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ENHANCING THE RESOLUTION OF IMAGING SYSTEMS BY SPATIAL SPECTRUM MANIPULATION

By

Wyatt Adams

A DISSERTATION

Submitted in partial fulfillment of the requirements for the degree of

DOCTOR OF PHILOSOPHY

In Electrical Engineering

MICHIGAN TECHNOLOGICAL UNIVERSITY

2019

 \bigodot 2019 Wyatt Adams

This dissertation has been approved in partial fulfillment of the requirements for the Degree of DOCTOR OF PHILOSOPHY in Electrical Engineering.

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Dedication

To my parents

for their love, guidance, and wisdom.

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Preface

This dissertation is mostly comprised of the four published articles in Chapters 2-5. Permissions to reuse the published articles have been obtained from the publishers, and letters of permission can be found in Appendix A. The majority of Chapter 6 is a draft of a paper which will be submitted for publication later.

The main contributions by the author to this research and to the body of knowledge are:

- 1. The theory and simulation of the proposed loss mitigation and resolution enhancement methods for incoherent light.
- 2. The implementation, design, construction, and analysis of the experiment in Chapter 6 to verify the proposed methods and extend them to far-field imaging.
- 3. The spectral SNR equation (Eq. 6.16) which quantitatively shows how to manipulate the pupil or transfer function of an imaging system to achieve a desired high spatial frequency SNR.

The individual contributions for each chapter are as follows:

¹Chapter 2: W. Adams performed the literature review, organized the paper, and wrote the body text. M. Sadatgol and D. Güney provided technical comments and edited the manuscript.

²Chapter 3: W. Adams and M. Sadatgol performed modeling and wrote data analysis code to obtain the primary results in the paper. W. Adams and M. Sadatgol organized the paper. W. Adams wrote the body text. X. Zhang and D. Güney provided technical comments. M. Sadatgol and D. Güney edited the manuscript.

³Chapter 4: W. Adams performed the modeling and data analysis to obtain the main results in the paper. W. Adams organized and wrote the paper. A. Ghoshroy and D. Güney provided technical comments. D. Güney edited the manuscript.

⁴Chapter 5: W. Adams performed the modeling and data analysis to obtain the main results in the paper. W. Adams organized and wrote the paper. A. Ghoshroy and D. Güney provided technical comments. D. Güney edited the manuscript.

Chapter 6: W. Adams, A. Ghoshroy, and D. Güney developed the underlying theory and core concepts. W. Adams constructed the experimental imaging system, collected

¹W. Adams, M. Sadatgol, and D. Ö. Güney, "Review of Near-field Optics and Superlenses for Sub-diffraction-limited Nano-imaging," *AIP Advances* 6, 100701 (2016).

²W. Adams, M. Sadatgol, X. Zhang, and D. Ö. Güney, "Bringing the 'perfect lens' into focus by near-perfect compensation of losses without gain media," *New Journal of Physics* 18, 125004 (2016).

³W. Adams, A. Ghoshroy, and D. Ö. Güney, "Plasmonic superlens image reconstruction using intensity data and equivalence to structured light illumination for compensation of losses," *Journal of the Optical Society of America B* 34, 2161-2168 (2017).

⁴W. Adams, A. Ghoshroy, and D. Ö. Güney, "Plasmonic Superlens Imaging Enhanced by Incoherent Active Convolved Illumination," *ACS Photonics* 5, 1294-1302 (2018).

experimental data, performed modeling, designed the superlensing device, performed all data and image processing, and wrote the chapter text. C. Kendrick printed the spatial filter transparencies for the experimental imaging system. C. Middlebrook provided the lab space and some of the optics equipment to obtain the experimental results.

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A special thank you is reserved for my PhD advisor, Dr. Durdu Güney, for his unrelenting support and guidance throughout the entirety of my graduate studies. He has taught me to not be too quick to doubt my ideas, even when they at first seem insignificant or far-fetched. I have never met such a forward-thinking person as him, a trait which I have found genuinely inspiring.

Finally, I owe a deep debt of gratitude to my family and friends. Their positive

contribution to my life, and consequently this work, is something I simply cannot even hope to repay.

Abstract

Much research effort has been spent in the 21st century on superresolution imaging techniques, methods which can beat the diffraction limit. Subwavelength composite structures called "metamaterials" had initially shown great promise in superresolution imaging applications in the early 2000s, owing to their potential for nearly arbitrary capabilities in controlling light. However, for optical frequencies they are often plagued by absorption and scattering losses which can decay or destroy their interesting properties. Similar issues limit the application of other superresolution devices operating as effective media, or metal films that can transfer waves with large momentum by supporting surface plasmon polaritons.

In this dissertation, new methods of mitigating the loss of object information in lossy and noisy optical imaging systems are presented. The result is an improvement in the upper bound on lateral spatial resolution. A concentration is placed on metamaterial and plasmonic imaging systems, and the same methods are subsequently adapted to more conventional far-field imaging systems.

First, through numerical simulation it is shown that a lossy metamaterial lens has degraded imaging performance which can be partially compensated by deconvolution post-processing of the resultant image. This post-processing procedure is then shown to emulate a physical process called plasmon injection, which has been previously implemented to effectively remove the losses in a plasmonic metamaterial.

Next, a more realistic scenario is considered; a thin film of silver acting as a near-field plasmonic "superlens." In this case, methods are implemented to model incoherent light propagation so that the image can be reconstructed using only intensity data, removing the need for phase measurement. The same procedure from above is followed, and the resolution is enhanced. To push the resolution further, a spatial filtering method called active convolved illumination is developed to overcome the resolution limit set by the noise floor of the system.

Finally, the spatial filtering methods are applied to more a more conventional far-field imaging system. Supported by experiment, the lateral resolution of a low numerical aperture imaging system is improved by blocking photons at the Fourier plane. For coherent light, a diffractive superlens is designed which uses the same principles from the above theory, except it encodes the high spatial frequency waves into propagating waves via a diffraction grating. The result is lateral resolution performance that surpasses similar previously published devices by 10 nm at a wavelength more than 80 nm longer.

Chapter 1

Introduction

1.1 Overview

Much of what humans have learned about the universe can be attributed to light. Points of visible light in the sky first notified humankind of the presence of stars and the entire cosmos. Perhaps the cell would not have been discovered without the compound light microscope. The discovery of the photoelectric effect led to the earliest quantum theories of matter and won Einstein his Nobel prize. The development of radio technology led to the telescopes which first viewed the cosmic microwave background. Fast forward to today, and it is hard to imagine what contemporary physics research would look like without the highly coherent light provided by lasers. It seems ironic that the humble massless photon, which lays unassuming in the standard model mostly consisting of massive particles, is the one which has revealed so much about the material world.

Consequently, it can be argued that photonics is one of the most important fields in all of science and engineering. As we increase our knowledge and control over light, we not only increase our understanding of the photon and its interactions, but also potentially further the reach of our senses to collect information about how the universe works.

Until the current millennium, we had manipulated light using literally the matter at hand, i.e., the elements readily available to us on Earth. A quick glance at the periodic table suggests that this offers a finite set of optical properties, and in turn a finite set of capabilities. However, arranging these given building blocks into carefully designed sub-wavelength composite structures can yield interesting effective properties which are difficult, if not impossible, to find in natural materials. So-called electromagnetic "metamaterials" have enabled a host of interesting phenomena and applications, including negative refraction[1], cloaking[2], and superlensing[3].

Perhaps the most exciting prospect of metamaterials is their potential to achieve high resolution imaging beyond the diffraction limit, the so-called "perfect lens" idea. This has inspired much research (and skepticism) in recent years, and has proved not without its unique challenges. There is not one single device which has given an all-encompassing solution to the superresolution problem, as plasmonic superlenses, hyperlenses, far-field superlenses, and many others, have appeared in the literature but all have their limitations.

1.2 Problem Statement and Goals

Metamaterials and plasmonics have been proposed as a means or achieving hypothetically unlimited image resolution[3]. However, all of the so-far developed superresolution devices have specific limitations than can be overcome to obtain images with higher spatial resolution. The goal of the research contained in this dissertation was to solve the following issues limiting the spatial resolution of metamaterial and plasmonic lenses: the absorption of electromagnetic waves which reduces image contrast and resolution, the reconstruction of images formed by plasmonic lenses using only intensity data, and the noise limit which imposes an upper bound on the superresolution capabilities.

A peripheral goal was later taken on to improve the resolution performance of a conventional far-field imaging system using the same principles that were applied to the metamaterial and plasmonic lenses. This would provide both experimental evidence for our theoretical models, as well as demonstrate the versatility of the methods we developed.

1.3 Summary of Research

In this dissertation, first provided is a review of near-field optics and superlenses demonstrated in the literature which can achieve superresolution. Then an explanation is systematically developed of how to improve the spatial resolution of imaging systems by performing operations on their spatial Fourier spectra. These systems can be metamaterials, plasmonic films, or even free-space optical systems.

Chapter 2 is a contemporary review of some near-field optical imaging techniques. In particular, a focus is placed on the improvements afforded by integrating plasmonic superlenses into scanning near-field optical microscopy (SNOM) systems. Also introduced is the concept of plasmon injection as a possible route for improving the imaging perfomance of such lenses and near-field optical systems.

In Chapter 3, a near-field imaging system is studied, which is composed of a homogenous slab of material with refractive index n = -1 + in''. The effect that n'' has on the imaging performance of the slab is characterized through numerical simulation, and it is shown that deconvolution post-processing is one means to compensate for the attenuation of the Fourier content in the object. It is then shown that this deconvolution is essentially emulating the physical process of injecting a high spatial frequency object into the lens at the outset, which can give the original object at the output with no post-processing required.

In Chapter 4, the concepts of Chapter 3 are applied to a more realistic scenario; a silver slab superlens which is excited by incoherent light. The incoherence of the light allows for proper reconstruction of the images with only intensity data, removing the need to perform delicate phase measurements in the near-field of the superlens. The same physical emulation idea as in Chapter 3 is demonstrated, but this time with incoherent intensity distributions and not coherent fields.

Chapter 5 is an expansion of the work in Chapter 4. Since the ultimate limit for deconvolving images is the noise level in the image spectrum, a method is developed to surpass this limit by spatially filtering the light emanating from the object with a band-pass filter that has transmission beyond the original noise limit of the superlens. By providing enough exposure to the object, an image with accentuated high spatial frequencies can be detected at the output, and then reconstructed to obtain resolution of an object with only 25 nm separation at an illumination wavelength of 365 nm. The technique is termed "active convolved illumination" and the spatial filter required for implementation in this plasmonic system can be realized by a hyperbolic metamaterial.

In Chapter 6, shot noise is identified as the main resolution limit for a conventional imaging system, and the work of Chapter 5 is adapted to fit this problem. Put briefly, the photons contributing to the image signal are reallocated to a portion of the Fourier spectrum near the original noise limit. In the image, there is increased contrast for the high spatial frequency objects. Experimental results are presented to support the theory. Then, the design and simulation of a metamaterial "superlens" is shown using the same Fourier principles. In contrast to many of the previous metamaterial and plasmonic flat lenses, this device can project the evanescent, high resolution image data into the far-field via a subwavelength diffraction grating. High fidelity and signal transfer for the desired diffraction orders is obtained by spatial filtering with an appropriately designed hyperbolic metamaterial. An imaging simulation shows that the device can resolve two slit objects separated by 45 nm at an illumination wavelength of 488 nm.

In short, the primary result for all these above applications is that an increase in the spatial resolution is demonstrated through both physical manipulation of the electromagnetic fields as well as through post-processing techniques.

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Chapter 2

Review of Near-field Optics and Superlenses for Nano-imaging 1

2.1 Introduction

The development of technologies such as scanning electron microscopy (SEM), scanning tunneling microscopy (STM), and atomic force microscopy (AFM) offered a dramatic new insight into physical structures at the nanometer scale. In contrast to conventional optical microscopy, avoiding the use of photons to gather image information at this length scale means that Abbe's diffraction limit[1] poses no problem.

¹Reproduced from W. Adams, M. Sadatgol, and D. Ö. Güney, "Review of Near-field Optics and Superlenses for Sub-diffraction-limited Nano-imaging," *AIP Advances* 6, 100701 (2016); doi:10.1063/1.4964498, with the permission of AIP Publishing.

However, the desire to do things such as imaging biomolecular processes in vivo and characterizing subsurface features in nanoelectronics requires us to return to optical methods and find a way to beat the diffraction limit. A simple analysis of optical imaging systems reveals that the Fourier components of an image greater than ω/c evanescently decay along the optical axis. Therefore, an obvious approach to imaging beyond the diffraction limit is to access the near-field within a distance $z \ll \lambda$ where the evanescent components containing high spatial frequency information are not yet fully attenuated. This is the basis for applying near-field optics to imaging beyond the diffraction limit.

The development of aperture scanning near-field optical microscopes (SNOM and many other acronyms) in the mid-1980s broke the diffraction limit barrier and achieved resolutions down to the 20-50 nm range. [2, 3, 4, 5, 6] This resolution was extended further by apertureless SNOM systems, which employed a nanoscale scattering probe in the near-field instead of a subwavelength aperture. [7, 8, 9, 10, 11, 12, 13, 14]

At the turn of the new millennium, imaginative new approaches for controlling electromagnetic waves began to appear for imaging, [15, 16, 17, 18, 19, 20, 21, 22, 23] photovoltaics, [24, 25, 26] quantum information processing and simulations, [27, 28, 29, 30, 31] wireless communications, [32, 33] and novel optical materials, [34, 35, 36, 37, 38, 39, 40] among many others. The advent of metamaterials with simultaneously negative permittivity and permeability [41] brought renewed interest in the properties of left-handed materials first proposed by Veselago, [42] which Pendry demonstrated could be applied to sub-diffraction-limited imaging with his perfect lens. [15] Pendrys seminal paper inspired the experimental verification of negative refractive index [43] and superlenses that demonstrate sub-diffraction-limited resolution in the near-field. [17, 44, 45] However, these superlenses only operate in the electrostatic limit, which restricts the operation to a single polarization of the light in the near field of the superlens. While there are designs for isotropic negative index metamaterials (NIMs), [46, 47] there exists no experimental 3D isotropic optical NIMs suitable for perfect imaging as envisioned by Pendry due to the constraints of nanofabrication and absorptive losses. [16, 48] Fortunately, a number of promising loss compensation methods have appeared, including the employment of gain media, [49, 50, 51, 52, 53, 54, 55, 56, 57, 58, 59] geometric tailoring, [60] and plasmon injection. [61, 62, 63, 64, 65]

Somewhat surprisingly, there is less detailed treatment given in the literature to the practical matter of extracting the image from a near-field superlens than one might expect. Since the image is focused into the near-field, one may wonder why there is any interest in superlensing at all if there is still a need for additional employment of a scanning microscope. Consequently, a goal of this review is to point out some recent exploration of this point and possible future directions, in addition to first providing a comprehensive and contemporary background in relevant near-field optics and superlenses for nano-imaging. We hope that this will provide both inspiration for

new research and a useful reference for researchers studying superlenses and related topics.

2.2 Aperture Scanning Near-Field Microscopy

2.2.1 Aperture SNOM Concept

In a clear example of being ahead of his time, Synge first proposed the idea of employing a small subwavelength aperture scanned across a sample for nano-imaging in 1928.[66] His idea was essentially to illuminate an opaque screen with a subwavelength hole to act as a light source, selectively illuminating small features of a specimen when brought within a distance smaller than the aperture diameter. Fig. 2.1 shows an illustration of his proposal. Synge even explored the application of piezoelectrics to his microscope.[67] Due to the experimental constraints of the time however, he was never able to realize his ideas in practice. A more developed proposal of the same concept was also published in 1972.[68]

In the 1980s, the development of the STM at IBM essentially provided all of the necessary micropositioning technology for Synge's original proposal.[69] Aperture SNOM systems were successfully demonstrated at IBM[2, 4, 70] and Cornell University[3, 5, 6, 71] shortly afterwards. Instead of an opaque screen however, these



Figure 2.1: Illustration of the scanning aperture concept for nano-imaging. (a) Two features of a specimen in free space are separated by a distance $\Delta < \lambda_0/2$. When illuminated by a plane wave, the resulting intensity pattern in the far-field is unresolved due to the diffraction limit. (b) A small aperture of diameter $a \ll \lambda_0$ under the same illumination is brought within a distance z < a to a specimen and scanned in the lateral direction. Since the near-field radiation pattern is localized to the aperture itself, each feature is illuminated separately provided that the spatial frequency characterizing the separation of the two features is below the cutoff spatial frequency of the aperture. In other words, a should be smaller than Δ . After the full scan is completed, the result is a resolved image.

microscopes utilized a tapered fiber probe coated with a metallic film which formed

a small aperture at the probe tip.

2.2.2 Near-field of a subwavelength aperture

To understand the aperture SNOM concept, it is illustrative to very briefly analyze the field distribution from an illuminated subwavelength aperture. As a simple example,

consider a slit aperture of subwavelength width $a \ll \lambda_0$ in an infinite opaque screen under monochromatic plane wave illumination with free space wavelength λ_0 as in Fig. 2.2. In this case, it is informative to analyze the fields in the spatial frequency domain. The wave vector component along the optical axis is

$$k_z = \sqrt{k_0^2 - k_\perp^2}$$
 (2.1)

where $k_0 = 2\pi/\lambda_0$ and k_{\perp} are the magnitude and transverse component of the wave vector, respectively. For an aperture of width a, it is evident that the highest spatial frequency which can be discriminated is given by $2\pi/a$, defining the cutoff spatial frequency. To collect the highest spatial frequencies possible as one would like to do with SNOM, it is appropriate to say that the aperture will function near cutoff such that $k_{\perp} \approx 2\pi/a$. Subsequently, $k_{\perp} > k_0$ for the subwavelength aperture and in accordance with Eq. 1, k_z becomes imaginary and the resulting field from the aperture is evanescent. Therefore, the resulting field distribution is highly localized to the aperture. This is an important point, since this confinement of the field near the aperture is what enables the aperture to locally probe features on a sample. The resolution of the aperture SNOM is then not solely limited by the incident wavelength, but rather is strongly dependent on the size of the aperture.

More rigorous electromagnetic calculations of the field distribution of a subwavelength



Figure 2.2: Illustration of the field distribution from a subwavelength aperture under plane wave illumination. In the near-field zone, the radiation is strongly confined to the aperture. In the far-field, the resulting diffraction pattern will become increasingly spread out as the aperture size is decreased, however this poses no problem for SNOM since as in Fig. 2.1 only small features of the sample are excited at each point in the image scan and highlyresolved image contrast can be achieved.

aperture are given by Bethe[72] and Bouwkamp.[73] Similar calculations for the specific case of a SNOM aperture probe were later carried out by Drezet et al.[74, 75] However, the brief analysis presented here is sufficient to understand the localization of the near-field radiation relevant to the general aperture SNOM concept.

