness was designed to place the bandgap near 1580 nm. This would allow the saturable absorber to operate with the lowest saturation intensity near the Cr$^{4+}$:YAG emission wavelengths.

After MBE growth, the structure must be oxidized to convert high-index AlAs into low-index Al$_x$O$_y$. This oxidation process converts high-index AlAs to low-index Al$_x$O$_y$ ($n = 1.6$) laterally from the edge of the structure [171, 172, 173, 174]. Therefore, only material near an exposed edge will oxidize. This could be cleaved edge of a semiconductor wafer, a surface defect, or a line scored onto the wafer surface. For the samples discussed we found the oxidized regions near a cleaved wafer edge to work best. The SBR is placed within an oxidation chamber at 400 degrees C. Nitrogen gas is bubbled through water in a nearby chamber and the resulting water vapor flows through the sample chamber. After ~ 9.5 hours, it is estimated that the resulting Al$_x$O$_y$ layers extended as far as 300 μm into the structure. This should be suitable to work for a cavity with a beam waist of ~ 50 μm on the SBR. Side-view scanning electron micrograph (SEM) images of an unoxidized and oxidized SBR structure are shown in Figure 4.11 (a) and (b). The oxidation process is extremely sensitive to tem-

Figure 4.11 Scanning electron micrograph images of an (a) unoxidized and (b) oxidized SBR structure. After oxidation, the AlAs layers are converted to polycrystalline Al$_x$O$_y$, which appears to be granular.

perature, and is generally not used to oxidize thick layers over a large lateral dimension. Sometimes, some layers would oxidize, while others would remain unoxidized. An SEM of such a structure is shown in Figure 4.12.
Silicon based SBRs are being fabricated in a different manner from the GaAs based devices. Si/SiO$_2$ quarter-wave layers are deposited onto a Si substrate by a high temperature CVD process. Because structural nonuniformity increases with the thickness of the layer deposited, only 4 Si/SiO$_2$ pairs were grown. On top of the Bragg mirror, a layer of InAs quantum dots will be deposited inside a half-wave layer of Sapphire with an electron beam sputtering system.

4.8 **Self-Starting Cr$^{4+}$:YAG Laser**

Using broadband saturable absorber mirrors, it was possible to self-start modelocking of an ultrafast Cr$^{4+}$:YAG laser. First, a mirror was placed as an end mirror in the laser cavity and lasing was attempted. A diagram of the laser cavity used with the SBRs is shown in Figure 4.13. The laser cavity consists of a 2 cm Cr$^{4+}$:YAG laser rod pumped by an 11 W Nd:YVO$_4$ laser at 1064 nm. The pump beam is focused onto the crystal by a 10 cm focal length lens. About 5 cm to either side of the laser crystal are two 10 cm radius of curvature double-chirped mirrors (DCMs) rotated 16 degrees from normal to compensate the astigmatism of the Brewster-Brewster cut laser rod. One arm of the cavity is 50 cm long and contains an output coupler (OC). The second cavity arm contains a 10 cm radius of curvature DCM focusing onto an SBR. The curved DCM focuses the cavity mode to a small spot size on the SBR. This is done to get a large enough energy density to saturate the absorber. Using a 10 cm radius of curvature mirror, the beamwaist on the SBR should be approximately 50 μm. The beamwaist could be changed by varying the M3 radius of curvature. The laser cavity is
designed to have 6 DCM reflections each cavity round-trip of the pulse. More details about this laser cavity design can be found in Chapters 2 and 3.

![Lasercavitydiagram](image.png)

Figure 4.13 Schematic of a Cr\(^{4+}\):YAG laser cavity consisting of 3 10 cm radius of curvature double-chirped mirrors (M1 - M3), an output coupler (OC), and an saturable Bragg reflector (SBR) end mirror.

The laser was first aligned with an ordinary flat laser mirror substituted for the SBR. In the configuration described above, the laser had a cw output power of 600 mW for 9 W of absorbed pump. The ordinary end mirror was then replaced by the SBR. Care was taken to focus the laser mode onto the ~ 300 µm oxidized region of the structure. Fluorescence from the laser cavity was monitored with a Ge photodetector. The SBR was translated until the fluorescence intensity increased near the sample edge oxidized region. The SBR spatial orientation was optimized to produce maximum feedback and hopefully lasing. Several SBRs were tried within the laser cavity to test the reflectivity of the Bragg mirror. SBR R859, consisting of an 8 pair GaAs/Al\(_x\)O\(_y\) mirror stack, had high enough reflectivity to allow lasing within the cavity. The quantum well of SBR R859 had PL centered at 1470 nm, indicating that the band-edge might have been too high to modelock the laser. Also, in this early design, the quantum well layer was close to a node rather than the anti-node of the electric field standing wave within the SBR. This quantum well placement results in a large saturation fluence and causes modelocking to be difficult.

Subsequent GaAs/Al\(_x\)O\(_y\) SBRs were designed with the quantum well layer in the antinode of the electric field pattern. SBR R885 was used to successfully lase and modelock the Cr\(^{4+}\):YAG laser. The average output power was 300 mW for 9 W of absorbed pump. Using a
birefringent filter [143], the laser was tunable from 1400 to 1525 nm, as shown in Figure 4.14.

![Graph showing intensity vs. wavelength](image)

Figure 4.14 Spectra of the saturable absorber modelocked Cr\(^{4+}\)-YAG laser tuned from 1400 to 1525 nm with a birefringent tuning plate.

The broadband SBR has a bandwidth at least as wide as the tuning range. The short wavelength limit of 1400 nm was due to the roll-off of the output coupler (OC3) reflectivity. The laser produced ~ 1 ps pulses over the entire tuning range.

The birefringent filter was removed, and the SBR was used to self-start KLM. First, the cavity was enclosed in plastic tubes and purged with dry nitrogen gas to remove water vapor from the air. This water vapor causes absorption and dispersion within the laser cavity, eliminating the possibility of KLM. Regardless of alignment, the laser modelocked with the SBR in the cavity. In order to achieve KLM, the curved mirror separation between mirrors M1 and M2 and the laser crystal position were varied. Self-starting KLM was easily found by maximizing the spectral width and minimizing the pulsewidth of the modelocked pulses.
A plot of the KLM pulse spectrum is shown below with linear and logarithmic scales in Figure 4.15. The pulse spectrum is centered at 1500 nm, and has a full-width half maximum of 68 nm. Spectral components are detected from 1200 to > 1700 nm.

The pulsewidth was measured by a fringe-resolved autocorrelator. The beam was split and recombined by broadband metallic beamsplitters and focused by an off-axis parabolic mirror onto a silicon p-i-n photodiode to perform the autocorrelation function by two-photon absorption. An autocorrelation trace is shown in Figure 4.16. The measured autocorrelation exhibits a good 8:1:0 ratio as expected for the interferometric geometry. The pulse fits to a 32 fs pulsewidth assuming a sech-shaped pulse. A sech-shaped pulse fit underestimates the true pulsewidth for non-sech-shaped pulses. With the measured spectrum, a bandwidth limited pulse would have a pulsewidth of 35 fs. There is no fundamental reason why the pulsewidth should be limited to 35 fs. By optimizing the cavity dimensions and dispersion, it should be possible to match or beat the pulsewidth measured from the Cr$^{4+}$:YAG laser without the SBR.
Figure 4.16 Interferometric autocorrelation of a self-started Cr\textsuperscript{4+}:YAG laser.

High-dielectric contrast Si/SiO\textsubscript{2} mirrors were also tested for suitability as laser cavity mirrors and as substrates for InAs absorber layers. The 4 pair Si/SiO\textsubscript{2} mirror stack was placed as an end mirror within the laser cavity. No saturable absorber had been deposited onto the mirror stack. It was easy to achieve lasing using this structure as an end mirror. Translating the mirror led to minor variations in the output power of the laser, indicating a high degree of mirror spatial uniformity. This mirror stack was also tested by replacing the 10 cm radius of curvature M3 with a flat DCM mirror. In this cavity configuration, the beamwaist on the mirror is \~ 250 \mu m, rather than 50 \mu m. The laser also worked in this configuration, once again indicating that the mirror has a high degree of spatial uniformity.
4.9 Future Work

Broadband GaAs/Al$_x$O$_y$ mirror-based SBRs successfully started modelocking 35 fs pulses from a Cr$^{4+}$:YAG laser. The SBR dynamics must be studied to determine if the pulsewidth is limited by absorber effects such as TPA or if the laser cavity alignment is limiting the pulsewidth. Concurrently, laser cavities with better GDD compensation should be assembled with the broadband SBR to generate a shorter pulsewidth.

The fabrication of Si/SiO$_2$ SBRs should be pursued further. For polycrystalline Si layers, the use of sputtered InAs quantum dots has been proposed, and could have properties suitable for use within Cr$^{4+}$:YAG. In the future, it may be possible to combine the Si/SiO$_2$ mirrors with higher quality saturable absorbers. For example, by creating a crystalline Si top layer on the mirror stack, it is possible that InGaAs quantum wells could be grown epitaxially on the mirror. Alternatively, wafer bonding techniques could be used to transfer an InGaAs/InP quantum well onto a Si/SiO$_2$ mirror. It is important to continue to try other potential materials as saturable absorbers. Because Germanium can be grown directly upon Silicon-based structures, Ge-quantum dots may be a promising material for an absorber. Antimonide-base materials such as GaSb may also have interesting absorption properties.

High quality Si/SiO$_2$ mirrors could be used for mirrors other than an SBR within a Cr$^{4+}$:YAG laser cavity. The mirror layers could be deposited onto flat or curved quartz substrates. By designing the mirror to have a reflectivity dip near 1064 nm, the mirrors could transmit the pump light and reflect the laser radiation. It would then be possible to build a laser entirely composed of broadband Si/SiO$_2$ mirrors. A single mirror set might have a large enough bandwidth to work with either a Cr$^{4+}$:YAG or Cr:Forsterite laser crystal. Alternatively the extremely broadband mirrors could be used to sustain extremely short pulses. Mirrors with fewer high/low index layers could serve the function of a broadband output coupler.

4.10 Conclusions

A broadband GaAs/Al$_x$O$_y$ mirror-based InGaAs/InP quantum well SBR was designed to self-start modelocking of an ultrafast Cr$^{4+}$:YAG laser. The SBRs were used to modelock the laser over a tuning range from 1400 to 1525 nm. In a double-chirped mirror dispersion compensated cavity, the SBRs successfully self-started 35 fs pulses at 1490 nm, the shortest self-started pulses from a Cr$^{4+}$:YAG laser to the best of our knowledge. These mirrors are epitaxially grown, and therefore can be made with precise layer thicknesses. Si/SiO$_2$ mirrors were investigated as a potentially easier to fabricate high-dielectric contrast material for
broadband SBRs. A 4 pair Si/SiO$_2$ mirror was shown to work well as a highly reflective mirror substrate, and could also serve as a broadband output coupler. The mirrors are easily grown by a high temperature chemical vapor deposition (CVD) technique, are not limited in size by oxidation, and are quite uniform for a small number of layers.
CHAPTER 5: PHOTONIC BANDGAP MICROCAVITY

5.1 Introduction

One-dimensional waveguide-based photonic crystals microcavities are studied in a GaAs/AlₓOᵧ III-V compound semiconductor material waveguide system, which has the potential for the future integration of active optical devices. Microcavities designed to support an optical cavity mode at 1.55 μm are fabricated in both monorail and airbridge geometries. Airbridge devices are shown to have higher Quality factors (Q's) than monorail devices for a given modal volume because of their increased optical confinement. Optical transmission spectra through the microcavities are studied experimentally, revealing cavity Q's as high as 360 near a wavelength of 1.55 μm. A modal volume as small as 0.026 μm³ has been calculated for a measured microstructure.

5.2 Background

The density of electromagnetic states available to light confined within optical cavities is dramatically different from the density of states for freely propagating light. While photons of any momentum are able to propagate in free-space, optical cavity photons can only exist at discrete momenta corresponding to cavity modes. By altering the optical density of states, the optical emission rate of atoms within the light confinement region can be significantly modified. For example, it is possible to enhance or inhibit the spontaneous emission of atoms in a
cavity by changing the overlap between the optical cavity modes and the atomic resonance. [175 - 182] With close enough overlap between the cavity mode and atomic resonance, it is even possible to create highly coupled atom-cavity modes, where energy oscillates back and forth between the atom and the optical cavity [183, 184].

In general, smaller optical cavities contain a larger free spectral range between adjacent optical states. Large free spectral ranges are ideal for applications where a single optical mode is desired. For example, a waveguide-coupled optical microcavity could perform the function of a wavelength division multiplexing (WDM) filter if the free spectral range is large enough so that unwanted adjacent optical modes lie outside of the WDM transmission band. Highly confined optical systems might reduce the size and power requirements of active integrated optical components. Single-mode cavities overlapped with the spectrum of light-emitting devices can reduce the lasing threshold of semiconductor lasers, and even allow higher modulation speeds of these devices [185]. Coupling losses from localized optical modes into free space radiation modes limit the minimum volume to which an optical state can be confined [186].

Photonic bandgap (PBG) materials, a class of photonic crystals, can be used to achieve strong photon confinement to volumes on the order of $(\lambda/2n)^3$, where $\lambda$ is the photon wavelength and $n$ is the refractive index of the host material [187 - 189]. Highly confined optical states arise from the introduction of local defects inside photonic crystals. In the high-dielectric-contrast material systems that are often necessary for achieving PBGs, the amplitude of the electromagnetic fields falls off sharply away from the defect, resulting in strong photon confinement [190].

This chapter presents the design, realization, and optical characterization of one-dimensional photonic crystal microcavities, with modal volumes as low as $2 (\lambda/2n)^3$. The enabling technology for this project is the remarkable semiconductor fabrication developed to create the photonic crystals devices. This fabrication is the subject of Dr. Kuo-Yi Lim’s Ph.D. Thesis [171]. The work presented here was done in close collaboration with the research groups of Professors Leslie Kolodziejski and John Joannopoulos.