2.2.3 Aperture SNOM probes

SNOM aperture probes have been realized using a variety of fabrication techniques. Ideally, the probe should combine small aperture size and high transmission to achieve both high resolution and signal-to-noise ratio, respectively. For the most part, aperture probes consist of a tapered optical fiber coated with a metal that forms an aperture at the very tip. The fiber taper can be realized by either chemical etching[76, 77, 78, 79, 80] or heating and pulling.[77, 81, 82] Fig. 2.3 shows the typical processing techniques for producing tapered fiber aperture probes. The chemical etching method was first performed by coating a fiber with an organic protective layer and subsequently dipping it into a HF etching solution.[76] The taper is formed by the meniscus between the organic layer and the HF, with the taper angle being varied by the use of different organic protective layers.[78] This method does result in some surface roughness on the sides of the taper, which leads to the formation of unwanted pinholes in the subsequently deposited metallic coating. This problem was solved by the so-called tube etching procedure, where the etching is done with the fiber cladding intact, leaving a tube within which the etching takes place.[79, 80]

Alternatively, the heating and pulling method has been used to produce SNOM fiber tips with a high degree of smoothness. In short, a fiber is heated by either a laser or filament and simultaneously mechanically stretched to the breaking point. Here, the transmission properties can be modified by the heating temperature, heating area, and the pulling parameters.[81, 82] This method often results in a flat tip at the aperture plane, but the resulting probe taper is relatively long compared to the chemical etching technique. To achieve high transmission, the taper angle should be as high as possible.[83] The ability to produce much shorter cones and high taper angle



Figure 2.3: Typical fabrication techniques for tapered fiber aperture probes. The two common approaches to producing tapered fibers are (a) chemical etching and (b) heating and pulling. After formation of the fiber tip, subsequent (c) metal deposition results in an aperture at the tip apex which provides light confinement.

is therefore an advantage of the chemical etching method, since higher transmission will lead to higher near-field signal during the microscope operation and in turn higher signal-to-noise ratio. Further engineering of the taper properties can also be beneficial, for example the triple-tapered probe[84] and the corrugated probe[85] which both demonstrated dramatically improved light throughput.

After forming the taper at the fiber tip, the next step is to create the aperture. This is usually done by evaporation of aluminum onto the sides of the fiber taper surface in a configuration such that the apex of the taper is minimally exposed to the evaporated metal. The result is commonly an aperture of diameter around 80-100 nm. There are limitations to the aperture diameter other than the evaporation conditions however.
The diameter can be no smaller than twice the skin depth of the metal coating, and sufficient light throughput for very small apertures can cause substantial heating which may adversely affect imaging results.[86, 87]

Aperture probes other than the tapered fiber have additionally been demonstrated. Hollow pyramidal/campanile metallic probes produced by conventional semiconductor fabrication methods were demonstrated in Ref. [88]. Also, batch fabrication of Si_3N_4 tips by plasma enhanced chemical vapor deposition was shown in Ref. [89]. These designs aimed to alleviate the problems associated with reproducible fabrication of high quality probes and the relatively small taper angles offered by fiber probes.

2.2.4 Implementation

Because it was the first form of scanning near-field optical microscopy, the earliest applications of SNOM to sub-diffraction-limited imaging were of the aperture probe type. An experimental setup diagram for a typical aperture SNOM system is shown in Fig. 2.4. In the configuration shown, a laser beam is coupled into the untapered end of an aperture probe. The probe is attached to an oscillated tuning fork or similar mechanical device to provide shear force feedback for control of the probesample distance.[90, 91] The probe then illuminates a sample, which is scanned and positioned with piezoelectrics. The light scattered from the sample is collected by a conventional microscope objective. The wavelength of interest is then selected by a monochromator, dichroic mirror, or filter before reaching the detector. Amplitude and phase contrast images can be formed by raster scanning the sample. The system can be operated in either illumination, collection, illumination/collection, reflection, or reflection/collection modes as outlined in Fig. 2.5.

2.3 Apertureless Scanning Near-Field Microscopy

2.3.1 Apertureless SNOM Concept

In the interest of historical completeness, it is worth mentioning that the concept in Fig. 2.1 was not Synge's only idea for sub-diffraction-limited imaging. He had also proposed the employment of a small scattering particle situated very close to a specimen in a correspondence to Einstein.[92] However, Einstein expressed some concern about the ability to differentiate the scattered signal from artifacts due to direct illumination of the specimen, leading to Synge's published aperture design.

Due to the lower limit on aperture size for aperture SNOM, such microscopes cannot achieve resolutions lower than about 20-30 nm. However, the development of aperture SNOM was in large part the basis for future near-field optical microscopes.



Figure 2.4: Experiment diagram of a typical aperture SNOM system. In this configuration, a laser beam is coupled into a tapered fiber aperture probe which illuminates a sample. The probe is attached to a shear force mechanism, often a tuning fork, to provide mechanical feedback for control of the probe-sample distance. Piezoelectric micropositioning hardware controls the sample in all three directions. The light scattered from the sample is collected with a conventional microscope objective, and then filtered before reaching a detector.

The observation of strong field lobes at the edges of aperture probes suggested that higher resolution is available by simply reducing the probe to a point.[83, 93, 94] In direct analogy to STM, the photon scanning tunneling microscope (PSTM) was demonstrated, which detected the evanescent waves on the boundary of a prism under total internal reflection with a bare, tapered optical fiber.[95, 96, 97, 98, 99, 100]



Figure 2.5: Operation modes for aperture type SNOM. Depicted from left to right in the figure is illumination, collection, illumination/collection, reflection, and reflection/collection modes.

Shortly after, researchers demonstrated similar apertureless SNOM systems that used a nanoscale tip probe instead of an aperture probe. [7, 8, 9, 10, 11, 12, 13, 14] This type of microscope is often constructed from an AFM with a cantilever tip that is formed of either a dielectric, semiconductor, or metal. Apertureless SNOM quickly enabled resolutions down to the 1 nm scale. [11] Additionally, there are a number of possible probe geometries which can be realized and exhibit distinct optical characteristics. [101] An illustration of the apertureless SNOM concept is shown in Fig. 62.6.

As with any microscopy technique, some mechanism of contrast must be present for apertureless SNOM. However, some of the earlier aperture and apertureless SNOM results in the literature were plagued by the appearance of topographic artifacts with the optical signal.[102] Generally, the desired contrast is resulting from a difference in optical properties of the constitutive regions on a sample. To address this, an improved interferometric detection technique was developed to achieve pure optical contrast images in the absence of topographical artifacts.[103] This method involved



Figure 2.6: Illustration of the apertureless SNOM concept. Two features of a sample separated by a distance $\Delta < \lambda_0/2$ can be resolved with local excitation by a sharp probe.

demodulating the detected signal at harmonics of the probe tapping frequency by lock-in amplification, which is discussed in more detail later.

2.3.2 Near-field of a sharp probe

Apertureless SNOM operates in general by introducing some optical perturbation in the near-field with a sharp probe. The probe can be either treated as a scattering center or a light source, depending on whether the fields resulting from sample illumination (scattering center) or tip illumination (light source) are dominant. The two conditions are not independent however, as the scattering efficiency in the former is related to field enhancement in the latter. Generally, the light source scheme is



Figure 2.7: Finite element simulation of the field distribution of an illuminated gold probe tip with polarization (a) parallel and (b) perpendicular to the probe axis. A strong field enhancement is shown at the tip apex in (a).

desired when the detected radiation is frequency-shifted from the excitation as in the case of spectroscopy. To treat the probe as a source, there must be some enhancement of the incident fields by the probe tip. It was discovered that the incident radiation should be polarized parallel to the probe axis for maximum field enhancement at the tip apex.[104, 105, 106] Field distributions of two different polarizations incident on a metallic probe are shown in Fig. 2.7.

Qualitatively, two effects are present in the case of the illuminated probe, first being the lightning rod effect due to small radius of curvature, [107, 108] and second being the plasmon resonance due to electron oscillations. [107, 109, 110] In reality the field enhancement depends strongly on the probe geometry and sample surface topography. [111] Naturally, numerical simulation provides a useful guide for design when analytical approximations begin to break down. However, care must be taken when simulating geometries with abrupt curvature and high field gradients. Detailed investigations of the field enhancement offered by different probe tips have been conducted using the multiple multipole,[112, 113] finite element time domain,[114] finite difference time domain,[101, 115] and finite element[108] techniques. The field enhancement at the apex of sharp probe tips has allowed for apertureless SNOM fluorescence imaging[116] down to single molecules.[117, 118]

2.3.3 Apertureless SNOM probes

The probe tip geometry and material is of great importance to the SNOM imaging performance. A fair amount of work has been published studying dielectric and semiconductor probes.[8, 11, 112, 118, 119, 120, 121, 122, 123] Metallic tips are usually preferred due to their more favorable scattering, confinement, and field enhancement characteristics.[14, 105, 116, 123, 124] Some study in the past decade however has shown that silicon tips are a promising solution for achieving the highest resolution with apertureless SNOM due to more capable semiconductor fabrication techniques.[113, 125]

One group combined the benefits of both the aperture and apertureless approaches by growing a thin metallic wire on an aperture probe tip.[126] This method directly coupled light from the aperture to the wire for sample excitation, and the background signal caused by direct tip illumination is reduced. Other geometries include a small metallic particle situated on a dielectric tip.[127, 128, 129, 130, 131, 132] Also, a microscope was demonstrated where a small metallic particle was trapped and positioned by optical forces.[133] Very recently, a novel corrugated probe geometry has been proposed.[134] This probe takes advantage of the additional degree of freedom in surface plasmon dispersion offered by engineering of so-called spoof SPPs first put forth by Pendry.[135]

2.3.4 Application

Apertureless SNOM systems are often built around an AFM operated in tapping mode. The amplitude of the probe tapping is controlled by a feedback mechanism monitored by displacement of a beam deflected off the AFM cantilever head to a position-sensitive photodiode. This tapping feedback control allows the tip to maintain a constant average tip-sample separation. A diagram of the typical experimental setup is shown in Fig. 2.8. Due to significant background radiation, there is some additional care needed for detection of the imaging signal of interest. Separation of the near-field signal from the background is usually done by means of either homodyne or heterodyne detection with a lock-in amplifier. In the homodyne detection scheme, the near-field signal is collected at the tapping frequency or some higher harmonic (2nd or 3rd). In contrast, the heterodyne scheme collects the signal at the tapping frequency mixed with a known offset frequency. Harmonic generation occurs due to the nonlinear dependence of the near-field signal on the tip-sample distance described by the signal Fourier series expansion.[103, 136, 137, 138] The signal-to-background ratio is increased by demodulation to higher order harmonics. Utilizing this technique allows for interferometric collection of pure optical contrast images.[103, 139, 140] Furthermore, heterodyne detection can improve the suppression of background fields leading to better imaging results.[138]

In addition to standard optical and fluorescence imaging, apertureless SNOM can be employed in high resolution infrared microscopy and Raman spectroscopy. Of considerable importance for many researchers in chemistry is the application to Raman spectroscopy, which gives information about the low frequency vibrational modes of molecules. The principle challenge of Raman spectroscopy is the low light output. A revolution came when surface-enhanced Raman scattering (SERS) was developed, which allowed for a large enhancement of the Raman light yield.[141, 142, 143, 144] In SERS, samples are adsorbed onto a rough metallic surface. The enhancement offered by SERS can be attributed to field enhancement by excitation of localized surface plasmons near the adsorbed sample molecules. Tipenhanced Raman spectroscopy (TERS), operates on a similar principle. Leveraging the increased Raman scattered light output of SERS with the highly localized field enhancement of apertureless SNOM, it is possible to create sub-diffractionlimited spectroscopic images with TERS that provide sufficient light output for



Figure 2.8: Experimental diagram for the implementation of a typical apertureless SNOM system. In principle, an apertureless SNOM can be constructed by the addition of an optical source and the corresponding detection hardware to an existing AFM. The configuration shown is excited by a laser beam focused by a microscope objective onto the apertureless probe on the opposite side of the sample. The resulting signal is collected by the same objective, and the signal is then filtered before reaching a sensitive photodetector. The detector signal is processed by lock-in amplification at some harmonic $n\Omega$ of the probe tapping frequency Ω (homodyne) or a mixed frequency $\Delta + n\Omega$ (heterodyne), where Δ is a controlled offset frequency. The tapping amplitude is controlled by the deflection of a laser beam off the probe head, which is detected with a position-sensitive photodiode. Positioning of the probe is provided by piezoelectrics.

detection. [145, 146, 147, 148, 149, 150, 151, 152, 153, 154, 155, 156, 157, 158, 159, 160]

2.4 Superlens for Sub-diffraction-limited Imaging

2.4.1 Development of Superlenses

In 2000, Pendry proposed his original idea of a perfect lens that provides both focusing of propagating modes and amplification of evanescent modes.[15] His original work achieved this using a slab of material with refractive index n = 1 that perfectly transfers the original fields at the object plane to the image plane opposite the lens, as shown in Fig. 2.9. Pendry also mentioned in his paper that sub-diffraction-limited imaging could be achieved in the electrostatic limit with a slab of silver. The reason for this proposal was that silver, like other noble metals, behaves approximately as a plasma at optical frequencies with dielectric function given by

$$\varepsilon(\omega) = 1 - \frac{\omega_p^2}{\omega^2} \tag{2.2}$$

where ω_p is the plasma frequency. Inspection of Eq. 2 reveals that ε becomes negative when $\omega < \omega_p$, and is exactly 1 when $\omega = \omega_p/\sqrt{2}$. This condition is where superlensing occurs in free space, and also where there exists high-k surface plasmons on the lens interface.[161] This suggests a link between sub-diffraction-limited imaging by evanescent wave amplification and excitation of surface plasmons. It should



Figure 2.9: (a) Focusing of propagating waves and (b) amplification of evanescent waves with Pendry's perfect lens.

be noted however that this simple model describes an ideal plasma and begins to fail in the practical case where loss in the metal is nonzero. In this case, the permittivity becomes complex and damping of the electron oscillations occurs. Nonetheless, the sub-diffraction-limited imaging capabilities of silver planar superlenses were first demonstrated experimentally in 2005.[17, 45, 162] It should be mentioned however that the cited resolutions with respect to the diffraction limit do not take into account the shortening of the incident wavelength by the media immediately surrounding the lens.

Other superlens designs include left-handed lenses for GHz frequencies, [163, 164, 165] photonic crystal lenses, [166, 167, 168, 169] and metal-dielectric layered lenses. [19, 170, 171] In the next section, other designs will be reviewed that consider the implementation of SNOM with the superlens.

2.4.2 Leveraging Superlensing and Near-field Optics for Imaging

Upon initial inspection, the practical use of a superlens may be questionable to some, since the image formed in the near-field must be read by a scanning microscope. This begs the question as to what a superlens adds to an imaging system, since the image could instead be directly read with SNOM. A principle objective of this review is to point out that, in fact, there is evidence to support that employment of superlenses in conjunction with SNOM provides benefits over SNOM alone.

In 2006, Taubner et al. demonstrated near-field microscopy with a SiC superlens.[172] Their experiment was configured as in Fig. 2.10. A SiC slab was coated on both sides with SiO₂. The object was a gold film with subwavelength holes evaporated on one side of the superlens structure. On the opposite side of the lens, the probe of an apertureless SNOM operating in the scattering configuration was placed. Illumination was provided by a frequency-tunable CO₂ laser operating in the mid-IR regime. By tuning the illumination wavelength to the superlensing condition around 11 μ m, the authors demonstrated sufficient amplitude and phase contrast with an interferometer to distinguish the holes on the opposite side of the lens. However, when the wavelength was tuned to around 9.25 μ m where no superlensing occurs, the contrast disappeared.



Figure 2.10: Experimental characterization of SiC superlens with near-field microscopy. (a) Diagram of the experiment showing an SNOM probe in the scattering configuration near a SiC superlens. Subwavelength holes in a gold film shown in the SEM micrograph in (b) were imaged by the SNOM-superlens system. (c) Amplitude and (d) phase contrast images were taken showing resolution of the hole structures. (e) Sufficiently away from the superlensing condition, the holes are not resolved. Reprinted with permission from T. Taubner, D. Korobkin, Y. Urzhumov, G. Shvets, and R. Hillenbrand, Science 313, 1595 (2006). Copyright 2006 AAAS.

In 2011, Kehr et al. reported on a similar experiment using a perovskite oxide superlens shown in Fig. 2.11.[173] As with the SiC superlens, an increase in contrast was observed near the superlensing condition. Comparison of the resulting superlens images with control images taken of bare objects with SNOM highlighted the performance enhancement of the superlens. Additionally, a superlens-enhanced probesample coupling was observed at some distance from the sample due to phonon polariton resonance at a wavelength near but away from the superlensing condition. Applications of this phenomenon include heat-assisted data recording, local thermal



Figure 2.11: Experimental characterization of perovskite superlenses with near-field microscopy. From left to right, each subfigure shows the experimental configuration, AFM images of the sample topography, and near-field images. (a) Bare SrRuO3 objects were imaged with SNOM. (b) Symmetric and (c) asymmetric superlenses were deposited onto the objects resulting in resolved images near the superlensing wavelength and an improvement of the results in (a). Away from the superlensing wavelength, the objects disappear. Reprinted with permission from S. C. Kehr et al., Nat. Commun. 2, 249 (2011). Copyright 2011 Macmillan Publishers Ltd.

sensors,[174] and metamaterial-based multifunctional circuits.[175, 176] Perovskites are also attractive because there are multiple options that can satisfy superlensing in the mid-IR and are ferroelectric so they may be tunable by an applied field. Consequently, layered perovskite superlenses could be used as bandpass filters for near-field spectroscopic imaging.[177]

2D materials like graphene are promising for superlens-type evanescent field growth

for continuously-tunable imaging of subsurface structures at IR and THz frequencies while additionally avoiding the practical challenges of depositing a high quality (i.e. low loss) superlens. [178] In 2012, a broadband layered graphene lens was proposed which displayed a simulated resolution of over $\lambda/10.[179]$ Particular applications of graphene superlenses include noninvasive imaging of nanowire doping concentrations, material growth defects, subcellular biological imaging, vibrational absorption microscopy, and material identification. The problem of depositing a large high-quality superlens can also be avoided by instead utilizing a local superlens. [180] Here a small BaTiO3 lamella with subwavelength dimensions was placed over an imaging area in free space. The resulting sub-diffraction-limited resolution underlines the localized nature of the excited phonon polariton modes, and also opens the possibility of depositing local superlenses on areas of interest on a sample instead of a larger lens that is likely to contain material defects or unwanted topographical features. Additionally, SNOM scan times would be reduced by scanning a smaller area. Obviously, this method inherently requires knowledge of where on the sample a useful image could be obtained, which is not always the case with microscopy.

Recent investigations of a GaAs superlens by Fehrenbacher et al.[181] have reached similar conclusions about superlens imaging enhancement as in Ref. [172] and Ref. [173]. A doped GaAs (doping concentration $n = 4 \times 10^{18} \text{ cm}^3$) superlens was deposited between two intrinsic GaAs layers on top of 2 μ m gold stripe objects. Comparison with images collected without the superlens showed a clear contrast enhancement of the sub-diffraction-limited image by the superlens configuration. Also, exploration of imaging a gold particle with SNOM and a 50 nm Si_3N_4 cap showed that the highest SNOM resolution can likely be obtained at the superlensing condition of the cap layer of any object (e.g. cap on a nanoelectronic device), provided that the imaginary part of the cap material permittivity is sufficiently small.[182]

To avoid the need for a near-field readout of the image from a superlens, there has been work on other classes of sub-diffraction-limited lenses which function by converting evanescent modes into propagating modes. In 2006, a so-called far-field superlens (FSL) was first developed which added a grating structure to a metallic film as in Fig. 2.12 (a).[183, 184, 185] This lens can be thought of as operating similarly to structured illumination microscopy,[186] where moiré fringes are formed at the spatial frequency difference between an object field and incident patterned illumination. In the case of the FSL, the patterned illumination is replaced by a grating with grating wavenumber Λ . The resulting spatial frequency content of the image is then downconverted to $k' = k - \Lambda$, and for $k \leq 2k_0$ the image content can be transferred to the far-field.

Another attempt to bring superlensing to the far-field was also first proposed in 2006, called the hyperlens due to the employment of metamaterials with hyperbolic dispersion.[18, 20, 22, 23, 187] Hyperbolic metamaterials are of interest for imaging because they can convert arbitrarily large-k modes into propagating modes by specific engineering of metal-dielectric interfaces. Fig. 2.12 (b) shows a schematic of an



Figure 2.12: (a) Far-field superlens and (b) hyperlens for sub-diffractionlimited imaging. (b) Reprinted with permission from Z. Liu, H. Lee, Y. Xiong, C. Sun, and X. Zhang, Science 315, 1686 (2007). Copyright 2007 AAAS.

experimental hyperlens.

While these lenses bring high spatial frequency content into the far-field, they are still near-field optical systems in the sense that the lens must be situated within the near-field of the object. In this way, far-field superlenses and hyperlenses are comparable to apertureless SNOM where some structure in the near-field of an object scatters evanescent waves to a detector in the far-field. This is in contrast to a conventional lens which only sees propagating waves. Other far-field optical super-resolution methods that rely on point spread function engineering or stochastic excitation like stimulated-emission-depletion microscopy[188] and stochastic optical reconstruction microscopy[189] do exist but are often subject to stringent illumination and sample fluorescence conditions.

There are a few obstacles to consider with the FSL and hyperlens as well. By inspection of the moiré fringe effect exploited by the FSL, it can be seen that there will be resolution enhancement only up to $2k_0$ unless the numerical aperture is increased or some nonlinear harmonic response is introduced, as appearing in saturated fluorophores. [190] Also, the lens grating must be reoriented rotationally multiple times to cover the entire bandwidth $|k'| = |k - \Lambda|$, adding complexity to a scheme which already requires substantial postprocessing. With the hyperlens, the object plane is constrained to a small size near the inner surface of the lens. Also, the problem of geometric aberration exists while under plane wave illumination, as in most practical cases. The highest image resolution that has been experimentally achieved so far using hyperlenses corresponds to about one third of the free space optical wavelengths. [23, 191] This is mainly due to the inherent losses[23, 60, 65] in the constitutive components of the hyperlens and finite discretization of the metamaterial. [192, 193] It is still worth pursuing the development of the hyperlens, since there may be methods found that alleviate these concerns. Another possibility is new ways to utilize the hyperlens such as illuminating it in reverse as a subwavelength focusing device for nanolithography or high-density optical data storage. [194, 195]

2.5 Future Outlook

2.5.1 Future Research for Superlens Imaging

Many of the superlenses demonstrated in the literature take the form of a planar slab of some conventional material, e.g. silver or silicon carbide, which can be deposited or grown by traditional microfabrication techniques. These lenses largely function in the electrostatic limit for a single polarization of the light in the near field, though for general imaging purposes the superlens ideally would exhibit both an electric and magnetic response so that arbitrarily polarized fields can be focused and resolved. This requires the development of bulk isotropic NIMs or suitable photonic crystals. Unfortunately, the problems of absorptive loss and the fabrication of miniaturized metamaterials have been major hurdles for the realization of isotropic NIMs operating in the optical spectrum. Methods for compensating losses in optical metamaterials have emerged in the last decade. Most notably, the application of gain media has appeared in the literature. [49, 50, 51, 52, 53, 54, 55, 56, 57, 58, 59] There are problems, however, with the employment of gain media in metamaterials as Stockman has pointed out. [196] Notably, full or overcompensation of losses with gain leads to instability that brings the metamaterial into a spaser state. Particularly for imaging it will be difficult to preserve in the metamaterial the required amplitude and phase relations of the underlying optical modes imposed by the object to be imaged.

Recently, a novel method for achieving full loss compensation in negative index metamaterials was developed, known as the plasmon injection (Π) scheme. [63] In contrast to many other loss compensation methods, the Π scheme requires no gain medium, thus eliminates its associated complexities. Instead, full compensation is achieved not by traditional amplification but by the coherent superposition of externally-driven "auxiliary" modes with the eigenmodes of a NIM. Surface plasmon polaritons are injected from auxiliary ports through a metallic grating structure, which are then superimposed with the native eigenmodes in the metamaterial. The Π scheme allows for diverging figure-of-merit (i.e., loss free metamaterial), and has been demonstrated to provide near-perfect loss compensation for superlenses [63, 64] and hyperlenses [65] applied to sub-diffraction-limited imaging. For imaging, applying the Π scheme is equivalent to a simple spatial filtering procedure, or superposition of the original source with an auxiliary source as shown in Fig. 2.13. Further development of the Π scheme concept could realize the goal of enhanced sub-diffraction-limited imaging with lower losses. There are also experimentally demonstrated imperfect negative index flat lenses [197, 198] at optical wavelengths which can benefit from the Π scheme for resolution beyond diffraction limit.

Based on the literature reviewed here, further exploration into leveraging the imaging



Figure 2.13: Loss compensation in a superlens with the Π scheme. The object to be imaged is superimposed with an auxiliary object, compensating the high spatial frequency components that are attenuated due to loss in the lens. The result on the image plane is a reconstruction of the original object. The same result can be achieved through post-processing by inverse filtering the raw image produced by the lens.

capabilities of SNOM and superlensing could lead to more readily available subdiffraction-limited imaging capabilities. One further key point to clarify would be if superlenses can provide any relaxation of the experimental constraints of SNOM, for example the mechanical wear on tapping-mode probes, the need for a feedback system to maintain constant tip-sample distance, and long scan times due to low nearfield signal. In particular, plasmonic and phononic interactions in SNOM-superlens systems has emerged as a non-trivial and highly complicated physical problem which should inspire more research. Applying loss compensation techniques such as the II scheme to existing superlens designs and novel optical metamaterials for imaging could be further investigated. Furthermore, integration of nano-optics and superlenses with microfluidic devices could provide lab-on-a-chip capabilities for imaging and characterization of biomolecular structures.[199]

2.6 Conclusion

In this paper, near-field optics and superlenses for imaging beyond Abbe's diffraction limit were reviewed. A main focus of the review is the integration of apertureless scanning near-field optical microscopy (SNOM) with superlenses for new sub-diffractionlimited imaging applications. It should be noted that there are some finer points of SNOM that were left out of this review. The reader is encouraged to refer whenever necessary to other reviews that cover some of these points in detail[200, 201, 202] and the textbook by Novotny and Hecht.[203]

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Chapter 3

Bringing the 'perfect lens' into focus by near-perfect compensation of losses without gain media¹

3.1 Introduction

Metamaterials provide unprecedented control of light for diverse applications such as wireless communications[1, 2], novel optical materials[3, 4, 5, 6, 7, 8], optical analog simulators[9, 10], photovoltaics[11, 12, 13], quantum manipulation of light[14, 15, 16],

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and imaging [17, 18, 19, 20, 21, 22, 23, 24, 25, 26], among many others. The extent to which an imaging system is capable of capturing high spatial frequency components of an incoming wave determines its resolution. Those components with spatial frequency greater than ω/c , where ω is the angular frequency of the wave and c is the speed of light in a medium, constitute evanescent modes that decay rather than propagate. In a conventional imaging system, the image detector is located far enough away from the source so that the evanescent modes are decayed beyond the sensitivity and noise level of the detector, i.e. in the far-field. Consequently, conventional imaging systems can only detect spatial frequencies up to ω/c . This is the so-called diffraction limit first discovered by Abbe[27]. In order to increase the resolution of imaging systems and retain spatial frequency components greater than ω/c , imaging with a slab of negative refractive index material was proposed [17]. This approach relies on the negative index material for focusing of propagating modes and amplification of evanescent modes incident on the slab. Unfortunately, current negative index metamaterial designs are not suitable for optical imaging due to the extreme sensitivity to absorptive losses in the constitutive components [28, 29]. A number of metamaterial loss compensation schemes using gain media have been proposed[30, 31, 32, 33, 34, 35]. However, the use of gain media for loss compensation can result in instability and spasing[36].

Previously, a loss compensation scheme that provides full compensation in negative index metamaterials without the need for a gain medium was proposed, called the plasmon injection or Π scheme[25]. In the Π scheme, loss compensation is achieved by

coherent excitation of the eigenmodes of a plasmonic negative index metamaterial by superimposing externally injected surface plasmon polaritons (SPPs) with the lossy domestic SPPs in the metamaterial [25, 37, 38]. Here, in analogy with optical amplifiers, the externally injected and domestic SPPs resemble the 'pump' and 'signal,' respectively. The plasmonic resonator structure presented in [25] is solely a proofof-concept device that functions only for normal incidence. However, the underlying loss compensation mechanism can be generalized to any negative index metamaterial structure or even homogeneous material and arbitrary angle of incidence, as long as the physical configuration is such that the injected fields can be superimposed coherently with the eigenmodes of the metamaterial or the homogeneous material. In the supplemental material of [25], a brief analytical calculation is carried out to demonstrate how the Π scheme could be applied to a flat silver superlens operating for a single polarization (i.e., so called 'poor man's superlens' under the electrostatic limit as considered in [17]) to compensate the absorption losses in the superlens. Interestingly, this purely physical phenomenon for loss compensation in the superlens has been shown to be equivalent to a simple spatial filtering post-processing algorithm. However, no imaging procedure is carried out explicitly, nor is any demonstration of imaging with the Π scheme intended in [25]. The question which naturally arises is 'can we use the Π scheme to enhance the performance of a metamaterial superlens, particularly 'Pendry's negative index flat lens' that is known to be extremely sensitive

to losses [28, 29, 30, 39]. In the current work, we exactly answer this important question, which is not considered in [25]. Therefore in the present work we demonstrate, for the first time, application of the Π scheme to sub-diffraction-limited imaging with a 'non-ideal Pendry's negative index flat lens' (referred to as NIFL for short in the rest of the paper). By applying this loss compensation scheme, we achieve resolution of a previously unresolved sub-diffraction-limited object.

Unlike previous near-field negative index flat lens imaging systems, this technique does not benefit from a lossless negative index material and has no gain requirements. The technique developed here is based on a NIFL with a practical value for loss. Recently, a similar spatial filtering approach to countering losses was proposed which also considered tuning the material parameters of the NIFL and surrounding media[40]. However, the optimum values for loss in the NIFL that were assumed are around one to two orders of magnitude lower than what is used to obtain the imaging results presented here. Also, there is little deviation between the optimized material parameters for different spatial frequencies, which suggests the results may be sensitive to any small changes in those values.



Figure 3.1: Block diagram of the Π loss compensation scheme for imaging with a non-ideal Pendry's negative index flat lens.

3.2 Methods

Figure 3.1 shows the block diagram of the II scheme applied to imaging with a NIFL. The procedure begins by producing an image with a NIFL. Then, a filter is applied to the image that compensates the attenuation of the high spatial frequency components. This compensation filter is the inverse of the NIFL transfer function, which can be calculated analytically or numerically. We should note that the method of inverse filtering is well known in the field of image processing, however there are two distinctions to be made between the work presented here and traditional inverse filtering. First, the method in this paper provides compensation for evanescent waves. Secondly, the compensation of these decayed evanescent waves is intimately related to a physical phenomenon for loss compensation in metamaterials as described in [25]. The remainder of this paper explains the methods used to perform the scheme loss compensation procedure and form the resulting resolved image.

As previously mentioned, the transfer function of the NIFL can be found through either analytical or numerical calculation. Here, a numerical approach for determining the NIFL transfer function is presented. For any spatial frequency component, the transfer function can be described by the relationship between the electric field at the object plane and image plane. Therefore, in order to find the transfer function it is sufficient to send known plane waves with different spatial frequency k_y and measure the electric field at the image plane. Figure 3.2 shows the geometry for the NIFL transfer function calculation using the finite element commercial software package COMSOL Multiphysics. Periodic boundary conditions (PBC) are imposed on the top and bottom boundaries, however the simulation domain itself has limited extent in the y-direction. As a result of the applied PBC in the y-direction, k_y becomes a discretized quantity which can have values of $k_y = \pm m \frac{2\pi}{W_y}$, where W_y is length of the simulation domain in the y-direction and integer m = 0, 1, 2, ... Therefore, an increase in W_y results in more accurate transfer function in terms of number of data points, but also increases the computational domain and in turn the simulation time.

After defining the geometric parameters of the NIFL transfer function simulation, the next step is to define the optical properties of the NIFL itself. Consider $\varepsilon_r = -\varepsilon' + j\varepsilon''$ to be the relative complex permittivity and $\mu_r = -\mu' + j\mu''$ to be the relative complex permeability of the NIFL, where $\mu', \mu'', \varepsilon', \varepsilon'' \ge 0$ and $j = \sqrt{-1}$. Then, the refractive



Figure 3.2: (a) COMSOL simulation geometry for calculation of the NIFL transfer function. Plane waves with spatial frequency k_y are sent from the input port on the left and measured at the image plane on the opposite side of the NIFL. A perfectly matched layer (PML) is added at the right boundary to suppress the transmitted waves. (b) Cross section of the simulation showing the location of the object and image planes with respect to the NIFL.

index n of the NIFL is

$$n = -\sqrt{\varepsilon_r \mu_r} = -(\varepsilon' \mu' - j(\mu' \varepsilon'' + \varepsilon' \mu'') + \varepsilon'' \mu'')^{1/2}.$$
(3.1)

Since $\varepsilon'', \mu'' \ll \varepsilon', \mu'$, the $\varepsilon''\mu''$ term can be neglected, and the expression for the refractive index is simplified to

$$n \approx -(\varepsilon'\mu' - j(\mu'\varepsilon'' + \varepsilon'\mu''))^{1/2} = -\sqrt{\varepsilon'\mu'} \left(1 - j\left(\frac{\varepsilon''}{\varepsilon'} + \frac{\mu''}{\mu'}\right)\right)^{1/2}.$$
 (3.2)

Using the binomial approximation, equation 3.2 can be further reduced to

$$n \approx \sqrt{\varepsilon' \mu'} \left(-1 + j \frac{1}{2} \left(\frac{\varepsilon''}{\varepsilon'} + \frac{\mu''}{\mu'} \right) \right).$$
(3.3)

If the real part of the relative permittivity and relative permeability are considered to be 1, the relations for the refractive index and impedance z of the NIFL can be written as

$$n = n' + jn'' = -1 + jn'' \approx -1 + j\frac{(\varepsilon'' + \mu'')}{2}$$
(3.4)

and

$$z = \sqrt{\frac{\mu_r}{\varepsilon_r}} = \left(\frac{-1+j\mu''}{-1+j\varepsilon''}\right)^{1/2} = \left(\frac{(-1+j\mu'')(-1-j\varepsilon'')}{1+\varepsilon''^2}\right)^{1/2} \\ \approx \frac{1+j\frac{1}{2}(\varepsilon''-\mu'')}{1+\frac{1}{2}{\varepsilon''}^2} \approx 1+j\frac{(\varepsilon''-\mu'')}{2}.$$
 (3.5)

Considering the result of equation 3.5, it can be seen that setting $\varepsilon'' = \mu''$ results in an impedance match with free space. However, the effect on the imaging performance of an impedance mismatch introduced when $\varepsilon'' \neq \mu''$ is small compared to the effect of the imaginary part of the refractive index n'' in equation 3.4, which characterizes the absorptive loss in the NIFL. As an example to illustrate this, consider the case of $\varepsilon'' = 0.2$ and $\mu'' = 0.1$. From equations 3.4 and 3.5, the resulting n'' would be 0.15, however the imaginary part of z would be only 0.05. Therefore, for simplicity of analysis the case of $\varepsilon'' = \mu''$ can be chosen without much consideration of the effect of impedance mismatch. By inspection of figure 3.3, it can be determined that ideally the loss in the NIFL would be small in order to preserve the higher spatial frequency components of an image. Unfortunately, fabrication of negative index metamaterials with low loss operating at optical frequencies is difficult. Therefore, an n'' of 10^{-1} is selected for the rest of the analysis, which is reasonable given current fabricated structures[24]. This corresponds to a figure-of-merit of |n'/n''| = 10. Although having such realistic loss levels in the base materials that form the NIFL is sufficient to benefit from the II scheme, any further improvement in the loss characteristics of the base materials using different techniques[12, 30, 41, 42, 43, 44] can have a profound effect on the II scheme results.

To conclude the methods used here for characterization of the NIFL, a discussion of the effect of the NIFL thickness on the performance of the imaging system is required. Figure 3.4 shows the transfer function as the NIFL thickness 2*d* changes from $\lambda_0/2$ to $2\lambda_0$. As the results suggest, a decrease in the NIFL thickness reduces the attenuation of high spatial frequency components, which in turn increases the resolution of the imaging system. However, the thickness of the NIFL cannot be decreased arbitrarily for two reasons. First, the NIFL would be constituted by a metamaterial structure, the minimum thickness of which would be constrained by the size of the corresponding unit cell. Secondly, as the thickness of the NIFL decreases, the working distance of the lens also decreases, making mechanical alignment of the imaging system more difficult.



Figure 3.3: Transfer function $H(k_y/k_0)$ of the NIFL imaging system for different values of ε'' with $\varepsilon_r = -1 + j\varepsilon''$, $\mu_r = -1 + j\varepsilon''$, $d = 0.25 \,\mu\text{m}$, and $\lambda_0 = 1 \,\mu\text{m}$.

After characterizing the NIFL itself, the next step is to numerically evaluate the imaging performance. Figures 3.5(a) and (b) show the simulation geometry and material settings used to produce an image of some arbitrary object with the NIFL using COMSOL Multiphysics. The object is formed by defining the z-component of the electric field E_z over the object plane, and image is produced by recording E_z on the image plane. In figure 3.5(c), an object with three Gaussian features separated by 1 μ m is defined on the object plane, and the corresponding electric field on the image plane is recorded. Figure 3.5(d) shows a surface plot of the resulting field distribution



Figure 3.4: Change in the transfer function around $k_y/k_0 = 1\left(k_0 = \frac{2\pi}{\lambda_0}\right)$ as the NIFL thickness 2*d* is changed from $\lambda_0/2 = 0.5 \,\mu\text{m}$ to $2\lambda_0 = 2 \,\mu\text{m}$. The results suggest the employment of a thinner NIFL will result in better imaging performance.

over the simulation domain.

This imaging simulation can be repeated to produce the image from any object with arbitrary feature size. Once the image is formed, the resolution can be improved by applying an inverse filter to emulate the Π scheme for compensation of losses in the NIFL.



Figure 3.5: (a) The geometry and material settings used to perform numerical simulation of the NIFL imaging system. The extent of the model in the y-direction is chosen to be $24 \,\mu$ m, though the figure shown here is compressed in the y-dimension to better fit the page. (b) A cross section of the imaging system showing the working distance of the object and image planes from the NIFL. (c) The z-component of electric field E_z on the object and image planes with incident wavelength $\lambda_0 = 1 \,\mu$ m and $2d = 0.5 \,\mu$ m. (d) Surface plot of the E_z distribution over the imaging system simulation domain.

3.3 Results

In order to improve the resolution of the image obtained by the NIFL, it is important to amplify the suppressed spatial frequency components. A compensation filter is required to undo this attenuation made by the imaging system. Obviously, a proper choice for the compensation filter would be the inverse of the imaging system transfer function. This corresponds to the Π scheme loss compensation technique for imaging, where a portion (i.e., pump or auxiliary object) of the total incident field in the object plane can be thought of as coherently exciting the underlying modes of the system in order to compensate the losses in the other portion (i.e., signal or actual object to be imaged) [25]. The equivalent is applying a filter in the spatial frequency domain that amplifies the components with $k_y > k_0$. Figure 3.6(a) shows the compensation filter for the NIFL imaging system described in figure 3.5. As an example, an object with features separated by a distance $\lambda_0/4$, twice beyond the diffraction limit, was imaged by the NIFL. The results of this procedure are shown in figure 3.6(b). It can be seen that the sub-diffraction-limit features of the object are not resolved in the raw image produced by the NIFL. However, after applying the compensation filter a perfect reconstruction of the original object is achieved.