5.3 Photonic Crystals

At various times in history, light has been thought to be in its essence either a particle or a wave [191]. There was no consensus as to light’s true nature among the early founders of optics. Sir Isaac Newton, for example, favored a corpuscular theory of light, while Christiaan
Huygens simultaneously advocated a wave theory of light. For many years, wave theory became commonly accepted, but debate was revived again when Albert Einstein described the photoelectric effect in terms of quantized particle-like radiation. In an inspiration to all graduate students, Louis deBroglie helped reconcile these outlooks with the idea of wave-particle duality in his 1924 Ph.D. thesis [192]. Not only did he unify the corpuscular and wave theories of light, but deBroglie also extended this wave-particle duality to all fundamental particles. According to his theory, all particles propagate according to a wave theory, but when measured, exhibit a particle like qualities. Based on wave-particle duality, Bohr, Schrodinger, Heisenberg, and others developed a quantum theory in the 1920's as a precise wave theory to describe the motion of electrons, previously only thought of as particles. Within the framework of quantum physics, both photons and electrons are examples of wave-particles. It is therefore not surprising then that there are many similarities in the theoretical treatment of photons and electrons. Both photons and electrons obey a wave equation; a high dielectric constant for photons is mathematically similar to a low electric potential for electrons.

The field of photonic crystals blossomed upon the realization that the formalism and intuition of solid state physics, traditionally exploring the electronic properties of crystalline (periodic) solid materials, can also be effectively applied to understand and design periodic dielectric materials. The connection was first exploited by Yablonovitch, who created a three-dimensional photonic crystal for microwaves [193]. It is important to remember that the field of photonic crystals did not invent the use of periodic dielectric structures. Photonic crystals have long been used in the form of dielectric stack mirrors, and some photonic crystals can even be naturally found in opal crystals and abalone shells. Rather, the field of photonic crystals applies mature electronic solid-state formulations to understand and design periodic dielectric structures. This approach is particularly well suited for the design of two or three dimensionally periodic photonic structures. In order to properly understand photonic crystals, it is important to recognize some similarities and differences between photonic crystals and their electronic counterparts.

A crystal is defined in the ideal case to be, “constructed by the infinite repetition of identical structural units in space” [194]. Electronic crystals are materials with a spatially periodic electronic potential. In most cases, electronic crystals are defined by a lattice of charged ions (assumed to be stationary to first order). A rough schematic of a typical elec-
Electronic crystal, the diamond structure, is shown in Figure 5.1(a). Each circle represents the location of a charged ion, and the bars connecting the ions schematically represent the electronic bonds between ions. The electron eigenstates in such a crystal obey an effective Schrödinger Equation:

$$\left[ -\frac{\hbar^2}{2m} + V_{\text{eff}}(\mathbf{r}) \right] |\psi(\mathbf{r})\rangle = E|\psi(\mathbf{r})\rangle$$  \hspace{1cm} (5.1)$$

where \( \mathbf{p} \) is the momentum operator, \( V_{\text{eff}}(\mathbf{r}) \) is the effective potential operator, \( E \) is the eigenenergy of a given state, and \( |\psi(\mathbf{r})\rangle \) is corresponding electron wavefunction. For a crystal, \( V_{\text{eff}}(\mathbf{r}) = V_{\text{eff}}(\mathbf{r} + \mathbf{R}) \) for all Bravias lattice vectors \( \mathbf{R} \).

Photonic crystals are materials with a spatially periodic dielectric constant. Whereas electronic crystals are commonly found in nature, photonic crystals almost always man-made. A depiction of an imagined 3-dimensional photonic crystal is shown in Figure 5.1(b). Each color represents a material with a different index of refraction. The blank space interlaced within the material represents vacuum or air. This particular structure was designed to be fabricatable by lithography [195]. The equation governing the allowed photon eigenstates in a generalized dielectric material is derived from Maxwell’s equations to be [190]:

$$\nabla \times \left( \frac{1}{\varepsilon(\mathbf{r})} \nabla \times \mathbf{H}(\mathbf{r}) \right) = \left( \frac{\omega}{c} \right)^2 \mathbf{H}(\mathbf{r})$$  \hspace{1cm} (5.2)$$
where $\varepsilon(\mathbf{r})$ is the position dependent dielectric constant, $\omega$ is the photon frequency, $c$ is the speed of light in a vacuum, and $\mathbf{H}(\mathbf{r})$ is the position dependent magnetic field. For a photonic crystal, $\varepsilon(\mathbf{r}) = \varepsilon(\mathbf{r}+\mathbf{R})$ for all Bravias lattice vectors $\mathbf{R}$.

Because of the mathematical similarities between the two eigenvalue equations (5.1) and (5.2), much of the intuition that has been developed in the study of solid state physics can be applied to understanding and designing photonic crystal structures. In particular, solutions to these two equations are often compiled in a plot of the bandstructure of the material. Typical electronic and photonic bandstructures are shown in Figures 5.2 (a,b). Figure 5.2 (a)

![Figure 5.2](image)

Figure 5.2 (a) The electronic bandstructure of Si. (b) The photonic band structure of a three-dimensional photonic crystal.

shows the bandstructure for bulk silicon, while Figure 5.2 (b) shows the photonic bandstructure for a three-dimensional photonic crystal. The bandstructure, or dispersion relation, of a structure relates the frequency of a given eigenstate to the momentum of the same eigenstate. In periodic structures, there is often a range of frequencies for which there are no eigenstates, called a bandgap. Bandgaps occur when the electron or photon wavefunction undergoes Bragg reflection for any propagation direction within the crystal at a given frequency. In each bandstructure plotted in Figure 5.2, the bandgaps are indicated by a gray bar.

Despite the remarkable similarities, there are many fundamental differences between electronic and photonic crystals. Electrons are fermions which propagate according to a complex scalar equation, whereas photons are bosons propagating according to a real vector equation. Also, there is a conservation of the total number of electrons in a system, while photons can be created and annihilated. Finally, the boundary conditions used for electrons in a Schroedinger equation will be different from those used for photons in Maxwell's equations.
5.4 \textit{Light Confinement}

In this work, two methods are used to localize light to a small volume: Bragg reflection from a photonic crystal and waveguiding from a large local index-of-refraction. Bragg reflection from a photonic crystal occurs when the photonic crystal has a PBG for a particular frequency of radiation. Because no optical modes are allowed within the material, the photonic crystal acts as a mirror for that frequency. By surrounding a defect region, which can support an optical mode, with the photonic crystal mirrors it is possible to confine light to a small volume. Waveguide confinement of light is mathematically equivalent to the confinement of electrons in finite potential wells. Light will be confined within regions of high index-of-refraction in much the same way as electrons are confined within regions of lower electronic potential. In the structures studied here, PBG crystals confine light in a one-dimensional Fabry-Perot geometry. The remaining two-dimensional confinement of the optical mode is provided by a high-dielectric-contrast semiconductor waveguide.

It is possible to achieve strong photon confinement using only a PBG crystal. Infinite three-dimensional photonic crystals can be designed to forbid photon propagation for a range of wavelengths in every direction in space. A defect placed in such a cavity could contain an optical mode confined to a small modal volume without any waveguide confinement. Three-dimensional crystals, however, are extremely difficult to fabricate. Furthermore, efficient coupling of light into and out of a three-dimensional PBG microcavity is non-obvious. Nevertheless, significant progress has been made in the fabrication of these structures. Several different three-dimensional photonic crystal geometries have been suggested [195 - 200] and successfully fabricated [201]. Three-dimensional PBGs have been measured optically in the near infrared for nanofabricated [202 - 206] and self-assembled [207] structures. More recently it has been shown that it is possible to create one-dimensionally periodic structures which reflect all incident radiation modes regardless of direction (called omnidirectional mirrors). [208 - 210]

Photonic crystals with one-dimensional periodicity embedded within an optical waveguide provide an attractive approach to achieve photon confinement in dielectric structures [211]. Strong three-dimensional confinement can be provided by the combination of a one-dimensional photonic crystal and a high-dielectric-contrast waveguide. The optical waveguide can be extended to both sides of the photonic crystal to couple light into and out of the photonic crystal microcavity. Waveguide platforms allow for the integration of such photonic crystals with other microphotronics devices.
There are some fundamental drawbacks to confining light with waveguiding. Unlike light confined within infinite three-dimensional photonic microcavities, light confined by a high index waveguide will couple into radiation modes when perturbed by the 1-D photonic crystal, and thus have a finite lifetime within the optical cavity [186, 212]. Smaller modal volume optical states within a index-guided structure generally have shorter lifetimes because of increased phasematching into radiation. The increased phase-matching into radiation is a result of the increase in the numerical aperture of light as the optical modal volume is decreased. It is important to design these structures to balance a high cavity quality factor (Q), small modal volume, and low radiation losses.

Small modal volume optical microcavities consisting of one-dimensional PBG and two-dimensional waveguide confinement have been studied in a variety of material systems and geometries. Microcavities in the monorail and airbridge waveguide geometries, described in the remainder of the chapter, are studied in the GaAs/Al$_x$O$_y$ material system in this work [213 - 216]. Monorail structures have been previously studied in Si/SiO$_2$ [217] and InGaAsP [218]. Microcavities with a different geometry one-dimensional photonic crystal have been studied using a periodic array of grooves by Krauss et. al. [219].

Several other approaches toward creating microcavities use combinations of waveguide and photonic crystal confinement. Very small vertical cavity emitting lasers (VCELs), for example, combine a one-dimensional photonic crystal Bragg stack with a truncated two-dimensional planar waveguide [220]. Microring disk [221] and microsphere [222 - 224] resonators confine light in the whispering gallery modes of a cylindrically symmetric dielectric structure. Ring resonators use a circular waveguide to confine light [225, 226]. Finally, microcavities using one-dimensional planar waveguide and two-dimensional photonic crystal confinement have been fabricated in many systems [227 - 231], and have been used as laser cavities [232, 233].
5.5  Microcavity Design

Photonic bandgap microcavities have been fabricated and studied in a monorail and airbridge geometry. A schematic of each structure is shown below in Figure 5.3. In both

![Schematic of microcavity design](image)

Figure 5.3  Schematic of a (a) monorail and (b) airbridge geometry photonic crystal microcavity. The microcavities consist of a defect region of length $a_d$ between one-dimensional photonic crystal lattice of air holes of spacing $a$ and diameter $d$ within a GaAs waveguide of width $w$ and thickness $T_{GaAs}$. The waveguides rest upon a $Al_xO_y$ cladding layer of thickness $T_{Oxide}$ on top of a GaAs substrate.

in cases, a dielectric waveguide is used to confine light along two dimensions, while a one-dimensional photonic crystal is used to confine light along the third. A GaAs waveguide of width $w$ and thickness $T_{GaAs}$ rests on top of an $Al_xO_y$ layer of thickness $T_{Oxide}$. GaAs has an
index-of-refraction of 3.39 at 1.5 μm, while Al₅O₃ has an index-of-refraction of 1.61. The typical waveguide dimensions are w = 600 nm and T_{GaAs} = 200 nm. These dimensions are chosen to allow a single TE-like waveguide mode, while cutting off both the first order TM-like mode and the second order TE-like mode. The waveguides extend 0.5 mm to 0.75 mm on both sides of the photonic crystal for efficient coupling into and out of the microcavity, and the waveguides' widths were flared to 3 μm over a distance of 100 μm near the input facet to improve coupling of light into the waveguides. A plot of the electric field intensity of the single optical mode in a 200 by 500 nm waveguide is plotted below in Figure 5.4. Because of the high index contrast between the GaAs waveguide and the Al₅O₃ and air cladding, the optical mode is tightly confined within the waveguide.

Both the Al₅O₃ layer, typically of thickness T_{oxide} = 3 μm, and the GaAs waveguide are grown on top of a GaAs substrate. The Al₅O₃ thickness was chosen to minimize loss from the waveguide into the substrate. The photonic crystal is defined by an array of eight holes through the waveguide. The holes have a diameter d and a center-to-center spacing of a. In a
typical sample, \( d = 200 \text{ nm} \) and \( a = 500 \text{ nm} \). A change in the lattice spacing between the two center holes creates an optical microcavity. The center-to-center spacing between these two holes is denoted by \( a_d \). In order to create an optical mode near 1.55 \( \mu \text{m} \), \( a_d \approx 650 \text{ nm} \). The photonic crystal dimensions were varied from device to device to control the PBG and resonance wavelengths, while maintaining small modal volumes and high peak transmission. In the monorail geometry, the microcavity completely rests upon the low refractive index \( \text{Al}_x\text{O}_y \) material, while in the airbridge geometry, the microcavity is suspended in air. Monorail geometry devices, with an oxide layer remaining intact under the photonic crystal hole, are easier to fabricate than airbridge structures, whereas airbridge structures are expected to exhibit better performance by further isolating the cavity mode from the substrate.

A schematic of GaAs waveguide, with no photonic crystal, and the guide's dispersion relations are shown in Figure 5.5 (a) and (b) [186]. The solid lines in Figure 5.5 (b) correspond with guided modes. The modes are labeled TE-like (electric field primarily parallel to the substrate) and TM-like (electric field perpendicular to the substrate). This labeling convention was chosen despite the fact that TE and TM modes are not strictly defined in strip waveguides with strong field confinement. Quantum wells placed at the center of the waveguide will couple to TE-like modes. The shaded region of the band diagram represents the existence of continuum radiation modes. The light line separating the region with and without radiation modes is defined by \( \omega = c/n \), where \( n \) is the index of refraction of the underlying cladding material. Light will preferentially couple into substrate over free space modes because of the substrate's the higher index of refraction.