This procedure can be replicated for any arbitrary object field, provided that enough of the spatial frequency components required to reproduce the field are available



Figure 3.6: (a) Fourier spectra for the object, raw image produced by the NIFL, and compensated image with the corresponding filter $H(k_y/k_0)^{-1}$ resulting from an object with three Gaussian features separated by a distance $\lambda_0/4$. (b) Electric field intensities for the original object, raw image produced by the NIFL, and the compensated image after applying the filter shown in (a).

to be compensated by the post-processing. Therefore, the limitation to the smallest feature size one could resolve with this technique would solely be the noise floor of the detection mechanism at the image plane, in this case the numerical simulation. Since inverse filtering is prone to noise amplification, it is required to roll off or truncate the filter at some spatial frequency where the noise floor is reached on the image plane. In figure 3.6(a), it can be seen that the raw image spectrum begins to flatten around $k_y = 2.5k_0-3k_0$. Therefore, simply truncating the filter at $3k_0$ gives a good compensated image that avoids noise amplification at high- k_y . It is important to note that truncating the filter in this way requires no a priori knowledge of the object; only the detected raw image is needed. While this noise limitation is present in practice, there is no theoretical limit imposed on the compensation scheme presented here. The compensated image shown in figure 3.6(b) results solely from post-processing the raw image with the inverse filter in figure 3.6(a) without using any auxiliary source. To explain the link between the II scheme and such inverse filtering step, it can be shown that providing the appropriate auxiliary source with the original object field is equivalent to the inverse filter compensation scheme with no auxiliary source. By adding the auxiliary source, the imaging system is essentially being 'pre-processed' to physically inject high spatial frequency components of the incident field at the necessary magnitudes to reconstruct the original object at the image plane. This is analogous to providing power to the auxiliary ports in the plasmonic structure presented in [25]. For the present imaging system, the total field incorporating the appropriate auxiliary field can be calculated from the compensated image spectra in figure 3.6(a) using the transfer function of the NIFL. The resultant auxiliary input field, which is simply the difference between the total field and the object field, is plotted in figure 3.7.

The images resulting from the inverse filter alone and total input are compared in figure 3.8. There is some small deviation in the two images, which can likely be attributed to the numerical methods used. In the case of the inverse filtered image with no auxiliary source, Maxwell's equations are solved with the finite element method, and then that result is processed with discrete Fourier transforms. These two steps are also used to calculate the superposition of the original object and auxiliary input (i.e. total input), however the finite element method is again applied to obtain the


Figure 3.7: Auxiliary input electric field source. If the original object is thought of as the 'signal,' the field plotted here can be thought of as the 'pump.' The superposition of the auxiliary input and original object then becomes the total input field. The inset shows the finer structure of the auxiliary source. Fortunately, there is no need to 'pre-process' the imaging system in this way, since excitation of the NIFL with the superposition of this auxiliary input and the original object is equivalent to simply applying the inverse filter alone to the raw image in post-processing as evidenced in figure 3.8.

resulting image. The accumulation of numerical error as the imaging system is solved multiple times could likely be the source of the small discrepancy between the sole inverse filter and total input images. Another source of error which is important to point out is the width of the auxiliary input field in the spatial domain. In figure 3.7, it can be seen that the auxiliary input has a width of 80 μ m. This large aperture



Figure 3.8: (a) Fourier spectra for the total input field incorporating the auxiliary field in figure 3.7 calculated for an object plane length of 15 μ m (same simulation setup as figure 3.5) and image resulting from the total input. Fourier spectra for the original object, raw image produced by the NIFL with no auxiliary source, and the corresponding compensated image resulting from the inverse filter are reproduced from figure 3.6(a) for comparison. (b) Electric field intensities of the original object, image resulting from the total input field and the inverse filter with no auxiliary source are compared. Electric field intensity of the raw image (same as figure 3.6(b)) is also shown. It can be seen that the images from inverse filtering with no auxiliary source and the total input are equivalent with some small discrepancy in the total input image likely resulting from accumulated numerical error.

auxiliary input is used to calculate the image in figure 3.8(b) in order to minimize the error resulting from truncating the aperture size, however this error was observed to not have a strong effect on the image resolution after the width was increased to approximately 15 μ m.

It can be hypothesized that the Π scheme could be applied to compensation of decayed evanescent components in the absence of absorptive loss or negative index. To test this, calculations were performed to determine if the same compensation

scheme can be applied to the loss of high spatial frequencies due to diffraction in free space. This was done by performing the same calculations as in figure 3.6, but with the NIFL replaced by free space. The results are shown in figure 3.9. In contrast to the NIFL imaging system where absorptive loss dominates the transfer function characteristics [28, 29, 30, 39], in this case diffraction dominates and the transfer function drops more steeply. Upon initial inspection of figure 3.9(a), the spatial frequencies $k_y/k_0 > 1$ which are lost due to diffraction in the raw image can be somewhat recovered, though only up to $k_y/k_0 \approx 1.9$ where the noise floor is reached. This is as expected, since the evanescent components of the image decay faster and the noise floor is reached at a smaller spatial frequency for free space compared to the NIFL imaging system, where despite some material absorption, the NIFL still provides amplification for the evanescent components with respect to the free space. Also, free space does not perfectly preserve the phase of the propagating field components as is the case for the NIFL imaging system. Consequently, the compensation scheme is less successful and works for a narrower band than with the NIFL as in figure 3.6, limiting the resulting resolution. This is evident in figure 3.9(b), where an attempt to recover the object with $\lambda_0/4$ feature size is unsuccessful. Therefore, the advantages of applying the compensation scheme with the NIFL instead of free space is that the NIFL preserves a larger band of spatial frequencies that can be recovered by the compensation scheme, and it also provides perfect phase compensation for the propagating field components with $k_y/k_0 < 1$. In other words, it is still important to



Figure 3.9: Results of simulations performed as in figures 3.2 and 3.5 with the NIFL replaced by free space. (a) Fourier spectra for the original object, raw image, and compensated image, along with the inverse filter calculated from the free space transfer function. In this case, the filter compensates the loss of high spatial frequencies due to diffraction in free space rather than absorptive loss. (b) Electric field intensities for the original object, raw image, and compensated image are shown. Applying the compensation scheme to free space for the $\lambda_0/4$ object is clearly unsuccessful since the high spatial frequency components in the raw image are too quickly attenuated to recover.

include the negative index slab in the imaging system in order to successfully reconstruct images with sub-diffraction-limited feature size. However, it would be possible to preserve more spatial frequency components of the free space image if the image plane is moved closer to the object plane.

3.4 Conclusion

In this paper, a study of the optical characteristics and near-field imaging performance of a NIFL with a practical value for loss was performed. The optical properties of the NIFL were investigated analytically, and a numerical calculation of the transfer function was performed and studied. The simulation results yielded an unresolved image from the NIFL, which subsequently underwent loss compensation using an inverse filter that emulates the II scheme from [25]. This involves the simple postprocessing step of multiplying the raw image produced by the NIFL with the inverse of the transfer function in the Fourier domain. There are no requirements for electric or magnetic gain in the NIFL and surrounding media. The demonstrated result is a perfect reconstructed image with sub-diffraction-limited feature size. Our findings decouple the more-than-a-decade-long loss problem from the general problem of how to realize a practical 'perfect lens' operating in the optical frequencies, and reduce the problem mainly to amenable design and fabrication issues[45, 46, 47, 48, 49, 50, 51, 52]. Further developments in metamaterials and the Π scheme approach can lead to advances in other applications besides ultra-high resolution imaging such as photolithography and optical storage technologies.

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Chapter 4

Plasmonic superlens image reconstruction using intensity data¹

4.1 Introduction

The wave nature of light intrinsically imposes an upper bound on the resolution of farfield imaging systems[1]. However, the so-called diffraction limit has proved not to be unsurmountable. Notably, metamaterials enabling extreme control of light[2, 3, 4, 5, 6, 7] entered the arena of superresolution imaging after Pendry's original proposal of

¹Adapted with permission from W. Adams, A. Ghoshroy, and D. Ö. Güney, "Plasmonic superlens image reconstruction using intensity data and equivalence to structured light illumination for compensation of losses," *Journal of the Optical Society of America B* 34, 2161-2168 (2017); doi:10.1364/JOSAB.34.002161, C The Optical Society.

a "superlens realized by a slab of negative index material or, in the electrostatic limit, by a film of metal excited below the plasma frequency[2]. A review of metamaterial super-resolution imaging and near-field microscopy can be found in [8].

Practical realization of a true negative index superlens at optical frequencies is largely limited by the constraints of modern 3D nanofabrication capabilities[9], the complexity of isotropic designs [10, 11, 12], and the inherent absorptive loss in the metal components of negative index metamaterials [13, 14]. Plasmonic superlenses operating for TM-polarized waves, however, have been fabricated [3, 4, 15, 16, 17, 18, 19, 20, 21], but their performance is still limited by losses. Our research has taken a winding path toward solving the loss issue by unconventional means, as illustrated in Fig. 4.1. In the literature, losses in metamaterials are often countered by implementation of a pumped gain medium [22, 23, 24, 25, 26, 27, 28, 29, 30, 31]. However, this is not a desirable approach because the metamaterial can transition to an unstable spaser state, and the negative refraction can be destroyed due to the requirement for obeying causality [32]. Also, the few options for gain media operating at specific wavelengths or limited lifetimes preclude the design of a robust, low-loss metamaterial incorporating a gain medium. Modifying the geometric structure of the base metamaterial to alter the induced current distribution also has shown to be beneficial but can only produce limited compensation[33].



Figure 4.1: Illustration of our journey in the field of metamaterial loss compensation for super-resolution imaging. To avoid the use of gain media, we proposed a new method for compensating losses in metamaterials (plasmon injection) by adding extra energy with an auxiliary beam and coherently superimposing the auxiliary light with the lossy surface plasmon polariton modes in a metamaterial. We have at this point considered a few imaging devices for which to apply the plasmon injection scheme (1). They include a hypothetical imperfect negative index flat lens (NIFL) modeled as an effective medium with refractive index n = -1 + i0.1, a hyperlens based on a cylindrical layered hyperbolic medium, and a single-negative plasmonic superlens (SNPSL) realized by a thin film of silver. The next step (2a) then became determining the physical field structure with which to excite the device to implement the compensation scheme. Interestingly, we found that the result (4) from the structured illumination is equivalent to the result from a simple image post-processing (2b). Demonstrating the equivalence (3) and the improved image result (4) for the practical SNPSL under incoherent illumination is the purpose of this work.

As an alternative to the above loss-compensation methods, we have recently demonstrated a novel scheme[34] relying on the coherent superposition of lossy surface plasmon polariton (SPP) modes in a base plasmonic metamaterial[35, 36] with an auxiliary beam. This design allowed for additional power to be introduced into the system to compensate absorption in the base metamaterial without destroying the effective negative refractive index. This was termed the plasmon injection scheme but, in principle, need not be applied to a plasmonic metamaterial design; the underlying superposition concept is generalizable to an arbitrary field configuration. Because the application of interest to us is imaging, we equated the physical superposition of an auxiliary field and the object field with a simple computational deconvolution of the raw image field produced by an imperfect negative index flat lens, up to a small numerical error [37]. In other words, we showed numerically that compensating the loss in the metamaterial lens can be reduced to a simple matter of post-processing the original image with a reconstruction algorithm determined by the known transmission properties of the lens system. This was also shown to improve the resolution of a realistic cylindrical hyperlenses constructed from concentric layers of dielectric and lossy metal[38, 39]. In fact, we have shown in [39] that this technique can compensate all major forms of loss in a metamaterial imaging system, including absorption loss, diffraction or propagation loss, impedance mismatch loss, and loss due to discretization of the metamaterial.

Our previous results in [37] have been based on a hypothetical optical effective medium with refractive index n = -1 + i0.1 illuminated with coherent light, which Fourier optics theory tells us is linear in complex field amplitude[40]. In practice, however, at optical wavelengths we can only produce a plasmonic superlens operating for TM polarization such as in [3], and detecting the phase information for visible and UV

wavelengths becomes difficult. In this work, we show that, by supplying incoherent illumination, such as that from a light-emitting diode (LED), the image from a silver plasmonic superlens can be reconstructed using only real-valued intensity data. Some inspiration for this approach was found from similar work applied to nearfield microscopy in the absence of a metamaterial lens[41, 42, 43, 44]. The practical consequence is that a high-quality sub-diffraction-limited image can be formed by a compact device, as suitable commercially available light sources are packaged at the millimeter scale [45]. Also, only the image intensity data needs to be detected for a proper reconstruction, relaxing the difficulties related to phase measurement at nanometer length scales or the complexity and proper convergence of various existing phase-retrieval algorithms [46]. This paper describes the underlying theory, procedure for simulating the incoherent light and imaging system, and the results for reconstructed super-resolution near-field images from metallic double-slit objects. Also, the reconstruction results are equated with illumination from a high spatial frequency structured light source to physically compensate the components attenuated by loss in the imaging system.

4.2 Imaging Theory and Simulation

The complex degree of coherence $\gamma_{12}(\tau)$, which quantifies the mutual coherence of waves at two points in space \mathbf{r}_1 and \mathbf{r}_2 , can be calculated by

$$\gamma_{12}(\tau) = \frac{\Gamma_{12}(\tau)}{[\Gamma_{11}(0)\Gamma_{22}(0)]^{1/2}},\tag{4.1}$$

where

$$\Gamma_{12}(\tau) = \langle u(\mathbf{r}_1, t+\tau) u^*(\mathbf{r}_2, t) \rangle \tag{4.2}$$

is the cross-correlation of the complex scalar waves $u(\mathbf{r}_1, t)$ at points \mathbf{r}_1 and \mathbf{r}_2 , and $\Gamma_{11}(0)$, $\Gamma_{22}(0)$ follow similarly as the autocorrelations of the waves at each point evaluated at time lag $\tau = 0[47]$. From Schwarz's inequality, it can be shown that

$$0 \le |\gamma_{12}(\tau)| \le 1. \tag{4.3}$$

Perfectly coherent light is characterized by $|\gamma_{12}(\tau)| = 1$ and can be realized by a wave with a perfectly temporally and spatially correlated light source. In contrast, perfectly incoherent light has $|\gamma_{12}(\tau)| = 0$ and can be produced by a random light source. Consider the case where we evaluate only the temporal coherence of the source, i.e., when $\mathbf{r}_1 = \mathbf{r}_2$, and we have $\gamma_{12}(\tau) \rightarrow \gamma(\tau)$. For a source with finite temporal bandwidth, the coherence time

$$\tau_c = \int_{-\infty}^{\infty} |\gamma_{12}(\tau)|^2 d\tau \tag{4.4}$$

characterizes the time interval over which the phase is predictable. For a source with Gaussian lineshape, the integral from Eq. 4.4 gives

$$\tau_c = \sqrt{\frac{2\ln 2}{\pi}} \frac{1}{\Delta\nu} \tag{4.5}$$

where $\Delta \nu$ is the half-power bandwidth of the source, and the pre-factor comes from normalization of the power spectral density. Therefore, from Eqs. 4.1-4.5, the phase correlation of a source and, in turn, its degree of coherence are decreased when it becomes less monochromatic.

Additionally, if we separate \mathbf{r}_1 and \mathbf{r}_2 , the spatial distribution of the source also becomes important to account for when characterizing its coherence. Directly analogous with temporal coherence, we can characterize the phase correlation of the waves at two distinct points with a spatial correlation function. Consider an LED, which emits light by spontaneous emission from a p-n junction, resulting from the radiative recombination of electronhole pairs. The radiative recombination in the junction volume can be interpreted intuitively as an ensemble of random emitters, which are separated by a distance much smaller than the center emission wavelength. Consequently, the spatial correlation is small such that it is likely negligible. Therefore, we can say that the LED is totally spatially incoherent in addition to having finite τ_c .

The coherence of light waves has a well-studied impact on imaging systems. Under coherent illumination, image formation can be described by the complex amplitude equation

$$i(\mathbf{r}) = h(\mathbf{r}) \otimes o(\mathbf{r}), \tag{4.6}$$

where $i(\mathbf{r})$ is the image, $h(\mathbf{r})$ is the point spread function (PSF) of the imaging system, $o(\mathbf{r})$ is the object, \mathbf{r} is a spatial coordinate, and \otimes denotes the convolution. For the incoherent case, Eq. 4.6 becomes an intensity equation

$$|i(\mathbf{r})|^2 = |h(\mathbf{r})|^2 \otimes |o(\mathbf{r})|^2.$$

$$(4.7)$$

In the presence of additive noise, Eq. 4.7 is modified to be

$$|i(\mathbf{r})|^2 = |h(\mathbf{r})|^2 \otimes |o(\mathbf{r})|^2 + n(\mathbf{r}), \qquad (4.8)$$

where we consider $n(\mathbf{r})$ a random variable of zero-mean Gaussian probability density with standard deviation σ_n determined by the signal-to-noise ratio (SNR) of the imaging system. From convolution theorem, Fourier transforming Eq. 4.8 gives

$$I(\mathbf{k}) = H(\mathbf{k})O(\mathbf{k}) + N(\mathbf{k}), \tag{4.9}$$

where the capital letters denote the Fourier transforms of the respective intensity signals and **k** is the spatial frequency. For small **k** such that $|I(\mathbf{k})| > |N(\mathbf{k})|$, by inspection of Eq. 4.9 it becomes obvious that we can reconstruct the object by performing the operation

$$H^{-1}(\mathbf{k})I(\mathbf{k}) = H^{-1}(\mathbf{k})(H(\mathbf{k})O(\mathbf{k}) + N(\mathbf{k}))$$

$$\approx H^{-1}(\mathbf{k})H(\mathbf{k})O(\mathbf{k}) = O(\mathbf{k}).$$
(4.10)

This is the well-known inverse filter for image deconvolution. For larger **k** such that $|I(\mathbf{k})| \leq |N(\mathbf{k})|$, however, we have

$$H^{-1}(\mathbf{k})I(\mathbf{k}) \le H^{-1}(\mathbf{k})N(\mathbf{k}),\tag{4.11}$$

which dominates the contribution from Eq. 4.10 to the total inverse filtered spectrum because $H(\mathbf{k}) \to 0$ for large \mathbf{k} . To avoid this noise amplification, instead of using a naïve inverse filter, we can use our knowledge of the system SNR as a regularization parameter to linearly find an estimate $\tilde{O}(\mathbf{k})$ of $O(\mathbf{k})$. A common choice for this estimation is the Wiener deconvolution given by

$$\tilde{O}(\mathbf{k}) = \left[\frac{H^*(\mathbf{k})S_o(\mathbf{k})}{|H(\mathbf{k})|^2 S_o(\mathbf{k}) + S_n(\mathbf{k})}\right] I(\mathbf{k})$$
$$= H^{-1}(\mathbf{k}) \left[\frac{|H(\mathbf{k})|^2}{|H(\mathbf{k})|^2 + \mathrm{SNR}^{-1}(\mathbf{k})}\right] I(\mathbf{k})$$
(4.12)
$$= D_W(\mathbf{k})I(\mathbf{k}),$$

where $\text{SNR}(\mathbf{k}) = S_o(\mathbf{k})/S_n(\mathbf{k})$ is the ratio of the power spectral densities of $o(\mathbf{r})$ and $n(\mathbf{r})$. It is worth noting here that, for a good realistic detector, a SNR corresponding to a random noise intensity with Gaussian probability density having $\sigma_n = 10^{-6}$ added to a normalized image intensity spanning the interval [0,1] is a reasonable estimate[48]. We use this figure for adding noise to the simulated images in the results section. Noise also can result from roughness on the silver film surface, and the image contrast may be reduced. However, high-contrast photolithography with experimental silver superlenses has been achieved[49, 50].

To conclude this discussion of coherence effects on imaging, by illuminating an imaging system with incoherent light from an LED, the image can be deconvolved or reconstructed using only knowledge of the intensity data, in contrast with the coherent case where the complex image field must be known.