Above some cut-off frequency, the waveguide modes in the band structure shown in Fig. 5.5 (b) do not have an upper bound to their wavevector. The introduction of a periodic array of holes into the waveguide folds the dispersion relations into the first Brillouin zone for wavevectors greater than \( \pi/a \), and splits the guided-mode bands at the fold of the Brillouin zone, as shown in Figs. 5.5 (c) and (d). The mode of the lower energy band, called the dielectric band, has electric field primarily in the dielectric material, while the mode of the higher energy band, called the air band, has a larger proportion of its electric field in the air holes. Both of these bands correspond with guided modes which exist in the waveguide despite the presence of holes. In between the air band and dielectric band is a region without any modes allowed in the structure. This frequency span without waveguide modes is known as the bandgap.
When a defect is introduced in the periodic array of holes, a resonant state can be created inside the bandgap. This defect state is strongly confined within the defect area, and has an exponentially decaying electric field into the photonic bandgap crystal to either side. The modal volume of the optical resonant state, $V_m$, is defined to be:

$$V_m = \int \frac{|\varepsilon (\mathbf{E}^*) \cdot \mathbf{E}| d^3 r}{|\varepsilon (\mathbf{E}^* \cdot \mathbf{E})|_{\text{max}}}$$

(5.3)
where $\varepsilon$ is the dielectric constant, $\varepsilon(E^*E)$ is the energy density in the electric field, and $|\varepsilon(E^*E)|_{\text{max}}$ is the peak value of $\varepsilon(E^*E)$. Using this definition, the minimum modal volume calculated for the airbridge and monorail structures is approximately $2(\lambda/2n)^3$.

The defect state has a finite lifetime determined by the coupling rate to guided modes in the waveguide (the desired decay mechanism) and coupling to radiation modes (considered loss). Minimizing the coupling between the defect mode and the radiation continuum is necessary to reduce the overall loss through the defect. The total quality factor of the resonant mode, $Q_{\text{tot}}$, is defined to be the ratio of the optical energy stored within the microcavity and the total cycle-average power radiated out of the cavity. It can be experimentally determined by the relationship $Q_{\text{tot}} = \lambda/\Delta\lambda$, where $\Delta\lambda$ is the width of the resonance and $\lambda$ is the peak wavelength of the resonance. $Q_{\text{tot}}$ can be broken down into waveguide and radiation coupling components with the relationship:

$$\frac{1}{Q_{\text{tot}}} = \frac{1}{Q_{\text{wv}}} + \frac{1}{Q_{\text{rad}}}$$  \hspace{1cm} (5.4)

where $Q_{\text{wv}}$ is the quality factor of coupling to waveguide modes and $Q_{\text{rad}}$ is the quality factor of coupling to radiation modes. A finite-difference time-domain computational scheme [234] was used by P. R. Villeneuve et. al. [235] to compute $Q_{\text{tot}}$ and the transmission through the structure. Typical results are shown in Figure 5.6 (a) for monorails and Figure 5.6 (b) for airbridge microcavities. In both cases, the computation shows a wide bandgap from 1400 nm to

![Figure 5.6](image_url)

**Figure 5.6** Finite-difference time-domain calculation of transmission through a (a) monorail and (b) airbridge geometry waveguide microcavity. In each case, a resonant defect state is designed to be at 1550 nm (*from reference [235]*).
1700 nm, and a single sharp resonant peak near 1550 nm. The resonant peak at 1550 nm is due to coupling into and out of the optical mode in the photonic bandgap microcavity. Transmission outside the bandgap is large which suggests that the modes remain guided as they propagate through the holes, and undergo minimal scattering. On resonance at 1550 nm, both the monorail and the airbridge microcavity devices exhibit high transmission. The coupling from the waveguide mode to the cavity resonant mode occurs via the evanescent field through the array of holes.

Electromagnetic propagation simulations have been performed to plot the electric field energy inside the waveguide and microcavity in different regions of the transmission curve [211]. The energy in the electric field above the bandgap and within the bandgap are shown below in Figure 5.7. Above the bandgap, good optical modes exist within the photonic crys-

![Electric field intensity](image)

Figure 5.7 Electric field intensity inside a waveguide microcavity outside and within the bandgap of the photonic crystal (*from reference [235]*).

tal, and incident light is transmitted through the crystal region. Light with a frequency within the bandgap, however, cannot propagate through the photonic crystal and is reflected backwards. A standing wave pattern is visible due to the interference of the incident and reflected beam. Amazingly, it is possible to reflect almost all of the incident light with only 3 or 4 periods of the photonic crystal because of the extremely high dielectric constant contrast between GaAs and air.
At Resonance

Figure 5.8 shows the electric field for incident light with a wavelength on resonance with the optical microcavity detector. In this case, an electric field mode builds up within the detector region. The electric field energy level for the incident light is the same intensity as for Figure 5.7, but has been scaled down to exhibit the high field energy of the resonant state. The electric field within the microcavity has a dipole nature, while the corresponding magnetic field has a lowest order mode profile with no nodes. The transmission through the photonic crystal cavity is not 100%, because of some loss to radiation and because the simulated incident optical pulse contains wavelengths that are not on resonance with the cavity.

In the structure modeled, the electric field is tightly confined within a small volume near the detector. One way to improve optical confinement is to increase the number of holes on either side of the microcavity. By increasing the number of holes, the reflectivity of the hole array is increased and $Q_{cav}$ is increased. However, because $Q_{det}$ is finite in the structure, $Q_{cav}$ does not increase with $Q_{det}$. Rather, as the modal confinement is increased, con...
pling to radiation modes will increase and eventually dominate over the preferred coupling with guided modes inside the waveguide. Coupling to radiation increases with increased spatial confinement because of the increased numerical aperture of the optical mode. The peak transmission through the cavity can be expressed as:

\[ T_{\text{max}} = \left( \frac{Q_{\text{tot}}}{Q_{\text{rad}}} \right)^2 \]  

(5.5)

A plot of the calculated maximum transmission and \( Q_{\text{tot}} \) of a monorail structures with different numbers of holes is shown in Figure 5.9 [235]. As the number of holes is increased, the cavity \( Q_{\text{tot}} \) increases. Correspondingly, the transmitted intensity decreases as loss to radiation begins to dominate.

Loss to radiation can be lowered by eliminating preferred radiative paths. The substrate material upon which the microcavity rests provides a favorable escape route for radia-
tion loss. Radiation loss can therefore be minimized in airbridge microcavities compared to monorail microcavities by removing the underlying material [236].

Altering the dimensions of the features in the microcavity structure affects the positions of the band edges and the resonant wavelengths in the transmission spectra. For example, lengthening the defect region will increase the resonant wavelength of a microcavity, and increasing in the thickness of the waveguide primarily shifts the band edges to longer wavelengths. The high-index waveguide provides strong field confinement in the vertical and lateral dimensions such that the guided modes extend only weakly outside the waveguide, allowing a large fraction of the guided modes to interact with the photonic crystal. Strong field attenuation through the array of holes is necessary to achieve small modal volumes.

It is useful to calculate the theoretical maximum spontaneous emission enhancement that an emitter would experience within the optical cavity. By coupling an optical emitter to the microcavity resonance, the spontaneous emission rate can be enhanced by a maximum factor of \( \eta \) compared to the rate without a cavity. The expression for \( \eta \) is given as [237]:

\[
\eta = \frac{Q}{4\pi V_m} \left( \frac{\lambda}{n} \right)^3
\]

(5.6)

where \( \lambda \) is the optical transition wavelength. A number of averaging factors must be taken further into account to find the actual spontaneous emission enhancement for a specific emitter placed within the cavity [238]. Other cavity geometries in one-dimensional PBG's could also realize spontaneous emission enhancements [181]. For the airbridge waveguide microcavity that has been measured, the maximum enhancement is calculated to be 72, which is significantly larger than any enhancement yet measured. This large spontaneous emission enhancement could lead to faster modulation of optical devices, and to the development of zero-threshold lasers.

### 5.6 Device Fabrication

The devices studied here are enabled by the remarkable fabrication that was necessary to produce them. Details of the device fabrication can be found in the Ph.D. thesis of Dr. Kuo-Yi Lim from the group of Professor Leslie Kolodziejski [171]. The basics of the device fabrication are discussed below.

AlGaAs and GaAs layers were deposited on a GaAs substrate by gas source molecular beam epitaxy (GSMBE). The thickness and Al content of the AlGaAs layer were measured to
be 3 μm and 93% respectively. The GaAs layer thickness was measured to be 185 nm for the monorail structure and 181 nm for the airbridge structure. An SiO₂ layer was deposited on the GaAs/Al₀.₉₃Ga₀.₀₇As heterostructure using PECVD, and was followed by a spun-on coating of polymethylmethacrylate (PMMA), an electron-beam resist. The devices were then patterned in the PMMA using direct-write electron-beam lithography. Sixteen electron-beam exposure fields, each approximately 100 μm by 100 μm, were stitched end-to-end in order to form the overall device structure. After the exposed resist was developed, a layer of nickel was deposited on the sample by electron-beam evaporation. A subsequent lift-off of the nickel film resulted in the image-reversal of the electron-beam written patterns. Thereafter, the nickel film acts as a mask for the etching of the SiO₂ layer by RIE (reactive ion etching). After the RIE process, the nickel mask was stripped off and the sample was backside-lapped down to approximately 150 μm.

An additional RIE process transferred the device patterns in the SiO₂ mask into the GaAs/Al₀.₉₃Ga₀.₀₇As heterostructure. The RIE process etches through the top GaAs layer and about 400 nm into the Al₀.₉₃Ga₀.₀₇As layer. The process sequences for the monorail and the airbridge microcavities diverged at this point. For the monorail microcavity, the SiO₂ mask layer was removed by RIE and the sample was oxidized for 30 min using the furnace and procedure described below. The holes are thought to extend into the oxidized layer.

For the airbridge microcavity, a photolithography step was performed to define a 10 μm wide trench pattern with the photonic crystal in the center of the trench. An RIE step then etches an additional 800 nm into the Al₀.₉₃Ga₀.₀₇As layer in order to define the trench region and remove the bulk of the sacrificial material. The resist and the SiO₂ mask were then removed using RIE. Thereafter, the Al₀.₉₃Ga₀.₀₇As layer was oxidized using the same procedure as for the monorail structure, discussed below. Another photolithography step was performed to redefine the same 10 μm wide trench pattern. The sample was then dipped into a diluted hydrofluoric acid (HF) solution that selectively etched the AlₓOᵧ beneath the photonic crystal to suspend the airbridge structure.

Finally, facets of the input and output waveguides in both the monorail and airbridge structures were cleaved to facilitate the coupling of light into the devices during optical characterization.

The thermal oxidization of AlGaAs layers with high Al content into AlₓOᵧ provides the necessary dielectric contrast for the operation of GaAs-based photonic crystals. The index of refraction of AlₓOᵧ is 1.61 at 1.55 μm. Oxides have been developed in particular for VCEL mirrors and other applications [172 - 174, 239 - 244]. The AlGaAs oxidization was per-
formed in a single zone, quartz tube furnace that was maintained at 435°C. Steam was introduced into the furnace by flowing N₂ through a water bubbler that was maintained at 90°C. The oxidization time was defined as the time that the sample was exposed to the steam. To improve the stability of the oxide, the AlGaAs composition was chosen to be Al₀.₉Ga₀.₁As as opposed to pure AlAs.

Scanning electron micrographs of both the monorail and airbridge microcavity structures are shown in Figure 5.10 (a) and (b) respectively.

(a)  
(b)  

Figure 5.10 Scanning electron micrograph (SEM) of a (a) monorail and (b) airbridge waveguide microcavity.
5.7 Experimental Setup

The fabricated PBG waveguide microcavities were studied using optical transmission spectroscopy. The primary challenges in performing this measurement were the difficulty of coupling light into waveguides with dimensions of only 500 by 200 nm (smaller than the free-space wavelength), the high transmission loss of the waveguides, and the large tuning range necessary to fully characterize a PBG device with a > 300 nm bandgap. An experimental set-up was developed to make waveguide transmission measurements despite these difficulties, and is shown in Figure 5.11.

![Experimental set-up designed to study the transmission through waveguide devices.](image)

A tunable continuous wave NaCl:OH⁻ laser was chosen as an optical source for its high output power and broadband tunability [23]. A schematic of the laser cavity is shown in Figure 5.12. The color center laser was pumped by a 4 W cw Nd:YAG laser operating at 1064 nm. A 300 mW Argon ion laser was also used to perform optical pumping, and prevent carrier buildup in dark states. The laser crystal itself lases due to transitions in F₂⁺ color center defects, electrons localized near a defect consisting of the absence of two negative ions. Because NaCl is hydroscopic and the F₂⁺ color-center defects will disappear at room temperature, the crystal must be mounted into a liquid nitrogen filled cryostat.
Figure 5.12  Schematic of a NaCl:OH\(^+\) color center laser used as an optical source for transmission measurements.

The laser wavelength was tuned with a birefringent plate on a rotary stepper stage, which provided a tuning range from 1500 nm to 1670 nm, a linewidth of approximately 0.1 nm, and a maximum average power ~ 250 mW. A typical calibration curve relating the laser wavelength to the birefringent tuning plate stepper motor position is shown in Figure 5.13. The relationship between wavelength and birefringent plate position is repeatable for multiple rotations of the plate. A tuning curve was obtained by monitoring the laser power for different wavelengths, and is plotted in Figure 5.14. The laser is tunable from 1500 to 1670 nm.
Figure 5.13  Calibration of the color center laser output wavelength to birefringent tuning plate stepper motor position.

Figure 5.14  Color center laser tuning curve.
During transmission measurements, fast wavelength tuning was used to acquire data at a rate of over 5 scans/minute, with a data point density of ~ 45 data points/nm. The laser light was coupled into an optical fiber, 3% of which was tapped off into a photodetector to monitor the laser power. The remaining 97% of light in the fiber was directed to a fiber-lens assembly used to couple into a waveguide device. Polarization was controlled with a polarizing beam splitter cube, a half-wave plate and a quarter-wave plate, all placed before the fiber-lens assembly. All measurements were taken with light exiting the fiber-lens assembly TE polarized, as TM light was not guided within the waveguide. After transmission through the waveguide, output light from the waveguide was imaged with either a microscope objective or an aspheric lens. The output is imaged through a 100 \( \mu \text{m} \) pinhole in a confocal geometry in order to suppress substrate guided light. A photodetector was placed after the pinhole to monitor the output power. Both the input and output power were measured with lock-in detection.

### 5.8 Waveguide Coupling and Loss

Efficient waveguide coupling was extremely difficult due to the sub-wavelength scale and asymmetric shape of the high-dielectric contrast waveguide. In these experiments, a fiber pig-tailed lens assembly with a 300 \( \mu \text{m} \) working distance was used to match the optical mode of the high-dielectric waveguide. Fibers with lensed ends were also found to be suitable for coupling into these waveguides. The waveguides themselves were flared from 500 nm in width to 3000 nm in order to more closely match the diffraction limited spot-size in air. Some samples had no flared input facets, and no discernible difference in the input coupling efficiency was observed.

Infrared vidicon images of both a bird’s-eye view of the waveguide and an end view of the output facet were monitored to optimize coupling. Typical bird’s-eye view infrared images of a coupled PBG-less waveguide are shown in Figure 5.15. Each picture images an area approximately 0.25 mm wide by 0.5 mm long. The top image views the input coupling facet of the waveguide. Light is significantly scattered at the air-waveguide interface due to the large mode mismatch between the lens-assembly focused beam and the high-dielectric contrast waveguide mode. Some light is then coupled into the waveguide, where it undergoes severe scattering due to surface roughness inside the waveguide. Optical coupling into the waveguide can be easily optimized by monitoring the intensity of light scattered from the waveguides. It was more difficult to monitor the output power of the waveguide because the large bandgap of the PBG microcavities limited the output power. Once light was initially coupled into the waveguides, the transmitted light power could then be found and optimized.
A typical infrared vidicon camera bird’s eye view of the waveguide-air output interface, as well as a head on view of the imaged output facet are shown in Figure 5.16. The output is imaged with a 10x microscope objective onto the infrared vidicon camera. A procedure has been developed for achieving and optimizing waveguide coupling while imaging the output end view. First, care must be taken to focus upon the output plane of the waveguide, rather than the input plane. Because the pump light is focused on the input plane, it is easy to accidentally image the coupling lens’ focus. Next, the input fiber is translated above and below the air-semiconductor interface. The semiconductor chip will act as a knife edge, and generate a characteristic diffraction patterned, that is used to place the optical input focus.
exactly at the semiconductor-air interface. The optical input beam is then displaced side to side, until a waveguide is found. Once the waveguide is found, the waveguide output is imaged through a pinhole. In the experiments reported below, a 100 μm diameter pinhole was used. The pinhole is used, as in confocal microscopy, to spatially filter the residual substrate guided light (imaged from a focal plane at the waveguide input) while passing the waveguided light (imaged from a focal plane at the waveguide output). There is a trade-off in signal to noise between using longer waveguides, which allow a high rejection of substrate guided light, and shorter waveguides which have lower total insertion losses.

Figure 5.16 Bird’s eye and output end views of light output from a high-dielectric contrast waveguide as imaged by an infrared sensitive vidicon camera.

The losses through straight waveguides without a photonic crystal were estimated to be 3-6 dB/mm. The intensity of the light scattered out of the waveguide as a function of position was used to estimate the linear loss. The dominant loss mechanisms appeared to be scattering from sidewall roughness and the GaAs/AlxOy interface, which are both accentuated by the tight optical confinement in the high-dielectric-contrast waveguide system. The sidewall roughness resulted from the roughness on the mask that was transferred into the GaAs waveguide. The GaAs/AlxOy interface roughness results from the AlGaAs layer growth and oxidation. Other possible loss mechanisms include the error in aligning two successive electron-beam lithography fields together and the intrinsic material loss of the waveguide material.
at 1.55 μm. Based on the total estimated loss, waveguide input coupling efficiencies of ~3% were calculated.

At high input powers, two-photon absorption within the GaAs waveguides has also been observed. Two-photon absorption excited photoluminescence from the GaAs waveguide was observed by Si-CCD array camera images shown in Figure 5.17. The GaAs semiconduc-

![Images showing two-photon absorption (TPA) pumped photoluminescence (PL) within a high-index contrast waveguide.](image)

Figure 5.17 Silicon CCD camera images of two photon absorption (TPA) pumped photoluminescence (PL) within a high-index contrast waveguide. Images are taken (a) without and (b) with light coupled into the waveguide. The CCD camera is sensitive to PL from GaAs, but not to the input light at 1.55 μm.

tor chip containing the optical waveguides is a black square in the middle of the picture. The metallic barrel coming from the bottom-left of the picture toward the waveguide chip is the fiber lens assembly used for input coupling. In the Figure 5.17 (a), no light is coupled into the optical waveguide, whereas in Figure 5.17 (b), intense 1.5 μm light is input into the waveguide. The waveguide is seen to be illuminated near the input interface. Because this light is detected by a Si CCD camera which is not sensitive at 1.5 mm, it is likely that the light is due to photoluminescence (PL) from GaAs. Placement of wavelength-dependent filters in front of the Si CCD camera confirmed this conclusion. TPA was made possible by the tight optical confinement of the high-dielectric contrast waveguide and could limit waveguide use for wavelength division multiplexing (WDM) applications.
5.9 Optical Microcavities

Optical transmission through several monorail and air-bridge microcavity devices has been measured. Information about each device is tabulated in Table 5.1. Transmission spec-

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<th>Monorail Devices</th>
<th>Airbridge Devices</th>
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<td>1 2 3 4</td>
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<td>(Hole center to Hole center)</td>
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<tr>
<td>Measured Q&lt;sub&gt;tot&lt;/sub&gt;</td>
<td>Q</td>
<td>136 142 117</td>
</tr>
<tr>
<td>Resonance Wavelength (nm)</td>
<td>λ&lt;sub&gt;res&lt;/sub&gt;</td>
<td>1522 1536 1566</td>
</tr>
</tbody>
</table>

* Values measured by scanning electron micrograph (SEM).
† Values measured by ex-situ variable angle spectroscopic ellipsometry.
All other dimensions are design parameters, and have not actually been measured.

Table 5.1 Dimensions of fabricated waveguide microcavity devices.

tra of three different monorail microcavities (monorail devices 1-3 in Table 5.1) are shown in Figure 5.18. Each device had a different defect length (defined as the distance between the centers of the holes neighboring the defect region). As expected, larger defect lengths support longer wavelength resonant modes. The cavity modes had cavity Q's of 136, 142, and 117 respectively. Due to the lack of a precise loss measurement and the lack of a full band edge within the experimental window, the maxima of the peaks in Figure 5.18 were normalized to unity rather than to the absolute transmission.
Figure 5.18 Transmission spectra through three monorail devices. Each monorail has a different defect length, resulting in a different transmission peak wavelengths. The maximum transmission of each peak is normalized to unity.

The normalized experimental transmissions through two distinct airbridge microcavities (airbridge devices 1 and 2 in Table 5.1), having defect center-to-center lengths of 632 nm and 703 nm respectively, are shown in Figure 5.19. The resonance states had Q's of 310 and 336 respectively. Resonance Q's as high as 360 have been recorded in airbridge device 2. As in the case of the monorail waveguides, the resonance shifted to longer wavelengths as the microcavity size was increased. Because the PBG was calculated to be over 300 nm wide, the band edges were not visible within the 180 nm tuning range of the laser. The plots in Figure 5.19 were normalized by the same method as the monorail structures in Figure 5.18. Both of the transmission spectra in Figure 5.19 were averages of 10 scans taken back-to-back, and had a wavelength resolution close to the laser bandwidth. These two devices had a modal volume of 0.031 μm³. Smaller modal volumes are predicted to lead to increased loss due to radiation
modes in waveguide based devices. However, the airbridge microcavities have higher Q's than monorail microcavities of the same modal volume, because of the increased isolation of the resonance mode from substrate radiation modes. Using a modal volume of 0.031 μm³, a Q of 310, and a peak wavelength of 1521 nm, the maximum enhancement factor for spontaneous emission from such a defect is computed to be η = 72.

![Transmission Spectra](image)

Figure 5.19 Transmission spectra measured through two different PBG air-bridge microcavities. The long and short wavelength resonances have defect center-to-center lengths of 632 nm and 703 nm respectively. The maxima of the resonance peaks are normalized to unity.

A top-view image of an illuminated microcavity is shown in Figure 5.20. This picture was taken with an infrared vidicon camera. Some of the light being transmitted by the waveguide was scattered by surface roughness, and was visible to the camera. Off resonance, the light stopped at the photonic crystal located in the center of the guide and was reflected as...
shown in Figure 5.20 (a). On resonance, the incident light was transmitted, and the portion of the waveguide after the photonic crystal was illuminated as shown in Figure 5.20 (b).

![Figure 5.20](image)

**Figure 5.20** Top view of an illuminated airbridge microcavity resonator (a). Off resonance and within the bandgap, the photonic crystal reflects the light. (b) On resonance, light is transmitted through the microcavity resonator. The light is scattered out of the waveguides primarily from surface roughness, and is viewed with an infrared Vidicon camera.

By changing the photonic crystal dimensions, a photonic band edge was observed within the laser tuning range. One device with a band edge at 1620 nm (airbridge device 3 in Table 1), is shown in Figure 5.21. This spectrum was an average of 10 scans taken back to back, and had a wavelength resolution close to the laser bandwidth. The bumpy features on the bands edge trace are real and repeatable. This band edge, plotted logarithmically, had 27 dB attenuation of transmission for wavelengths within the bandgap compared to outside of the bandgap which agrees well with what is expected from calculations for the long wavelength band edge. The transmission floor within the bandgap was a true feature of the device measured, and was consistent with theoretical simulations.
Figure 5.21 Transmission spectrum of a PBG air bridge sample where the long wavelength band edge has been shifted into the experimentally accessible wavelength range. The band edge is shown on a log plot, and exhibits a 27 dB suppression of transmission within the bandgap compared to outside the bandgap.

A second device (airbridge device 4 in Table 5.1) exhibited both a band edge and the resonance within the laser tuning window Figure 5.22. Because the band edge and the resonance were both fully visible at the same time, it is possible to determine the transmission of the microcavity resonance relative to that outside of the bandgap. This yielded a maximum relative transmission of 72% on resonance, with a Q of 230. The modal volume of device 4 was calculated to be 0.026 μm³, which is only 2 (λ/2n)³. The lower Q is result of the resonance being shifted closer to the band edge.
Figure 5.22  Transmission spectrum of an airbridge microcavity exhibiting both a resonance and a band edge within the experimentally accessible wavelength range. Transmission outside the bandgap is normalized to unity, revealing 72% relative transmission on resonance, with a cavity Q of 230.

5.10  Future Work

Photonic bandgap microcavities with small modal volumes and large Q's have been demonstrated in a GaAs/Al_xO_y waveguide system. Theoretically, spontaneous emission enhancements as high as 72 should be possible from the measured cavities. By constructing a similar microcavity in an active material such InP with an InGaAs quantum well, it should be able to create a low threshold microlaser. The primary difficulty in demonstrating such a
device is the large amount of nonradiative loss introduced by surface states from the holes and waveguide sides. The electric field within the microcavity has been modeled using computation. A direct measurement of the electric field can not be obtained by conventional microscopy because of the sub-wavelength size of the structure. Near-field scanning optical microscopy (NSOM) could be used to directly measure the field inside a sub-wavelength microresonator [245]. Because the microcavity resonance wavelength is more sensitive to the device dimensions then the fabrication precision, it is difficult to successfully fabricate a device for a specific wavelength. A tunable structure could allow a device to be fabricated, and then tuned to the desired wavelength. Micro-electro-mechanical (MEMs) tuning elements could be constructed near the structure to achieve such tuning of the operating wavelength. Furthermore, some form of MEMs might be used to switch the device on and off.

5.11 Conclusions

One-dimensional photonic bandgap crystals have the potential to control the propagation and spontaneous emission of light, and have applications for efficient, sub-micron scale opto-electronics. Resonant microcavities with small modal volumes will serve as the basis of a low threshold laser cavity with fast modulation rates. Air bridge and monorail geometry microcavities with modal volumes as low as 0.026 μm³ have been designed, fabricated, and measured in one-dimensional photonic bandgap structures. The structures were designed to have high Q's, taking into account of the trade-off between increased optical confinement and increased radiation loss. Waveguide microcavity structures were measured to have Q's as high as 360 in the 1.55 μm regime. Furthermore, the photonic crystals have been fabricated in the GaAs/AlₓOᵧ III-V compound semiconductor system, which provides the potential for fabrication of active devices and allows integration of photonic crystal and active semiconductor technologies.
CHAPTER 6: PHOTONIC CRYSTAL LIGHT EMITTING DIODE

6.1 Introduction

Enhanced coupling to vertically emitted radiation modes was obtained from a light emitting diode (LED) using a two-dimensional photonic crystal that lies entirely inside the upper cladding layer of an asymmetric quantum well structure. For a particular photonic crystal quantum well structure emitting at 980 nm, an eightfold integrated enhancement and a hundredfold single wavelength spectral enhancement in collected photoluminescence were observed. Similar photonic crystals are also used to enhance optical pump input coupling, leading to an overall enhancement in the LED quantum efficiency.

6.2 Background

Light emitting diodes (LEDs) are used for a variety of low-cost lighting applications. Equipment panel indicators, traffic lights and scrolling signs are all commonly illuminated by LEDs. Because of their low cost, long lifetime, and brightness, semiconductor LEDs might become useful for some surprising applications. A blend of multiple color LEDs could be used as an alternative to traditional incandescent light bulbs, especially for large buildings. The basement of a skyscraper might contain a ‘light generator room’, with optical fibers channeling LED emitted white light to each floor of the building. One might also imagine a dining
room illuminated by multiple color LEDs, where the owner is not only able to dim the light level for ‘atmosphere,’ but could even adjust the hue and color of the room lighting by changing relative proportions of red, blue, and green. In addition to lighting, LEDs have potential applications for inexpensive optical communications. Offices might employ short distance fiber optical communications systems based on LEDs as an alternative to more expensive laser diodes.