The finite-difference time-domain (FDTD) technique is a widely used numerical technique for simulating the scattering of electromagnetic waves at subwavelength scales. An advantage of a time-domain computation such as FDTD is that a short simulation in time can give the system response to a wide range of temporal frequencies after the time-series data are Fourier transformed. Because we are interested in modeling incoherent illumination, FDTD is desirable over a frequency domain technique because it requires only one simulation to obtain a broadband result. However, directly simulating realistic incoherent light by defining a random wave is not practical because the coherence time for most real physical sources is long compared with the simulated time, which for optical frequencies is often on the order of 10^{-15} s. Additionally, ensemble averaging of a distribution of random emitters would require tens or hundreds of simulations to converge to a statistically robust result to mimic spatial incoherence. These difficulties are especially apparent when simulating plasmonic structures that exhibit large field enhancements and require a high mesh resolution because the required memory and computing time for a 2D simulation scales approximately by the square and the cube, respectively, as the Yee cell size is decreased.

Fortunately, these problems can be sidestepped rather conveniently. To efficiently model the temporal and spatial incoherence, we can use mathematically the results from a FDTD simulation of a short coherent light pulse, which is easy to simulate. First, calculating the response due to the finite coherence time of a real physical light source, such as an LED, can be done by sampling the portion of the broadband FDTD data in the frequency domain, obtained by a Fourier transformation, corresponding to the LED band. The resulting spectral simulation data are then only nonzero within the nonzero portion of the LED spectrum. Second, calculating the response due to spatial incoherence can be done by simulating each slit of the double slit object we wish to individually image in Fig. 4.2. Because the object is large compared with the spatial correlation length of an LED source, the single-slit simulation results on the image plane can be added in intensity due to the linearity from Eq. 4.7. The same technique could be applied to simulate a more arbitrary slit object with different slit sizes and separations and results in a similar image quality. Increasing the slit size while keeping the separation constant can slightly increase the image contrast due to an increase in the center-to-center separation; however, for simplicity we will consider slit sizes equal to the separation. This technique has an effect similar to [44], where interference fringes from an image formed with coherent light are reduced when the spatial coherence of the source is destroyed. The equation describing the total incoherent simulation procedure used here is given by

$$\langle |H_{z}(\omega_{0})|^{2} \rangle = \int W(\omega;\omega_{0})|H_{z,1}(\omega_{0})|^{2}d\omega + \int W(\omega;\omega_{0})|H_{z,2}(\omega_{0})|^{2}d\omega$$

$$= \int W(\omega;\omega_{0})[|H_{z,1}(\omega_{0})|^{2} + |H_{z,2}(\omega_{0})|^{2}]d\omega,$$

$$(4.13)$$

where $\langle |H_z(\omega_0)|^2 \rangle$ is the total time-average intensity result for the double-slit object under TM-polarized incoherent illumination with center frequency ω_0 , $W(\omega;\omega_0)$ is the selected physical source line function with center frequency ω_0 , and $|H_{z,1}(\omega_0)|^2$ and $|H_{z,2}(\omega_0)|^2$ are the broadband intensity results from each single-slit simulation.

For this work, we employed MIT electromagnetic equation propagation (MEEP), a free, open-source FDTD software package developed at MIT[51]. The simulation geometry can be found in Fig. 4.2. A $t_{Ag} = 35$ nm thick silver superlens is placed in the computational domain bounded by perfectly matched layers (PML). A $t_{Cr} =$ 100 nm thick chromium mask is situated half the silver thickness, $t_{Ag}/2$, above the superlens interface. The slit width and separation is Δx . In the actual simulation, we



Figure 4.2: Two-dimensional FDTD simulation geometry for a double-slit chromium mask object (green) imaged by a silver superlens (gray). In our real simulation, only one slit is open, and then the results are flipped along x and added in intensity to the original to obtain the double-slit result. Perfectly matched layer (PML) boundaries are implemented to truncate the computational domain. A TM-polarized magnetic field source (purple line) is applied to the inner boundary of the upper PML layer. A nondispersive $\epsilon_{\text{background}} = 1.88$ is added to the background to match $-\text{Re}[\hat{\epsilon}_{\text{lens}}]$ at the center wavelength $\lambda = 365$ nm. The thicknesses of the Cr and Ag layers are 100 and 35 nm, respectively. The image and object planes are situated half of the silver thickness away from the lens surfaces. The Yee cell size is set to 1 nm to accurately resolve the fields near the mask interfaces.

only set one of the slits open, then flipped the image plane results along x and added the intensity result to the unflipped intensity result in accordance with Eq. 4.13 to mimic spatial incoherence. The object and image planes are defined at $\pm t_{Ag}/2$ relative to the superlens interfaces. A TM-polarized source is placed at the inner boundary of the +y PML. The source introduces a Gaussian light pulse in time with an FWHM of about 0.33 fs. The total simulation time is about 100 times the pulse



Figure 4.3: Lorentz-Drude complex permittivity $\hat{\epsilon}_r(\lambda) = \epsilon'(\lambda) + i\epsilon''(\lambda)$ as a function of wavelength λ for silver and chromium using the model from Eq. 4.14. Fit parameters were taken from [52].

width to ensure full decay of all the reflected and absorbed waves.

To account for dispersive permittivity in the silver superlens and chromium mask, the Lorentz-Drude model[52] was used, given by

$$\hat{\epsilon}_r(\omega) = 1 - \frac{f_0 \omega_p^2}{\omega(\omega - i\Gamma_0)} + \sum_{j=1}^k \frac{f_j \omega_p^2}{(\omega_j^2 - \omega^2) + i\omega\Gamma_j}$$
(4.14)

where ω_p is the plasma frequency, f_0 is the oscillator strength for intraband or freeelectron effects, Γ_0 is the intraband damping constant, ω_j is the *j*th interband or bound-electron oscillator frequency, f_j is the *j*th interband oscillator strength, and Γ_j is the *j*th interband damping constant. A plot of the fitted model from Eq. 4.14 with the parameters given in [52] is shown in Fig. 4.3.

At $\lambda = 365$ nm, the silver is excited below its plasma frequency and exhibits relatively low loss, conditions that are amenable for superlensing. Also, there exist real physical sources with this center wavelength 3. For these reasons, we select 365 nm as the center wavelength for the FDTD simulation source. The permittivity of silver at $\lambda = 365$ is calculated from Eq. 4.14 to be $\hat{\epsilon} = -1.88 + i0.51$. Because ideal superlensing occurs when $\operatorname{Re}[\hat{\epsilon}_{\text{background}}] = -\operatorname{Re}[\hat{\epsilon}_{\text{lens}}]$, a background medium with nondispersive $\epsilon_{\text{background}} = 1.88$ was applied to the entire simulation domain. Changing this parameter to a more practical value should not detrimentally affect the imaging results; rather, it is likely that typical dielectrics to be integrated with a real fabricated silver superlens would have larger permittivity. This would increase the numerical aperture and make the effective wavelength shorter. We performed a PSF calculation with $\epsilon_{\text{background}} = 2.4$ to verify the performance with a more realistic dielectric^[3, 49, 50] and determined that the FWHM was only 1-2 nm larger than the PSF with $\epsilon_{\text{background}} = 1.88$. The results for this simulation procedure and the image reconstructions using the deconvolution in Eq. 4.12 are given in the following section.

4.3 Results

After performing the simulations, the time-series data extracted from the image plane is Fourier transformed in MATLAB, and the magnitude is squared to give $|H_{z,1}(\omega)|^2 +$ $|H_{z,2}(\omega)|^2$ from Eq. 4.13. To represent a real physical LED source such as given by [45], $W(\omega; \omega_0)$ from Eq. 4.13 is chosen as a Gaussian line function with a center wavelength of 365 nm and a bandwidth of 9 nm. The FDTD source spectrum, which excited the simulation geometry, is given in Fig. 4.4(a), along with $W(\omega; \omega_0)$. The image plane result $|H_{z,1}(\omega)|^2 + |H_{z,2}(\omega)|^2$ for a $\Delta x = 60$ nm double-slit object is given in Fig. 4.4(b). This is then multiplied with $W(\omega; \omega_0)$ for all x to give the result in Fig. 4.4(c).

After integrating the result from Fig. 4.4(c) according to Eq. 4.13, the time-average image plane intensity from the incoherent illumination is obtained. We repeated this process for two other double-slit objects with $\Delta x = 30$ nm and $\Delta x = 20$ nm. Using the incoherent PSF measured from a single 20 nm slit simulation using the same procedure as above, the images were then reconstructed using the linear Wiener deconvolution from Eq. 4.12. The results are shown in Fig. 4.5, plotted with the intensity on the object and image planes as indicated in Fig. 4.2. The spikes in the object plane data are a result of field enhancements at the edges of the Cr mask and are clearly not transferred to the image plane. It can be seen in Fig. 4.5 that the reconstructions easily achieve super-resolution of the double-slit objects. Particularly in Fig. 4.5(c), the distinct slits, which were previously unresolved, are now visible in the deconvolved image, with a resolution better than $\lambda/18$.

Further inspection of Fig. 4.5 reveals that the reconstructed images have smaller FWHM than the slits themselves, which can be explained by the properties of the



Figure 4.4: Frequency domain results for the FDTD simulation. (a) FDTD simulation (black line) and real LED source $W(\omega; \omega_0)$ (purple line) spectra. The center wavelength for the simulation source is set to 365 nm, but a blueshift resulted from some numerical error introduced by turning on the source. The final calculated image is not adversely affected, however, because the portion of the spectrum sampled by $W(\omega; \omega_0)$ does not vary significantly. (b) Frequency domain image plane data from the simulation in Fig. 4.2 of a $\Delta x = 60$ double-slit object. This is the bracketed quantity of the integrand in Eq. 4.13. (c) Result from (b) multiplied by the physical source band $W(\omega; \omega_0)$ in (a). This is the integrand of Eq. 4.13, which, when integrated over the frequencies, gives the time-average intensity on the image plane.

PSF, $|h(x)|^2$. We determined the PSF by illuminating a single 20 nm slit and measuring the image plane result. The FWHM of this measurement is comparable with the FWHM of the image from a single 60 nm slit; thus, it is reasonable to say that slits below 60 nm in size are described approximately by a delta function. Mathematically,



Figure 4.5: Linear deconvolution results for image data from (a) $\Delta x = 60$ nm, (b) $\Delta x = 30$ nm, and (c) $\Delta x = 20$ nm double-slits with noise added. The object and image plane intensities are also plotted and normalized for comparison. The reconstructions have negative values due to the loss of information to random noise in the imaging process. Therefore, a physical result cannot be ensured by applying the filter $D_W(k_x)$ because the time-reversal symmetry is broken.

it can be written as

$$|i_{\Delta x}(x)|^{2} = |h(x)|^{2} \otimes |o_{\Delta x}(x)|^{2}$$

$$\approx |h(x - \Delta x)|^{2} + |h(x + \Delta x)|^{2}$$

$$= |h(x)|^{2} \otimes [\delta(x - \Delta x) + \delta(x + \Delta x)],$$
(4.15)

where $|i_{\Delta x}(x)|^2$ is the image resulting from double-slit object $|o_{\Delta x}(x)|^2$ with characteristic dimensions Δx as defined in Fig. 4.2. From the first and third lines of Eq. 4.15, we then have

$$|o_{\Delta x}(x)|^2 \approx \delta(x - \Delta x) + \delta(x + \Delta x).$$
(4.16)

In other words, having a smaller FWHM in the reconstruction is allowed because small slits, such as the ones simulated here, produce the same images as delta functions, so the reconstruction FWHM is only limited by the noise level in the system and not the slit dimensions.

Another curiosity of Fig. 4.5 is the appearance of negative values in the reconstruction. Obviously, this result is nonphysical because we cannot measure a negative number of photons. The origin of this is solely the fact that we apply no non-negativity constraints in the linear Wiener deconvolution $D_W(k_x)$ from Eq. 4.12. Because the forward process in time of imaging the object $O(k_x)$ with the optical transfer function $H(k_x)$ includes a loss of information due to added random Gaussian noise, the timereversal symmetry is broken. Therefore, trying to reverse the process in time with the linear filter $D_W(k_x)$ allows a nonphysical result. We can, however, apply a nonlinear iterative reconstruction that rejects negative values at each iteration. A common such approach is given by the Richardson-Lucy (RL) algorithm[53, 54]. Employing the RL algorithm to our double-slit images using the built-in MATLAB function (deconvlucy) gives the results in Fig. 4.6. After 20 iterations, the algorithm offers



Figure 4.6: Nonlinear Richardson-Lucy deconvolution results for noisy data from (a) $\Delta x = 60$ nm, (b) $\Delta x = 30$ nm, and (c) $\Delta x = 20$ nm double-slits. To obtain these results, the algorithm was run for 20 iterations.

satisfactory reconstructions of the images, which exclude negative values. As in Fig. 4.5, the deconvolved images in Fig. 4.6 show resolution well beyond the diffraction limit, again better than $\lambda/18$ in Fig. 4.6(c) for a previously unresolved object.

Aside from reconstructing superlens images with intensity data, another objective of this work is to show the equivalence between the linear post-processing technique in Fig. 4.5 with loss compensation realized by structured light illumination. An obvious way to do so analytically is to show the total structured object $O_{\text{total}}(\mathbf{k})$, which, after passing through the system, gives rise to a total image equal to the reconstruction result $\tilde{O}(\mathbf{k})$. This can be determined mathematically by

$$O_{\text{total}}(\mathbf{k}) = D_W(\mathbf{k})\tilde{O}(\mathbf{k}), \qquad (4.17)$$

as propagating through the lens system gives

$$H(\mathbf{k})O_{\text{total}}(\mathbf{k}) = I_{\text{total}}(\mathbf{k})$$
$$= H(\mathbf{k})D_W(\mathbf{k})\tilde{O}(\mathbf{k}) \qquad (4.18)$$
$$\approx \tilde{O}(\mathbf{k}).$$

We call them the total object and images to draw comparison with our previous work[37, 38, 39]. In those cases, an auxiliary coherent structured field is superimposed with the original object field to give a total object, which is the imaging equivalent to the plasmon injection scheme from [34], where an auxiliary beam supplies additional energy to the metamaterial to combat absorption. The total object intensity distribution, which satisfies Eq. 4.17 for the result in Fig. 4.5(a), is shown in Fig. 4.7.

As previously in Fig. 4.5, because we applied a linear filter with no non-negativity constraints, the result of Eq. 4.17 before any further manipulation exhibits negative


Figure 4.7: Total object distribution $o_{\text{total}}(x)$ calculated from Eq. 4.17 for $\Delta x = 60$ nm. A DC offset is applied to make all the intensity levels positive. (b) Image $i_{\text{total}}(x)$ calculated from Eq. 4.18 resulting from the object distribution in (a) with the DC offset removed for comparison with the deconvolved image $\tilde{o}(x)$.

values in the spatial domain, which are obviously nonphysical. However, a DC offset simply can be added to the total object to make it non-negative, as in Fig. 4.7(a), and the resulting image would just be offset accordingly. This procedure is not directly useful in a practical imaging scenario since it requires knowledge of the reconstruction $\tilde{O}(\mathbf{k})$ but could be useful for photolithography, where the desired intensity on the image plane is known. For this work, however, it is important to emphasize that Fig. 4.7 shows the equivalence of image reconstruction by post-processing and the physical injection of a high-spatial frequency object.

4.4 Discussion and Conclusion

We have recently shown that a similar active reconstruction scheme for metamaterial lenses can recover even noise-obscured spatial frequency components of an image by injecting an appropriate high-spatial frequency source covering a finite band convolved with the object field of interest using coherent light[48]. Practical implementation is then only dependent on the development of a near-field spatial filter or metasurface that can physically produce the structured light convolution. This procedure could then be extended to incoherent light simply by appropriately configuring the physical filter, in accordance with the intensity linearity from Eq. 4.8. Such an active implementation of our reconstruction method could dramatically improve the resolution of metamaterial lenses even in the presence of high absorption loss. Also, it would be much more convenient to implement with incoherent light because only the intensity information is required. We plan to perform this implementation in future work.

In summary, we showed the simulation of incoherent imaging of subwavelength double-slit objects with a plasmonic superlens using the FDTD method. The resulting images were reconstructed using the linear Wiener deconvolution and the nonlinear Richardson-Lucy deconvolution. Both algorithms gave reconstructions with improved contrast and decreased FWHM. Resolution better than $\lambda/18$ was achieved with both algorithms. The linear deconvolution results were unified with the physical injection of structured light, corresponding to the compensation of the losses in the superlens imaging system. The results of this work show that our previous loss compensation methods can be carried out using only intensity data if incoherent light is provided. The practical implication is that a high-quality super-resolution image can be formed by a thin film plasmonic lens illuminated by a compact, inexpensive light source such as a light-emitting diode and subsequently reconstructed with only intensity data.

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Chapter 5

Plasmonic Superlens Imaging Enhanced by Incoherent Active Convolved Illumination¹

5.1 Introduction

The theory and experimental demonstration of metamaterials has inspired interesting avenues of imaging, [1, 2, 3, 4, 5, 6, 7, 8, 9, 10, 11, 12] lithography, [13, 14] and beam generation [15] beyond the diffraction limit, particularly motivated by the prospect of

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a perfect lens[16]. Researchers have quickly realized superlenses[1, 2, 3, 4, 5, 6, 13, 14], hyperlenses[7, 8, 9, 10], integrated metalenses,[11, 15] and non-resonant elliptical lenses,[12] which can either amplify or propagate evanescent waves carrying the precious high spatial frequency information of an object. However, these efforts have stopped short of creating a truly perfect lens, since their resolution is limited by losses and they cannot focus light of arbitrary polarization especially at optical wavelengths. Metasurfaces[17, 18, 19, 20, 21, 22, 23, 24, 25, 26, 27, 28] with unprecedented wavefront manipulation capabilities have also emerged to overcome challenging fabrication and attenuation issues[29, 30, 31, 32, 33, 34] relevant to three-dimensional metalenses and other metadevices.

We have previously shown a technique for compensating the loss in a plasmonic metamaterial[35] by injecting additional surface plasmon polaritons into the metamaterial via a coupled external beam[36]. This allows the introduction of additional energy into the system to compensate for absorption loss in the metallic structures without detrimentally altering the effective negative refractive index. In contrast to traditional loss compensation methods which often implement gain media[37, 38, 39, 40, 41, 42, 43, 44, 45, 46], the so-called "plasmon injection" (Π) scheme does not suffer from the instability introduced by gain or any issues of causality[47]. The desired effective parameters of the metamaterial can then be preserved, and the many practical problems of implementing gain media can be avoided, including the selection of suitable materials and pump sources for specific wavelengths and their limited lifetimes. We subsequently translated the Π scheme to superresolution imaging with a homogeneous negative index flat lens (NIFL) with nonzero loss[48], a hyperlens[49, 50], and a silver superlens[51]. Our findings showed that linear deconvolution of the image produced by a lossy metamaterial lens is equivalent to physically compensating the loss with injection of an additional structured source, as originally demonstrated with the Π [36]. However, passive post-processing can only recover spatial frequency components which are not lost to random noise in the imaging process.

To push the performance of our compensation method further, we developed an active version which relies on the coherent convolution of a high spatial frequency function with an object field focused by a lossy NIFL[52]. Selective amplification of a small band of spatial frequencies by spatially filtering the object under a strong illumination beam can favorably alter the transfer function so that spatial frequency components which would originally be lost to noise can be successfully transferred to the image plane. In this article, we present a more advanced and versatile method that alternatively employs a simple plasmonic superlens structure illuminated by incoherent UV light, avoiding the complexity and practical difficulties related to phase retrieval or phase detection of coherent fields. Numerical imaging results with this method can resolve point dipole objects separated by a few tens of nanometers, an improvement over passive post-processing. Additionally, we identify that signal-dependent noise provides a limitation to this method and discuss potential strategies for experimental implementation.