In order to make these dreams an economical reality, the wallplug efficiency of semiconductor LEDs must be improved. The net wallplug efficiency of an LED is a product of the quantum efficiency (the fraction of pump photons converted into LED emitted photons) and the extraction efficiency (the fraction of emitted photons radiated out of the device into a useful direction). Semiconductor LEDs are limited in wallplug efficiency because of their poor light extraction efficiencies. Much of the light generated within a semiconductor LED is trapped by slab waveguide modes, rather than coupled out of the semiconductor material into free space [246]. Once trapped within the high index waveguide, this light is often reabsorbed by the LED, absorbed by free carriers, or emitted into an unwanted direction.

The extraction efficiency of an LED is highly dependent on geometric factors. Several approaches have been used to increase the emission efficiency of semiconductor LEDs. Geometric configurations such as hemispherical dome structures [247], random corrugation of the LED surface [248, 249], and truncated inverted pyramid structures [250] have all been used to increase the extraction efficiency by reducing total-internal-reflection. One-dimensional microcavities [251], two-dimensional photonic crystal microcavities [252, 253], and two-dimensional photonic crystal Bragg scattering [253 - 257] all improve light output coupling by modifying the photonic density of states. In this work, photonic crystal microstructures are used to increase the extraction efficiency of LEDs. Photonic crystals are patterned into the LED substrate to couple light from the waveguide modes into free space [258]. Finally, an approach using photonic bandgaps is discussed which increases the extraction efficiency by eliminating lossy waveguide modes entirely [259].

The LED experiments described in this chapter were made possible by state-of-the-art fabrication techniques developed by Alexei Erchak in the lab of Professor Leslie Kolodziejski. Dr. Shanhui Fan of Professor John Joannopoulos’ group developed and applied the computational techniques used to model our structures. Many of the experiments have been performed with the help of ingenuity of Peter Rakich. The conclusions drawn are the product of a collaboration among all of the parties involved.
6.3 980 nm Light Emitting Diode

Several InGaAs/InGaP quantum well LEDs were designed to emit at 980 nm in order to study the effect of various geometrical and photonic crystal microstructures on the extraction efficiency of semiconductor LEDs. In the absence of microstructuring, most of the light emitted by such an LED will be trapped by total-internal-reflection in waveguide modes of the high-index InGaP substrate. A schematic of this process is shown in Figure 6.1. Assuming a spherical radiation pattern, and that the semiconductor substrate's index-of-refraction $n >> 1$, the extraction efficiency into a cone of angle $\theta$ from such a structure is:

$$\eta = \frac{NA^2}{4n^2}$$ (6.1)

where $\eta$ is the extraction efficiency, $NA$ is the numerical aperture of the light collection cone given by $\sin \theta$, and $n$ is the substrate index-of-refraction. An InGaP index of $n = 3.2$ and a maximum $NA$ value of 1 yield an extraction efficiency of only 2.5% from the top surface. A mirror fabricated underneath the LED can increase the light extraction efficiency by a factor of 2 by reflecting light escaping out of the bottom surface.

In order to study methods for improving this low extraction efficiency, an LED emitting at 980 nm has been designed and fabricated. A schematic of the test LED is shown in Figure 6.2. The LED consists of an asymmetric In$_{0.51}$Ga$_{0.49}$P/In$_{0.20}$Ga$_{0.80}$As/In$_{0.51}$Ga$_{0.49}$P quantum well on top of an Al$_x$O$_y$ spacer layer, a broadband GaAs/Al$_x$O$_y$ distributed Bragg reflector (DBR), and a GaAs substrate. The quantum well consists of an 8 nm InGaAs quantum well surrounded by a 32 nm InGaP lower cladding layer and an upper InGaP cladding layer with a thickness of either 95 nm or 158 nm, depending on the structure. The quantum well structure is designed to emit light centered at 980 nm at room temperature, with a full-width at half-maximum of around 65 nm.
Figure 6.2  Schematic of a light emitting diode (LED) designed to emit light at 980 nm. A single InGaAs quantum well is sandwiched between InGaP cladding layers. A low index $\text{Al}_x\text{O}_y$ spacer layer separates the high index quantum well region from a 6 period GaAs/$\text{Al}_x\text{O}_y$ distributed Bragg reflector (DBR) mirror. A photonic crystal, or other geometric structure, is embedded within the LED device.

In the photonic crystal structures, holes are defined and etched into the top InGaP region. The quantum well layer is offset asymmetrically between the InGaP cladding layers so that it is possible to etch holes which are deep enough to cause strong Bragg scattering of the guided light without penetrating through the quantum well layer. Holes penetrating through the InGaAs layer would introduce nonradiative recombination surface states, and reduce the overall quantum efficiency of the quantum well. The 0.5 $\mu$m low-index $\text{Al}_x\text{O}_y$ layer is placed below the quantum well structure to optically isolate the quantum well structure from the distributed Bragg reflector (DBR) mirror below and suppress coupling of light from the LED into laterally guided modes in the GaAs layers of the DBR. The GaAs/$\text{Al}_x\text{O}_y$ DBR is very broadband due to the high-dielectric contrast between the two materials. The index of refraction of GaAs is 3.4, while $\text{Al}_x\text{O}_y$ is 1.6. The DBR is designed to exhibit broadband reflection between 800 nm - 1300 nm. A plot showing the emission of the LED and the theoretical reflection of the broadband DBR is shown in Figure 6.3. Each LED is fabricated into a square or circular pedestal, called an LED mesa. The mesas are either 50 $\mu$m by 50 $\mu$m, or 100 $\mu$m by 100 $\mu$m depending on the particular sample. These LEDs are constructed as
mesas to allow side oxidation of AlAs into Al$_x$O$_y$. Each LED mesa is either left unpatterned, or fabricated with a 30 µm by 30 µm microstructured region in the middle of the mesa.

![Figure 6.3](image)

Figure 6.3 Emission of light emitting diode (LED) near 980 nm and calculated reflectance of GaAs/Al$_x$O$_y$ high index contrast distributed Bragg reflector (DBR).

### 6.4 LED Fabrication

The LED structures designed require high-index contrast material systems as well as sub-micron feature sizes. The enabling technologies for the experiments described below are the sophisticated microfabrication techniques that are used to create the LED devices. A brief overview of the fabrication process follows. The LED structure is grown on a GaAs substrate using gas-source molecular beam epitaxy (GSMBE). The Al$_x$O$_y$ layers in the DBR and spacer layer are initially grown as AlAs and AlGaAs respectively. The photonic crystal region is patterned using electron beam lithography. The pattern is transferred into the upper InGaP layer by reactive-ion etching (RIE) using a SiO$_2$ hard mask. Square mesas are defined using photolithography and formed by a plasma RIE step. Finally, the AlGaAs and AlAs layers are thermally oxidized from the edges of the mesa inwards.
The oxidation of AlGaAs and AlAs into $\text{Al}_x\text{O}_y$ is a critical step in the fabrication process. The high index contrast between GaAs and $\text{Al}_x\text{O}_y$ is necessary for the proper functioning of the optical device. $\text{Al}_x\text{O}_y$, however, can not be directly grown by MBE. The approach is to grow the AlAs/AlGaAs by GSMBE, and then change the material composition by oxidation. The oxidation process is highly sensitive to both the water vapor temperature and Ga concentration of $\text{Al}_x\text{Ga}_{1-x}\text{As}$. Plots of the oxidation rates for different temperatures and Ga concentrations are shown below in Figure 6.4 [260]. At high temperatures, the rate of oxidation increases, but the chance of layers lifting off increases as well.

Scanning electron micrographs (SEMs) of LED mesas are shown in Figure 6.5. Figure 6.5(a) shows a view of a 50 $\mu$m diameter circular LED mesa, with a 12.5 by 12.5 $\mu$m square photonic crystal region fabricated in the middle. The image in Figure 6.5(b) shows a side view of a cleaved LED mesa. The patterned active quantum well layer is on top of an $\text{Al}_x\text{O}_y$ spacer layer and GaAs/$\text{Al}_x\text{O}_y$ Bragg mirror stack. The bottom two images (Figure 6.5(c,d)) show close-up angled top view and side view images of a triangular photonic crystal lattice microfabricated into an LED cladding layer.
Figure 6.5 Scanning electron micrographs (SEMs) of fabricated light emitting diode (LED) structures
(a) A 50 μm diameter circular mesa, with a 12.5 by 12.5 μm square photonic crystal region
(b) Side view of LED layers  (c) Angled top view of photonic crystal lattice  (d) Side view of photonic crystal lattice holes

6.5 Experimental Setup

An experimental apparatus was developed to image the photoluminescence (PL) of the LEDs, as well as to measure the PL spectrum at different spatial location on the LEDs. A schematic of the set-up is shown in Figure 6.6. For most experiments, a Spectra Physics Tsu-
nami Ti:Sapphire cw laser emitting light at 810 nm with a power up to 100 mW was used to excite photoluminescence. The pump wavelength is constrained by a couple of factors. The pump wavelength should be significantly shorter than the emission wavelength so that it can be easily spectrally filtered for PL measurements. On the other hand, wavelengths shorter than 650 nm will excite the InGaP cladding layer that is lattice matched to GaAs. While the overall pump absorption efficiency would be increased by InGaP absorption, the pump absorption efficiency would vary from mesa to mesa depending on the volume of cladding material removed to create the photonic crystal pattern. The variable pump absorption would make comparison of the emission from two different photonic crystal LEDs extremely difficult.

![Schematic of experimental set-up](image)

Figure 6.6  Schematic of experimental set-up used to image photoluminescence (PL) from an LED sample and measure PL spectrum as a function of position on the sample.

It was found that pumping at 810 nm limited the study of emission enhancements at the short wavelength end of the spectrum near the pump. To get around this, a tunable Ti:Sapphire laser was built to shift the pump wavelength to shorter wavelengths. A schematic of the tunable Ti:Sapphire laser is shown in Figure 6.7. The laser is constructed in an X-fold cavity, that can be designed to compensate for the astigmatism of the Ti:Sapphire plate. The Ti:Sap-
The pump light from the Ti:Sapphire laser is directed through a beamsplitter, and focused onto the LED sample by a 10x microscope objective. This microscope objective has a numerical aperture of \( NA = 0.25 \), which corresponds with a collection angle of 14.5 degrees in air. By varying the offset of the LED sample position with respect to the focal plane of the microscope objective, different pump spot-sizes can be achieved. The pump light is absorbed by the InGaAs quantum well layer of the LED, and excites electrons hole pairs.

The PL and some residually reflected pump light are collected by the same 10x microscope objective, and reflected off of the beamsplitter toward a measurement apparatus. The light can either be imaged onto a Si charged coupled device (CCD) camera or spectrally resolved by an optical spectrum analyzer. To image the PL from the LED, chromatic filters are placed in the beam path to separate the PL from the residual pump. Images can be obtained with either a broadband filter, rejecting the pump light while transmitting all PL wavelengths, or by narrowband filters that only pass PL within a narrow wavelength range. By changing a flip mirror, the PL can be focused into an optical fiber, and directed into an optical spectrum analyzer (OSA) to measure the spectral content of the PL. A pinhole is placed in the image plane (before focusing into an optical fiber). By moving the pinhole across the image plane of the LED, the spectrum of the PL can be isolated from different spatial locations on the LED.
6.6 Unpatterned LED

An unpatterned LED mesa was studied as a reference sample that can be compared to micropatterned LEDs. A Si CCD camera image of such an unpatterned LED is shown in Figure 6.8. The LED mesa is in the shape of a square of dimensions 50 µm by 50 µm. A broadband spectral filter centered at 1014 nm, with a 103 nm full-width half-maximum, and RG 850 filters are used to suppress the reflected pump light, and transmit PL signal centered at 980 nm. The mesa is optically pumped by the Ti:Sapphire laser at 810 nm with a ~75 mm spot size, to uniformly illuminate the LED.

The emission from the LED is relatively uniform, with some notable internal structure. The border of the mesa is slightly brighter than the center of the mesa, because some guided light from the structure is output at the mesa facets. Inside the mesa, three square regions, progressively dimmer, are visible in the mesa. The interfaces between regions correspond to water refills in the oxidation process. The oxidation process is very sensitive to temperature [260]. When new water is added to the water boiler, the temperature of the water vapor changes, resulting in an interface in the oxidized layer. It is believed, however, that the entire mesa is oxidized from studies of the mesa's pump reflection.

Figure 6.8 Silicon CCD image of photoluminescence (PL) from an unpatterned light emitting diode (LED) mesa.
With all spectral filters removed from the optical beam path, residual pump light at 810 nm swamps the image. It is thereby possible to image the pump reflection as a function of position on the LED mesa. Such an image is shown in Figure 6.9. As can be seen, the 810 nm pump light is reflected more or less uniformly by the LED mesa. The pump wavelength is reflected by the broadband reflecting DBR mirror below the LED mesa. The uniform reflection indicates that the AlGaAs layers have been completely oxidized into Al₂O₃.

Figure 6.10 shows the LED photoluminescence (PL), as measured and recorded by an optical spectrum analyzer (OSA). The LED spectrum is peaked at 990 nm, close to the design goal of 980 nm, with a full-width half-maximum of 32 nm. A small side-band, on the short wavelength side, is centered at 930 nm. This smaller side-band is thought to be a result of a second excited quantum well state transition.
6.7 **LED with Random Holes**

Before studying the effectiveness of ordered photonic crystal microstructures to enhance the output coupling of LEDs, it is interesting to explore the ability of disordered microstructures to achieve similar enhancements. Figure 6.11 shows a schematic of the use of disordered surface roughness to increase the light extraction efficiency of LEDs. In a ray optics picture, surface roughness will increase the probability that optical rays will hit the semiconductor-air interface at an angle which will be emitted rather than undergo total-internal-reflection. This picture holds in the regime where $\lambda \ll L_{\text{rough}}$, where $L_{\text{rough}}$ is a characteristic length-scale of the surface roughness. For $\lambda \gg L_{\text{rough}}$, Raleigh scattering dominates,
and gives a characteristic $\lambda^{-4}$ scattering dependence. Yablonovitch et. al. have shown that a disordered array of wavelength scale features can be used to increase the extraction efficiency of LEDs [248]. In this intermediate regime, where $\lambda \sim L_{\text{rough}}$, a combination of Mie and other scattering process combine. Disordered arrays of wavelength-scale holes can increase the extraction by generating large amounts of incoherent scattering.

![Figure 6.11 Schematic of extraction enhancement process from an a disordered microstructure on top of a planar waveguide.](image)

In order to study the effects of disordered scattering on LED extraction, a microstructured region consisting of randomly placed holes was fabricated by electron-beam lithography. The microstructured region was a 30 µm by 30 µm area within the larger 50 µm by 50 µm LED mesa. The extraction efficiency of the microstructured region can be directly compared to the extraction efficiency of the non-patterned border region. A scanning electron microscope image of the disordered hole array is shown in Figure 6.12. A 5 µm by 5 µm pattern of holes was created, and repeated in a grid to create the entire 30 µm by 30 µm microstructured region.
A Si CCD image of the PL from a LED with random hole arrays is shown in Figure 6.13. The 30 μm by 30 μm microstructured region contained within the larger LED mesa is clearly brighter than the unpatterned LED border region. The repeating structure of the disordered array of holes is visible in the image of the LED mesa.
Figure 6.13  Silicon CCD image of photoluminescence (PL) from a 50 by 50 μm light emitting diode (LED) mesa. A significant enhancement in PL is visible from a 30 by 30 μm region with a disordered array of holes.

By using a pinhole placed in the image plane of the optical path, it was possible distinguish between spectrum from the patterned region and the unpatterned region. Figure 6.14 shows the spectrum from the patterned region of the LED, and of an adjacent LED with no patterning. The emission from the patterned region exhibits an enhanced LED emission compared to the LED mesa without holes. Furthermore, the patterned region exhibits the same general spectral shape as the unpatterned LED in Figure 6.10. The enhancement is therefore appears to be relatively uniform over a broad range of wavelengths.
A spectral enhancement efficiency is calculated by taking the ratio of the PL spectra of a patterned LED and a reference unpatterned LED for a constant optical pumping power, and plotted against wavelength in Figure 6.15. The spectral enhancement efficiency is a smooth function across the emission wavelengths, and is as high as 23 for at 925 nm. The wavelength dependence could either be a result of wavelength dependent output scattering or an increase in emission for shorter wavelengths as the pump power is increased by scattering enhanced input coupling. An integrated enhancement factor is determined by taking the ratio of the patterned and unpatterned PL spectra. The region with a disordered array of holes is enhanced by a total of 8.9 times over the unpatterned LED.
Figure 6.15  Spectral enhancement, the ratio between the photoluminescence (PL) intensity of the enhanced patterned region and a reference unpatterned region, as a function of wavelength.

An image of the pump reflection is shown in Figure 6.16. Unlike for the LED mesa with no hole patterning, the pump reflection is nonuniform. The wavelength-scale disordered array of holes not only scatters the waveguided emission inside the LED mesa, but also the 810 nm pump light. This is further evidence that the random array of holes produces incoherent, broadband scattering.
6.8 Triangular Photonic Crystal

While a disordered array of holes can incoherently scatter light from the waveguide modes of an LED into radiation modes, it is possible to design an ordered array of holes to coherently scatter light from LED waveguide modes into radiation modes. One approach to this end is to choose a photonic crystal lattice that will interact with the propagating light through Bragg scattering. In order to study the effects of Bragg scattering on the extraction efficiency, LED mesas with triangular photonic crystals were fabricated and studied. Triangular lattice photonic crystals have been shown to exhibit superprism effects [261] and to support PBG waveguiding with one-dimensional defects [262 - 265].

Figure 6.16: Silicon CCD image of the pump reflection from a 50 by 50 μm light emitting diode (LED) mesa, with a 30 by 30 μm patterned region with a disordered array of holes.
Each triangular lattice consisted of a regular array of circular holes, defined by electron-beam lithography into the semiconductor material. A scanning electron micrograph (SEM) of a typical triangular lattice photonic crystal is shown in Figure 6.17(a).

The reciprocal (k-space) lattice of a triangular lattice is another triangular lattice, rotated 30 degrees from the real-space lattice, with reciprocal lattice vectors of magnitude $(2\pi/a)$. A diagram of the real-space and reciprocal-space lattices are shown in Figure 6.17(b). The black circles represent the real-space lattice points, while the gray circles represent the location of reciprocal lattice points. The center circle is part of both the real-space and reciprocal-space lattices. By drawing a line perpendicular to each reciprocal lattice vector, a vector from a reciprocal lattice point to its nearest neighbors, it is possible to construct the first Brillouin zone. Any vector in reciprocal space can then be transformed into the first Brillouin zone by addition or subtraction of reciprocal lattice vectors. The dispersion relations of a photon in a triangular lattice can therefore be described completely by considering only the first Brillouin zone.

Calculated dispersion relations for the particular photonic crystal shown in Figure 6.17 are shown in Figure 6.18. Each line represents a waveguided mode within the photonic crystal semiconductor slab. The white region contains only purely guided waveguide modes, while the gray region contains radiating free-space modes in addition to leaky waveguide modes. The waveguide modes are leaky because they phase match to radiation with some lifetime. The light line separating the region with radiation modes and without is the disper-
sion relationship for freely propagating light in the waveguide plane. It is possible for radia-
tion to exist above the light line by propagating out of the waveguide plane. Because the
waveguideing layer is asymmetrical, purely TE (electric field parallel to the substrate) and
TM (electric field perpendicular to the substrate) modes do not exist. Only modes that are at
least 80% TE-like are plotted in the dispersion diagram. TE-like modes will dominate the
quantum well-photonic crystal coupling. A schematic of the first Brillouin zone with its sym-
metry directions labeled is inset. The nearest neighbor reciprocal lattice points are located in
the six M directions from the \( \Gamma \) point.

Figure 6.18  Calculated dispersion relations for a triangular lattice photonic crystal. Waveguided modes are
indicated by black lines. Radiation modes exist in a region above the light line (the dispersion
line for freely propagating light) because radiating modes can carry energy propagating
orthogonal to the waveguide plane. Waveguides in the region of radiation modes are called
leaky guided modes, and can phase match radiation with a lifetime.

It is interesting to look more closely at the bandstructure in the \( \Gamma-\text{M} \) direction. There
is a waveguiding mode that exists, after some low frequency cut-off. At the M point, the band
is folded back toward the \( \Gamma \)-point, and a local bandgap is created. In this particular band dia-
gram, there is no global photonic bandgap; for each frequency, there exists at least one guided optical mode. Folding back from the M-point toward the Γ-point, the mode crosses the lightline dividing allowed and forbidden radiation modes. The waveguide mode continues on toward the Γ-point as a leaky guided mode. At the Γ-point, we see that 3 optical modes converge. Actually, there are 6 modes including TM modes which are almost degenerate.

The electric field profile of the optical mode at the degenerate point has been calculated, and is shown in Figure 6.19. As can be seen by the side view, the electric field is concentrated within the dielectric. It is therefore legitimate to call this optical mode a leaky guided mode. The electric field forms a pattern that is similar to the interference of 6 plane-waves traveling along the 6 M symmetry directions. At the Γ-point, group theory indicates that the degenerate mode in a triangular lattice will be composed of a superposition of those plane-waves. LEDs contained within the semiconductor waveguide emitting at the frequency of degenerate mode will efficiently couple light to the leaky mode. Because the mode is leaky, the light will eventually be lost to radiation modes. This radiated light will be emitted in a directional manner normal to the waveguideing plane, because the mode couples to free-space radiation traveling at the Γ-point.

It is possible to design a single LED wafer to study several locations within the photonic bandstructure by creating photonic crystals with a variety of hole diameters (d), hole lattice center-to-center spacings (a), and hole depths (t). Each processed wafer contained
photonic crystals with the same hole depth, but a range of hole diameters and lattice spacings. Figure 6.20 shows a schematic of a photonic crystal LED mesa array. Each LED mesa in a row contain triangular photonic crystal lattices with the same hole diameter, but increasing lattice constant from mesa to mesa. Each column contains photonic crystal lattices with the same lattice constant but increasing hole diameter. The photonic crystal mesas are surrounded by LED mesas with no microstructuring, which are used as references with which to compare extraction efficiencies.

6.9 Output Coupling with Photonic Crystal

By designing a microstructured photonic crystal lattice to be the proper dimensions, it is possible to overlap the $\Gamma$-point degenerate leaky guided mode with the emission wavelength of the LED. Light emitted by the LED would primarily couple into the leaky guided mode, and then slowly radiate out of the structure. Such a structure was fabricated and studied, and shown to exhibit a significant extraction efficiency enhancement compared to non-microstructures LEDs. A scanning electron micrograph of a photonic crystal lattice containing a degenerate state overlapping with the LED emission is shown in Figure 6.17(a). This microstructure from sample R713d4, has a hole diameter $d = 100$ nm and a lattice spacing of $a = 380$ nm. The holes do not penetrate through the top InGaP cladding layer, in order to prevent the removal of active InGaAs material and the creation of surface recombination states.
Unfortunately, quantum well damage from other sources was generated below the microstructured region in some mesas which inhibiting quantum well emission in the photonic crystal region. The quantum well damage is thought to be a result of oxidation and annealing during the processing.

Despite some quantum well damage, it was possible to observe an enhancement in the extraction efficiency of the LED. A Si CCD camera image of the photonic crystal microstructured LED is shown in Figure 6.21 (a,b). The Ti:Sapphire pump spot size is approximately 75 μm. The left image (Figure 6.21(a)) was taken with a broadband spectral filter passing all PL wavelengths. A clear extraction efficiency enhancement is visible in the central photonic crystal region compared with the outer unpatterned region. The center square is an artifact of changing water during the oxidation process, as discussed above. The right-hand CCD image was taken without a spectral filter, and therefore images the reflected pump light at 810 nm. The reflection is seen to be roughly uniform across the photonic crystal and unpatterned regions.

Spectrally resolved CCD images were taken with a chromatic filter in the beam path with a bandwidth of 10 nm full width at half-maximum centered at 925 nm, 950 nm, 975 nm, and 1000 nm respectively, and are shown in Figure 6.22. At 925 and 950 nm a particularly strong extraction efficiency enhancement is observed from the central photonic crystal region. As the wavelength of the chromatic filter is increased, the PL intensity from the unpatterned region increases and peaks for the 975 nm filter (closest to the 980 nm peak of the PL), while...
the extraction efficiency enhancement from the photonic crystal region compared with the neighboring unpatterned region decreases in strength.

Figure 6.22 Spectrally resolved Silicon CCD images of photoluminescence (PL) from a light emitting diode (LED) mesa. Each image was taken with a 10 nm full width at half maximum chromatic filter to isolate PL of a particular wavelength. Particularly strong extraction efficiency enhancements from the central photonic crystal region are observed near 925 and 950 nm.

Measurements were performed with small pump spot sizes to determine if the increased light extraction in the photonic crystal region was a result of an increase in local emission or increased output coupling and to examine the coupling of light from the purely
guiding unpatterned region of the LED into the leaky guided mode of the photonic crystal region. Si CCD images are shown in Figure 6.23. In these measurements, a pump beam with

![Figure 6.23](image)

Figure 6.23 Spectrally resolved Silicon CCD images of photoluminescence (PL) from a light emitting diode (LED) mesa. Each image was taken with a 10 nm full width at half maximum chromatic filter to isolate PL of a particular wavelength. In each image, the PL was optically excited by a small pump spot on the left-hand side of the unpatterned border region of each mesa.

a spot size < 5 μm was focused onto the LED mesa. The optical pump was used to excited the LED in the unpatterned border region. An image of the LED mesa was then taken with 10 nm full-width at half-maximum filters, centered at 925 nm, 950 nm, 975 nm, and 1000 nm. At 1000 nm, PL from the LED was not observed from the photonic crystal region. This indicates that light was not efficiently coupled out of the photonic crystal at this wavelength. Use of the 975 nm filter revealed roughly similar results. Placement of the 950 nm and 925 nm filters, however, revealed light extracted from the photonic crystal region. It is believed that light guided in the unpatterned LED region coupled into the leaky guided mode in the photonic crystal. The light was then efficiently coupled out of the LED in the normal direction by the degenerate optical state at the Γ-point.
A second small spot size measurement was performed on a larger 100 μm by 100 μm LED mesa, with a 50 μm by 50 μm photonic crystal region. The additional size allowed the observation of interesting qualitative features in the coupling between the purely guided mode of the unpatterned LED region, and the leaky guided mode of the photonic crystal. A concatenation of four Si CCD images, taken from a single LED mesa, is shown in Figure 6.24. The

![100 μm square mesa](image)

![pump](image)

![photonic crystal](image)

Figure 6.24 Concatenation of Silicon CCD images of photoluminescence (PL) from a 100 by 100 μm light emitting diode (LED) mesa. In each of the 4 images, the mesa was excited in the unpatterned border region close to the central 50 by 50 μm photonic crystal region. Light is coupled from the unpatterned region guided modes into the patterned region's leaky guided modes. The propagation directions of the light coupled into the photonic crystal region reflect the 6-fold symmetry of the triangular lattice.