5.2 Theoretical Description and Noise Characterization

Incoherent linear shift invariant imaging systems can be conveniently described by the intensity convolution relation

$$i(\mathbf{r}) = h(\mathbf{r}) * o(\mathbf{r}),\tag{5.1}$$

where $i(\mathbf{r})$ is an observed intensity image, $h(\mathbf{r})$ is the incoherent point spread function (PSF) of the system, $o(\mathbf{r})$ is the object intensity, \mathbf{r} is a position coordinate, and *denotes the convolution operation. Here we will treat the intensities as normalized quantities. Since convolution in position space is equivalent to multiplication in frequency space, Fourier transformation of eq 5.1 gives the resulting spatial frequency content on the image plane as

$$I(\mathbf{k}) = H(\mathbf{k})O(\mathbf{k}),\tag{5.2}$$

where \mathbf{k} is the spatial frequency and the capital letters denote the respective Fourier transforms. Observing the theoretical and experimental transmission properties of

lossy near-field superlenses shows that $H(\mathbf{k})$ has a low-pass filtering effect on the image[16, 53, 54, 55, 56]. Consequently, many of the high spatial frequency components of the object are not transferred to the image plane. This is problematic for the imaging of nanometric objects, since the absence of the high- \mathbf{k} information reaching the detection plane can result in an indiscernible blurry image.

To combat the attenuation of this information, we propose a method to recover it by adding additional energy to the system, which we call active convolved illumination (ACI). Consider an "active" object $o_{ACI}(\mathbf{r})$ obtained by the convolution

$$o_{ACI}(\mathbf{r}) = o(\mathbf{r}) * a(\mathbf{r}) + a_0, \tag{5.3}$$

where $o(\mathbf{r})$ is the object distribution we wish to obtain, $a(\mathbf{r})$ is a function that passes high spatial frequencies which we use to inject extra energy to compensate the decaying transmission, and a_0 is a "DC offset" to ensure that $\forall \mathbf{r} : o_{ACI}(\mathbf{r}) \geq 0$. The choice of a_0 will be evidently dependent on the term $o(\mathbf{r}) * a(\mathbf{r})$, and we select it as a constant for simplicity. However, in a real imaging system, the non-negativity of $o_{ACI}(\mathbf{r})$ will automatically enforced by the physical propagation, meaning that no knowledge of the object $o(\mathbf{r})$ would be required. The spatial frequency content is then

$$O_{ACI}(\mathbf{k}) = O(\mathbf{k})A(\mathbf{k}) + a_0\delta(\mathbf{k}).$$
(5.4)

The relation in eq 5.4 is not very informative until we specify the mathematical form of $A(\mathbf{k})$. Therefore, let us define

$$A(\mathbf{k}) = 1 + P(\mathbf{k}) \tag{5.5}$$

where $P(\mathbf{k})$ is of a form convenient in terms of mathematical simplicity and practical considerations, such as a Gaussian function. To perform the ACI compensation, we can then define a series of Gaussian $P(\mathbf{k})$ and subsequently convolve $A(\mathbf{k})$ with $O(\mathbf{k})$. Explicitly, $P(\mathbf{k})$ can be written as

$$P(\mathbf{k}) = \sum_{j} P_{j} \exp\left[-\frac{(\mathbf{k} - \mathbf{k}_{j})^{2}}{2\sigma_{j}^{2}}\right],$$
(5.6)

where j is the Gaussian number, σ_j is a parameter proportional to the spatial frequency bandwidth of the *j*th Gaussian, and P_j and \mathbf{k}_j are the amplitudes and center spatial frequencies for the *j*th Gaussian, respectively. Propagation of the ACI object through the superlens system to the image plane gives

$$H(\mathbf{k}) [O(\mathbf{k})A(\mathbf{k}) + a_0\delta(\mathbf{k})]$$

= $H(\mathbf{k})O(\mathbf{k}) + H(\mathbf{k}) [O(\mathbf{k})P(\mathbf{k}) + a_0\delta(\mathbf{k})]$ (5.7)
= $I(\mathbf{k}) + I_{ACI}(\mathbf{k}).$

The term $H(\mathbf{k}) [O(\mathbf{k})P(\mathbf{k}) + a_0\delta(\mathbf{k})]$ represents the ACI contribution to the image plane spatial frequency content, $I_{ACI}(\mathbf{k})$. In order to collect deterministic information on the image plane for a selected \mathbf{k} within the bandwidth of $P(\mathbf{k})$, we then must satisfy the inequality

$$|H(\mathbf{k})O(\mathbf{k})A(\mathbf{k})| > |N(\mathbf{k})|, \tag{5.8}$$

where $|N(\mathbf{k})|$ is the noise level. In summary, the reason this method can provide more spatial frequency content than simple passive propagation is that we can control the "effective" transfer function so that high-**k** components of $O(\mathbf{k})$ can reach the image plane without being fully attenuated below the noise level, provided that P_j is sufficiently large.

To evaluate the prospects of our ACI method in a practical imaging scenario, the effects of noise must be taken into account. Particularly in the case of intensity measurements, since the optical power applied with the active convolution will become large, understanding the signal-dependent nature of the noise is crucial, since the resulting image will possess a substantial mean pixel value. In turn, the signal-dependent noise level will be inherently increased compared with passive imaging. To obtain a noisy image $i_n(\mathbf{r})$ we assume a parametric noise addition of the form

$$i_n(\mathbf{r}) = i(\mathbf{r}) + i(\mathbf{r})^{\gamma} u(\mathbf{r}) + v(\mathbf{r}), \qquad (5.9)$$

where $i(\mathbf{r})$ is the noiseless image, $u(\mathbf{r})$ and $v(\mathbf{r})$ are independent zero-mean Gaussian random variables, and γ is a parameter satisfying $|\gamma| \leq 1$. In this case, we take each term of eq 5.9 to represent the corresponding photon counts read out by the detector. Additionally, in all following calculations, the photon counts and their corresponding standard deviations are normalized to the maximum count value of the original object distribution that we wish to image. We maintain the same notation for $i(\mathbf{r})$ as eq 5.1 since we treat both the intensities and photon counts are normalized. Due to the independence of $u(\mathbf{r})$ and $v(\mathbf{r})$, we can then write the standard deviation of $i_n(\mathbf{r})$ as

$$\sigma_n = \sqrt{i(\mathbf{r})^{2\gamma} \sigma_u^2 + \sigma_v^2},\tag{5.10}$$

where σ_u^2 and σ_v^2 are the variances of $u(\mathbf{r})$ and $v(\mathbf{r})$, respectively. For our ACI method, $i(\mathbf{r})$ will become large such that $v(\mathbf{r})$ is negligible, assuming σ_u and σ_v have similar orders of magnitude. Therefore, we can simplify eq 5.10 to be

$$\sigma_n = \sigma_u i(\mathbf{r})^{\gamma}. \tag{5.11}$$

Let us then define a signal-to-noise ratio (SNR),

$$SNR = \frac{i(\mathbf{r})}{\sigma_n} = \frac{i(\mathbf{r})}{\sigma_u i(\mathbf{r})^{\gamma}} = \frac{i(\mathbf{r})^{1-\gamma}}{\sigma_u}.$$
 (5.12)

Optical detectors can reach a SNR of around 60 dB[52, 57]. If we select $\gamma = 1$ in eq 5.12, σ_u becomes 10⁻⁶ for a 60 dB SNR. However, for most realistic detectors $\gamma = 0.5$ due to the Poisson distribution of photon noise [58]. In this case, eq 5.12 becomes

$$SNR = \frac{\sqrt{i(\mathbf{r})}}{\sigma_u},\tag{5.13}$$

and the SNR is evidently dependent on the signal.

5.3 Imaging Simulation

To model a realistic silver superlens structure, we consider a geometry similar to an already experimentally realized superlens which transfers image intensity data onto a photoresist (PR) layer[1]. In order to obtain the PSF for this superlens structure, we simulated the point dipole response with the commercial finite-difference time-domain solver Lumerical FDTD Solutions. The simulation geometry and calculated PSF can be found in Figure 5.1. In principle, the image plane for a flat silver superlens should lie at the z-position where the phase is matched to the object plane. However, there is no such well-defined image plane in Figure 5.1 since the object lies further than a lens thickness from the lens interface on the object side, and the developed PR on the imaging side will have a topographical distribution with varying z. Therefore, we somewhat arbitrarily selected an image plane 40 nm from the lens interface in the PR layer for our calculations. This actually highlights that even a defocusing of the image can be overcome with our ACI method. The dipole source is defined as

a y-oriented magnetic dipole with center wavelength $\lambda_0 = 365$ nm and bandwidth $\Delta \lambda = 9$ nm to mimic the spectrum of a commercially available UV light-emittingdiode (LED) [59]. Dispersion in the silver is modeled with a fit to experimental data [60]. The time-average intensity signal in the simulation is finally obtained by an average of the squared modulus of the y-component of the magnetic field, $|\mathbf{H}_y(x,y)|^2$, over the bandwidth of the LED source[51]. In Figure 5.1, the PSF is asymmetric and narrower along x. This is expected, since the Ag superlens has negative permittivity at $\lambda_0 = 365$ nm and can only effectively focus the magnetic field in the direction perpendicular to the dipole orientation. Therefore, illumination with an unpolarized source may slightly worsen the achievable spatial resolution.

To obtain the image resulting from a spatially-incoherent distributed object, we can define a distribution of dipole sources on the object plane. Simulating the spatiallyincoherent object is then easily performed by adding the contributions of each dipole to the image plane time-average intensity separately. For example, if we consider an object consisting of n dipoles each with intensity a_j and located at position \mathbf{r}_j , to obtain the resulting image plane distribution $i(\mathbf{r})$ we perform the summation

$$i(\mathbf{r}) = \sum_{j=1}^{n} \left[h(\mathbf{r}) * a_j \delta(\mathbf{r} - \mathbf{r}_j) \right], \qquad (5.14)$$

following from the imaging theory in eq 5.1. An example simulation is shown in Figure 5.2 with n = 4, $a_1 = a_2 = a_3 = a_4$, $\mathbf{r}_1 = (0,0)$ nm, $\mathbf{r}_2 = (25,0)$ nm, $\mathbf{r}_3 = (-100,0)$



Figure 5.1: Three-dimensional FDTD superlens simulation. A single magnetic dipole embedded in a polymethyl methacrylate (PMMA) dielectric layer is oriented along y and situated 40 nm above the superlens. The image plane (dashed line) in the photoresist (PR) layer is chosen to lie 40 nm below the superlens. The lower plot shows the resulting image plane intensity distribution, which is the PSF of the superlens structure for y-polarization of the magnetic field.

nm, and $\mathbf{r}_4 = (-25, 50)$ nm. It can be seen that the two sources near the origin, which are separated by a distance of 25 nm, are unresolved even after deconvolution with the iterative Richardson-Lucy algorithm [61, 62]. The low-pass filtering due to $H(\mathbf{k})$ therefore cuts off some of the spatial frequencies that are required for reconstruction of the object. Using our active convolution method, we can recover the lost spatial frequencies that are beyond the cutoff of the passive imaging system.



Figure 5.2: Superlens incoherent imaging simulation example of four arbitrarily positioned magnetic point dipole sources. (a) The source distribution at the object plane. (b) Intensity distribution on the image plane. (c) Passive deconvolution of the image in (b) using the Richardson-Lucy algorithm. The two sources near the origin separated by 25 nm are clearly unresolved after deconvolution.

5.4 Results and Discussion

The imaging simulation from Figure 5.2 (a) and (b) was used as an example to implement the ACI method. The parameters of $P(\mathbf{k})$ from eq 5.6 were chosen to be $P_1 = P_2 = 10^4$, $|\mathbf{k}_1| = 7nk_0$, $|\mathbf{k}_2| = 8nk_0$, and $\sigma_1 = \sigma_2 = 1.5nk_0/2\sqrt{2\log 2}$, where n = 1.6099 is the refractive index of the PR imaging medium and $k_0 = 2\pi/\lambda_0$ is the free space wave number. The resulting ACI object was then propagated through the system using the transfer function $H(k_x, k_y)$ calculated by Fourier transformation of the simulated PSF h(x, y). The resulting spatial frequency content is shown in Figure 5.3. In Figure 5.3 (c), it can be seen that a larger band of spatial frequencies are recovered on the image plane compared to (b) due to propagation of the ACI object. After reconstruction with the Richardson-Lucy algorithm and the "active"



Figure 5.3: Active convolved illumination in the spatial frequency domain. (a) Fast Fourier Transform (FFT) magnitude of the object distribution in Figure 5.2 (a). (b) FFT magnitude of the image distribution in Figure 5.2 (b). (c) FFT magnitude of the image using the ACI method. (d) FFT magnitude of the image from (c) after deconvolution with the Richardson-Lucy algorithm. All plots are on a logarithmic scale.

 \mathbf{PSF}

$$h_{ACI}(x,y) = h(x,y) * a(x,y),$$
(5.15)

the object spectrum is mostly recovered in Figure 5.3 (d). Note that the DC component introduced by the ACI procedure is excluded in eq 5.15 since it only contributes to the $\mathbf{k} = 0$ component and in turn has no effect on the the reconstruction. The calculation in Figure 5.3 was performed with $\sigma_u = 0$ and $\sigma_v = 10^{-6}$ to better explicate the impact of the ACI on the imaging process. However, it is crucial to evaluate the ACI imaging performance in the presence of signal-dependent noise. To do so, we added simulated noise with $\gamma = 1$, corresponding to a constant SNR, and $\gamma = 0.5$ for a signal-dependent SNR. We have chosen these values of γ to represent both the expected Poissonian counting statistics ($\gamma = 0.5$) and a "worst case" scenario ($\gamma = 1$) to both compare the effects of signal-dependent and signal-independent SNR as well as evaluate the robustness of our method to a variety of noise conditions. Figure 5.4



Figure 5.4: Active convolved illumination imaging in the presence of signaldependent noise. (a)-(c) The calculated images after applying the ACI method for $\gamma = 1$ and different values of σ_u from eq 5.11. These are the "measured" images that would be detected in an experiment. (d)-(f) The final reconstructed images after deconvolution with the Richardson-Lucy algorithm and the active PSF defined in eq 5.15. (a) and (d) are the ideal results with $\sigma_u = 0$. Qualitatively, it can be seen that the image in (c) is corrupted by noise, and the reconstructed image in (f) consequently suffers. However, the noise in (b), corresponding to a 60 dB SNR, is small enough to achieve a good reconstructed image in (e).

shows the imaging results for $\gamma = 1$ with varied σ_u . Figure 5.4 (a) and (d) respresent the ideal ACI image and corresponding reconstruction with $\sigma_u = 0$. The ACI image in (b) and the corresponding successful reconstruction in (e) consider a 60 dB SNR attainable with modern photodetectors. Unfortunately, decreasing the SNR to 50 dB in (c) leads to an image almost fully corrupted by noise that cannot be reconstructed in (f). In contrast, Figure 5.5 considers the case of Poisson-distributed noise with $\gamma = 0.5$ and varied σ_u . As shown in eq 5.13, when $\gamma = 0.5$ the SNR becomes dependent on the signal level. Therefore, the σ_u value we can define as a "realistic" noise in



Figure 5.5: Active convolved illumination with $\gamma = 0.5$ and three different values of σ_u . (a)-(c) The noisy ACI images and (d)-(f) the corresponding reconstructions. (d) and (e) successfully resolve the object, however the image in (c) is too noisy to obtain a good reconstruction in (f).

Figure 5.5 is not explicit. However, if we inspect eq 5.13 and take $i(\mathbf{r}) \approx \overline{i(\mathbf{r})}$, where $\overline{i(\mathbf{r})}$ is the mean pixel value, we can solve for the σ_u corresponding approximately to a 60 dB SNR. We can make this approximation since the variations in $i(\mathbf{r})$ are four orders of magnitude smaller than $\overline{i(\mathbf{r})}$ for this specific imaging example. The ACI imaging results for these parameters are shown in Figure 5.6, and it can be clearly seen in (b) that the object can again be successfully resolved.

There are a few aspects of the ACI method that require some qualitative discussion, in particular the limitations of incoherent ACI for increasing the resolution of an imaging system. So-called "perfect" imaging exhibiting a flat effective transfer function could in principle be approached with this method by iteratively applying multiple $P(\mathbf{k})$





Figure 5.6: Active convolved illumination imaging with $\gamma = 0.5$ and SNR \approx 60 dB. (a) The noisy ACI image and (b) the corresponding reconstruction. In (b) the object is successfully resolved, despite the addition of realistic noise in (a).

passing distinct spatial frequency bands so that the full spectrum of the object can be recovered from the noise. However, since the main objective of ACI is to add energy to a narrow band of the object's spatial spectrum in order to overcome the attenuation of those spatial frequencies by the imaging system, it becomes evident that shifting this band to higher \mathbf{k} will require larger intensities. Therefore, we are met with a trade-off between increasing the detectable spatial frequencies and reducing the noise to an acceptable level. We have identified this trade-off as the theoretical limit of the spatial resolution achievable with the ACI method. Using the simulations and parameters described above, we found that the best resolution we could obtain was about 20 nm while considering a 60 dB SNR. However, this could be improved by better optimizing the phase and impedance match between the object and image planes by appropriately tuning the refractive indices of the dielectrics surrounding the superlens along with the locations of each plane. As an example, in our simulation from Figure 5.1, it is reasonable to expect a better resolution when the dipole is moved slightly closer to the Ag layer, since more of the evanescent components from the source will reach the superlens. We expect these steps would better optimize the results in terms of spatial resolution. Our ACI method could be just as effectively applied to any of these different geometries.

The most pressing obstacle for implementation of the ACI method into an experimental system is creating the physical convolution in eq 5.4. One way to do this would be to illuminate the object with a high-intensity beam and then spatially filter the near-field intensity distribution. We have shown the design of a near-field spatial filter[63] based on hyperbolic dispersion for a similar function to approximate the behavior of $P(\mathbf{k})$ in eq 5.6 and used the filter under "coherent" convolved illumination for enhanced superlens imaging[64]. We can use a similar configuration for realization of ACI even when we do not have access to the phases of the fields and the light is not strictly perfectly coherent. Suppose we want to image the intensity pattern formed by a periodic Chromium (Cr) grating object illuminated with TM-polarized light as shown in Figure 5.7. Here we consider the same LED light with $\lambda_0 = 365$ nm and $\Delta \lambda = 9$ nm as in the previous simulations. To realize the convolution, we place below

the grating a hyperbolic spatial filter we have designed which passes a small band of the spatial frequencies near $6k_0$ which are present in the selected grating. The filter is formed by alternating layers of Aluminum (Al) and Titanium dioxide (TiO_2) with thicknesses of 16 nm and 15 nm, respectively. Each pair of metal-dielectic layers constitutes a unit cell of the hyperbolic metamaterial (HMM). In this case, we only need to use 4 unit cells to construct the filter since the low spatial frequencies which would otherwise tunnel through the filter are automatically rejected by the grating. This also has the positive side effect of better transmission compared to the corresponding filter we previously designed [64]. The red dashed line in Figure 5.7 (a) at the exit interface of the HMM is the plane at which the active convolution exists. To the the physical system in with our theory, the HMM spatial filter essentially performs the operation in eq. 5.4, and the amplitude of $P(\mathbf{k})$ can be controlled by simply modulating the intensity of the illumination incident on the grating object. The ACI image is then formed at the image plane (white dashed line) within the PR layer. Using FDTD solutions, we performed simulations of this geometry, and also the passive configuration in Figure 5.7 (b), in order to provide a physical proof-of-concept for our ACI method and its potential for enhancing the resolution. In this case we increased the illumination intensity by selecting $P_1 = 10^8$ to fully overcome the added 60 dB signal-independent noise, but decreasing this value by about two orders can still give good results. A 10-20 dB signal-dependent SNR was also found to be tolerable using these parameters. The simulation results can be found in Figure 5.8. The intensity

profile induced on the Cr grating mask shows "hot spots" that occur every 60 nm at the sharp edges of the grating (see black solid line in Figure 5.8 (a)). The spatial frequency corresponding to a 60 nm period is within the passband of the HMM spatial filter, and in Figure 5.8 (a) this frequency is accentuated in the ACI image (blue solid line) compared to the passive image (turquoise solid line). The magnitudes of this frequency for the data in (a) are shown in (b) for comparison. Finally, the deconvolution of the ACI image with the PSF (calculated by removing the grating and placing a point source on the object plane) gives a better representation of the intensity induced on the grating than the passive deconvolution (i.e., compare the red solid line with the purple). This result can in principle be improved by tuning the spatial filter to higher spatial frequencies, provided that it passes one of the primary grating frequencies.