PL from the small pump spot size is seen at 4 positions around the outside of the photonic crystal region. A circle of PL is visible in the region of the pump excitation. While some light is extracted in this small region, the remainder couples into waveguided modes in the LED mesa slab. When incident upon the photonic crystal region, some of the light from the waveguided slab couples into the leaky guided mode of the photonic crystal region. The light propagating in the leaky guided modes of the photonic crystal region is visible because of coupling into radiating modes. Interestingly, the light streaks are seen to propagate along the 6 plane-wave directions that make up the degenerate modes at the Γ-point.
In order to quantify the extraction efficiency enhancement of the LED in the photonic crystal region, the PL was spectrally resolved with an optical spectrum analyzer. Figure 6.25 shows PL of the photonic crystal mesa with a degenerate Γ-point mode at 925 nm, compared with the PL from a reference unpatterned LED mesa. The PL from the photonic crystal mesa

![Graph showing PL with Triangular Lattice and No Holes](image)

**Figure 6.25** Photoluminescence (PL) from an unpatterned light emitting diode (LED) mesa (no holes), and from a region patterned with a photonic crystal (with triangular lattice). There is a significant enhancement in the intensity of extracted light from the patterned region.

is dramatically enhanced as compared with the PL of the reference unpatterned mesa. A particularly large enhancement is visible near 925 nm. This enhancement is thought to be due to the existence of a degenerate Γ-point leaky mode at near 925 nm.

The ratio between the photonic crystal PL and the reference LED PL is computed, and plotted in Figure 6.26. A large spectral enhancement of the extraction efficiency is seen at the shorter wavelength end of the LED emission spectrum. In particular, a spectral enhancement of 100 is observed at 925 nm. At first glance, a spectral enhancement of over 100 times may appear to be unphysical, as it has been already noted that the LED mesa without patterning...
should extract 5% of the emitted light by the LED. If there is a 100 times spectral enhancement at 925 nm, this would mean that 500% of the light emitted would be extracted. This apparent difficulty is resolved by the fact that the Γ-point degenerate state will directionally couple light normal to the LED surface. The PL experiment described above collects light with a 0.25 NA microscope objective. The light enhancement observed is therefore a result of both an increase in the extracted light and an increase in the directionally of the output coupled light. This directional characteristic of the light extracted by the photonic crystal is in stark contrast with the nondirectional enhancement due to incoherent light scattering from a random array of holes.

![Graph showing spectral enhancement](image)

Figure 6.26  Spectral enhancement, the ratio between the photoluminescence (PL) intensity of the enhanced patterned region and a reference unpatterned region, as a function of wavelength. A spectral intensity over 100 is observed at 925 nm.
6.10 **Power Dependence**

The PL spectrum of the LED mesa with a strong enhancement at 925 nm exhibits a qualitatively different spectral profile, characterized by a sharp peak at 925 nm, than the reference LED mesas. It was noticed that at low excitation powers, this enhanced LED mesa had a spectrum that looked more like the reference LED mesa. The pump power dependence of the PL emitted from mesa (4,4) was studied. A plot of the PL for a variety of pump powers is shown in Figure 6.27. Because the short wavelength features are attributed to excited states of

![Graph showing the power dependence of the photoluminescence (PL) from mesa (4,4).](image)

Figure 6.27 Power dependence of the photoluminescence (PL) from mesa (4,4).

the quantum well which are above the ground state, the change in spectral shape can be attributed to saturation of the ground state. For all pump powers studies, the enhancement remained roughly constant. Further study of the power dependence could help differentiate non-radiative and radiative recombination effects in the LED.
6.11 **Wavelength Tuning**

Enhancements of LED extraction by the Γ-point degenerate leaky mode are spectrally narrow and highly directional. By changing the dimensions of the photonic crystal, it is possible to tune the Γ-point mode over the emission spectrum of the LED. The devices discussed until this point enhanced the extraction from the LED mesa at 925 nm. This wavelength corresponds with the short wavelength shoulder of the PL spectrum. Despite the relatively small fraction of PL emitted near the enhancement wavelength, the 100-fold spectral enhancement led to an 8-fold enhancement integrated over all wavelengths. By tuning the photonic crystal to have a Γ-point mode at 980 nm at the peak of the PL spectrum, it should be possible to realize much larger integrated enhancements.

Unfortunately, other effects have limited the upper wavelength of the degenerate mode in our samples. The wavelength of the leaky mode can be increased by either increasing the lattice constant (a) of the photonic crystal, or decreasing the size of the hole diameter (d). The photonic crystal lattice with the resonance at 925 nm already had the largest lattice constant on the fabricated LED chip. Therefore, it was decided to look at neighboring LED mesas fabricated with larger or smaller hole diameters.

A series of PL spectra taken from photonic crystal LED mesas with varying hole diameters are shown in Figure 6.28. The LED mesa with a calculated degenerate mode at 925 nm (mesa (4,4)) exhibits the largest spectral enhancement.
Figure 6.28  Photoluminescence (PL) from several unpatterned light emitting diode (LED) mesas patterned with photonic crystals of with different hole diameters.
Figure 6.29 shows the calculated spectral enhancement for the LED mesas. As the hole diameter increases, the spectral enhancement should shift to shorter wavelengths. 925 nm, however, is already on the edge of the emitted spectrum, and therefore no major spectral shift is observed. Rather, the peak at 925 nm is seen to diminish as the resonance wavelength shifts to lower wavelengths. As the hold diameter is decreased, the spectral enhancement should shift further into the PL spectrum to higher wavelengths. Instead, the enhancement disappears. SEM pictures of the LED mesas with smaller hole diameters show that the holes were underdeveloped. Basically, the holes were too small, and the RIE etching process was unable to remove the material. Therefore, no holes were actually created.

The PL from each mesa was integrated, and the total light emission was tabulated in Figure 6.30. Mesa (4,1) is a reference mesa, with no photonic crystal patterning. Mesas (4,2)
through (4,9) contain a photonic crystal lattice with increasing hole diameter. Mesa (1,4) is the LED with a disordered array of holes. The incoherent random scattering mesa exhibited the largest total enhancement because it enhanced emission over a broader wavelength range than the photonic crystal.

![Bar Chart](image)

**Figure 6.30** Integrated enhancement of photoluminescence (PL) from several light emitting diode (LED) mesas. Mesa (4,1) is unpatterned, mesas (4,2) through (4,9) have a photonic crystal of increasing hole diameter, and mesa (1,4) contains a disordered array of holes.

### 6.12 Angular Dependence

The extraction efficiency from photonic crystal LEDs is thought to be a result of radiation from leaky guided degenerate modes at the Γ-point of the photonic band structure. If this is the case, the enhanced radiation should be to free-space modes at the Γ-point in the...
plane of the LED. These free-space modes correspond to radiation traveling normal to the LED mesa’s plane, and are highly directional. The angular dependence of the PL from the enhanced LED mesa is therefore investigated to determine if the LED extraction efficiency enhancements are directional.

The angular dependence of PL from an unpatterned LED mesa, studied as a reference, is plotted in Figure 6.31. The LED mesa is rotated with respect to the collection optics to an angle of 0 (normal), 10, 15, and 25 degrees. The spectral lineshape of the PL does not change as the angle is increased, but a slight increase in the total collected PL intensity is observed. This increase in the collected PL is expected because the imaged field of view of the collection optics also increases with angle.
The PL of an enhanced LED mesa was also measured while rotated at different angles with respect to the collection optics’ optical axis, and is shown in Figure 6.32. The strong spectral enhancement at the wavelength of the Γ-point degenerate mode is only present in the PL from the normally (0 degrees) collected LED mesa. The PL spectra taken at various angles also exhibit a broadband enhancement, but have spectral lineshapes similar to those of an unpatterned mesa. This indicates that at larger angles, there are no dramatic narrow frequency spectral enhancements.

A plot of the spectral extraction enhancement efficiencies for PL collected at different angles are shown in Figure 6.33. The spectral enhancement is calculated by taking the ratio between the PL of the photonic crystal LED mesa and an unpatterned reference LED mesa. The 0 degree PL exhibits the strong spectral enhancement at 925 nm, corresponding to the degenerate leaky mode at the Γ-point. The PL from this mesa at larger angles also exhibits...
enhancements that are not spectrally localized. These enhancements may be due to leaky guided modes throughout the bandstructure spread out through many wavelengths.

![Graph showing spectral enhancement vs. wavelength](image)

Figure 6.33 Angular dependence of spectral enhancement of photoluminescence (PL) from a light emitting diode (LED) mesa patterned with a photonic crystal. PL enhancements are measured with the LED mesa rotated 0, 10, 15, and 25 degrees from normal.

6.13 **Pump Input Coupling with Photonic Crystal**

A photonic crystal lattice with the proper dimensions can coherently couple light from waveguided modes into radiation modes through Bragg scattering. In a bandstructure picture, a degenerate state is created at the $\Gamma$-point which corresponds to a leaky guided mode. Light in the leaky guided mode will couple to radiation directionally in the $\Gamma$-direction, normal to
the LED plane. The extraction efficiency of a semiconductor LED was increased by overlapping the degenerate optical mode with the emission wavelengths of the LED.

The total wallplug efficiency of an LED, however, is the product of the extraction efficiency and the quantum efficiency of the LED. The quantum efficiency can be increased by increasing the optical pumping efficiency. Using the same degenerate optical state as is used to increase output extraction efficiency, the optical pumping efficiency of a photonic crystal LED can be increased. A photonic crystal with a degenerate Γ-point optical mode at 810 nm is created. Light from the 810 nm pump will then be Bragg scattered into the planar waveguide mode. Once in the waveguide mode, the pump light will have a larger interaction length with the quantum well, leading to greater optical excitation, and greater PL.

A scanning electron micrograph (SEM) of a photonic crystal designed for optical pumping enhancement is shown in Figure 6.34. The photonic crystal has holes of diameter \( d = 210 \text{ nm} \) and hole lattice spacing of \( a = 340 \text{ nm} \). The hole diameter is increased, while the lattice spacing has decreased compared to the photonic crystal with a Γ-point mode at 980 nm. It is straightforward to see that a smaller lattice spacing will correspond to bandstructure shifting to smaller wavelengths. The larger hole size results in more of the electric field being in air, rather than the higher index dielectric. This will increase the energy, or decrease the wavelength, of the optical mode because the energetically favorable location of electric field is in the dielectric. The calculated bandstructure for the photonic crystal shown in Figure 6.34

![Photonic Crystal SEM Image](image-url)
is shown in Figure 6.35. A resonant degenerate optical state is created near the pump wave-

![Diagram of band structure of photonic crystal]

Figure 6.35  Band structure of photonic crystal designed to increase pumping efficiency at 810 nm

length of 810 nm. Otherwise, the bandstructure remained similar to the bandstructure for the 980 nm emission enhancement case.

Si CCD images of the LED PL and reflected pump light are shown in Figure 6.36. In these images the LED mesa is 50 µm by 50 µm, with a 30 µm by 30 µm photonic crystal region in the middle. The LEDs were pumped with a Ti:Sapphire spot size of about 75 µm. The left image was recorded with a broadband chromatic filter covering the emission range of the LED. In this image, a dark spot is observed in the photonic crystal region. It is thought that the larger diameter holes, in conjunction with the oxidation process, resulted in quantum well damage that reduced the PL in this area. This is despite the fact that the hole depth was deliberately designed to stop short of the active quantum well layer to avoid the creation of surface recombination states. A small bright spot is observed in the upper right-hand corner of the photonic crystal region. This is a defect of no holes placed intentionally within the photonic crystal region. Around the edge of the photonic crystal region is a bright halo which is particularly pronounced on the right-hand side. This bright halo is thought to be enhanced optical coupling due to the degenerate state at the pump wavelength. Further evidence of the
pump coupling is visible in the 810 nm reflection image. A reduced reflection is visible in the photonic crystal region. The reflection in the normal direction is reduced because some of the pump beam is Bragg scattered, either into the waveguiding slab, or off axis with the imaging system, by the photonic crystal.

980 nm emission

810 nm reflection

50 µm square mesa

Photonic crystal

Figure 6.36 Silicon CCD images of a light emitting diode (LED) with a photonic crystal designed to increase optical pumping efficiency at 810 nm. Because 810 nm light is coupled into the waveguide modes, there is a dark spot in the pump reflection in the presence of the photonic crystal.

In order to elucidate the nature of the enhanced pumping of the bright halo region, small pump spot size measurements were performed on a larger 100 µm by 100 µm photonic crystal mesa. A composite of 8 images taken with a pump spot size < 5 mm is shown in Figure 6.37. Each of the 8 images are separate images taken with the optical pump directed just inside the photonic crystal region. The pump is Bragg scattered by the photonic crystal lattice into the waveguiding modes, where it efficiently pumps the LED. The LED does not illuminate inside the photonic crystal because of quantum well damage. In the unpatterned border region surrounding the photonic crystal, however, the LED shines brightly. In fact, much like the small spot size measurements in the efficient extraction sample, the efficient pumping can be seen directionally along the 6 symmetry axes of the photonic crystal. This is because the degenerate Γ-point mode that is coupled to by the pump beam can be written as a superposition of the 6 plane waves in the 6 M symmetry directions.
Figure 6.37 Composite of 8 Silicon CCD images of a small pump spot size measurement of photoluminescence (PL) from a photonic crystal light emitting diode (LED). Each image was taken with a small pump spot within the photonic crystal region. Bright patterns of PL are visible with the 6-fold symmetry corresponding with the triangular lattice photonic crystal.

6.14 Future Work

Photoluminescence excitation (PLE) measurements could be used to study in detail the optical pumping enhancement due to the photonic crystal effect. Photoluminescence (PL) experiments are performed by looking at the emission versus wavelength for a sample optically pumped at a fixed wavelength. This gives information about the sample’s emission characteristics. PLE measurements monitor the PL at a fixed wavelength, and adjust the optical pump wavelength. As the optical pump wavelength is scanned in wavelength across an absorption feature, the PL will increase and then decrease with the absorption. Thus, PLE gives absorption information about a sample.

Once the input pump and LED output coupling is completely understood, it should be possible to combine the two effects to create an optically pumped LED with extremely high
efficiency. An LED could, for example, be fabricated in the 'soccer ball' geometry shown in Figure 6.38. In this geometry, regions of photonic crystal designed to enhance input coupling

![Soccer ball geometry for combining pump input and light emitting diode (LED) output photonic crystal regions.](image)

for the pump and output coupling for the LED placed next to each other. The pump input coupling regions are placed to distribute the pump light evenly throughout the LED. Alternatively, it may be possible to design a single photonic crystal that has enhancements for both optical pump input and LED output.

Initial experiments have observed lasing from a quantum well structure in the presence of a triangular lattice photonic crystal with suitable pump power. The lasing occurs despite the absence of a laser cavity because of distributed feedback from photonic crystal scattered light [266]. This effect has been observed previously in active organic [267, 268] and semiconductor [269] materials. The previous semiconductor lasing results were in a structure where the active region was wafer-bonded onto the separately fabricated photonic crystal. In the structures studied here, the holes were fabricated into the active region. Further study of the lasing could help elucidate the exact nature of the lasing mechanism and the photonic crystal bandstructure.

It has been shown that a photonic crystal with a Γ-point resonances can enhance the extraction efficiency in a directional manner. By using a photonic crystal with a photonic
bandgap (PBG), the extraction efficiency of an LED could be increased in a non-directional manner [259]. The PBG enhances the forward emitted light by totally eliminating the waveguided modes. A PBG is a range of frequencies for which no optical modes exist. An example of a one-dimensional PBG is a periodic quarter wave dielectric stack mirror with a stopband. For the wavelengths of the stopband, no optical modes exist inside the dielectric stack. If a two-dimensional photonic crystal LED could be made with a PBG over the emission wavelengths of an LED, the LED emission will have literally no place to go but to the free space modes. Unfortunately, this effect could not be observed due to quantum well damage from the fabrication process. The Bragg enhancement effects observed do not require that the light is generated in the region of the holes. Light could be generated in the unpatterned border region of the mesa, and then output coupled by the photonic crystal. To observe the PBG enhancement effect, however, light must be generated in the photonic crystal region.

6.15 Conclusions

Extraction efficiency enhancements as large as 8-fold have been observed from LEDs with photonic crystal degenerate states near the $\Gamma$-point. At the wavelength of the degenerate states, spectral enhancement of 100-fold were observed. In the inverse effect, LED efficiencies were increased by using Bragg scattering to increase the input pump coupling.
CHAPTER 7: CONCLUSION

7.1 Conclusions

This thesis has reported on several developments in the manipulation of near infrared light.

Using double chirped mirrors (DCMs) for dispersion compensation, 20 fs pulses were generated by a prismless Cr\textsuperscript{4+}:YAG laser. The pulses had a corresponding spectral bandwidth of 190 nm fwhm and observable spectral intensity from 1140 to >1700 nm.

Broadband saturable absorber mirrors with high-dielectric contrast GaAs/Al\textsubscript{x}O\textsubscript{y} mirrors and InGaAs/InP quantum well absorbers were developed to self-start ultrafast Cr\textsuperscript{4+}:YAG laser pulses. Self-starting of 35 fs pulses was observed, and longer self-started pulses were tunable from 1400 to 1525 nm with a birefringent filter. Si/SiO\textsubscript{2} mirrors were shown to be of high enough quality to operate within a Cr\textsuperscript{4+}:YAG laser, and could be used as future SBR substrates or broadband output couplers.

Transmission through a photonic bandgap (PBG) waveguide microcavity was studied near 1550 nm. The resonant states had a modal volume as small as 2 (\lambda/2n)\textsuperscript{3} and cavity quality factors (Qs) as high as 360 were observed.

Two-dimensional photonic crystals were used to enhance the extraction efficiency of InGaAs/InGaP light emitting diodes (LEDs) at 980 nm by as much as 8 times. At particular wavelengths, the collected photoluminescence (PL) was increased by as much as 100 times.
In addition, the same photonic crystal was used to enhance the input coupling of pump light into the LED.

Building on these achievements, a variety of more exciting research seems possible in the near future.

7.2 Shorter Cr$^{4+}$:YAG Pulses

Pulses as short as 20 fs have been generated by a prismless Cr$^{4+}$:YAG laser. While prismless lasers are easy to align, they do not necessarily produce the shortest pulses. Intracavity prisms allow fine-tuning of the net cavity group delay dispersion (GDD). BaF$_2$ has been identified as a prism material that may allow generation of even shorter pulses. BaF$_2$ has a low water content and extremely low third-order dispersion. Furthermore, the small third-order dispersion complements the DCMs that have already been fabricated for the 20 fs prismless laser.

Shorter Cr$^{4+}$:YAG laser crystals would allow the fold-mirror angle, used to compensate laser crystal astigmatism, to be reduced. This reduction of angle would lead to a corresponding reduction in the GDD oscillation amplitude of the DCMs. Furthermore, it is possible that the Kerr lens modelocking (KLM) strength of cavities with short crystals is larger than for cavities with long crystals. The general area of KLM cavity design and optimization should be studied in more detail, and should include effects such as thermal lensing and space-time effects. Novel cavity geometries should also be given further thought.

Alternatively, a new set of DCMs could be designed to work in conjunction with BaF$_2$ prisms and a 2 cm laser crystal. To use the BaF$_2$ prisms, 8 to 10 reflections off of the current DCMs are required for proper GDD compensation. New mirrors that could compensate dispersion with fewer bounces would be desirable. Furthermore, π-phase shifted DCM pairs could be used to lower the GDD oscillation amplitude [135].

In the case of Ti:Sapphire, the pulsewidth is partially limited by the bandwidth of the intracavity optical elements. It is likely that at some point in the future the situation will be the same for Cr$^{4+}$:YAG. It is therefore important to develop ultrabroadband laser optics. A single wide area Si/SiO$_2$ mirror was demonstrated to be suitable for use within a Cr$^{4+}$:YAG laser. By polishing the backside of the Si wafer, a slightly less reflective mirror could be used as a broadband output coupler with flat GDD characteristics over a broad wavelength range. Alternatively, a complete set of Si/SiO$_2$ laser optics could be deposited on glass substrates and
used to construct a broadband laser cavity. This laser could be tunable over a wide range and potentially supports short pulses.

Efficient broadband saturable absorbers should continue to be studied and developed in a wide range of material systems. GaAs/Al₅O₇ oxidized saturable Bragg reflectors (SBRs) have been demonstrated to support at least 35 fs pulses. Pump-probe spectroscopy should be used to determine if the pulsewidth is limited by absorber effects such as two-photon absorption (TPA) and if the saturation fluence is an ideal value. If the absorbers are limiting, further design should be performed to eliminate the limitations. Otherwise, Cr⁴⁺:YAG laser cavities with better GDD compensation should be constructed to generate shorter pulses.

At the same time, absorber materials that can be integrated with Si/SiO₂ mirrors should be studied. Deposition of InAs quantum dot absorber layers onto the mirror substrates should be tested. New techniques to deposit GaAs or InP-based layers onto Si substrates should be explored. Alternative materials such as germanium and antimonides should be investigated to determine their suitability as saturable absorbers.

7.3 Toward Single-Cycle Pulses

The shortest pulses generated directly from a laser are the two-cycle pulses generated by a Ti:Sapphire laser. At two optical cycles, the corresponding spectral bandwidth already spans an octave (λ to 2λ). To generate shorter pulses, it becomes increasingly difficult to create laser mirrors with a large enough bandwidth.

Theoretically, pulses can be produced as short as a single optical cycle. One approach to generate single-cycle pulses is to lock together the phases and repetition-rate of two spectrally adjacent two-cycle laser pulses, and interferometrically combine the two pulse-trains. It has been shown recently that two Ti:Sapphire lasers can be locked together in such a manner to generate shorter pulses than either of the pulses separately [155]. More recently, a Ti:Sapphire and Cr:Forsterite laser were synchronized by passing the Cr:Forsterite through the Ti:Sapphire Kerr medium overlapped with the Ti:Sapphire pulse [156]. In both cases, the individual pulses were not short enough to produce single-cycle pulses once locked. It is possible that Ti:Sapphire and Cr⁴⁺:YAG together could be locked to generate a single-cycle pulse. A plot of the spectrum from a Ti:Sapphire and Cr⁴⁺:YAG laser is shown in Figure 3.49. Spectrum from the two lasers is visible from 600 to >1700 nm. Assuming a square spectrum over this range, it should be possible to generate pulses with widths on the order of 3 fs, which
is close to a single-cycle pulse at the carrier frequency. For this reason, multi-laser locking remains a particularly exciting area of current research.

7.4  Pump-Probe Spectroscopy

The temporal resolution of ultrafast spectroscopic techniques such as pump-probe are limited by the finite pulsewidth of the laser source. Near 1.5 μm, optical parametric oscillators with a pulsewidth of ~ 100 fs are commonly used to study the dynamics of semiconductor materials and devices. With a 100 fs source, pump-probe spectroscopy can often resolve the time-scale of carrier-lattice interactions but not carrier-carrier interactions. In many materials, these carrier-carrier scattering times and intervalley scattering times are believed to occur on the 10 - 100 fs time-scale. With the 20 fs Cr₄⁺:YAG laser, it should be possible to resolve some of these fast effects and set a new upper limit on the time-scale of others. As a start, bulk material films of materials with bandgaps of slightly lower energy than the excitation pulses would be relatively simple to study and possibly to interpret. InGaAs, Ge, and GaSb films all have fast effects that could possibly be resolved by ultrafast pump-probe.

7.5  Frequency Chains

The spectrum of modelocked laser pulsetrain consists of a comb of modes separated by the longitudinal mode spacing. For a longitudinal mode spacing of Δν = c/2L, where 2L is the cavity round-trip length, the frequency of the mth mode is:

\[ ν_m = m(Δν) + δ \] (7.1)

where δ is the phase-slip between the center of the pulse envelope and the underlying carrier wave that results from pulse to pulse due to the difference between the phase and group velocity of the pulses [9]. By measuring and stabilizing δ, it has become possible to use the ultrafast laser frequency comb as an ultrastable frequency reference [7]. In order to stabilize δ, a broadband spectrum is required. If the spectrum spans from λ to 2λ, it is possible to frequency double the light and interfere spectral line m with the line 2m of the fundamental light. The interference will beat at the frequency:

\[ ν_{beat} = 2ν_m - ν_{2m} = δ \] (7.2)
Change in the beat frequency $\delta$ is then used to feedback an error signal to piezoelectric elements, which in turn changes the laser cavity length and dispersion to stabilize $\delta$. Alternatively, the second and third harmonic generation signals from a laser can be beat against each other but might have smaller signals for octave-spanning spectra.

Carrier envelope phase control has been achieved in Ti:Sapphire lasers in two ways. In the first approach, pulses without enough bandwidth were broadened spectrally by propagation through a highly nonlinear photonic crystal fiber with low dispersion for the Ti:Sapphire wavelength [6]. Enough spectrum was generated to permit stabilization of the phase $\delta$. The second approach generated 5 fs pulses directly from a Ti:Sapphire laser with enough spectral bandwidth for direct stabilization of the phase [8].

The frequency combs that were generated cover the Ti:Sapphire spectral range. By stabilizing the carrier-envelope phase of the Cr$^{4+}$:YAG laser, it should be possible to produce a frequency reference in the telecommunication wavelength range. With a 20 fs pulse, there is already spectrum from 1140 nm to $>1700$ nm. Doubled 1200 nm and tripled 1800 nm light could be interfered to stabilize $\delta$. The signal is likely to be small for these frequencies at the edge of the spectrum. Shorter Cr$^{4+}$:YAG pulses could possibly extend the spectrum such that fundamental and doubled light could be used. It should also be possible to generate a broader spectrum by propagating the Cr$^{4+}$:YAG pulses through a telecommunication fiber with low dispersion such as TrueWave fiber. The low dispersion fiber could allow the pulses to propagate far enough to experience significant nonlinearities. Finally, the Cr$^{4+}$:YAG could simply be doubled and phase-locked to a stabilized Ti:Sapphire pulse. This step is required to generate the single-cycle pulses discussed in section 7.3.

### 7.6 Tunable, Switchable, and Active Photonic Crystals

Photonic crystals in high-dielectric contrast semiconductor systems have been shown to be capable of manipulating light on a wavelength scale. One of the primary difficulties in demonstrating the effectiveness of the photonic crystal devices has been that the wavelength of operation is often highly sensitive to lattice dimensions such as hole diameter, lattice spacing, and hole depth. The development of tunable devices is therefore of high importance. A device could then be designed for a particular wavelength and then fine-tuned to the proper location if the fabricated device is off. Furthermore, a single device could be mass produced, and the individual devices tuned to the proper wavelength. Tunability would give optical communication networks the ability to reconfigure the transmission path of light in real time.
System flexibility is also a driver for optical switching. Optical communication system designers would like to use devices that can turn on and off an effect, such as dropping a wavelength to an individual user. Micro-electro-mechanical machines (MEMs) can be used in conjunction with photonic crystals to tune or switch devices.

The development of active photonic crystal devices such as LEDs and microlasers has proven to be difficult for a couple of reasons. Nonradiative recombination at surface states increases as devices are made smaller. Furthermore, the oxidation process used to make Al$_x$O$_y$ could possibly cause damage to active quantum well layers in the current fabrication process. Further exploration should be done to reduce the detrimental effects of oxidation on active layers. Approaches to grow active GaAs or InP-based layers on Si should be pursued as an alternative to develop active semiconductor photonic crystal devices.

7.7 Near-field Scanning Optical Microscopy

Near-field scanning optical microscopy (NSOM) uses a sub-wavelength aperture to image light with sub-wavelength resolution. NSOM could be used to directly measure the modes inside of waveguide devices. Techniques to generate short pulses of light or manipulate light over wavelength-scale distances have been studied in this thesis. NSOM could be used to deliver short pulses of light to sub-wavelength areas. Light would then be shrunk to its fundamental time and length-scales.
REFERENCES


[44] Fast Kerr Effect