It is interesting to note that such spatial filters integrated with a superlens cavity and type I hyperbolic metamaterial have also been recently proposed for nanofocusing of Bessel beams[15] and an implementation of a hyperbolic dark-field lens[11], respectively. Additionally, spatial filtering has been shown to reduce the line edge roughness of photolithographic exposures in the presence of surface roughness[65]. Therefore, one natural extension of our work would be the studying of these intriguing highresolution imaging systems from the ACI method perspective to not only experimentally confirm our theoretical predictions but also improve their performances. The actual transmission properties of the filter may not exactly replicate $P(\mathbf{k})$ from eq 5.6,



Figure 5.7: (a) Simulation geometry for the physical realization of ACI superlens imaging with a hyperbolic metamaterial spatial filter. The black, red, and white dashed lines indicate the object, active convolution, and image planes, respectively. The image plane is set 5 nm below the Ag superlens in order to make the total propagation distance in the PMMA and PR equal to the lens thickness. (b) The passive simulation geometry used for comparison with the results from (a). The PMMA layer separating the Cr mask and the Ag superlens is set to 30 nm in order to match the total propagation distance in the PMMA in (a).

but the important property is the ability to selectively amplify a band of high spatial frequencies relative to the spatial frequencies in the original passband of the superlens. We showed in Ghoshroy *et al.*[52] that simply amplifying the entire object spectrum (the "strong illumination" case) will only result in a deleterious amplification of the



Figure 5.8: (a) Electric field intensities at the object and image planes for the ACI and superlens imaging configurations from Figure 5.7 considering 60 dB SNR, along with the reconstructions calculated with the Richardson-Lucy deconvolution algorithm. Deconvolution of the ACI image (red solid line) better matches the sharp peaks in the object plane (see black solid lines) as compared to deconvolution of the image formed by Figure 5.7 (b) (see purple solid line). (b) FFT magnitudes of the data in (a) within the passband of the HMM spatial filter. A clear enhancement of the spatial frequency content near $k_x/k_0 = 6.1$ can be seen for the ACI image (blue solid line) as compared to the image formed by the silver superlens alone (turquoise solid line). For the ACI image, the ratio of this spectral component to the DC component is increased by more than 18% over the unfiltered superlens image.

signal-dependent noise in the image. This is why the selective amplification of a finite portion of the spatial spectrum is required. A secondary obstacle for experimental implementation is the near-field detection of the subwavelength intensity distributions produced by this method. Near-field images produced by silver superlenses operating in the UV are often read out by exposing a negative tone PR layer on the imaging side of the lens, developing the PR, then characterizing the developed PR topography with an atomic force microscope. The ACI imaging system we have shown in Figure 5.7 (a) is within the capabilities of modern nanofabrication. However, it is likely that scaling the experiment to infrared, terahertz, or microwave frequencies would be more convenient in terms of both fabrication and detection. There is no theoretical restriction to scaling our ACI method to other frequencies. The images could then be directly read out with a subwavelength near-field probe[2, 6, 66, 67, 68, 69, 70, 71, 72, 73, 74] or detector[75, 76, 77, 78, 79, 80, 81, 82, 83, 84] small enough to resolve the important features.

5.5 Conclusion

We developed a loss compensation method to improve the resolution of a near-field silver superlens using incoherent active convolved illumination. A theoretical description of the imaging method for incoherent light is developed and implemented in numerical simulations using a combination of the finite-difference time-domain method and linear shift-invariant imaging theory. The presence of signal-dependent noise is taken into account to represent a realistic imaging scenario. The imaging method presented can achieve a resolution of around $\lambda_0/15$ or better under optimal phase and impedance matching conditions even when corrupted by realistic noise. The theory was then implemented in the design and simulation of a superlens imaging system that uses a hyperbolic metamaterial spatial filter to perform the required convolution operation physically to improve the imaging performance. The results do not only indicate the power of superlenses for enhanced sub-diffraction imaging but also the efficacy of the II loss compensation scheme by decently connecting and attempting to resolve two grand issues of optics, namely loss compensation and imaging beyond diffraction limit. The experimental implementation of the imaging method was also discussed.
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Chapter 6

Enhanced Lateral Spatial Resolution in Far-field Imaging

6.1 Introduction

Imaging is an indispensable tool in the toolbox of nearly every field of science, engineering, technology, and medicine. Unfortunately, encoding the desired information into electromagnetic waves imposes a limit to the performance of imaging systems at the outset – the detection of the fields by the interaction of photons (the light signal) and matter (the light detector) means that the signal-to-noise ratio (SNR) for long exposures will always be limited physically by shot noise. A naïve analysis would reveal that adding up more photons in the detector would lead to higher SNR. This is true, however for incoherent light, the magnitude of the optical transfer function for an imaging system with an unobstructed pupil monotonically decreases with increasing spatial frequency[1]. It follows that the spectral SNR then also decreases with increasing spatial frequency, since the shot noise variance is constant in the spatial frequency domain, as will be shown later. Consequently, adding up more photons does not lead to much increase for the SNR of high spatial frequencies. In the lab, one also does not have the freedom to arbitrarily increase the number of photons collected, since at some point the detector will become saturated. Each specific imaging modality will also have its specific limitations. For example, in fluorescence imaging only a certain exposure can be obtained before photobleaching occurs.

Our research in loss compensation for metamaterial and plasmonic imaging systems[2, 3, 4, 5, 6, 7, 8, 9, 10] led us to obtain a unique perspective on the noisy imaging problem. We first determined that the fundamental resolution limit to superresolving lenses is not determined by the diffraction limit, but rather by a shot noise limit, i.e. where the shot noise overcomes the transfer function in the spatial frequency domain. We then found that by rejecting the detection of low spatial frequency harmonics in an object field, the resulting image can have large SNR for high spatial frequencies, due to reduction of the overall noise level and amplification of the high spatial frequencies by increased illumination intensity or exposure. These previous works, along with some inspiration from other research in far-field imaging[11, 12], led us to construct the current work.

In this paper, a theory is presented that shows how to improve the SNR for the high spatial frequencies of an image obtained from an object illuminated with incoherent light. The theory is implemented in numerical simulations to predict the resolution enhancement, and experimental images are collected using a low numerical aperture imaging system to confirm the predictions. Richardson-Lucy deconvolution is used to reconstruct the images. Point spread function measurements are also performed in order to further support the theory. Then end result is an image with higher resolution and improved contrast for the high spatial frequencies as compared to the control image. It is also shown that this method can help prevent pixel saturation for longer exposures.

6.2 Theory

6.2.1 Incoherent Imaging

Consider a uniform beam of spatially incoherent, narrowband light with photon flux density Φ_0 [photons/ $m^2 \cdot s$] striking a planar, transmissive object. After passing through the object, the light has spatial variations and the resulting transmitted photon flux density is given by $O(\mathbf{r})$, with $\mathbf{r} \in \mathbb{R}^2$ denoting the position coordinate on the plane. The process of mapping this light distribution with an imaging system can be represented by a convolution

$$I(\mathbf{r}) = H(\mathbf{r}) * O(\mathbf{r}), \tag{6.1}$$

where $I(\mathbf{r})$ is the photon flux density on the image plane, $H(\mathbf{r})$ is the point spread function of the imaging system, and * denotes the convolution. Here it should be noted that $I(\mathbf{r}), O(\mathbf{r}) \in \mathbb{R}_{\geq 0}$ since they represent flux densities and not complex fields. Fourier transforming Eq. 6.1 gives

$$\tilde{I}(\mathbf{k}) = \tilde{H}(\mathbf{k})\tilde{O}(\mathbf{k}) \tag{6.2}$$

using the corresponding Fourier transform pair

$$\tilde{S}(\mathbf{k}) = \int_{\mathbb{R}^2} S(\mathbf{r}) e^{-i2\pi \mathbf{k} \cdot \mathbf{r}} d^2 \mathbf{r}, \qquad (6.3a)$$

$$S(\mathbf{r}) = \int_{\mathbb{R}^2} \tilde{S}(\mathbf{k}) e^{i2\pi\mathbf{k}\cdot\mathbf{r}} d^2\mathbf{k},$$
(6.3b)

and it becomes evident that in general $\tilde{I}(\mathbf{k})$, $\tilde{H}(\mathbf{k})$, $\tilde{O}(\mathbf{k}) \in \mathbb{C}$. Here, $\tilde{I}(\mathbf{k})$ and $\tilde{O}(\mathbf{k})$ are the image and object spectra, respectively, and $\tilde{H}(\mathbf{k})$ is the optical transfer function (OTF) of the imaging system. For later use, we define $|\tilde{H}(\mathbf{k})|$ as the modulation transfer function (MTF).

6.2.2 Detection and Noise

Eqs. 6.1 and 6.2 assume continuous signals in position and reciprocal space. To incorporate practical detection of the deterministic signal $I(\mathbf{r})$, let us consider the case where we collect an image on the image plane using a photoelectric detector with n_p pixels. The center of the p^{th} pixel ($p \in \mathbb{Z}_+$) is at $\mathbf{r}_p = (x_i, y_j)$, and the pixels are rectangular with side lengths Δx and Δy along the x and y dimensions, respectively. Then, at the p^{th} pixel, the expected number of detected photons is given by

$$\bar{I}_{p} = \eta T \int_{A_{p}} I(\mathbf{r}) d^{2}\mathbf{r}
= \eta T \int_{y_{j}-\Delta y/2}^{y_{j}+\Delta y/2} \int_{x_{i}-\Delta x/2}^{x_{i}+\Delta x/2} I(x,y) dx dy
\approx \eta T \Delta x \Delta y I(x_{i},y_{j}),$$
(6.4)

where η is the quantum efficiency of the pixel and T is the exposure time or integration time. The approximation in the third line of Eq. 6.4 assumes that the signal $I(\mathbf{r})$ is slowly varying across the area of the pixel, i.e. the signal is well sampled. Since Eq. 6.4 describes a sampling of a spatial distribution of discrete particles (photons), there will be an intrinsic randomness due to shot noise in the photon signal $I_{p,\gamma}$. In this case, the probability mass function (PMF) is

$$\mathcal{P}\left(I_{p,\gamma}|\bar{I}_p\right) = \frac{\left(\bar{I}_p\right)^{I_{p,\gamma}}}{I_{p,\gamma}!}e^{-\bar{I}_p}.$$
(6.5)

This is of course coming from the fact that the counting of discrete particles at a constant rate follows a Poisson distribution, for which

$$\operatorname{Var}(I_{p,\gamma}) = \bar{I}_p. \tag{6.6}$$

In most photoelectronic imaging detectors, such as complementary metal-oxidesemiconductor (CMOS) or charge-coupled device (CCD) cameras, there are primarily two sources of noise. The first is due to the statistics of Eq. 6.5, the shot noise which is dependent on the photon signal. The second is noise from the readout electronics, which is independent of the photon signal. We can then write the detected signal as

$$I_p = \bar{I}_p + N_{p,\gamma} + N_{p,e}, \tag{6.7}$$

where $N_{p,\gamma}$ is a discrete random variable representing the shot noise with PMF described by Eq. 6.5, and $N_{p,e}$ is a discrete random variable representing the electronic readout noise.

In order to show how the noise addition in Eq. 6.7 affects the image spectrum, we

compute the discrete Fourier transform

$$\tilde{I}_{q} = \sum_{p} I_{p} e^{-i2\pi \mathbf{k}_{q} \cdot \mathbf{r}_{p}}$$

$$= \sum_{p} (\bar{I}_{p} + N_{p,\gamma} + N_{p,e}) e^{-i2\pi \mathbf{k}_{q} \cdot \mathbf{r}_{p}}$$

$$= \sum_{p} \bar{I}_{p} e^{-i2\pi \mathbf{k}_{q} \cdot \mathbf{r}_{p}} + \sum_{p} N_{p,\gamma} e^{-i2\pi \mathbf{k}_{q} \cdot \mathbf{r}_{p}} + \sum_{p} N_{p,e} e^{-i2\pi \mathbf{k}_{q} \cdot \mathbf{r}_{p}}$$

$$= \tilde{I}_{q} + \tilde{N}_{q,\gamma} + \tilde{N}_{q,e},$$
(6.8)

where $q \in \mathbb{Z}_+$ and $\{\mathbf{k}_q = (k_{x,l}, k_{y,m}) \mid 1 \leq q \leq n_p\} \subset \{\mathbf{k} \in \mathbb{R}^2\}$ is the Fourier space corresponding to the pixelated position space $\{\mathbf{r}_p \mid 1 \leq p \leq n_p\}$. We then consider the statistical properties of $\tilde{N}_{q,\gamma}$ and $\tilde{N}_{q,e}$. Since the shot noise variance is known from Eq. 6.6 (replacing $I_{p,\gamma}$ with $N_{p,\gamma}$) and we can assume the pixels are statistically independent, we can write [10, 11, 12]

$$\operatorname{Var}(\tilde{N}_{q,\gamma}) = \operatorname{Var}\left(\sum_{p} N_{p,\gamma} e^{-i2\pi \mathbf{k}_{q} \cdot \mathbf{r}_{p}}\right) = \sum_{p} \operatorname{Var}(N_{p,\gamma}) |e^{-i2\pi \mathbf{k}_{q} \cdot \mathbf{r}_{p}}|^{2}$$
$$= \sum_{p} \bar{I}_{p} = n_{\gamma}, \quad (6.9)$$

where n_{γ} is the total number of photons in the entire image. In words, Eq. 6.9 states that the variance in Fourier space is constant, and is controlled by the total number of photons collected on the detector. We can write a similar equation for the readout noise[12],

$$\operatorname{Var}(\tilde{N}_{q,e}) = \operatorname{Var}\left(\sum_{p} N_{p,e} e^{-i2\pi \mathbf{k}_{q} \cdot \mathbf{r}_{p}}\right) = \sum_{p} \operatorname{Var}(N_{p,e}) |e^{-i2\pi \mathbf{k}_{q} \cdot \mathbf{r}_{p}}|^{2} = \sum_{p} \sigma_{p,e}^{2}, \quad (6.10)$$

where $\sigma_{p,e}^2$ is the readout noise variance at pixel p and again the assumption is made that the pixels are statistically independent. Let us also assume that $\sigma_{p,e}^2 = \sigma_e^2$, meaning every pixel has similar electrical performance in terms of noise. Then Eq. 6.10 becomes

$$\operatorname{Var}(\tilde{N}_{q,e}) = n_p \sigma_e^2. \tag{6.11}$$

Therefore, the spectral readout noise variance is also a constant, and scales linearly with the number of pixels. For modern cameras, the readout noise is usually minimal such that it is neglected, though we keep it here for completeness.

6.2.3 Tailoring the Spectral SNR

From imaging theory, we know that the optical transfer function of an incoherent imaging system is given by the autocorrelation[1] of the system's pupil function $\tilde{P}(\mathbf{k})$, or

$$\tilde{H}(\mathbf{k}) = \frac{\int \tilde{P}(\boldsymbol{\kappa})\tilde{P}^*(\boldsymbol{\kappa} - \mathbf{k})d^2\boldsymbol{\kappa}}{\int |\tilde{P}(\boldsymbol{\kappa})|^2 d^2\boldsymbol{\kappa}} = \tilde{P}(\mathbf{k}) \star \tilde{P}(\mathbf{k}).$$
(6.12)

In the discrete notation described in the previous section, we can write the autocorrelation as

$$\tilde{H}_q = \frac{\sum_{\kappa} \tilde{P}_{\kappa} \tilde{P}_{\kappa-k_q}^*}{\sum_{\kappa} |\tilde{P}_{\kappa,o}|^2} = \tilde{P}_q \star \tilde{P}_q, \qquad (6.13)$$

where $\tilde{P}_{\kappa,o}$ is a reference pupil. To normalize 6.13, the reference pupil is conventionally chosen as the pupil itself, making the DC pixel of $\tilde{H}_q = 1$, similar to Eq. 6.12. However we define a reference pupil in Eq. 6.13 in order to later directly compare two different pupil functions. Consider an incoherent imaging system in air with maximum resolvable spatial frequency $k = 2\text{NA}k_0$, where NA is the numerical aperture, $k_0 = 1/\lambda_0$, and λ_0 is the center free space wavelength of the illumination source. The pupil function is assumed to have circular symmetry about the optical axis (which is along *z*-direction) and we define it as

$$\tilde{P}_q = \begin{cases} 1 & \text{if } k_- \le |\mathbf{k}_q| \le k_+ \\ 0 & \text{otherwise,} \end{cases}$$
(6.14)

where $k_{-} \geq 0$ and $k_{-} < k_{+} \leq k/2$. From Eq. 6.14 we can see that setting $k_{-} = 0$ and $k_{+} = k/2$ gives a typical diffraction-limited imaging system with open pupil. We choose this case as our reference pupil $\tilde{P}_{q,o}$. However, if we make k_{-} nonzero, we introduce an obstruction in the central portion of the pupil, which has the effect of lowering the overall transmission with respect to the reference pupil, and also preferentially passing larger spatial frequencies with respect to the smaller ones in comparison to the reference pupil.

An important metric for an imaging system is its ability to discern image spectrum content from noise, or its spectral SNR. To relate the pupil function to the spectral SNR, we first use Eqs. 6.2 and 6.13 to rewrite the expected image spectrum as

$$\tilde{\bar{I}}_q = \left[\tilde{P}_q \star \tilde{P}_q\right]\tilde{O}_q.$$
(6.15)

Using a standard definition of SNR, we then can write

$$SNR_q = \frac{|\tilde{I}_q|}{\sqrt{n_\gamma + n_p \sigma_e^2}} = \frac{\left| \left[\tilde{P}_q \star \tilde{P}_q \right] \tilde{O}_q \right|}{\sqrt{n_\gamma + n_p \sigma_e^2}}$$
(6.16)

as the image spectrum SNR. In Eq. 6.16, the numerator is the expected image spectrum, which can be engineered by manipulation of the pupil function, and the denominator is the total standard deviation of the signal from the photonic and electronic noise terms in Eqs. 6.9 and 6.10. An obvious consequence of 6.16 is that reducing n_{γ} will decrease the constant noise floor in the image spectrum. The signal in the numerator will also decrease similarly, but can be engineered through \tilde{P}_q to enhance different portions of the spectrum. This is the central idea of this chapter, engineering the system to preferentially pass a certain band of the image spectrum, reducing the overall noise and improving SNR_q for that band. To put into analogy, you are given a "photo-budget" n_{γ} , and you can freely decide how to spend that budget



Figure 6.1: Imaging experiment configuration. A diffused LED source illuminates an imaging target, which is then processed by a 4f system consisting of two achromatic doublet lenses and a transparency in the Fourier plane. The images are detected with a CMOS camera.

in the image spectrum via \tilde{P}_q so that you achieve an improved SNR for certain spatial frequencies.

6.3 Experiment

In order to implement the above spectral SNR engineering into an experimental imaging system, we simply need access to the Fourier plane in order to manipulate the pupil function. Therefore, we chose to construct a typical 4f system with no magnification as shown in Fig. 6.1 in order to simplify the analysis and experiment. However, it should be emphasized that the concepts presented should be generally applicable to any imaging system which is linear and shift-invariant, provided there is a mechanism with which to manipulate the Fourier content of the light.

The experiment in Fig. 6.1 images an object (target) illuminated by a narrowband



Figure 6.2: (a) Open (reference) and (b) annular pupils with (c) corresponding OTFs calculated from 6.13 as a function of radial spatial frequency $k_r = \sqrt{k_x^2 + k_y^2}$ used in the experiment in Fig. 6.1. Both pupils have $k_+ = k/2$. In (a), $k_- = 0$ and in (b), $k_- = k/2\sqrt{2}$. In the actual experiment, we chose an outer pupil diameter of 5 mm for (a) and (b), making NA= 0.0066.

incoherent light source (Thorlabs LIU525B). Before hitting the target, the light is focused onto a diffuser in order to decrease the spatial coherence and to avoid imaging the light source onto the Fourier plane. Then the light is roughly collimated by a second lens before hitting the target. The light distribution exiting the target is Fourier transformed by an achromatic-doublet lens (Space Optics Research Labs) with a focal length of f = 38.1 cm. On the Fourier plane, a pupil transparency is placed that has either a circular or annular opening, as shown in Fig. 6.2. For all the images, we chose an outer pupil diameter of d = 5 mm, making NA= 0.0066 using the formula NA= d/2f. The transparencies were printed with a photoplotter onto transparent plastic sheets, then cut out and mounted in standard optical mounts. After passing through the pupil on the Fourier plane, the light is again Fourier transformed by a second identical achromatic-doublet which then focuses the resulting image onto a CMOS camera (Thorlabs DCC1645C).

The goal of this experiment was to show directly an enhancement in image resolution by modifying the pupil to improve SNR_q for the largest spatial frequencies in accordance with 6.16. This led to the annular pupil configuration in Fig. 6.2 (b). To show quantitatively the improvement in the resolution performance afforded by the annular pupil configuration over the open pupil, we replaced the resolution target with a 10 μ m diameter pinhole. Since this diameter is below the diffraction limit for the chosen numerical aperture defined by k_+ and f, the resulting image of the pinhole is the PSF of the imaging system. These PSF images were taken with both the open and annular pupil transparencies with varying exposure times. From these images, the transfer functions and corresponding object-independent SNR_q can be computed for each exposure.

6.4 Results

6.4.1 Point Spread Functions and Transfer Functions

Examples of the experimentally measured PSFs and MTFs are presented in Figs. 6.3 and 6.4 for the open and annular pupils, respectively. Also shown are the theoretical PSFs and MTFs determined from scalar diffraction theory. A good agreement can be seen with both the PSF and MTF between theory (black lines) and experiment (red and blue lines), indicating that the imaging system is well aligned and not inducing any unwanted aberrations. Also, the calculated standard deviation (black dashed line) seems to accurately predict where the MTF is overcome by the shot noise, providing evidence supporting our noise theory. Since $\sigma_e = 0$ in this calculation, it is apparent that the readout noise is in fact likely negligible. Setting $\sigma_e = 1$ (and $n_p = 2.5 \times 10^5$) leads to a similar agreement with the experimental spectra, but any larger value begins to deviate from the observed noise level.

6.4.2 Spectral SNR

From the OTFs computed from the measured PSFs in Figs. 6.3 and 6.4, it is then straightforward to compute SNR_q for each exposure using Eq. 6.16, assuming that



Figure 6.3: Measured PSF and MTF for full, unobstructed pupil with NA= 0.0066 and exposure time of T = 2 s. (a) The PSF collected from the setup in Fig. 6.1. (b) The MTF calculated by fast Fourier transformation of (a). (c) Cross-sections of (a) through the origin along x (red line) and y (blue line). The theoretical prediction is given by the black line. (d) Cross-sections of (b) through the origin along k_x (red line) and k_y (blue line). The theoretical prediction is given by the black line. The dashed black line denotes the calculated shot noise standard deviation using Eq. 6.16.

the experimental pixel values and the number of photons at each pixel are about linearly related. These are plotted in Fig. 6.5, where the solid lines indicate the open pupil SNR_q and the dashed lines indicate the annular pupil SNR_q .

For a direct comparison of the SNR_q for the two pupil configurations, we define a



Figure 6.4: Measured PSF and MTF for annular pupil with $k_{-} = k/2\sqrt{2}$, NA= 0.0066, and exposure time of T = 4 s. The subfigures are defined in the same manner as Fig. 6.3.

spectral SNR improvement metric

$$SNRi_q = \frac{SNR_{q,a,T}}{SNR_{q,o,T_0}},$$
(6.17)

where $\text{SNR}_{q,a,T}$ is the annular pupil spectral SNR for exposure time T and SNR_{q,o,T_0} is the open pupil spectral SNR for exposure time T_0 . Plotted in Fig. 6.6 is the SNRi_q for $T_0 = 2$ s and three different values of T. It can be seen that a clear enhancement in the SNR for spatial frequencies near $0.83k_0$ can be obtained using an annular pupil



Figure 6.5: Measured spectral SNR for open (solid lines) and annular (dashed lines) pupils for different exposure times. The crossover by the black solid line over the red and blue lines is due to distortion of the PSF by pixel saturation at T = 3 s for the open pupil.

provided a sufficient exposure.

6.4.3 Test Images

To verify that the high spatial frequency improvement in spectral SNR with sufficient exposure manifests as improved image resolution, we imaged a USAF-1951 resolution test target (Thorlabs R1DS1N) using the setup in 6.1. The collected images for the open and annular pupil are shown in Fig. 6.7 for three values of T. Also shown is the corresponding reconstructions obtained by deconvolving the images with the



Figure 6.6: Spectral SNR improvement SNRi_q corresponding to the data in Fig. 6.5 with $T_0 = 2$ s. The black dashed line indicates $\text{SNRi}_q = 1$.

Richardson-Lucy algorithm as implemented in MATLAB. In Fig. 6.7, it can be seen that three bar target in element 5 is always blurred together by the open pupil, however even for low exposure (e.g. T = 10 ms), the bars become qualitatively visible in the annular pupil case. After reconstruction for 30 iterations, element 5 remains unresolved in the open pupil images, but the annular pupil images of element 5 are further improved. Particularly for elements 2-4, the annular pupil image is improved greatly by the reconstruction since those spatial frequencies were originally attenuated with respect to the open pupil. The spectral SNR for those frequencies remained above 1, so they could still be easily reconstructed computationally.



(a) T = 10 ms

(b) T = 30 ms



(c) T = 50 ms

Figure 6.7: Experimental images of a USAF-1951 resolution test target collected from the setup in Fig. 6.1 with (a) T = 10 ms, (b) T = 30 ms, and (c) T = 50 ms exposure times. The individual images in each subfigure correspond to the following: i. Raw image collected with open pupil in the Fourier plane and NA=0.0066. ii. Raw image with annular pupil in the Fourier plane and same NA as i. iii. Image from part i deconvolved by the Richardson-Lucy algorithm after 30 iterations. iv. Image from part ii deconvolved by the Richardson-Lucy algorithm after 30 iterations.

6.5 Discussion

6.5.1 Pixel Saturation

Upon viewing the open pupil SNR_q from Fig. 6.5, it would seem that simply increasing T would lead to an improvement in SNR_q itself, without having to modify the pupil. However, the pixels in typical digital cameras only have finite well depth and dynamic range, meaning they can experience saturation for long exposures and/or intense illumination. The saturation causes a nonlinear response of the pixel as a function of the input photon signal. Therefore, one cannot arbitrarily increase T or illumination intensity to increase SNR_q . In terms of spatial resolution, saturation can manifest as an effective blurring due to clipping of the pixel values and blooming of photoelectrons to adjacent pixels.

Along with provided improved resolution and contrast for larger spatial frequencies, the proposed method of engineering SNR_q for high spatial frequencies can also provide resistance to pixel saturation in cases when long exposure or intense illumination *and* high resolution is required. We collected images in which the pixels become saturated for the open pupil system, and compared them to images of similar exposure in the annular pupil system. They are shown in Fig. 6.8. Since a portion of the Fourier



Figure 6.8: Images demonstrating the resistance of annular pupil to deleterious effects caused by detector saturation. Parts i-iv are defined similar to Fig. 6.7. For the open pupil, T = 150 ms, and for the annular pupil, T = 300ms. It can be seen that even for twice as long exposure time, the annular pupil image quality is mostly maintained compared to the open pupil image, which is severely blurred due to pixel saturation.

plane is blocked by the annular pupil, the total number of photons reaching the detector is decreased. As evidenced by Fig. 6.2, the blocked photons correspond to lower spatial frequencies which are more likely to contribute to pixel saturation, since the transmission for these portions of the object will be high due to the larger local photon flux.

6.5.2 Relationship to Previous Works

We have previously published some articles in which we utilize a high spatial frequency passband function, in conjunction with an increased exposure, to enhance the resolution performance of thin metal films acting as near-field plasmonic "superlenses." [5, 7, 8, 9, 10] This passband function was used to not only compensate
the losses inherent in the metal film, but also to improve the image spectrum SNR similar to what we have shown in this chapter. We have called this method 'active convolved illumination' for a couple reasons. First, the term 'active' was chosen since an increased illumination intensity or exposure time is needed to obtain enhancement over the control (a bare superlens). Secondly, the term 'convolved' was chosen since the applied passband function is physically convolved with the fields emanating from the object.

For coherent illumination, this passband function can easily be realized by a type-II hyperbolic metamaterial (HMM). While there is no "Fourier plane" in the near-field configuration, the HMM can modify the Fourier components of the incident evanescent waves by its dispersion. In this case, the "pupil function" can then simply be thought of as the OTF of the cascaded HMM-superlens system, which for the p-polarization passes evanescent waves with large spatial frequency and rejects low spatial frequencies. Subsequently, the SNR for the large spatial frequencies is enhanced with respect to the control. In fact, a principal goal of the experimental work of this chapter was to show that our theories both accurately represent real noisy optical signals, and can be straightforwardly extended to conventional imaging systems. Put briefly, the theory and experiment we have applied to far-field incoherent imaging systems in this chapter is for the most part an extension of our previously published methods for near-field superresolution enhancement. In the next section, we present another such enhanced superresolution system, though the image detection can now be performed in the far-field.

6.6 An application to Superresolution Imaging

6.6.1 Far-field Superlensing

To explore the versatility of this idea, we chose an extreme realm of imaging in which sub-diffraction-limited objects can be imaged with a specially designed device which transfers their spatial content to a far-field detector. There have been numerous studies of so-called "far-field superlenses" published previously[13, 14, 15]. These lenses usually operate by resonantly exciting surface plasmons on a thin metal film, and then diffracting the surface waves into free-space propagating waves with a subwavelength grating structure. However, there are some issues with these structures which preclude their application to more arbitrary scenarios. First, the working wavelength cannot be chosen freely, since the frequency at which the surface plasmon resonance occurs is essentially fixed for a silver slab in a given dielectric medium. Second, the signal strength of the propagating Fourier components in the object must be small compared to the evanescent components in order to properly reconstruct the object. In other words, if the waves from the zeroth and first diffraction orders have similar strength, the relative values of the two orders cannot be determined after detection, since they will be overlapped in Fourier space. These problems were partially solved by a new design which replaced the single metal film with a metal-dielectric multilayered structure[16]. In this design, the dispersion of the multilayered material, essentially a HMM, can be easily tuned by changing the working wavelength, constituent materials, and geometric parameters of the layers. We briefly present here a similar implementation of a HMM-superlensing device that can be thought of as applying the SNR engineering idea from the previous sections for coherent light.

6.6.2 Theory

First, we explore the relevant theory for multilayered HMMs which can guide the design of our our high-pass "pupil function." The perpendicular (with respect to the optical axis along z) and parallel components of the permittivity tensor for a layered HMM are given by[17]

$$\epsilon_{\perp} = \frac{\epsilon_m t_m + \epsilon_d t_d}{t_m + t_d} \tag{6.18}$$

and

$$\epsilon_{\parallel} = \frac{t_m + t_d}{t_m / \epsilon_m + t_d / \epsilon_d},\tag{6.19}$$

where t_m , t_d are the metal and dielectric layer thicknesses, respectively, and ϵ_m , ϵ_d and the metal and dielectric permittivities. From the permittivity tensor, the dispersion relation for the TM-polarization can be derived as

$$\frac{k_{\perp}^2}{\omega^2 \mu \epsilon_{\parallel}} + \frac{k_z^2}{\omega^2 \mu \epsilon_{\perp}} = 1.$$
(6.20)

Inspection of Eq. 6.20 reveals that if $\epsilon_{\parallel} > 0$ and $\epsilon_{\perp} < 0$, for small k_{\perp} , k_z becomes imaginary and the waves in the HMM are evanescent. If we set $k_z = 0$, we can find the k_{\perp} at which the waves are no longer evanescent inside the material. The result is a "cutoff" k_{\perp} given by

$$k_{\perp,c} = \pm \omega \sqrt{\mu \epsilon_{\parallel}},\tag{6.21}$$

which tells the minimum value of k_{\perp} propagating waves can have inside the HMM. However, real metallic films will introduce an imaginary part to the permittivity of the HMM, making the ϵ_{\parallel} and ϵ_{\perp} complex. Then the dispersion relation becomes

$$\frac{k_{\perp}^2}{\omega^2 \mu(\epsilon'_{\parallel} + i\epsilon''_{\parallel})} + \frac{k_z^2}{\omega^2 \mu(\epsilon'_{\perp} + i\epsilon''_{\perp})} = 1$$
(6.22)

with corresponding cutoff condition

$$k_{\perp,c,\text{lossy}} = \pm \omega \sqrt{\mu(\epsilon'_{\parallel} + i\epsilon''_{\parallel})}.$$
(6.23)

From Eq. 6.23, it can be seen that $k_{\perp,c,\text{lossy}}$ is actually a complex quantity. The physical consequence of this is that there is some allowed transmission of the waves with $|k_{\perp}| < |k_{\perp,c}|$. Plotted in Fig. 6.9 is the dispersion for a lossy HMM along with



Figure 6.9: Dispersion plots for a multilayered HMM with $t_m = t_d = 15$ nm, $\epsilon_m = -7.241 + i0.248$, and $\epsilon_d = 4.285$. The black solid lines in the right plot show $\text{Re}[k_z]$ and red dashed lines indicate $\text{Re}[k_{\perp,c,\text{lossy}}]$.

the calculated $\operatorname{Re}[k_{\perp,c,\text{lossy}}]$. It can be seen that the complex value of ϵ_m introduces not only absorption in the HMM, but also changes the nature of it's dispersion.

6.6.3 Design and Simulation

An example HMM far-field superlensing device is shown in Fig. 6.10. This design gives a huge parameter space for us to explore, as the materials, layer thicknesses, grating parameters, and working wavelength can all be modified. To narrow down the design space, we first chose a wavelength of $\lambda_0 = 488$ nm as a starting point for the device. This blue wavelength is the same as that emitted by argon ion lasers and commercially available solid state lasers. At this wavelength, the permittivity for silver was determined to be $\epsilon_m = -7.241 + i0.248$ using the most accurate available ellipsometry measurements of pure silver films[18, 19]. As the dielectric material,



Figure 6.10: Hyperbolic diffractive "far-field superlens" design. The waves emanating from the object plane are filtered by a HMM consisting of 20 unit cells. The filtered waves are then coupled into the far-field by a diffraction grating with thickness t_g , period d, and metal fill ratio f. The gray portions indicate silver and the blue portions indicate Si₃N₄.

we chose Si₃N₄, which at $\lambda_0 = 488$ nm has a permittivity of $\epsilon_d = 4.285[20]$. Using rigorous coupled-wave analysis[21], we calculated the transmittance of the p- and spolarizations in Fig. 6.11 for a HMM with 20 unit cells embedded in a background medium with permittivity $\epsilon_b = 2.25$. After designing the HMM which defines our "pupil function," the next step is to design a grating which will efficiently transfer the Fourier components passed by the HMM to the far-field via the -1 diffraction order. Since Re[$k_{\perp,c,\text{lossy}}$] = $4.57k_0$, from the grating equation we can determine that a grating period of d = 110 nm should give minimal overlap of the transmitted Fourier components while still maximizing the amount of information contained in the freespace passband. Then the metal fill ratio f = 0.3 and grating thickness $t_g = 85$



Figure 6.11: HMM transmittance for (a) p-polarization and (b) spolarization. The red dashed lines in (a) denote the predicted cutoff value of Re[$k_{x,c,\text{lossy}}$]. The number of unit cells (metal-dielectric layer pairs) in the HMM is 20 and the background permittivity is $\epsilon_b = 2.25$.

nm were chosen to give good overall transmission. Again using rigourous coupledwave analysis, we calculated the transfer functions of the superlens for the p- and s-polarizations. These are shown in Fig. 6.12. In Fig. 6.12 (a), the zeroth order transfer function magnitude for p-polarized light is plotted, giving the ratio of the input wave to output wave as a function of k_x and k_y . Similarly in (b), the transfer function is shown for p-polarized light, except the output waves are shifted so the output $k_x = k_{x,in} \pm (2\pi/d)$, giving the first order transfer function. For $k_{x,in} < 0$, the + is taken, while for $k_{x,in} > 0$ the - is taken. These wave components can be directly calculated from the RCWA code. In (c) and (d), the zeroth and first order transfer functions are plotted for s-polarized light. It can be seen that no transmission occurs beyond the free space passband for s-polarization.

Using the transfer functions in Fig. 6.12, we performed an imaging simulation in which a double-slit object with 45 nm separation and 45 nm slit opening size is



Figure 6.12: (a) Zeroth and (b) first order transfer functions of the designed superlens for p-polarization. (c) Zeroth and (d) first order transfer functions for s-polarization.

illuminated and imaged by the designed superlens. From Fig. 6.13, it can be seen that the object with characteristic size smaller than $\lambda_0/10$ can be reconstructed from the data the superlens transfers to the detector through the p- and s-polarizations. Through the s-polarized zeroth order transfer function, the free space passband can be detected as indicated in Fig. 6.12 (c). This alone gives the diffraction-limited result indicated by the blue lines in Fig. 6.13 after deconvolution. To obtain the higher spatial frequencies, p-polarized light is provided, and the object information is encoded into the first order transfer function in Fig. 6.12 (b). After obtaining this



Figure 6.13: (a) Double-slit object, diffraction-limited image, and reconstructed superlens image spectra, (b) reconstructed superlens image, and (c) 1D object and image cross-sections. The legend in (c) also applies to (a).

diffraction-limited image, the detected Fourier components must be shifted back to their original position by a factor of $\pm 2\pi/d$ and deconvolved, essentially reversing the effects of diffraction and attenuation through the superlens. The result of this procedure, after being added with the low frequency waves obtained from the spolarization, is given by the dashed red lines in Fig. 6.13, showing reconstruction of the high- k_x Fourier components and successful resolution of the double-slit object.

6.7 Conclusion

In this chapter, a resolution enhancement method for far-field imaging systems was presented and demonstrated experimentally. This method relies on manipulation of the Fourier plane in the imaging system in order to engineer the image spectrum SNR. Placing an annular pupil transparency in the Fourier plane leads to improved SNR for larger spatial frequencies, which gives improved resolution over an imaging system with unobstructed pupil and the same numerical aperture. Also, the annular pupil increases resistance to pixel saturation, since a portion of the Fourier plane is blocked and the number of photons at the detector is decreased. Also, the connection between the proposed method and previous metamaterial implementations is discussed. These methods are then implemented into the design of a new far-field superlensing device which can achieve resolution of an object with edge-to-edge separation less than $\lambda/10$.

In summary, the contributions of this dissertation to the body of knowledge by the author are the theory and simulation of the proposed loss mitigation and resolution enhancement methods for incoherent light, the experiment to verify and expand upon the proposed methods, and the spectral SNR equation (Eq. 6.16) which quantitatively shows how to manipulate the pupil or transfer function of an imaging system to achieve a desired high spatial frequency SNR.

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